Научном већу Института за физику у Београду

Предмет: Покретање поступка за избор у звање научни сарадник

Молим Научно веће Института за физику у Београду да покрене поступак за мој избор у звање научни сарадник.

У прилогу достављам:

1. мишљење руководиоца пројекта са предлогом комисије за избор у звање;
2. стручну биографију.
3. преглед научне активности;
4. елементе за квалитативну и квантитативну оцену научног доприноса;
5. списак и копије објављених научних радова и других публикација;
6. податке о цитираности;
7. уверење о одбрањеној докторској дисертацији.

С поштовањем,

Ана Худомал
истраживач сарадник

Ana Hudomal

Научном већу Института за физику у Београду

Предмет: Мишљење руководиоца лабораторије о избору др Ане Худомал у звање научни сарадник

Др Ана Худомал је запослена у Лабораторији за примену рачунара у науци, у оквиру Националног центра изузетних вредности за изучавање комплексних система Института за физику у Београду. У истраживачком раду се бави динамиком ултрахладних бозонских система у оптичким решеткам. С обзиром да испуњава све предвиђене услове у складу са Правилником о поступку, начину вредновања и квантитативном исказивању научноистраживачких резултата истраживача МПНТР, сагласан сам са покретањем поступка за избор др Ане Худомал у звање научни сарадник.

За састав комисије за избор др Ане Худомал у звање научни сарадник предлажем:

(1) др Ивана Васић, виши научни сарадник, Институт за физику у Београду,
(2) др Антун Балаж, научни саветник, Институт за физику у Београду,
(3) проф. др Божидар Николић, ванредни професор Физичког факултета Универзитета у Београду.

др Антун Балаж
научни саветник
руководилац Лабораторије за
примену рачунара у науци

2. БИОГРАФСКИ ПОДАЦИ

Ана Худомал је рођена 8. 3. 1991. године у Београду, где је завршила основну школу и Математичку гимназију. Основне академске студије на Физичком факултету Универзитета у Београду, смер Теоријска и експериментална физика, започела је 2010. године и завршила јула 2014. године са просечном оценом 10,0. Мастер академске студије на истом факултету, смер Теоријска и експериментална физика, завршила је октобра 2015. године са просечном оценом 10,0, одбранивши мастер рад на тему **„***New Periodic Solutions to the Three-Body Problem and Gravitational Waves***”** (наслов на српском: **„**Нова периодична решења проблема три тела и гравитациони таласи**”**). Мастер рад је урађен под руководством др Вељка Дмитрашиновића, научног саветника Института за физику у Београду. Новембра 2015. године уписала је докторске академске студије на Физичком факултету Универзитета у Београду, ужа научна област физика кондензоване материје. Докторску дисертацију под насловом „*Numerical study of quantum gases in optical lattices and in synthetic magnetic fields*” (наслов на српском: „Нумеричко проучавање квантних гасова у оптичким решеткама и у синтетичким магнетним пољима”) урадила је под менторством др Иване Васић, вишег научног сарадника Института за физику у Београду, и одбранила децембра 2020. године.

Од марта 2016. године Ана Худомал је запослена у Институту за физику у Београду као истраживач приправник у Лабораторији за примену рачунара у науци, у оквиру Националног центра изузетних вредности за изучавање комплексних система. Априла 2019. године изабрана је у звање истраживач сарадник. Од марта 2016. до децембра 2019. године била је ангажована на пројекту основних истраживања **„***Моделирање и нумеричке симулације сложених вишечестичних система***”** **(**ОН171017) Министарства просвете, науке и технолошког развоја Републике Србије, којим је руководио др Антун Балаж, а од јануара 2020. је ангажована институционално. Поред тога, учествовала је и на билатералним пројектима са Немачком (Гете универзитет у Франкфурту) и са Хрватском (Универзитет у Загребу), под руководством др Иване Васић. До сада је похађала неколико школа за докторанде, међу којима су International School on Thermal, Quantum, and Topological Phase Transitions – Bad Endorf, Germany (2016), Winter School on Topological Matter in Artificial Gauge Fields – Dresden, Germany (2018), ICAP2018 Summer School – Barcelona, Spain (2018), International School and Workshop Anyon Physics of Ultracold Atomic Gases – Kaiserslautern, Germany (2018), New Developments in Topological Condensed Matter – Les Houches, France (2019), Virtual Winter School on Strongly Correlated Quantum Matter (2020).

Од 2007. до 2011. године Ана Худомал је била стипендиста Републичке фондације за развој научног и уметничког подмлатка, затим од 2011. до 2013. стипендиста града Београда, док је од 2013. до 2015. била стипендиста Фонда за младе таленте Републике Србије. Током академске 2014/2015. године учествовала је у извођењу наставе на Физичком факултету Универзитета у Београду, као сарадник у настави на предмету Квантна теоријска физика (предметни наставник доц. др Душко Латас).

До сада, Ана Худомал има један рад објављен у часопису категорије М21а, четири рада објављена у часописима категорије М21, као и 11 саопштења са међународних скупова штампаних у изводу (категорије М34). Из теме доктората, Ана Худомал је објавила 3 рада категорије М21 и 11 саопштења категорије М34.

3. ПРЕГЛЕД НАУЧНЕ АКТИВНОСТИ

На докторским студијама Ана Худомал се у свом научном раду бавила проучавањем особина ултрахладних квантних гасова у оптичким решеткама. Оптичке решетке су стојећи светлосни таласи који представљају периодични потенцијал за атоме и омогућавају квантне симулације физичких модела налик моделима физике кондензоване материје. Основни бозонски модел који се реализује на овај начин је Бозе-Хабард модел и рана истраживања су била посвећена проучавању квантног фазног дијаграма овог модела. У новијим експериментима фокус је на уопштењу физичких модела који се могу проучавати у оптичким решеткама и на разумевању њихове неравнотежне динамике.

Посебно, једна од веома важних и актуелних тема је реализација синтетичких магнетних поља у оптичким решеткама. Почев од 2013. године, јака синтетичка магнетна поља се успешно реализују на посредан начин, у периодично вођеним оптичким решеткама. Коришћењем овог метода реализован је један од основних модела физике кондензованог стања, Харпер-Хофштетер модел. Важна карактеристика овог модела су тополошке инваријанте енергетских зона, тзв. Чернови бројеви, који су у основи целобројног квантног Холовог ефекта. Поврх тога, познато је да при одређеним густинама и за довољно јаке интеракције овај модел има фракционо Холово стање као основно стање. Реализација фракционих Холових стања у систему ултрахладних атома је прекретница која је дефинисана још на самом настанку читаве области, а на коју се и даље чека.

Подстакнута поменутим експерименталним напретком, Ана Худомал је најпре истраживала улогу слабих атомских интеракција при реализацији Харпер-Хофштетер модела у вођеним оптичким решеткама. Недавно је први пут измерен Чернов број тополошке енергетске зоне у систему хладних атома и за успешну интерпретацију резултата било је неопходно извршити додатна мерења како би ефекти интеракција били узети у обзир. Применом апроксимативног аналитичког развоја по инверзној фреквенцији вођења, кандидаткиња је прво показала присуство додатних чланова у ефективном Хамилтонијану и испитала њихов утицај на тополошке карактеристике ефективног Хамилтонијана. Потом су доприноси интеракција урачунати употребом теорије средњег поља. Користећи нумеричке симулације и аналитичке увиде, кандидаткиња је утврдила да интеракције доприносе атомским прелазима између различитих енергетских зона ефективног модела, чиме се усложњава експериментална процедура у складу са очекивањима. Међутим, добијени резултати такође показују да слабе атомске интеракције олакшавају мерење Черновог броја на неколико начина. Како се очекује да мерење Черновог броја постане рутински алат у блиској будућности - први корак у припреми занимљивијих тополошких фаза - досада добијени резултати о ефектима слабих интеракција су од значаја и за будуће експерименте.

У следећој фази истраживања, кандидаткиња је разматрала сличан модел, али у режиму јаких интеракција при густинама на којима можемо очекивати појаву фракционих Холових стања. Најпре је разматран неизбежан процес загревања оваквог система, тј. одређене су реалистичне фреквенције вођења у зависности од јачине локалних интеракција при којима је загревање на бесконачну температуру довољно споро. У овом разматрању почетно стање је припремљено као основно стање идеалног ефективног модела, које је потом изложено пуном вођењу. Поред тога, за довољно мале системе, нумерички је конструисан стробоскопски еволуциони оператор и одређено је при којим најнижим фреквенцијама вођења нека од својствених стања овог оператора одговарају фракционим Холовим стањима. Од највећег експерименталног значаја је реалистична припрема ових стања. Показано је да се полазећи од једнодимензионалних низова, постепеним укључивањем тунелирања и вођења, фракциона Холова стања могу припремити и мерити на релевантним временским скалама реда десетина милисекунди. Резултати ових истраживања представљени су у публикацијама

- **A. Hudomal**, I. Vasić, H. Buljan, W. Hofstetter, and A. Balaž,
  Dynamics of weakly interacting bosons in optical lattices with flux,

Phys. Rev. A **98**, 053625 (2018).
DOI: 10.1103/PhysRevA.98.053625; ISSN 2469-9926; **IF(2018)=2.907**

- **A Hudomal,** N. Regnault, and I. Vasić,
  Bosonic fractional quantum Hall states in driven optical lattices,
  Phys. Rev. A **100**, 053624 (2019).
  DOI: 10.1103/PhysRevA.100.053624; ISSN 2469-9926; **IF(2019)=2.777**

Други важан фокус у актуелним истраживањима ултрахладних атома је на разумевању неравнотежне динамике квантних система. Ултрахладни атоми су веома добро изоловани квантни системи, па је њихова динамика заиста описана унитарним оператором на експерименталним временским скалама. Добијени експериментални резултати и теоријски аргументи указују да се интеграбилни системи типично брзо термализују. Међутим, у скорашњим експериментима изведеним на квантном симулатору ултрахладних Ридбергових атома примећена је неочекивано спора термализација одређених почетних стања, што је објашњено појавом специјалних својствених стања ефективног модела, тзв. квантних ожиљака (quantum scars). Са циљем проналажења и разумевања режима споре термализације, кандидаткиња је разматрала уопштене бозонске Хамилтонијане у којима амплитуда тунелирања из међу два суседна чвора оптичке решетке зависи од њихове насељености. Анализом корелационих функција и ентропије увезаности показала је да се одређене почетне конфигурације веома споро термализују. Ова појава је објашњена појавом квантних ожиљака, чије особине су потом детаљно анализиране семианалитичком апроксимацијом. За неке од ових модела показано је и присуство експоненцијално великог броја нултих мода. Резултати су представљени у раду:

- **A. Hudomal,** I. Vasić, N. Regnault, and Z. Papić,
  Quantum scars of bosons with correlated hopping,
  Commun. Phys. 3, 99 (2020).
  DOI: 10.1038/s42005-020-0364-9; ISSN 2399-3650; **IF(2019)=4.684**

Током мастер студија, истраживање Ане Худомал било је фокусирано на проналажење и карактеризацију периодичних решења за класичан проблем три тела која међусобно интерагују гравитационом силом. Њутнов проблем три тела је један од најстаријих нерешених проблема физике и математике, формулисан у 17. веку, а до првих решења за периодичне орбите дошли су Ојлер и Лагранж у 18. веку. У последњим деценијама употреба модерних рачунара је омогућила интензиван нумерички приступ и систематизацију нових класа решења. У сарадњи са др Вељком Дмитрашиновићем и др Милованом Шуваковом са Института за физику у Београду, Ана Худомал је испитивала гравитационе таласе који се могу придружити новооткривеним периодичним орбитама, што је од директног значаја за детекцију предвиђених решења. Поред тога, Ана је радила на потврди и објашњењу линеарне везе између периода орбите и њеног одговарајућег тополошког индекса, која представља уопштену верзију трећег Кеплеровог закона. Резултати ових истраживања представљени су у публикацијама

- V. Dmitrašinović, M. Šuvakov, and **A. Hudomal**,
  Gravitational Waves from Periodic Three-Body Systems,
  Phys. Rev. Lett. **113**, 101102 (2014).
  DOI: 10.1103/PhysRevLett.113.101102; ISSN 1079-7114; **IF(2014)=7.512**
- V. Dmitrašinović, **A. Hudomal**, M. Shibayama, and A. Sugita,
  Linear Stability of Periodic Three-body Orbits with Zero Angular Momentum and Topological Dependence of Kepler's Third Law: a Numerical Test,
  J. Phys. A: Math. Theor. 51, 315101 (2018).
  DOI: 10.1088/1751-8121/aaca41; ISSN 1751-8121; **IF(2014)=2.110**

4. ЕЛЕМЕНТИ ЗА КВАЛИТАТИВНУ ОЦЕНУ НАУЧНОГ ДОПРИНОСА

4.1. Квалитет научних резултата

4.1.1 Значај научних резултата

У току докторских студија, кандидаткиња др Ана Худомал је проучавала бозонску динамику у оптичким решеткама, са циљем проналажења режима споре термализације у којима је могуће припремити и испитати тополошки нетривијална стања.

Најважнији рад кандидаткиње је

- **A. Hudomal**, I. Vasić, N. Regnault, and Z. Papić, Quantum scars of bosons with correlated hopping, Commun. Phys. **3**, 99 (2020), цитиран 5 пута

У овом раду, разматрањем класе бозонских модела у којима амплитуда тунелирања између два суседна чвора оптичке решетке зависи од њихове насељености, кандидаткиња је показала да се спора термализација и тзв. кванти ожиљци могу појавити у систему са слабим кинетичким ограничењима.

У другом делу истраживања кандидаткиња је испитала улогу атомских интеракција на мерење Черновог броја у вођеној оптичкој решетки. Добијени резултати показују да, сем што доводе до загревања система, интеракције на одређене начине такође поспешују процедуру мерења. Како су слабе интеракције између атома скоро увек присутне, а очекује се да овакво мерење Черновог броја постане рутински корак у будућим експериментима, од великог је значаја разумети могуће последице њиховог дејства.

Поред тога, кандидаткиња је одредила оптималне микроскопске параметре за припрему и мерење бозонских Лафлинових стања у периодично вођеном Бозе-Хабард моделу. Ово је једно од главних питања у области, на које још увек не постоји дефинитиван одговор у научној литератури, па су резултати кандидаткиње од великог значаја.

За време мастер студија, кандидаткиња др Ана Худомал се бавила истраживањима везаним за класични проблем три тела. У фокусу ових истраживања било је проналажење нових периодичних решења и њихова карактеризација. Кандидаткиња је такође испитивала гравитационе таласе емитоване из периодичних система три тела, што је резултирало публикацијом у престижном часопису Physical Review Letters.

4.1.2 Параметри квалитета часописа

Кандидаткиња др Ана Худомал је објавила укупно пет радова у међународним часописима и то:

- 1 рад у међународном часопису изузетних вредности *Physical Review Letters* (ISSN 1079-7114; IF(2014)=7.512, SNIP(2014) = 2.65)
- 1 рад у врхунском међународном часопису *Communications Physics* (ISSN 2399-3650; IF(2019)=4.684, SNIP(2019) =1.28)
- 2 рада у врхунском међународном часопису *Physical Review A* (ISSN 2469-9926; IF(2019)=2.777 , SNIP(2019) = 0.94, IF(2018)=2.907, SNIP(2018) = 0.99)
- 1 рад у врхунском међународном часопису *Journal of Physics A: Mathematical and Theoretical* (ISSN 1751-8121; IF(2018)=2.110, SNIP(2018) = 0.93)

Библиометријски показатељи су сумирани у наредној табели.

	ИФ	М	СНИП
Укупно	19,99	42	6,79
Усредњено по чланку	4,00	8,4	1,36
Усредњено по аутору	5,71	11,6	1,95

4.1.3 Позитивна цитираност научних радова

Према бази *Web of Science* на дан 9. децембар 2020. године, радови кандидаткиње су цитирани укупно 20 пута, од чега 19 пута без аутоцитата. Према истој бази, Хиршов индекс кандидаткиње је 2.

4.1.4 Међународна сарадња

Међународне активности др Ане Худомал обухватају:

- учешће на билатералним проејктима Србије и Немачке, (BKMH, QDDB, BEC-L) у периоду од 2016. до 2020. године,
- учешће на билатералном пројекту Србије и Хрватске TOP-FOP у периоду од 2016. до 2017. године,
- учешће у COST акцији CA16221 Quantum Technologies with Ultra-Cold Atoms.

Такође, др Ана Худомал је остварила сарадњу са др Николом Рењоом са Универзитета Принстон, САД и са др Златком Папићем са Универзитету у Лидсу, Велика Британија.

4.2. Нормирање броја коауторских радова, патената и техничких решења

Радови кандидаткиње припадају категорији радова са комплексним нумеричким симулацијама и имају пет или мање коаутора, па се признају са пуним бројем М бодова.

4.3. Учешће у пројектима, потпројектима и пројектним задацима

Кандидаткиња је учествовала на следећим пројектима:

- пројекат Министарства просвете, науке и технолошког развоја Републике Србије ОН171017 “Моделирање и нумеричке симулације сложених вишечестичних система” (март 2016. - децембар 2019.).

Кандидаткиња је члан Центра изузетних вредности за изучавање комплексних система, где је ангажовна на потпројекту Динамика ултрахладних атома, којим руководи др Ивана Васић.

4.4. Педагошки рад

Ана Худомал је током академске 2014/2015. године учествовала у извођењу наставе на Физичком факултету Универзитета у Београду, као сарадник у настави на предмету Квантна

теоријска физика (предметни наставник доц. др Душко Латас). У току школске 2016/17. године учествује у раду Државне комисије за такмичења ученика средњих школа из физике као аутор задатака за четврти разред.

4.5. Утицај научних резултата

Утицај научних резултата кандидаткиње се огледа у броју цитата који су наведени у тачки 4.1.3 овог прилога, као и у прилогу о цитираности. Значај резултата кандидаткиње је такође описан у тачки 1.

4.6. Конкретан допринос кандидата у реализацији радова у научним центрима у земљи и иностранству

Кандидаткиња је све своје истраживачке активности реализовала у Институту за физику у Београду. Кандидаткиња је први аутор три публикације, други и трећи аутор по једног рада. Кандидаткиња је дала кључни допринос у свим објављеним радовима који се огледа у изради потребних нумеричких симулација, интерпретацији и презентацији нумеричких резултата, развијању аналитичких апроксимација, писању радова и комуникацији са уредницима и рецензентима часописа.

5. ЕЛЕМЕНТИ ЗА КВАНТИТАТИВНУ ОЦЕНУ НАУЧНОГ ДОПРИНОСА КАНДИДАТА

Остварени резултати у периоду након одлуке Научног већа о предлогу за стицање претходног научног звања:

Категорија	**М бодова по раду**	**Број радова**	**Укупно М бодова***
M21a	10	1	10
M21	8	4	32
M34	0.5	11	5.5
M70	6	1	6

Поређење са минималним квантитативним условима за избор у звање научни сарадник:

Минимални број М бодова		Остварено, М бодова без нормирања*
Укупно	16	**53.5**
M10+M20+M31+M32+M33+M41+M42	10	**42**
M11+M12+M21+M22+M23	6	**42**

Радови кандидаткиње припадају категорији радова са комплексним нумеричким симулацијама и имају пет или мање коаутора, па се признају са пуним бројем М бодова. Према бази *Web of Science* на дан 9. децембар 2020. године, радови кандидаткиње су цитирани укупно 20 пута, од чега 19 пута без аутоцитата. Према истој бази, Хиршов индекс кандидаткиње је 2.

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  DOI: 10.1103/PhysRevLett.113.101102; ISSN 1079-7114; **IF(2014)=7.512**

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  DOI: 10.1103/PhysRevA.100.053624; ISSN 2469-9926; **IF(2019)=2.777**
- **A. Hudomal**, I. Vasić, N. Regnault, and Z. Papić,
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Gravitational Waves from Periodic Three-Body Systems

V. Dmitrašinović and Milovan Šuvakov
Institute of Physics, University of Belgrade, Pregrevica 118, Zemun, Post Office Box 57, 11080 Beograd, Serbia

Ana Hudomal
Fizički fakultet, University of Belgrade, Studentski Trg 12, 11000 Belgrade, Serbia

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Three bodies moving in a periodic orbit under the influence of Newtonian gravity ought to emit gravitational waves. We have calculated the gravitational radiation quadrupolar waveforms and the corresponding luminosities for the $13+11$ recently discovered three-body periodic orbits in Newtonian gravity. These waves clearly allow one to distinguish between their sources: all $13+11$ orbits have different waveforms and their luminosities (evaluated at the same orbit energy and body mass) vary by up to 13 orders of magnitude in the mean, and up to 20 orders of magnitude for the peak values.

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Direct detection of gravitational waves [1,2] ought to come about in the foreseeable future, due to the substantial effort made at the operational and/or pending detectors. One of the most promising candidates for astrophysical sources of gravitational waves are the coalescing, i.e., inspiraling and finally merging binary compact stars [3,4]. Binary coalescence is the only source for which there is a clear prediction of the signal and an estimate of the detection distance limit, as general relativists have completed numerical simulations of mergers of compact binaries, such as neutron stars and/or black holes, Refs. [5–7].

Slowly changing, quasiperiodic two-body orbits are weak sources of gravitational radiation, Refs. [8,9]—only accelerated collapse leads to an increase in energy loss. The major part of the emitted energy in a binary coalescence comes from the final merger of two neutron stars, or black holes, that produces an intense burst of gravitational radiation. Of course, such mergers are one-off events, never to be repeated in the same system, so their detection is subject to their (poorly known) distribution in our Galaxy. It is therefore interesting to look for periodic sources of intense gravitational radiation.

There is now a growing interest in three-body systems as astrophysical sources of gravitational waves, Refs. [10–12]. These early works did not find a substantial increase in the luminosity (emitted power) from representative three-body orbits belonging to three families that were known at the time, Refs. [13–22], over the luminosity from a comparable periodic two-body system [23]. The luminosity of a (quadrupolar) gravitational wave is proportional to the square of the third time derivative of the quadrupole moment, see Refs. [8,9], which, in turn, is sensitive to close approaches of two bodies in a periodic orbit [24]. Thus, getting as close as possible to a two-body collision without actually being involved in one, is a desirable property of the radiating system.

Recently 13 new distinct periodic orbits belonging to 12 (new) families have been discovered in Ref. [25], as well as 11 "satellite orbits" in the figure-eight family [26]. Some of these three-body orbits pass very close to binary collisions and yet avoid them, so they are natural candidates for periodic sources of intense gravitational radiation.

In this Letter we present our calculations of quadrupolar waveforms, Fig. 1, and of luminosities, see Table I and Fig. 2 of gravitational radiation emitted by the $13+11$ recently discovered periodic three-body gravitating orbits, Refs. [25,26]. We have also calculated waveforms of all published Broucke-Hadjidemetriou-Henon (BHH) orbits [14–20], which we omit from this Letter for the sake of brevity, and because they are closely related to Henon's "criss-cross" one, studied in Ref. [10]. The waves of the $13+11$ new orbits show clear distinctions in form and luminosity, thus ensuring that they would be distinguishable (provided their signals are strong enough to be detected).

We consider systems of three equal massive particles moving periodically in a plane under the influence of Newtonian gravity. The quadrupole moment I_{ij} of three bodies with equal masses $m_n = m$, $(n = 1, 2, 3)$ is expressed as $I_{ij} = \sum_{n=1}^{3} m x_n^i x_n^j$, where x_n^i is the location of nth body, and the spatial dimension indices i and j run from 1 to 3 (with $x^1 = x$, $x^2 = y$, $x^3 = z$). The reduced quadrupole Q_{ij} is defined as $Q_{ij} = I_{ij} - \frac{1}{3}\delta_{ij}\sum_{k=1}^{3} I_{kk}$. The gravitational waveforms denoted by h_{ij}^{TT} are, asymptotically,

$$h_{ij}^{TT} = \frac{2G}{rc^4}\frac{d^2 Q_{ij}}{dt^2} + \mathcal{O}\left(\frac{1}{r^2}\right), \tag{1}$$

where r is the distance from the source, Refs. [8,9]. Here, TT means (i) transverse ($\sum_{i=1}^{3} h_{ij}^{TT}\hat{n}^i = 0$) and (ii) traceless ($\sum_{i=1}^{3} h_{ii}^{TT} = 0$), where $\hat{n}_i$ denotes the unit vector of the gravitational wave's direction of propagation. The two

TABLE I. Initial conditions and periods of three-body orbits. $\dot{x}_1(0)$, $\dot{y}_1(0)$ are the first particle's initial velocities in the x and y directions, respectively, T is the period of the (rescaled) orbit to normalized energy $E=-1/2$, Θ is the rotation angle (in radians) and $\langle P\rangle$ is the mean luminosity (power) of the waves emitted during one period. Other two particles' initial conditions are specified by these two parameters, as follows: $x_1(0)=-x_2(0)=-\lambda$, $x_3(0)=0$, $y_1(0)=y_2(0)=y_3(0)=0$, $\dot{x}_2(0)=\dot{x}_1(0)$, $\dot{x}_3(0)=-2\dot{x}_1(0)$, $\dot{y}_2(0)=\dot{y}_1(0)$, $\dot{y}_3(0)=-2\dot{y}_1(0)$. The Newtonian coupling constant G is taken as $G=1$ and the masses are equal $m_{1,2,3}=1$.

Name	$\dot{x}_1(0)$	$\dot{y}_1(0)$	λ	T	Θ(rad)	$\langle P\rangle$
Moore's figure eight	0.216 343	0.332 029	2.574 29	26.128	0.245 57	1.35×10^{0}
Simo's figure eight	0.211 139	0.333 568	2.583 87	26.127	0.277 32	1.36×10^{0}
(M8)[7]	0.147 262	0.297 709	3.008 60	182.873	0.269 21	2.46×10^{0}
I.A.1 butterfly I	0.147 307	0.060 243	4.340 39	56.378	0.034 78	1.35×10^{5}
I.A.2 butterfly II	0.196 076	0.048 690	4.016 39	56.375	0.066 21	5.52×10^{6}
I.A.3 bumblebee	0.111 581	0.355 545	2.727 51	286.192	−1.090 4	1.01×10^{5}
I.B.1 moth I	0.279 332	0.238 203	2.764 56	68.464	0.899 49	5.25×10^{2}
I.B.2 moth II	0.271 747	0.280 288	2.611 72	121.006	1.138 78	1.87×10^{3}
I.B.3 butterfly III	0.211 210	0.119 761	3.693 54	98.435	0.170 35	3.53×10^{5}
I.B.4 moth III	0.212 259	0.208 893	3.263 41	152.330	0.503 01	7.48×10^{5}
I.B.5 goggles	0.037 785	0.058 010	4.860 23	112.129	−0.406 17	1.33×10^{4}
I.B.6 butterfly IV	0.170 296	0.038 591	4.226 76	690.632	0.038 484	1.23×10^{13}
I.B.7 dragonfly	0.047 479	0.346 935	2.880 67	104.005	−0.406 199	1.25×10^{6}
II.B.1 yarn	0.361 396	0.225 728	2.393 07	205.469	−1.015 61	2.33×10^{6}
II.C.2a yin-yang I	0.304 003	0.180 257	2.858 02	83.727	0.659 242	1.31×10^{5}
II.C.2b yin-yang I	0.143 554	0.166 156	3.878 10	83.727	−0.020 338	1.31×10^{5}
II.C.3a yin-yang II	0.229 355	0.181 764	3.302 84	334.877	0.472 891	7.19×10^{10}
II.C.3b yin-yang II	0.227 451	0.170 639	3.366 76	334.872	0.254 995	7.19×10^{10}

independent waveforms $h_{+,\times}$ of a quadrupolar gravitational wave propagating along the z axis, Refs. [8,9] can be expressed as

$$h_{+}=\frac{2G}{c^4 r}\sum_{i=1}^{3} m_i(\dot{x}_i^2+x_i\ddot{x}_i-\dot{y}_i{}^2-y_i\ddot{y}_i), \tag{2}$$

$$h_{\times}=\frac{2G}{c^4 r}\sum_{i=1}^{3} m_i(\ddot{x}_i y_i+2\dot{x}_i\,\dot{y}_i+x_i\ddot{y}_i), \tag{3}$$

where r denotes the distance from the source to the observer. We set the units of $G=c=m=1$ throughout this Letter.

Here the coordinate axes x and y are chosen so that they coincide with the orbits' two (reflection) symmetry axes, when they exist, i.e., when the orbits are from class I, as defined in Ref. [25]. Otherwise, e.g., when only a single point reflection symmetry exists, as in class II orbits, the x, y axes are taken to be the eigenvectors of the moment-of-inertia tensor. The rotation angle necessary for each orbit to be aligned with these two axes is given in Table I [27].

The first gravitational radiation waveforms for periodic three-body systems were studied in Refs. [10–12]. They calculated the quadrupole radiation waveforms for three periodic orbits of the following three-equal-mass systems: (i) of the Lagrange "equilateral triangle" orbit [13], (ii) of Henon's "criss cross" [19], and (iii) of Moore's "figure eight" [21]. These three orbits are characteristic representatives of the (only) three families of periodic three-body orbits known at the time. Reference [10] found distinct gravitational waveforms for each of the three families, thus suggesting that one might be able to distinguish between different three-body systems as sources of gravity waves by looking at their waveforms [28].

In the meantime 13 + 11 new orbits belonging to 12 new families have been found, Refs. [25,26]. The families of three-body orbits can be characterized by their topological properties viz. the conjugacy classes of the fundamental group, in this case, the free group on two letters (a, b), Ref. [29]. The free group element tells us the number of times the system's trajectory on the shape sphere passes around one or another (prechosen) two-body collision point within one period. Every time the system is close to a two-body collision the (relative) velocities, accelerations, and the third derivatives of relative coordinates increase, so that the luminosity of gravitational radiation also increases; i.e., there is a burst of gravitational radiation. This argument can be made more quantitative by appealing to two-body results of Ref. [8], as is shown in footnote [30].

We show the gravitational radiation waveforms $h_{+,\times}$ in Fig. 1, emitted by three massive bodies moving according to the orbits from Refs. [25,26] belonging to these families, where Eqs. (2) and (3) are used as the definitions of the two waveforms.

First, we note that all of the calculated three-body orbits' waveforms are distinct [31], thus answering (in the

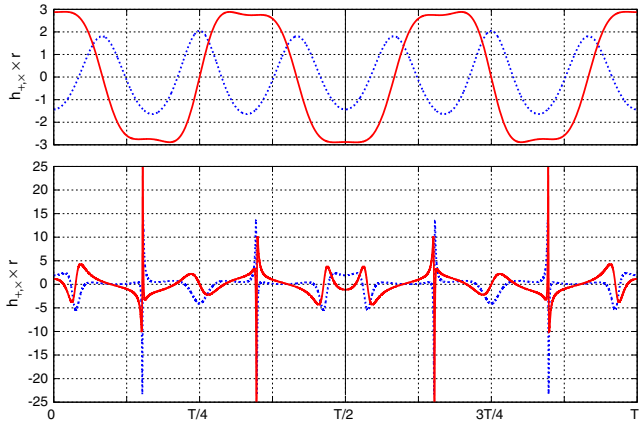

FIG. 1 (color online). The gravitational radiation quadrupolar waveforms $h_{+,\times} \times r$ as functions of the elapsed time t in units of the period T, for two periodic three-body orbits (in units of Gm/c^2; we have set $G = m = c = 1$ throughout this Letter) and r is the radial distance from the source to the observer. Dotted (blue) and solid (red) curves denote the $+$ and $\times$ modes, respectively. Top: Simo's figure eight, Ref. [22]; and bottom: orbit I.B.1 Moth I. Note the symmetry of these two graphs under the (time-)reflection about the orbits' midpoint $T/2$ during one period T.

positive) the question about their distinguishability posed in Ref. [10]. In Fig. 1 we also show the gravitational waveform of one "old" orbit: Simo's figure eight, (discovered in 2002) belonging to the figure-eight family. Simo's figure eight is an important example, as it is virtually indistinguishable from Moore's one, and yet the two have distinct gravitational waveforms, see our Fig. 1, and Fig. 2 in Ref. [10]. That is so because these two figure-eight solutions have distinct time dependences of the hyperradius R, where $R^2 \sim (1/m)\delta_{ij}\sum_{k=1}^{3} I_{kk}$, so that the two orbits have different quadrupolar waveforms.

Note, moreover, the symmetry of the waveforms in Fig. 1 with respect to reflections of time about the midpoint of the period $T/2$: this is a consequence of the special subset of initial conditions (vanishing angular momentum and passage through the Euler point on the shape sphere) that we used. There are periodic three-body orbits, such as those from the BHH family, that do not have this symmetry.

The gravitational waveforms' maxima range from 20 to 50 000 in our units, with the energy fixed at $E = -1/2$. This large range of maximal amplitudes is due to the differences in the proximity of the approach to two-body collisions in the corresponding orbits. One can explicitly check that the bursts of gravitational radiation during one period correspond to close two-body approaches.

As stated above, the (negative) mean power loss $\langle dE/dt \rangle$ of the three-body system, or the (positive) mean luminosity (emitted power) of quadrupolar gravitational radiation $\langle P \rangle$, averaged over one period, is proportional to the square of the third time derivative of the (reduced) quadrupole moment $Q^{(3)}_{jk}$, $\langle dE/dt \rangle = -\langle P \rangle = -\frac{1}{5}(G/c^5)\sum_{j,k=1}^{3}\langle Q^{(5)}_{jk}\dot{Q}_{jk}\rangle = -\frac{1}{5}(G/c^5)\sum_{j,k=1}^{3}\langle Q^{(3)}_{jk}Q^{(3)}_{jk}\rangle$, (for an original derivation see Refs. [8,9], for pedagogical ones, see Refs. [1,2]). But,

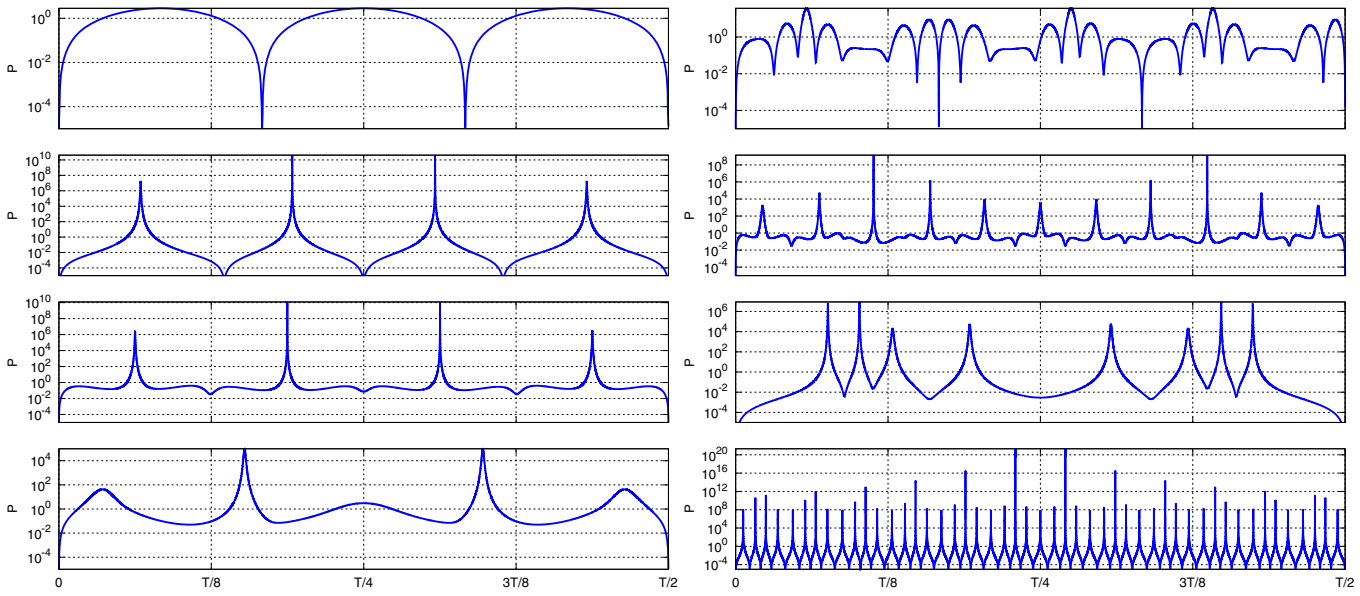

FIG. 2 (color online). The instantaneous (time unaveraged) luminosity P of quadrupolar gravitational radiation emitted from periodic three-body orbits as a function of the elapsed time t in units of the period T. Note the logarithmic scale for the luminosity P (y axis). Top left: Moore's figure eight; second from top left: I.A.2 butterfly II; third from top left: II.B.7 dragonfly; bottom left: I.B.1 moth I; top right: $(M8)^7$, second from top right: I.A.3 bumblebee; second from bottom right: I.B.5 goggles; bottom right: II.B.6 butterfly IV.

$Q^{(3)}_{jk}$ are proportional to the first time derivatives of the gravitational waveforms $Q^{(3)}_{jk} = (d/dt)Q^{(2)}_{jk} \propto (d/dt)h_{+,\times}$. The peak amplitudes of gravitational waveforms $h_{+,\times}$, in turn, grow in the vicinity of two-body collisions [30], which explains the burst of gravitational radiation as one approaches a two-body collision point.

The mean and instantaneous luminosities, expressed in our units, of these orbits, normalized to $E = -1/2$, are shown in Table I and Fig. 2, respectively. Note that in Table I we show only three orbits belonging to the figure-eight family: Moore's, Simo's, and the stable choreography $(M8)^7$; they have all the same order of magnitude of the mean luminosity [32], whereas the butterfly I and butterfly II orbits, which belong to the same topological family, have mean luminosities that differ by more than a factor of 40.

Generally, the mean luminosities of these orbits cover 13 orders of magnitude, ranging from 1.35 (Moore's figure eight) to 1.23×10^{13} (I.B.6 butterfly IV) in our units; see Table I. The peak instantaneous luminosities have an even larger range: 20 orders of magnitude; see Fig. 2. Here, the symmetric form of the instantaneous (time unaveraged) power $P = \frac{1}{5}(G/c^5)\sum_{j,k=1}^{3} Q^{(3)}_{jk} Q^{(3)}_{jk}$ was used. This gives us hope that at least some of these three-body periodic orbits can, perhaps, lead to detectable gravitational radiation signals.

It is a different question if some or all of these sources of gravitational radiation would be observable by the present-day and the soon-to-be-built gravitational wave detectors: that strongly depends on the absolute values of the masses, velocities, and the average distances between the three celestial bodies involved, as well as on the distribution of such sources in our Galaxy.

Moreover, note that all of the newly found and analyzed three-body orbits have zero angular momentum, and many of them are unstable. It is well known [16–20] that by changing the angular momentum within the same family of three-body orbits, the stability of an orbit changes as well. So, it may happen that a previously stable orbit turns into an unstable one, and vice versa. For this reason it should be clear that a careful study of gravitational-radiation-induced energy- and angular-momentum dissipation is necessary for these orbits [33]. Moreover, if realistic results are to be obtained, post-Newtonian approximations will have to be applied in the future. Such relativistic corrections are most important at large velocities, i.e., precisely near close approaches that are so crucial for large gravitational radiation. Thus, the present Letter is meant only to highlight the possibilities in this field, and should be viewed as an invitation to join in the more realistic future studies.

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[31] We do not show these waveforms here, except for the two in Fig. 1, for brevity's sake, and because many are fairly similar to the second waveform in Fig. 1—regular sequences of spikes.
[32] Note that the figure-eight family members have, on the average, the lowest luminosity among the orbits considered here.
[33] We plan to do such a study, which cannot be completed, however, without an extension of each orbit to a family of orbits with nonvanishing angular momenta. So far, only the BHH family has been extended in such a way, but even that one case is not complete [34].
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Linear stability of periodic three-body orbits with zero angular momentum and topological dependence of Kepler's third law: a numerical test

V Dmitrašinović[1,5], Ana Hudomal[2], Mitsuru Shibayama[3] and Ayumu Sugita[4]

[1] Institute of Physics Belgrade, University of Belgrade, Pregrevica 118, Zemun, PO Box 57, 11080 Beograd, Serbia
[2] Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia
[3] Department of Applied Mathematics and Physics, Graduate School of Informatics, Kyoto University, Yoshida-honmachi, Sakyo-ku, Kyoto 606-8501, Japan
[4] Department of Applied Physics, Osaka City University, 3-3-138 Sugimoto, Sumiyoshi-ku, Osaka 558-8585, Japan

E-mail: dmitrasin@ipb.ac.rs

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Abstract

We test numerically the recently proposed linear relationship between the scale-invariant period $T_{\text{s.i.}} = T|E|^{3/2}$, and the topology of an orbit, on several hundred planar Newtonian periodic three-body orbits. Here T is the period of an orbit, E is its energy, so that $T_{\text{s.i.}}$ is the scale-invariant period, or, equivalently, the period at unit energy $|E| = 1$. All of these orbits have vanishing angular momentum and pass through a linear, equidistant configuration at least once. Such orbits are classified in ten algebraically well-defined sequences. Orbits in each sequence follow an approximate linear dependence of $T_{\text{s.i.}}$, albeit with slightly different slopes and intercepts. The orbit with the shortest period in its sequence is called the 'progenitor': six distinct orbits are the progenitors of these ten sequences. We have studied linear stability of these orbits, with the result that 21 orbits are linearly stable, which includes all of the progenitors. This is consistent with the Birkhoff–Lewis theorem, which implies existence of infinitely many periodic orbits for each stable progenitor, and in this way explains the existence and ensures infinite extension of each sequence.

[5] Author to whom any correspondence should be addressed.

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S Supplementary material for this article is available online

(Some figures may appear in colour only in the online journal)

1. Introduction

There is no general solution to the Newtonian three-body problem [1], so particular solutions, such as periodic orbits, are of special interest. Up until five years ago, only three topologically distinct families of periodic orbits were known [2–5], with the latest two discoveries being received with some fanfare. No theorem guaranteeing the existence of further periodic solutions was known at the time. Indeed contradictory claims [6], and counterclaims [7] in the 1950s and 1960s led to some confusion, which was (only partially) resolved by subsequent numerical discoveries—the corresponding formal existence proofs for these orbits are still being sought, and only in a few rare examples, have been supplied—for a brief history of this problem up to mid 1970's see section 1[6] in Broucke [8], and for subsequent developments, see section I in [9].

The questions of existence, density and distribution of *stable* orbits is of some importance for astronomy: stable orbits have at least a fighting chance of being produced in astrophysical processes and, therefore, of being subsequently observed. These questions can only be addressed by explicit discovery, or construction of new stable orbits[7]. Therefore any reliable new source of information about periodic orbits, even if it is (only) empirical and incomplete, ought to be welcomed by the community and subjected to further tests.

Several hundred demonstrably distinct families of periodic orbits have been found by numerical means over the past few years [9–14]. This progress in numerical studies has led to a new, wholly unexpected insight into the distribution of periodic orbits, that was, at first, rather tentative: soon after the papers [5, 10] appeared a relationship between an orbit's period and its topology was observed—at first just in one class of orbits [10], and then more generally [15]. The initial set of orbits was fairly 'sparse', consisting of only about 45 orbits, so the observed regularities had large gulfs yet to be filled. In the meantime we have continued our search for new orbits, as well as tests of their stability, amounting to more than 200 orbits, this time with a clear indication that their number grows without bounds as the scale-invariant period increases, and still following the linear dependence of an orbit's period on its topology [9].

Here we present a new, detailed numerical test of the previously observed regularities, based on more than 200 orbits, as well as several new regularities regarding (probably) infinite sequences of orbits. Moreover, we present a semi-empirical observation about the relation between stability of certain orbits and the existence of infinite sets of periodic orbits, as

[6] 'However, the existence of periodic solutions for the general three-body problem has been considered a somewhat controversial question in the last few years. Vernić (1953) has published a detailed study containing a mathematical proof of the non-existence of periodic solutions other than the Lagrange solutions. Later it is seen that Merman (1956) and Leimanis (1958) have questioned Vernić's non-existence proof. More recently Arenstorf (1967) has published a new existence proof for periodic solutions of the general problem, although his work contains no examples, whereas Kolenkiewicz and Carpenter have numerically computed a periodic solution with masses and configuration of the Sun–Earth–Moon system. Jefferys and Moser (1966) have also published existence proofs for almost periodic solutions in the three-dimensional case. However, the most convincing explicit examples of periodic solutions have recently been obtained numerically by Szebehely and Standish (1969), and Peters (1967). Their publications definitely settle the question of whether the general problem has non-trivial periodic solutions, although all of their examples are rather specialized; i.e. collision orbits or zero total angular momentum orbits'.

[7] Only roughly one out of ten of the newly discovered orbits are linearly stable [9, 14].

related by the Birkhoff–Lewis theorem [16], as well as some analytic arguments about the causes of the linear relation between the period and topology, that still remain without rigorous proofs. These arguments have evolved from the study [44] of the three-body system in the so-called strong Jacobi–Poincaré potential, which system is simpler than the Newtonian one, and therefore allows certain theorems about the existence of solutions to be proven and analytical arguments to be made. The extension of these analytic arguments to the Newtonian three-body system may seem straightforward at first, but a closer inspection might prove more complicated. We have tried and pointed out lacunae in our arguments, in the hope that experts will either complete the proofs, or definitely disprove the conjectures.

If our numerical and empirical arguments withstand a more rigorous mathematical scrutiny, they should have: (1) significant implications for the distribution of periodic three-body orbits in all homogeneous potentials with singularities at the two-body collision points: at least one such potential (the Coulomb one) is of direct physical interest; and (2) ready generalizations for 4-, 5-, ... n-body periodic orbits in the Newtonian potential.

In this paper, after the present Introduction, in section 2 we provide the necessary preliminaries for our work. Then in section 3 we provide more than 200 periodic zero-angular-momentum orbits and identify their topologies using two integers, n_w and $\bar{n}_w$, defined in section 2. There we test their $T_{\text{s.i.}}$ versus $(n_w + \bar{n}_w)$ relationship(s) and refine the quasi-linear rule, equation (2), by classifying the new orbits into ten algebraically well-defined sequences. In section 4 we study the linear stability of three-body orbits, which leads us to the identification of six orbits as progenitors of ten sequences of orbits. There, we offer a possible explanation for the existence of infinitely many orbits in each sequence, in terms of the Birkhoff–Lewis theorem, which we do not prove in this case, however. In section 5 we offer a possible explanation of the observed linear regularities, using the virial theorem and the analyticity of the action. Finally, in section 6 we summarize and discuss our results, as well as present some open questions. Appendices A–E are devoted to various necessary technical topics, that would distract the flow of our arguments, if they were kept in the main text.

2. Preliminaries: topology and period of periodic three-body orbits

For a quantitative relationship between topology and period to be possible one has to have an algebraic method for the description of an orbit's topology. There are several such methods in the literature, variously based on the braid group B_2, [2], on the free group F_2 on two elements [17], and on three symbols [18], see appendices B and C.

The original discovery of the linear relationship between period and topology was based on Montgomery's free group method [17], which was used to identify and label periodic orbits.

The topology of a periodic three-body orbit O can be algebraically described by a finite sequence of symbols, e.g. letters $(\mathtt{a}, \mathtt{b})$ and $(\mathtt{A}, \mathtt{B})$, that we shall call 'word' w_{O}[8], as defined in [17], and presented in detail in [19], and briefly reviewed in appendix B. For an alternative method of assigning symbols to a topology, see appendix C.

With such an algebraic description one could, for the first time, search for relations between topological and dynamical properties of orbits. At first, the curious approximate linear functional relation

$$\frac{T_{\text{s.i.}}(w_8^k)}{T_{\text{s.i.}}(w_8)} \equiv \frac{T(w_8^k)|E(w_8^k)|^{3/2}}{T(w_8)|E(w_8)|^{3/2}} \simeq k = 1, 2, 3, \ldots \quad , \tag{1}$$

[8] More precisely, the conjugacy class of the free group element.

was noticed between the periods T, energies E and the free-group elements $w_8 = (\mathtt{ab})(\mathtt{AB})$ for the figure-eight orbit [3] and their topological-power satellite orbits with topologies $w^k = [(\mathtt{ab})(\mathtt{AB})]^k$, $(k = 1, 2, 3, \cdots)$. We define 'topological-power satellite' orbits as those whose topologies can be described as k times repeated topology, i.e. integer powers w^k of the simplest ('progenitor') orbit described by the word w [10]. Here $\simeq$ means equality within the estimated numerical precision of [10]. In the meantime, with improved numerics, several cases have been found where this relation breaks down at higher decimal places.

Initially, only the 'topological-power satellites' of the figure-eight orbit were known[9], but, in the meantime new examples of topological-power satellites[10] have been found to obey equation (1) within their respective numerical errors. This naturally raises the question: why do only some orbits have topological-power satellites and not others? We shall argue below that the linear stability of the shortest-period ('progenitor') orbit plays a crucial role in this regard.

Following this observation, [15] investigated all of the 45 orbits known at the time and not just the topological-power satellites, and observed the following more general[11] quasi-linear relation

$$\frac{T_{\text{s.i.}}(w)}{T_{\text{s.i.}}(w_p)} \simeq \frac{N_w}{N_{w_p}} = \frac{n_w + \bar{n}_w}{n_{w_p} + \bar{n}_{w_p}} \tag{2}$$

for three-body orbits with zero angular momentum. Here $N_w = n_w + \bar{n}_w$ is one half of the minimal total number of letters[12], in the free group element $w = w(O)$ characterizing the (family of) orbit O, and similarly for $w_p = w(\text{progenitor})$, the word describing the progenitor orbit in a sequence, where n_w is the number $n_w = \frac{1}{2}(n_{\mathtt{a}} + n_{\mathtt{b}})$, of small letters $\mathtt{a}$, or $\mathtt{b}$, and $\bar{n}_w = \frac{1}{2}(n_{\mathtt{A}} + n_{\mathtt{B}})$ is the number of capital letters $\mathtt{A}$, or $\mathtt{B}$.

Equation (2) suggested 'at least four and at most six' distinct sequences among the 45 orbits considered in [15]. Precise algebraic definitions of these sequences, analogous to the definition w^k of the topological-power satellites, were not known at the time, again due to the dearth of distinct orbits[13]. This clearly demanded further, finer searches to be made.

Equation (2) predicts (infinitely) many new, as yet unobserved orbits together with their periods; if true, even approximately, equation (2) would be a spectacular new and unexpected property of three-body orbits, that would open new insights into the Newtonian three-body problem, as well as provide help in practical searches to find new orbits. Therefore equation (2) merits a thorough investigation, which we shall attempt below. The scope, of course, is limited by the number and type of available orbits.

3. Classification of orbits in sequences

Using equation (2) we predicted the periods and numbers of letters of new orbits, and then searched for them, with the results first reported in [9]. We did so by first identifying the linearly stable orbits among the original 13 orbits, and then by 'zooming in' our search on smaller windows around the stable orbits. Thus we found new periodic orbits that have 'filled'

[9] With one exception: the yarn orbit $w_{\text{yarn}} = (\mathtt{babABabaBA})^3 = w^3_{\text{moth I}}$, where $w_{\text{moth I}} = \mathtt{babABabaBA}$ in [5].
[10] E.g. of the 'moth I' orbit, as well as several topological-power satellites of three other orbits, see [9, 20, 39].
[11] Equation (1) is manifestly a special case of equation (2).
[12] Here, by 'minimal total number of letters' we mean the number of letters after all pairs of adjacent identical small and capital letters, such as $\mathtt{aA}$, have been eliminated, as explained in [9].
[13] Many distinct satellite orbits' points almost overlapped on the $T_{\text{s.i.}} - N_w$ graph, due to identical values of N_w and similar periods, which further reduced the number of distinct data points. Moreover, there were significant 'gaps' between the data points, as well as one 'outlier point' (orbit), in figure 1 in [15], that was roughly 8% off the conjectured straight line.

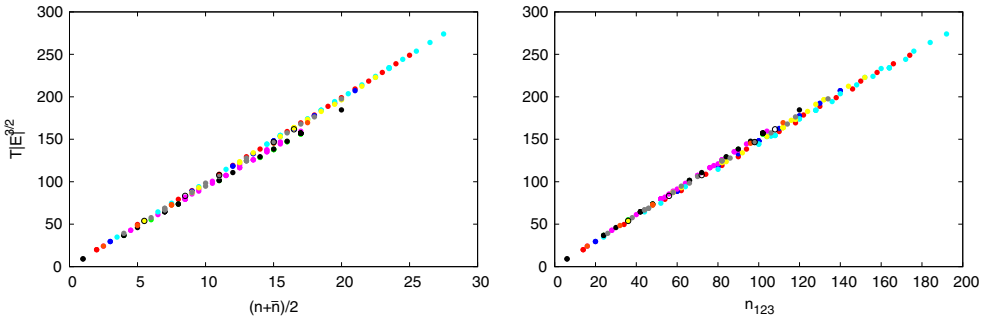

Figure 1. (a) Left panel: the scale-invariant periods $|E|^{3/2}T(w)$ of more than 200 presently known zero-angular-momentum three-body orbits versus one half of the number of all letters in the free-group word w describing the orbit, $N_w = n_w + \bar{n}_w$, where n_w is the number of small letters `a`, or `b`, and $\bar{n}_w$ is the number of capital letters `A`, or `B` in the word w. (b) Right panel: same as (a), only in terms of the number of symbols n_{123} in the sequence of symbols (1,2,3) describing the topology of the orbit, see appendix C. Color code: (1) red = sequence I—butterfly I; (2) green = sequence II—dragonfly; (3) dark blue = sequence III—yin-yang; (4) pink = sequence IVa—moth I; (5) light blue = sequence IVb—butterfly III; (6) yellow = sequence IVc—moth III; (7) black = sequence V—figure-eight; (8) orange = sequence VI—yarn; (9) grey = sequence VII—moth; (10) empty circles = other.

Table 1. Typical (non-minimal) free group elements' w structure for orbits in various sequences, their progenitors, the line parameters c_1, c_2, where the $T_{\text{s.i.}}(N_w)$ dependence is fitted as $f(x) = c_1 x + c_2$. Not all words $w(n_i)$ in any particular sequence need have the presented structure, however, see supplementary notes.

Sequence number and name	Free group element$w(n)$	progenitor	c_1	c_2
I butterfly I (n,n)	$(\mathtt{AB})^2(\mathtt{abaBAB})^n(\mathtt{ab})^2(\mathtt{ABAbab})^n$	Schubart	9.957 ± 0.011	-0.2 ± 0.2
II dragonfly (n,n)	$\mathtt{bA}(\mathtt{baBA})^n\mathtt{aB}(\mathtt{abAB})^n$	isosceles	9.194 ± 0.004	0.04 ± 0.06
III yin-yang (n,n)	$(\mathtt{abaBAB})^n\mathtt{a}(\mathtt{babABA})^n\mathtt{A}$	S-orbit	9.8667 ± 0.0003	0.002 ± 0.004
IVa moth I $(n,n+1)$	$(\mathtt{abAB})^n\mathtt{A}(\mathtt{baBA})^n\mathtt{B}$	moth I	9.34 ± 0.06	0.7 ± 0.7
IVb butterfly III $(n,n+1)$	$[(\mathtt{ab})^2(\mathtt{AB})^2]^n\mathtt{b}[(\mathtt{ba})^2(\mathtt{BA})^2]^n\mathtt{a}$	butterfly III	9.967 ± 0.012	-0.3 ± 0.3
IVc moth III $(n,n+1)$	$(\mathtt{babABA})^n\mathtt{A}(\mathtt{abaBAB})^n\mathtt{B}$	Schubart	9.94 ± 0.04	-1.2 ± 0.7
V figure-eight (n,n)	$(\mathtt{abAB})^n$	figure-8	9.2377 ± 0.0014	-0.03 ± 0.02
VI moth I—yarn $(2n,3n)$	$[(\mathtt{abAB})\mathtt{A}(\mathtt{baBA})\mathtt{B}]^n$	moth I	9.683 ± 0.002	0.01 ± 0.02
VIIa moth (n,n)	$(\mathtt{abAB})^{(n+1)}\mathtt{a}(\mathtt{baBA})^n\mathtt{b}$	Schubart	9.61 ± 0.07	-0.2 ± 0.7
VIIb moth (n,n)	$(\mathtt{abaBAB})^{(n+1)}\mathtt{b}(\mathtt{babABA})^n\mathtt{a}$	Schubart	9.88 ± 0.04	-0.7 ± 0.5

many of the 'gaps' in the older versions of the $T_{\text{s.i.}} - N_w$ graph, see figure 1(a), the website [20] and the supplementary notes (stacks.iop.org/JPhysA/51/315101/mmedia). The 'outlier' point, in figure 1 in [15], has become just another orbit in a new sequence with a slightly steeper slope on the same graph. The totality of the $T_{\text{s.i.}} - N_w$ points is shown in figure 1.

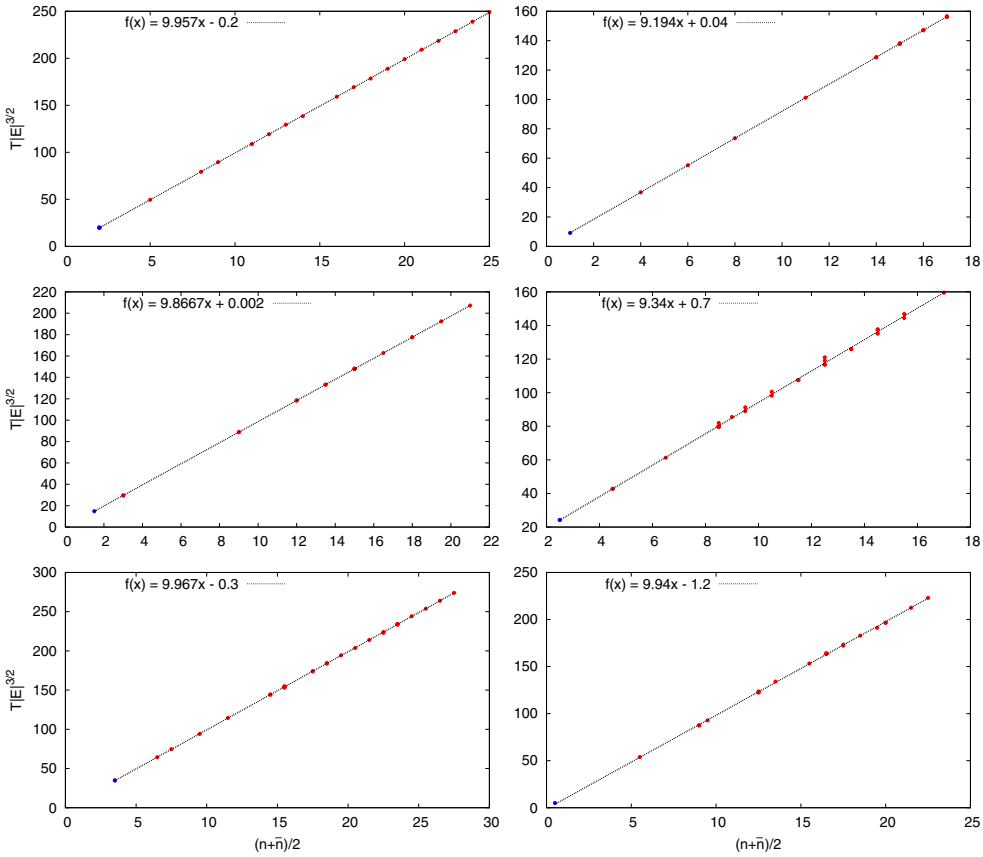

Figure 2. The scale-invariant periods $|E|^{3/2}T(w)$ of zero-angular-momentum three-body orbits versus one half of the number of all letters in the free-group word w describing the orbit, $N_w = n_w + \bar{n}_w$, where n_w is defined as in figure 1. (a) Top left: sequence I—butterfly I, ; (b) top right: sequence II—dragonfly; (c) center left: sequence III—yin-yang; (d) center right: sequence IVa—moth I; (e) bottom left: sequence IVb—butterfly III; (f) bottom right: sequence IVc—moth III. The blue points at the lower ends of sequences are the progenitors of the respective sequences, see the text. Progenitors of sequences II, III and IVc, that involve collisions were not used in the fitting procedure, so the validity of the linear Ansatz for these sequences can be evaluated by inspection.

It is clear that the scale-invariant periods $T_{\text{s.i.}}$ do not lie on one straight line, but rather on several lines with slightly different slopes, emerging from a small 'vertex' area, forming a (thin) wedge-like structure in figure 1. All the newly found orbits passing through an Euler configuration, see supplementary notes, fit into one of ten sequences, where the fourth ('moth I') sequence in [15] has now been divided into three: (a) 'moth I $(n, n+1)$'; (b) 'butterfly III–IV $(n, n+1)$'; (c) 'moth III $(n, n+1)$'. Moreover, we found two entirely new sequences: (1) 'VIIa moth (n, n)' and (2) 'VIIb moth (n, n)', and one sequence of pure 'topological-power satellites' of the moth I orbit.

Each of these ten sequences has an algebraic pattern of free-group elements, see table 1, associated with it. Here we use the sequence label (n, m) to denote the general form of $(n_w, \bar{n}_w)$ in that sequence: for example (n, n) means that n_w and $\bar{n}_w$ are equal integers: $n = n_w = \bar{n}_w = 1, 2, 3, \ldots$. Then, n can be used to label orbits within the sequence, see supplementary notes. By setting $n = 0$, or $n = 1$, in the second column of table 1, in each sequence,

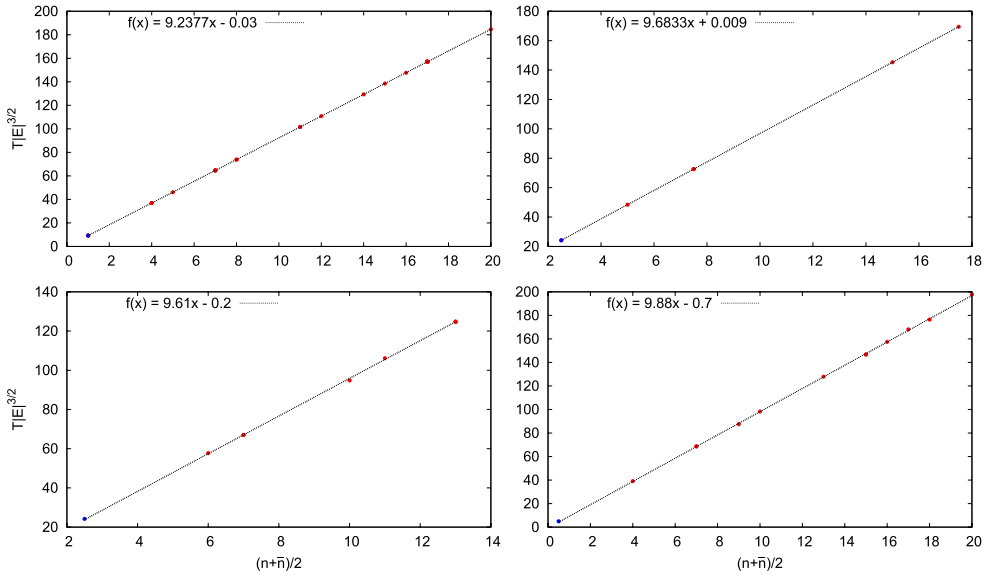

Figure 3. Same as in figure 2, except for the following sequences: (a) top left: sequence V—figure-eight; (b) top right: sequence VI—yarn; (c) bottom left: sequence VIIa—moth III (n, n); (d) bottom right: sequence VIIb—moth III (n, n). The progenitors of sequence VIIa and VIIb were not used in the fitting procedure.

we can read off the topology of their respective progenitor, which is shown in the third column of table 1.

The individual $T_{\text{s.i.}} - N_w$ graphs are shown in figures 2 and 3, and their free-group patterns are in table 1. The agreement of separate sequences with the linear functional Ansatz, equation (2), see figures 1(b)–(d), is much better than for the aggregate of all orbits, Figure 1, but the (root-mean-square) variations of line parameters (c_1, c_2) reported in table 1 are generally larger than the estimates numerical errors, thus indicating that equation (2) is still *approximate, and not exact*, even in these sequences.

Whereas the approximate empirical rule equation (2) now appears established, and its extension to ever-longer periods just a technical difficulty, some deeper questions remain open. For example, the raison d'être of so many periodic orbits remains obscure, let alone the linear relation among their periods.

4. Linear stability and progenitor orbits

Perhaps the first hint at a solution to this puzzle was given in [39], where it was noticed that the topological satellite orbits in the Broucke–Hadjidemetriou–Hénon (BHH), [8, 35, 36, 40–43], family of orbits with non-zero angular momentum, exist only when their progenitor is linearly stable. There is a theorem, due to Birkhoff and Lewis [16], see also section 3.3 (by Jürgen Moser) in [25], which holds for systems with three degrees-of-freedom and implies the existence of infinitely many periodic orbits[14]. So, whereas the Birkhoff–Lewis theorem might solve

[14] In [12], it was conjectured that the topological-power satellites of the figure-eight orbit are a consequence of the Poincaré–Birkhoff theorem [22], see also section 24 in [23] and section 2.7 in [24], as applied to the figure-eight orbit. That conjecture is incorrect, however, because the Poincaré–Birkhoff theorem applies only to systems with two degrees-of-freedom, to which class the planar three-body problem does not belong.

Table 2. The Floquet exponents ν_j, where $\lambda_j = \exp(\pm 2\pi \mathrm{i}\nu_j)$ define the linear stability coefficients of linearly stable periodic three-body orbits.

Label	ν_1	ν_2
S-orbit	0.131 093	0.470 591
Moore 8	0.298 093	0.00 842 275
NC1 (8^7)	0.27 216	0.158 544
V.17.H (O13 $= 8^{17}$)	0.31 573	0.0002 988
V.17.I (O14 $= 8^{17}$)	0.0435 411	0.00 262 681
V.17.J (O15 $= 8^{17}$)	0.0435 411	0.00 262 681
II.11.A (bumblebee)	0.137 149	0.0325 135
IVa.2.A (moth I)	0.159 013	0.491 881
IVa.4.A (moth II)	0.108 451	0.0886 311
IVb.3.A (butterfly III)	0.378 728	0.00 173 642
I.5.A	0.170 764	0.001 476
I.14.A	0.443 006	0.000 121 435
II.17.B	0.138 698	0.0335 924
III.13.A.β	0.175 816	0.000 655 417
IVb.9.A	0.194 186	0.000 561 819
IVc.12.B	0.0863 933	0.00 394 124
IVc.17.A	0.0442 047	0.00 206 416
VIIa.11.A	0.416 228	0.0088 735
VIIb.7.A	0.27 753	0.0360 425
VIIb.9.A	0.216 455	0.0584 561
VIIb.13.A	0.0621 421	0.0141 894

one part of the puzzle, it does not say anything about the relation of topologies and periods. There is, however, another (the so-called 'twist') condition underlying this theorem, which we shall not try to check here—we simply conjecture that the Birkhoff–Lewis theorem holds for the linearly stable periodic three-body orbits. Linear stability of periodic orbits is tested numerically, see below, and thus the conjecture of Birkhoff–Lewis theorem can be falsifed.

We have analyzed linear stability of all zero-angular-momentum three-body orbits and tabulated the linearly stable ones in table 2. The Floquet exponents ν_j, and the linear stability coefficients $\lambda_j = \exp(\pm 2\pi \mathrm{i}\nu_j)$, are the standard ones, as defined in [9]. We note that two orbits, 'butterfly III' and 'moth I', lie at the origins of two 'linear sequences[15]' of 'non-topological-power satellite' orbits observed among the original 13 orbits [15].

Thus, the manifest candidates for progenitors are: (1) 'figure-eight' for the sequence V 'figure-eight (n, n)'; (2) 'butterfly III' for the sequence IVb 'butterfly III $(n, n+1)$'; and (3) 'moth I' for the sequences IVa 'moth I $(n, n+1)$' and VI 'moth I—yarn $(2n, 3n)$'. These three progenitors are collisionless orbits with three degrees-of-freedom, that are linearly stable.

Next we extend this reasoning to sequences of periodic three-body orbits with collisional progenitors.

(1) The parent orbit of sequence II 'dragonfly (n, n)' is Broucke's isosceles triangle orbit [37, 38], that involves two-body collisions. This orbit always stays in an isosceles triangle configuration, thus eliminating one degree-of-freedom, and is linearly stable [37, 38], so it also satisfies the Poincaré–Birkhoff theorem.
(2) The parent orbit of the 'yin-yang' sequence is the collisional 'S-orbit' of [4, 11][16].

[15] The orbits 'moth I' and 'moth II' have different topologies, but belong to the same sequence.
[16] See the initial condition #20 in table I in [11].

(3) The Schubart orbit [34] is the progenitor of four sequences: I, IVc, VIIa and VIIb, see table 1 and supplementary notes. The Schubart orbit is linearly stable in three spatial dimensions, [35, 36], but due to its collinear nature, it has only two degrees-of-freedom. As it has two degrees-of-freedom, it satisfies the Poincaré–Birkhoff theorem [22–24], which also predicts the existence of infinitely many orbits[17].

Thus, we have shown a definite correlation between the sequences in table 1 and linear stability of the progenitor orbit in each sequence.

5. Virial theorem and analyticity of the action

The remaining mysteries are: (i) why are the $T_{\text{s.i.}}(N_w)$ graphs linear, and (ii) why are the slopes of different sequences so close to each other?

Our answers to these questions are still not proven in a sufficiently rigorous way. Therefore, we shall present them here in the same, or similar way, as they were discovered; otherwise the motivation, and the weak points of our arguments might be lost.

It should be clear that the mere formulation of $T_{\text{s.i.}} = T|E|^{3/2}$ depends crucially on the homogeneity of the Newtonian potential: the exponent 3/2 follows from the Newtonian potential's degree of homogeneity $\alpha = 1$, see [15, 19]. So, one may ask if the same, or similar behaviour occurs in other homogeneous potentials? A (partial) answer to this question was provided in [44], where periodic three-body orbits in the so-called strong potential $V^{\alpha=2}(r) \simeq -1/r^2$ and their relation to topology were studied, which has led to our proposed answer to question (i). The strong potential $V^{\alpha=2}(r) \simeq -1/r^2$, is also homogeneous, see appendix D.

It was shown in [44] that the periodic solutions to the three-body problem in the strong potential form sequences, very much like those in the Newtonian potential shown in section 3, but their periods do *not* increase linearly with the topological complexity N_w of the orbit. Rather, it is the action integral, $S_{\text{min}} \simeq N_w$, that rises linearly with N_w, which fact can be understood using Cauchy's residue theorem, which is based on the analyticity of the action integral,

$$S_{\text{min}}^{\alpha=2} = -2\int_0^T V^{\alpha=2}(\mathbf{r}(t))\mathrm{d}t,$$

where $\mathbf{r}(t)$ is a periodic solution to the equations-of-motion (e.o.m.) at fixed energy $E = 0$, see appendix E.

But, in the Newtonian potential the action of (any) periodic orbit is proportional to its period $S_{\text{min}}^{\alpha=1}(T) = 3|E|T$, see equation (D.5), derived in appendix D.2. So, the scale-invariant period $T_{\text{s.i.}}$ must depend in the same way on the topological complexity N_w of the orbit as the corresponding action $S_{\text{min}}^{\alpha=1}(T)$. The question now arises if the same argument as in [44], about the analyticity of the action $S_{\text{min}}^{\alpha=1}(T)$ can be extended to the Newtonian potential?

In the Newtonian potential this argument becomes more complicated because the hyperradius $R = |Z|$ is not constant in Newtonian three-body orbits, and the problem becomes one in the calculus of two complex variables, see appendices A and E. This leads to new possibilities that have not been considered thus far. Indeed, the second complex variable in the Newtonian potential immediately leads to the possibility that there is a pole in the second complex variable Z, which could lead to non-zero contributions to the integral, and thus change the $T_{\text{s.i.}}(N_w)$ functional dependence, under right conditions.

Assuming that the variation of periodic orbits in the second complex variable Z is limited such that no new poles arise in the action integral, see appendix E, we may conclude that

[17] We see that one colliding orbit is the progenitor of more than one sequence of collisonless orbits.

$$S^{\alpha}_{\min} = \left(\frac{\alpha+2}{\alpha-2}\right) E\, T \simeq N_w.$$

This cannot be true in general, however: a moment's thought shows that the linear dependence cannot hold in the harmonic oscillator, as all harmonic oscillatory motions have the same period there. More formally, equation $S^{\alpha}_{\min} = \left(\frac{\alpha+2}{\alpha-2}\right) E\, T$, implies that the action of a periodic orbit in the harmonic oscillator always vanishes $S^{\alpha=-2}_{\min} = 0$. Moreover, we note that the action integral equation (D.4) must have (at least one) pole if the residue theorem should hold. Consequently, there is an upper bound on the exponent: $\alpha \geqslant 0$, for which this kind of action-topology dependence can exist.

These arguments provide also a (possible) answer to question (ii) above, as the slope of of the $T_{\text{s.i.}}(N_w)$ graph depends on the residue(s) at the same poles in all sequences, the main difference being the ordering of circles around the poles, i.e. of the Riemann sheet(s) one is on ('crossings of branch cuts'), see appendix E.

Of course, the foregoing arguments do not constitute a mathematical proof—the missing dots on the i's and crosses on the t's, or, perhaps more interestingly, counter-arguments/proofs—ought to be supplied by the interested reader.

6. Summary, discussion and outlook

We have shown that:

(1) The presently known periodic three-body orbits with vanishing angular momentum and passing through an Euler configuration, can be classified into 10 sequences according to their topologies. Each sequence probably extends to infinitely long periods, and emerges from one of six linearly stable (shortest-period) progenitor orbits.
(2) Numerically, the scale-invariant periods of orbits in each sequence obey linear dependences on the number of symbols in the algebraic description of the orbit's topology.
(3) There is a possible explanation for the existence of this infinity of periodic orbits, in the form of Birkhoff–Lewis theorem, provided that each progenitor orbit also satisfies the 'twist' condition [16].
(4) Some of the longer-period orbits are linearly stable: (a) the seventh satellite of 'figure-eight' orbit[18]; (b) moth II, which lies in, but is not the progenitor of the 'moth I' sequence; and (c) the 'bumblebee' orbit, which lies in, but is not the progenitor of the 'dragonfly' sequence.

We note that in 1976 [35], Hénon established the linear stability of many orbits with non-vanishing angular momenta ($L \neq 0$) in the Broucke–Hadjidemetriou–Hénon family. The topological-power satellites of these linearly stable BHH orbits were discovered only recently [39], where an $L \neq 0$ version of the period-topology linear dependence equation (2) was checked numerically, as well. The agreement there is also (only) approximate, as a small, but numerically significant discrepancy exists.

Furthermore, [44] indicates that a linear dependence of the action, but not of the period, on the topology exists also in the case of periodic three-body orbits in the so-called strong Jacobi–Poincaré potential, which is in agreement with the virial theorem, see appendix D. The argument in [44] can be extended to the Newtonian potential, but it becomes a complicated question in the calculus of two complex variables[19].

[18] The stability of 'figure-eight' orbit was established in [32, 33].
[19] Indeed, the second complex variable in the Newtonian potential immediately leads to new possibilities: there is a pole in the second variable, which could lead to non-zero contributions, and thus change the $T_{\text{s.i.}}(N_w)$ function, under right conditions.

Our results are generic, so they imply that similar linear relations may be expected to hold for 3-body orbits in the Coulombian[20], and in all other homogeneous potentials containing poles.

Moreover, similar functional dependences might also hold for 4-, 5-, 6-body etc orbits in the Newtonian potential.

Our results also raise new questions:

(1) Each of the six progenitors generates a family of orbits, at different masses and non-vanishing angular momenta, e.g. the Schubart colliding orbit [34], generates the BHH family of collisionless orbits with non-zero angular momenta, that describe the majority of presently known triple-star systems. The remaining five progenitors may now be viewed as credible candidates for astronomically observable three-body orbits, provided that their stability persists under changes of mass ratios and of the angular momentum. Those dependences need to be explored in detail.
(2) Checking the 'twist' condition of the Birkhoff–Lewis theorem, for each progenitor orbit, is a task for mathematicians, as is the explanation of the topologies of the so-predicted orbits: why do these sequences exist and not some others?
(3) The question of existence of other stable two-dimensional colliding orbits, and of new sequences of periodic orbits that they (may) generate. Rose's new linearly stable colliding orbits [13] are particularly interesting in this regard. Turning the foregoing argument around, one can use any newly observed sequence of orbits to argue for the the existence of its, perhaps as yet unknown, progenitor.
(4) A remaining mystery is why are the slopes of different sequences so close to each other?

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Appendix A. Three-body variables

The graphical representation of the three-body system can be simplified with the use of translational and rotational invariance—by changing the coordinates to the Jacobi ones [30]. Jacobi or relative coordinates are defined by two relative coordinate vectors, see figure A1:

[20] Several such periodic orbits have been found in [45, 46], but their topological classification was not considered.

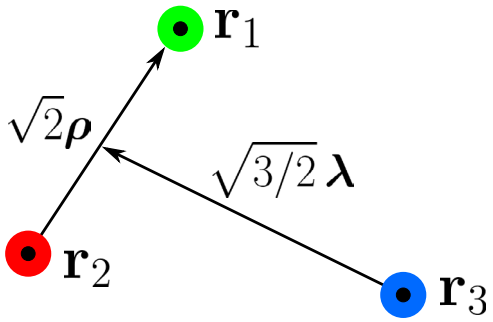

Figure A1. The two three-body Jacobi coordinates $\boldsymbol{\rho}, \boldsymbol{\lambda}$.

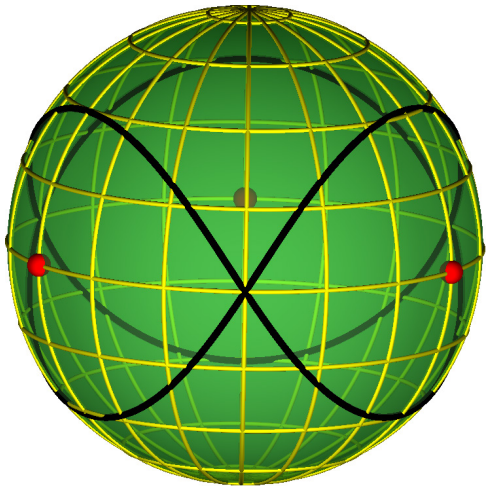

Figure A2. The shape-space sphere: the figure-eight orbit (solid black curve); three two-body collision points (red), singularities of the potential, lie on the equator.

$$\boldsymbol{\rho} = \frac{1}{\sqrt{2}}(\mathbf{r}_1 - \mathbf{r}_2), \quad \boldsymbol{\lambda} = \frac{1}{\sqrt{6}}(\mathbf{r}_1 + \mathbf{r}_2 - 2\mathbf{r}_3). \tag{A.1}$$

Three independent scalar variables can be constructed from Jacobi coordinates: $\boldsymbol{\rho}^2$, $\boldsymbol{\lambda}^2$ and $\boldsymbol{\rho} \cdot \boldsymbol{\lambda}$. The overall size of the orbit is characterized by the hyperradius $R = \sqrt{\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2}$. These scalar variables are connected to the unit vector with Cartesian components [17]:

$$\hat{\mathbf{n}} = \left(\frac{2\boldsymbol{\rho} \cdot \boldsymbol{\lambda}}{R^2}, \frac{\boldsymbol{\lambda}^2 - \boldsymbol{\rho}^2}{R^2}, \frac{2(\boldsymbol{\rho} \times \boldsymbol{\lambda}) \cdot \mathbf{e}_z}{R^2} \right). \tag{A.2}$$

Therefore, every configuration of three bodies (shape of the triangle formed by them, independent of size) can be represented by a point on a unit sphere. This sphere is called the shape-sphere.

Every relatively periodic orbit of a three-body system is therefore represented on the shape-sphere by a closed curve (collisionless solutions), a finite open section of a curve (free-fall and colliding solutions), or a point (Lagrange–Euler solutions). One example, the figure-eight orbit, is illustrated in figure A2.

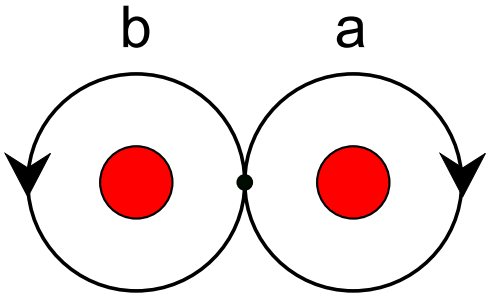

Figure B1. The two elements (a, b) of the free group.

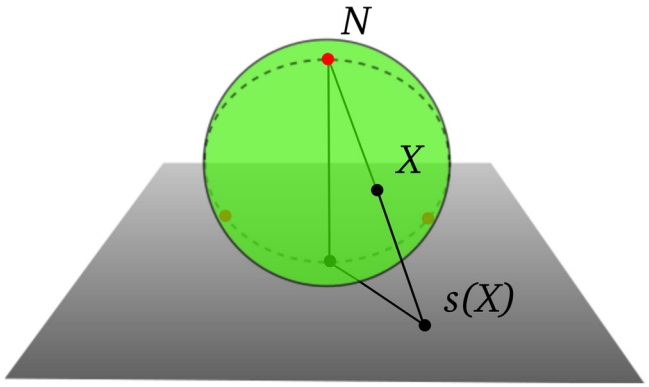

Figure B2. Stereographic projection of a sphere onto a plane. Three two-body collision points (solid red) lie on a meridian (dashed circle), with one of them being at the north pole (denoted by the letter *N*).

The north and the south pole of the shape-sphere correspond to equilateral triangles, while the equator corresponds to degenerate triangles, where the bodies are in collinear configurations (syzygies). There are three points on the equator that correspond to two-body collision points—the singularities of the potential, see figure A2.

Two orbits with identical representations on the shape-sphere are considered to be the same solution. For example, periodic orbits subjected to symmetry transformations, such as translations, rotations, dilations, reflections of space and time, all have identical curves on the shape-sphere and are counted as one.

Size or energy scaling, $\mathbf{r} \to \alpha\mathbf{r}$, and the equations of motion imply $t \to \alpha^{3/2}t$ [31]. Therefore, the velocity scales as $\mathbf{v} \to \mathbf{v}/\sqrt{\alpha}$, the total energy scales as $E \to \alpha^{-1}E$, and the period T as $T \to \alpha^{3/2}T$. Consequently, the combination $|E|^{3/2}T$ is invariant under scale transformations and we call it scale invariant period $T_{\text{s.i.}} = |E|^{3/2}T$. It is always possible to remove one of the three scalar variables by changing the hyper-radius to the desired value by means of these scaling rules.

Appendix B. Montgomery's topological identification method

A curve corresponding to a collisionless periodic orbit can not pass through any one of the three two-body collision points. Stretching this curve across a collision point would therefore change its topology. The classification problem of closed curves on a sphere with three

punctures is given by the conjugacy classes of the fundamental group, which is in this case the free group on two letters (a, b), see figure B1.

This abstract notation has a simple geometric interpretation: it classifies closed curves in a plane with two punctures according to their topologies. The shape sphere can be mapped onto a plane by a stereographic projection using one of the punctures as the north pole, see figure B2. The selected puncture is thusly removed to infinity, which leaves two punctures in the (finite) plane. Any closed curve on the shape sphere (corresponding to a periodic orbit) can now be classified according to the topology of its projection in the plane with two punctures. Topology of a curve can be algebraically described by a 'word'—a sequence of letters a, b, A and B—which is, more formally, an element of the free group F_2. Here a denotes a clockwise full turn around the right-hand-side puncture, b the counter-clockwise full turn around the left-hand-side puncture (see figure B1), and the upper case letters denote their inverse elements $\mathtt{a}^{-1} = \mathtt{A}$ and $\mathtt{b}^{-1} = \mathtt{B}$.

A specific periodic orbit can be equally well described by several different sequences of letters. As there is no preferred starting point of a closed curve, any other word that can be obtained by a cyclic permutation of the letters in the original word represents the same curve.

The conjugacy class of a free group element (word) contains all cyclical permutations of the letters in the original word. For example, the conjugacy class of the free group element aB also contains the cyclically permuted word Ba. The class of topologically equivalent periodic orbits therefore corresponds not merely to one specific free group element, but to the whole conjugacy class.

Time-reversed orbits are represented by the inverse elements of the original free group elements. Naturally, they correspond to physically identical solutions, but they generally form different words (free group elements) with different conjugacy classes.

Another ambiguity is related to the choice of the puncture to be used as the north pole of the stereographic projection (of the sphere onto the plane). A single loop around any one of the three punctures on the original shape sphere (denoted by a or b) must be equivalent to the loop around either of the two remaining punctures. But as can be seen in figure B2, a simple loop around the third ('infinite') puncture on the shape sphere corresponds to aB, a loop around both poles in the plane. Therefore, aB must be equivalent to a and b.

Some periodic solutions have free group elements that can be written as $w^k = w^k(\mathtt{a}, \mathtt{b}, \mathtt{A}, \mathtt{B})$, where $w = w(\mathtt{a}, \mathtt{b}, \mathtt{A}, \mathtt{B})$ is a word that describes some solution, and k is an integer. Such orbits will be called topological-power satellites. For example, the orbits with free group element $(\mathtt{abAB})^k$ are called figure-eight (k) satellites, and are all free from the stereographic projection ambiguity.

Appendix C. Tanikawa and Mikkola's (syzygy) method of topological identification

There is an alternative method of assigning a sequence of three symbols, in this case three digits (1,2,3), to any given 'word' in the free group F_2. It has been proposed for collisionless orbits, by [18], see also [21], to use the sequence of syzygies (collinear configurations) as a symbolic dynamics for the 3-body problem.

The rules for converting 'words' consisting of letters a, b, A, B into 'numbers' consisting of three digits—(1, 2, 3)—are as follows: (i) make the substitution $\mathtt{a} = 12$, $\mathtt{A} = 21$, $\mathtt{b} = 32$, $\mathtt{B} = 23$; (ii) $11 = 22 = 33 =$ empty sequence ('cancellation in pairs rule'). So, for example:

(1) The symbolic sequence corresponding to the BHH family of orbits, $\mathtt{aB} = 1223 = 13$ is equivalent, by way of cyclic permutations, to: $\mathtt{a} = 12$ and to $\mathtt{B} = 23$, as one would expect intuitively. Thus we see that the 'lengths' N_n, i.e. the number of symbols in a sequence are identical for all three symbolic sequences representing the BHH family, $N_n(13) = N_n(12) = N_n(23)$, unlike the Montgomery's method, where $N_w(\mathtt{aB}) \neq N_w(\mathtt{a}) = N_w(\mathtt{B})$. This indicates that the 'lengths' $N_n(w)$ are good algebraic descriptors of the complexity of an orbit's topology.
(2) The symbolic sequence $\mathtt{abAB} = \mathtt{(12)(32)(21)(23)} = 12322123 = 123123 = \mathtt{(123)}^2$ corresponding to the figure-eight orbit is now manifestly invariant under cyclic permutations, $1 \to 2 \to 3$ and $1 \to 3 \to 2$, whereas it is so only in a non-manifest way in the two-letter scheme. Here, also, the 'length' $N_n(w)$ is also a good algebraic descriptor of the complexity of an orbit's topology.

Note that:

1. As stated above, the numbers 1, 2, and 3 can be viewed as denoting syzygies, i.e. crossings of the equator on the shape sphere, in one of three corresponding segments on the said equator, where the index of the body passing between the other two is used as a symbol.
2. Each symbol is its own inverse, which accounts for the 'cancellation in pairs' rule[21]. This circumstance leads to the reduction (by a factor of two) of the number of symbolic sequences denoting one topology, as the time-reversed orbit has an identical symbolic sequence to the original one (which is not the case in the two-letter scheme); and
3. That the cyclic permutation symmetry indicates irrelevance of which syzygy is denoted by which digit.

In this way, we have restored the three-body permutation symmetry of the problem into the algebraic notation describing the topology of a periodic three-body orbit, albeit at the price of having three symbols, rather than two. This restoration of permutation symmetry also implies an absence of the 'automorphism ambiguity' [15]. Such three-symbol sequences have been used e.g. in [18, 21] to identify the topology of periodic three-body orbits.

The length of a sequence of symbols necessary to describe any given topology generally increases by a factor close to 1.5 as one switches from two letters N_w to three digits N_s, as symbols used, i.e. $N_s \simeq 1.5N_w$. The precise value of this proportionality factor ($\simeq 1.5$) is not important for our purposes, as we shall be concerned with the length(s) of symbolic sequences with a well-defined algebraic form, such as $\mathtt{w}_1(\mathtt{w}_2)^{\mathtt{n}}\mathtt{w}_3(\mathtt{w}_4)^{\mathtt{n}}$, where $n = 1, 2, 3, \cdots$. In such a case, the following relation holds $N[\mathtt{w}_1(\mathtt{w}_2)^{\mathtt{n}}\mathtt{w}_3(\mathtt{w}_4)^{\mathtt{n}}] \simeq N[\mathtt{w}_1\mathtt{w}_3] + nN[\mathtt{w}_2\mathtt{w}_4]$ using either set of symbols for $\mathtt{w}_\mathtt{i}$. Only the value of the slope parameter changes as one switches from one set to another. Of course, it is an additional mystery if and when the slopes of different sequences happen to coincide.

Appendix D. Virial theorem and the action of periodic orbits in homogeneous potentials

D.1. The Lagrange–Jacobi identity and the virial theorem

We know that the Lagrange–Jacobi identity [30],

[21] This is only possible for periodic orbits that form closed loops on the shape sphere; otherwise one would have to define one symbol for crossing the equator from above and another one for crossing from below.

$$\frac{1}{2}\frac{\mathrm{d}G}{\mathrm{d}t} = 2K_{\text{total}} + \alpha V^{\alpha}_{\text{total}}, \tag{D.1}$$

where $G = \sum_{i=1}^{N} \mathbf{q}_i \cdot \mathbf{p}_i$ is the so-called virial, gives a relation between kinetic $K_{\text{total}} = \sum_i K_i$ and potential energy $V^{\alpha}_{\text{total}}$, for homogeneous potentials with homogeneity degree $-\alpha$. One example of such a homogeneous potential is the sum of two-body terms $\sum_{i<j} V_\alpha(r_{ik})$, where $V_\alpha(r_{ik}) \simeq -1/r^{\alpha}_{ik}$ is a power-law interaction . Here r_{ik} is the distance between two particles, and α is a positive real number.

For periodic motions, with period T, this identity can be integrated to yield

$$\begin{aligned}\frac{1}{2}\int_0^T \mathrm{d}t \frac{\mathrm{d}G}{\mathrm{d}t} &= \frac{1}{2}\left(G(T) - G(0)\right) = 0 \\ &= \int_0^T (2K_{\text{total}} + \alpha V^{\alpha}_{\text{total}})\mathrm{d}t \end{aligned} \tag{D.2}$$

which tells us that the time integral of the kinetic energy is related to the time integral of the potential energy:

$$\int_0^T \mathrm{d}t K_{\text{total}} = -\frac{\alpha}{2}\int_0^T \mathrm{d}t V^{\alpha}_{\text{total}}.$$

Energy conservation

$$E = K_{\text{total}} + V^{\alpha}_{\text{total}}$$

implies

$$E = \frac{1}{T}\int_0^T (K_{\text{total}} + V^{\alpha}_{\text{total}})\mathrm{d}t = \frac{1}{T}\int_0^T (-\frac{\alpha}{2}V^{\alpha}_{\text{total}} + V^{\alpha}_{\text{total}})\mathrm{d}t$$

which leads to the equipartition of energy (or 'virial') theorem:

$$E = \left(\frac{\alpha-2}{-2}\right)\frac{1}{T}\int_0^T V^{\alpha}(r(t))\mathrm{d}t \equiv \left(\frac{\alpha-2}{-2}\right)\langle V^{\alpha}(r)\rangle \tag{D.3}$$

$$E = \left(\frac{\alpha-2}{\alpha}\right)\frac{1}{T}\int_0^T K(\dot{r}(t))\mathrm{d}t \equiv \left(\frac{\alpha-2}{\alpha}\right)\langle K(\dot{r}(t))\rangle \tag{D.4}$$

which holds exactly for periodic orbits. This, in turn, reduces the action S to one or another time integral.

D.2. The action for three-body orbits in a homogeneous potential

The (minimized) action of a periodic n-body orbit in a homogeneous potential $V^{\alpha}(r) \simeq -1/r^{\alpha}$ is

$$S_{\text{min}} = \int_0^T L(q(t), \dot{q}(t))\mathrm{d}t = \int_0^T \left(T(\dot{r}(t)) - V^{\alpha}(r(t))\right)\mathrm{d}t,$$

leads to

$$S^{\alpha}_{\text{min}}(T) = \left(\frac{\alpha+2}{\alpha-2}\right)E\,T, \tag{D.5}$$

which depends only on the energy E and period T of the orbit. Note the singularity on the right-hand-side of equation (D.5) at $\alpha = 2$, which demands that $E = 0$ in that case. For the Newtonian case, $\alpha = 1$, equation (D.5) leads to

$$S_{\min}^{\alpha=1}(T) = -3ET = 3|E|T,$$

as claimed in [15].

Appendix E. Complex variables and analytic properties of the action

Here we follow closely appendix C in [44]. The minimized action $S_{\min}^{\alpha} = \int_0^T L(q(t), \dot{q}(t))\mathrm{d}t$ of a periodic orbit $q(t)$ in the homogeneous (power) potential $V^{\alpha}(r)$, written as a time integral of twice the kinetic energy K over period T,

$$S_{\min}^{\alpha}(T) = \left(\frac{\alpha+2}{\alpha}\right)\sum_{i=1}^{3}\int_0^T \frac{\mathbf{p}_i^2}{2m}\mathrm{d}t = \left(\frac{\alpha+2}{\alpha}\right)\sum_{i=1}^{3}\int_{\mathbf{r}_i(0)}^{\mathbf{r}_i(T)} \mathbf{p}_i \cdot \mathrm{d}\mathbf{r}_i \tag{E.1}$$

where $m = 1$, can be expressed as a closed-contour integral of two complex variables. After shifting to the relative-motion variables, $(\boldsymbol{\rho}, \boldsymbol{\lambda})$, one finds

$$S_{\min}^{\alpha}(T) = \left(\frac{\alpha+2}{\alpha}\right)\left(\int_{\boldsymbol{\rho}(0)}^{\boldsymbol{\rho}(T)} \mathbf{p}_\rho \cdot \mathrm{d}\boldsymbol{\rho} + \int_{\boldsymbol{\lambda}(0)}^{\boldsymbol{\lambda}(T)} \mathbf{p}_\lambda \cdot \mathrm{d}\boldsymbol{\lambda}\right).$$

The real Jacobi two-vectors $\boldsymbol{\rho}$ and $\boldsymbol{\lambda}$ may be replaced with two complex variables

$$z_\rho = \rho_x + \mathrm{i}\rho_y,\ z_\lambda = \lambda_x + \mathrm{i}\lambda_y,$$

so that the action $S_{\min}^{\alpha}$, can be rewritten as a (double) closed contour integral in two complex variables:

$$S_{\min}^{\alpha}(T) = \left(\frac{\alpha+2}{\alpha}\right)\left(\int_{z_{\rho(0)}}^{z_{\rho(T)}} \dot{z}_\rho^* \mathrm{d}z_\rho + \int_{z_\lambda(0)}^{z_\lambda(T)} \dot{z}_\lambda^* \mathrm{d}z_\lambda\right).$$

Note that the periodicity of motion $\boldsymbol{\rho}(0) = \boldsymbol{\rho}(T)$, $\boldsymbol{\lambda}(0) = \boldsymbol{\lambda}(T)$ implies $z_\rho(T) = z_\rho(0)$ and $z_\lambda(T) = z_\lambda(0)$, which makes this integral a closed contour one

$$S_{\min}^{\alpha} = \left(\frac{\alpha+2}{\alpha}\right)\left(\oint_{C_\rho} \dot{z}_\rho^* \mathrm{d}z_\rho + \oint_{C_\lambda} \dot{z}_\lambda^* \mathrm{d}z_\lambda\right).$$

If there were only one complex variable, then the so-defined function would be analytic. Indeed, the action of two-body elliptic motion in the Newtonian potential has been evaluated using Cauchy's residue theorem in section 18.16 of [26], and in section 11.8 in [27]. With two complex variables, there is no such guarantee, however. Moreover, the residue theorem for functions of two complex variables is a more complicated matter, see [48–51].

The existence and positions of poles in this (double) contour integral are not manifest in its present form; the same integral is given by equation (D.3) in appendix D.2, $S_{\min}^{\alpha}(T) = \left(\frac{\alpha+2}{-2}\right)\int_0^T V^{\alpha}(r(t))\mathrm{d}t$, due to the virial theorem, however, where the potential $V^{\alpha}(r(t))$ is known to have three singularities (simple poles) at three binary collisions and the time-evolution dependence $r(t)$ of the periodic orbit, which parametrizes the contour. For the Newtonian potential $\alpha = 1$ the binary collisions are regularizable, and this integral has been studied by Sundman [28] with the result that the functions $r_k(u), 1 \leqslant k \leqslant 3$, are holomorphic

in a strip $|\text{Im } u| < \delta$ of the complex plane $u \in C$ which contains the real axis, see section 2.3 in [29]. Since $S_{\min}^{\alpha=1}(T) = S(T) = -\left(\frac{3}{2}\right) u(T)$, we know that the trajectories $r_k(S), 1 \leqslant k \leqslant 3$ are holomorphic functions of the action S in a strip $|\text{Im } S| < \delta$ of the complex plane $S \in C$ which contains the real axis.

Note the following implications of this result: (1) for non-singular potentials ($\alpha < 0$) there are no poles in the potential, and consequently no poles encircled by the contour, so the residue vanishes; (2) for singular potentials ($2 > \alpha > 0$) there are poles in the potential, but the residue depends on the integration contour, i.e. on the trajectory on the shape sphere and its topology w; (3) if the integration contour, i.e. the trajectory on the shape sphere repeats k times the topologically equivalent path, then, for singular potentials ($2 > \alpha > 0$), the residue equals k times the single path residue.

Next, we switch from the real $(\boldsymbol{\rho}, \boldsymbol{\lambda})$, or complex (z_ρ, z_λ) Cartesian Jacobi variables to the curvilinear hyper-spherical variables: the real hyper-radius R and the overall rotation angle $\Phi = \frac{1}{2}(\varphi_\rho + \varphi_\lambda)$, and the two angles parametrizing the shape-sphere, e.g. $(\theta = (\varphi_\rho - \varphi_\lambda), \chi = 2\text{Tan}^{-1}(\frac{\rho}{\lambda}))$. Here $(\varphi_\rho, \varphi_\lambda)$ are the angles subtended by the vectors $(\boldsymbol{\rho}, \boldsymbol{\lambda})$ and the x-axis. Equivalently, we may use the complex variables Z, defined by (R, Φ) and z, defined by way of a stereographic projection from the shape-sphere parametrized by (θ, χ).

The variable Z has limited (bounded) variation for all periodic orbits (with zero angular momentum) studied in this paper. Indeed, the value of $R = |Z| = 0$ occurs only in the 'triple collision' ('der Dreierstoss') orbits, which does not happen in our case. The condition $\Phi = \text{const.}$ is trickier, however, because there are 'relatively periodic' solutions with vanishing angular momentum ($L = 0$) and a non-zero change $\Delta\Phi \neq 0$ of angle Φ over one period. All of the orbits considered in this paper are absolutely periodic, i.e. they have $\Delta\Phi = 0$ over one period, so this *caveat* does not apply. Therefore one may eliminate the complex variable Z from further consideration, at least for the orbits considered here, and the problem becomes (much) simpler.

Thus, we see that the complex integration contour C_z relevant to Cauchy's theorem, $S_{\min} = 2i\pi \sum \text{Res}$, for the considered periodic orbits, is determined solely by the orbit's trajectory on the shape sphere: the only poles relevant to this contour integral are the two-body collision points on the shape sphere. Consequently, the periodic orbits' minimized action (integral) is determined (predominantly) by the topology of the closed contour on the shape sphere, i.e. by the homotopy group element of the periodic orbit, unless there is a closed contour in the $Z = (R, \Phi)$ variable, as well.

Repeated k-fold loops of the contour lead to k times the initial integral, i.e. $S_{\min}(w^k) = 2ki\pi \sum \text{Res} = kS_{\min}(w)$, or, equivalently $T_{\text{s.i.}}(w^k) = kT_{\text{s.i.}}(w)$, as observed in topological satellite orbits in section 3. Crossings of branch cuts[22] provide for the change of residue(s) of the pole(s) at different values of k, which may account for the different values of Res, i.e. for different slopes of $T_{\text{s.i.}}(N_w)$ graphs in different sequences.

Detailed study of analytic properties of the action should be a subject of interest to pure mathematicians, however, [47].

ORCID iDs

V Dmitrašinović https://orcid.org/0000-0003-0192-921X

[22] We have shown in [44] that in the strong potential each of the three poles is also a logarithmic branch cut, which implies a complicated structure of branch cuts and different residues. Similar situation ought to be expected in the Newtonian potential as well.

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Dynamics of weakly interacting bosons in optical lattices with flux

Ana Hudomal,[1] Ivana Vasić,[1] Hrvoje Buljan,[2] Walter Hofstetter,[3] and Antun Balaž[1]
[1]*Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, 11080 Belgrade, Serbia*
[2]*Department of Physics, Faculty of Science, University of Zagreb, 10000 Zagreb, Croatia*
[3]*Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität, 60438 Frankfurt am Main, Germany*

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Realization of strong synthetic magnetic fields in driven optical lattices has enabled implementation of topological bands in cold-atom setups. A milestone has been reached by a recent measurement of a finite Chern number based on the dynamics of incoherent bosonic atoms. The measurements of the quantum Hall effect in semiconductors are related to the Chern-number measurement in a cold-atom setup; however, the design and complexity of the two types of measurements are quite different. Motivated by these recent developments, we investigate the dynamics of weakly interacting incoherent bosons in a two-dimensional driven optical lattice exposed to an external force, which provides a direct probe of the Chern number. We consider a realistic driving protocol in the regime of high driving frequency and focus on the role of weak repulsive interactions. We find that interactions lead to the redistribution of atoms over topological bands both through the conversion of interaction energy into kinetic energy during the expansion of the atomic cloud and due to an additional heating. Remarkably, we observe that the moderate atomic repulsion facilitates the measurement by flattening the distribution of atoms in the quasimomentum space. Our results also show that weak interactions can suppress the contribution of some higher-order nontopological terms in favor of the topological part of the effective model.

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I. INTRODUCTION

Ultracold atoms in optical lattices provide a perfect platform for quantum simulations of various condensed-matter phenomena [1]. Yet, since charge-neutral atoms do not feel the Lorentz force, a big challenge in this field was realization of synthetic magnetic fields. After years of effort, artificial gauge potentials for neutral atoms were implemented by exploiting atomic coupling to a suitable configuration of external lasers [2,3]. These techniques were further extended to optical lattices, leading to the realization of strong, synthetic, magnetic fields. As a result, important condensed-matter models—the Harper-Hofstadter [4] and the Haldane model [5]—are nowadays available in cold-atom setups [6–9]. The key property of these models is their nontrivial topological content. In the seminal TKNN paper [10] it was shown that the quantization of the Hall conductivity observed in the integer Hall effect can be directly related to the topological index of the microscopic model—the Chern number.

Cold-atom realizations of topological models exploit periodic driving, either through laser-assisted tunneling [6,7] or by lattice shaking [8]. Using Floquet theory [11,12], a periodically driven system can be related to the time-independent effective Hamiltonian that corresponds to a relevant condensed-matter system. The mapping is known as Floquet engineering and its important features in the context of optical lattices are discussed in Refs. [13–20]. Because of important differences of cold-atom setups and their condensed-matter counterparts, new quench protocols for probing topological features were proposed [21–25]. Following up on these studies, the deflection of an atomic cloud as a response to external force was used to experimentally measure the Chern number in a nonelectronic system for the first time [26].

While Floquet engineering is a highly flexible and powerful technique, it poses several concerns. One of the main open questions is related to the interplay of driving and interactions which can heat up the system to a featureless, infinite-temperature regime according to general considerations [27,28]. In particular, it is shown that an initial Bose-Einstein condensate in a periodically driven optical lattice may become unstable due to two-body collisions [29] or through the mechanism of parametric resonance [28,30–36]. The preparation protocol, stability and a lifetime of strongly correlated phases, expected in the regime of strong interactions under driving is a highly debated topic at the moment [28,37,38].

In order to further explore the role of weak atomic interactions in probing topological features, here we consider the dynamics of weakly interacting incoherent bosons in a driven optical lattice exposed to an external force. The setup that we consider includes all basic ingredients for the Chern-number measurement [22,26]—the Chern number of the topological band can be extracted from the center-of-mass motion of atomic cloud in the direction transverse to the applied force. We assume an ideal initial state where the lowest topological band of the effective model is almost uniformly populated. The optimal loading sequence necessary to reach this state is considered in Refs. [39,40]. Following the recent experimental study [26], we assume that atoms are suddenly released from the trap and exposed to a uniform force. We perform numerical simulations for the full time-dependent Hamiltonian and take into account the effects of weak repulsive interactions between atoms within the mean-field approximation. We make a comparison between the dynamics governed by the effective and time-dependent Hamiltonian and delineate the

contribution of interactions to the center-of-mass response and to the overall cloud expansion dynamics. Our results show that interactions lead to the undesirable atomic transitions between topological bands [41], but we also find that a weak atomic repulsion can facilitate the Chern-number measurements in several ways.

The paper is organized as follows. In Sec. II we describe the model and introduce a method that we apply for the description of incoherent bosons. In Sec. III we address the dynamics of noninteracting incoherent bosons, and then in Sec. IV we address the regime of weak repulsive interactions. Finally, we summarize our results in Sec. V. Appendixes A to F provide further details.

II. MODEL AND METHOD

In this section, we first present the driven model introduced in Ref. [26], and then derive the corresponding effective model and discuss its basic characteristics. At the end, we explain our choice of the initial state and outline the method that we use to treat the dynamics of weakly interacting incoherent bosons.

A. Effective Floquet Hamiltonian

Interacting bosons in a two-dimensional optical lattice can be described by the Bose-Hubbard Hamiltonian

$$\begin{aligned}\hat{H}_{\mathrm{BH}} = &-J_x \sum_{l,m} (\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}) \\ &- J_y \sum_{l,m} (\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}) \\ &+ \frac{U}{2} \sum_{l,m} \hat{n}_{l,m}(\hat{n}_{l,m} - 1), \end{aligned} \tag{1}$$

where $\hat{a}^\dagger_{l,m}$ and $\hat{a}_{l,m}$ are creation and annihilation operators that create and annihilate a particle at the lattice site $(l, m) = la\mathbf{e}_x + ma\mathbf{e}_y$ (a is the lattice constant), $\hat{n}_{l,m} = \hat{a}^\dagger_{l,m}\hat{a}_{l,m}$ is the number operator, J_x and J_y are the hopping amplitudes along $\mathbf{e}_x$ and $\mathbf{e}_y$, and U is the on-site interaction. In the derivation of the model (1) we use the single-band tight-binding approximation [1]. Although the experimental setup [26] is actually three dimensional, with an additional confinement in the third direction, our study is simplified to a two-dimensional lattice.

In order to engineer artificial gauge field in the experiment [26], hopping along $\mathbf{e}_x$ was at first inhibited by an additional staggered potential

$$\hat{W} = \frac{\Delta}{2} \sum_{l,m} (-1)^l \hat{n}_{l,m}, \tag{2}$$

and then restored using resonant laser light. The experimental setup can be described by a time-dependent Hamiltonian

$$\tilde{H}(t) = \hat{H}_{\mathrm{BH}} + \hat{V}(t) + \hat{W}, \tag{3}$$

where $\hat{V}(t)$ is a time-dependent modulation

$$\begin{aligned}\hat{V}(t) = \kappa \sum_{l,m} \hat{n}_{l,m} \Bigg[&\cos\left(\frac{l\pi}{2} - \frac{\pi}{4}\right) \cos\left(\omega t - \frac{m\pi}{2} + \phi_0\right) \\ &+ \cos\left(\frac{l\pi}{2} + \frac{\pi}{4}\right) \cos\left(-\omega t - \frac{m\pi}{2} + \frac{\pi}{2} + \phi_0\right)\Bigg], \end{aligned} \tag{4}$$

κ is the driving amplitude, and $\omega = \Delta$ is the resonant driving frequency. We set the relative phase ϕ_0 between the optical-lattice potential and the running waves used for laser-assisted tunneling to $\phi_0 = \pi/4$.

Using Floquet theory, the time-evolution operator corresponding to the Hamiltonian (3) can be represented as

$$\hat{U}(t, t_0) = e^{-i\hat{W}t} e^{-i\hat{K}(t)} e^{-i(t-t_0)\hat{H}_{\mathrm{eff}}} e^{i\hat{K}(t_0)} e^{i\hat{W}t_0}, \tag{5}$$

where $\hat{H}_{\mathrm{eff}}$ is the full time-independent effective Hamiltonian that describes slow motion and $\hat{K}(t)$ is the time-periodic kick operator that describes micromotion [13,14].

For the moment, in this subsection we first consider the noninteracting model $U = 0$. We also assume that the driving frequency ω is the highest energy scale, but that it is still low enough that the lowest-band approximation used in deriving Eq. (1) is still valid. In the leading order of the high-frequency expansion, the effective Hamiltonian $\hat{H}_{\mathrm{eff}}$ is given by

$$\begin{aligned}\hat{H}_{\mathrm{eff},0} = &J'_x \sum_{l,m} [e^{i((m-l-1)\pi/2 - \pi/4)} \hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \mathrm{H.c.}] \\ &- J'_y \sum_{l,m} (\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}), \end{aligned} \tag{6}$$

where the renormalized hopping amplitudes are $J'_x = \frac{J_x\kappa}{\sqrt{2}\omega} = J_y$ and $J'_y = J_y(1 - \frac{1}{2}\frac{\kappa^2}{\omega^2})$. A schematic representation of this model is presented in Fig. 1(a). The unit cell is shaded and the full lattice is spanned by the vectors $\mathbf{R}_1 = (4, 0)$ and $\mathbf{R}_2 = (1, 1)$. Particle hopping around a plaquette in the counterclockwise direction acquires a complex phase $-\frac{\pi}{2}$ and the model is equivalent to the Harper-Hofstadter Hamiltonian [4] for the case $\alpha = 1/4$ [4]. The explicit form of the kick operator $\hat{K}(t)$ from Eq. (3) is given in Appendix A.

Following Refs. [13,14], we find that additional corrections of the order J_x^2/ω contribute to the system's dynamics and we introduce another approximation for the effective Hamiltonian

$$\begin{aligned}\hat{H}_{\mathrm{eff},1} = \hat{H}_{\mathrm{eff},0} + \frac{J_x^2}{\omega} \sum_{l,m} (-1)^l (&2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} \\ &+ \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m}). \end{aligned} \tag{7}$$

The derivation of Hamiltonian (7) is given in Appendix A and its schematic representation is given in Fig. 1(b). The J_x^2/ω correction introduces next-nearest-neighbor hopping along x direction with opposite signs for lattice sites with either even or odd x-coordinate l. This term does not change the total complex phase per plaquette, but the unit cell is now doubled and thus the first Brillouin zone is halved. A similar term was engineered on purpose in order to implement the Haldane model [8].

In the next subsection we investigate properties of energy bands of both effective Hamiltonians, $\hat{H}_{\mathrm{eff},0}$ and $\hat{H}_{\mathrm{eff},1}$. We use the units where $\hbar = 1$ and $a = 1$. Unless otherwise stated, we set the parameters to the following values: lattice size 100×100 sites, hopping amplitudes $J'_x = J_y = 1 \equiv J$, and the driving amplitude $\kappa = 0.58\omega$. This value of the driving amplitude was chosen to be the same as in the experiment [26]. In order to set the renormalized hopping amplitude along $\mathbf{e}_x$ to $J'_x = 1$, the initial hopping amplitude has to be

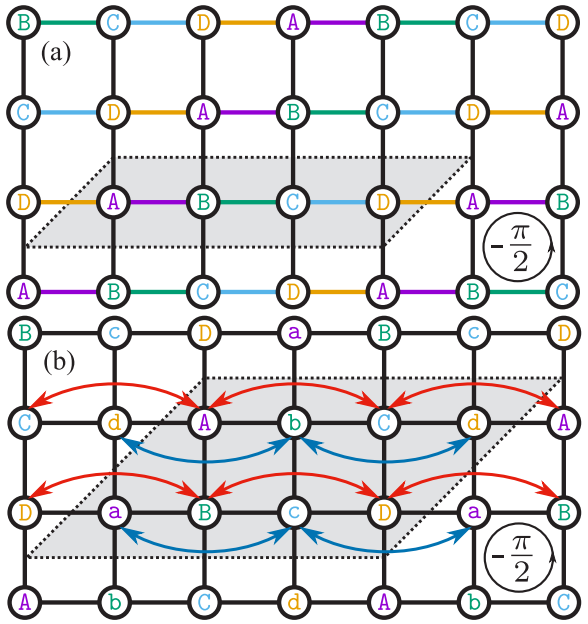

FIG. 1. Schematic representation of the model. The unit cells are shaded. (a) Effective Hamiltonian without correction, $\hat{H}_{\text{eff},0}$ (6). Vertical links correspond to real hopping amplitudes (along $\mathbf{e}_y$ direction), while the horizontal links to the right of lattice sites labeled A, B, C, and D correspond to complex hopping amplitudes with phases $\frac{3\pi}{4}$, $\frac{\pi}{4}$, $-\frac{\pi}{4}$, and $-\frac{3\pi}{4}$, respectively (when hopping from left to right). (b) Effective Hamiltonian with correction, $\hat{H}_{\text{eff},1}$ (7). Red lines represent positive next-nearest-neighbor hopping amplitudes (connecting uppercase letters), while the blue lines represent negative next-nearest-neighbor hopping amplitudes (connecting lowercase letters). Nearest-neighbor hopping amplitudes are the same as in (a).

$J_x = \sqrt{2}\omega/\kappa = 2.44$, and the correction term is therefore proportional to $J_x^2/\omega = 5.95/\omega$, so it cannot be safely neglected unless the driving frequency is very high.

B. Band structure

Momentum-space representations of the effective Hamiltonians $\hat{H}_{\text{eff},0}$ and $\hat{H}_{\text{eff},1}$, denoted by $\hat{\mathcal{H}}_{\text{eff},0}(\mathbf{k})$ and $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$, respectively, are derived in Appendix B. Band structures for the effective Hamiltonian $\hat{\mathcal{H}}_{\text{eff},0}$ without the J_x^2/ω correction, Eq. (B1), as well as for the effective Hamiltonian $\hat{\mathcal{H}}_{\text{eff},1}$ including the correction term, Eq. (B2), are shown in Fig. 2 for the two values of driving frequencies $\omega = 20$ and $\omega = 10$.

The Hamiltonian $\hat{H}_{\text{eff},0}$ is the Harper-Hofstadter Hamiltonian for the flux $\alpha = 1/4$. It has four energy bands, where the middle two bands touch at $E = 0$ and can therefore be regarded as a single band; see Fig. 2(a). The topological content of these bands is characterized by the topological index called the Chern number. The Chern number is the integral of the Berry curvature [42] over the first Brillouin zone divided by 2π,

$$c_n = \frac{1}{2\pi}\int_{\text{FBZ}} \mathbf{\Omega}_n(\mathbf{k})\cdot d\mathbf{S}, \tag{8}$$

where n denotes the band number and the Berry curvature is $\mathbf{\Omega}_n(\mathbf{k}) = i\nabla_{\mathbf{k}} \times \langle u_n(\mathbf{k})|\nabla_{\mathbf{k}}|u_n(\mathbf{k})\rangle$, expressed in terms of eigenstates of the effective Hamiltonian $|u_n(\mathbf{k})\rangle$. The Chern numbers of the three well-separated bands are $c_1 = 1$, $c_2 = -2$, and $c_3 = 1$.

Because the correction from Eq. (7) includes next-nearest-neighbor hopping terms, the elementary cell in real space is doubled [see Fig. 1(b)] and, as a consequence, the first Brillouin zone for the Hamiltonian $\hat{\mathcal{H}}_{\text{eff},1}$ is reduced by a factor of 2 compared to $\hat{\mathcal{H}}_{\text{eff},0}$. There are now eight lattice sites in the unit cell and eight energy bands, but the number of gaps depends on the driving frequency. The new bands touch in pairs, in such a way that there are always maximally three well-separated bands. When the driving frequency is high enough, the correction is small and the gaps between the three bands remain open; see Fig. 2(b). The original band structure of $\hat{\mathcal{H}}_{\text{eff},0}$ is recovered in the limit $\omega \to \infty$. The Berry curvature and the Chern number can be calculated using the efficient method presented in Ref. [43]. Our calculations confirm that the Chern numbers of $\hat{\mathcal{H}}_{\text{eff},1}$ are equal to those of $\hat{\mathcal{H}}_{\text{eff},0}$ ($c_1 = 1$, $c_2 = -2$, and $c_3 = 1$), as long as the gaps between the energy bands are open. The gaps close when the driving frequency is too low, see Fig. 2(c), and the Chern numbers of the subbands can no longer be properly defined.

C. Dynamics of incoherent bosons

We need to take into account a contribution of weak, repulsive interactions. Full numerical simulations of an interacting many-body problem are computationally demanding, so we need a reasonable, numerically tractable approximation. To this end we will use the classical field method [44], which

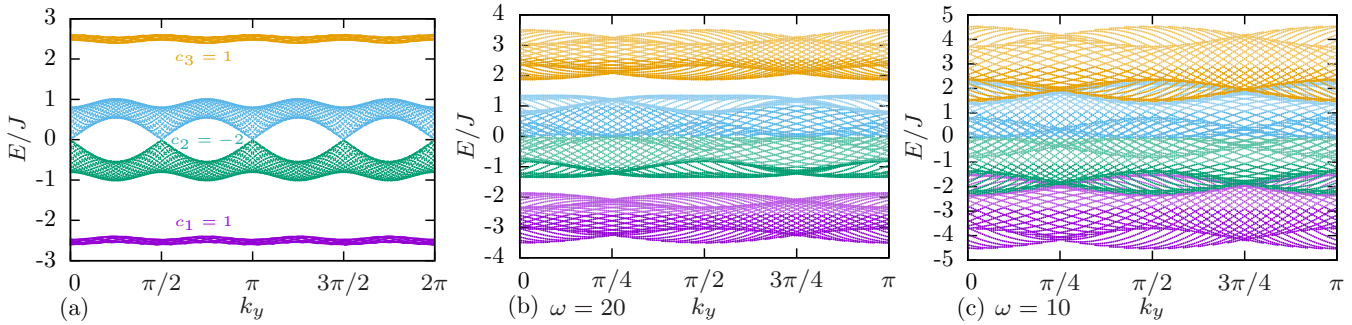

FIG. 2. Energy bands of the effective Hamiltonians. (a) $\hat{\mathcal{H}}_{\text{eff},0}(\mathbf{k})$ Eq. (B1), which is without the J_x^2/ω correction term. (b) $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$ Eq. (B2), which includes the correction term. Driving frequency $\omega = 20$; gaps are open. (c) Same as (b), but with $\omega = 10$. Gaps are closed.

belongs to a broader class of truncated Wigner approaches [45]. This method is similar to the approach used to treat incoherent light in instantaneous media [46,47], known in optics as the modal theory.

The underlying idea of the method is to represent the initial state as an incoherent mixture of coherent states $|\psi\rangle$, $\hat{a}_{l,m}|\psi\rangle = \psi_{l,m}|\psi\rangle$ [44]. This is explained in more detail in Appendix C. In our study, we sample initial configurations of these coherent states with

$$|\psi(t=0)\rangle = \sum_{k=1}^{N_m} e^{i\phi_k}|k\rangle, \tag{9}$$

where $\phi_k \in [0, 2\pi)$ are random phases and the states $|k\rangle$ correspond closely to the lowest-band eigenstates of $\hat{H}_{\text{eff}}$. Each of N_{samples} initial states is time evolved and physical variables can be extracted by averaging over an ensemble of different initial conditions.

The time evolution of each of these coherent states is governed by

$$i\frac{d\psi_{l,m}(t)}{dt} = \sum_{ij} H_{lm,ij}(t)\psi_{i,j}(t) - F\,m\,\psi_{l,m}(t) + U|\psi_{l,m}(t)|^2\psi_{l,m}(t), \tag{10}$$

where $H_{lm,ij}(t) = \langle l, m|\hat{H}(t)|i, j\rangle$ are matrix elements of $\hat{H}(t)$ from Eq. (3), F is the external force, and interactions U contribute with the last, nonlinear term. Formally, Eq. (10) takes the form of the Gross-Pitaevskii equation [48–50]. The performances and limitations of the method are discussed and reviewed in Ref. [51].

For comparison, we also consider the related time evolution governed by the effective Hamiltonian

$$i\frac{d\psi_{l,m}(t)}{dt} = \sum_{ij} h^{\text{eff}}_{lm,ij}\psi_{i,j}(t) - F\,m\,\psi_{l,m}(t) + U|\psi_{l,m}(t)|^2\psi_{l,m}(t), \tag{11}$$

where $h^{\text{eff}}_{lm,ij} = \langle l, m|\hat{h}^{\text{eff}}|i, j\rangle$, with $\hat{h}^{\text{eff}}$ being either $\hat{H}_{\text{eff},0}$ from Eq. (6), or $\hat{H}_{\text{eff},1}$ from Eq. (7). Equation (11) should be considered only as a tentative description of the system: the mapping between $\hat{H}(t)$ and $\hat{H}_{\text{eff}}$ is strictly valid only in the noninteracting regime and the interaction term may introduce complex, nonlocal, higher-order corrections [27]. However, we expect their contribution to be small in the limit $U \to 0$, and for time scales which are not too long [52–55].

In the following we use $N_m = 300$ modes and accommodate $N_p = 300$ particles per mode, so in total in the simulations we have $N = N_m N_p = 90\,000$ bosons. Typical densities in real space are up to 100 particles per site and we choose the values of U in the range $U \in [0, 0.05]$. Other parameters: $J_x' = J_y = 1$, $\kappa/\omega = 0.58$, $\omega = 10, 20$, and $F = 0.25J/a$. The correction terms are non-negligible in this frequency range. In practice, we first numerically diagonalize the Hamiltonian (C2) from Appendix C and set our parameters in such a way that the lowest N_m modes have high overlap with the lowest band of the effective model. In the next step, we sample initial configurations (9). For each of $N_{\text{samples}} = 1000$ sets of initial conditions we then time evolve Eq. (10) and extract quantities of interest by averaging over resulting trajectories. This value of N_{samples} is chosen to be high enough, so that the fluctuations are weak. We present and discuss results of our numerical simulations in the following sections.

TABLE I. Four different cases: the same effective Hamiltonian is always used for the initial state and band definitions, either with or without the correction. The evolution is governed either by the time-dependent Hamiltonian or by the same effective Hamiltonian as the one that was used for the initial state and calculation of band populations.

Case	Initial state	Band populations	Evolution
1	$\hat{H}_{\text{eff},1}$	$\hat{H}_{\text{eff},1}$	$\hat{H}_{\text{eff},1}$
2	$\hat{H}_{\text{eff},1}$	$\hat{H}_{\text{eff},1}$	$\hat{H}(t)$
3	$\hat{H}_{\text{eff},0}$	$\hat{H}_{\text{eff},0}$	$\hat{H}_{\text{eff},0}$
4	$\hat{H}_{\text{eff},0}$	$\hat{H}_{\text{eff},0}$	$\hat{H}(t)$

III. NONINTERACTING CASE

We start by addressing the dynamics of noninteracting bosons. In this case we set $U = 0$ in Eq. (10) and numerically solve the single-particle Schrödinger equation without further approximations. Our aim is to numerically validate and compare the two approximations, Eqs. (6) and (7), for the effective Hamiltonian. To this purpose, we juxtapose results of the two approximative schemes with the numerically exact results obtained by considering the full time evolution governed by $\hat{H}(t)$. For clarity, the four different time evolutions that we consider in this section are summarized in Table I. We calculate the center-of-mass position $x(t)$ and plot the results in Fig. 3. In this way we also find the regime of microscopic

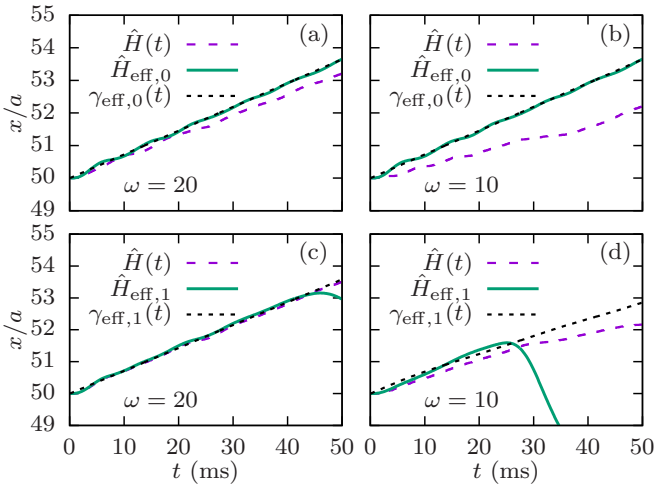

FIG. 3. Anomalous drift $x(t)$. Dashed purple lines: numerical simulations using the time-dependent Hamiltonian $\hat{H}(t)$ (cases 2 and 4 from Table I). Solid green lines: effective Hamiltonians $\hat{H}_{\text{eff},1}$ (c) and (d) and $\hat{H}_{\text{eff},0}$ (a) and (b) (cases 1 and 3). Dotted black lines: theoretical prediction (14) from $\gamma_{\text{eff},1}(t)$ or $\gamma_{\text{eff},0}(t)$. (a) Initial states and band populations obtained using the effective Hamiltonian $\hat{H}_{\text{eff},0}$ without the correction (cases 3 and 4). Driving frequency $\omega = 20$. (b) $\omega = 10$. (c) Hamiltonian $\hat{H}_{\text{eff},1}$ with the J_x^2/ω correction (cases 1 and 2). Driving frequency $\omega = 20$. (d) $\omega = 10$.

parameters where the Chern-number measurement can be optimally performed.

First, we consider the basic Harper-Hofstadter Hamiltonian (6) and select the occupied modes $|k\rangle$ of the initial state (C1) as eigenstates of the model from Eq. (9) for $\hat{h}_{\text{eff}} = \hat{H}_{\text{eff},0}$. As explained in the previous section, at the initial moment $t_0 = 0$, the confinement is turned off and the force $\mathbf{F} = -F\mathbf{e}_y$ is turned on. As a consequence of the applied external force and the nonzero Chern number of the lowest band of the model (6), the particles exhibit an anomalous velocity in the direction perpendicular to the force [56]. In the ideal case, when the lowest band is fully populated, the theoretical prediction for the center-of-mass position in the $\mathbf{e}_x$ direction is [26]

$$x(t) = x(t_0) + c_1 \frac{2Fa^2}{\pi\hbar} t, \tag{12}$$

where $c_1 = 1$ is the Chern number (8) of the lowest band. However, even in the ideal case, due to the sudden quench of the linear potential, a fraction of particles is transferred to the higher bands. To take this effect into account, the authors of Ref. [26] introduced a filling factor $\gamma(t)$

$$\gamma(t) = \eta_1(t) - \eta_2(t) + \eta_3(t), \tag{13}$$

where $\eta_i(t)$ are populations of different bands of Hamiltonian (6) from Eq. (C4) in Appendix C and the plus and minus signs in Eq. (13) are defined according to the Chern numbers $c_1 = 1$, $c_2 = -2$, and $c_3 = 1$. The final theoretical prediction is then [26]

$$x(t) = x(t_0) + c_1 \frac{2Fa^2}{\pi\hbar} \int_0^t \gamma(t')dt'. \tag{14}$$

In Fig. 3(a) we consider the anomalous drift for a high value of the driving frequency $\omega = 20$, where we expect the expansion in $1/\omega$ to be reliable. We find an excellent agreement between the prediction (14) (dotted black line) and numerical calculation based on $\hat{H}_{\text{eff},0}$ (solid green line). However, some deviations between the full numerical results (dashed purple line) and the results of the approximation scheme (solid green line) are clearly visible. These deviations are even more pronounced for $\omega = 10$, Fig. 3(b).

Now we turn to the effective model (7). In this case we select the modes of the initial state as eigenstates of Eq. (9) for $\hat{h}_{\text{eff}} = \hat{H}_{\text{eff},1}$. Moreover, we also consider band populations (C4) of the same model. In the case when $\omega = 20$, Fig. 3(c), the anomalous drift obtained using the effective Hamiltonian (7) (solid green line) closely follows the theoretical prediction (14). Moreover, from the same figure we can see that the effective Hamiltonian $\hat{H}_{\text{eff},1}$ reproduces the behavior of the time-dependent Hamiltonian very well. All three curves almost overlap for intermediate times (5–40 ms); see Fig. 3(c). We attribute the long-time (>45 ms) deviations to the finite-size effects introduced by the next-nearest-neighbor hopping terms, which cause the atomic cloud to reach the edge of the lattice faster. This effect is explained in more detail in Sec. IV B.

For a lower driving frequency $\omega = 10$, the effective and the time-dependent Hamiltonians do not agree so well anymore; see Fig. 3(d). The finite-size effects can be observed even earlier in this case (around 25 ms). This happens because the next-nearest-hopping terms are inversely proportional to the driving frequency. It is interesting to note that the prediction (14) is close to numerical data for short times even in this case when the gaps of the effective model are closed, see Fig. 2(c), and the Chern number of the lowest band is not well defined. In fact, it is surprising that the anomalous drift even exists in this case, as all subbands are now merged into a single band. We attribute this effect to our choice of the initial state. When the gaps are closed, it is hard to set the parameters in such a way that the lowest band is completely filled. The top of this band usually remains empty and the particles thus do not "see" that the gap is closed.

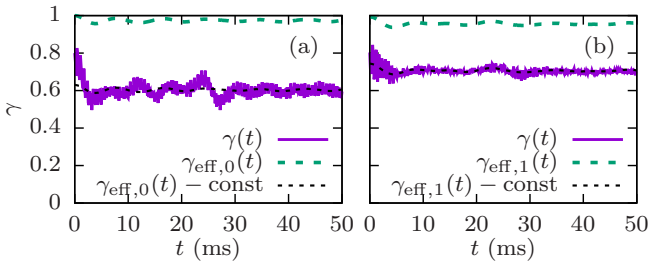

FIG. 4. Time evolution of the filling factor $\gamma(t)$ for driving frequency $\omega = 20$. Solid purple lines: evolution governed by the time-dependent Hamiltonian $\hat{H}(t)$ (cases 2 and 4 from Table I). Dashed green lines: evolution governed by the effective Hamiltonian $\hat{H}_{\text{eff},1}$ or $\hat{H}_{\text{eff},0}$ (cases 1 and 3). Dotted black lines: green lines shifted in order to compare them with purple lines. Shift is chosen so that the two lines approximately overlap. (a) Initial states and band populations obtained using the effective Hamiltonian $\hat{H}_{\text{eff},0}$, which is without the J_x^2/ω correction term (cases 3 and 4). (b) Hamiltonian $\hat{H}_{\text{eff},1}$ which is with the correction term (cases 1 and 2).

Time evolution of the filling factor $\gamma(t)$ is plotted in Fig. 4 for four different cases from Table I—evolution using the effective Hamiltonian without correction $\hat{H}_{\text{eff},0}$ [$\gamma_{\text{eff},0}(t)$, case 3, dashed green line in Fig. 4(a)], the effective Hamiltonian with correction $\hat{H}_{\text{eff},1}$ [$\gamma_{\text{eff},1}(t)$, case 1, dashed green line in Fig. 4(b)], or the time-dependent Hamiltonian $\hat{H}(t)$ [$\gamma(t)$, cases 2 and 4, solid purple lines]. At the initial moment $\gamma(t_0 = 0) < 1$, because the initial state was multiplied by the operator $e^{-i\hat{K}(0)}$. This introduces a shift between $\gamma(t)$ and $\gamma_{\text{eff},1}(t)$. Apart from the shift, these two curves behave similarly, unlike the $\gamma_{\text{eff},0}(t)$ curve that exhibits completely different behavior. Because of this, we use only $\gamma_{\text{eff},1}(t)$ to estimate the value of the prediction (14).

We find that the values of $\gamma_{\text{eff},1}(t)$ for $\omega = 20$ are high: $\geqslant 0.95$; see Fig. 4. For this reason, up to 50 ms the center-of-mass position $x(t)$ exhibits roughly linear behavior with some additional oscillations. Interestingly, the anomalous drift $x(t)$ exhibits quadratic behavior on short time scales in all cases from Fig. 3. In Appendix D, we explain this feature using the time-dependent perturbation theory and Fermi's golden rule.

IV. INTERACTING CASE

We now investigate the effects of weak repulsive interactions. We work in the high-frequency regime and set $\omega = 20$. As shown in Sec. II B, for $U = 0$ the effective Hamiltonian with correction, $\hat{H}_{\text{eff},1}$, is in this case equivalent to the Harper-Hofstadter Hamiltonian with flux $\alpha = 1/4$. Moreover, the

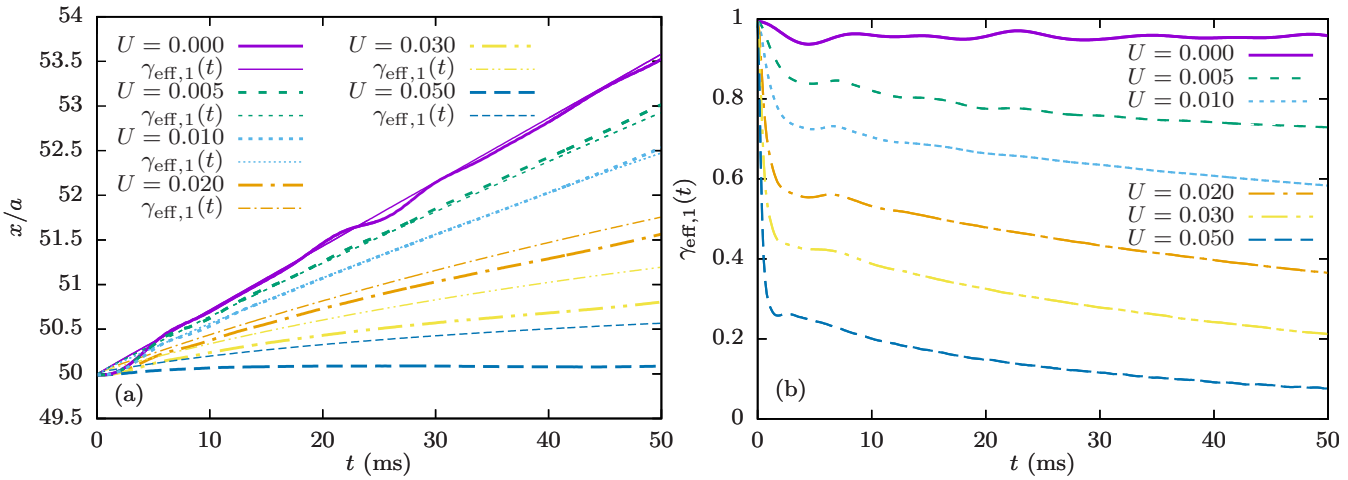

FIG. 5. Effects of interactions. (a) Anomalous drift $x(t)$ for several different values of the interaction coefficient U. U is given in units where $J = 1$. Thick lines: numerical simulations using the time-dependent Hamiltonian $\hat{H}(t)$. Thin lines: theoretical prediction (14) from $\gamma_{\mathrm{eff},1}(t)$. (b) Corresponding $\gamma_{\mathrm{eff},1}(t) = \eta_1(t) - \eta_2(t) + \eta_3(t)$, obtained from simulations using the effective Hamiltonian $\hat{H}_{\mathrm{eff},1}$.

same approximative form of the full effective model accurately reproduces the behavior of the time-dependent Hamiltonian up to 50 ms and thus provides a good starting point for the study of weakly interacting particles. We first consider the anomalous drift of the center of mass of the atomic cloud and then we inspect the expansion dynamics more closely in terms of atomic density distributions in real and momentum space.

A. Anomalous drift and dynamics of band populations

To simulate the dynamics of many incoherent bosons, we use the classical field method presented in Sec. II C and propagate Eq. (10) in time. We assume that at $t_0 = 0$ atoms are uniformly distributed over the lowest band of $\hat{H}_{\mathrm{eff},1}$. For this reason, the initial state is the same as the one that we use in the noninteracting regime. In this way, the dynamics is initiated by an effective triple quench: at $t_0 = 0$ the confining potential is turned off, atoms are exposed to the force $\mathbf{F} = -F\mathbf{e}_y$, and also the interactions between particles are introduced. The total number of particles is set to $N = 90\,000$, which amounts to approximately 100 particles per lattice site in the central region of the atomic cloud. We consider only weak repulsion $U \leqslant 0.05$.

The anomalous drift $x(t)$ obtained using the full time-dependent Hamiltonian is shown in Fig. 5(a) for several different values of the interaction strength U. In comparison to the noninteracting regime, we find that the weak repulsive interactions inhibit the response of the center of mass to the external force. In particular, at $t = 50$ ms the drift is reduced by about 15% for $U = 0.005$ and it is further lowered by an increase in U. Finally, at $U = 0.05$, the anomalous drift is barely discernible. Interestingly, for weak $U \in (0.001, 0.01)$ we find that the drift $x(t)$ in the range of $t \in (10, 50)$ ms looks "more linear" as a function of time in comparison to the noninteracting result.

We now analyze the anomalous drift in terms of the filling factor $\gamma(t)$ and compare the results of Eq. (10) with the description based on Eq. (11). By solving Eq. (11) we obtain the filling factor $\gamma_{\mathrm{eff},1}(t)$ following Eq. (C4) and present our results in Fig. 5(b). Whenever the results of Eq. (10) reasonably agree with the results obtained from Eq. (11), we are close to a steady-state regime with only small fluctuations in the total energy, as Eq. (11) preserves the total energy of the system. In this regime, during the expansion dynamics the interaction energy is converted into the kinetic energy and atoms are transferred to higher bands of the effective model. Consequently, the filling factor $\gamma_{\mathrm{eff},1}(t)$ is reduced. Typically, we find three different stages in the decrease of $\gamma_{\mathrm{eff},1}(t)$.

In an early stage, $t \leqslant t_1 = 5$ ms, a fast redistribution of particles over the bands of the effective model sets in due to the sudden quench of U. The factor $\gamma_{\mathrm{eff},1}(t)$ decays quadratically as a function of time down to $\gamma_{\mathrm{eff},1}(t_1) \approx 0.75$ for $U = 0.01$, and $\gamma_{\mathrm{eff},1}(t_1) \approx 0.25$ for $U = 0.05$. In this process the interaction energy of the system is quickly lowered as described in Appendix E. At later times $t > 5$ ms, we observe a linear decay of the filling factor $\gamma_{\mathrm{eff},1}(t)$ as a function of time, that finally turns into an exponential decay at even later times ($t > 10$ ms). Similar regimes are observed in other dynamical systems. For example, a decay rate of an initial state suddenly coupled to a bath of additional degrees of freedom exhibits these three stages [57]. The initial quadratic decay is often denoted as "the Zeno regime." For longer propagation times, Fermi's golden rule predicts the linear decay. At even longer time scales, when the repopulation of the initial state is taken into account, the time-dependent perturbation theory yields the exponential regime, known under the name of the Wigner-Weisskopf theory [57].

We now investigate this last regime in more detail. For the population of the lowest band $\eta_1(t)$, an exponential decay function $f(t) = a + b\,e^{-ct}$ provides high quality fits for $t \in (10, 50)$ ms; see Fig. 6(a) for an example. Similarly, the populations of two higher bands can also be fitted to exponential functions. The obtained exponential decay coefficients c for the lowest band population are plotted as a function of the interaction strength U in Fig. 6(b). The resulting dependence is approximately quadratic: $c(U) = \alpha_0 + \alpha_1 U + \alpha_2 U^2$. For small values of U, the exponents $c(U)$ obtained for the dynamics governed by $\hat{H}(t)$ and $\hat{H}_{\mathrm{eff},1}$ agree very well and

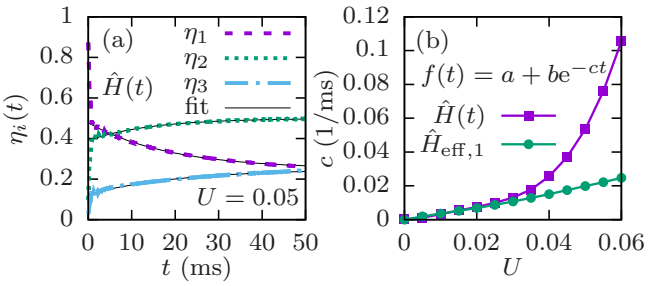

FIG. 6. (a) Evolution of the band populations $\eta_i(t)$. Dashed lines: numerical results obtained using the time-dependent Hamiltonian $\hat{H}(t)$. Solid black lines: exponential fit using $f(t) = a + b\,e^{-ct}$. The coefficient a was fixed to $a_1 = 0.25$, $a_2 = 0.50$, and $a_3 = 0.25$ for the first, second, and third band, respectively. (b) Dependence of the exponential decay coefficients for the lowest band population $\eta_1(t)$ on the interaction strength. U is given in units where $J = 1$.

exhibit linear behavior. At stronger interaction strengths $U \geqslant 0.03$, the approximation of Eq. (11) becomes less accurate as it omits the quadratic contribution in $c(U)$ found in the full time evolution. In addition, the values of the exponents c are affected by the force strength F and driving frequency ω.

As we now understand some basic features of $\gamma_{\mathrm{eff},1}(t)$, we make an explicit comparison between the numerical results for the anomalous drift and the expectation (14). The dashed lines in Fig. 5(a) correspond to the theoretical prediction (14) calculated from $\gamma_{\mathrm{eff},1}(t)$. For the intermediate interaction strengths $U \leqslant 0.01$, we find a very good agreement between the two. From this we conclude that the interaction-induced transitions of atoms to higher bands are the main cause of the reduced anomalous drift $x(t)$ as a function of U. When the interactions become strong enough ($U \sim 0.02$), the numerical results start to deviate from the theoretical prediction (14) with $\gamma_{\mathrm{eff},1}(t)$. In this regime, Eq. (11) does not provide a reliable description of the dynamics, as higher-order corrections need to be taken into account.

B. Real and momentum-space dynamics

So far we have considered the averaged response of the whole atomic cloud. We now inspect the expansion dynamics in a spatially resolved manner. The real-space probability densities at the initial moment and after 50 ms (75 driving periods) are shown in Figs. 7 and 8, and the corresponding momentum-space probability densities in Appendix F.

At the initial moment, the atomic cloud is localized in the center of the lattice. By setting $r_0 = 20$ in the confining potential of Eq. (C2) and populating the lowest-lying states, we fix the cloud radius to $r = 20$, Fig. 7(a). The cloud density is of the order of 100 atoms per lattice site and a weak density modulation is visible along x direction. After the confining potential is turned off, and the external force in the $-\mathbf{e}_y$ direction is turned on, the cloud starts to expand and move in the $+\mathbf{e}_x$ direction. As shown in the previous subsection, the band populations and therefore the anomalous drift are significantly altered by the interaction strength, and this is also the case with the expansion dynamics; see Figs. 7 and 8.

In the noninteracting case, Fig. 7(b), the atomic cloud nearly separates into two parts moving in opposite directions

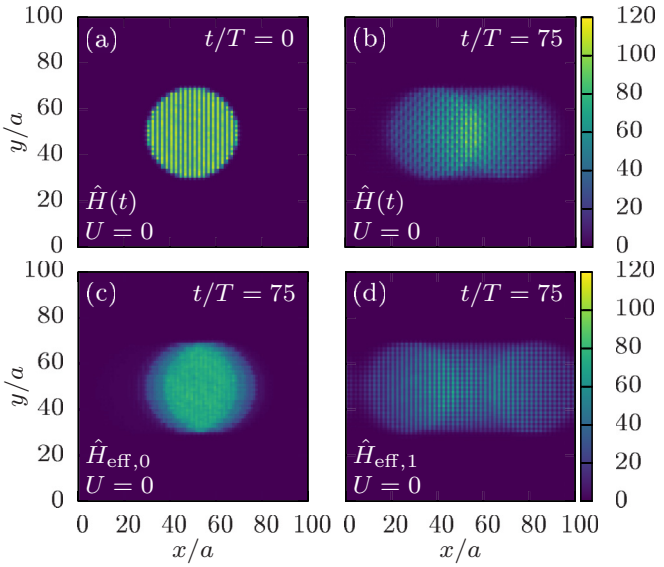

FIG. 7. Real-space density distribution, noninteracting case $U = 0$. (a) Initial state. (b) After 50 ms (75 driving periods), evolution using the time-dependent Hamiltonian $\hat{H}(t)$. (c) Evolution using effective Hamiltonian without correction $\hat{H}_{\mathrm{eff},0}$. (d) Evolution using effective Hamiltonian with correction $\hat{H}_{\mathrm{eff},1}$.

along x axes (while the center of mass still moves in the $+\mathbf{e}_x$ direction). By comparing Fig. 7(c) and Fig. 7(d), we conclude that this effect stems from the next-nearest-neighbor hopping along x present in the effective Hamiltonian (7), as it does not happen in the effective model without the correction term (6). This type of separation was already observed in Ref. [22], where the next-nearest-neighbor hopping terms were also present. When the interactions between particles are included,

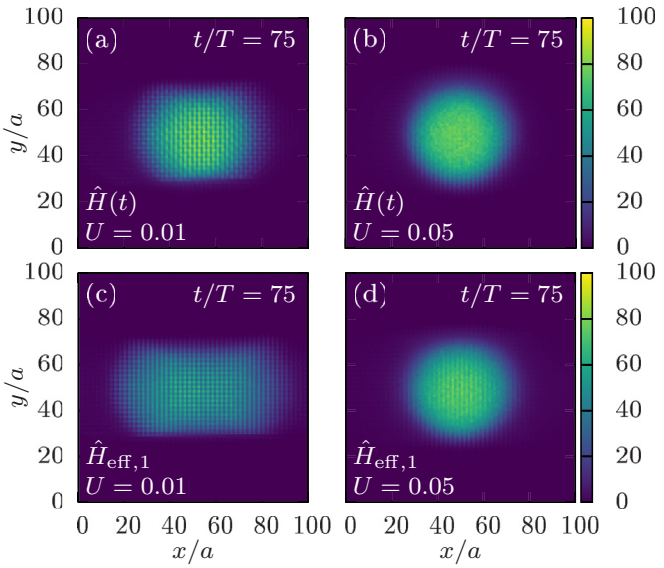

FIG. 8. Real-space density distribution after 50 ms (75 driving periods), interacting case. U is given in units where $J = 1$. (a) Evolution using the time-dependent Hamiltonian $\hat{H}(t)$, $U = 0.01$. (b) Same with $U = 0.05$. (c) Evolution using the effective Hamiltonian $\hat{H}_{\mathrm{eff},1}$, $U = 0.01$. (d) Same with $U = 0.05$.

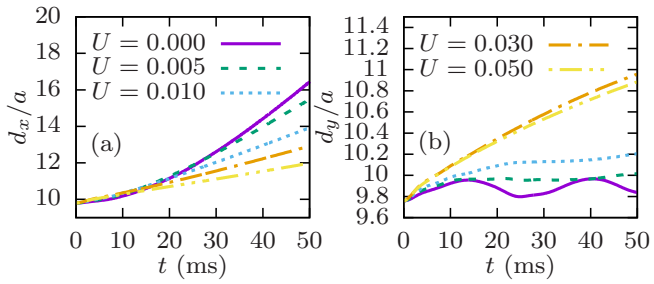

FIG. 9. Atomic cloud width for different interaction strengths, evolution using the time-dependent Hamiltonian $\hat{H}(t)$. U is given in units where $J = 1$. (a) $d_x = \sqrt{\langle x^2\rangle - \langle x\rangle^2}$. (b) $d_y = \sqrt{\langle y^2\rangle - \langle y\rangle^2}$.

this separation is not so prominent [Fig. 8(a), $U = 0.01$], and it almost completely disappears when the interactions are strong enough [Fig. 8(b), $U = 0.05$]. This is also the case when the evolution is governed by the effective Hamiltonian $\hat{H}_{\text{eff},1}$; see Figs. 8(c) and 8(d). Atomic cloud widths $d_x = \sqrt{\langle x^2\rangle - \langle x\rangle^2}$ during the expansion are plotted in Fig. 9. We observe a slow expansion of the cloud in y direction, Fig. 9(b), and much faster expansion along x direction, Fig. 9(a), which comes about as a consequence of the cloud separation. On top of this, we observe that the interactions enhance expansion along y. Surprisingly, the opposite is true for the dynamics along x. This counterintuitive effect is often labeled as self-trapping and its basic realization is known for the double-well potential [58,59]. In brief, strong repulsive interactions can preserve the density imbalance between the two wells, as the system cannot release an excess of the interaction energy. In our case, the situation is slightly more involved as the cloud splitting is inherent (induced by the corrections of the ideal effective Hamiltonian). Apart from this, due to the driving the total energy is not conserved. However, our numerical results indicate that the interaction energy is slowly released in the second expansion stage, Fig. 14. Effectively, in this way the interactions cancel out the contribution of the next-nearest-neighbor hopping and favor the measurement of the properties of the model (6). In Fig. 10(a) we show that deviations between different approximations based on $\hat{H}(t)$, $\hat{H}_{\text{eff},1}$, and $\hat{H}_{\text{eff},0}$ in the anomalous drift $x(t)$ nearly vanish at $U = 0.01$.

Another desirable effect might be that the interactions make the momentum-space probability density more homogeneous, see Appendix F, so that the real-space probability density becomes more localized. We can quantify momentum-space homogeneity using the inverse participation ratio $R(t) = \frac{1}{\sum_i P_i^2(t)}$, where $P_i(t) = |\psi_i(t)|^2$ is the probability that the state ψ_i is occupied at time t. Minimal value of the inverse participation ratio (IPR) is 1 and it corresponds to a completely localized state, while the maximal value is equal to the total number of states (in our case 10 000) and corresponds to the completely delocalized state, where the particles have the same probability of being at any quasi-momentum $\mathbf{k}$. As stated before, the first Brillouin zone of the lowest band has to be as homogeneously populated as possible in order to properly measure the lowest band Chern number. From Fig. 10(b), we see that IPR increases in time when the interaction coefficient U is large enough, so we can conclude that the interactions are actually beneficial for measuring the Chern number, as they can "smooth out" the momentum-space probability density. In Fig. 10(c) we give estimates for the Chern number that can be extracted from our numerical data for different values of U. We find the best estimate $c_1 \sim 0.99$ for the intermediate interaction strength $U \sim 0.01$.

C. Staggered detuning

Here we briefly consider the effects of staggered detuning that was introduced in the experimental study [26] during the loading and band mapping sequences. This detuning can be described by an additional term

$$\frac{\delta}{2}\sum_{l,m}[(-1)^l + (-1)^m]\hat{n}_{l,m} \tag{15}$$

in the Hamiltonians $\hat{H}(t)$ and $\hat{H}_{\text{eff},1}$. We will ignore the higher-order [at most $O(\frac{1}{\omega^2})$] corrections that this term introduces to the effective Hamiltonian. Staggered detuning does not break the symmetry of the effective Hamiltonian $\hat{H}_{\text{eff},1}$, but if δ is large enough, it can cause a topological phase

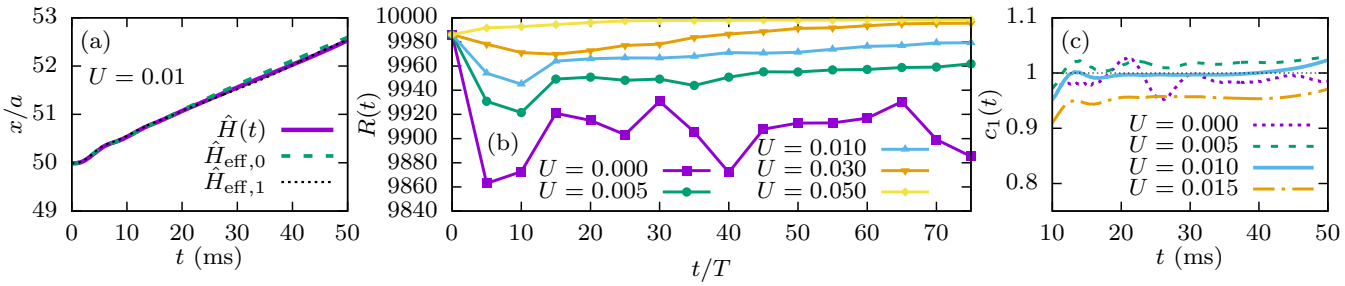

FIG. 10. (a) Comparison of anomalous drifts obtained from evolution using the time-dependent Hamiltonian $\hat{H}(t)$ (solid purple line), effective Hamiltonian without correction $\hat{H}_{\text{eff},0}$ (dashed green line) and effective Hamiltonian with correction $\hat{H}_{\text{eff},1}$ (dotted black line). Intermediate interaction strength $U = 0.01$. U is given in units where $J = 1$. (b) Time evolution of the inverse participation ratio in momentum space for several different values of U. Evolution is performed using the time-dependent Hamiltonian $\hat{H}(t)$. When the interactions are strong enough, IPR approaches the maximal possible value (10 000 in this case), which is equal to the total number of states and corresponds to the completely delocalized state. U is given in units where $J = 1$. (c) Chern number of the lowest band obtained for different interaction strengths as the ratio of the theoretical prediction for the anomalous drift and numerical results: $c_1(t) = (\frac{2Fa^2}{\pi\hbar}\int_0^t \gamma_{\text{eff},1}(t')dt')/[x(t) - x(t_0)]$.

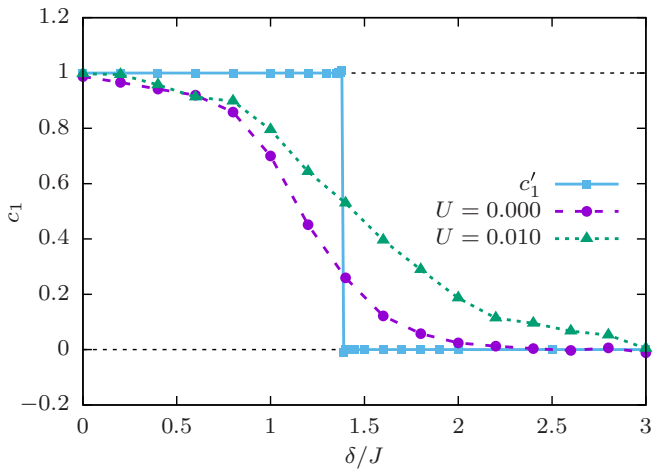

FIG. 11. Lowest band Chern numbers extracted from numerical data for several different values of detuning δ. Purple circles: noninteracting case, $U = 0$. Green triangles: $U = 0.01$. Blue squares: theoretical values of the lowest band Chern number c_1'. A topological phase transition is visible at $\delta_c \approx 1.38$. The lines between points are only a guide to the eye.

transition and make all bands topologically trivial. By numerically calculating the Berry curvature and Chern numbers c_i', we find that this transition occurs at $\delta_c \approx 1.38\,J$; see Fig. 11. This value is lower than the one for the ordinary Harper-Hofstadter Hamiltonian for $\alpha = 1/4$, which is $\delta_c = 2\,J$ [26], due to the different hopping amplitudes J_x' and J_y', and due to the additional J_x^2/ω correction that we consider.

We now investigate how this topological transition can be probed through the dynamical protocol used in the experiment. We again numerically calculate the anomalous drift and the evolution of the filling factor, but now with staggered detuning (15) included in the Hamiltonian $\hat{H}_{\text{initial}}$ (C2) used to obtain the initial state, in the equations of motion (10) and (11), and in the definitions of the band populations $\eta_i(t)$ (C4). Using these results, we repeat the procedure for the extraction of the lowest band Chern number from numerical data that was carried out in the previous section. The Chern number obtained by comparing the anomalous drift to the prediction calculated from the filling factor is then averaged over the time interval $t \in (20, 40)$ ms. This interval was chosen in order to avoid the initial quadratic regime and the finite-size effects at later times. The resulting lowest band Chern numbers for several different values of detuning δ in both the noninteracting case and the case of intermediate interaction strength $U = 0.01$ are presented in Fig. 11.

We can see that the calculated value of the Chern number decreases from $c_1 = 1$ to $c_1 = 0$ with increasing detuning δ. The obtained value of the Chern number is lower than 1 even before the phase transition occurs. This is due to our choice of the initial state, which is not perfectly homogeneous in momentum space. Close to the phase transition, both the energy bands and the Berry curvature have pronounced peaks at the same regions of the first Brillouin zone, and these regions are initially less populated. Because of this, the Berry curvature at these regions contributes less to the anomalous drift, which lowers the measured Chern number. This effect is somewhat reduced by the interactions, as they smooth out the momentum-space probability density, and might also cancel out the detuning term. Similar interplay of interactions and staggering was observed in the fermionic Hofstadter-Hubbard model [60]. The obtained results are in line with experimental measurements [26].

V. CONCLUSIONS

Motivated by the recent experimental results reporting the Chern numbers of topological bands in cold-atom setups, we studied numerically bosonic transport in a driven optical lattice. The considered driving scheme and the range of microscopic parameters were chosen to be close to those in a recent experimental study [26]. The driving frequency was set to be high enough in order to avoid strong energy absorption for the relevant time scales. Additionally, the system was restricted to a two-dimensional lattice, even though the actual experimental setup had continuous transverse degrees of freedom. This restriction stabilizes the system [29,31,41] and leads to lower heating rates than those in the experiment. It corresponds to the case of strongly confined third dimension.

We investigated bosonic dynamics for the full time-dependent Hamiltonian, the effective Floquet Hamiltonian, and included the effects of weak repulsive interactions between atoms using the mean-field approximation. In the noninteracting case, we found that the effective Hamiltonian and its band structure depend on the frequency of the drive ω through an additional J_x^2/ω correction term. The initial state was set as a mixture of incoherent bosons homogeneously populating the lowest band, but a possible direction of future research could be to simulate the full loading sequence of an initial Bose-Einstein condensate and to try to obtain the incoherent state through driving, as it was done in the experiment.

The main focus of this work is on the effects of weak interactions. For a weak atomic repulsion, atomic transitions to higher effective bands obtained in our simulations mainly occur due to a release of the initial interaction energy during the atomic-cloud expansion. Although the effect is undesirable, it can be properly taken into account in the extraction of the Chern number. At larger interaction strengths, the transitions are more pronounced as the system absorbs energy from the drive. In this regime the good agreement between the full and effective description is lost and the measurement should become more complicated. In addition to causing redistribution of atoms over bands, our results show that weak interactions can also be beneficial in measuring the Chern number. Their desirable effect comes about due to smoothening the atomic distribution over the topological band and due to canceling out the contribution of some less relevant terms to the bosonic dynamics.

ACKNOWLEDGMENTS

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APPENDIX A: EFFECTIVE MODEL

After a unitary transformation into the rotating frame $\tilde{\psi} = e^{-i\hat{W}t}\psi$, where $\tilde{\psi}$ and ψ are the old and the new wave functions, and $\hat{W}$ is the staggered potential, the new time-dependent Hamiltonian that describes the experimental setup is given by [26]

$$\hat{H}(t) = J_y \sum_{l,m} (\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}) + \hat{V}^{(+1)}e^{i\omega t} + \hat{V}^{(-1)}e^{-i\omega t} + \frac{U}{2}\sum_{l,m}\hat{n}_{l,m}(\hat{n}_{l,m}-1), \tag{A1}$$

where

$$\hat{V}^{(+1)} = \kappa/2 \sum_{l,m} \hat{n}_{l,m} g(l,m) - J_x \sum_{l_{\text{odd}},m} (\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}), \tag{A2}$$

$$\hat{V}^{(-1)} = \kappa/2 \sum_{l,m} \hat{n}_{l,m} g^*(l,m) - J_x \sum_{l_{\text{even}},m} (\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}), \tag{A3}$$

$$g(l,m) = \cos(l\pi/2 - \pi/4)e^{i(\phi_0 - m\pi/2)} + \cos(l\pi/2 + \pi/4)e^{i(m\pi/2 - \phi_0 - \pi/2)}. \tag{A4}$$

The kick operator is given by

$$\hat{K}(t) = \frac{1}{i\omega}\left(\hat{V}^{(+1)}e^{i\omega t} - \hat{V}^{(-1)}e^{-i\omega t}\right) + O\left(\frac{1}{\omega^2}\right) \tag{A5}$$

and the effective Hamiltonian by

$$\hat{H}_{\text{eff}} = \underbrace{\hat{H}_0}_{\hat{H}^{(0)}_{\text{eff}}} + \underbrace{\frac{1}{\omega}[\hat{V}^{(+1)}, \hat{V}^{(-1)}]}_{\hat{H}^{(1)}_{\text{eff}}} + \underbrace{\frac{1}{2\omega^2}([[\hat{V}^{(+1)}, \hat{H}_0], \hat{V}^{(-1)}] + [[\hat{V}^{(-1)}, \hat{H}_0], \hat{V}^{(+1)}])}_{\hat{H}^{(2)}_{\text{eff}}} + O\left(\frac{1}{\omega^3}\right). \tag{A6}$$

If we assume that the driving frequency is high and interactions are weak, the interaction term and almost all $O(\frac{1}{\omega^2})$ terms can be neglected. After substituting Eqs. (A1), (A2), and (A3) into Eq. (A6) we obtain

$$\hat{H}^{(0)}_{\text{eff}} = -J_y \sum_{l,m} (\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}), \tag{A7}$$

$$\begin{aligned}\hat{H}^{(1)}_{\text{eff}} = \frac{1}{\omega}\Bigg[&\frac{\kappa}{2}\sum_{l,m}\hat{a}^\dagger_{l,m}\hat{a}_{l,m}g(l,m) - J_x\sum_{l_{\text{odd}},m}(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}),\\ &\frac{\kappa}{2}\sum_{l,m}\hat{a}^\dagger_{l,m}\hat{a}_{l,m}g^*(l,m) - J_x\sum_{l_{\text{even}},m}(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m})\Bigg] = \hat{H}_1 + \hat{H}_2 + \hat{H}_3 + \hat{H}_4.\end{aligned} \tag{A8}$$

We will now separately calculate each term:

$$\begin{aligned}\hat{H}_1 &= -\frac{J_x\kappa}{2\omega}\sum_{l_{\text{odd}},m,l',m'} g^*(l',m')[\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}, \hat{a}^\dagger_{l',m'}\hat{a}_{l',m'}]\\ &= -\frac{J_x\kappa}{2\omega}\sum_{l_{\text{odd}},m}[(g^*(l,m) - g^*(l+1,m))\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + (g^*(l,m) - g^*(l-1,m))\hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}],\end{aligned} \tag{A9}$$

$$\begin{aligned}\hat{H}_2 &= -\frac{J_x\kappa}{2\omega}\sum_{l_{\text{even}},m,l',m'} g(l',m')[\hat{a}^\dagger_{l',m'}\hat{a}_{l',m'}, \hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}]\\ &= \frac{J_x\kappa}{2\omega}\sum_{l_{\text{even}},m}[(g(l,m) - g(l+1,m))\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + (g(l,m) - g(l-1,m))\hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}],\end{aligned} \tag{A10}$$

$$\begin{aligned}\hat{H}_3 &= \frac{J_x^2}{\omega}\sum_{l_{\text{odd}},m,l'_{\text{even}},m'}[\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}, \hat{a}^\dagger_{l'+1,m'}\hat{a}_{l',m'} + \hat{a}^\dagger_{l'-1,m'}\hat{a}_{l',m'}]\\ &= \frac{J_x^2}{\omega}\sum_{l_{\text{odd}},m}(2\hat{a}^\dagger_{l+1,m}\hat{a}_{l+1,m} + \hat{a}^\dagger_{l+3,m}\hat{a}_{l+1,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l+1,m} - 2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} - \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} - \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m})\\ &= \frac{J_x^2}{\omega}\sum_{l,m}(-1)^l(2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m}),\end{aligned} \tag{A11}$$

$$\hat{H}_4 = \frac{\kappa^2}{4\omega} \sum_{l,m,l',m'} g(l,m)g^*(l',m')[\hat{a}^\dagger_{l,m}\hat{a}_{l,m}, \hat{a}^\dagger_{l',m'}\hat{a}_{l',m'}] = 0. \tag{A12}$$

Using trigonometric identities and

$$g(l,m) - g(l \pm 1, m) = \pm\sqrt{2}(\sin[(2l \pm 1 - 1)\pi/4]e^{i(\pi/4 - m\pi/2)} + \sin[(2l \pm 1 + 1)\pi/4]e^{i(m\pi/2 - 3\pi/4)}), \tag{A13}$$

we can rewrite the sum of terms (A9) and (A10) in a more convenient form:

$$\hat{H}_1 + \hat{H}_2 = \frac{J_x\kappa}{\sqrt{2}\omega} \sum_{l,m} (e^{i[(m-l)\pi/2 - \pi/4]}\hat{a}^\dagger_{l,m}\hat{a}_{l-1,m} + e^{-i[(m-l-1)\pi/2 - \pi/4]}\hat{a}^\dagger_{l,m}\hat{a}_{l+1,m}). \tag{A14}$$

The only $O(\frac{1}{\omega^2})$ $(\hat{H}^{(2)}_{\text{eff}})$ term that cannot be neglected in the parameter range that we use is [26]

$$\frac{J_y}{2}\frac{\kappa^2}{\omega^2} \sum_{l,m} (\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}). \tag{A15}$$

Finally, the effective Hamiltonian becomes

$$\hat{H}_{\text{eff},1} = \frac{J_x\kappa}{\sqrt{2}\omega} \sum_{l,m} (e^{i[(m-l-1)\pi/2 - \pi/4]}\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + e^{-i[(m-l)\pi/2 - \pi/4]}\hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}) - J_y\left(1 - \frac{1}{2}\frac{\kappa^2}{\omega^2}\right) \sum_{l,m} (\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}) \tag{A16}$$

$$+ \frac{J_x^2}{\omega} \sum_{l,m} (-1)^l (2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m}), \tag{A17}$$

with the renormalized nearest-neighbor hopping amplitudes $J'_x = \frac{J_x\kappa}{\sqrt{2}\omega} = J_y$ and $J'_y = J_y(1 - \frac{1}{2}\frac{\kappa^2}{\omega^2})$, and a next-nearest-neighbor along $\mathbf{e}_x$ hopping term proportional to $\frac{J_x^2}{\omega}$ in (A17).

APPENDIX B: EFFECTIVE HAMILTONIAN IN MOMENTUM SPACE

If we choose the unit cell as in Fig. 1(a) [lattice sites $\mathtt{A} = (1,0)$, $\mathtt{B} = (2,0)$, $\mathtt{C} = (3,0)$, and $\mathtt{D} = (4,0)$], the momentum-space representation of the effective Hamiltonian without correction $\hat{H}_{\text{eff},0}$ (6) is given by a 4×4 matrix

$$\hat{\mathcal{H}}_{\text{eff},0}(\mathbf{k}) = \begin{pmatrix} 0 & J'_x e^{-i\frac{3\pi}{4}} - J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_2} & 0 & J'_x e^{-i\frac{3\pi}{4} - i\mathbf{k}\cdot\mathbf{R}_1} - J'_y e^{i\mathbf{k}\cdot(\mathbf{R}_2 - \mathbf{R}_1)} \\ J'_x e^{i\frac{3\pi}{4}} - J'_y e^{i\mathbf{k}\cdot\mathbf{R}_2} & 0 & J'_x e^{-i\frac{\pi}{4}} - J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_2} & 0 \\ 0 & J'_x e^{i\frac{\pi}{4}} - J'_y e^{i\mathbf{k}\cdot\mathbf{R}_2} & 0 & J'_x e^{i\frac{\pi}{4}} - J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_2} \\ J'_x e^{i\frac{3\pi}{4} + i\mathbf{k}\cdot\mathbf{R}_1} - J'_y e^{i\mathbf{k}\cdot(\mathbf{R}_1 - \mathbf{R}_2)} & 0 & J'_x e^{-i\frac{\pi}{4}} - J'_y e^{i\mathbf{k}\cdot\mathbf{R}_2} & 0 \end{pmatrix}, \tag{B1}$$

where $\mathbf{R}_1$ and $\mathbf{R}_2$ are the lattice vectors $\mathbf{R}_1 = (4,0)$ and $\mathbf{R}_2 = (1,1)$ and $\mathbf{k}$ is in the first Brillouin zone, which is given by the reciprocal lattice vectors $\mathbf{b}_1 = \frac{\pi}{2}(1,-1)$ and $\mathbf{b}_2 = 2\pi(0,1)$.

When the $\frac{J_x^2}{\omega}$ correction is included in the effective Hamiltonian, $\hat{H}_{\text{eff},1}$ (7), the unit cell is doubled, see Fig. 1(b), and the first Brillouin zone is therefore halved. If we now choose the lattice sites $\mathtt{a} = (1,0)$, $\mathtt{B} = (2,0)$, $\mathtt{c} = (3,0)$, $\mathtt{D} = (4,0)$, $\mathtt{A} = (2,1)$, $\mathtt{b} = (3,1)$, $\mathtt{C} = (4,1)$, and $\mathtt{d} = (5,1)$ for the unit cell, the momentum-space representation of the effective Hamiltonian will be an 8×8 matrix

$$\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k}) = \begin{pmatrix} -\frac{2J_x^2}{\omega} & J'_x e^{-i\frac{3\pi}{4}} & -\frac{J_x^2}{\omega}(1 + e^{i\mathbf{k}\cdot\mathbf{R}_1}) & J'_x e^{-i(\frac{3\pi}{4} - \mathbf{k}\cdot\mathbf{R}_1)} & 0 & -J'_y e^{i\mathbf{k}\cdot\mathbf{R}_2} & 0 & -J'_y e^{i\mathbf{k}\cdot\mathbf{R}_1} \\ J'_x e^{i\frac{3\pi}{4}} & \frac{2J_x^2}{\omega} & J'_x e^{-i\frac{\pi}{4}} & \frac{J_x^2}{\omega}(1 + e^{i\mathbf{k}\cdot\mathbf{R}_1}) & -J'_y & 0 & -J'_y e^{i\mathbf{k}\cdot\mathbf{R}_2} & 0 \\ -\frac{J_x^2}{\omega}(1 + e^{-i\mathbf{k}\cdot\mathbf{R}_1}) & J'_x e^{i\frac{\pi}{4}} & -\frac{2J_x^2}{\omega} & J'_x e^{i\frac{\pi}{4}} & 0 & -J'_y & 0 & -J'_y e^{i\mathbf{k}\cdot\mathbf{R}_2} \\ J'_x e^{i(\frac{3\pi}{4} - \mathbf{k}\cdot\mathbf{R}_1)} & \frac{J_x^2}{\omega}(1 + e^{-i\mathbf{k}\cdot\mathbf{R}_1}) & J'_x e^{-i\frac{\pi}{4}} & \frac{2J_x^2}{\omega} & -J'_y e^{-i\mathbf{k}\cdot(\mathbf{R}_1 - \mathbf{R}_2)} & 0 & -J'_y & 0 \\ 0 & 0 & 0 & -J'_y e^{i\mathbf{k}\cdot(\mathbf{R}_1 - \mathbf{R}_2)} & \frac{2J_x^2}{\omega} & J'_x e^{-i\frac{3\pi}{4}} & \frac{J_x^2}{\omega}(1 + e^{i\mathbf{k}\cdot\mathbf{R}_1}) & J'_x e^{-i(\frac{3\pi}{4} - \mathbf{k}\cdot\mathbf{R}_1)} \\ -J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_2} & 0 & -J'_y & 0 & J'_x e^{i\frac{3\pi}{4}} & -\frac{2J_x^2}{\omega} & J'_x e^{-i\frac{\pi}{4}} & -\frac{J_x^2}{\omega}(1 + e^{i\mathbf{k}\cdot\mathbf{R}_1}) \\ 0 & -J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_2} & 0 & -J'_y & \frac{J_x^2}{\omega}(1 + e^{-i\mathbf{k}\cdot\mathbf{R}_1}) & J'_x e^{i\frac{\pi}{4}} & \frac{2J_x^2}{\omega} & J'_x e^{i\frac{\pi}{4}} \\ -J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_1} & 0 & -J'_y e^{-i\mathbf{k}\cdot\mathbf{R}_2} & 0 & J'_x e^{i(\frac{3\pi}{4} - \mathbf{k}\cdot\mathbf{R}_1)} & -\frac{J_x^2}{\omega}(1 + e^{-i\mathbf{k}\cdot\mathbf{R}_1}) & J'_x e^{-i\frac{\pi}{4}} & -\frac{2J_x^2}{\omega} \end{pmatrix}, \tag{B2}$$

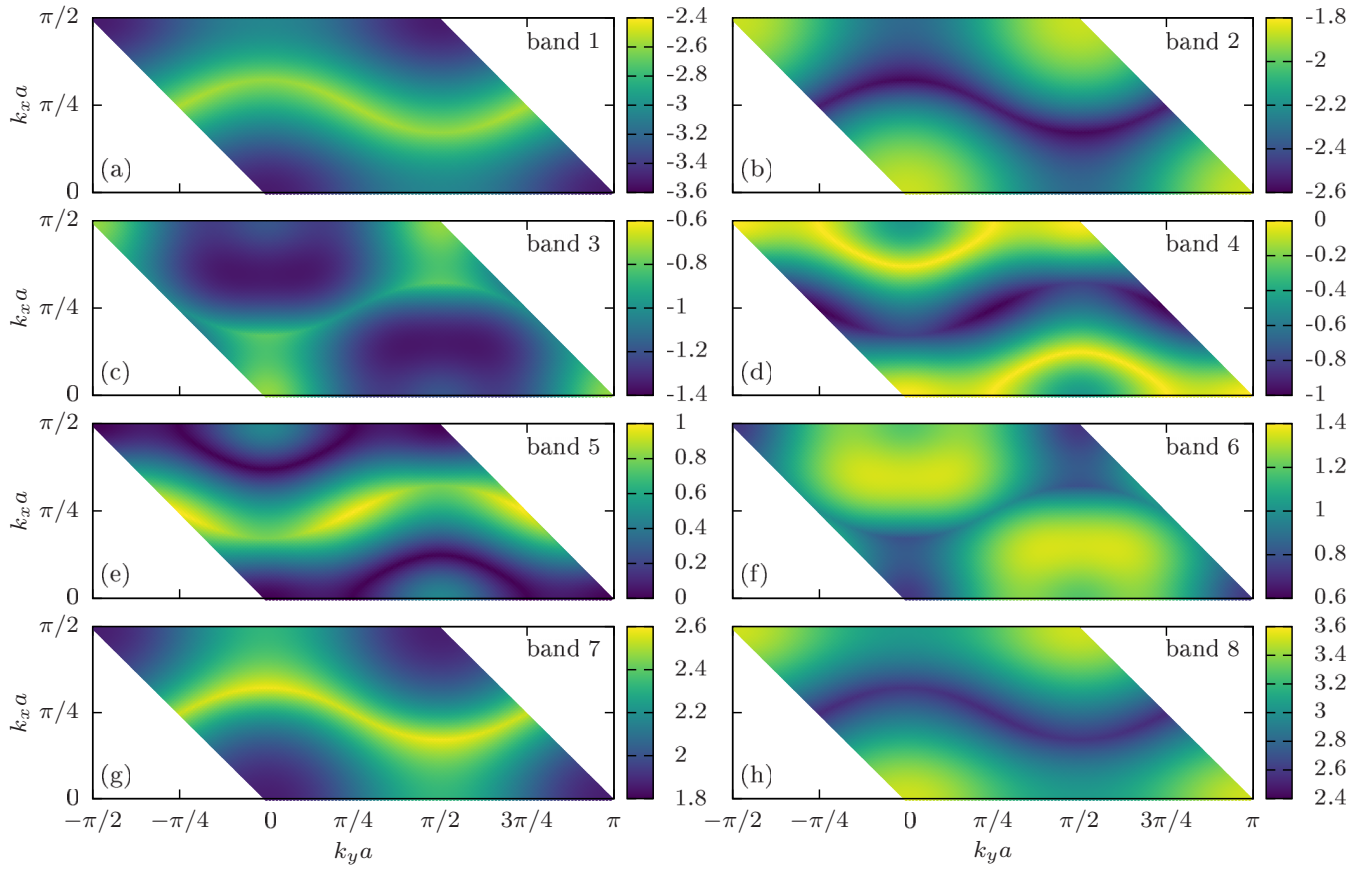

FIG. 12. Eight energy subbands of $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$ for the driving frequency $\omega = 20$. Subbands 1 and 2 form the lowest band with Chern number $c_1 = 1$, subbands 3, 4, 5, and 6 form the middle band with $c_2 = -2$, and subbands 7 and 8 form the highest band with $c_3 = 1$.

with the lattice vectors $\mathbf{R}_1 = (4, 0)$ and $\mathbf{R}_2 = (2, 2)$. The reciprocal lattice vectors are then $\mathbf{b}_1 = \frac{\pi}{2}(1, -1)$ and $\mathbf{b}_2 = \pi(0, 1)$.

The energy bands of $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$ are shown in Figs. 2 and 12.

APPENDIX C: DESCRIPTION OF INCOHERENT BOSONS

In a typical condensed-matter system constituent particles are electrons. Due to their fermionic statistics, at low enough temperatures, and with Fermi energy above the lowest band, that band of the topological model is uniformly occupied, and consequently the transverse Hall conductivity can be expressed in terms of the Chern number (8) [10]. In contrast, weakly interacting bosons in equilibrium form a Bose-Einstein condensate in the band minima and only probe the local Berry curvature [21].

Yet in the experiment [26] the Chern number was successfully measured using bosonic atoms of ^{87}Rb. This was possible because in the process of ramping up the drive (4), the initial Bose-Einstein condensate was transferred into an incoherent bosonic mixture. Conveniently, it turned out that the bosonic distribution over the states of the lowest band of the effective Floquet Hamiltonian was nearly uniform. Motivated by the experimental procedure, we model the initial bosonic state by a statistical matrix

$$\rho(t = 0) = \prod_{k=1}^{N_m} |k, N_p\rangle\langle k, N_p|, \tag{C1}$$

where the states $|k\rangle = a_k^\dagger|0\rangle$ approximately correspond to the lowest-band eigenstates of $\hat{H}_{\text{eff}}$ and each of these N_m states is occupied by N_p atoms $|k, N_p\rangle = \mathcal{N}(a_k^\dagger)^{N_p}|0\rangle$.

A procedure for selecting the states $|k\rangle$ is described in Refs. [22,26]. In order to probe the Chern number of the lowest band, the states $|k\rangle$ should correspond closely to the lowest-band eigenstates of $\hat{H}_{\text{eff}}$. At the same time, in the experiment in the initial moment the atomic cloud is spatially localized. According to Refs. [22,26] the optimal approach is to consider a steep confining potential and to use the low-lying eigenstates of

$$\hat{H}_{\text{initial}} = \hat{h}^{\text{eff}} + \left(\frac{r}{r_0}\right)^{\zeta}, \tag{C2}$$

where in our calculations $\hat{h}^{\text{eff}}$ is either $\hat{H}_{\text{eff},0}$ from Eq. (6) or $\hat{H}_{\text{eff},1}$ from Eq. (7) and the parameters of the confining potential are set to $r_0 = 20$, $\zeta = 20$.

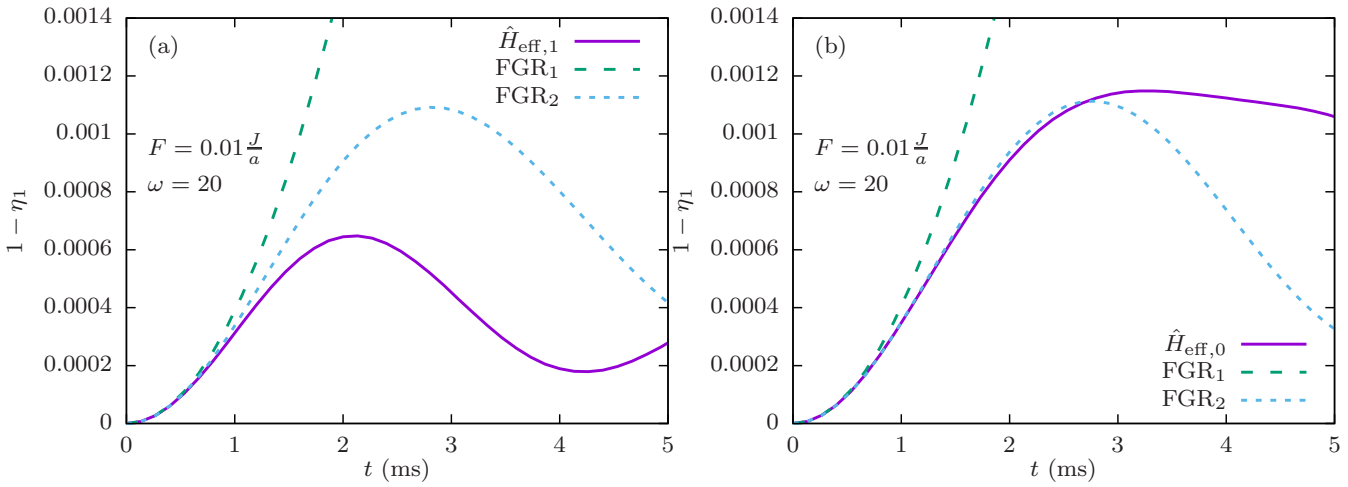

FIG. 13. Population in higher bands, comparison of numerical results (solid line) with the Fermi's golden rule in the first and second approximation (dashed lines). Band populations are calculated for an initial BEC in an eigenstate of the effective Hamiltonian and then averaged over (approximately) all states in the first band. (a) Initial state and evolution from the effective Hamiltonian with correction $\hat{H}_{\text{eff},1}$, Eq. (7). (b) Without the correction, $\hat{H}_{\text{eff},0}$, Eq. (6).

The dynamics of the initial state (C1) is induced by a double quench: at $t_0 = 0$ the atomic cloud is released from the confining potential and exposed to a uniform force of intensity F along the y direction. During the whole procedure the driving providing the laser-assisted tunneling, defined in Eq. (4), is running.

The main observables of interest are the center-of-mass position along x direction

$$x(t) = \left\langle \sum_{l,m} l |\psi_{l,m}(t)|^2 \right\rangle \tag{C3}$$

and the population of the ith band of the effective model

$$\eta_i(t) = \left\langle \sum_{|k\rangle \in i\text{th band}} \left| \sum_{l,m} \alpha_{lm}^{k*} \psi_{lm}(t) \right|^2 \right\rangle, \tag{C4}$$

where the states $|k\rangle = \sum_{l,m} \alpha_{lm}^k |l,m\rangle$ correspond to the eigenstates of the effective model. Here, angle brackets $\langle\,\rangle$ denote averaging over N_{samples} sets of initial conditions.

In the case of noninteracting particles, these and other quantities can be numerically accessed by solving the single-particle time-dependent Schrödinger equation for N_m different initial states $|k\rangle$. This is equivalent to sampling the initial state according to Eq. (9).

In the end, we give two technical remarks. First, all our calculations are done in the rotating frame; see Eq. (A1) in Appendix A. The staggered potential (2) is removed in this way. Second, in the case when the evolution is governed by the time-dependent Hamiltonian (10), the initial state is multiplied by the operator $e^{-i\hat{K}(0)}$ in order to properly compare these results to the ones obtained from the evolution governed by the effective Hamiltonian (11); see Eq. (5).

APPENDIX D: INITIAL QUADRATIC REGIME

For simplicity, we will consider only the case without the confining potential and with very weak force $F = 0.01$. The initial state is a Bose-Einstein condensate in one of the eigenstates of the effective Hamiltonian. The results are later averaged over all first band eigenstates.

Fermi's golden rule predicts that the probability for transition from an initial state ψ_i to a final state ψ_f, induced by a perturbation $\Delta\hat{H}$, is proportional to the square of matrix elements $|\langle\psi_i|\Delta\hat{H}|\psi_f\rangle|^2$. In this case, the perturbation is $\Delta\hat{H} = F\hat{y}$. If we assume that the probability of a particle being in the initial state is always $P_i(t) = |\psi_i(t)|^2 \approx 1$, Fermi's golden rule predicts [61]

$$P_{i\to f}^{\text{FGR}_1}(t) = \frac{1}{\hbar^2} |\langle\psi_i|\Delta\hat{H}|\psi_f\rangle|^2 t^2. \tag{D1}$$

If we now also consider transitions from the other states to the initial state, but keep the assumption that the populations in other states are small $P_{j\neq i}(t) = |\psi_{j\neq i}(t)|^2 \ll 1$, the time-dependent perturbation theory then predicts [61]

$$P_{i\to f}^{\text{FGR}_2}(t) = |\langle i|\Delta\hat{H}|f\rangle|^2 \frac{1 - 2\,e^{-\frac{\Gamma}{2\hbar}t}\cos\left(\frac{E_f - E_i}{\hbar}t\right) + e^{-\frac{\Gamma}{\hbar}t}}{(E_f - E_i)^2 + \frac{\Gamma^2}{4}}, \tag{D2}$$

where $\Gamma = \frac{2\pi}{\hbar}|\langle i|\Delta\hat{H}|f\rangle|^2$ and E_i (E_f) is the energy of the initial (final) state.

We plot the numerical results and both theoretical predictions from Fermi's golden rule in Fig. 13. Here we can see that all three curves agree well for short times, the second approximation longer remains close to the numerical results, and that the initial quadratic regime is reproduced by theory. This is the so-called quantum Zeno regime [57].

APPENDIX E: ENERGY

Time evolution of kinetic and interaction energy per particle for different interaction strengths is plotted in Fig. 14. Here we define the kinetic energy per particle as the expectation value of the time-dependent Hamiltonian (A1) divided by the total number of particles $E_{\text{kin}}(t) = \frac{1}{N}\langle\sum_{l,m,i,j} \psi_{l,m}^*(t)$

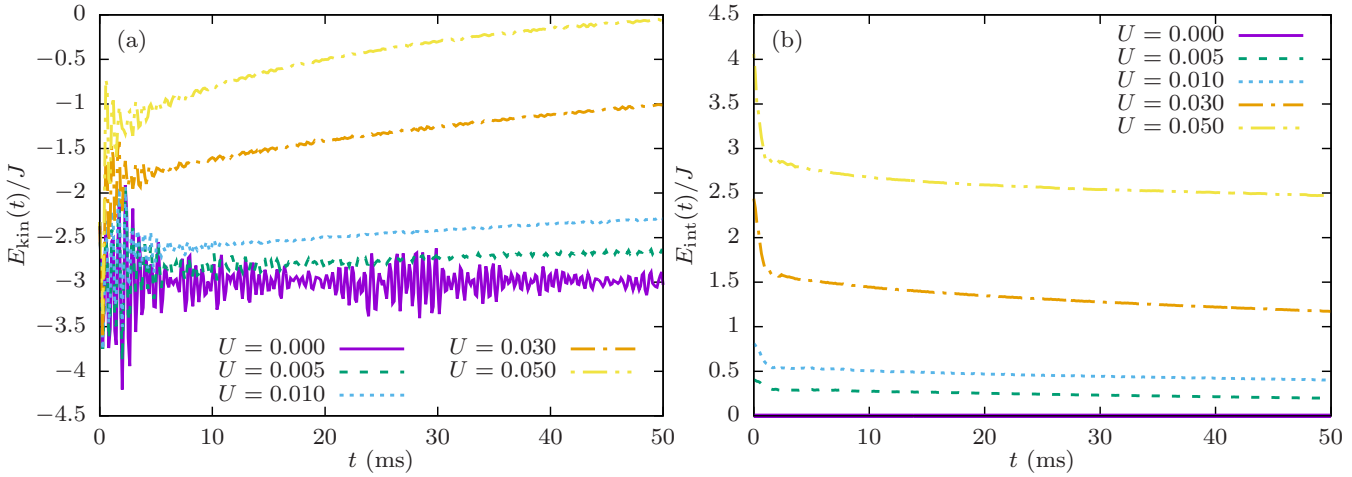

FIG. 14. (a) Kinetic energy per particle (expectation value of the time-dependent Hamiltonian $E_{\rm kin}(t) = \frac{1}{N}\langle \sum_{l,m,i,j} \psi^*_{l,m}(t) H_{lm,ij}(t)\psi_{i,j}(t)\rangle$ divided by the total number of particles N) for several different interaction strengths. (b) Interaction energy per particle $E_{\rm int}(t) = \frac{1}{N}\frac{U}{2}\langle \sum_{l,m} |\psi_{l,m}(t)|^2[|\psi_{l,m}(t)|^2 - 1]\rangle$. U is given in units where $J = 1$.

$H_{lm,ij}(t)\psi_{i,j}(t)\rangle$, while the interaction energy per particle is $E_{\rm int}(t) = \frac{1}{N}\frac{U}{2}\langle \sum_{l,m} |\psi_{l,m}(t)|^2[|\psi_{l,m}(t)|^2 - 1]\rangle$. Both energies grow with increasing interaction coefficient U.

When the interactions are strong enough and after long enough time, the atoms become equally distributed between the eigenstates of the Hamiltonian $\hat{H}(t)$. As the energy spectrum of $\hat{H}(t)$ is symmetric around zero, the expectation value of $\hat{H}(t)$ (kinetic energy) should be zero when all bands are equally populated. We can see this in Fig. 14(a), where the kinetic energy approaches zero at $t \approx 50$ ms for the case $U = 0.05$.

The interaction energy at first rapidly decreases, as the cloud rapidly expands after turning off the confinement potential $\hat{V}_{\rm conf}$, and after that continues to slowly decrease as the cloud slowly expands; see Fig. 14(b).

These considerations also provide a possibility to discuss the applicability of the approximative method introduced in Sec. IV. As we work in the regime of high frequency $\omega = 20$, we find that for weak interaction, at short enough times of propagation, the energy is approximately conserved. At stronger values of $U \geqslant 0.01$ we observe a slow increase in the total energy on the considered time scales. In both cases we do not find the onset of parametric instabilities [31]. If present, these instabilities are signaled by an order of magnitude increase in energy on a short time scale, that we do not find.

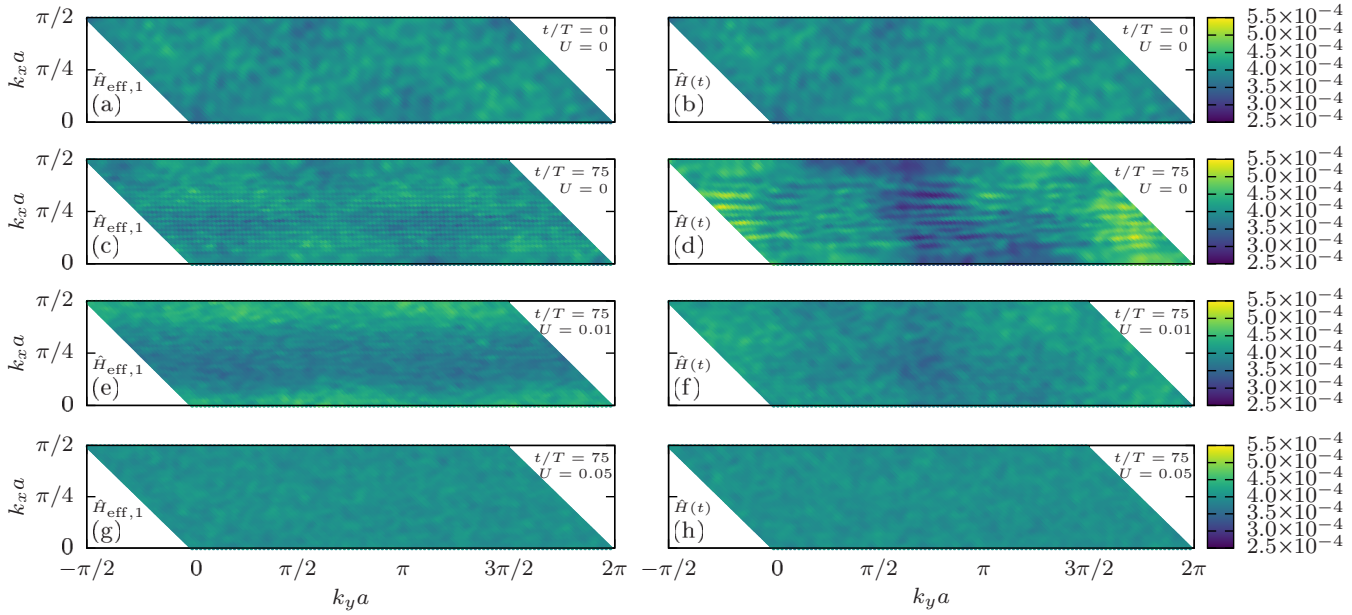

FIG. 15. Momentum-space density distribution in all bands, $\eta_1(\mathbf{k}) + \eta_2(\mathbf{k}) + \eta_3(\mathbf{k})$. U is given in units where $J = 1$. Left: evolution using the time-dependent Hamiltonian $\hat{H}_{\rm eff,1}$. Right: evolution using the time-dependent Hamiltonian $\hat{H}(t)$. (a), (b) Initial state. (c), (d) Final state after 50 ms (75 driving periods), noninteracting case $U = 0$. (e), (f) $U = 0.01$. (g), (h) $U = 0.05$.

In addition, the two-body interaction can deplete the occupancies of initial coherent modes [29,41] and limit the validity of our approach. In principle, these types of processes can be addressed by including quantum fluctuations along the lines of the full truncated Wigner approach [45]. Yet, we set our parameters in such a way that these additional contributions are small.

APPENDIX F: MOMENTUM-SPACE DENSITY DISTRIBUTION

The momentum-space probability densities at the initial moment and after 75 driving periods (50 ms) are shown in Fig. 15. The interactions deplete the lowest band, but also smooth out the density distribution.

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Bosonic fractional quantum Hall states in driven optical lattices

Ana Hudomal,[1] Nicolas Regnault,[2,3] and Ivana Vasić[1]

[1]*Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Pregrevica 118, 11080 Belgrade, Serbia*

[2]*Joseph Henry Laboratories and Department of Physics, Princeton University, Princeton, New Jersey 08544, USA*

[3]*Laboratoire de Physique de l'École Normale Supérieure, ENS, Université PSL, CNRS, Sorbonne Université, Université Paris-Diderot, Sorbonne Paris Cité, 75005 Paris, France*

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Strong synthetic magnetic fields have been successfully implemented in periodically driven optical lattices. However, the interplay of the driving and interactions introduces detrimental heating, and for this reason it is still challenging to reach a fractional quantum Hall state in cold-atom setup. By performing a numerical study, we investigate stability of a bosonic Laughlin state in a small atomic sample exposed to driving. We identify an optimal regime of microscopic parameters, in particular interaction strength U and the driving frequency ω, such that the stroboscopic dynamics supports the basic $\nu = 1/2$ Laughlin state. Moreover, we explore slow ramping of a driving term and show that the considered protocol allows for the preparation of the Laughlin state on experimentally realistic time-scales.

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I. INTRODUCTION

Cold atoms in optical lattices provide a highly tunable platform for quantum simulations of relevant many-body Hamiltonians [1,2]. Since early experiments with quantum gases, there has been a strong interest in the realization of fractional quantum Hall (FQH) states in these setups [3–17]. Despite numerous experimental achievements and a variety of theoretical proposals, FQH physics has still not been reached in cold-atom experiments.

A milestone in the field has been recently achieved by the realization of artificial gauge potentials [18–28]. In particular, the topological index of a resulting energy band of an optical lattice featuring a strong synthetic magnetic field has been directly probed [22]. At first glance, both key requirements for the emergence of FQH states—atomic interactions and strong synthetic magnetic fields—are now experimentally available. However, there are several specific details in the implementation of strong synthetic magnetic fields for cold atoms that make the realization of FQH states still challenging.

The most advanced recent realizations of artificial gauge potentials exploit periodically driven optical lattices [19–28]. Using Floquet theory, the stroboscopic dynamics of a noninteracting driven system can be related to an effective time-independent Hamiltonian [29–32]. This approach, Floquet engineering, enriches the set of quantum models that can be simulated in cold-atom experiments. However, general arguments and numerical studies [33–35] suggest that the interplay of interactions and driving in a thermodynamically large system introduces heating, leading to a featureless infinite-temperature state in the long-time limit.

Although this general result might sound discouraging, the heating process can be very slow in some driven systems for specific regime of microscopic parameters. There, the system can be described by a physically interesting "prethermal" Floquet state on experimentally relevant time-scales [36–42]. Moreover, the onset of thermalization in a finite-size interacting system may exhibit unexpected features, not found in the thermodynamic limit [43,44]. Heating rates and resulting instabilities have been recently investigated both theoretically and experimentally for the driven Bose-Hubbard model in the weakly interacting regime [38,45–47]. Moreover, experimental studies of the driven Fermi-Hubbard model in a honeycomb lattice have established a timescale of the order of 100 tunneling times for the regime where the effective-model description applies [48,49].

In this paper, we consider small systems of several interacting bosonic atoms in a periodically driven optical lattice featuring synthetic magnetic flux. The focus of our study is on finding optimal microscopic parameters that would allow to prepare and probe the basic bosonic Laughlin state in this setup. To this end, we employ exact numerical simulations of the driven Bose-Hubbard model [50] for small system sizes.

From one point of view, it is expected that a small driven system exhibits low heating rates for a driving frequency set above a finite bandwidth of an effective model [33]. However, driving a system with such a high frequency may lead to undesirable effects, such as coupling of the lowest band to higher bands of the underlying optical lattice, thus making the initial description based on the lowest-band Hubbard model inapplicable. These effects have been addressed in a recent study [51] where an optimal intermediate frequency window for Floquet engineering has been established.

In our study, we go a step further in the search for the optimal regime that might allow for the bosonic Laughlin states under driving. In particular, for a realistic, intermediate value of a driving frequency, the interaction term complicates the effective model by introducing several higher-order terms. Their effect on the topological states has been addressed only recently [52,53] and it has been found that typically these

terms work against the topological state. For this reason, the stability of the Laughlin state at intermediate driving frequency requires a separate study, that we perform here. Moreover, we numerically investigate an experimentally relevant preparation protocol for the Laughlin state in a driven system [54]. For a reference, we note that a simpler but closely related question concerning the static (undriven systems) has gained lot of attention [6,7,15,55].

The paper is organized as follows: in Sec. II we introduce the model under study and briefly review key features of the particle-entanglement spectra that we will exploit in the identification of the Laughlin-like state. Then, in Sec. III A we investigate general heating effects of interacting bosons exposed to the driving. By extending this approach, in Sec. III B we construct the stroboscopic time-evolution operator and inspect its eigenstates in order to identify possible FQH states. Finally, in Sec. IV we address the possibility of accessing these states in an experiment through a slow ramp of the driving term.

II. MODEL AND METHOD

In this section we first introduce the driven model and explain the basis of Floquet engineering. Then we summarize several key features of the particle-entanglement spectra that we use to characterize the bosonic Laughlin states.

A. Driven model

Properties of bosonic atoms in a deep optical lattice can be realistically described within the framework of the Bose-Hubbard model [1]. We consider a basic driving scheme [50] that introduces a uniform, synthetic magnetic flux into a *square* optical lattice here spanned by the two vectors $\mathbf{e_x}$ and $\mathbf{e_y}$. The corresponding Hamiltonian is given by the driven Bose-Hubbard model

$$\begin{aligned}\hat{H}(t) = &-J_x \sum_{m,n} (\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \text{H.c.}) \\ &- J_y \sum_{m,n} (e^{i\omega t}\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} + \text{H.c.}) \\ &+ \frac{\kappa}{2}\sum_{m,n} \sin\left[\omega t - (m+n-1/2)\,\phi\right]\hat{n}_{m,n} \\ &+ \frac{U}{2}\sum_{m,n} \hat{n}_{m,n}(\hat{n}_{m,n}-1), \end{aligned} \tag{1}$$

where operators $\hat{a}_{m,n}$ ($\hat{a}^\dagger_{m,n}$) annihilate (create) a boson at lattice position (m,n), and local density operators are $\hat{n}_{m,n} = \hat{a}^\dagger_{m,n}\hat{a}_{m,n}$. J_x and J_y are tunneling amplitudes and U is the on-site local repulsive interaction. We use the units where $\hbar = 1$ and the lattice constant $a = 1$. The driving scheme is defined by the driving frequency ω, the driving amplitude κ and by a phase ϕ. In the following we set $\phi = \pi/2$ and $\kappa/\omega = 0.5$. These values were recently used in an experimental realization of the Harper-Hofstadter model [22]. The derivation of this model is briefly reviewed in Appendix. We assume periodic boundary conditions implemented using the vectors $\mathbf{R}_1 = 4\,\mathbf{e}_x$, $\mathbf{R}_2 = -\mathbf{e}_x + \mathbf{e}_y$, as presented in Fig. 1. This choice is compatible with the driving term and it allows us to exploit translational symmetry by working in the fixed quasimomentum basis.

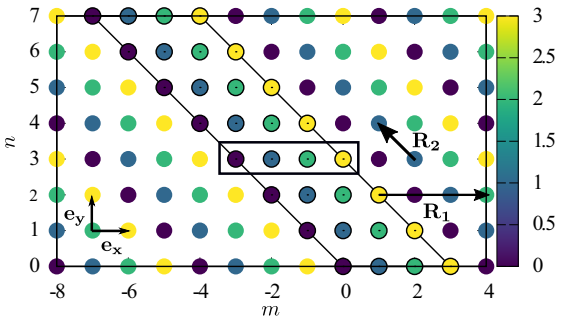

FIG. 1. Lattice geometry used throughout the paper. The parallelogram gives the exemplary lattice size $(L_x, L_y) = (4, 8)$. The color scale is defined by mod $(m+n, 4)$, in accordance with the driving term from Eq. (1). The vectors $\mathbf{R}_1 = 4\,\mathbf{e}_x$, $\mathbf{R}_2 = -\mathbf{e}_x + \mathbf{e}_y$ are used to implement periodic boundary conditions. The small rectangle gives the magnetic unit cell for the effective model in Eq. (3).

Formally, by using the Floquet theory [29,30,56], it can be shown that the full time-evolution operator corresponding to this model is given by

$$\hat{U}(t, t_0) = e^{-i\hat{K}(t)} e^{-i(t-t_0)\hat{\mathcal{H}}_{\text{eff}}} e^{i\hat{K}(t_0)}, \tag{2}$$

where $\hat{K}(t)$ is a periodic "kick" operator $\hat{K}(t) = \hat{K}(t + 2\pi/\omega)$ and $\hat{\mathcal{H}}_{\text{eff}}$ is a time-independent effective Hamiltonian. The full-time evolution operator is periodic as well and consequently the (quasi)eigenenergies of $\hat{\mathcal{H}}_{\text{eff}}$ are defined up to modulo ω. The last equation gives formal mapping of a periodically driven system to an effective model that captures the stroboscopic time evolution of the model.

In the noninteracting regime, $U = 0$, there are several well controlled approximations to obtain the effective Hamiltonian. These techniques are the essence of Floquet engineering, an approach where the driving scheme is implemented in such a way to yield a sought-after effective model. However, according to general analytical arguments and numerical insights, the corresponding effective model of a driven interacting many-body system in the thermodynamic limit exhibits nonphysical features [33,34]. In particular, the system thermalizes and in the long-time limit its steady state is a featureless, infinite-temperature state, independent of the initial state.

Here we consider small samples of several bosonic atoms. Due to a finite spectrum bandwidth, we expect the high-frequency expansion to be relevant for a finite range of the driving frequency. Within these assumptions, the leading-order (in $1/\omega$) effective Hamiltonian is

$$\begin{aligned}\hat{H}_{\text{eff}} = &-J_x \sum_{m,n} (\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \text{H.c.}) \\ &- J'_y \sum_{m,n} (e^{i(m+n)\phi}\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} + \text{H.c.}) \\ &+ \frac{U}{2}\sum_{m,n} \hat{n}_{m,n}(\hat{n}_{m,n}-1). \end{aligned} \tag{3}$$

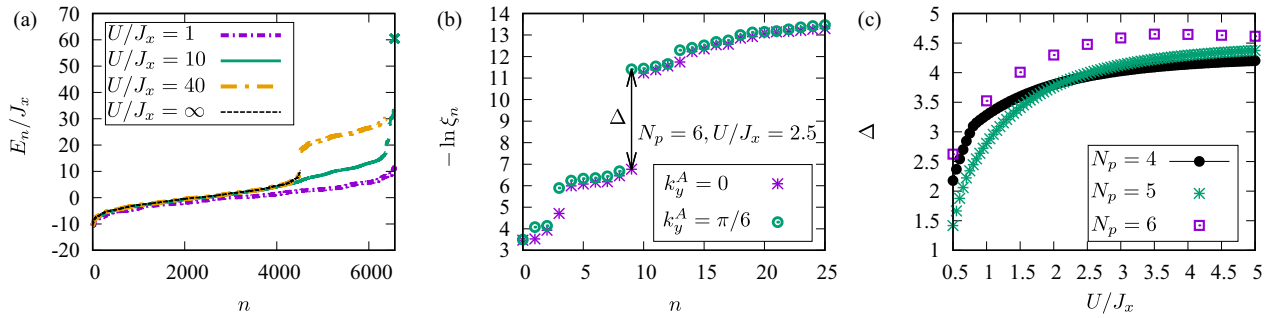

FIG. 2. (a) The energy spectrum E_n of the model from Eq. (3) in the $k_x = 0$, $k_y = 0$ sector for $N_p = 4$ and different values of interaction $U/J_x = 1, 10, 40$ and $U/J_x = \infty$ (hard-core bosons). The top part of the spectrum is at $\approx (U/J_x)N_p(N_p - 1)/2$. (Not shown for $U/J_x = 40$.) For a high ratio U/J_x the spectrum splits into bands. The lowest band corresponds to hard-core bosons. (b) The low-lying part of the particle-entanglement spectrum $-\ln \xi_n$ of the ground-state incoherent superposition, Eq. (6), in the region A momentum sectors $k_y^A = 0$ and $k_y^A = \pi/6$, and for $N_p = 6$, $U/J_x = 2.5$. (c) The particle-entanglement gap Δ of the incoherent superposition Eq. (6) as a function of interaction strength U for $N_p = 4, 5, 6$.

The Hamiltonian (3) features complex hopping phases $e^{i(m+n)\phi}$ that result in a uniform synthetic magnetic flux ϕ per lattice plaquette. Due to the driving, the renormalized hopping amplitude along the y direction turns into

$$J_y' \equiv \frac{\kappa}{2\omega} \sin(\phi/2) J_y. \tag{4}$$

For the values $\phi = 2\pi\alpha$, where the flux density α is set to $\alpha = 1/4$, and $\kappa/\omega = 0.5$, the tunneling amplitude along y direction in the effective model is $J_y' \approx J_y \times 0.1768$.

In a certain regime of microscopic parameters, the ground state of the model defined in Eq. (3) is given by the lattice version of the Laughlin state [7,9,57–59]. The Laughlin state is stabilized for the filling factor $\nu = N_p/N_\phi = 1/2$, where $N_\phi = \alpha L_x \times L_y$ is the total number of fluxes (N_ϕ being an integer) and N_p is the number of bosons, and for a strong-enough repulsion U. Another important requirement for the Laughlin state is to avoid the strong hopping anisotropy and to keep $J_x \approx J_y'$, so we set $J_x = 0.2 J_y$. We consider system sizes $N_p = 4, 5, 6$ and the respective lattices sizes $(L_x, L_y) = (4, 8)$, $(4, 10)$, and $(4, 12)$, see Fig. 1, where we expect the ground state to correspond to the $\nu = 1/2$ Laughlin state. The Hilbert space sizes for $k_x = k_y = 0$ are $\dim \mathcal{H} = 6564$, 108 604, and 1 913 364 respectively. For this choice of microscopic parameters, the model ground state of Eq. (3) is approximately twofold degenerate. The two ground-states are found in the sectors $k_x = 0$, $k_y = 0$ and $k_x = 0$, $k_y = \pi$. We denote them by $|\psi_{\mathrm{LGH}}^{0,0}\rangle$ and $|\psi_{\mathrm{LGH}}^{0,\pi}\rangle$.

As we are mainly interested in the driven regime, not only the ground state, but the full spectrum of the model from Eq. (3) plays a role. A rough argument is that the system does not absorb energy provided that the driving frequency ω is set above the bandwidth of the effective model. Several spectra of the model from Eq. (3) for $k_x = 0$, $k_y = 0$ are presented in Fig. 2(a). It can be seen that the ground-state energy is weakly affected by the value of $U \geqslant J_x$, while the top part of the spectrum with few states is found at $UN_p(N_p - 1)/2$. For higher values of U the spectrum splits into bands where the lowest band corresponds to the hard-core bosons and higher bands include double and higher occupancies.

B. Particle-entanglement spectra

There are several ways to characterize the ground states of the model from Eq. (3) as the Laughlin states. Usually, the starting point in this direction is the identification of the twofold degeneracy expected in the implemented torus geometry for $\nu = 1/2$. Another relevant quantity is the overlap of the numerically obtained state with the Laughlin analytical wave function in the torus geometry [9,59]. More direct evidence can be obtained through the calculation of the relevant topological index (Chern number) or the quantized Hall conductance. An additional convincing approach, that we pursue here, is based on the analysis of the entanglement spectra of the relevant states.

In the following we will use the particle-entanglement spectrum (PES) [59,60] to distinguish possible topologically nontrivial states. In order to obtain this type of entanglement spectrum, we partition N_p particles into two sets of N_A and $N_B = N_p - N_A$ particles. For a given mixed state ρ, we construct a reduced density matrix $\rho_A = \mathrm{tr}_B \rho$ by performing a partial trace over N_B particles. The resulting PES is given by $-\ln \xi_n$, where ξ_n are eigenvalues of ρ_A. The related particle-entanglement entropy is given by [61,62]

$$S_A = -\mathrm{tr}(\rho_A \ln \rho_A). \tag{5}$$

By partitioning particles, we keep the geometry of the system unchanged. For this reason, we will inspect the PES for the different momentum sectors k_y^A of the remaining N_A particles. An example of a PES is presented in Fig. 2(b). As proposed in Refs. [59,60], we have considered the incoherent superposition of the almost twofold degenerate ground state of Eq. (3) as the density matrix

$$\rho_{\mathrm{GS}} = \tfrac{1}{2}\left(|\psi_{\mathrm{LGH}}^{0,0}\rangle\langle\psi_{\mathrm{LGH}}^{0,0}| + |\psi_{\mathrm{LGH}}^{0,\pi}\rangle\langle\psi_{\mathrm{LGH}}^{0,\pi}|\right). \tag{6}$$

For simplicity, we only present the PES for the two momenta $k_y^A = 0$ and $k_y^A = \pi/6$. We observe a clear particle-entanglement gap Δ. In addition, the counting of low-lying modes below this gap (ten modes for $k_y^A = 0$ and nine modes for $k_y^A = \pi/6$, at $N_A = 3$, $N_p = 6$) corresponds to the Laughlin state [59,60]. In this way the PES encodes topological features of the state ρ in the form of well defined number of excitations per momentum sector k_y^A [59,60]. This type of

TABLE I. Counting of modes $\mathcal{N}_L(k_y^A)$ in the PES of the Laughlin state for several system sizes and particle partitions. The last column lists the $\mathcal{N}_L(k_y^A)$ values for each momentum sector $k_y^A = 2\pi i/L_y$, $i = 0, \ldots, L_y - 1$.

N_p	(L_x, L_y)	N_A	PES: $\mathcal{N}_L(k_y^A)$
4	(4, 8)	2	3, 2, 3, 2, 3, 2, 3, 2
5	(4, 10)	2	4, 3, 4, 3, 4, 3, 4, 3, 4, 3
6	(4, 12)	3	10, 9, 9, 10, 9, 9, 10, 9, 9, 10, 9, 9

analysis is useful as it can identify topological features even without model states, as done for the case of fractional Chern insulators [63,64].

In the following we will consider specific particle partitions $N_A = 2$, $N_p = 4$; $N_A = 2$, $N_p = 5$; and $N_A = 3$, $N_p = 6$. For these cases the counting of excitations $\mathcal{N}_L(k_y^A)$ per momentum sector k_y^A is well established and given in Table I. In Fig. 2(c) we show the particle-entanglement gap of the mixtures, Eq. (6), obtained at different values of U. Numerical results for the obtained PES indicate that a reasonably large gap is found starting at $U \sim 0.5J_x$ and the characteristic features of the Laughlin state persist with a further increase in U. We note that at lower values of the flux density, $\alpha < 1/4$, the Laughlin state can be found at even lower values of the repulsion U [9,59].

By analyzing the effective model from Eq. (1), we have obtained a guidance for the regime of microscopic parameters and for the geometry of the small system that can give rise to Laughlin states. In the next sections our aim is to go beyond the effective model from Eq. (3) and to identify topological states supported by the full driven dynamics as captured by the model given in Eq. (1).

III. DRIVEN DYNAMICS

In this section we discuss the full driven dynamics as captured by the model given in Eq. (1).

A. Heating

First we address the onset of heating following the standard procedure discussed in Refs. [42,65]. The initial state of the system is prepared using the ground state of the effective model

$$|\psi(t=0)\rangle = e^{-i\hat{K}(t=0)}\left|\psi_{\mathrm{LGH}}^{0,0}\right\rangle \tag{7}$$

and we monitor the stroboscopic time-evolution $t = NT$, $T \equiv 2\pi/\omega$ governed by the full driven model defined in Eq. (1). In our numerical simulations, we approximate the micromotion operator $\hat{K}(t=0)$ using the leading-order high-frequency expansion; see Eq. (A12). The quantity of interest is the expectation value of the effective Hamiltonian (3):

$$\langle \hat{H}_{\mathrm{eff}}(t=NT)\rangle_K = \langle\psi(t)|e^{-i\hat{K}(t=0)}\hat{H}_{\mathrm{eff}}e^{i\hat{K}(t=0)}|\psi(t)\rangle. \tag{8}$$

We expect this quantity to reasonably correspond to the ground-state energy of the effective model E_0 in the regime of very high frequency. On the other hand, for a "low" driving frequency we expect the system to quickly reach the infinite-temperature $\beta \to 0$ regime defined by

$$\lim_{\beta\to 0}\langle \hat{H}_{\mathrm{eff}}\rangle = \frac{1}{\dim\mathcal{H}}\mathrm{tr}(\hat{H}_{\mathrm{eff}}). \tag{9}$$

For this reason we monitor the normalized total energy

$$Q(t=NT) = \frac{\langle \hat{H}_{\mathrm{eff}}(t=NT)\rangle_K - E_0}{\lim_{\beta\to 0}\langle \hat{H}_{\mathrm{eff}}\rangle - E_0} \tag{10}$$

and we present it in Fig. 3(a), for $U/J_x = 10$. In agreement with the known results [65], we find that the thermalization is quick for both a "high" driving frequency $\omega/J_x \geqslant 20$ and for a "low" driving frequency $\omega/J_x \leqslant 10$. For the intermediate values of ω, the heating process is slow [65] and the total energy exhibits a slow exponential growth captured by $Q(t = NT) \approx 1 - b\exp(-ct)$, $t \gg 1$. An example of this behavior is given for $\omega/J_x = 15$ in Fig. 3(a). The heating process can also be monitored through the particle-entanglement entropy S_A as a function of time. In Fig. 3(b) for $N_p = 5$ and low driving frequency we find that this quantity quickly saturates to its maximal value. Indeed, for a thermal state at infinite temperature, S_A is given by

$$S_A^{\max} \approx \ln\binom{L_xL_y + N^A - 1}{N^A}, \tag{11}$$

marked by the horizontal, dot-dashed line in Fig. 3(b). Except for the highest frequency considered ($\omega/J_x = 50$), we find that, in the process of heating, the particle-entanglement gap of the initial state quickly closes (not shown in the plots).

Here we briefly discuss finite-size effects by comparing numerical results for the normalized total energy for $N_p = 4$, $N_p = 5$, and $N_p = 6$. In line with the known results [33,34,38], the "high-frequency" regime with low heating rates moves toward higher ω as the system size increases. However, we find that the estimates obtained in this section ($\omega/J_x \geqslant 20$ for the high- and $\omega/J_x \leqslant 10$ for the low-frequency regime, for $U/J_x = 10$) apply to all the three sizes $N_p = 4, 5, 6$, at least for the time-scales that we consider. A comprehensive study of the leading finite-size effects in driven systems can be found in Refs. [33,42,65].

B. The stroboscopic time-evolution operator

In order to better understand the limitations of the effective model, here we time evolve all relevant basis states for a single driving period $T = 2\pi/\omega$ and construct the stroboscopic time-evolution operator

$$\hat{U}_F \equiv \hat{U}(t_0 + T,\ t_0 = 0) \tag{12}$$

such that $\hat{U}(NT + t_0) = \hat{U}_F^N$. In the next step, for a system size $N_p = 4$, $(L_x, L_y) = (4, 8)$ we fully diagonalize this operator and inspect its eigenstates $|n\rangle$. Following the described procedure, we obtain the long-time limit

$$\lim_{N\to\infty}\langle \hat{H}_{\mathrm{eff}}(NT)\rangle_K = \sum_n |\langle n|\psi(t=0)\rangle|^2\langle n|\hat{H}_{\mathrm{eff}}|n\rangle_K, \tag{13}$$

where we define

$$\langle n|\hat{H}_{\mathrm{eff}}|n\rangle_K = \langle n|e^{-i\hat{K}(t=0)}\hat{H}_{\mathrm{eff}}e^{i\hat{K}(t=0)}|n\rangle. \tag{14}$$

Results for $Q(t = NT)$ from Eq. (10) obtained in this way are summarized in Fig. 3(c), where we make a comparison

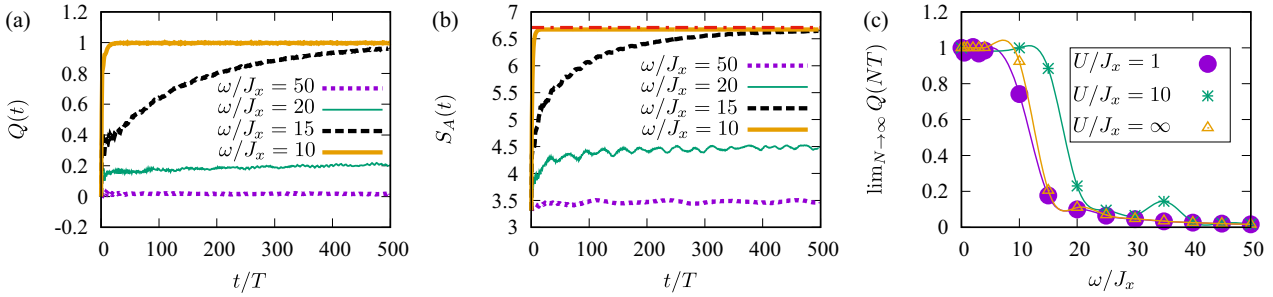

FIG. 3. (a) The normalized total energy $Q(t = NT)$ from Eq. (10), and the (b) particle-entanglement entropy $S_A(t = NT)$, Eq. (5), during the time evolution governed by Eq. (1) for several driving frequencies $\omega/J_x = 50,\ 20,\ 15,\ 10$. Parameters: $N_p = 5$, $U/J_x = 10$. Note that the asymptotic value of S_A for $\omega/J_x = 10$ and $\omega/J_x = 15$ matches the one given in Eq. (11), as presented by the horizontal, dot-dashed line. (c) The long-time limit $\lim_{N\to\infty} Q(NT)$ for $N_p = 4$ and the on-site interactions $U/J_x = 1$, 10 and $U/J_x = \infty$ (hard-core bosons). The lines are only guides to the eye.

between the long-time energies for the case of hard-core bosons ($U \to \infty$) and soft-core bosons (finite values of U). The obtained results indicate that heating rates of hard-core bosons are closer to the case of $U/J_x = 1$ in comparison to $U/J_x = 10$, which is expected from the bandwidths shown in Fig. 2(a). Overall we observe that the "high-frequency regime" is wider for lower ratios U/J_x.

In Fig. 4, we make a comparison between the exact driven model captured by $\hat{U}_F$ and $\hat{H}_{\text{eff}}$. In Figs. 4(a) and 4(b) we inspect the distribution of expectation values $\langle n|\hat{H}_{\text{eff}}|n\rangle_K$. By comparing these values to the eigenenergies of the effective model, Eq. (3), we get an insight into the pertinence of the effective description [33,34]. In particular, for an interacting system in the thermodynamic limit, the distribution is flat and the effective description is useless. We state again that we consider only small atomic samples. For this reason, it is expected that for high values of ω the full stroboscopic description nicely matches to the effective model values. Such an example is given in Fig. 4(a) for $U/J_x = 1$ and $\omega/J_x = 20$. As the value of ω gets lower the distribution becomes flatter, as can be seen in Fig. 4(b) for $U/J_x = 10$ by comparing results for $\omega/J_x = 50$ and $\omega/J_x = 10$.

The intermediate regime of frequencies, e.g., $\omega/J_x = 20$ for $U/J_x = 10$, is of the main experimental relevance [51]. We now investigate whether the driven stroboscopic dynamics supports some Laughlin-like states, by calculating the PES of the mixture

$$\rho_F = \tfrac{1}{2}(|n_0(0,0)\rangle\langle n_0(0,0)| + |n_0(0,\pi)\rangle\langle n_0(0,\pi)|), \quad (15)$$

where $|n_0(k_x, k_y)\rangle$ is the state from the k_x, k_y sector with the lowest expectation value $\langle n|\hat{H}_{\text{eff}}|n\rangle_K$. The results are presented in Fig. 4(c). We find that the states with a well defined gap and the Laughlin-like PES can be found down to $\omega/J_x \geqslant 10$ for $U/J_x = 1$, and down to $\omega/J_x \geqslant 20$ for $U/J_x = 10$. Having established existence of these states for small samples of $N_p = 4$ particles, in the next section we discuss dynamical protocol which can be exploited to prepare these states.

IV. SLOW RAMP

The question about an optimal adiabatic protocol that can be used to prepare the Laughlin state in a cold-atom setup has gained lot of attention [6,7,15,55]. The situation becomes even more complex once the full driving process is taken into account. A general wisdom is that, by starting from a topologically trivial state, the topological index of a thermodynamically large system cannot be changed adiabatically. We consider a small atomic sample and follow the proposal of Ref. [15]. Our main contribution is that we extend this protocol to the case of the driven, interacting Bose-Hubbard model.

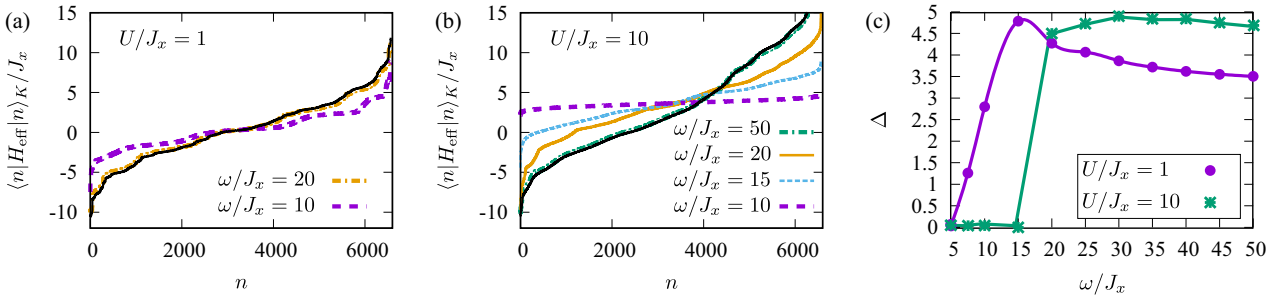

FIG. 4. Properties of the eigenstates $|n\rangle$ of the stroboscopic time-evolution operator $\hat{U}_F$, Eq. (12), in the $k_x = 0$, $k_y = 0$ sector for $N_p = 4$. Expectation values $\langle n|\hat{H}_{\text{eff}}|n\rangle_K$ defined in Eq. 14 for (a) $U/J_x = 1$, $\omega/J_x = 10$, 20 and (b) $U/J_x = 10$, $\omega/J_x = 10$, 15, 20, 50. The black solid lines mark eigenenergies of $\hat{H}_{\text{eff}}$, Eq. (3). Note that in (b) we do not include few states from the top of the spectrum of $\hat{H}_{\text{eff}}$, Eq. (3), for clarity. (c) The particle-entanglement gap Δ of the incoherent superposition ρ_F, Eq. (15), for $U/J_x = 1$ and $U/J_x = 10$, $N_p = 4$. The lines are only guides to the eye.

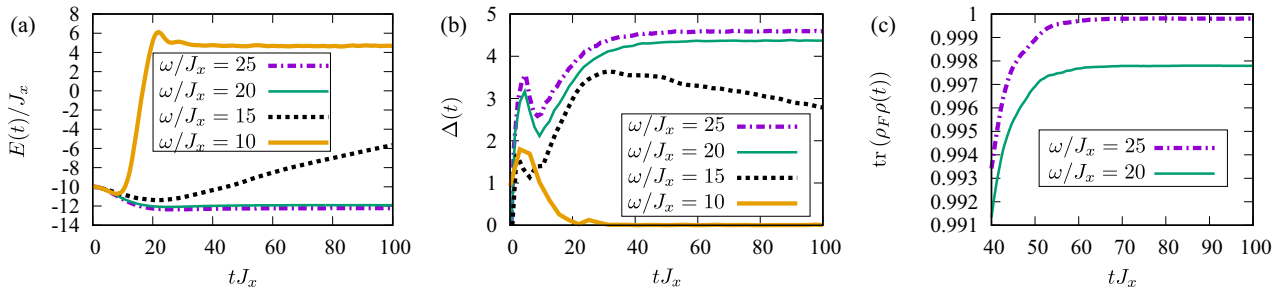

FIG. 5. (a) The expectation value $E(t)$ defined in Eq. (19) and (b) the particle-entanglement gap $\Delta(t)$ of $\rho(t)$, Eq. (18), during the time evolution governed by Eq. (17) for several driving frequencies $\omega/J_x = 25,\ 20,\ 15,\ 10$. Parameters: $N_p = 5,\ U/J_x = 10,\ \eta/J_x = 0.05$. (c) The overlap $\mathrm{tr}\,[\rho(t)\rho_F]$ of the time evolved state with the target eigenstates of $\hat{U}_F$ for $\omega/J_x = 25,\ 20$. Parameters: $N_p = 4,\ U/J_x = 10,\ \eta/J_x = 0.05$.

A. Model

Following results of Ref. [15], we consider a slow ramp of the tunneling amplitude along y direction, $J_y(t)$, as well as a slow ramp of the driving amplitude $\kappa(t)$. Namely, we start from a series of decoupled wires along the x direction and start coupling them. More precisely, initial states are selected as the ground states of $\hat{H}_{\mathrm{ini}}$:

$$\hat{H}_{\mathrm{ini}} = -J_x \sum_{m,n} (\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \mathrm{H.c.}) + \frac{U}{2}\sum_{m,n}\hat{n}_{m,n}(\hat{n}_{m,n}-1). \tag{16}$$

For the filling factors that we consider, the ground states of the $\hat{H}_{\mathrm{ini}}$ are simple noninteracting states with the ground state energy $E_{0,\mathrm{ini}} = -2J_xN_p$. Out of the several degenerate ground states, we select those where atoms occupy every second wire. There are two such states and we label them as $|\psi_+\rangle$ (even wires occupied) and $|\psi_-\rangle$ (odd wires occupied). These states have finite projections only onto the sectors $k_x = 0,\ k_y = 0$ and $k_x = 0,\ k_y = \pi$ of the driven model from Eq. (1). Therefore we may expect the two initial states $|\psi_\pm(t=0)\rangle$ to be transformed into the two Laughlin states during the ramp.

Having prepared the initial state, we slowly restore the tunneling amplitude along the y direction, $J_y(t)$, and slowly ramp up the driving amplitude $\kappa(t)$. The time-evolution is governed by

$$\hat{H}_{\mathrm{sr}}(t) = -J_x\sum_{m,n}(\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \mathrm{H.c.}) - J_y(t)\sum_{m,n}(e^{i\omega t}\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} + \mathrm{H.c.}) + \frac{\kappa(t)}{2}\sum_{m,n}\sin\left[\omega t - (m+n-1/2)\,\phi\right]\hat{n}_{m,n} + \frac{U}{2}\sum_{m,n}\hat{n}_{m,n}(\hat{n}_{m,n}-1), \tag{17}$$

where $J_y(t) = J_y\tanh(\eta t)$, $\kappa(t) = \kappa\tanh(\eta t)$, η being the ramping rate. In the long-time limit, we recover the original Hamiltonian from Eq. (1). During the ensuing time evolution we construct the mixture

$$\rho(t) = \tfrac{1}{2}(|\psi_+(t)\rangle\langle\psi_+(t)| + |\psi_-(t)\rangle\langle\psi_-(t)|). \tag{18}$$

We monitor stroboscopically the energy expectation value

$$E(t) = \mathrm{tr}(\rho(t)\hat{H}_{\mathrm{eff}}) \tag{19}$$

and the PES of $\rho(t)$.

B. Results

In Fig. 5(a) we present the energy expectation value from Eq. (19) for $U/J_x = 10$ and several driving frequencies $\omega/J_x = 25,\ 20,\ 15,\ 10$. Our numerical results indicate that ramps with the rates up to $\eta/J_x \sim 0.1$ work reasonably well. Slower ramps give better results, but are less practical [15]. By construction, the initial state is a noninteracting state with particles delocalized along the x direction and therefore the initial energy is $E(t=0) = -2\,N_p\,J_x$. During the ramp with the rate $\eta/J_x = 0.05$, for the regime of high driving frequencies, down to approximately $\omega/J_x = 20$, we find that the energy initially decreases and reaches an almost constant value at around $tJ_x \sim 20$. On the other hand, for $\omega/J_x = 15$, the system slowly heats up during the ramping process, and for $\omega/J_x = 10$ the system quickly reaches the infinite-temperature state.

One of our main results is summarized in Fig. 5(b), where we plot the particle-entanglement gap of $\rho(t)$, from Eq. (18), as a function of time. In the high-frequency regime $\omega/J_x \geqslant 20$, starting around $tJ_x \sim 20$ we find a persistent particle-entanglement gap, marking the onset of a topologically nontrivial state. It is even more interesting that, even for $\omega/J_x \sim 15$, the state seems to exhibit a finite gap on intermediate time-scales. This is not the case for $\omega/J_x \leqslant 10$, where the gap quickly vanishes. In Fig. 5(c), we present the value of the overlap $\mathrm{tr}\,[\rho(t)\rho_F]$, of the time-evolved mixed state with the relevant state from Eq. (15) for $N_p = 4$. Clearly, the slow ramp of the type given in Eq. (17) allows for the preparation of the relevant eigenstates of $\hat{U}_F$ with high fidelity (better than 1%).

In Figs. 6(a) and 6(b) we show the time evolution of the PES in the two momentum sectors $k^A_y = 0$ and $k^A_y = \pi/6$ for $N_p = 6$, $U/J_x = 5$, and $\eta/J_x = 0.05$. The PES of the initial state is easy to understand. As the $L_y/2$ wires are occupied by single atoms, the reduced density matrix is proportional to the identity matrix with the proportionality factor yielding $-\ln\xi_n = \ln\left(2\binom{Ly/2}{N_A}\right) \approx 3.69$. During the ramp we find that additional modes in PES are gaining weight and moving down in the spectrum. Finally, the state $\rho(t)$ reached around $t \approx 50T$ exhibits a well defined gap and the correct counting of the low-lying modes: there are ten low-lying modes for $k^A_y = 0$

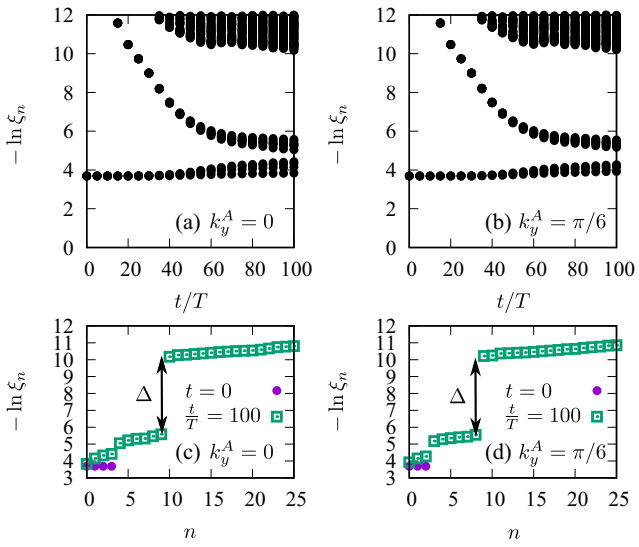

FIG. 6. The low-lying part of the particle-entanglement spectra $-\ln \xi_n$ of $\rho(t)$, Eq. (18), during the time evolution governed by Eq. (17) in the (a) $k_y^A = 0$ and (b) $k_y^A = \pi/6$ momentum sectors. The low-lying part of the PES in the sectors (c) $k_y^A = 0$ and (d) $k_y^A = \pi/6$, at two instances of time $t = 0$ and $t/T = 100$. Parameters: $N_p = 6$, $U/J_x = 5$, $\omega/J_x = 15$, $\eta/J_x = 0.05$.

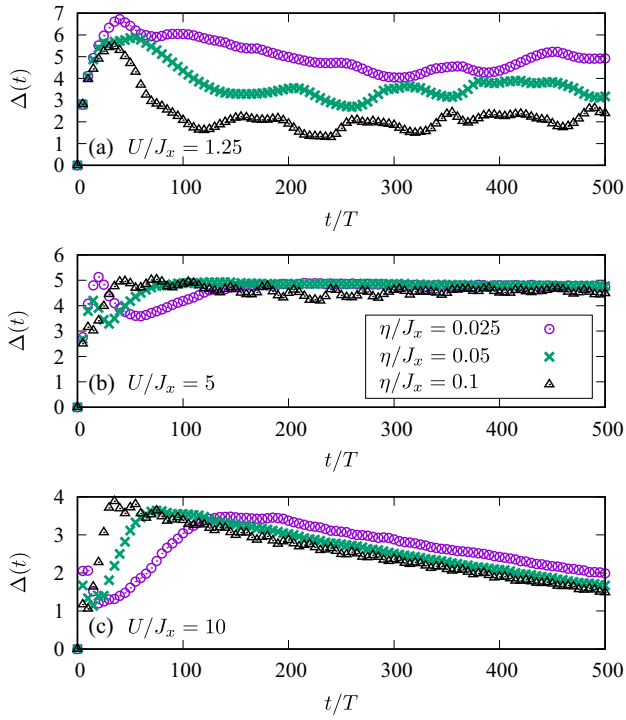

FIG. 7. The particle-entanglement gap $\Delta(t)$ as a function of time during the time evolution governed by Eq. (17), for several interaction strengths (a) $U/J_x = 1.25$ (b) $U/J_x = 5$ and (c) $U/J_x = 10$, and several ramping rates $\eta/J_x = 0.025, 0.05, 0.1$. Other parameters: $N_p = 5$, $\omega/J_x = 15$.

and nine low-lying modes for $k_y^A = \pi/6$; see Figs. 6(c) and 6(d) and also Table I.

In Fig. 7 we discuss a satisfactory range of ramping rates η for a given interaction strength U and a given driving frequency ω that we fix at $\omega/J_x = 15$. The obtained numerical results suggest that at weaker interaction strengths $U/J_x \leqslant 2$, slower ramping rates are needed. One way to explain this behavior is by using the effective model and arguing that the gap protecting the Laughlin state is smaller at weaker U. On the other hand, for stronger interaction strengths $U/J_x \geqslant 8$ the particle-entanglement gap closes at later stages as the heating process becomes dominant. Finally, in the intermediate range $U/J_x \sim 5$, faster ramps with $\eta/J_x = 0.1$ lead to the sought-after state $\rho(t)$ from Eq. (18), with persistent features in the PES up to $t = 500T$. These results indicate that, when optimizing the ramping protocol in an actual experiment, there will be a tradeoff between the unfavorable heating and a faster ramping into the desired state, as both of these processes are promoted by interactions.

V. CONCLUSIONS

The technique of Floquet engineering has been successfully exploited for the implementation of synthetic magnetic fields in driven optical lattices. Following up on these achievements and on a long-standing pursuit for the FQH states in cold-atom setups, in this paper we have addressed possible realization of the bosonic Laughlin state in a small atomic sample in a periodically driven optical lattice. While a thermodynamically large interacting system generally heats up into an infinite-temperature state under driving, the heating process can be controlled to some extent in a few-particle system.

We have assumed a realistic driving protocol and finite on-site interactions, and we have identified the FQH state based on analysis of its particle-entanglement spectra. Results of our numerical simulations show that the stroboscopic dynamics of $N_p = 4, 5, 6$ particles supports the topological $\nu = 1/2$ Laughlin state down to $\omega/J_x = 20$ for $U/J_x = 10$, and down to $\omega/J_x = 15$ for $U/J_x = 1$, for the driving amplitude $\kappa/\omega = 0.5$. These results are in reasonable agreement with the recent estimates of the optimal heating times [51] that take into account the contribution of the higher bands of the underlying optical lattice. In addition, we have investigated slow ramping of the driving term and found that it allows for the preparation of the Laughlin state on experimentally realistic time-scales of the order of 20 $\hbar/J_x$, where $\hbar/J_x$ is the tunneling time. Interestingly, we find that some topological features persist during an intermediate stage even in the regime where the system exhibits a slow transition into the infinite-temperature state (e.g., $\omega/J_x = 15$ for $U/J_x = 10$).

In the future, we plan to address the preparation scheme for the relevant correlated states in a driven honeycomb lattice, which exhibits lower heating rates in comparison to a cubic lattice according to the recent experiments [48,49]. Another highly relevant question, that we have not tackled and that we postpone to future investigation, concerns suitable experimental probes of topological features. The recent progress in the field has led to the development of several detection protocols specially suited for the cold-atom systems [66–71]. For the type of systems considered in this paper, the most promising are results of the recent study [71] showing that fractional excitations can be probed even in small systems of several bosons.

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APPENDIX: DRIVEN OPTICAL LATTICES

In this Appendix we review the derivation of the model given in Eq. (1). The system is described by

$$\hat{H}_{\rm lab}(t) = \hat{H}_{BH} + \hat{H}_{\rm drive}(t) + \omega\,\hat{V}, \tag{A1}$$

where we start with the Bose-Hubbard model

$$\begin{aligned}\hat{H}_{BH} = &-J_x \sum_{m,n}(\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \text{H.c.}) \\ &- J_y \sum_{m,n}(\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} + \text{H.c.}) \\ &+ \frac{U}{2}\sum_{m,n}\hat{n}_{m,n}(\hat{n}_{m,n}-1),\end{aligned} \tag{A2}$$

and we introduce an offset $\omega\hat{V}$:

$$\hat{V} = \sum_{m,n} n\,\hat{n}_{m,n}. \tag{A3}$$

This shifted Bose-Hubbard model is exposed to a suitable resonant driving scheme:

$$\begin{aligned}\hat{H}_{\rm drive}(t) &= \frac{\kappa}{2}\sum_{m,n}\sin\left(\omega t - \phi_{m,n} + \frac{\phi}{2}\right)\hat{n}_{m,n}, \\ \phi_{m,n} &= (m+n)\,\phi.\end{aligned} \tag{A4}$$

We assume periodic boundary conditions compatible with the driving term (A4) in the laboratory frame. To this purpose we use vectors $\mathbf{R}_1 = 4\,\mathbf{e}_x$ and $\mathbf{R}_2 = -\mathbf{e}_x + \mathbf{e}_y$ as presented in Fig. 1. For simplicity, we work in the rotating frame

$$|\psi_{\rm rot}(t)\rangle = e^{i\omega t\hat{V}}|\psi_{\rm lab}(t)\rangle \tag{A5}$$

and derive the Schrödinger equation

$$i\frac{d|\psi_{\rm rot}(t)\rangle}{dt} = \hat{H}_{\rm rot}(t)|\psi_{\rm rot}(t)\rangle, \tag{A6}$$

where

$$\hat{H}_{\rm rot}(t) = (e^{i\omega t\hat{V}}\hat{H}_{\rm lab}(t)e^{-i\omega t\hat{V}} - \omega\hat{V}). \tag{A7}$$

Now we calculate $\hat{H}_{\rm rot}(t)$ explicitly. The only nontrivial action of this rotation on $\hat{H}_{\rm lab}$ comes from the nearest-neighbor hopping along y direction. Indeed, we have

$$e^{i\omega t\hat{V}}\hat{a}^\dagger_{m,n}\hat{a}_{m,n'}e^{-i\omega t\hat{V}} = e^{i\omega t(n-n')}\hat{a}^\dagger_{m,n}\hat{a}_{m,n'}. \tag{A8}$$

In total we obtain

$$\begin{aligned}\hat{H}_{\rm rot}(t) = &-J_x\sum_{m,n}(\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \text{H.c.}) + \frac{U}{2}\sum_{m,n}\hat{n}_{m,n}(\hat{n}_{m,n}-1) \\ &+ e^{i\omega t}\hat{H}_1 + e^{-i\omega t}\hat{H}_{-1} + e^{-i\omega t(L_y-1)}\hat{H}_{L_y-1} \\ &+ e^{i\omega t(L_y-1)}\hat{H}_{-L_y+1},\end{aligned} \tag{A9}$$

with

$$\begin{aligned}\hat{H}_1 &= -J_y\sum_{m,n}^{\rm OBC}\left(\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} - \frac{i}{4}\kappa e^{i(-\phi_{m,n}+\frac{\phi}{2})}\hat{n}_{m,n}\right), \\ \hat{H}_{-1} &= \hat{H}_1^\dagger,\end{aligned} \tag{A10}$$

$$\hat{H}_{-L_y+1} = -J_y\sum_m \hat{a}^\dagger_{m,0}\hat{a}_{m-L_y,L_y-1}, \quad \hat{H}_{L_y-1} = \hat{H}^\dagger_{-L_y+1}. \tag{A11}$$

In the terms $\hat{H}_{-L_y+1}$ and $\hat{H}_{L_y-1}$ we take into account periodic boundary conditions along the direction parallel to $\mathbf{R_2}$ as imposed in the laboratory frame. In order to limit the complexity of the numerical calculation, we keep translational invariance and impose the periodic boundary conditions in both directions in the rotating frame. This implies that we will neglect "phasors" $e^{-i\omega t(L_y-1)}$ and $e^{i\omega t(L_y-1)}$. Under these assumptions, we can recast Eq. (A9) into the time-dependent Hamiltonian given in Eq. (1). In practice, this would require engineering additional non-trivial terms in the laboratory frame.

The leading order of the kick operator is given by

$$\hat{K}(t=0) \approx -\frac{\kappa}{2\omega}\sum_{m,n}\cos(\phi_{m,n} - \phi/2)\hat{n}_{m,n}. \tag{A12}$$

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Quantum scars of bosons with correlated hopping

Ana Hudomal [1✉], Ivana Vasić [1], Nicolas Regnault[2,3] & Zlatko Papić[4]

Recent experiments on Rydberg atom arrays have found evidence of anomalously slow thermalization and persistent density oscillations, which have been interpreted as a many-body analog of the phenomenon of quantum scars. Periodic dynamics and atypical scarred eigenstates originate from a "hard" kinetic constraint: the neighboring Rydberg atoms cannot be simultaneously excited. Here we propose a realization of quantum many-body scars in a 1D bosonic lattice model with a "soft" constraint in the form of density-assisted hopping. We discuss the relation of this model to the standard Bose-Hubbard model and possible experimental realizations using ultracold atoms. We find that this model exhibits similar phenomenology to the Rydberg atom chain, including weakly entangled eigenstates at high energy densities and the presence of a large number of exact zero energy states, with distinct algebraic structure.

[1] Institute of Physics Belgrade, University of Belgrade, 11080 Belgrade, Serbia. [2] Joseph Henry Laboratories and Department of Physics, Princeton University, Princeton, NJ 08544, USA. [3] Laboratoire de Physique de l'Ecole Normale Supérieure, ENS, Université PSL, CNRS, Sorbonne Université, Université Paris-Diderot, Sorbonne Paris Cité, 75005 Paris, France. [4] School of Physics and Astronomy, University of Leeds, Leeds LS2 9JT, UK. ✉email: ana.hudomal@ipb.ac.rs

Semiclassical studies of chaotic stadium billiards have revealed the existence of remarkable non-chaotic eigenfunctions called "quantum scars"[1]. Scarred eigenfunctions display anomalous enhancement in regions of the billiard that are traversed by one of the periodic orbits in the classical limit when $\hbar \to 0$. It was shown that quantum scars lead to striking experimental signatures in a variety of systems, including microwave cavities[2], quantum dots[3], and semiconductor quantum wells[4].

A recent experiment on a quantum simulator[5], and subsequent theoretical work[6,7], have shown that quantum many-body scars can occur in strongly interacting quantum systems. The experiment used a one-dimensional Rydberg atom platform in the regime of the Rydberg blockade[5,8,9], where nearest-neighbor excitations of the atoms were energetically prohibited. The experiment observed persistent many-body revivals of local observables after a "global quench"[10] from a certain initial state. In contrast, when the experiment was repeated for other initial configurations, drawn from the same type of "infinite" temperature ensemble, the system displayed fast equilibration and no revivals. These observations pointed to a different kind of out-of-equilibrium behavior compared to previous studies of quantum thermalization in various experimental platforms[11–15].

In both single-particle and many-body quantum scars, the dynamics from certain initial states leads to periodic revivals of the wave function. In the former case, this happens when the particle is prepared in a Gaussian wave packet initialized along a periodic orbit[1], while in the latter case the revivals can be interpreted as a nearly-free precession of a large emergent su(2) spin degree of freedom[16,17]. Another similarity between single- and many-body quantum scars is the existence of non-ergodic eigenstates. In the single-particle case, such eigenstates are easily identified by their non-uniform probability density that sharply concentrates along classical periodic orbits. In the many-body case, non-ergodic eigenstates are broadly defined as those that violate eigenstate thermalization hypothesis (ETH)[18,19]. Scarred eigenstates violate the ETH in a number of ways: for example, they appear at evenly spaced energies throughout the spectrum[6,20,21], they have anomalous expectation values of local observables compared to other eigenstates at the same energy density, and their entanglement entropy obeys a sub-volume law scaling[20].

In recent works, the existence of atypical eigenstates has been taken as a more general definition of quantum many-body scaring. For example, highly excited eigenstates with low entanglement have previously been analytically constructed in the non-integrable AKLT model[22,23]. A few of such exact eigenstates are now also available for the Rydberg atom chain model[24]. The collection of models that feature atypical eigenstates is rapidly expanding, including perturbations of the Rydberg atom chain[20,25,26], theories with confinement[27–29], Fermi–Hubbard model beyond one dimension[30,31], driven systems[32], quantum spin systems[33,34], fractional quantum Hall effect in a one-dimensional limit[35], and models with fracton-like dynamics[36–39]. In a related development, it was proposed that atypical eigenstates of one Hamiltonian can be "embedded" into the spectrum of another, thermalizing Hamiltonian[40], causing a violation of a "strong" version of the ETH[41,42]. This approach allows to engineer scarred eigenstates in models of topological phases in arbitrary dimensions[43]. From a dynamical point of view, it has been shown that models with scarred dynamics can be systematically constructed by embedding periodic on-site unitary dynamics into a many-body system[44].

A feature shared by many scarred models is the presence of some form of a kinetic constraint. In the Rydberg atom chain, the constraint results from strong van der Waals forces, which project out the neighboring Rydberg excitations[45]. Such Hilbert spaces occur, for example, in models describing anyon excitations in topological phases of matter[46–50] and in lattice gauge theories[51–53], including the Rydberg atom system[54,55]. Recent works on periodically driven optical lattices have started to explore such physics[56,57]. On the other hand, kinetic constraints have been investigated as a possible pathway to many-body localization without disorder[58]. In classical systems, non-thermalizing behavior without disorder is well known in the context of structural glasses[59–61]. The mechanism of this type of behavior is the excluded volume interactions that impose kinetic constraints on the dynamics[62,63]. Similar type of physics has recently been explored in quantum systems where a "quasi many-body localized" behavior was proposed to occur in the absence of disorder[64–74].

In this paper, we investigate the relation between kinetic constraints, slow dynamics and quantum many-body scars. In contrast to previous work, which focused on models of spins and fermions that are closely related in one dimension due to the Jordan–Wigner mapping, here we study one-dimensional models of bosons with density-assisted hoppings, which realize both "hard" and "soft" kinetic constraints, whilst being non-integrable. Depending on the form of the hopping term, we demonstrate that the models encompass a rich phenomenology, including regimes of fast thermalization, the existence of periodic revivals and many-body scars, as well as the Hilbert space fragmentation that has been found in recent studies of fractonic models[36–39]. Unlike the experimentally realized Rydberg atom system, we find evidence of many-body scars in a bosonic model without a hard kinetic constraint, i.e., with a fully connected Hilbert space. We identify initial states that give rise to periodic many-body revivals in the quantum dynamics, and we introduce a "cluster approximation" that captures the scarred eigenstates that are responsible for periodic revivals. We discuss possible experimental realizations of these models using ultracold atoms.

Results

Models and their Hilbert spaces. A fundamental ingredient of kinetically constrained models is "correlated hopping": a particle can hop depending on the state of its neighbors. In this paper we consider a system of N_p bosons on a one-dimensional lattice with L sites. We consider models where the total filling factor, $\nu = N_p/L$, is conserved, and we will mainly present results in the dense regime, $\nu = 1$. We have studied models with $\nu < 1$ and $\nu > 1$, but we found them to be either too constrained or not constrained enough, and therefore less interesting. We emphasize that the bosons in our study are not hard-core, i.e., the occupancy of any lattice site can take any value from 0 to N_p.

We study three different models, defined by the Hamiltonians:

$$H_1 = -J\sum_{j=1}^{L}\left(b_j^\dagger b_{j+1} n_j + n_{j-1} b_j^\dagger b_{j-1}\right), \quad (1)$$

$$H_2 = -J\sum_{j=1}^{L}\left(n_j b_j^\dagger b_{j+1} + b_j^\dagger b_{j-1} n_{j-1}\right), \quad (2)$$

$$H_3 = -J\sum_{j=1}^{L}\left(n_{j+1} b_j^\dagger b_{j+1} n_j + n_{j-1} b_j^\dagger b_{j-1} n_j\right). \quad (3)$$

All three models contain a free-boson hopping term, $b_j^\dagger b_{j+1}$, which is dressed in various ways by density operators, $n_j = b_j^\dagger b_j$. We will show that the position of the density operator n_j completely changes the behavior of these models, ranging from fast thermalization to the breakup of the Hamiltonian into

disconnected, exactly solvable sectors. For example, note that H_1 and H_2 are related to each other via free-boson hopping,

$$H_2 = H_1 - J\sum_j \left(b_j^\dagger b_{j+1} + b_j^\dagger b_{j-1}\right), \quad (4)$$

which can be easily proven using bosonic commutation relations. We will see below that this innocuous free-boson hopping leads to surprisingly different dynamical properties of the two models.

The motivation behind introducing three different models in Eqs. (1)–(3) can be summarized as follows. Hamiltonian H_1 describes a model where a particle cannot hop to the left if that site is not already occupied by at least one particle, and cannot hop to the right if it is the only particle left on its initial site. This introduces constraints to the system. Conversely, there are no such constraints in the case of H_2. Indeed, the hopping coefficients are only modified in intensity by the particle-number operator. Hamiltonian H_3 introduces additional constraints compared to H_1. The number of unoccupied sites and their positions remain constant under the action of this Hamiltonian. This leads to different connectivity of the Hilbert space in each of the models, as we explain in the next Section.

We consider periodic boundary conditions ($L+1 \equiv 1$) and set $\hbar = J = 1$. With periodic boundary conditions, all three Hamiltonians H_1, H_2 and H_3 have translation symmetry, thus their eigenstates can be labeled by momentum quantum number, k, quantized in units of $2\pi/L$. In addition, H_3 has inversion symmetry. We denote by $I = 0$ and $I = 1$ the sectors that are even and odd under inversion, respectively.

Without restrictions on the boson occupancy, the Hilbert space of H_1, H_2 and H_3 grows very rapidly. For $L = N_p = 12$, the Hilbert space size of the $k = 0$ sector is 112720 (the largest one we will consider for H_1 and H_2). As previously mentioned (see also the next Section), the Hilbert space of H_3 splits into many disconnected components, thus it is possible to consider only one connected component at a time and disregard the unoccupied sites whose positions do not change. This is more relevant when looking at properties such as thermalization, than fixing the filling factor. However, the boundary conditions are in that case no longer periodic, and the system does not have translation symmetry. Considering only a system with the size $L/2$, filling factor $\nu = 2$, open boundary conditions and minimal number of particles per site equal to 1 is completely equivalent to considering the largest component of the full system which has the size L, filling factor $\nu = 1$, periodic boundary conditions and no restrictions on the occupancies. The Hilbert space size of the symmetric invariant sector of the largest connected component of $L = N_p = 22$ is 176484 and this is the largest sector that we will consider for H_3.

Graph structure of the models. Since we will be interested in the dynamical properties, it is convenient to first build some intuition about the structure of the Hamiltonians of the three models in Eqs. (1)–(3). A Hamiltonian can be viewed as the adjancency matrix of a graph whose vertices are Fock states of bosons, $|n_1, n_2, \ldots, n_L\rangle$. If the Hamiltonian induces a transition between two Fock states, the corresponding vertices of the graph are connected by a link. The graphs that show how the configuration space is connected have very different structure for the three Hamiltonians H_1, H_2, and H_3, as can be observed in Fig. 1.

The entire graph of H_2 is well connected and it has the same structure as the graph of the standard Bose-Hubbard model: the particle-number operators in H_2 do not introduce any constraints, but only affect the magnitude of the hopping coefficients. In contrast, the H_1 graph shows several clusters of configurations that are weakly connected to the rest of the graph. "Weakly connected" means that there is a small number of connections leading outside the cluster and that their respective hopping coefficients are smaller in magnitude than those of the surrounding connections within the cluster. A state that is initially located inside a cluster is therefore more likely to stay inside during an initial stage of the time evolution, which increases the probability of revivals and slows down the growth of

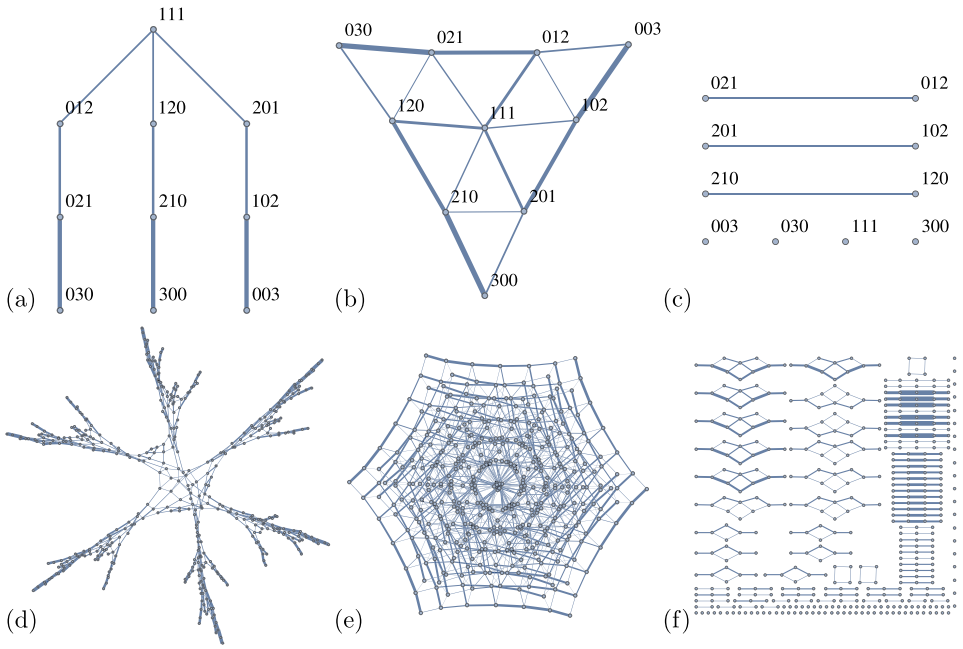

Fig. 1 Connectivity of the Hilbert space. Adjacency graph for **a** H_1, **b** H_2, **c** H_3, all for $L = N_p = 3$. **d**–**f** same as **a**, **b** and **c** but for $L = N_p = 6$. To avoid clutter, we do not label the vertices in **d**–**f**. All graphs are weighted, i.e., the line thickness is proportional to the magnitude of the corresponding hopping coefficient. Several different clusters of configurations are visible in the case of H_1. The clusters start to form already for $L = 3$ (for example, the configurations 012-021-003 in **a**) and become more prominent for $L = 6$ **d**. In the case of H_2, almost all configurations are well connected to the rest of the graph. The graphs for H_3 show that the Hilbert space is highly reducible: its graph splits into many disconnected components.

entanglement entropy. We will provide a more quantitative description and examples that illustrate this in Section "Quantum scars in H_1 and H_3 models". Finally, the graph of H_3, due to even stronger constraints, is actually disconnected, which is an example of Hilbert space fragmentation that was previously shown to cause non-ergodic behavior in fracton-like models[37,38]. This predicts that thermalization and dynamics in the three models will be very different, which we will confirm in the following Section. However, we note that the number of connections and the topology of the graph is not the only relevant factor for the dynamics. The magnitude of the hopping coefficients between different configurations is also important (Supplementary Note 1).

We note that the relation between H_1 and H_3 is reminiscent of the relation between the quantum East model[75] and the "PXP" model describing the atoms in the Rydberg blockade regime[6,20,45]. Like H_3, the PXP model is doubly constrained and inversion symmetric, while H_1 and the quantum East model are asymmetric versions of those two models with only a single constraint. The graph of the quantum East model is similar to that of H_1, in that it contains bottlenecks which slow down the growth of entanglement entropy[75].

Dynamics and entanglement properties. We now investigate the phenomenology of the models introduced in Eqs. (1)–(3). We use exact diagonalization to obtain the complete set of energy eigenvalues and eigenvectors, from which we evaluate the level statistics and the distribution of entanglement entropies for the three models. Furthermore, we probe dynamical properties of the models by studying a global quench, simulated via Krylov iteration.

The energy level statistics is a standard test for thermalization of models that cannot be solved exactly. A convenient way to probe the level statistics is to examine the probability distribution $P(r)$[76] of ratios between consecutive energy gaps $s_n = E_{n+1} - E_n$,

$$r = \frac{\min(s_n, s_{n+1})}{\max(s_n, s_{n+1})}. \tag{5}$$

The advantage of studying $P(r)$, instead of $P(s_n)$, is that there is no need to perform the spectrum unfolding procedure—see ref. [77]. For standard random matrix theory ensembles, both $P(r)$ and the mean $\langle r \rangle$ are well known[78]. When computing the same quantities in a microscopic physical model, it is crucial to resolve all the symmetries of the model.

The probability distribution $P(r)$ of the ratios of two consecutive energy gaps is shown in Fig. 2a–c for the three Hamiltonians H_1, H_2, and H_3 respectively, and two momentum or inversion sectors. In all three cases, the energy levels repel, i.e., the distribution tends to zero as $r \to 0$. For H_2, the distribution is particularly close to the Wigner–Dyson (non-integrable) line. For H_1, the distribution is also consistent with Wigner–Dyson when we restrict to the middle 1/3 of the spectrum (and after removing special states with $E = 0$). We exclude the edges of the spectrum because they contain degeneracies which are not symmetry-related. However, such states do not appear to have a major effect on the level statistics distribution, which is still closer to the Wigner–Dyson than the Poisson distribution even if they are included. The level statistics of H_3 within the largest connected component of the Hilbert space is shown in Fig. 2c and is also consistent with the Wigner–Dyson distribution without restricting the spectrum. However, we will demonstrate below that the dynamics in some smaller connected components of H_3 can be exactly solved.

As a complementary diagnostic of thermalization, we next compute the entanglement entropy of all eigenstates. We divide the lattice into two sublattices, A and B, of lengths L_A and

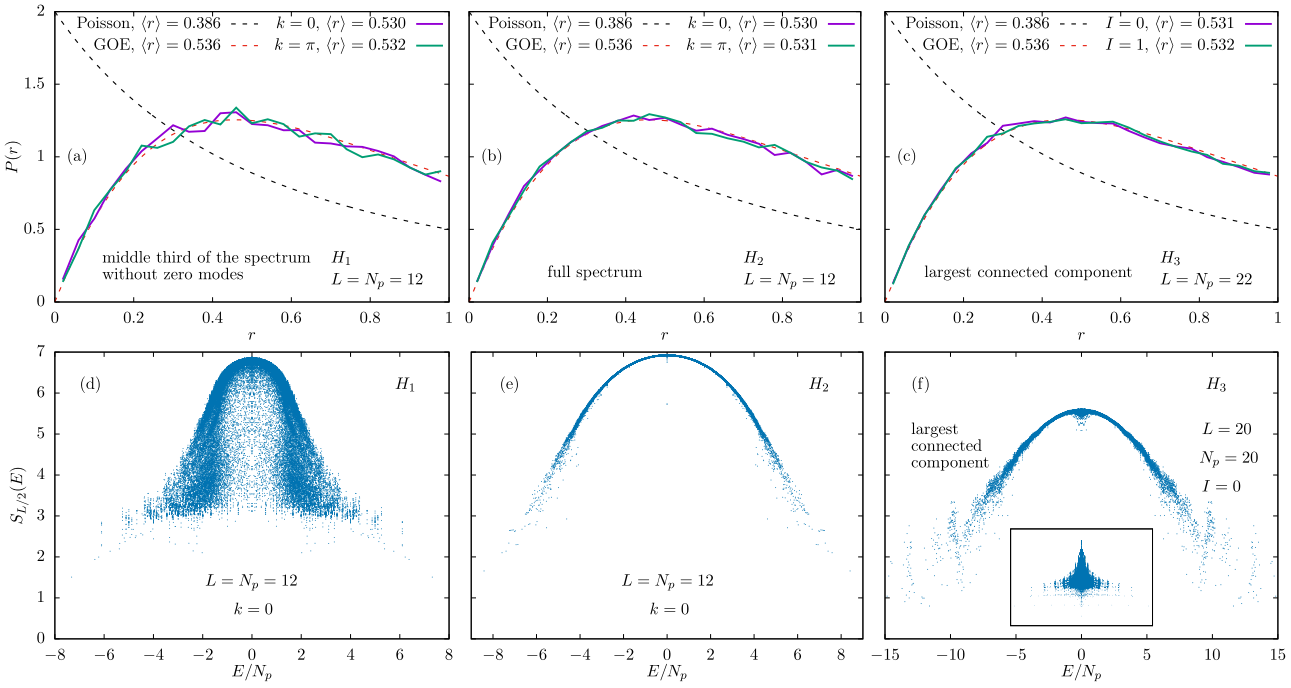

Fig. 2 Level statistics and entanglement. a–c Probability distribution of the ratios of two consecutive energy gaps. **a** H_1 (middle third of the spectrum without $E = 0$ states, $L = N_p = 12$), **b** H_2 (full spectrum, $L = N_p = 12$) and **c** H_3 (largest connected component of $L = N_p = 22$). The black dashed line shows the Poisson distribution, which corresponds to the integrable case, while the red dashed line is the distribution of the Gaussian orthogonal ensemble, which corresponds to the thermalizing case. **d–f** Entanglement entropies $S_{L/2}$ of all eigenstates plotted as a function of the eigenstate energy per particle, E/N_p. **d** H_1 ($L = N_p = 12$, $L_A = 6$, $k = 0$), **e** H_2 (same) and **f** H_3 in the largest connected component of $L = N_p = 20$, $L_A = 10$, $I = 0$. The inset shows all connected components for $L = N_p = 12$, $L_A = 6$, $k = 0$.

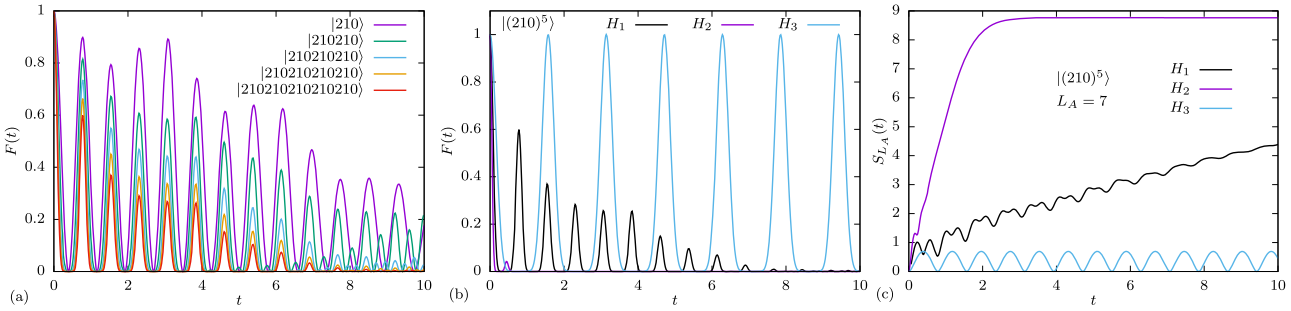

Fig. 3 Dynamics of quantum fidelity and entanglement entropy for initial configurations in Eq. (9). a Time evolution of fidelity $F(t)$ in Eq. (8) for system sizes $L = 3n$. The evolution is governed by the Hamiltonian H_1, different colors represent different system sizes L. **b** Fidelity evolution $F(t)$ for the Hamiltonians H_1, H_2 and H_3 and system size $L = 15$. **c** Entanglement entropy evolution $S_{L_A}(t)$ for the same cases as in **b**.

$L_B = L - L_A$. For a given pure state $|\psi\rangle$, the entanglement entropy is defined as

$$S_A = -\mathrm{tr}_A(\rho_A \ln \rho_A), \tag{6}$$

where $\rho_A = \mathrm{tr}_B|\psi\rangle\langle\psi|$ is the reduced density matrix of the subsystem A. The scatter plots, showing entanglement entropy of all eigenstates $|E_n\rangle$ as a function of their energy E_n, are displayed in Fig. 2d–f. Here we take into account the translation symmetry of the system and work in the momentum sector $k = 0$ for H_1 and H_2, and consider only the largest connected component and the inversion sector $I = 0$ for H_3. The results for other sectors are qualitatively similar.

Entanglement entropy distribution in Fig. 2d, e reveals a striking difference between the Hamiltonians H_1 and H_2, even though they only differ by a free-boson hopping term, Eq. (4). The model H_1 is constrained, which leads to a large spread of the entropy distribution and many low-entropy eigenstates including in the bulk of the spectrum. From this perspective, H_1 is reminiscent of PXP model[20,25]. By contrast, H_2 has no such constraints and in this case the entanglement entropy is approximately a smooth function of the eigenstate energy. The Hamiltonian H_3 is doubly constrained, and this is reflected in its entanglement distribution, which also shows a large spread and several disconnected bands, reminiscent of an integrable system like the XY model[79].

Global quenches. The constraints in the models in Eqs. (1), (2), and (3) have significant effects on the dynamics governed by these Hamiltonians. We probe the dynamics by performing a global quench on the system. We assume the system is isolated and prepared in one of the Fock states, $|\psi_0\rangle$, at time $t = 0$. We restrict to $|\psi_0\rangle$ being product states which are not necessarily translation-invariant, as such states are easier to prepare in experiment. However, our results remain qualitatively the same if we consider translation-invariant $|\psi_0\rangle$. After preparing the system in the state $|\psi_0\rangle$, which is not an eigenstate of the Hamiltonian, the system is let to evolve under unitary dynamics,

$$|\psi(t)\rangle = \exp\left(-\frac{i}{\hbar}Ht\right)|\psi_0\rangle. \tag{7}$$

where H is one of the Hamiltonians of interest. From the time-evolved state, we evaluate the quantum fidelity,

$$F(t) = |\langle\psi_0|\psi(t)\rangle|^2, \tag{8}$$

i.e., the probability for the wave function to return to the initial state. In a general many-body system, fidelity is expected to decay as $F(t) \sim \exp(-L(Jt)^2)$. It thus becomes exponentially suppressed in the system size for any fixed time t^*, i.e., $F(t^*) \sim \exp(-cL)$, where c is a constant. In scarred models, such as the Rydberg atom chain, fidelity at the first revival peak occurring at a time T still decays exponentially, but exponentially slower, i.e., $F(T) \sim \exp(-c'L)$, with $c' \ll c$. In ref. [20], for a finite system with $L \lesssim 32$ atoms, the fidelity at the first revival can be as high as ~70%, and several additional peaks at times nT are also clearly visible.

We first consider the Hamiltonian H_1. Several configurations exhibit periodic revivals of the fidelity $F(t)$, which can in some cases be higher than 90%. Most of these configurations involve a very dense cluster of bosons such as $|...0N10...\rangle$. In contrast, a completely uniform configuration $|...111...\rangle$ thermalizes very quickly. Here we focus on periodically-reviving configurations with density being as uniform as possible. One family of such reviving configurations involves n unit cells made of three lattice sites:

$$|210210\ldots\,210\rangle \equiv |(210)^n\rangle. \tag{9}$$

Time evolution of the fidelity for the initial state $|(210)^n\rangle$ for different system sizes $L = 3n$ is shown in Fig. 3a. The initial state is assumed to be the product state, e.g., $|\psi_0\rangle = |210\rangle$ for $L = 3$. The frequency of the revivals in Fig. 3 is approximately the same for all system sizes. We emphasize that similar results are obtained for a translation-symmetric initial state, e.g., $|\psi_0\rangle = \frac{1}{\sqrt{3}}(|210\rangle + |021\rangle + |102\rangle)$. Both cases converge in the large system limit, and the differences are only significant for $L = 3$ when the revival frequency of the initial state with transition symmetry differs from the frequencies of other system sizes.

In Fig. 3b we compare the fidelity for the initial state in Eq. (9) when it is evolved by all three Hamiltonians in Eqs. (1)–(3). The initial state is fixed to be $|(210)^5\rangle$. We observe that the dynamics with H_3 has very prominent revivals; in fact as we will later show, these revivals are perfect and their period is approximately twice the revival period for H_1. In contrast, for H_2 the fidelity quickly drops to zero without any subsequent revivals.

Finally, in Fig. 3c we plot the time evolution of entanglement entropy. As expected from the fast decay of the fidelity, the entropy for H_2 rapidly saturates to its maximal value. Moreover, as expected from the perfect revivals in H_3, the entropy in that case oscillates around a constant value close to zero. For H_1, we observe a relatively slow growth of entropy, with oscillations superposed on top of that growth, again similar to PXP model[6]. For the initial state that is not translation-invariant, it is important how we cut the system, e.g., $|...210|210...\rangle$ versus $|...2102|10...\rangle$. In the first case, the entanglement entropy remains zero for H_3 because no particle can hop from one subsystem to the other, while in the second case the entropy oscillates around a constant value, which is the case in Fig. 3c.

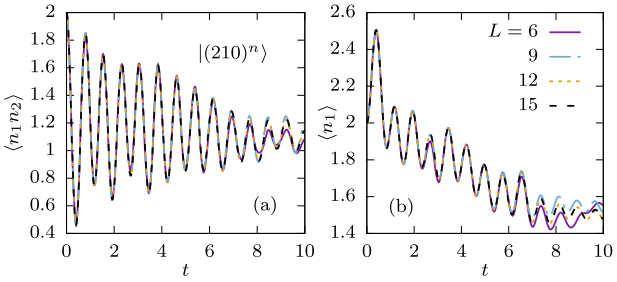

Fig. 4 Evolution of local observables for the Hamiltonian H_1. a Correlations between adjacent sites $\langle n_1 n_2(t)\rangle$ for different system sizes and the initial state $|(210)^n\rangle$. **b** Density on one site $\langle n_1(t)\rangle$.

In Fig. 4 we show the H_1 evolution of two local observables, density correlations between two adjacent sites $\langle n_1 n_2(t)\rangle$ and density on the first site $\langle n_1(t)\rangle$, starting from the initial state $|(210)^n\rangle$. Unlike fidelity and entanglement entropy, these observables can be easily measured in experiment. Both observables robustly oscillate with approximately the same frequency as the fidelity. The heights of the first few revival peaks are approximately converged for the system sizes ranging from $L = 6$ to $L = 15$, which suggests that revivals in such local observables can be observed in the thermodynamic limit. In the following Section, we will show that the oscillations observed in the dynamics from $|(210)^n\rangle$ state in Eq. (9) and their frequency can be explained using a tractable model that involves only a small subset of all configurations in the Hilbert space, thus providing a realization of quantum scars in a correlated bosonic system. Our starting point will be the model H_3, whose graph explicitly separates into disconnected subsets which makes the toy model exact, hence we can analytically calculate the revival frequency. Based on these results, we then introduce an approximation scheme that describes the dynamics from the same initial state under the H_1 Hamiltonian.

Quantum scars in H_1 and H_3 models. The quench dynamics of fidelity and entanglement entropy in Fig. 3 suggest that H_1 and H_3 models are candidate hosts for many-body scarred eigenstates that can be probed by initializing the system in product states $|(210)^n\rangle$. We now analyze the structure of these states using our approach called "cluster approximation" that is introduced in detail in Methods.

The dynamics of H_3 within the sector containing the state $|(210)^n\rangle$ can be solved exactly, as shown in Methods. The connected component of the state $|(210)^n\rangle$ consists of all possible combinations of patterns 210 and 120. This means that triplets of sites evolve independently, and dynamically the system behaves as a collection of independent two level systems (spins-1/2). From this observation, it can be shown that revivals will be perfect with a period $T_3 = \pi/2$. The same period is obtained for initial product state $|(210)^n\rangle$ and its translation-invariant version; if the initial state is both translation-invariant and inversion-symmetric, the period is doubled.

In contrast to the free dynamics in H_3, the H_1 model exhibits decaying revivals and does not admit an exact description. In order to approximate the quench dynamics and scarred eigenstates in H_1, we project the Hamiltonian to smaller subspaces of the full Hilbert space. These subspaces contain clusters of states which are poorly connected to the rest of the Hilbert space and thereby cause dynamical bottlenecks. As explained in Methods, the clusters can be progressively expanded to yield an increasingly accurate description of the dynamics from a given initial state.

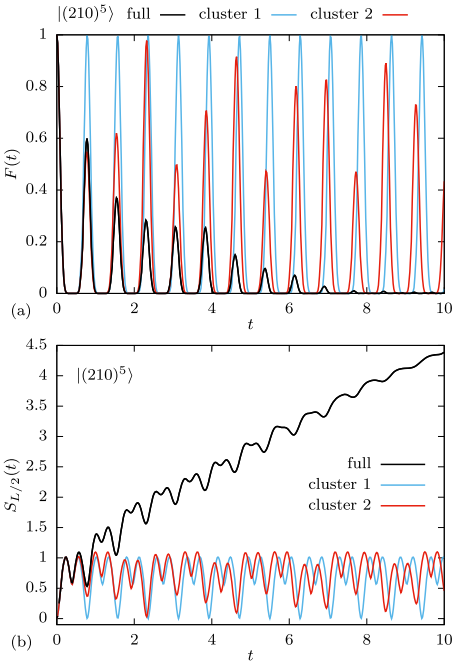

Fig. 5 Comparison of the full dynamics against the minimal cluster (1) and extended cluster (2) approximation schemes. We consider the system size $L = 15$ with the initial state $|(210)^5\rangle$. **a** Time evolution of the fidelity. The frequency of revivals is approximately the same in both cases, but the results for the extended cluster show better agreement with the results for the full Hilbert space. **b** Time evolution of the entanglement entropy.

For our initial state $|(210)^n\rangle$, the minimal cluster is defined as one that contains all the states given by tensor products of 210, 120 and 300 patterns. Similar to the H_3 case, within this approximation, triplets of sites again evolve independently, and the dimension of the reduced Hilbert space is $\dim \mathcal{H}^c = 3^{L/3}$. The time-evolved state within the cluster is given by

$$|\psi_n^c(t)\rangle = \cos^n(4t)|(210)^n\rangle + \dots , \tag{10}$$

where the dots denote other configurations. The fidelity is

$$F_n^c(t) = |\langle\psi_n^c(0)|\psi_n^c(t)\rangle|^2 = |\cos(4t)|^{2n}. \tag{11}$$

As in the case of H_3, this result is also valid for the translation-invariant initial state. We see that the period of revivals is $T_1 = \pi/4$, which is the same as for H_3 with a translation and inversion symmetric initial state.

The result of the cluster approximation is compared against the exact result for system size $L = 15$ in Fig. 5. The frequency of the fidelity revival, shown by the blue line in Fig. 5a, is accurately reproduced in this approximation, however the approximation does not capture the reduction in the magnitude of $F(t)$. Similarly, the dynamics of entanglement entropy, blue line in Fig. 5b, is only captured at very short times. In particular, we observe that the maximum entanglement within the cluster remains bounded even at long times $t \sim 10$, while the exact entropy continues to increase and reaches values that are several times larger.

To obtain a more accurate approximation, we can expand the minimal cluster with several neighboring configurations on the graph. We define the extended cluster as a set of all states which can be obtained using tensor products of the configurations 210, 120, 300, and 111. The enlarged cluster clearly contains the minimal cluster studied above, but it also includes additional configurations, resulting in a much better prediction for the first revival peak height, while still allowing for analytical treatment.

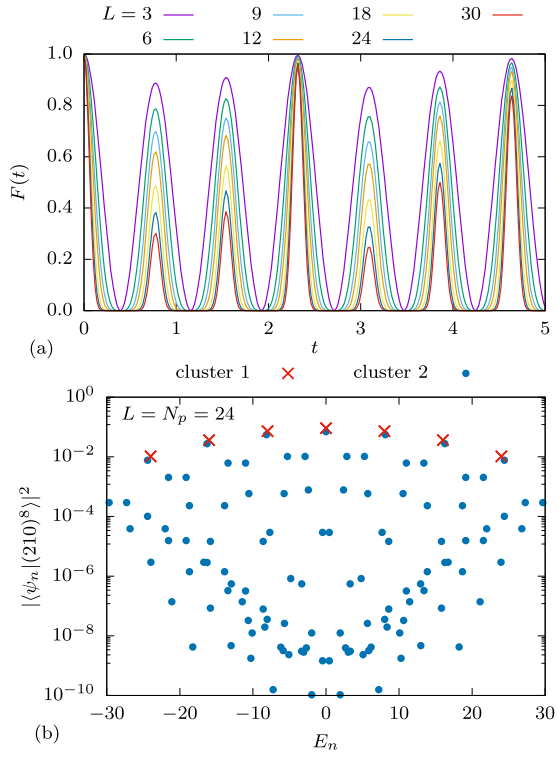

Fig. 6 Cluster approximations. a Fidelity $F(t)$, for the Hamiltonian H_1 and initial states $|(210)^n\rangle$, in the extended cluster approximation for various system sizes. **b** Eigenstate overlap with the initial state $|(210)^8\rangle$ plotted on a log scale, for both cluster approximations. In the case of degenerate eigenstates the sum of their overlaps is shown.

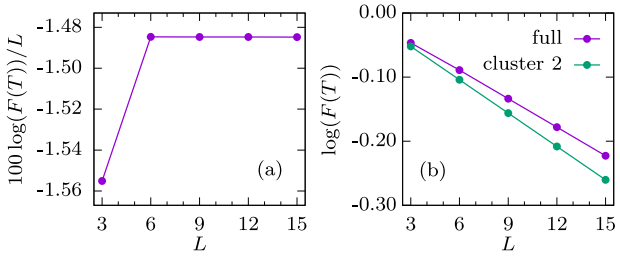

Fig. 7 First peak height. a Logarithm (base 10) of the first revival peak divided by the system size, $\log(\mathrm{F(T)})/\mathrm{L}$, seems to saturate at a finite value in the thermodynamic limit. **b** Comparison of the logarithm of the first revival peak height for the full dynamics and the improved cluster approximation. The approximation serves as a lower bound.

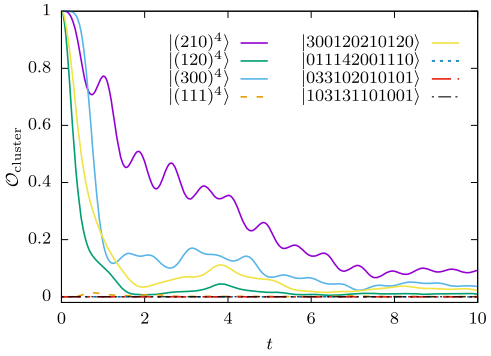

Fig. 8 Evolution of the probability to remain inside the minimal cluster. $\mathcal{O}_{\mathrm{Cluster}}$, as defined in Eq. (14). Initial configurations are indicated in the legend. Solid lines: configurations initially inside the cluster. Dashed lines: configurations initially outside the cluster (all except $|(111)^4\rangle$ are randomly chosen). Similar results are obtained for the extended cluster (not shown). System size $L = 12$.

The dimension of the extended cluster grows as $\dim \mathcal{H}^{\tilde{c}} = 4^{L/3}$, and is thus exponentially larger than the minimal cluster approximation. Nevertheless, the extended cluster dimension is still exponentially smaller compared to the full Hilbert space, and within this approximation it is possible to numerically simulate the dynamics of larger systems, $L \lesssim 30$—see Fig. 6a. The revivals are no longer perfect, while their frequency is independent of the system size and closer to the frequency of revivals for the full Hilbert space compared to to the minimal cluster approximation in Fig. 5. The overlap between the eigenstates of the Hamiltonian H_1 reduced to both the minimal and extended cluster and the state $|(210)^8\rangle$ is given in Fig. 6b. The eigenstates that correspond to the minimal cluster approximately survive in the extended cluster, where they form a band with the highest overlap.

For the initial product state $(210)^n$, it is possible to analytically obtain the fidelity within the improved approximation for arbitrary system size. Similar to the previous methods, it can be shown (see Supplementary Note 2)

$$F^{\tilde{c}}_{L=3n}(t) = 4^n |b^2 \cos(\alpha t) + d^2 \cos(\beta t)|^{2n}, \qquad (12)$$

where $\alpha = \sqrt{9+\sqrt{57}} \approx 4.06815$, $\beta = \sqrt{9-\sqrt{57}} \approx 1.20423$, $b \approx 0.694113$ and $d \approx 0.134933$. Eq. (12) is in excellent agreement with the numerical results in Fig. 6a. It was also found to be a very good approximation for the translation-invariant initial state when $L \geq 9$ (data not shown).

Figure 7a shows that the logarithm of the fidelity per site, $\log(\mathrm{F(T)})/\mathrm{L}$, at the first peak, saturates at a finite value for large L. In the improved cluster approximation, the first peak height decays as $\mathrm{e}^{-0.04L}$ (Supplementary Note 2). For a completely random state, the fidelity would be $F \sim 1/\dim_{\mathcal{H}}$. In the case $v = 1$ and large L, the Hilbert space dimension grows with the system size as

$$\dim_{\mathcal{H}} = \binom{2L-1}{L} \approx \binom{2L}{L} \approx \frac{4^L}{\sqrt{\pi L}}. \qquad (13)$$

This back-of-the-envelope estimate suggests the fidelity of a random state is $F \sim \mathrm{e}^{-1.39L}$, which decays considerably faster than the first peak height in Fig. 7. The improved cluster approximation correctly reproduces the short-time dynamics, including the first revival peak, and sets a lower bound for the first peak height – see Figs. 5 and 7b.

The evolution of the entanglement entropy for the extended cluster approximation is shown in Fig. 5b. Inside the cluster, entropy remains approximately constant with periodic oscillations that have the same frequency as the wave function revivals. Any further growth of the entanglement entropy can be attributed to the spreading of the wave function outside the cluster. To illustrate the "leakage" of the wave function outside the cluster, in Fig. 8 we compute the time evolution of the overlap with a cluster, i.e., the probability to remain inside a cluster at time t,

$$\mathcal{O}_{\mathrm{Cluster}} = \sum_{a \in \mathrm{Cluster}} |\langle a|\psi(t)\rangle|^2. \qquad (14)$$

We consider several initial configurations that lie inside or outside the cluster. The configurations initially inside the cluster mostly stay there, and the configuration $|(210)^4\rangle$ that has the highest revivals also has the highest overlap. Similarly, configurations initially outside the cluster continue to have negligible overlaps. The overlap starting from the configuration $|(210)^4\rangle$

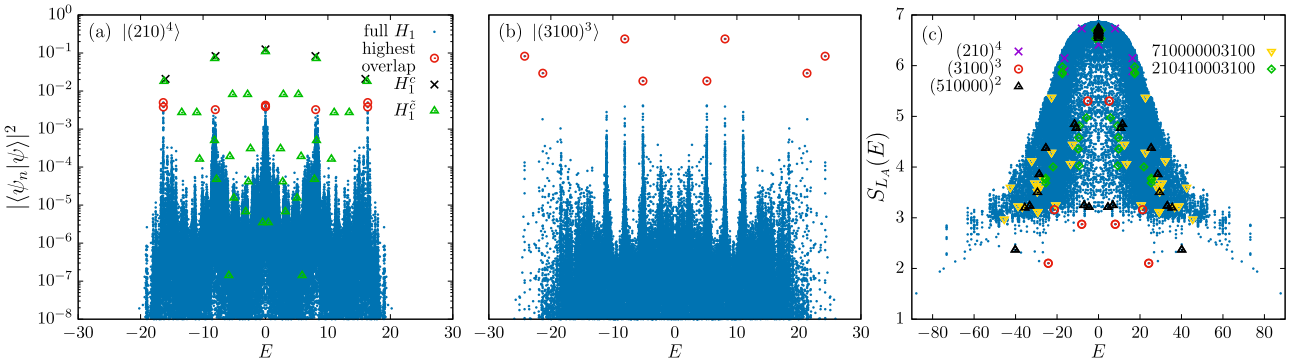

Fig. 9 Non-ergodic eigenstates. a Overlap of the configuration $|(210)^4\rangle$ with all the eigenstates of H_1, H_1^c and $H_1^{\tilde{c}}$ versus the eigenstate energy for sector $k = 0$ and system size $L = N_p = 12$. **b** Same for $|(3100)^3\rangle$. **c** Entanglement entropy, eigenstates which have the highest overlap with some product states are marked in different colors.

approximately predicts the revival peak heights for the full dynamics.

We now summarize the relation between H_3 and H_1 from the point of view of the cluster approximation. For the initial state $|(210)^n\rangle$, the two models yield similar dynamics, compare Eq. (23) and Eq. (12). The only difference is that the revival frequency is doubled in the latter case, which can be easily explained by the symmetry of the initial state and that of the Hamiltonian. Hamiltonian H_3 is inversion-symmetric. If the initial state is also chosen to be inversion-symmetric, the frequency of the revivals doubles. The period is then $T_3^{\rm inv} = \pi/4$, which is equal to the period of revivals T_1 of H_1 in the cluster approximation. This is also proven analytically in Methods, see Eq. (27). For comparison, the revival period for the full Hilbert space is approximately 0.77, which is slightly less than $\pi/4 \approx 0.79$. The Hamiltonian H_1 is not inversion-symmetric, so the frequency does not double for an inversion-symmetric initial state, but the revivals are lower in that case. This shows that it is important for the symmetry of the initial state to match the symmetry of the Hamiltonian.

Finally, the eigenstates of H_1, projected to the subspace of the minimal cluster approximation, form several degenerate bands whose eigenenergies are equally spaced in integer multiples of 4. Interestingly, some of these eigenstates approximately survive in the full H_1 model, and they are precisely the eigenstates that have the highest overlap with the configurations $|(210)^n\rangle$, see Fig. 9a. In small system sizes, such as $L = 6$, the surviving eigenstates are also the lowest entropy eigenstates in the middle of the spectrum, which is reminiscent of quantum scars in the PXP model[20]. In larger systems, e.g., $L = 12$, the surviving eigenstates are slightly lower in entropy than their neighbors, but are far from being the lowest entropy eigenstates, as can be seen in Fig. 9. The lowest entropy eigenstates have high overlaps with other configurations, such as $|(3100)^3\rangle$, as shown in Fig. 9b, c. In the case of $|(210)^n\rangle$, the eigenstates surviving in the full system belong to every other band of eigenstates in the cluster approximation and the number of the surviving eigenstates is $n + 1$. For even system sizes this counting includes a zero-energy eigenstate. In Methods we discuss in more details the generalization of the cluster approximations to the states of the form $|(N10...0)^n\rangle$, which were also found to have robust revivals and high overlaps with some low-entropy eigenstates.

Discussion

In this paper, we have introduced three models of bosons with "soft" kinetic constraints, i.e., density-dependent hopping. We have demonstrated that some of these models exhibit similar phenomenology to other realizations of quantum many-body scars, for example the Rydberg atom system[5]. We have studied quantum dynamics of these systems by performing global quenches from tensor-product initial states. We have shown that both the connectivity of the Hilbert space and the relative magnitude of the hopping coefficients have dramatic effects on the dynamics. For certain initial configurations, the constraints can lead to slow thermalization and revivals in the quantum fidelity. The revival frequency can be predicted by considering an exponentially reduced subset of the Hilbert space. For a family of initial configurations of the form $|(210)^n\rangle$, we have derived analytical expressions for the evolution of quantum fidelity within this approximation, which accurately capture the revival frequency obtained from exact numerical data. One notable difference between scarred dynamics in the present bosonic models and the PXP model is that the revivals exist in the absence of a hard kinetic constraint, i.e., in the fully connected Hilbert space. Our cluster approximation also explains the structure of some low-entropy eigenstates in the middle of the many-body spectrum. In addition, we have calculated the evolution of two local observables which are experimentally measurable, density correlations between two neighboring sites and density on a single site, and both of them show robust oscillations over a range of system sizes. We have also shown that the introduced models contain additional special properties, like the exponentially large zero-energy degeneracy which is related to the bipartite structure of the model.

We now comment on the possible experimental realizations of the models we studied. The implementation of a correlated hopping term $(n_k b_i^\dagger b_j)$ in optical lattices has attracted lot of attention due to a possible onset of quantum phases related to high-Tc superconductivity[80]. An early theoretical proposal exploits asymmetric interactions between the two atomic states in the presence of a state-dependent optical lattice[80]. As a result, the obtained effective model corresponds to the inversion-symmetric form of H_1. In addition, the same term has been found to feature as a higher-order correction of the standard Bose-Hubbard model[81–84]. Although in this case the term typically represents a modification of the regular hopping term of the order of several percent, its contribution was directly measured[85,86]. More recently, the set of quantum models accessible in cold-atom experiments has been enriched through the technique of Floquet engineering[87]. As a notable example, a suitable driving scheme can renormalize or fully suppress the bare tunneling rate[88]. On top of that, by modulating local interactions an effective model with the density-dependent tunneling term has been engineered[89]. For the models considered in this paper the most promising is a more recent driving scheme exploiting a double modulation of a local potential and on-site interactions[90]. Related

sophisticated driving schemes have already enabled a realization of dynamical gauge fields[56,57,91] where both the amplitude and the phase of the effective tunneling are density-dependent. Although these experimental proposals explain how to realize some of the correlated hopping terms present in our models using ultracold atoms in optical lattices, finding a scheme that exactly realizes our models requires further study. We emphasize that other models which would exhibit non-ergodic dynamics and scarred eigenstates as a result of the same mechanism that was explained in this work could be built, for example a linear combination of H_1 and H_2.

Note added: During the completion of this work, we became aware of ref. [92] which identified non-thermal eigenstates and slow dynamics in the quantum East model. Moreover, a recent study[93] proposed a Floquet scheme for a bosonic model with density-assisted hopping, finding signatures of quantum many-body scars.

Methods

In order to more efficiently describe the dynamics of our models, we introduce a method—"cluster approximation", that is based on Hilbert space truncation inspired by the bipartite graph structure of H_1. Before providing details about the cluster approximation for H_1 and its generalizations, we present an exact solution for the perfect revivals in H_3 model, which serves as a motivation for the more complicated case of H_1.

Bipartite lattice and zero modes. The graph of H_1 is bipartite, i.e. all the basis configurations can be divided into two disjoint sets, and the action of the Hamiltonian connects configurations in one set only to the configurations in the other and vice-versa (the Hamiltonian is off-diagonal). One way to sort configurations into these two sets is by parity of the quantity

$$\Delta_a = \frac{|n_{\rm even} - n_{\rm odd} + C|}{2}, \tag{15}$$

where $C = 0$ if L is even and $C = 1$ if L is odd. We define $n_{\rm even}$ and $n_{\rm odd}$ as the total numbers of particles at even and odd sites, respectively,

$$n_{\rm even} = \sum_{l=1}^{L_1} n_{2l},\ n_{\rm odd} = \sum_{l=1}^{L_2} n_{2l-1}, \tag{16}$$

where $L_1 = L_2 = L/2$ if L is even, and $L_1 = (L-1)/2$, $L_2 = (L+1)/2$ if L is odd. If only nearest-neighbor hoppings are allowed and if no two odd sites are coupled (if the system has open boundary conditions for any L or periodic boundary conditions for L-even), each hopping either increases $n_{\rm even}$ by one and decreases $n_{\rm odd}$ by one, or vice-versa. This means that each hopping can change Δ_a only by ± 1.

In special cases, like H_1 at filling factor $\nu = 1$, it is also possible to define quantities like Δ_a for odd system sizes and periodic boundary conditions. This is a consequence of the constraints imposed by H_1, i.e., the fact that a particle cannot hop to an empty site to its left (Supplementary Note 3). Note that H_2 in the same geometry is not bipartite.

Another way to sort configurations into two sets is by parity of the distance from the configuration $|111...111\rangle$, which we define as

$$d_a = \min_n\{\langle 111...111|H_1^n|a\rangle \neq 0\}. \tag{17}$$

In this case, the two sets are the configurations with even and with odd distances d_a. One hopping can change d_a only by ± 1 or 0. Changes by other values are not possible by definition if the Hamiltonian is Hermitian (all hoppings are reversible). Both d_a and Δ_a have the same parity, thus d_a must always change after one hopping in even system sizes or in systems with open boundary conditions. As a consequence, d_a cannot change by 0 if Δ_a can only change by ± 1.

The graphs of bipartite systems do not contain any loops of odd dimension (triangles, pentagons, heptagons and so on). Moreover, the energy spectra of bipartite systems are symmetric around zero. Their Hamiltonians anticommute with the operator $(-1)^{\Delta_a}$. The presence of such an operator in a bipartite lattice leads to exact zero energy states in the spectrum[94,95]. It can be shown that the exponentially growing number of zero modes of H_1 is related to the difference between the numbers of elements in the two sets of its bipartite graph (Supplementary Note 4). Additionally, the algebraic structure of zero energy eigenstates can be explained by the structure of the graph – such eigenstates can be constructed as superpositions of configurations from only one of the sets. Similar properties are found for H_2 for even L, as its graph is also bipartite in that case. The properties of the zero-energy manifold are discussed in more detail in Supplementary Note 4.

Perfect revivals in the H_3 model. We start with a warmup calculation for H_3 acting on $L = 3$ sites. The connected subspace of 210 contains only two configurations, 120 and 210. The Hamiltonian reduced to this subspace is

$$H_3' = -\begin{pmatrix} 0 & 2 \\ 2 & 0 \end{pmatrix}, \tag{18}$$

where the basis vectors are

$$\begin{pmatrix} 1 \\ 0 \end{pmatrix} = |210\rangle,\ \begin{pmatrix} 0 \\ 1 \end{pmatrix} = |120\rangle. \tag{19}$$

The eigenvalues of H_3' are $E_1 = -2$ and $E_2 = 2$. The initial state $|\psi_1(t=0)\rangle = |210\rangle$ evolves as

$$|\psi_1(t)\rangle = \cos(2t)|210\rangle - i\sin(2t)|120\rangle, \tag{20}$$

and the state $|\psi_2(t=0)\rangle = |120\rangle$ evolves as

$$|\psi_2(t)\rangle = -i\sin(2t)|210\rangle + \cos(2t)|120\rangle. \tag{21}$$

Previous results can be straightforwardly generalized to larger systems. Let the length of the system be $L = 3n$ for simplicity. The connected component of the state $|(210)^n\rangle$ consists only of combinations of patterns 210 and 120, which means that triplets of sites evolve independently. From Eq. (20), the initial state $|\psi_n(t=0)\rangle = |(210)^n\rangle$ evolves as

$$\begin{aligned}|\psi_{L=3n}(t)\rangle =& \cos^n(2t)|(210)^n\rangle \\ &+ (-i)^n\sin^n(2t)|(120)^n\rangle + ...\end{aligned} \tag{22}$$

where ". . . " denotes contributions of the basis configurations other than $|(210)^n\rangle$ or $|(120)^n\rangle$. The fidelity is

$$F_{L=3n}(t) = |\langle\psi_n(0)|\psi_n(t)\rangle|^2 = |\cos 2t|^{2n}. \tag{23}$$

It follows that the revivals are perfect, with a period $T_3 = \pi/2$. This result is also valid for the translation-invariant initial state $|(210)^n\rangle_{\rm T}$,

$$|(210)^n\rangle_{\rm T} \equiv \frac{1}{\sqrt{3}}(|(210)^n\rangle + |(021)^n\rangle + |(102)^n\rangle), \tag{24}$$

as the connected subspaces of 210, 021 and 102 do not overlap and therefore evolve independently.

However, an initial state that is both translation symmetric and inversion symmetric has different dynamics. The inverse of the configuration $|(210)^n\rangle$ is the configuration $|(012)^n\rangle$, which is a translation of the state $|(120)^n\rangle$ that belongs to the connected subspace of $|(210)^n\rangle$. The initial state

$$|\psi_n^{\rm inv}(t=0)\rangle = \frac{1}{\sqrt{2}}|(210)^n\rangle_{\rm T} + \frac{1}{\sqrt{2}}|(120)^n\rangle_{\rm T} \tag{25}$$

evolves as

$$|\psi_n^{\rm inv}(t)\rangle = (\cos^n 2t + (-i)^n\sin^n 2t)|\psi_n^{\rm inv}(t=0)\rangle + ... \tag{26}$$

and the fidelity is

$$\begin{aligned}F_n^{\rm inv}(t) =& |\langle\psi_n^{\rm inv}(0)|\psi_n^{\rm inv}(t)\rangle|^2 \\ =& |\cos^n 2t + (-i)^n\sin^n 2t|^2.\end{aligned} \tag{27}$$

The frequency of the revivals is now doubled, so the period is $T_3^{\rm inv} = \pi/4$.

Cluster approximations for the H_1 model. Here we introduce a scheme for approximating the dynamics from initial states $(210)^n$ in the H_1 model. As can be observed in Fig. 3, the revival periods are approximately the same for different system sizes. We first focus on the non-trivial case $L = 6$. Figure 10 shows part of the graph that contains the initial state, $|210210\rangle$. Configurations labeled inside the ellipses denote representatives of an orbit of translation symmetry, i.e., the configurations are translation-invariant such as the one in Eq. (24).

The minimal subcluster of the graph is highlighted in blue color in Fig. 10. This cluster is indeed weakly connected to the rest of the configuration space, as it has only 3 connections that lead outside this cluster (dashed lines) and their hopping coefficients are slightly lower in magnitude than those inside the cluster, meaning that the probability is higher to stay inside the cluster than to leave. The hopping coefficients leading outside are not significantly smaller than the coefficients staying inside, but in combination with the relatively small number of connections this has significant effects on the dynamics. This effect is even more pronounced when the difference in magnitudes is further increased by squaring the particle-number operators (see Supplementary Note 1).

The minimal cluster from Fig. 10 contains all the states given by tensor products of 210, 120 and 300 configurations. The set of configurations belonging to this cluster could have been chosen differently, but this particular choice has at least two advantages. Firstly, inside this cluster, the evolution of the configuration $|210210\rangle$ can be thought of as two subsystems 210 evolving separately. The evolution of all such configurations at different system sizes can be reduced to the evolution of $L = 3$ subsystems 210, similar to the case of H_3 in the connected subspace of $(210)^n$. Secondly, this definition allows easy generalization to different system sizes $L = 3n$ with initial states $(210)^n$. We would like to emphasize that the cluster was not chosen arbitrarily. The calculations of the probability density distribution starting from the initial configuration $|210210\rangle$ and evolving with H_1

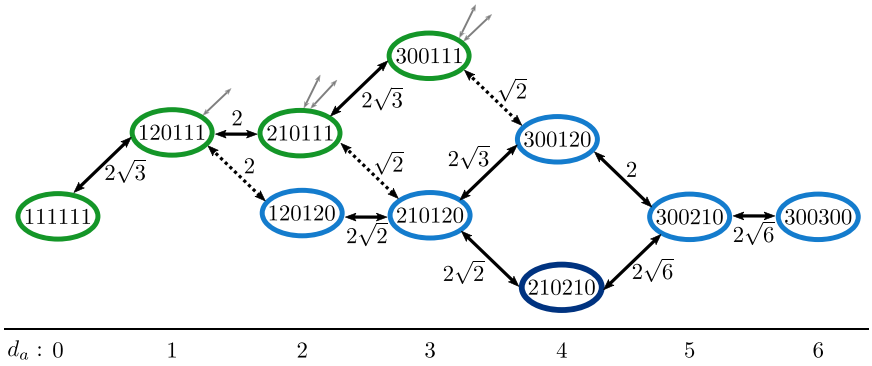

Fig. 10 Minimal and extended clusters. Hamiltonian H_1 and system size $L = N_p = 6$. Configurations labeled inside the ellipses are representatives of an orbit of translation symmetry. The minimal cluster is defined by the blue configurations, while green configurations represent the additional components of the extended cluster. Gray arrows connect to configurations outside the extended cluster. The numbers bellow the graph show the distance d_a from the configuration 111111 evaluated using Eq. (17).

have shown that the probability density stays high in this region of the Hilbert space as long as the revivals in fidelity are visible. The configurations important for the dynamics were then identified by analyzing the structure of the graph around the initial configuration.

As an example, consider system size $L = 3$. The reduced Hilbert space of the cluster $\mathcal{H}^c$ is spanned by the (non-translation-invariant) configurations

$$\begin{pmatrix}1\\0\\0\end{pmatrix} = |300\rangle, \quad \begin{pmatrix}0\\1\\0\end{pmatrix} = |210\rangle, \quad \begin{pmatrix}0\\0\\1\end{pmatrix} = |120\rangle. \tag{28}$$

The Hamiltonian reduced to this subspace is

$$H_1^c = -\begin{pmatrix}0 & 2\sqrt{3} & 0\\ 2\sqrt{3} & 0 & 2\\ 0 & 2 & 0\end{pmatrix}, \tag{29}$$

and its eigenvalues are $E_1 = -4$, $E_2 = 4$, $E_3 = 0$. The initial configuration $|210\rangle$ evolves according to

$$|\psi_1^c(t)\rangle = -\frac{i}{2}\sin(4t)\left(\sqrt{3}|300\rangle + |120\rangle\right) + \cos(4t)|210\rangle. \tag{30}$$

By generalizing this result to larger systems, it is easy to prove Eqs. (10) and (11).

The minimal clusters can be expanded by adding several neighboring configurations. For similar reasons as in the case of minimal clusters, the extended clusters are defined as sets of all states which can be obtained using tensor products of the configurations 210, 120, 300 and 111. In the case of $L = 6$, the enlarged cluster can be observed in Fig. 10. It contains the minimal cluster studied previously, but it also includes additional configurations shown in green ellipses. Again, the approximation could be improved by including more configurations, but this particular choice is well suited for analytical treatment (Supplementary Note 2) and, as shown above, it gives a good prediction for the first revival peak height.

Generalization to other clusters. Building on the previous observation that some of the low-entropy eigenstates have large weight on $|(3100)^3\rangle$ product state, we have investigated periodic revivals from such a larger class of initial states. We find that robust revivals are associated with initial product states of the form

$$\left|((N-1)1\underbrace{0\ldots0}_{N-2})^n\right\rangle, \tag{31}$$

where N is the length of the unit cell ($L = Nn$). For example, some of these configurations are $|(3100)^n\rangle$, $|(41000)^n\rangle$ and $|(510000)^n\rangle$. Combinations of those patterns such as $|310041000\rangle$ also exhibit similar properties, but we will restrict ourselves to the simpler former cases.

We can construct a generalization of the cluster approximation for configurations of the form in Eq. (31). As in the case of $|(210)^n\rangle$, the dynamics inside one unit cell explains the dynamics of the full system. The generalized clusters can be chosen in such a way that their Hilbert spaces are spanned by N configurations

$$|i\rangle = \left|((N+1-i)(i-1)\underbrace{0\ldots0}_{N-2})^n\right\rangle, \tag{32}$$

where i takes values 1, 2, …N. If we consider only one unit cell ($n = 1$), the graph that connects these configurations has a linear structure without any loops, i.e., each configuration $|i\rangle$ is solely connected to the configurations $|i \pm 1\rangle$, except the two configurations at the edges, $|1\rangle$ and $|N\rangle$, which are only connected to $|2\rangle$ and $|N-1\rangle$, respectively.

The projection of the Hamiltonian H_1 to this cluster, which we denote by H_1^c, has a very simple structure: it has the form of a tight-binding chain with the only nonzero matrix elements on the upper and lower diagonals:

$$H^c_{1;i,i+1} = H^c_{1;i+1,i} = (N-i)\sqrt{i(N+1-i)}. \tag{33}$$

The dynamics within a single unit cell under H_1^c corresponds to density fluctuations between the first and the second site. Following the same procedure as previously, we can now diagonalize H_1^c and compute the fidelity time series for the initial configuration $|(N-1)10\ldots0\rangle$. This result can be directly generalized to configurations of the form $|((N-1)10\ldots0)^n\rangle$. The derivation is valid for both translation-invariant and non-translation-invariant initial configurations, as the cluster in Eq. (32) is disconnected from its translated copies. We stress that this disconnection, namely the absence of a hopping term between $|1(N-1)0\ldots0\rangle$ and $|0N0\ldots0\rangle$, is a consequence of the constraints imposed by H_1 and it would not hold for H_2. In this way, we have calculated the time evolution of the fidelity starting from the configurations $|(3100)^n\rangle$ (for $n = 1, 2, 3, 4$), $|(41000)^n\rangle$ ($n = 1, 2, 3$) and $|(510000)^n\rangle$ ($n = 1, 2$), and compared it with the exact numerical results for the full H_1. The cluster approximation captures both the revival frequency and the height of the first peak. Similar to the $|(210)^n\rangle$ case, the approximation can be improved by adding further configurations to the clusters. Moreover, if we want to consider translation-invariant initial states, we can follow the same recipe for $|(210)^n\rangle$ by summing translated patterns with the required phase factors given in terms of momenta in multiples of $2\pi/N$. We have checked that revivals appear in these momentum sectors, with roughly the same frequency.

We note that the configurations with larger units cells thermalize more quickly on shorter timescales, but slower at long times. Initially, the states starting from configurations with smaller N have lower entanglement entropies than those with larger N. The Hilbert spaces of large N unit cells are larger, so the entanglement entropy starting from these configurations rapidly grows to the maximal value for that unit cell. However, the only way for the wave function to spread through the entire Hilbert space is that a unit cell reaches a state close to 111. . . 111, so that particles can hop to the other unit cells. This is less likely for large N, and therefore such configurations need long times to fully thermalize. As a result, smaller N entanglement entropies grow faster and after long enough time they overtake those for larger N. For example, in the case of $L = 12$ and translation-invariant initial states, $(210)^4$ overtakes $(3100)^3$ and $(510000)^2$ around $t \sim 2$, and $(3100)^3$ overtakes $(510000)^2$ around $t \sim 80$ (Supplementary Fig. 3).

Finally, we note that non-thermal behavior reminiscent of the one studied here was previously observed in an unconstrained Bose-Hubbard model, for example in the context of "arrested expansion"[96,97] and quenches from superfluid to Mott insulator phase[98,99]. In these cases, the main ingredient is the strong on-site interaction, which causes the energy spectrum to split into several bands. Due to the large energy differences between bands, the dynamics of an initial state from a particular band is at first limited only to the eigenstates that belong to the same band. Additionally, these energy bands are approximately equally spaced, which can lead to revivals in fidelity if several bands are populated. In contrast, our models do not feature on-site interaction, and the mechanism which slows down the spread of the wave function is correlated hopping, which suppresses connections between certain configurations and modifies the hopping amplitudes between others, thus creating bottlenecks that separate different clusters of states.

Data availability

The data that support the plots within this paper and other findings of this study are available at https://doi.org/10.5518/810[100].

Code availability

Code is available upon reasonable request.

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Author contributions

A.H., I.V., N.R., and Z.P. contributed to developing the ideas, analyzing the results and writing the manuscript. A.H. performed the calculations and designed the figures.

Competing interests

The authors declare no competing interests.

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Correspondence and requests for materials should be addressed to A.H.

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the Feshbach resonance.

Poster Q 53.10 Thu 17:00 P OG2

Spin and Charge Correlation Measurements in the 2D Hubbard Model — Jan Drewes[1], Luke Miller[1,2], Eugenio Cocchi[1,2], Chun Fai Chan[1], Nicola Wurz[1], •Marcell Gall[1], Daniel Pertot[1], Ferdinand Brennecke[1], and Michael Köhl[1] — [1]Physikalisches Institut, University of Bonn, Wegelerstrasse 8, 53115 Bonn, Germany — [2]Cavendish Laboratory, University of Cambridge, JJ Thomson Avenue, Cambridge CB3 0HE, United Kingdom

We experimentally study the emergence of correlations in an ultracold, fermionic 2D lattice system, representing a realisation of the Hubbard model. Our ability to precisely tune the system parameters over a large range and the possibility to simultaneously detect the density distribution of both spin components in-situ enables us to examine the emergence of density and spin correlations as a function of doping interaction strength and temperature. In addition we gain from the measurement of the equation of state insight into the full thermodynamics of the 2D Hubbard model. To improve our preparation and detection capabilities, we use a spin spiral technique which allows us to detect the spin structure factor at arbitrary wave vectors. Further we employ a spatial light modulator to reshape the underlying trapping potential of the optical lattice to realize the homogeneous Hubbard model and reach lower temperatures by redistributing entropy between different spatial regions.

Poster Q 53.11 Thu 17:00 P OG2

BEC of ^{41}K in a Fermi Sea of ^{6}Li — Rianne S. Lous[1,2], Isabella Fritsche[1,2], •Fabian Lehmann[1,2], Michael Jag[1,2], Emil Kirilov[1,2], Bo Huang[1], and Rudolf Grimm[1,2] — [1]IQOQI, Austrian Academy of Science, Innsbruck, Austria — [2]Inst. for Experimental Physics, University of Innsbruck, Innsbruck, Austria

We report on the production of a double-degenerate Fermi-Bose mixture of ^{6}Li and ^{41}K. In our experimental sequence the potassium atoms are sympathetically cooled by the lithium atoms, which are evaporatively cooled in an optical dipole trap. We obtain 10^4 ^{41}K atoms with a BEC fraction close to 1 and a $T/T_F \approx 0.05$ with 10^5 ^{6}Li atoms in each spin state. To measure the temperature of our fermionic sample we use the ^{41}K BEC as a tool for thermometry. As the system is in thermal equilibrium we evaluate the condensed fraction of our ^{41}K atoms and extract the temperature of the atoms. To investigate the properties of the ^{6}Li-^{41}K mixture near the inter-species Feshbach resonance at 335.8 G we use another scheme of evaporation around 300 G which enables us to achieve similar temperatures. We explore both the repulsive side and attractive side of the Feshbach resonance and observe phase separation for strong repulsive interactions and collapse for attractive interactions. This work is supported by the Austrian Science Fund FWF within the SFB FoQuS.

Poster Q 53.12 Thu 17:00 P OG2

Probing Many-body physics with an ultra-narrow clock transition in an Ytterbium quantum gas — •Bodhaditya Santra[1], Benjamin Abeln[1], Bastian Hundt[1], André Kochanke[1], Thomas Ponath[1], Anna Skottke[1], Klaus Sengstock[1,2], and Christoph Becker[1,2] — [1]Zentrum für Optische Quantentechnologien, Universität Hamburg, Luruper Chaussee 149, 22761 Hamburg, Germany — [2]Institut für Laserphysik, Universität Hamburg, Luruper Chaussee 149, 22761 Hamburg, Germany

During the last decade ultracold fermionic alkaline earth quantum gas attracted a lot of attention due to their unique properties such as long-lived meta-stable state, an ultra-narrow optical clock transition, SU(N) symmetric interactions as well as the existence of an interorbital Feshbach resonance. In particular fermionic Yb quantum gas allow for quantum simulation of lattice systems with orbital degrees of freedom, like the Kugel-Khomskii model or the Kondo lattice model (KLM).

We will present recent progress of the Hamburg Yb experiment towards realizing the KLM and correlated KLM, including measurements on spin polarized as well as on interacting Fermi gases with an improved clock laser setup.

This work is supported by the DFG within the SFB 925 and the Marie Curie Initial Training Network QTea.

Poster Q 53.13 Thu 17:00 P OG2

Local control of transport in an atomic quantum wire: from one scanning gate to a finite size lattice — •Samuel Häusler[1], Martin Lebrat[1], Dominik Husmann[1], Laura Corman[1], Sebastian Krinner[1], Shuta Nakajima[2], Jean-Philippe Brantut[1], and Tilman Esslinger[1] — [1]Institute for Quantum Electronics, ETH Zürich, 8093 Zürich, Switzerland — [2]Department of Physics, Graduate School of Science, Kyoto University, Kyoto 606-8502, Japan

Building on the holographic shaping of optical potentials and a high-resolution microscope, we demonstrate the local control of fermionic lithium atoms flowing through a one-dimensional structure. We first image the transport through a quantum wire, in a way similar to the scanning gate technique applied to solid state devices. By scanning the position of a sharp, repulsive optical gate over the wire and measuring the subsequent variations of conductance, we spatially map the transport at a resolution close to the transverse wavefunction inside the wire. The control of the gate at the scale of the Fermi wavelength makes it sensitive to quantum tunnelling. Furthermore, our knowledge of the optical potential allows a direct comparison of the experimental maps with a numerical and an analytical model for non-interacting particles.

The flexibility offered by our setup makes it relatively simple to imprint more complex structures. By projecting several consecutive scatterers, a lattice of variable length can be built inside the quantum wire. This opens the path to study metal-insulator physics with strong attractive interactions.

Poster Q 53.14 Thu 17:00 P OG2

Interacting Anyons in a One-Dimensional Optical Lattice — •Martin Bonkhoff, Kevin Jägering, Sebastian Eggert, and Axel Pelster — State Research Center OPTIMAS and Fachbereich Physik, Technische Universität Kaiserslautern, 67663 Kaiserslautern, Germany

We analyze in detail the properties of the one-dimensional Anyon-Hubbard model, which can be mapped to a corresponding Bose-Hubbard model with a density-dependent Peierls phase via a generalized Jordan-Wigner transformation [1]. At first we extend the modified version of the classical Gutzwiller-mean-field ansatz of Ref. [2] in order to obtain the pair-correlation function for both the bosonic and the anyonic system. A comparison of the resulting quasi-momentum distributions with high-precision DMRG calculations reveals in general a parity breaking, which is due to anyonic statistics. Afterwards, we determine how the boundary of the superfluid-Mott quantum phase transition changes with the statistical parameter. We find in accordance with Ref. [1] that the statistical interaction has the tendency to destroy superfluid coherence.

[1] T. Keilmann, S. Lanzmich, L. McCulloch, and M. Roncaglia, Nat. Commun. **2**, 361 (2011)

[2] G. Tang, S. Eggert, and A. Pelster, New J. Phys. **17**, 123016 (2015)

Poster Q 53.15 Thu 17:00 P OG2

Creating topological interfaces and detecting chiral edge modes in a two-dimensional optical lattice — •Frederik Görg[1], Nathan Goldman[2], Gregor Jotzu[1], Michael Messer[1], Kilian Sandholzer[1], Rémi Desbuquois[1], and Tilman Esslinger[1] — [1]Institute for Quantum Electronics, ETH Zurich, Zurich, Switzerland — [2]CENOLI, Université Libre de Bruxelles, Brussels, Belgium

The appearance of topological properties in lattice systems caused by a non-trivial topological band structure in the bulk is closely related to the existence of chiral edge modes via the bulk-edge correspondence. These edge states appear at the interface of two spatial regions with a distinct topology, which for example naturally arise at the boundaries of a sample surrounded by vacuum. In cold atom systems, these edge modes are difficult to detect, since the underlying harmonic trapping potential does not feature sharp boundaries. Therefore, we propose a different method to design topological interfaces within the bulk of the system. We illustrate this scheme by an optical lattice realization of the Haldane model, where a spatially varying lattice beam leads to the appearance of distinct topological phases in separated regions of space. The versatility of the method allows to tune the position, the localization length and the chirality of the edge modes. We numerically study the propagation of wave packets in such a system and demonstrate the feasibility to experimentally detect chiral edge states. Finally, we show that the edge modes, unlike the bulk states, are topologically protected against the effects of disorder, which makes a random potential a powerful tool to detect edge states in cold atom setups.

Poster Q 53.16 Thu 17:00 P OG2

Transport dynamics in optical lattices with flux — •Ana Hudomal[1], Ivana Vasić[1], Walter Hofstetter[2], and Antun Balaž[1] — [1]Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia — [2]Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität, Frankfurt am Main, Germany

Recent cold atom experiments have realized artificial gauge fields in periodically modulated optical lattices [1,2]. We study the dynamics of atomic clouds in these systems by performing numerical simulations using the full time-dependent Hamiltonian and comparing these results to the semiclassical approximation. Under constant external force, atoms in optical lattices with flux exhibit an anomalous velocity in the transverse direction. We investigate in detail how this transverse drift is related to the Berry curvature and Chern number, taking into account realistic experimental conditions.

[1] G. Jotzu et al., Nature **515**, 237 (2014).
[2] M. Aidelsburger et al., Nature Phys. **11**, 162 (2015).

Poster Q 53.17 Thu 17:00 P OG2

Towards the investigation of collective scattering in nanofiber-trapped atomic ensembles — •Adarsh S. Prasad, Jakob Hinney, Samuel Rind, Philipp Schneeweiss, Jürgen Volz, Christoph Clausen, and Arno Rauschenbeutel — TU Wien - Atominstitut, Stadionallee 2, 1020 Wien, Austria

We realize an efficient optical interface between guided light and laser-cooled atoms which are arranged in two linear arrays in a two-color evanescent-field dipole trap created around an optical nanofiber [1]. In this configuration, the probability of a nanofiber-guided photon being absorbed and then re-emitted into free space by a trapped atom is as high as 10%. For a periodic array of atoms, interference of the fields scattered by different atoms result in a collective emission into a cone with a well-defined angle with respect to the fiber axis. We plan to study this collective emission and its dependence on various experimental parameters. The next step will be to adjust the periodicity of the atomic array to fulfill the Bragg condition such that fiber-guided light is strongly back-reflected [2]. Here, the interaction between the atomic array and the fiber-guided light depends strongly on the polarization of the light field. In particular, light that is polarized in (orthogonal to) the plane of atoms will be weakly (strongly) reflected. We want to implement such highly reflecting atomic arrays, which could then be used to implement cavity quantum electrodynamics experiments in which the resonator itself is made of quantum emitters.

[1] E. Vetsch et al., Phys. Rev. Lett. 104, 203603 (2010).
[2] Fam Le Kien et. al., Phys. Rev. A 90, 063816 (2014).

Poster Q 53.18 Thu 17:00 P OG2

Setup of a new micro-structured linear Paul trap with integrated solenoids and reduced axial micromotion — •H. Siebeneich, D. Kaufmann, T. Gloger, P. Kaufmann, M. Johanning, and ch. Wunderlich — Department Physik, Universität Siegen, 57068 Siegen, Germany

We present the status of a new 3d segmented ion trap setup with integrated solenoids, in which an improved design allows for a substantial reduction of axial micromotion and for an increased magnetic gradients. Our trap consists of three layers of gold plated alumina, where the segmented outer layers provide the trapping potentials [1], and the middle layer contains solenoids that are used to create a magnetic field gradient [2]. The gradient gives rise to coupling between the ions' internal and motional states. The trap is mounted on a ceramic chip carrier that, at the same time, acts as an ultra-high vacuum interface, featuring about 100 thick-film printed current and voltage feedthroughs. The thick film interface has been improved by replacing previously used Ag-Pd layers by Au layers which reduced their resistivity by a factor of eight. The previously high resistivity used to be a a bottleneck for achieving high solenoid currents and thus a magnetic gradient. The shape of the solenoids was redesigned, leading to an expected reduction of axial micromotion by four orders of magnitudes.

[1] S.A. Schulz et al.: Sideband cooling and coherent dynamics in a microchip multi-segmented ion trap, New Journal of Physics, Volume 10, April 2008 [2] D. Kaufmann et al.: Thick-film technology for ultra high vacuum interfaces of micro-structured traps, Appl Phys B (2012) 107:935-943

Poster Q 53.19 Thu 17:00 P OG2

Design and construction of a Perpetual Atom Laser Machine — •Chun-Chia Chen, Shayne Bennetts, Benjamin Pasquiou, and Florian Schreck — Institute of Physics, University of Amsterdam, Amsterdam, The Netherlands

We have developed a machine aimed at producing a perpetual atom laser, a long standing goal within atomic physics. Continuous production of Bose-Einstein condensate (BEC) or an atom laser requires two incompatible cooling processes, laser cooling a gas sample, then cooling evaporatively until degeneracy is reached. In order to produce a perpetual output these stages take place simultaneously in different parts of our machine. To protect the condensate from scattered photon heating we use a combination of physical separation, baffles and a "transparency" beam. Our machine has now demonstrated a perpetual MOT of 2×10^9 ^{88}Sr atoms with temperatures as low as 20μK on a 7.4-kHz wide laser cooling transition with a continuous loading rate of 7×10^8 atoms/s. Using a different set of parameters and location we have also demonstrated a perpetual MOT of 2×10^8 ^{88}Sr at 2μK with a loading rate of 9×10^7 atoms/s which we have successfully loaded into a dipole trap. By switching to the 0.5% abundance ^{84}Sr isotope we are able to evaporate to BECs of 3×10^5 ^{84}Sr atoms. Critically, for the second location we have validated the effectiveness of our architecture in protecting a BEC from scattered broad-linewidth laser cooling light, which is used in the first cooling stages. We will describe our design and the performance demonstrated so far.

Poster Q 53.20 Thu 17:00 P OG2

Optical trapping of neutral mercury — •Holger John and Thomas Walther — Technische Universität Darmstadt, Institut für Angewandte Physik, Schlossgartenstraße 7, 64289 Darmstadt

Laser-cooled mercury constitutes an interesting starting point for various experiments, in particular in light of the existence of bosonic and fermionic isotopes. On the one hand the fermionic isotopes could be used to develop a new time standard based on an optical lattice clock employing the $^1S_0$ - $^3P_0$ transition. Another interesting venue is the formation of ultra cold Hg-dimers employing photo-association and achieving vibrational cooling by employing a special scheme.

The laser system is based on an interference-filter stabilized external cavity diode laser with excellent spectral properties combined with a home built non-cryogenic fiber amplifier for the 1015nm fundamental wavelength with a slope-efficiency of more than 35 % delivering up to 4W of pump limited output power. The fundamental wavelength is frequency doubled twice to reach the cooling transition at 253.7 nm. The challenging requirements meeting the natural linewidth of 1.27 MHz are mastered by use of a ULE reference resonator.

After integrating a 2D-MOT as an atom source to the vacuum system the first measurements of ultra-cold atoms with the new laser system will be reported.

Poster Q 53.21 Thu 17:00 P OG2

Diffusion of Single Atoms in Bath — •Daniel Adam, Farina Kindermann, Tobias Lausch, Daniel Mayer, Felix Schmidt, Steve Haupt, Michael Hohmann, Nicolas Spethmann, and Artur Widera — TU Kaiserslautern, Department of Physics, Kaiserslautern, Germany

Diffusion is an essential phenomenon occurring in various systems such as biological cells, traffic models or stock markets. While most systems are well described by standard Brownian motion, anomalous diffusion can lead to markedly different dynamical properties.

Experimentally, we study the diffusion of individual atoms illuminated by near-resonant light and trapped in a periodic potential. All relevant parameters such as damping coefficient and potential hight can be controlled in order to realize different diffusive regimes.

We explore the amount of information contained in the Kramers rate, i. e. the rate at which a diffusing atom can escape from a potential well. Furthermore we exploit the excellent control over the optical trapping potential and study the diffusion of the atom in a time-varying periodic trap, complemented by numerical simulations of the dynamics.

Poster Q 53.22 Thu 17:00 P OG2

Kinetic Monte Carlo simulation of percolation in driven-dissipative Rydberg gases — •Stephan Helmrich, Philipp Fabritius, Graham Lochead, and Shannon Whitlock — Physikalisches Institut, Universität Heidelberg, Im Neuenheimer Feld 226, 69120 Heidelberg

Directed percolation is perhaps the most prominent example of a unique class of phenomena which exhibit genuine non-equilibrium phase transitions and non-trivial critical behaviour. We explore whether highly tunable gases of ultracold atoms excited to long-range interacting Rydberg states can serve as a clean experimental realisation of percolation phenomena in two and three dimensions. The mechanism investigated is the cooperative excitation of Rydberg atoms triggered when the excitation laser is resonant for atoms within a characteristic distance of another Rydberg atom (facilitated excitation). To simulate the dynamics of this system we use a kinetic Monte Carlo algorithm which is able to reproduce many of the experimental features of laser excited Rydberg gases. We investigate the scaling behavior for the fraction of Rydberg excitations (active sites) and their spa-

Transport dynamics in optical lattices with flux

A. Hudomal[1], I. Vasić[1], H. Buljan[2], W. Hofstetter[3], and A. Balaž[1]
[1]*Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia*
[2]*Department of Physics, University of Zagreb, Croatia*
[3]*Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität, Frankfurt am Main, Germany*
e-mail:ana.hudomal@ipb.ac.rs

Recent cold atom experiments have realized artificial gauge fields in periodically modulated optical lattices [1,2]. We study the dynamics of atomic clouds in such systems by performing numerical simulations using the full time-dependent Hamiltonian and compare results with the semiclassical approximation. Under constant external force, atoms in optical lattices with flux exhibit an anomalous velocity in the transverse direction. We investigate in detail how this transverse drift is related to the Berry curvature and Chern number, taking into account realistic experimental conditions.

REFERENCES

[1] G. Jotzu, M. Messer, R. Desbuquois, M. Lebrat, T. Uehlinger, D. Greif, T. Esslinger, Nature **515**, 237 (2014).
[2] M. Aidelsburger, M. Lohse, C. Schweizer, M. Atala, J. T. Barreiro, S. Nascimbène, N. R. Cooper, I. Bloch, N. Goldman, Nat. Phys. **11**, 162 (2015).

For each poster contribution there will be one poster wall (width: 97 cm, height: 250 cm) available. Please do not feel obliged to fill the whole space. Posters can be put up for the full duration of the event.

POSTER

Sonic Landau levels and synthetic gauge fields in mechanical metamaterials
Abbaszadeh, Hamed

Probing quantum turbulence in He II by quantum evaporation measurements
Amelio, Ivan

Topology and dynamics in driven hexagonal lattices
Asteria, Luca

Spin-orbitcoupling in a Bose-Einstein condensate: Triple-well in momentum space
Cabedo Bru, Josep

Tailoring the Fermi velocity in 2D Dirac Materials
Díaz Fernández, Álvaro

A new machine for dysrosium experiment
Du, Li

Exact Edge and Bulk States of Topological Models and their Robustness Against an Impurity
Duncan, Callum

Enhanced chiral anomaly in Floquet Schwinger model
Ebihara, Shu

Synthetic dimensions and chiral currents with spin-orbit-coupled two-electron ultracold fermions
Franchi, Lorenzo

Realizing and detecting a topological insulator in the AIII symmetry class
García Velasco, Carlos

Real and imaginary part of conductivity of strongly interacting bosons in optical lattices
Grygiel, Barbara

Topological Phases in Ultracold Fermionic Ladders
Haller, Andreas

Characterzing interacting topological states of matter via charge pumps and single-particle topological invariants
Hayward, Andrew

Transport in optical lattices with flux
Hudomal, Ana

Recent cold atom experiments have realized artificial gauge fields in periodically modulated optical lattices [1,2]. We study the dynamics of atomic clouds in these systems by performing numerical simulations using the full time-dependent Hamiltonian and comparing these results to the semiclassical approximation. Under constant external force, atoms in optical lattices with flux exhibit an anomalous velocity in the transverse direction. We investigate in detail how this transverse drift is related to the Berry curvature and Chern number, taking into account realistic experimental conditions. [1] G. Jotzu et al., Nature 515, 237 (2014). [2] M. Aidelsburger et al., Nature Phys. 11, 162 (2015).

Time-periodic driving of spinor condensates in a hexagonal optical lattice
Ilin, Alexander

Local topological invariant of the Interacting Hofstadter Interface

Q 35: Quantum Gases (Fermions) II

Tuesday 14:00–16:15 K 1.022

Talk Q 35.1 Tue 14:00 K 1.022

Artificial gauge potentials in periodically driven optical lattices: numerical simulations of atomic transport — •Ana Hudomal[1], Ivana Vasić[1], Hrvoje Buljan[2], Walter Hofstetter[3], and Antun Balaž[1] — [1]Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia — [2]Department of Physics, University of Zagreb, Croatia — [3]Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität, Frankfurt am Main, Germany

Artificial gauge potentials have been recently realized in cold-atom experiments with periodically driven optical lattices [1,2]. In such systems, atoms subjected to a constant external force gain an anomalous velocity in the direction transverse to the direction of the applied force. Taking into consideration realistic experimental conditions, we perform numerical simulations in order to investigate the dynamics of atomic clouds and relate it to the Chern number of the effective model. We use the full time-dependent Hamiltonian and take into account the effects of weak repulsive interactions between atoms. The results are compared to the semiclassical approximation.

[1] G. Jotzu et al., Nature **515**, 237 (2014).
[2] M. Aidelsburger et al., Nature Phys. **11**, 162 (2015).

Talk Q 35.2 Tue 14:15 K 1.022

Experimental characterization and control of Floquet states in a periodically driven two-body quantum system — •Kilian Sandholzer, Rémi Desbuquois, Michael Messer, Frederik Görg, Joaquín Minguzzi, Gregor Jotzu, and Tilman Esslinger — Institute for Quantum Electronics, ETH Zürich, Zürich, Switzerland

Floquet engineering is a powerful tool to modify properties of a static system such as opening topological gaps or controlling magnetic order. The versatility of cold atom experiments offers the possibility to implement many of these schemes. Nonetheless, preparing a certain Floquet state that has the desired properties in this out-of-equilibrium situation is a more difficult task, especially when the driving frequency is close to a characteristic energy scale of the system. In this work, we prepare fermionic atoms in a driven optical lattice such that the system can be described by two interacting particles on a double well potential with a periodically modulated tilt. In the case of near-resonant driving we achieve to enter adiabatically individual Floquet states by using a two-step ramping protocol. In addition, the fast coherent dynamics inherently connected to the drive are studied in detail. Finally, an analytical derivation of the effective time-independent Hamiltonian of the realized system is presented and then compared to numerical studies and experimental data.

Talk Q 35.3 Tue 14:30 K 1.022

Dynamics of driven interacting many-body systems — •Michael Messer, Frederik Görg, Kilian Sandholzer, Joaquín Minguzzi, Rémi Desbuquois, and Tilman Esslinger — Institute for Quantum Electronics, ETH Zurich, 8093 Zurich, Switzerland

Periodic driving can be used to coherently control the properties of a many-body state and to engineer new phases which are not accessible in static systems. The successful implementation of a periodically driven Fermi-Hubbard model on a 3D hexagonal lattice offers the possibility to explore the intriguing dynamics of Floquet many-body systems. A theoretical analysis of driven many-body Hamiltonians is inherently challenging, however, in combination with our experiments a deeper understanding seems feasible.

By controlling the detuning between shaking frequency and interactions, and setting a variable strength of the periodic drive, we achieve independent control over the single particle tunneling and the magnetic exchange energy. This control allows us to investigate the dynamics and build-up of nearest-neighbor spin-spin correlations. Furthermore, we explore possible mechanisms behind the formation of correlations in interacting Floquet systems. In addition, we can analyze the creation of double occupancies, as one mechanism to form excitations.

Talk Q 35.4 Tue 14:45 K 1.022

Enhancement and sign change of magnetic correlations in a driven quantum many-body system — •Frederik Görg[1], Michael Messer[1], Kilian Sandholzer[1], Joaquín Minguzzi[1], Gregor Jotzu[1,2], Rémi Desbuquois[1], and Tilman Esslinger[1] — [1]Institute for Quantum Electronics, ETH Zurich, 8093 Zurich, Switzerland — [2]Max Planck Institute for the Structure and Dynamics of Matter, 22761 Hamburg, Germany

Strong periodic driving can be used to control the properties of interacting quantum systems. In solid state experiments, ultrashort laser pulses are employed to tune the charge order as well as magnetic and superconducting properties of materials. At the same time, continuous driving has been used in cold atom experiments to engineer novel effective Floquet-Hamiltonians which feature for example a topological bandstructure. We realize a strongly interacting Fermi gas in a periodically driven hexagonal optical lattice and investigate its charge and magnetic properties. We first demonstrate that in the high-frequency regime, the effective description of the many-body system by a renormalized tunnelling amplitude remains valid by comparing our results to an equivalent static system. When driving at a frequency close to the interaction energy, we show that anti-ferromagnetic correlations can be enhanced or even switched to ferromagnetic ordering. Our observations can be explained by a microscopic model, in which the particle tunnelling and magnetic exchange energies can be controlled independently. Therefore, Floquet engineering constitutes an alternative route to experimentally investigate unconventional pairing.

Talk Q 35.5 Tue 15:00 K 1.022

Manipulating and probing excitations of a Chern insulator by Floquet engineering an optical solenoid — •Botao Wang, Nur Ünal, and André Eckardt — Max Planck Institute for the Physics of Complex Systems, Dresden, Germany

The realization of artificial gauge fields in optical lattice systems has paved a way to the experimental investigation of various topological quantum effects. Here we propose a realistic scheme for realizing tunable local (solenoid type) artificial magnetic fields by means of Floquet engineering. We show that such an optical solenoid field can be used to coherently manipulate and probe Chern insulator states of the Hofstadter Hamiltonian. In particular, we investigate the possibility to create local quasiparticle and quasihole excitations, to coherently populated edge modes, and to achieve quantized charge pumping. All these effects are manifested on the spatial density distributions, which can be measured directly in quantum-gas microscopes.

Talk Q 35.6 Tue 15:15 K 1.022

Characterizing topology by dynamics: Chern number from linking number — •Matthias Tarnowski[1,2], Nur Ünal[3], Nick Fläschner[1,2], Benno Rem[1,2], André Eckardt[3], Klaus Sengstock[1,2,4], and Christof Weitenberg[1,2] — [1]Institut für Laserphysik, Universität Hamburg, 22761 Hamburg, Germany — [2]The Hamburg Centre for Ultrafast Imaging, 22761 Hamburg, Germany — [3]Max-Planck-Institut für Physik komplexer Systeme, Nöthnitzer Straße 38, 01187 Dresden, Germany — [4]Zentrum für Optische Quantentechnologien, Universität Hamburg, 22761 Hamburg, Germany

Topology plays an important role in modern solid state physics describing intriguing quantum states such as topological insulators. It is an intrinsically non-local property and therefore challenging to access, often studied only via the resulting edge states. Here, we report on a new approach by connecting the Chern number with the dynamical evolution of highly excited states of the system and demonstrate it experimentally with cold atoms in hexagonal optical lattices. We study the contour of dynamically created vortex pairs in momentum space following a sudden quench into the system of interest and infer the Chern number of the post-quench Hamiltonian from the topology of the contour, quantified by the linking number with the static vortices. Our work exploits a direct mapping between two topological indices and allows detecting topology by the naked eye.

Talk Q 35.7 Tue 15:30 K 1.022

1D fermionic Floquet topological insulators with Hubbard interaction — •Haixin Qiu[1] and Johann Kroha[1,2] — [1]Physikalisches Institut and Bethe Center for Theoretical Physics, Universität Bonn, Nussallee 12, 53115 Bonn, Germany — [2]Center for Correlated Matter, Zhejiang University, Hangzhou, Zhejiang 310058, China

The fermionic Rice-Mele model is a standard model for quantum ratchet transport in periodically driven, one-dimensional, bipartite chains. In the adiabatic limit, this model exhibits quantized transport (Thouless pump), while in the limit of fast drive quasistatic approximations with effective hopping parameters are possible. Here we study the Rice-Mele model with periodic drive of both, the hopping amplitudes and the onsite energy modulation,in the intermediate regime

Transport in optical lattices with flux

A. Hudomal[1], **I. Vasić**[1], **H. Buljan**[2], **W. Hofstetter**[3], **and A. Balaž**[1]
[1] *Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia*
[2] *Department of Physics, University of Zagreb, Croatia*
[3] *Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität, Frankfurt am Main, Germany*
e-mail: *hudomal@ipb.ac.rs*

Different condensed matter systems, such as electrons in a crystal lattice, can be simulated using ultracold atoms in optical lattices. Unlike electrons, atoms are electrically neutral and therefore do not feel the effects of magnetic field. Artificial gauge potentials have been recently realized in cold-atom experiments with periodically driven optical lattices [1, 2]. In such systems, atoms subjected to a constant external force gain an anomalous velocity in the direction transverse to the direction of the applied force.

Taking into consideration realistic experimental conditions, we perform numerical simulations in order to investigate the dynamics of atomic clouds and relate it to the Chern number of the effective model. We consider incoherent bosons and the full time-dependent Hamiltonian. The effects of weak repulsive interactions between atoms are taken into account using the mean-field approximation.

Our results show that driving, external force and interactions all cause heating and transitions to higher bands, which have significant effects on the dynamics. It turns out that weak interactions can be beneficial, because they make the momentum-space probability density more homogeneous. In the future, we also plan to study the details of the atomic-cloud expansion dynamics, and to simulate the full loading sequence of an initial Bose-Einstein condensate, as it was done in the experiment [2].

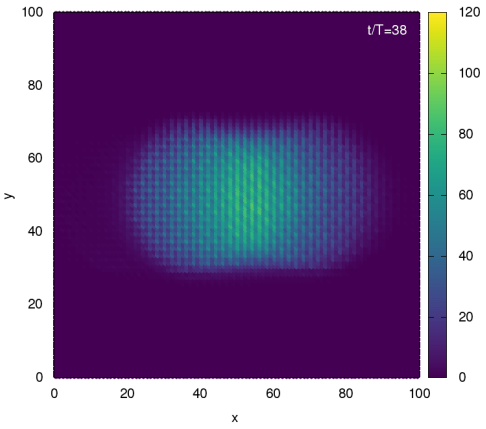

Figure 1: Density profile of an atomic cloud during expansion dynamics after release from a trap in the presence of an artificial gauge field and external force.

[1] G. Jotzu, M. Messer, R. Desbuquois, M. Lebrat, T. Uehlinger, D. Greif, and T. Esslinger, Nature **515**, 237 (2014).

[2] M. Aidelsburger, M. Lohse, C. Schweizer, M. Atala, J. T. Barreiro, S. Nascimbène, N. R. Cooper, I. Bloch, and N. Goldman, Nat. Phys. **11**, 162 (2015).

Transport in Optical Lattices with Flux

A. Hudomal[1], I. Vasić[1], H. Buljan[2], W. Hofstetter[3], and A. Balaž[1]

[1]Scientific Computing Laboratory, Center for the Study of Complex Systems,

Institute of Physics Belgrade, University of Belgrade, Serbia
[2] Department of Physics, Faculty of Science, University of Zagreb, Croatia
[3]Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität,
Frankfurt am Main, Germany

Different condensed matter systems, such as electrons in a crystal lattice, can be simulated using ultracold atoms in optical lattices. Unlike electrons, atoms are electrically neutral and therefore do not feel the effects of magnetic field. Artificial gauge potentials have been recently realized in cold-atom experiments with periodically driven optical lattices [1,2]. In such systems, atoms subjected to a constant external force gain an anomalous velocity in the direction transverse to the direction of the applied force. Taking into consideration realistic experimental conditions, we perform numerical simulations in order to investigate the dynamics of atomic clouds and relate it to the Chern number of the effective model [3]. We consider incoherent bosons and the full time-dependent Hamiltonian. The effects of weak repulsive interactions between atoms are taken into account using the mean-field approximation. Our results show that driving, external force and interactions all cause heating and transitions to higher bands, which have significant effects on the dynamics. It turns out that weak interactions can be beneficial, because they make the momentum-space probability density more homogeneous.

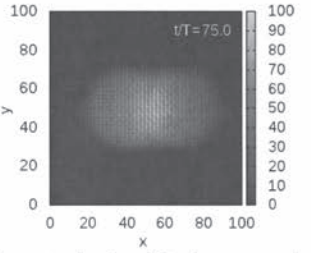

Fig. 1: Density profile of an atomic cloud during expansion dynamics after release from a trap in the presence of an artificial gauge field and external force.

References

[1] G. Jotzu et al., Nature **515**, 237 (2014).
[2] M. Aidelsburger et al., Nat. Phys. **11**, 162 (2015).
[3] A. Hudomal, I. Vasić, H. Buljan, W. Hofstetter, and A. Balaž, arXiv:1809.05125

Q 25: Poster: Quantum Optics and Photonics I

Time: Tuesday 16:30–18:30 Location: S Atrium Informatik

Q 25.1 Tue 16:30 S Atrium Informatik

Unequal-time correlations in Bose-Einstein condensates — •Linda Shen[1,2] and Martin Gärttner[2] — [1]Institut für Theoretische Physik, Philosophenweg 16, 69120 Heidelberg, Germany — [2]Kirchhoff-Institut für Physik, Im Neuenheimer Feld 227, 69120 Heidelberg, Germany

We develop measurement schemes for unequal-time correlation functions in a Bose-Einstein condensate (BEC). Both the spectral and statistical components of the two-point correlation function are investigated out of equilibrium. Thereby, the time-evolution of a BEC is computed numerically using classical-statistical simulation methods based on the Gross-Pitaevskii equation.

The spectral correlation function is approached by linear response methods, which are in principle applicable to both numerical computations as well as experimental measurements. The statistical correlation function can be computed directly in the classical-statistical approximation. Extracting the unequal-time statistical function experimentally, however, requires involved techniques in order to avoid quantum back action effects. We propose to use a non-invasive measurement protocol where the system is weakly coupled to an ancillary system.

In thermal equilibrium, the spectral and statistical components are related by the fluctuation-dissipation theorem. Measuring both will allow a better understanding of how the fluctuation-dissipation theorem builds up as the system approaches equilibrium.

Q 25.2 Tue 16:30 S Atrium Informatik

Quantum Droplets with Tilted Dipoles — •Manuel Schmitt[1], Vladimir Veljić[2], Antun Balaž[2], and Axel Pelster[1] — [1]Research Center OPTIMAS and Department of Physics, Technische Universität Kaiserslautern, Germany — [2]Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia

Since 2005 there have been many striking advancements in Bose-Einstein condensates (BECs) with dipolar interactions, the most recent one being the discovery of quantum droplets, which are stabilized due to quantum fluctuations [1, 2]. With a variational approach we investigate the influence of a tilted dipole axis on quantum droplets in a wave guide-like setup [3]. At first we generalize for one quantum droplet the energy functional for the extended Gross-Pitaevskii theory to tilted dipoles and determine the resulting deformation of the cloud as well as its stability as a function of the tilting angle. Furthermore, we consider two quantum droplets in a trap and calculate how their equilibrium distance depends on the tilting of the dipole axis. With this we gain new insight into the emergence of filaments of dipolar BECs.

[1] M. Schmitt et al., Nature **539**, 259 (2016)
[2] L. Chomaz et al., Phys. Rev. X **6**, 041039 (2016)
[3] I. Ferrier-Barbut et al., Phys. Rev. Lett. **116**, 215301 (2016)

Q 25.3 Tue 16:30 S Atrium Informatik

Many-body Multifractality in Fock space for Interacting Bosons — Jakob Lindinger, Andreas Buchleitner, and •Alberto Rodríguez — Physikalisches Institut, Albert-Ludwigs-Universität Freiburg, Hermann-Herder-Straße 3, D-79104 Freiburg, Germany

We analyse the many-body multifractality of the Bose-Hubbard Hamiltonian's eigenstates in Fock space, for arbitrary values of the interparticle interaction. For the ground state, generalized fractal dimensions unambiguously signal, even for small system sizes, the emergence of a Mott insulator. We show that the scaling of the derivative of any generalised fractal dimension with respect to the interaction strength encodes the critical point of the superfluid to Mott insulator transition, and we establish that the transition can be quantitatively characterized by one single wavefunction amplitude from the exponentially large Fock space [1]. Furthermore, multifractality of the excited eigenstates is investigated and the possible existence of localization in Fock space is thoroughly studied.

[1] J. Lindinger, A. Buchleitner, A. Rodríguez, arXiv:1810.06369

Q 25.4 Tue 16:30 S Atrium Informatik

Dynamics in multi-species bosonic systems — Tobias Brünner, •Gabriel Dufour, Alberto Rodríguez, and Andreas Buchleitner — Physikalisches Institut, Albert-Ludwigs-Universität Freiburg, Hermann-Herder-Straße 3, D-79104 Freiburg, Germany

The dynamics of bosons in multimode systems is determined by an involved interplay between interactions and indistinguishability-induced many-particle interference. We construct a formalism to investigate systematically the dynamics of multiple bosonic species, distinguishable by an internal degree of freedom which is insensitive to the time evolution. We unveil how interparticle interactions lead to a hierarchy of interaction-induced interference processes, such that even the dynamics of single-particle observables is influenced by the degree of indistinguishability (DOI). Time-averaged expectation values of observables dominated by two-particle interference are shown to correlate with a measure of the DOI for initial Fock states [1]. Time-resolved features of the dynamics, such as the frequency content of the signals, are also influenced by the DOI and reveal the interacting or non-interacting nature of the system. We show that this can be understood from the symmetry properties of the Hamiltonian based on group-theoretical arguments [2].

[1] T. Brünner, G. Dufour, A. Rodríguez, A. Buchleitner, Phys. Rev. Lett. 120, 210401 (2018)
[2] T. Brünner, PhD Thesis, Albert-Ludwigs-Universität Freiburg (2018). https://doi.org/10.6094/UNIFR/16683

Q 25.5 Tue 16:30 S Atrium Informatik

Rotational cooling of molecules in a BEC — •Martin Will, Tobias Lausch, and Michael Fleischhauer — University of Kaiserslautern, 67663 Kaiserslautern, Germany

We discuss the rotational cooling of homonuclear diatomic molecules in a Bose-Einstein-condensate (BEC) . For typical molecules there is no frictionless rotation since the dominant cooling occurs via emission of particle-like phonons. Only for macro-dimers, whose size becomes larger than the condensate healing length, a Landau-like, critical angular momentum exists below which phonon emission is suppressed. We find that the phonon-induced angular momentum relaxation is much faster than the cooling of linear motion of impurities in a BEC. This also leads to a finite lifetime of angulons, quasi-particles of rotating molecules coupled to orbital angular-momentum phonons. The lifetimes are however still smaller than typical angulon binding energies. We analyze the dynamics of rotational cooling for homo-nuclear diatomic molecules based on a quantum Boltzmann equation including single- and two-phonon scattering and discuss the effect of thermal phonons. For typical molecules two-phonon scattering becomes relevant at finite temperature.

Q 25.6 Tue 16:30 S Atrium Informatik

Coexistence of phase transitions and hysteresis near the onset of Bose-Einstein condensation — Michael Maennel[3] and •Klaus Morawetz[1,2] — [1]Münster University of Applied Sciences,Stegerwaldstrasse 39, 48565 Steinfurt, Germany — [2]International Institute of Physics- UFRN,Campus Universitário Lagoa nova,59078-970 Natal, Brazil — [3]Informatik DV,Petersstr. 14,04109 Leipzig, Germany

Multiple phases occurring in a Bose gas with finite-range interaction are investigated [2]. In the vicinity of the onset of Bose-Einstein condensation (BEC), the chemical potential and the pressure show a van der Waals-like behavior indicating a first-order phase transition for weak interactions like Hartree-Fock or Popov approximation. However, for strong interactions there remains a multivalued region for the T-matrix approximation even after the Maxwell construction, which is interpreted as a density hysteresis [1]. This unified treatment of normal and condensed phases becomes possible due to the recently found scheme to eliminate self-interactions in the T-matrix approximation, which allows one to calculate properties below and above the critical temperature [3,4]. [1] Phys. Rev. A 87 (2013) 053617, [2] New J. Phys. 12 (2010) 033013, [3] J. Stat. Phys. 143 (2011) 482, [4] Phys. Rev. B 84 (2011) 094529

Q 25.7 Tue 16:30 S Atrium Informatik

Dynamics of weakly interacting bosons in optical lattices with flux — •Ana Hudomal[1], Ivana Vasić[1], Hrvoje Buljan[2], Walter Hofstetter[3], and Antun Balaž[1] — [1]Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics

Belgrade, University of Belgrade, Serbia — [2]Department of Physics, Faculty of Science, University of Zagreb, Croatia — [3]Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität, Frankfurt am Main, Germany

Realization of strong synthetic magnetic fields in driven optical lattices has enabled implementation of topological bands in cold-atom setups [1,2]. A milestone has been reached by a recent measurement of a finite Chern number based on the dynamics of incoherent bosonic atoms [2]. Motivated by these recent developments, we investigate the dynamics of weakly interacting incoherent bosons in a two-dimensional driven optical lattice exposed to an external force, which provides a direct probe of the Chern number [3]. We find that interactions lead to the redistribution of atoms over topological bands both through the conversion of interaction energy into kinetic energy during the expansion of the atomic cloud and due to an additional heating. Remarkably, we observe that the moderate atomic repulsion facilitates the measurement by flattening the distribution of atoms in the quasimomentum space.

[1] G. Jotzu et al., Nature **515**, 237 (2014).
[2] M. Aidelsburger et al., Nature Phys. **11**, 162 (2015).
[3] A. Hudomal et al., Phys. Rev. A **98**, 053625 (2018).

Q 25.8 Tue 16:30 S Atrium Informatik

Quench dynamics and boundary condition dependence of the one-dimensional extended Bose Hubbard model — •Sebastian Stumper, Junichi Okamoto, and Michael Thoss — Insitute of Physics, University of Freiburg, Freiburg, Germany

The one-dimensional extended Bose Hubbard model exhibits a variety of quantum phases due to its competing interactions. For large on-site interactions, a Mott insulating (MI) phase exists, while a charge density wave (CDW) phase becomes dominant for large nearest-neighbour interactions. In between these phases, there exists a topologically nontrivial phase of a Haldane insulator (HI), which is characterized by a non-local string order (Phys. Rev. Lett. 97, 260401 (2006)). Ground state properties and low energy spectra are, however, very sensitive to the treatment of boundary conditions (arXiv:1403.2315 (2014)). We study an open chain of the extended Bose Hubbard model for various configurations of chemical potentials applied at the edges using the density matrix renormalization group method (Comput. Phys. Commun. 225, 59 (2018)). Without edge potentials, the CDW and HI phases show a non-degenerate ground state, and the order parameters change signs in the middle of the chain. This feature is robust against finite size scaling and is explained by a simple effective picture for the low energy states. On the other hand, with large edge potentials, the sign change of the order parameters disappears, and we recover uniform bulk ground states. Furthermore, we simulate quenched dynamics with initial states from MI, HI and CDW phases and discuss the results in terms of our findings on the equilibrium cases.

Q 25.9 Tue 16:30 S Atrium Informatik

Staggered-immersion cooling of a quantum gas in optical lattices — •Bing Yang[1,2,3], Hui Sun[1,2,3], Chun-jiong Huang[2,3], Han-yi Wang[1,2,3], You-jin Deng[2,3], Han-ning Dai[1,2,3], Zhen-sheng Yuan[1,2,3], and Jian-wei Pan[1,2,3] — [1]Physikalisches Institut, Ruprecht-Karls-Universität Heidelberg, Im Neuenheimer Feld 226, 69120 Heidelberg, Germany — [2]Hefei National Laboratory for Physical Sciences at Microscale and Department of Modern Physics, University of Science and Technology of China, Hefei, Anhui 230026, China — [3]CAS Centre for Excellence and Synergetic Innovation Centre in Quantum Information and Quantum Physics, University of Science and Technology of China, Hefei, Anhui 230026, China

Here we realize efficient cooling of ten thousand ultracold bosons in staggered optical lattices. By immersing Mott-insulator samples into removable superfluid reservoirs, thermal entropy is extracted from the system. Losing less than half of the atoms, we lower the entropy of a Mott insulator by 65-fold, achieving a record-low entropy per particle of 0.0019 k_B (k_B is the Boltzmann constant). We further engineer the sample to a defect-free array of isolated single atoms and successfully transfer it into a coherent many-body state. The present staggered-immersion cooling opens up an avenue for exploring novel quantum matters and promises practical applications in quantum information science.

Q 25.10 Tue 16:30 S Atrium Informatik

Simulation of the Quantum Rabi Model with Ultracold Rubidium Atoms in the Deep Strong Coupling Regime — •Geram Hunanyan[1], Johannes Koch[1], Martin Leder[1], Enrique Rico[2,3], Carlos Sabin[4], Enrique Solano[2,3], and Martin Weitz[1] — [1]Institut für Angewandte Physik Bonn, Wegelerstr. 8, D-53115 Bonn, Germany — [2]Department of Physical Chemistry, University of the Basque Country UPV/EHU, Apartado 644, E-48080 Bilbao, Spain — [3]IKERBASQUE, Basque Foundation for Science, Maria Diaz de Haro 3, E-48013 Bilbao, Spain — [4]Instituto de Fisica Fundamental, CSIC, Serrano 113-bis, E-28006 Madrid, Spain

The Quantum Rabi Model (QRM) has been applied to describe the dynamics of a two-level quantum system interacting with a single bosonic mode. Although a fair quantity of experiments explore the strong coupling regime of the QRM, where due to the still limited coupling strength the system can be transformed to the widely known Jaynes-Cummings Model, researchers are just beginning to exploit the regime where the full QRM must be considered. Our experimental implementation to simulate the QRM uses ultracold rubidium atoms in an optical lattice potential, with the effective two-level quantum system being simulated by different Bloch bands in the first Brillouin zone. The bosonic mode is represented by the oscillations of the atoms in an optical dipole trapping potential. We experimentally observe the atomic dynamics in the deep strong coupling regime. The present status of results will be presented.

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Probing the mott-insulator state in optical lattices with photoassociation collisions — •Hui Sun, Bing Yang, Zhen-sheng Yuan, and Jian-wei Pan — Physikalisches Institut, Ruprecht-Karls-Universität Heidelberg, Im Neuenheimer Feld 226, 69120 Heidelberg, Germany

The photoassociation collision is a process two colliding atoms form an excited molecular state after absorbing a photon, which can be used to remove doublons in optical lattices. In this work, we present the detection of a bosonic Mott-insulator state in optical lattices via photoassociation collisions. The photoassociation frequency and collision strength in the 0_g^- molecular channel are calibrated in ultracold quantum gases of Rb^{87}. Then we measure the density distributions of two-dimensional Mott-insulator states in optical lattices after illuminated by a photoassociation light, which is 13.6 cm^{-1} red detuned to the D2 line. From the density profiles, we extract the temperatures of the Mott-insulators and demonstrate an improvement of the measurement precision. This new method extends our ability to probe this ultracold strongly correlated systems.

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Probing Equilibration of Isolated Quantum Systems in a Spinor Bose-Einstein Condensate — •Stefan Lannig, Rodrigo Rosa-Medina Pimentel, Maximilian Prüfer, Philipp Kunkel, Alexis Bonnin, Helmut Strobel, and Markus K. Oberthaler — Kirchhoff-Institut für Physik, Im Neuenheimer Feld 227, 69120 Heidelberg

If and how isolated quantum systems eventually reach thermal equilibrium is still an open question. To address this we experimentally investigate the spin dynamics of a Bose-Einstein condensate of ^{87}Rb. In particular, we focus on the long-time dynamics in the $F = 1$ hyperfine manifold, which realises a spin-1 system. We prepare the system in different out-of-equilibrium states and probe its subsequent evolution by applying a new readout technique which allows to simultaneously extract multiple spin projections. We observe that the kinetic temperature, leading to a finite non-condensed fraction, impacts the coherent evolution and relaxation of the spin observables.

Using local control of the spin orientation and atomic density we aim at further exploring and understanding the relaxation processes involved in the temporal evolution of a 1-d spinor system. We investigate the response of the system to controlled local perturbations which can be connected to spatial and temporal correlations offering new observables for characterisation of general many-particle quantum dynamics.

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Non-equilibrium dynamics of interacting Bosons in an optical lattice — •Jens Benary[1], Christian Baals[1,2], Jian Jiang[1], and Herwig Ott[1] — [1]Department of Physics and OPTIMAS research center, Technische Universität Kaiserslautern, 67663 Kaiserslautern, Germany — [2]Graduate School Materials Science in Mainz, 55128 Mainz, Germany

We study the non-equilibrium dynamics of ultracold Bose gases using a scanning electron microscope. In our latest setup an optical system

Searching for quantum scars in constrained bosonic models

A. Hudomal[1], I. Vasić[1], N. Regnault[2] and Z. Papić[3]
[1]*Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia*
[2]*Laboratoire de Physique de l'Ecole Normale Supérieure, ENS, Université PSL, CNRS, Sorbonne Université, Université Paris-Diderot, Sorbonne Paris Cité, Paris, France*
[3]*School of Physics and Astronomy, University of Leeds, United Kingdom*
e-mail: hudomal@ipb.ac.rs

Recent experiments on arrays of Rydberg atoms have shown that preparing a system in a certain initial state can lead to unusually slow thermalization and persistent density oscillations [1]. This type of non-ergodic behavior has been attributed to the existence of "quantum many-body scars", i.e., atypical, weakly-entangled eigenstates of the system that have high overlaps with a small subset of vectors in the Hilbert space. Periodic dynamics and many-body scars are believed to originate from a "hard" kinetic constraint: due to strong interactions, no two neighbouring atoms are both allowed to be in an excited Rydberg state. Here we investigate quantum many-body scars in a 1D bosonic lattice model with a "soft" constraint: there are no restrictions on the allowed boson states and the particles can hop freely, but the amplitude of a hop depends on the occupancy of the hopping site. We find that this model exhibits similar phenomenology to the Rydberg atom chain, including weakly entangled eigenstates at high energy densities and the presence of a large number of exact zero energy states, with distinct algebraic structure. We discuss the relation of this model to the standard Bose-Hubbard model and possible experimental realizations using ultracold atoms.

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Probing fractional Hall states in driven optical lattices

I. Vasić[1], A. Hudomal[1] and N. Regnault[2]
[1]*Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia*
[2]*Laboratoire de Physique de l'Ecole Normale Supérieure, ENS, Université PSL, CNRS, Sorbonne Université, Université Paris-Diderot, Sorbonne Paris Cité, Paris, France*
e-mail: ivana.vasic@ipb.ac.rs

Driven optical lattices enrich the set of quantum models that can be simulated in cold-atom experiments [1, 2]. General arguments suggest that the interplay of strong interactions and driving in a thermodynamically large system introduces heating, leading to a featureless infinite-temperature state in the long-time limit [3, 4]. Recently, several papers have focused on a possibility of prethermalization, arguing that some strongly correlated states can be probed on experimentally relevant timescales, before reaching the infinite-temperature limit [5, 6]. We investigate ways to prepare and probe fractional Hall states in a few-particle bosonic sample in a driven optical lattice.

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Searching For Quantum Scars In Constrained Bosonic Models

Ana Hudomal[a], Ivana Vasić[a], Nicolas Regnault[b] and Zlatko Papić[c]

[a]*Scientific Computing Laboratory, Center for the Study of Complex Systems, Institute of Physics Belgrade, University of Belgrade, Serbia*
[b]*Laboratoire de Physique de l'Ecole Normale Supérieure, ENS, Université PSL, CNRS, Sorbonne Université, Université Paris-Diderot, Sorbonne Paris Cité, Paris, France*
[c]*School of Physics and Astronomy, University of Leeds, United Kingdom*

Abstract. Recent experiments on arrays of Rydberg atoms have shown that preparing a system in a certain initial state can lead to unusually slow thermalization and persistent density oscillations [1]. This type of non-ergodic behavior has been attributed to the existence of "quantum many-body scars", i.e., atypical, weakly-entangled eigenstates of the system that have high overlaps with a small subset of vectors in the Hilbert space. Periodic dynamics and many-body scars are believed to originate from a "hard" kinetic constraint: due to strong interactions, no two neighbouring atoms are both allowed to be in an excited Rydberg state. Here we investigate quantum many-body scars in a 1D bosonic lattice model with a "soft" constraint: there are no restrictions on the allowed boson states and the particles can hop freely, but the amplitude of a hop depends on the occupancy of the hopping site. We find that this model exhibits similar phenomenology to the Rydberg atom chain, including weakly entangled eigenstates at high energy densities and the presence of a large number of exact zero energy states, with distinct algebraic structure. We discuss the relation of this model to the standard Bose-Hubbard model and possible experimental realizations using ultracold atoms.

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Recent experiments on Rydberg atom arrays have found evidence of anomalously slow thermalization and persistent density oscillations, which have been interpreted as a many-body analog of the phenomenon of quantum scars. Periodic dynamics and atypical scarred eigenstates originate from a "hard" kinetic constraint: the neighboring Rydberg atoms cannot be simultaneously excited. Here we propose a realization of quantum many-body scars in a 1D bosonic lattice model with a "soft" constraint in the form of density-assisted hopping. We discuss the relation of this model to the standard Bose-Hubbard model and possible experimental realizations using ultracold atoms. We find that this model exhibits similar phenomenology to the Rydberg atom chain, including weakly entangled eigenstates at high energy densities and the presence of a large number of exact zero energy states, with distinct algebraic structure.

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Ana Hudomal; Ana Hudomal; Ivana Vasic; Hrvoje Buljan; Walter Hofstetter; Antun Balaz

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Barcelona, July 27th, 2018

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University of Belgrade
Faculty of Physics

Ana Hudomal

NUMERICAL STUDY OF QUANTUM GASES IN OPTICAL LATTICES AND IN SYNTHETIC MAGNETIC FIELDS

Doctoral Dissertation

Belgrade, 2020

Univerzitet u Beogradu
Fizički fakultet

Ana Hudomal

NUMERIČKO PROUČAVANJE KVANTNIH GASOVA U OPTIČKIM REŠETKAMA I U SINTETIČKIM MAGNETNIM POLJIMA

Doktorska disertacija

Beograd, 2020.

Thesis defense committee

Thesis advisor:

Dr. Ivana Vasić
Associate Research Professor
Institute of Physics Belgrade
University of Belgrade

Committee members:

Dr. Antun Balaž
Research Professor
Institute of Physics Belgrade
University of Belgrade

Dr. Milan Knežević
Full Professor
Faculty of Physics
University of Belgrade

Dr. Božidar Nikolić
Associate Professor
Faculty of Physics
University of Belgrade

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During the work on this thesis I was employed at the Scientific Computing Laboratory (SCL) of the Institute of Physics Belgrade. I had many opportunities to participate in international collaboration and to attend schools and conferences. As many of my colleagues from SCL work in different fields, I could hear about a wide variety of research topics. I would like to thank Dr. Antun Balaž, the head of SCL, for giving me the opportunity to work in this stimulating environment, as well as for all his help and useful advice.

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Numerical study of quantum gases in optical lattices and in synthetic magnetic fields

Abstract

Theoretical and experimental advances in the past few decades have resulted in the development of a new research field – quantum simulations with ultracold atoms. The main idea is to create a clean and precisely controllable quantum system which can then be used to simulate another system of interest – one that is not as easy to study experimentally. Although there are several different experimental platforms for quantum simulations, ultracold atoms are often used because they are especially well suited for this role. Progress in cooling and trapping techniques has enabled experimentalists to cool down atomic gases to quantum degeneracy and to easily manipulate them. These systems are well isolated from their environment and do not contain any impurities or defects. In particular, cold atoms in optical lattices – periodic potentials made by interference of laser beams – have shown to be an excellent platform for the study of various condensed matter systems. It is possible to realize different lattice geometries in the desired number of dimensions. The lattice size, number of atoms, temperature of the system and even the strength of interactions between atoms can be precisely tuned. Some of these parameters would be impossible to change in a real condensed matter system. The set of models which can be realized with ultracold atoms can be further extended by the use of Floquet engineering. In this approach, the system is exposed to a suitable time-periodic modulation. The resulting stroboscopic dynamics of this driven system can be related to a corresponding static model through Floquet theory. In particular, Floquet engineering was used to realize synthetic magnetic fields in systems of neutral cold atoms.

In this thesis we use numerical simulations and analytical insights to study dynamics of several relevant systems which have been the focus of recent experiments with ultracold atoms in optical lattices. One of the fundamental open questions that has gained lot of attention recently is related to the thermalization of a general isolated quantum system. Such systems are typically shown to thermalize in experiments, meaning that they lose all memory of their initial state. However, there are several notable counterexamples. In particular, a new class of systems which exhibit unusual thermaization has been recently discovered – the systems with special eigenstates called quantum many-body scars. Another long-standing problem is

realization of topologically nontrivial models with ultracold atoms. These systems usually require magnetic fields, whose effects can be mimicked by driving in a cold atom system. However, interactions between atoms are always present in a realistic system. The interplay of driving and interactions typically leads to the thermalization and additional considerations are necessary in order to identify regimes where this process is slow and allows for the preparation and measurement of an interesting topological state.

Recent experiments on Rydberg atom arrays have found evidence of anomalously slow thermalization and persistent density oscillations, which have been interpreted as a many-body analog of the phenomenon of quantum scars. Periodic dynamics and atypical scarred eigenstates have been obtained in a model with a "hard" kinetic constraint: the neighboring Rydberg atoms cannot be simultaneously excited. In the first part of this thesis we propose a realization of quantum many-body scars in a one-dimensional bosonic lattice model with a "soft" constraint in the form of density-assisted hopping. We discuss the relation of this model to the standard Bose-Hubbard model and possible experimental realizations using ultracold atoms. We find that this model exhibits similar phenomenology to the Rydberg atom chain, including weakly entangled eigenstates at high energy densities and the presence of a large number of exact zero energy states, with distinct algebraic structure.

Realization of strong synthetic magnetic fields in driven optical lattices has enabled implementation of topological bands in cold-atom setups. A milestone has been reached by a recent measurement of a finite Chern number based on the dynamics of incoherent bosonic atoms. The measurements of the quantum Hall effect in semiconductors are related to the Chern-number measurement in a cold-atom setup; however, the design and complexity of the two types of measurements are quite different. Motivated by these recent developments, in the second part of this thesis we investigate the dynamics of weakly interacting incoherent bosons in a two-dimensional driven optical lattice exposed to an external force, which provides a direct probe of the Chern number. We consider a realistic driving protocol in the regime of high driving frequency and focus on the role of weak repulsive interactions. We find that interactions lead to the redistribution of atoms over topological bands both through the conversion of interaction energy into kinetic energy during the expansion of the atomic cloud and due to an additional heating. Remarkably, we observe that the moderate atomic repulsion facilitates the measurement by flattening the distribution of atoms in the quasimomentum space. Our results also show that weak interactions can suppress the contribution of some higher-order nontopological terms in favor of the topological part of the effective model.

Strong interactions and strong synthetic magnetic fields, the main ingredients for the realization of fractional quantum Hall states, are already available in experiments on cold atom gases in periodically driven optical lattices. However, the interplay of the driving and interactions introduces detrimental heating, and for this reason it is still challenging to reach a fractional quantum Hall state in cold-atom setup. By performing a numerical study, in the third part of

this thesis we investigate stability of a bosonic Laughlin state in a small atomic sample exposed to driving. We identify an optimal regime of microscopic parameters, in particular interaction strength U and the driving frequency ω, such that the stroboscopic dynamics supports the basic $\nu = 1/2$ Laughlin state. Moreover, we explore slow ramping of a driving term and show that the considered protocol allows for the preparation of the Laughlin state on experimentally realistic time-scales.

Keywords: ultracold gases, quantum simulations, nonequilibrium dynamics, quantum scars, topological phases of matter, synthetic gauge fields, Floquet systems, exact diagonalization
Research field: Physics
Research subfield: Condensed matter physics
UDC number: 538.9

Numeričko proučavanje kvantnih gasova u optičkim rešetkama i u sintetičkim magnetnim poljima

Sažetak

Teorijski i eksperimentalni napredak u poslednjih nekoliko decenija doveo je do razvoja nove oblasti istraživanja – kvantnih simulacija sa ultrahladnim atomima. Osnovna ideja je da se napravi čist i precizno podesiv kvantni sistem koji se zatim može koristiti za simuliranje nekog drugog sistema od intresa koji nije tako jednostavan za eksperimentalno ispitivanje. Iako postoji nekoliko različitih platformi za kvantne simulacije, ultrahladni atomi se često koriste jer su izuzetno pogodni za ovu ulogu. Napredak u tehnikama za hlađenje i zarobljavanje atoma omogućio je eksperimentalnim fizičarima da ohlade atomske gasove do kvantne degeneracije i lako manipulišu njima. Ovi sistemi su dobro izolovani od svog okruženja i ne sadrže nikakve nečistoće ili defekte. Hladni atomi u optičkim rešetkama, periodičnim potencijalima nastalim interferencijom laserskih zraka, pokazali su se kao odlična platforma za proučavanje raznovrsnih sistema kondenzovane materije. Moguće je realizovati različite geometrije rešetke u željenom broju dimenzija. Veličina rešetke, broj atoma, temperatura sistema, čak i jačina interakcija između atoma mogu se precizno podešavati. Neke od ovih parametara ne bi bilo moguće promeniti u sistemu kondenzovane materije. Skup modela koji se mogu realizovati pomoću ultrahladnih atoma može se dalje proširiti korišćenjem Floke inženjeringa. U ovom pristupu, sistem je izložen pogodnoj vremenski periodičnoj modulaciji. Rezultirajuća stroboskopska dinamika ovog vođenog sistema može se povezati sa odgovarajućim statičkim modelom putem Floke teorije. Floke inženjering je korišćen i za realizaciju sintetičkih magnetnih polja u sistemima neutralnih hladnih atoma.

U ovoj tezi koristićemo numeričke simulacije i analitičke metode u cilju proučavanja dinamike nekoliko relevantnih sistema koji su bili u fokusu skorašnjih eksperimenata sa ultrahladnim atomima u optičkim rešetkama. Jedno od osnovnih otvorenih pitanja koje je nedavno privuklo dosta pažnje povezano je sa termalizacijom izolovanog kvantnog sistema u opštem slučaju. Takvi sistemi se u eksperimentima tipično termalizuju, što znači da gube svu memoriju o svom početnom stanju. Ipak, postoji nekoliko značajnih kontraprimera. Jedna nova klasa sistema sa atipičnom termalizacijom je nedavno otkrivena – to su sistemi sa posebnim svojstvenim stanjima nazvanim kvantni višečestični ožiljci. Još jedan dugogodišnji problem

predstavlja realizacija topološki netrivijalnih modela sa ultrahladnim atomima. Takvi sistemi obično zahtevaju magnetna polja, čiji efekti mogu da se oponašaju vođenjem u sistemu hladnih atoma. Međutim, interakcije između atoma su uvek prisutne u realističnim sistemima. Uzajamno dejstvo dejstvo vođenja i interakcija dovodi do termalizacije, pa su dodatna razmatranja neophodna da bi se identifikovali režimi u kojima je ovaj proces spor i dozvoljava pripremu i merenje interesantnih topoloških stanja.

Nedavni eksperimenti na nizovima Ridbergovih atoma pronašli su dokaze o anomalno sporoj termalizaciji i dugotrajnim oscilacijama gustine, što je bilo interpretirano kao višečestični analogon fenomena kvantnih ožiljaka. Periodična dinamika i atipična svojstvena stanja sa ožiljcima dobijeni su u modelu sa "jakim" kinetičkim ograničenjem: susedni Ridbergovi atomi ne mogu da budu istovremeno pobuđeni. U prvom delu ove teze predstavljena je realizacija kvantnih višečestičnih ožiljaka u jednodimenzionalnom bozonskom modelu na rešetki sa "slabim" ograničenjem u formi tunelovanja potpomognutim gustinom. Diskutovana je veza ovog modela sa standardnim Boze-Habard modelom i mogućnost njegove eksperimentalne realizacije sa ultrahladnim atomima. Ovaj model ispoljava sličnu fenomenologiju kao i lanac Ridbergovih atoma, uključujući slabo kvantno uvezana svojstvena stanja na velikim gustinama energije, kao i prisustvo velikog broja egzaktnih nultih energetskih stanja sa posebnom algebarskom strukturom.

Realizacija jakih sintetičkih magnetnih polja u vođenim optičkim rešetkama je omogućila implementaciju topoloških energetskih zona u sistemima hladnih atoma. Značajno postignuće predstavlja skorašnje merenje nenultog Černovog broja bazirano na dinamici nekoherentnih bezonskih atoma. Merenja kvantnog Holovog efekta u poluprovodnicima povezana su sa merenjem Černovog broja u eksperimentima sa hladnim atomima, ali se dizajn i kompleksnost ove dve vrste eksperimenata dosta razlikuju. Motivisani skorašnjim eksperimentima, u drugom delu ove teze ispitujemo dinamiku slabo interagujućih nekoherentnih bozona u dvodimenzionalnoj vođenoj optičkoj rešetki pod dejstvom spoljašnje sile, što omogućava direktno merenje Černovog broja. Razmaramo realističan protokol vođenja u režimu visoke frekvencije i fokusiramo se na ulogu slabih odbojnih interakcija. Pokazujemo da interakcije dovode do redistribucije atoma na topološke energetske zone kroz konverziju interakcione energije u kineticku energiju u toku širenja atomskog oblaka, kao i zbog dodatnog zagrevanja. Primećujemo da umereno odbijanje između atoma olakšava merenje putem poravnanja distribucije atoma u kvazi-impulsnom prostoru. Naši rezultati takođe pokazuju da slabe interakcije mogu da ponište doprinos nekih netopoloških članova višeg reda u korist topološkog dela efektivnog modela.

Jake interakcije i jaka sintetička magnetna polja, glavni sastojci za realizaciju frakcionog kvantnog Holovog efekta, već su dostupni u eksperimentima sa hladnim atomskim gasovima u periodično vođenim optičkim rešetkama. Međutim, uzajamno dejstvo vođenja i interakcija izaziva neželjeno zagrevanje, pa je iz tog razloga ostvarivanje frakcionog kvantnog Holovog stanja u sistemu hladnih atoma i dalje veliki izazov. Uz pomoć numeričkih simulacija, u trećem

delu ove teze istražujemo stabilnost bozonskih Laflinovih stanja u malom atomskom uzorku pod dejstvom periodičnog vođenja. Pronalazimo optimalan režim mikroskopskih parametara, jačinu interakcija U i frekvenciju vođenja ω, takvih da stroboskopska dinamika podržava osnovno $\nu = 1/2$ Laflinovo stanje. Pored toga, istražujemo postepeno uključivanje člana koji opisuje vođenje i pokazujemo da razmatrani protokol dozvoljava pripremu Lafinovog stanja na eksperimentalno relevantnim vremenskim skalama.

Ključne reči: ultrahladni gasovi, kvantne simulacije, neravnotežna dinamika, kvantni ožiljci, topološke faze materije, sintetička gejdž polja, Floke sistemi, egzaktna dijagonalizacija
Naučna oblast: Fizika
Uža naučna oblast: Fizika kondenzovanog stanja
UDK broj: 538.9

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Chapter 1

Introduction

When atoms are cooled down to sufficiently low temperatures, their quantum statistics starts to play an important role and the differences between bosons and fermions become apparent. In the case of bosons, cooling leads to macroscopic occupation of a single quantum state – the Bose-Einstein condensate (BEC). More than 70 years after its first theoretical prediction in 1924, BEC was finally experimentally achieved in 1995, when the atoms were cooled down to 170 nK [1]. This was made possible by advances in laser and evaporative cooling techniques during the previous decades, as well as magnetic and optical trapping. These techniques were later also applied to fermionic atoms, molecules, and mixtures of different atomic species, successfully cooling down these systems to quantum degeneracy. During the last two decades since BEC was first attained, significant progress has been made in both theoretical and experimental research on quantum gases. There are many possible applications of ultracold quantum gases. They can be used for precision measurements, and can also serve as a platform for quantum computing and quantum simulations.

1.1 Quantum simulations

Quantum simulators were first envisioned by Feynman in 1982 [2, 3]. The idea was to create a clean and highly controllable quantum system that can be used to simulate another complex quantum system described by an equivalent Hamiltonian. The need for quantum computing and quantum simulations has arisen due to computational limits of classical computers. While it is theoretically possible to find the exact eigenvalues and eigenstates for any finite-dimensional interacting quantum many-body system, the Hilbert space dimensions of such systems grow exponentially with the system size, which quickly makes it impossible to store the necessary data in the memory of even the most advanced modern supercomputers. Exact numerical calculations are therefore limited to the smallest systems consisting only of tens of particles and

lattice sites, which sometimes makes it very hard to extrapolate the results to the thermodynamic limit. Feynman's proposal was to use a computer which operates on quantum-mechanical laws to solve quantum-mechanical problems. Unlike a classical computer, in this case the number of necessary computer elements – quantum bits – scales linearly with the size of the system of interest. Additionally, a quantum computer is probabilistic – there is no unique output for each input. Instead, such a computer returns several outputs with different probabilities, which is expected as it is actually performing a quantum measurement.

Quantum simulators are a similar yet distinct concept. Unlike universal quantum computers, which could in theory be programmed to execute any possible algorithm, but are still years or even decades away from practical applications [4], quantum simulators are designed for a specific task and are therefore easier to build. There is already a variety of their experimental realizations on different platforms, for example using neutral cold atoms [5, 6], trapped ions [7], superconducting circuits [8], photonic systems [9] or nitrogen-vacancy centers in diamond [10, 11]. Some of the phenomena simulated in this way are quantum magnetism [12], strange metal phase of high-temperature superconductors [13, 14], decay of a Higgs particle [15], black-hole radiation [16] and photosynthesis [17]. Quantum simulations have shown to be useful in a wide range of different research fields, including condensed-matter physics, high-energy physics, cosmology and quantum chemistry. They could also be applied to classical problems which require large computing power. While universal quantum computers are very prone to computation errors, as each particle needs to be set to a precisely defined quantum state, this is not the case for quantum simulators where the resulting quantity is typically an average over the whole system and thus less sensitive to the exact state of an individual particle.

In particular, ultracold atoms in optical lattices provide a perfect platform for quantum simulations of various condensed-matter phenomena [5, 18]. Here, the optical lattice plays the role of the crystal lattice in solid-state systems, while the atoms play the role of the electrons. Optical lattice is a spatially periodic potential created by interference of two counter-propagating laser beams. Schematic representations of two-dimensional and three-dimensional optical lattices are shown in Fig. 1.1. Cold atoms can be trapped inside the lattice by the optical dipole force. Various optical-lattice geometries can be realized depending on the intensity and relative angles between interfering laser beams. The strength of interactions between the atoms can be tuned using Feshbach resonances. The possibility to precisely adjust all microscopic parameters of the system can be used to create exotic phases of matter that are not observed in nature, or to experimentally probe quantities that are otherwise not accessible in condensed-matter systems.

Two seminal condensed-matter models which were realized with ultracold atoms are the Bose-Hubbard and the Fermi-Hubbard model. The experimental setup is similar in both cases; a dilute atomic gas, typically consisting of alkali atoms (Li, Na, K, Rb, Cs), is placed inside a magneto-optical trap, cooled down to low temperatures and exposed to an optical lattice po-

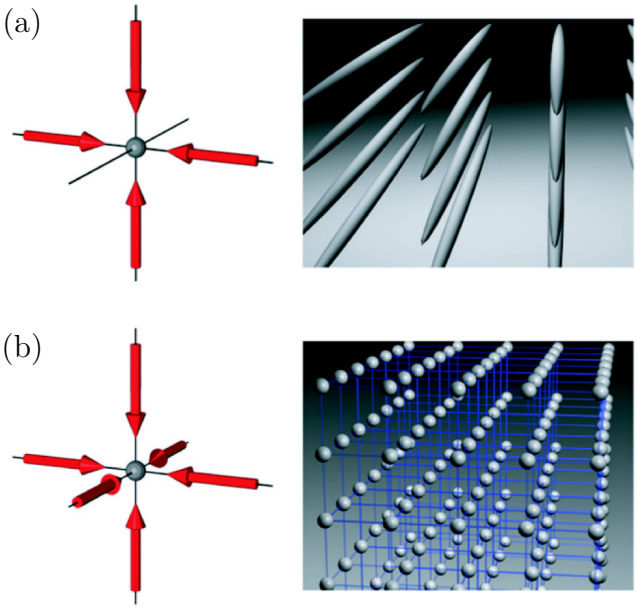

Figure 1.1: Schematic of a two-dimensional (a) and three-dimensional (b) square optical lattice. Adapted from Ref. [5].

tential. The exact model which is realized, Bose- or Fermi-Hubbard, depends on the bosonic or fermionic nature of the atomic species which is used. Bosonic atoms are used more often as they are easier to cool down to quantum degeneracy, given that their ground state is a Bose-Einstein condensate. Early experiments have studied quantum phase transitions and characterized different phases of matter in these two models. Another topic that was experimentally explored with these models is their nonequilibrium dynamics. In recent experiments, Bose-Hubbard and Fermi-Hubbard models serve as a foundation for the realization of other generalized and more complex models.

1.2 Bose-Hubbard model

Interacting spinless bosons in a periodical potential are described by the Bose-Hubbard Hamiltonian [19]. This model has only two parameters; hopping amplitude J which determines the probability for a particle to tunnel to a neighbouring site, and the strength of interactions U between two particles on the same site which can be either repulsive ($U > 0$) or attractive ($U < 0$). The Bose-Hubbard Hamiltonian can be written as

$$\hat{H}_{\mathrm{BH}} = -J \sum_{\langle l,m \rangle} \left(\hat{a}_l^\dagger \hat{a}_m + \mathrm{H.\ c.} \right) + \frac{U}{2} \sum_l \hat{n}_l \left(\hat{n}_l - 1 \right), \tag{1.1}$$

where $\hat{a}_l^\dagger$ and $\hat{a}_l$ are creation and annihilation operators that create and annihilate a particle at the lattice site with index l, $\hat{n}_l = \hat{a}_l^\dagger \hat{a}_l$ is the particle number operator, and the label $\langle l, m \rangle$ stands for nearest neighbors. This model was derived using the single-band tight-binding ap-

proximation [5]. The Hamilatonian shown here is given by a general expression for an arbitrary number of dimensions. The equation for a specific number of dimensions can be written in a similar manner, with potentially different hopping amplitudes in different directions. The Bose-Hubbard Hamiltonian has a global $U(1)$ gauge symmetry and all the symmetries of the underlying lattice. As the occupancy of a single lattice site is not limited for bosonic particles, the Hilbert space dimension grows with the system size as

$$\dim\mathcal{H} = \binom{N+L-1}{N}, \tag{1.2}$$

where N is the number of particles and L is the number of lattice sites.

Although the Bose-Hubbard Hamiltonian [19] was first formulated for solid state systems as generalization of the (Fermi-)Hubbard model [20], the closest experimental realization of this model is provided by ultracold bosonic atoms in an optical lattice. The necessary conditions for the approximation to be valid are that the atoms are cooled down to a low enough temperature, so that only the lowest energy band is significantly occupied, and that the lattice is deep enough, thus making the Wannier functions sufficiently localized at each lattice site. As the atoms interact by short-range Van der Waals forces and there are no long-range Coulomb interactions between neutral atoms, the onsite interaction term is sufficient to describe the interactions in this system. Long-range dipolar interactions can be avoided by choosing an atomic species without a dipolar moment.

The phase diagram of the Bose-Hubbard model consists of the superfluid phase and the Mott insulator phase, depending on the ratio J/U and the filling factor $\nu = N/L$ [19]. In the superfliuid phase where the hopping term is dominant over the interaction term, the particles are completely delocalized across all lattice sites. The ground state in this phase is a BEC and can be described by a single Bloch wavefunction. It is characterized by long-range phase coherence. This state can be described by the Gross-Pitaevskii equation – a nonlinear Schrödinger-like equation. In contrast, the interaction term is dominant in the Mott insulator phase and in the ground state an integer number of particles is localized at each lattice site. This is a strongly-correlated quantum many-body state. In the limit $U \to 0$, it is a product of local Fock states at each site. Unlike the superfluid state, the Mott insulator state is incompressible, as particle number fluctuations are energetically very costly. There is no more phase coherence in the Mott insulating phase. Instead, there are long-range particle number correlations.

The phase transition between these two phases has been experimentally observed in an ultracold quantum gas [21]. In this experiment, the ratio J/U was varied by increasing the lattice potential depth, which resulted in decreased hopping amplitude J and increased onsite interaction strength U. After setting the desired ratio, the confining potential was turned off and the atomic cloud was left to expand freely in the optical lattice potential. The phase coherence between different lattice sites was visible in the resulting interference pattern. The

absence of interference maxima marked the onset of the Mott insulating phase. A schematic of the two states and the corresponding interference patterns are given in Fig. 1.2. Experiments like this one and their excellent agreement with theoretical predictions have confirmed that the Bose-Hubbard model is indeed realized with cold atoms in optical lattices.

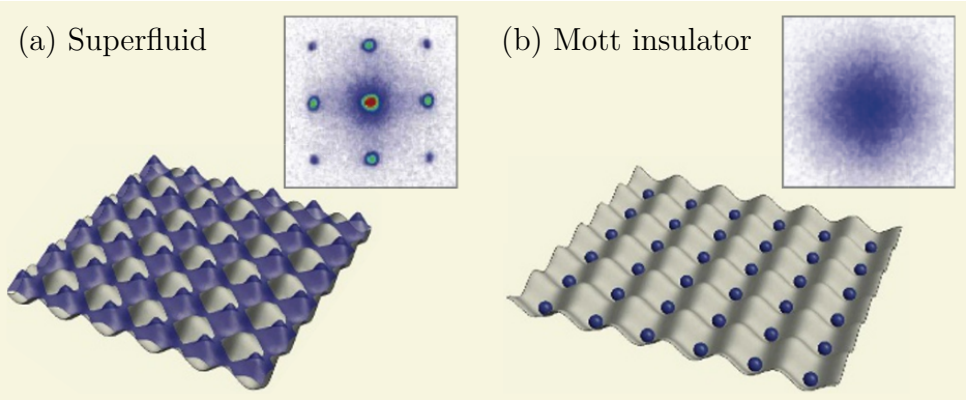

Figure 1.2: Schematic of a superfluid state (a) and a Mott insulator state (b) in an optical lattice and corresponding interference patterns after expansion. Adapted from Ref. [21].

There are several generalizations of this model which include additional potentials, long-range density-density or dipolar interactions, next-nearest-neighbour and other hopping terms, interaction- or density-induced hopping terms, particles with spin, as well as the Bose-Fermi-Hubbard which describes a mixture of bosonic and fermionic particles. The Bose-Hubbard Hamiltonian will be the starting model for all the studies in this thesis. We will be interested in systems with density-dependent hopping terms and systems with additional time-dependent potentials – the so-called driven systems.

1.3 Nonequilibrium dynamics

While the properties of quantum systems in equilibrium are generally well understood, nonequilibrium dynamics of such systems is still an active field of research. A system can be taken out of equilibrium by changing its Hamiltonian, either by a sudden quench or by continuously ramping one or more parameters. Systems in contact with a thermal reservoir are generally expected to thermalize, but it is not clear whether a general isolated quantum system should ever reach thermal equilibrium. Ultracold atoms provide an excellent experimental platform for the study of quantum many-body physics out of equilibrium, as they are both precisely tunable and well isolated from their environment [22]. In such experiments, the system is typically first prepared in the ground state $|\psi_0\rangle$ of some initial Hamiltonian $\hat{H}_0$ and then evolved under a different Hamiltonian $\hat{H}$, which may or may not be time-dependent.

The main question is through which mechanism and under which conditions isolated quantum systems initially out of equilibrium evolve towards a state in thermal equilibrium. This

question has been studied in a variety of experiments on quantum gases [23–29]. Complete thermalization has been observed in some of these experiments [24, 28]. The final thermalized state is independent of the initial state – all memory is lost during the process of thermalization. However, this loss of information is somewhat surprising because of the fact that the time-evolution operator $\hat{U}(t) = \mathrm{e}^{-i\hat{H}t}$ is unitary and therefore preserves the system in a pure quantum state. A potential answer is provided by the eigenstate thermalization hypothesis (ETH) [30–33]. ETH states that thermalization happens already at the eigenstate level and that this becomes visible in the expectation values of certain observables during time evolution, due to dephasing between eigenstates which constitute the initial state. In other words, even though the system stays in a pure quantum state indefinitely, the results of measurements will appear thermalized after long enough time.

Let the initial state be a superposition of eigenstates of the Hamiltonian $\hat{H}$ from an energy range $E_{\mathrm{min}} \leq E \leq E_{\mathrm{max}}$. If ETH is obeyed for a certain observable $\hat{\mathcal{O}}$, its diagonal matrix elements in the basis of $\hat{H}$, $\mathcal{O}_{ii} = \langle i|\hat{\mathcal{O}}|i\rangle$, will be approximately a smooth function of energy $\mathcal{O}_{ii}(E)$ over the energy range of interest. ETH then predicts that the expectation value of the observable $\hat{\mathcal{O}}$ after thermalization will be equal to the value predicted by the microcanonical ensemble over the same energy range. A single eigenstate is sufficient to predict the thermal expectation value, as all the eigenstates from the energy range in question would lead to the same prediction. ETH typically holds for few-body observables, both local and nonlocal, which are the ones usually measured in experiments. Although ETH has not been analytically proven for a general interacting quantum system, it was shown to be valid in a number of numerical studies [33]. Analytical proofs exist for several specific cases.

However, there are several types of systems which do not obey ETH. Some of them completely avoid thermalization. Two well known examples are integrable systems [34] and systems which exhibit many-body localization [35]. Integrable systems typically have a large number of conserved quantities which constrain their evolution and prevent them from eventually thermalizing. Many-body localization (MBL) requires random disorder in the system and is characterized by the emergence of local integrals of motion. These local integrals of motion preserve the information about the initial state. MBL was shown to be robust with respect to perturbations. Signatures of MBL have been observed in multiple experiments [36–39].

An useful diagnostic of thermalization is the bipartite entanglement entropy [40, 41], a quantity which measures the degree of entanglement between two complementary subsystems labeled A and B:

$$S_{\mathrm{A}} = -\mathrm{Tr}\left(\rho_{\mathrm{A}} \ln \rho_{\mathrm{A}}\right). \tag{1.3}$$

The entanglement entropy is defined for a certain pure state $|\psi\rangle$. In the preceding equation, $\rho_{\mathrm{A}} = \mathrm{Tr}_{\mathrm{B}}|\psi\rangle\langle\psi|$ is the reduced density matrix of the subsystem A. In contrast to thermalizing systems where the entanglement entropy of highly excited eigenstates typically scales as the

volume of the subsystem A, in the MBL phase it obeys the so-called "area law" – it is proportional to the area of the boundary between the subsystems [35, 41, 42]. Additionally, during the nonequilibrium dynamics the entanglement entropy grows logarithmically with time, instead of the usual linear or power-law growth [43–45].

Another type of systems with unusually slow thermalization has emerged in recent years – quantum many-body scarred systems. It was shown that these systems are neither integrable nor many-body localized. This phenomenon was first observed experimentally in a quantum simulator with an array of Rydberg atoms [46]. While the system exhibited the expected thermalizing behavior for most initial states, preparing it in certain initial states surprisingly resulted in persistent density oscillations – the system kept returning to its initial state. Theoretical explanation for such behavior has followed soon after the experimental results [47, 48]. It was shown that the states with unusual dynamics are superpositions of atypical eigenstates dubbed "qunatum many-body scars". These eigenstates have significantly smaller entanglement entropies than the other eigenstates at the same energy scale. Their name was inspired by the phenomenon of quantum scars in noninteracting systems, where some eigenstates of a quantum system have enhanced probability density in vicinity of unstable periodic orbits of the corresponding classical system [49]. In the case of quantum many-body scars, the eigenstates are concentrated in certain regions of the Hilbert space instead of real space. Such atypical eigenstates have been subsequently found in a variety of different systems [50–67]. It is however still not completely clear what are the necessary conditions for their existence. A particular realization of quantum many-body scars in a one-dimensional model of bosons with correlated hopping will be the topic of the first part of this thesis.

Systems which slowly thermalize or do not thermalize at all could have several practical applications. Most importantly, they could be useful for quantum computing, where it is of crucial importance to preserve the system in a particular quantum state. Another application would be to extend the available timescales in experiments on exotic quantum phases which are currently hindered by rapid thermalization.

1.4 Synthetic magnetic fields

Strong magnetic fields are required for many phenomena in condensed matter physics, such as integer and fractional quantum Hall effects. The Harper-Hofstadter model [68] is a two-dimensional model which describes a charged particle in a periodic potential and perpendicular magnetic field. Its Hamiltonian can be written as

$$\hat{H}_{\mathrm{HH}} = -J \sum_{l,m} \left(\mathrm{e}^{-2i\pi m\alpha} \hat{a}^{\dagger}_{l+1,m} \hat{a}_{l,m} + \hat{a}^{\dagger}_{l,m+1} \hat{a}_{l,m} + \mathrm{H.\ c.} \right), \tag{1.4}$$

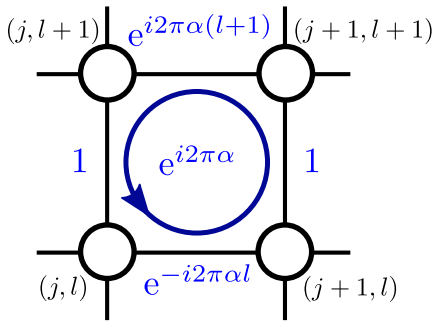

Figure 1.3: Schematic representation of the Harper-Hofstadter model.

where the parameter α is the magnetic flux through one plaquette in the units of flux quanta h/q. The effects of magnetic field are encoded in the position-dependent complex hopping coefficients $e^{-2i\pi m\alpha}$. A particle hopping in counterclockwise direction along the smallest closed loop acquires a phase $2\pi\alpha$, as shown in Fig. 1.3, in analogy to a charged particle in magnetic field which would acquire the Aharonov-Bohm phase. The energy spectrum of the Harper-Hofstadter model has an interesting fractal structure, as can be seen in Fig. 1.4. The number of energy bands depends on α; when α is a rational number the number of bands is equal to its denominator. For example, there are three energy bands for $\alpha = 1/3$, see Fig. 1.4(a), as well as for $\alpha = 2/3$. The plot of energy versus α shown in Fig. 1.4(b) is called the "Hofstadter butterfly".

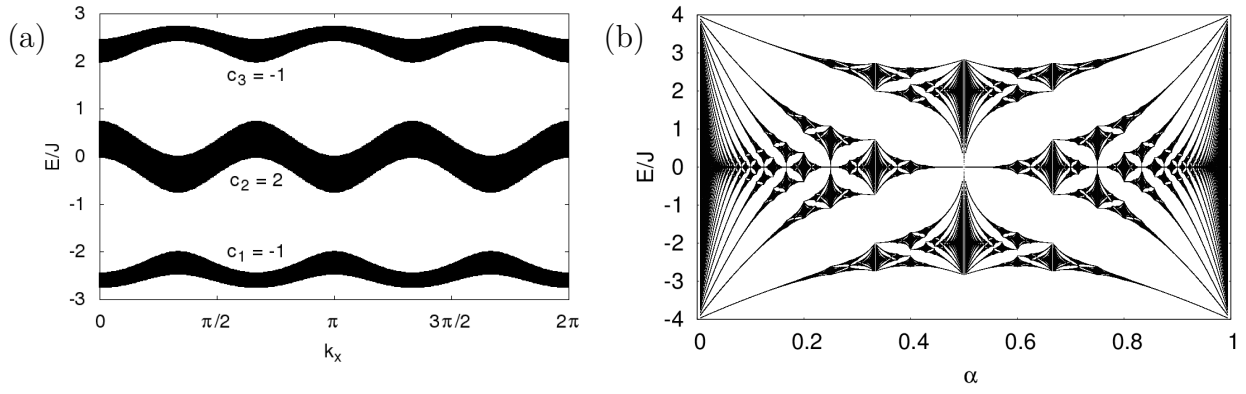

Figure 1.4: (a) Energy spectrum for $\alpha = 1/3$. (b) "Hofstadter butterfly" [68] – energy spectrum of the Harper-Hofstadter model.

In the seminal TKNN paper [69] it was shown that the quantization of the Hall conductivity observed in the integer Hall effects can be directly related to the topological index of the microscopic model (1.4) – the Chern number. The Chern number is defined for a single energy band and is always an integer. It is also a topological invariant, which means that it is insensitive to local deformations and disorder, and changes only when the global topological properties of the system change, i.e. when the band gap closes. The Chern numbers for the three energy bands of the Harper-Hofstadter model with $\alpha = 1/3$ are shown in Fig. 1.4(a). Topological concepts in physics have become an increasingly important field of research in the last few

decades. In 2015, the deflection of an atomic cloud in an optical lattice with synthetic magnetic field as a response to external force was used to experimentally measure the Chern number in a non-electronic system for the first time [70].

The Chern number is closely related to the Berry curvature and Berry phase [71]. It is defined as the integral of the Berry curvature [72] over the first Brillouin zone divided by 2π,

$$c_n = \frac{1}{2\pi} \int_{\mathrm{FBZ}} \mathbf{\Omega}_n(\mathbf{k}) \cdot d\mathbf{S}, \tag{1.5}$$

where n denotes the band number and the Berry curvature is $\mathbf{\Omega}_n(\mathbf{k}) = i\nabla_{\mathbf{k}} \times \langle u_n(\mathbf{k}) | \nabla_{\mathbf{k}} | u_n(\mathbf{k}) \rangle$, expressed in terms of eigenstates $|u_n(\mathbf{k})\rangle$. The Berry phase is a geometric phase that a particle acquires when it makes a closed loop in some parameter space and can be thought of as the analogue of the Aharonov-Bohm phase, while the Berry curvature is then analogous to the magnetic field.

If one wants to simulate systems with magnetic fields using cold atoms, the problem arises from the fact that atoms are charge-neutral and therefore do not feel the Lorentz force. These effects have to be included in some artificial way, by engineering the so-called synthetic magnetic fields. One of the first ideas was to take advantage of the similarity between the Lorentz force and the Coriolis force. This was experimentally realized in 2000 using rotating quantum gases [73], where the appearance of quantized vortices was observed, which is a property of superfluids and superconductors in magnetic field. However, this experimental realization poses several difficulties. Maximal rotation velocity achievable in the experiment is technically limited and it is difficult to realize a stable rotating optical lattice. All of this has made it impossible to obtain strong synthetic magnetic fields required for quantum Hall effects using rotational approach. In 2009, after years of efforts, synthetic magnetic potentials for neutral atoms were implemented by exploiting atomic coupling to a suitable configuration of external lasers [74, 75]. These techniques were further extended to optical lattices, leading to the realization of strong, synthetic, magnetic fields. As a result, important condensed-matter models – the Harper-Hofstadter (1.4) and the Haldane model [76] – are nowadays available in cold-atom setups [77–80]. The key property of these models is their non-trivial topological content.

1.5 Floquet engineering

One way to enrich the set of models which can be realized using cold atoms in optical lattices is by adding time-periodic perturbations to the system. Such systems are called driven systems. This can be done either by lattice shaking – periodic modulation of the lattice position, or through laser-assisted tunneling, which results in a periodic modulation of the lattice potential. Both methods have been used in present-day cold-atom realizations of important topological models [77–79], where they were necessary in order to engineer synthetic magnetic fields. A

schematic representation of the experimental setup [77] which realized the famous Harper-Hofstadter model [68] given by Eq. (1.4) can be seen in Fig. 1.5. In this thesis we will mostly consider laser-assisted tunneling. The resulting time-dependent Hamiltonian can in that case be written as

$$\hat{H}(t) = \hat{H}_0 + \hat{V}(t), \tag{1.6}$$

where $\hat{H}_0$ is typically the underlying Bose-Hubbard Hamiltonian $\hat{H}_{\mathrm{BH}}$ and $\hat{V}(t)$ is a time-periodic modulation of the lattice potential with frequency ω and amplitude κ

$$\hat{V}(t) = \sum_i \hat{n}_i V_i(t). \tag{1.7}$$

Here i is the lattice site index and $\hat{n}_i$ is the particle number operator for that lattice site.

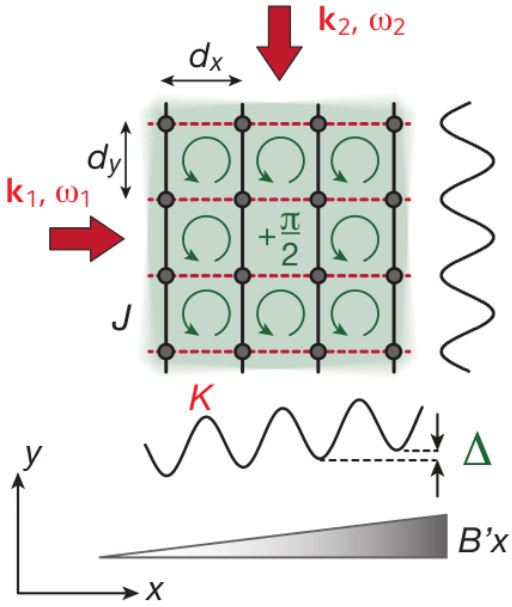

Figure 1.5: Schematic of the experimental setup which implements the Harper-Hofstadter model using ultracold atoms in an optical lattice. Synthetic magnetic field is realized through Floquet engineering. Hopping along the x-axis is first inhibited by a magnetic field gradient B' which generates an offset Δ between neighboring sites. The hopping is then restored using two additional pairs of laser beams with wave vectors $\mathbf{k}_1$, $\mathbf{k}_2$ and frequencies ω_1, ω_2. Laser-assisted tunneling adds a complex phase to the hopping coefficients. The bare hopping amplitude is denoted by J and the complex hopping amplitude by K. This setup mimics the complex phase acquired by a charged particle moving in a magnetic field. In this model, a particle hopping in counterclockwise direction around a single plaquette acquires a phase $\pi/2$, which makes it equivalent to the Harper-Hofstadter model with $\alpha = 1/4$. Adapted from Ref. [77].

Using Floquet theory [81], a periodically driven system can be related to a time-independent effective Hamiltonian. A properly chosen periodic modulation can result in an effective Hamiltonian which corresponds to a relevant condensed-matter system. The mapping is known as Floquet engineering and its important features in the context of optical lattices are discussed in

Refs. [82–88]. The basis of this mapping is Floquet's theorem [89], which is applicable to time-periodic Hamiltonians, analogously to the more famous Bloch's theorem for systems periodic in space. Ultracold quantum gases provide an optimal platform for Floquet engineering, as these systems are usually isolated from their environment, which inhibits dissipative processes, and the highly-controllable parameters of the system can be easily periodically modulated.

Analogously to the Bloch wave states in spatially periodic systems, Floquet states can be written as

$$|\psi_n(t)\rangle = \mathrm{e}^{-i\epsilon_n t}|u_n(t)\rangle, \tag{1.8}$$

where ϵ_n are the quasienergies and $|u_n(t)\rangle = |u_n(t+T)\rangle$ are the time-periodic Floquet modes which have the same period $T = 2\pi/\omega$ as the Hamiltonian $\hat{H}(t)$. As in the rest of this thesis, here we work in the units where $\hbar = 1$. Quasienergies ϵ_n are only defined up to integer multiples of the driving frequency ω, similarly to the quasimomenta which are only defined inside the first Brillouin zone. The Floquet states are the eigenstates of the evolution operator over one driving period $\hat{U}(T, 0)$. When a system governed by a time-periodic Hamiltonian is observed stroboscopically – at times that are integer multiples of the period, it behaves as if it was governed by some time-independent effective Hamiltonian. The dynamics during one period – micromotion – is described by another periodic operator, the so-called kick operator that has the same periodicity as the original time-dependent Hamiltonian. The time-evolution operator corresponding to the Hamiltonian (1.6) can be represented as

$$\hat{U}(t, t_0) = \mathrm{e}^{-i\hat{K}(t)}\mathrm{e}^{-i(t-t_0)\hat{H}_{\mathrm{eff}}}\mathrm{e}^{i\hat{K}(t_0)}, \tag{1.9}$$

where $\hat{H}_{\mathrm{eff}}$ is the full time-independent effective Hamiltonian that describes slow motion and $\hat{K}(t)$ is the time-periodic kick-operator that describes micromotion [82, 83].

The effective Hamiltonian corresponding to a certain time-dependent Hamiltonian always exists, but in the general case it cannot be analytically calculated. However, there are several approximation schemes that allow computation of the leading terms of the effective Hamiltonian and the kick-operator. In cases where the driving frequency ω is large enough compared to the matrix elements of the Hamiltonian, two most often used approximations are the high-frequency expansion [82, 83, 90] and the Magnus expansion [91–93], which is in terms of the driving period $T = 2\pi/\omega$. The advantage of the former method over the latter is that the effective Hamiltonian does not depend on the initial driving phase. For this reason, we will only use the high-frequency expansion in this work. A general time-periodic modulating potential can be written in the form

$$\hat{V}(t) = \sum_{j=1}^{\infty}\left(\hat{V}^{(j)}\mathrm{e}^{ij\omega t} + \hat{V}^{(-j)}\mathrm{e}^{-ij\omega t}\right). \tag{1.10}$$

In the high-frequency expansion, the first few terms of the effective Hamilatonian corresponding

to the time-dependent Hamilatonian (1.6) are then [82]

$$\begin{aligned}\hat{H}_{\text{eff}} &= \hat{H}_0 + \frac{1}{\omega}\sum_{j=1}^{\infty}\frac{1}{j}\left[\hat{V}^{(j)},\hat{V}^{(-j)}\right] + \frac{1}{2\omega^2}\sum_{j=1}^{\infty}\frac{1}{j^2}\left(\left[\left[\hat{V}^{(j)},\hat{H}_0\right],\hat{V}^{(-j)}\right] + \text{H.c.}\right) \\ &+\frac{1}{3\omega^2}\sum_{j,l=1}^{\infty}\frac{1}{jl}\left(\left[\hat{V}^{(j)},\left[\hat{V}^{(l)},\hat{V}^{(-j-l)}\right]\right] - \left[\hat{V}^{(j)},\left[\hat{V}^{(-l)},\hat{V}^{(l-j)}\right]\right] + \text{H.c.}\right) \\ &+\mathcal{O}\left(\frac{1}{\omega^3}\right). \end{aligned} \tag{1.11}$$

This is a general equation that can be applied to a variety of different setups.

Floquet engineering provides a powerful tool for modern cold-atom experiments. However, the interactions between atoms can never be completely avoided and the combination of interactions and driving leads to unwanted heating. Thermalization to infinite temperature in the long-time and thermodynamic limit is practically unavoidable in such systems [94, 95]. Nevertheless, it might still be possible to find some parameter regime where thermalization is slow enough, for example if the system stays in the prethermalized state [96–98] on experimentally relevant timescales. This would allow experimental measurements in driven systems to be performed before full thermalization occurs.

1.6 This thesis

Throughout this thesis, we will use numerical simulations to study several interesting models which can be realized using ultracold bosonic atoms in optical lattices. In particular, we will be interested in systems which exhibit unusual nonergodic dynamics and systems with nontrivial topological properties. Both of these topics represent very active fields of research which have rapidly developed in recent years.

In Chapter 2 we will present a realization of quantum many-body scars in a bosonic model with density-dependent hopping. Since the first experimental observation in a Rydberg atom quantum simulator [46] and the subsequent theoretical explanation using the PXP model [47, 48], quantum many-body scars have been shown to exist in a variety of different systems [50–67]. However, most of the previous realizations have relied on the presence of "hard" kinetic constraints which restrict the available Hilbert space. For example, in the original PXP model this constraint arises from the fact that two neighbouring atoms cannot be simultaneously excited into Rydberg states due to an infinite energy penalty. Another question that was raised in the literature is proximity of quantum-scarred models to integrability, as adding certain perturbations to the PXP model can make it integrable. The main goal of Chapter 2 will be to show that quantum many-body scars can exist in manifestly nonintegrable systems with "soft" constraints only. To this end we will compare several similar models with different types of constraints. We will also formulate an analytically tractable approximation which

can explain and qualitatively predict revivals of quantum fidelity in the model which exhibits quantum many-body scars. An advantage of this model is that it could be easily realized in an optical lattice under a suitable Floquet scheme, thus providing a new experimental platform for quantum many-body scars. This would also allow future probes of this phenomenon to go beyond one-dimensional systems.

The main objective of the second part of this doctoral dissertation will be to study the interplay of topological features and interactions by investigating the dynamics of bosons in driven optical lattices. The approximations necessary for the computation of the effective Hamiltonian in a Floquet engineering scheme are mostly applied to noninteracting systems in the high-frequency limit, even though interactions usually cannot be realistically neglected and the interplay of driving and interactions can heat up the system to a featureless, infinite-temperature regime according to general considerations [94, 95]. One of the main open questions is whether it is possible to find some parameter regimes where the system is in the so-called prethermalized state [96–98] on intermediate timescales and can be described by some model of interest.

In Chapter 3 we will examine the effects of weak on-site interactions in relation to the recent Chern number measurement [70]. We take into account a realistic driving scheme and experimental parameters. Although the interactions are generally thought to complicate the experimental procedure, we will show that weak repulsive interactions can be beneficial for the measurement in at least two ways. Firstly, the interactions make the probability distribution in momentum space more homogeneous, which is important because the Chern number captures the contributions of Berry curvature from the whole Brillouin zone. Secondly and somewhat surprisingly, the interactions can cancel-out some unwanted higher-order therms which are a result of the driving protocol but are not related to the topological model of interest.

In Chapter 4 we will study a similar system as in Chapter 3, but now focusing on strong interactions. We will investigate the possibility of realizing a strongly correlated phase – fractional quantum Hall state – in driven systems on experimentally relevant timescales. Two main ingredients for the realization of fractional quantum Hall states, strong interactions and strong synthetic magnetic fields, are already available in present-day cold-atom experiments. However, even after years of experimental progress and numerous theoretical proposals, these states have still not been achieved, mainly due to problems related to heating caused by driving in the strongly interacting regime. Experimental realization of fractional quantum Hall states in cold-atom setups would be of particular interest due to their anyonic excitations, which could be of use for topological quantum computing [99]. Here we identify an optimal regime of microscopical parameters for the preparation of these highly sought-after states.

Finally, we will summarize all the results from this doctoral dissertation in Chapter 5. Additional derivations and technical details are provided in Appendices A-G.

Chapter 2

Quantum scars of bosons with correlated hopping

Semiclassical studies of chaotic stadium billiards have revealed the existence of remarkable non-chaotic eigenfuctions called "quantum scars" [49]. Scarred eigenfunctions display anomalous enhancement in regions of the billiard that are traversed by one of the periodic orbits in the classical limit when $\hbar \to 0$. It was shown that quantum scars lead to striking experimental signatures in a variety of systems, including microwave cavities [100], quantum dots [101], and semiconductor quantum wells [102].

A recent experiment on a quantum simulator [46], and subsequent theoretical work [47, 103], have shown that quantum many-body scars can occur in strongly interacting quantum systems. The experiment used a one-dimensional Rydberg atom platform in the regime of the Rydberg blockade [46, 104, 105], where nearest-neighbour excitations of the atoms were energetically prohibited. The experiment observed persistent many-body revivals of local observables after a "global quench" [106] from a certain initial state. In contrast, when the experiment was repeated for other initial configurations, drawn from the same type of "infinite" temperature ensemble, the system displayed fast equilibration and no revivals. These observations pointed to a different kind of out-of-equilibrium behavior compared to previous studies of quantum thermalization in various experimental platforms [23, 28, 36, 38, 107].

In both single-particle and many-body quantum scars, the dynamics from certain initial states leads to periodic revivals of the wave function. In the former case, this happens when the particle is prepared in a Gaussian wave packet initialized along a periodic orbit [49], while in the latter case the revivals can be interpreted as a nearly-free precession of a large emergent SU(2) spin degree of freedom [108, 109]. Another similarity between single- and many-body quantum scars is the existence of non-ergodic eigenstates. In the single-particle case, such eigenstates are easily identified by their non-uniform probability density that sharply concentrates along classical periodic orbits. In the many-body case, non-ergodic eigenstates are broadly

defined as those that violate Eigenstate Thermalization Hypothesis (ETH) [30, 31]. Scarred eigenstates violate the ETH in a number of ways: for example, they appear at evenly spaced energies throughout the spectrum [47, 48, 60], they have anomalous expectation values of local observables compared to other eigenstates at the same energy density, and their entanglement entropy obeys a sub-volume law scaling [48].

In recent works, the existence of atypical eigenstates has been taken as a more general definition of quantum many-body scaring. For example, highly-excited eigenstates with low entanglement have previously been analytically constructed in the non-integrable AKLT model [50, 51]. A few of such exact eigenstates are now also available for the Rydberg atom chain model [52]. The collection of models that feature atypical eigenstates is rapidly expanding, including perturbations of the Rydberg atom chain [48, 53, 54], theories with confinement [55, 56, 110], Fermi-Hubbard model beyond one dimension [57, 111], driven systems [58], quantum spin systems [59, 61], fractional quantum Hall effect in a one-dimensional limit [62], and models with fracton-like dynamics [63–65]. In a related development, it was proposed that atypical eigenstates of one Hamiltonian can be "embedded" into the spectrum of another, thermalizing Hamiltonian [112], causing a violation of a "strong" version of the ETH [33, 41]. This approach allows to engineer scarred eigenstates in models of topological phases in arbitrary dimensions [66]. From a dynamical point of view, it has been shown that models with scarred dynamics can be systematically constructed by embedding periodic on-site unitary dynamics into a many-body system [67].

A feature shared by many scarred models is the presence of some form of a kinetic constraint. In the Rydberg atom chain, the constraint results from strong van der Waals forces, which project out the neighboring Rydberg excitations [113]. Such Hilbert spaces occur, for example, in models describing anyon excitations in topological phases of matter [114–118] and in lattice gauge theories [119–121], including the Rydberg atom system [122, 123]. Recent works on periodically driven optical lattices have started to explore such physics [124, 125]. On the other hand, kinetic constraints have been investigated as a possible pathway to many-body localization without disorder [35]. In classical systems, non-thermalizing behavior without disorder is well-known in the context of structural glasses [126–128]. The mechanism of this type of behavior is the excluded volume interactions that impose kinetic constraints on the dynamics [129, 130]. Similar type of physics has recently been explored in quantum systems where a "quasi many-body localized" behavior was proposed to occur in the absence of disorder [131–141].

In this Chapter we investigate the relation between kinetic constraints, slow dynamics and quantum many-body scars. In contrast to previous work, which focused on models of spins and fermions that are closely related in one dimension due to the Jordan-Wigner mapping, here we study one-dimensional models of bosons with density-assisted hoppings, which realize both "hard" and "soft" kinetic constraints, whilst being non-integrable. In Section 2.1 we

introduce the models and discuss properties of their Hamiltonians when viewed as adjacency matrices of graphs in the Fock space. In Section 2.2 we investigate thermalization properties of these models by studying their energy level statistics, entanglement entropy of eigenstates, and dynamics under global quench. Depending on the form of the hopping term, we demonstrate that the models encompass a rich phenomenology, including regimes of fast thermalization, the existence of periodic revivals and many-body scars, as well as the Hilbert space fragmentation that has been found in recent studies of fractonic models [63–65]. Unlike the experimentally realized Rydberg atom system, we find evidence of many-body scars in a bosonic model without a hard kinetic constraint, i.e., with a fully connected Hilbert space. In Section 2.3 we identify initial states that give rise to periodic many-body revivals in the quantum dynamics, and we introduce a "cluster approximation" that captures the scarred eigenstates that are responsible for periodic revivals. In Section 2.4 we discuss zero-energy eigenstates of our models and their algebraic structure. Finally, in Section 2.5 we present our conclusions and discuss possible experimental realizations of these models using ultracold atoms.

2.1 Models and their Hilbert spaces

A fundamental ingredient of kinetically constrained models is "correlated hopping": a particle can hop depending on the state of its neighbors. In this Chapter we consider a system of N_p bosons on a one-dimensional lattice with L sites. We consider models where the total filling factor, $\nu = N_\mathrm{p}/L$, is conserved, and we will mainly present results in the dense regime, $\nu = 1$. We have studied models with $\nu < 1$ and $\nu > 1$, but we found them to be either too constrained or not constrained enough, and therefore less interesting. We emphasize that the bosons in our study are not hard-core, i.e., the occupancy of any lattice site can take any value from 0 to N_p.

2.1.1 Models

We study three different models, defined by the Hamiltonians:

$$\hat{H}_1 = -J\sum_{j=1}^{L}\left(\hat{b}_j^\dagger\hat{b}_{j+1}\hat{n}_j + \hat{n}_{j-1}\hat{b}_j^\dagger\hat{b}_{j-1}\right), \tag{2.1}$$

$$\hat{H}_2 = -J\sum_{j=1}^{L}\left(\hat{n}_j\hat{b}_j^\dagger\hat{b}_{j+1} + \hat{b}_j^\dagger\hat{b}_{j-1}\hat{n}_{j-1}\right), \tag{2.2}$$

$$\hat{H}_3 = -J\sum_{j=1}^{L}\left(\hat{n}_{j+1}\hat{b}_j^\dagger\hat{b}_{j+1}\hat{n}_j + \hat{n}_{j-1}\hat{b}_j^\dagger\hat{b}_{j-1}\hat{n}_j\right). \tag{2.3}$$

All three models contain a free-boson hopping term, $\hat{b}_j^\dagger\hat{b}_{j+1}$, which is dressed in various ways by density operators, $\hat{n}_j = \hat{b}_j^\dagger\hat{b}_j$. We will show that the position of the density operator $\hat{n}_j$ completely changes the behavior of these models, ranging from fast thermalization to the breakup

of the Hamiltonian into disconnected, exactly solvable sectors. For example, note that $\hat{H}_1$ and $\hat{H}_2$ are related to each other via free boson hopping,

$$\hat{H}_2 = \hat{H}_1 - J\sum_j \left(\hat{b}_j^\dagger \hat{b}_{j+1} + \hat{b}_j^\dagger \hat{b}_{j-1}\right), \tag{2.4}$$

which can be easily proven using bosonic commutation relations. We will see below that this innocuous free-boson hopping leads to surprisingly different dynamical properties of the two models.

The motivation behind introducing three different models in Eqs. (2.1)-(2.3) can be summarized as follows. Hamiltonian $\hat{H}_1$ describes a model where a particle cannot hop to the left if that site is not already occupied by at least one particle, and cannot hop to the right if it is the only particle left on its initial site. This introduces constraints to the system. Conversely, there are no such constraints in the case of $\hat{H}_2$. Indeed, the hopping coefficients are only modified in intensity by the particle-number operator. Hamiltonian $\hat{H}_3$ introduces additional constraints compared to $\hat{H}_1$. The number of unoccupied sites and their positions remain constant under the action of this Hamiltonian. This leads to different connectivity of the Hilbert space in each of the models, as we explain in the next Section.

We consider periodic boundary conditions ($L+1 \equiv 1$) and set $\hbar = J = 1$. With periodic boundary conditions, all three Hamiltonians $\hat{H}_1$, $\hat{H}_2$ and $\hat{H}_3$ have translation symmetry, thus their eigenstates can be labelled by momentum quantum number, k, quantized in units of $2\pi/L$. In addition, $\hat{H}_3$ has inversion symmetry. We denote by $I = 0$ and $I = 1$ the sectors that are even and odd under inversion, respectively.

Without restrictions on the boson occupancy, the Hilbert space of $\hat{H}_1$, $\hat{H}_2$ and $\hat{H}_3$ grows very rapidly. For $L = N_\mathrm{p} = 12$, the Hilbert space size of the $k = 0$ sector is 112720 (the largest one we will consider for $\hat{H}_1$ and $\hat{H}_2$). As previously mentioned (see also the next Section), the Hilbert space of $\hat{H}_3$ splits into many disconnected components, thus it is possible to consider only one connected component at a time and disregard the unoccupied sites whose positions do not change. This is more relevant when looking at properties such as thermalization, than fixing the filling factor. However, the boundary conditions are in that case no longer periodic, and the system does not have translation symmetry. Considering only a system with the size $L/2$, filling factor $\nu = 2$, open boundary conditions and minimal number of particles per site equal to 1 is completely equivalent to considering the largest component of the full system which has the size L, filling factor $\nu = 1$, periodic boundary conditions and no restrictions on the occupancies. The Hilbert space size of the symmetric invariant sector of the largest connected component of $L = N_\mathrm{p} = 22$ is 176484 and this is the largest sector that we will consider for $\hat{H}_3$.

2.1.2 Graph structure and bipartite lattice

Since we will be interested in the dynamical properties, it is convenient to first build some intuition about the structure of the Hamiltonians of the three models in Eqs. (2.1)-(2.3). A Hamiltonian can be viewed as the adjancency matrix of a graph whose vertices are Fock states of bosons, $|n_1, n_2, \ldots, n_L\rangle$. If the Hamiltonian induces a transition between two Fock states, the corresponding vertices of the graph are connected by a link. The graphs that show how the configuration space is connected have very different structure for the three Hamiltonians $\hat{H}_1$, $\hat{H}_2$ and $\hat{H}_3$, as can be observed in Fig. 2.1.

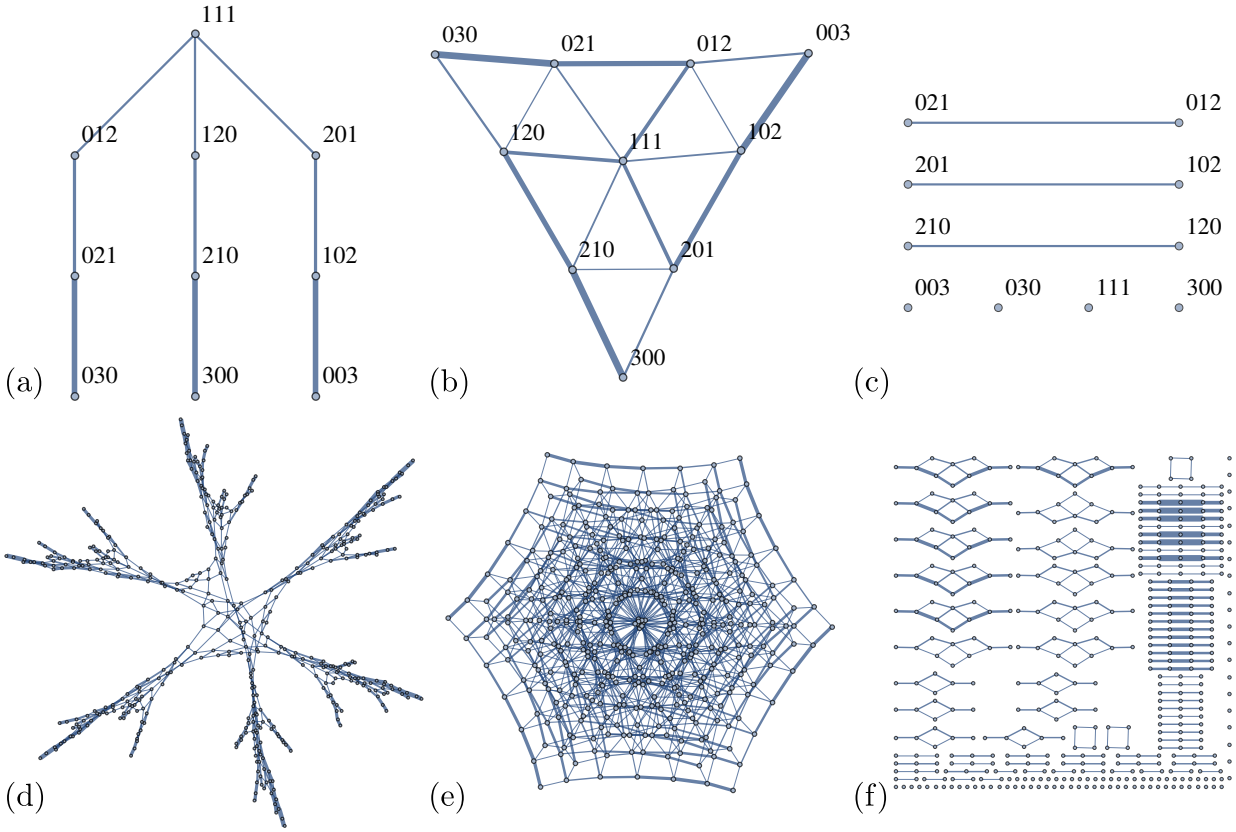

Figure 2.1: Connectivity of the Hilbert space. Adjacency graph for (a) $\hat{H}_1$, (b) $\hat{H}_2$, (c) $\hat{H}_3$, all for $L = N_\mathrm{p} = 3$. (d), (e) and (f): same as (a), (b) and (c) but for $L = N_\mathrm{p} = 6$. To avoid clutter, we do not label the vertices in (d), (e) and (f). All graphs are weighted, i.e., the line thickness is proportional to the magnitude of the corresponding hopping coefficient. Several different clusters of configurations are visible in the case of $\hat{H}_1$. The clusters start to form already for $L = 3$ (for example, the configurations 012–021–003 in (a)) and become more prominent for $L = 6$ (d). In the case of $\hat{H}_2$, almost all configurations are well-connected to the rest of the graph. The graphs for $\hat{H}_3$ show that the Hilbert space is highly reducible: its graph splits into many disconnected components.

The entire graph of $\hat{H}_2$ is well-connected and it has the same structure as the graph of the standard Bose-Hubbard model: the particle-number operators in $\hat{H}_2$ do not introduce any constraints, but only affect the magnitude of the hopping coefficients. In contrast, the $\hat{H}_1$ graph shows several clusters of configurations that are weakly connected to the rest of the

graph. "Weakly connected" means that there is a small number of connections leading outside the cluster and that their respective hopping coefficients are smaller in magnitude than those of the surrounding connections within the cluster. A state that is initially located inside a cluster is therefore more likely to stay inside during an initial stage of the time evolution, which increases the probability of revivals and slows down the growth of entanglement entropy. We will provide a more quantitative description and examples that illustrate this in Section 2.3. Finally, the graph of $\hat{H}_3$, due to even stronger constraints, is actually disconnected, which is an example of Hilbert space fragmentation that was previously shown to cause non-ergodic behavior in fracton-like models [64, 65]. This predicts that thermalization and dynamics in the three models will be very different, which we will confirm in the following Section. However, we note that the number of connections and the topology of the graph is not the only relevant factor for the dynamics. The magnitude of the hopping coefficients between different configurations is also important, see Appendix A.

We note that the relation between $\hat{H}_1$ and $\hat{H}_3$ is reminiscent of the relation between the quantum East model [142] and the "PXP" model describing the atoms in the Rydberg blockade regime [47, 48, 113]. Like $\hat{H}_3$, the PXP model is doubly constrained and inversion symmetric, while $\hat{H}_1$ and the quantum East model are asymmetric versions of those two models with only a single constraint. The graph of the quantum East model is similar to that of $\hat{H}_1$, in that it contains bottlenecks which slow down the growth of entanglement entropy [142].

The graph of $\hat{H}_1$ is bipartite, i.e. all the basis configurations can be divided into two disjoint sets, and the action of the Hamiltonian connects configurations in one set only to the configurations in the other and vice-versa (the Hamiltonian is off-diagonal). One way to sort configurations into these two sets is by parity of the quantity

$$\Delta_a = \frac{|n_{\text{even}} - n_{\text{odd}} + C|}{2}, \tag{2.5}$$

where $C = 0$ if L is even and $C = 1$ if L is odd. We define n_{even} and n_{odd} as the total numbers of particles at even and odd sites, respectively,

$$n_{\text{even}} = \sum_{l=1}^{L_1} n_{2l}, \;\; n_{\text{odd}} = \sum_{l=1}^{L_2} n_{2l-1}, \tag{2.6}$$

where $L_1 = L_2 = L/2$ if L is even, and $L_1 = (L-1)/2$, $L_2 = (L+1)/2$ if L is odd. If only nearest neighbor hoppings are allowed and if no two odd sites are coupled (if the system has open boundary conditions for any L or periodic boundary conditions for L-even), each hopping either increases n_{even} by one and decreases n_{odd} by one, or vice-versa. This means that each hopping can change Δ_a only by ± 1.

Another way to sort configurations into two sets is by parity of the distance from the

configuration $|111...111\rangle$, which we define as

$$d_a = \min_n\{\langle 111...111|\hat{H}_1^n|a\rangle \neq 0\}. \tag{2.7}$$

In this case, the two sets are the configurations with even and with odd distances d_a. One hopping can change d_a only by ± 1 or 0. Changes by other values are not possible by definition if the Hamiltonian is Hermitian (all hoppings are reversible). Both d_a and Δ_a have the same parity, thus d_a must always change after one hopping in even system sizes or in systems with open boundary conditions. As a consequence, d_a cannot change by 0 if Δ_a can only change by ± 1.

We have shown above that the $\hat{H}_1$ model is bipartite for open boundary conditions irrespective of the system size L parity or for periodic boundary conditions when L is even. Now we prove that this property of $\hat{H}_1$ holds true when L is odd and filling factor $\nu = 1$. Due to the constraints imposed by $\hat{H}_1$, a particle cannot hop to an empty site to its left. At the filling factor $\nu = 1$, all configurations except 111...111 contain at least one empty site. These configurations can be connected to 111...111 by hoppings only to the right, which is also the shortest possible path d_a, defined in Eq. (2.7). The empty sites can be filled only with particles that come from the site on their left, as hopping from the other side is forbidden. This implies that at least for one pair of adjacent sites (an empty site and the filled one on its right) there will be no particles hopping between them on the path to 111...111. We can then redefine the numbering of sites to start from the filled site in this pair. This is equivalent to setting the right (filled) site as the first and the left (empty) site as the last site in the chain and imposing open boundary conditions. In this way, no two odd sites will be coupled and the argument that the absolute difference between the numbers of particles on even and odd sites can only change by ± 1 will still be valid.

Unlike $\hat{H}_1$, $\hat{H}_2$ in the same geometry is not bipartite. The reason for this is that there are no constraints in the case of $\hat{H}_2$, so the shortest path to 111...111 can include hoppings both to the right and to the left, which means that it is not always possible to choose the numbering in such a way that no two odd sites are coupled. Because of the open boundary conditions, the Hamiltonian $\hat{H}_3$ in its largest connected component is also bipartite for all system sizes.

The graphs of bipartite systems do not contain any loops of odd dimension (triangles, pentagons, heptagons and so on). Moreover, the energy spectra of bipartite systems are symmetric around zero. Their Hamiltonians anticommute with the operator $(-1)^{\Delta_a}$. The presence of such an operator in a bipartite lattice leads to exact zero-energy states in the spectrum [143, 144]. It can be shown that the exponentially growing number of zero modes of $\hat{H}_1$ is related to the difference between the numbers of elements in the two sets of its bipartite graph, as explained in Section 2.4. Additionally, the algebraic structure of zero-energy eigenstates can be explained by the structure of the graph – such eigenstates can be constructed as superpositions of config-

urations from only one of the sets. Similar properties are found for $\hat{H}_2$ for even L, as its graph is also bipartite in that case. The properties of the zero-energy manifold are discussed in more detail in Section 2.4.

2.2 Dynamics and entanglement properties

We now investigate the phenomenology of the models introduced in Eqs. (2.1)-(2.3). We use exact diagonalization to obtain the complete set of energy eigenvalues and eigenvectors, from which we evaluate the level statistics and the distribution of entanglement entropies for the three models. Furthermore, we probe dynamical properties of the models by studying a global quench, simulated via Krylov iteration.

2.2.1 Level statistics and entanglement entropy

The energy level statistics is a standard test for thermalization of models that cannot be solved exactly. A convenient way to probe the level statistics is to examine the probability distribution $P(r)$ [145] of ratios between consecutive energy gaps $s_n = E_{n+1} - E_n$,

$$r = \frac{\min(s_n, s_{n+1})}{\max(s_n, s_{n+1})}. \tag{2.8}$$

The advantage of studying $P(r)$, instead of $P(s_n)$, is that there is no need to perform the spectrum unfolding procedure – see Ref. [94]. For standard random matrix theory ensembles, both $P(r)$ and the mean $\langle r \rangle$ are well-known [146]. When computing the same quantities in a microscopic physical model, it is crucial to resolve all the symmetries of the model.

The probability distribution $P(r)$ of the ratios of two consecutive energy gaps is shown in Figs. 2.2(a), (b) and (c) for the three Hamiltonians $\hat{H}_1$, $\hat{H}_2$ and $\hat{H}_3$ respectively, and two momentum or inversion sectors. In all three cases, the energy levels repel, i.e., the distribution tends to zero as $r \to 0$. For $\hat{H}_2$, the distribution is particularly close to the Wigner-Dyson (non-integrable) line. For $\hat{H}_1$, the distribution is also consistent with Wigner-Dyson when we restrict to the middle 1/3 of the spectrum (and after removing special states with $E = 0$). We exclude the edges of the spectrum because they contain degeneracies which are not symmetry-related. However, such states do not appear to have a major effect on the level statistics distribution, which is still closer to the Wigner-Dyson than the Poisson distribution even if they are included. The level statistics of $\hat{H}_3$ within the largest connected component of the Hilbert space is shown in Fig. 2.2(c) and is also consistent with the Wigner-Dyson distribution without restricting the spectrum. However, we will demonstrate below that the dynamics in some smaller connected components of $\hat{H}_3$ can be exactly solved.

As a complementary diagnostic of thermalization, we next compute the entanglement en-

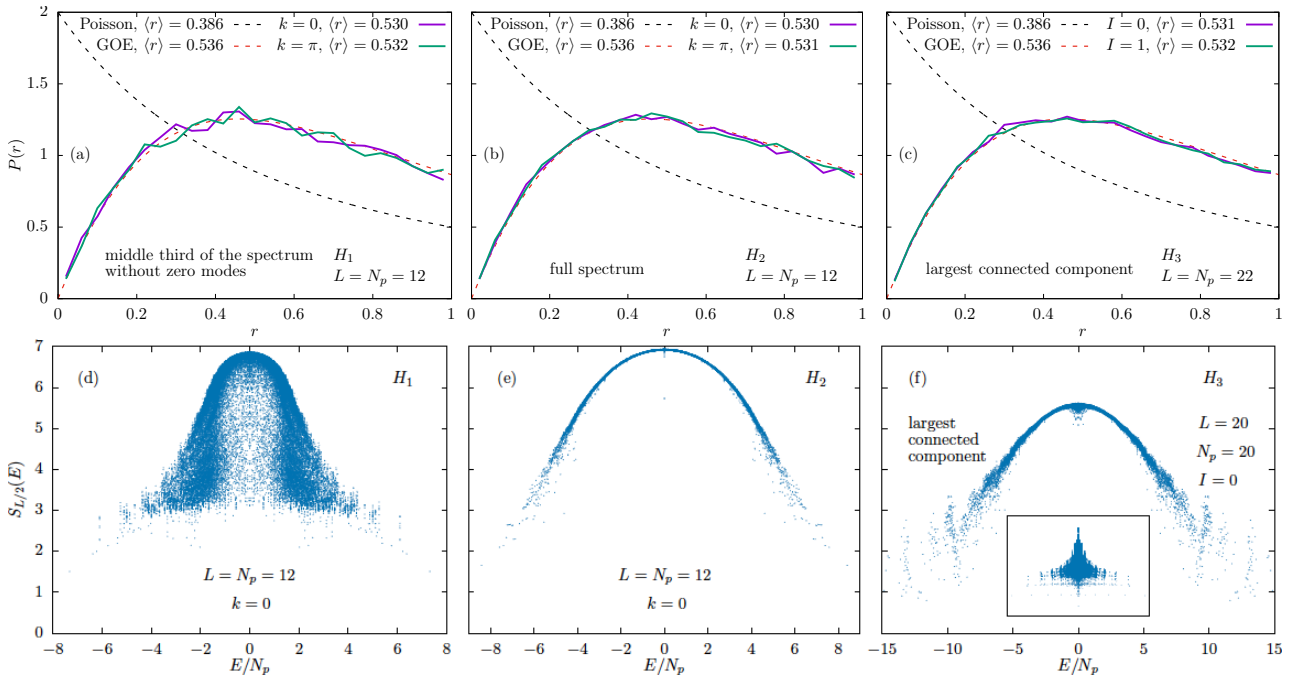

Figure 2.2: Level statistics and entanglement. (a), (b) and (c): Probability distribution of the ratios of two consecutive energy gaps. (a) $\hat{H}_1$ (middle third of the spectrum without $E = 0$ states, $L = N_\mathrm{p} = 12$), (b) $\hat{H}_2$ (full spectrum, $L = N_\mathrm{p} = 12$) and (c) $\hat{H}_3$ (largest connected component of $L = N_\mathrm{p} = 22$). The black dashed line shows the Poisson distribution, which corresponds to the integrable case, while the red dashed line is the distribution of the Gaussian orthogonal ensemble, which corresponds to the thermalizing case. (d), (e) and (f): Entanglement entropies $S_{L/2}$ of all eigenstates plotted as a function of the eigenstate energy per particle, E/N_p. (d) $\hat{H}_1$ ($L = N_\mathrm{p} = 12$, $L_\mathrm{A} = 6$, $k = 0$), (e) $\hat{H}_2$ (same) and (f) $\hat{H}_3$ in the largest connected component of $L = N_\mathrm{p} = 20$, $L_\mathrm{A} = 10$, $I = 0$. The inset shows all connected components for $L = N_\mathrm{p} = 12$, $L_\mathrm{A} = 6$, $k = 0$.

tropy of all eigenstates. We divide the lattice into two sublattices, A and B, of lengths L_A and $L_\mathrm{B} = L - L_\mathrm{A}$. For a given pure state $|\psi\rangle$, the entanglement entropy is defined as

$$S_\mathrm{A} = -\mathrm{Tr}_\mathrm{A}(\rho_\mathrm{A} \ln \rho_\mathrm{A}), \tag{2.9}$$

where $\rho_\mathrm{A} = \mathrm{Tr}_\mathrm{B}|\psi\rangle\langle\psi|$ is the reduced density matrix of the subsystem A. The scatter plots, showing entanglement entropy of all eigenstates $|E_n\rangle$ as a function of their energy E_n, are displayed in Figs. 2.2(d), (e) and (f). Here we take into account the translation symmetry of the system and work in the momentum sector $k = 0$ for $\hat{H}_1$ and $\hat{H}_2$, and consider only the largest connected component and the inversion sector $I = 0$ for $\hat{H}_3$. The results for other sectors are qualitatively similar.

Entanglement entropy distribution in Figs. 2.2(d) and (e) reveals a striking difference between the Hamiltonians $\hat{H}_1$ and $\hat{H}_2$, even though they only differ by a free-boson hopping term, Eq. (2.4). The model $\hat{H}_1$ is constrained, which leads to a large spread of the entropy distribution and many low-entropy eigenstates including in the bulk of the spectrum. From this perspective, $\hat{H}_1$ is reminiscent of PXP model [48, 53]. By contrast, $\hat{H}_2$ has no such constraints

and in this case the entanglement entropy is approximately a smooth function of the eigenstate energy. The Hamiltonian $\hat{H}_3$ is doubly constrained, and this is reflected in its entanglement distribution, which also shows a large spread and several disconnected bands, reminiscent of an integrable system like the XY model [147].

2.2.2 Global quenches

The constraints in the models in Eqs. (2.1), (2.2) and (2.3) have significant effects on the dynamics governed by these Hamiltonians. We probe the dynamics by performing a global quench on the system. We assume the system is isolated and prepared in one of the Fock states, $|\psi_0\rangle$, at time $t = 0$. We restrict to $|\psi_0\rangle$ being product states which are not necessarily translation-invariant, as such states are easier to prepare in experiment. However, our results remain qualitatively the same if we consider translation-invariant $|\psi_0\rangle$. After preparing the system in the state $|\psi_0\rangle$, which is not an eigenstate of the Hamiltonian, the system is let to evolve under unitary dynamics,

$$|\psi(t)\rangle = \exp\left(-\frac{i}{\hbar}\hat{H}t\right)|\psi_0\rangle. \tag{2.10}$$

where $\hat{H}$ is one of the Hamiltonians of interest. From the time-evolved state, we evaluate the quantum fidelity,

$$F(t) = |\langle\psi_0|\psi(t)\rangle|^2, \tag{2.11}$$

i.e., the probability for the wave function to return to the initial state. In a general many-body system, fidelity is expected to decay as $F(t) \sim \exp(-L(Jt)^2)$. It thus becomes exponentially suppressed in the system size for any fixed time t^*, i.e., $F(t^*) \sim \exp(-cL)$, where c is a constant. In scarred models, such as the Rydberg atom chain, fidelity at the first revival peak occurring at a time T still decays exponentially, but exponentially slower, i.e., $F(T) \sim \exp(-c'L)$, with $c' \ll c$. In Ref. [48], for a finite system with $L \lesssim 32$ atoms, the fidelity at the first revival can be as high as $\sim 70\%$, and several additional peaks at times nT are also clearly visible.

We first consider the Hamiltonian $\hat{H}_1$. Several configurations exhibit periodic revivals of the fidelity $F(t)$, which can in some cases be higher than 90%. Most of these configurations involve a very dense cluster of bosons such as $|...0N10...\rangle$. In contrast, a completely uniform configuration $|...111...\rangle$ thermalizes very quickly. Here we focus on periodically-reviving configurations with density being as uniform as possible. One family of such reviving configurations involves n unit cells made of 3 lattice sites:

$$|210210\ldots210\rangle \equiv |(210)^n\rangle. \tag{2.12}$$

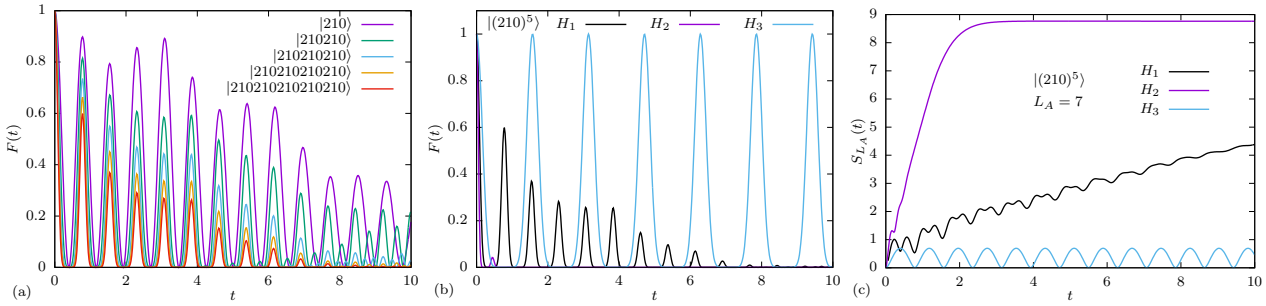

Figure 2.3: Dynamics of quantum fidelity and entanglement entropy for initial configurations in Eq. (2.12). (a) Time evolution of fidelity $F(t)$ in Eq. (2.11) for system sizes $L = 3n$. The evolution is governed by the Hamiltonian $\hat{H}_1$, different colors represent different system sizes L. (b) Fidelity evolution $F(t)$ for the Hamiltonians $\hat{H}_1$, $\hat{H}_2$ and $\hat{H}_3$ and system size $L = 15$. (c) Entanglement entropy evolution $S_{L_A}(t)$ for the same cases as in (b).

Time evolution of the fidelity for the initial state $|(210)^n\rangle$ for different system sizes $L = 3n$ is shown in Fig. 2.3(a). The initial state is assumed to be the product state, e.g., $|\psi_0\rangle = |210\rangle$ for $L = 3$. The frequency of the revivals in Fig. 2.3 is approximately the same for all system sizes. We emphasize that similar results are obtained for a translation-symmetric initial state, e.g., $|\psi_0\rangle = \frac{1}{\sqrt{3}}(|210\rangle + |021\rangle + |102\rangle)$. Both cases converge in the large system limit, and the differences are only significant for $L = 3$ when the revival frequency of the initial state with transition symmetry differs from the frequencies of other system sizes.

In Fig. 2.3(b) we compare the fidelity for the initial state in Eq. (2.12) when it is evolved by all three Hamiltonians in Eqs. (2.1)-(2.3). The initial state is fixed to be $|(210)^5\rangle$. We observe that the dynamics with $\hat{H}_3$ has very prominent revivals; in fact as we will later show, these revivals are perfect and their period is approximately twice the revival period for $\hat{H}_1$. In contrast, for $\hat{H}_2$ the fidelity quickly drops to zero without any subsequent revivals.

Finally, in Fig. 2.3(c) we plot the time evolution of entanglement entropy. As expected from the fast decay of the fidelity, the entropy for $\hat{H}_2$ rapidly saturates to its maximal value. Moreover, as expected from the perfect revivals in $\hat{H}_3$, the entropy in that case oscillates around a constant value close to zero. For $\hat{H}_1$, we observe a relatively slow growth of entropy, with oscillations superposed on top of that growth, again similar to PXP model [47]. For the initial state that is not translation-invariant, it is important how we cut the system, e.g., $|...210|210...\rangle$ versus $|...2102|10...\rangle$. In the first case, the entanglement entropy remains zero for $\hat{H}_3$ because no particle can hop from one subsystem to the other, while in the second case the entropy oscillates around a constant value, which is the case in Fig. 2.3(c).

In Fig. 2.4 we show the $\hat{H}_1$ evolution of two local observables, density correlations between two adjacent sites $\langle \hat{n}_1 \hat{n}_2(t) \rangle$ and density on the first site $\langle \hat{n}_1(t) \rangle$, starting from the initial state $|(210)^n\rangle$. Unlike fidelity and entanglement entropy, these observables can be easily measured in experiment. Both observables robustly oscillate with approximately the same frequency as the fidelity. The heights of the first few revival peaks are approximately converged for the system

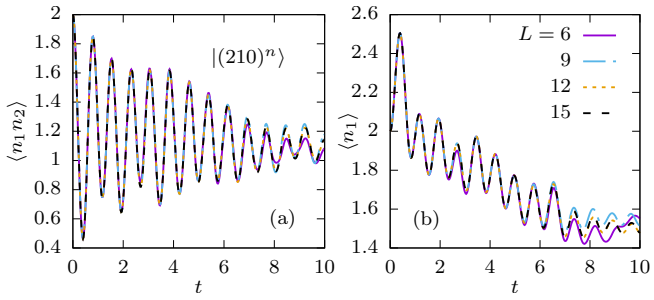

Figure 2.4: Evolution of local observables for the Hamiltonian $\hat{H}_1$. (a) Correlations between adjacent sites $\langle \hat{n}_1 \hat{n}_2(t) \rangle$ for different system sizes and the initial state $|(210)^n\rangle$. (b) Density on one site $\langle \hat{n}_1(t) \rangle$.

sizes ranging from $L = 6$ to $L = 15$, which suggests that revivals in such local observables can be observed in the thermodynamic limit. In the following Section, we will show that the oscillations observed in the dynamics from $|(210)^n\rangle$ state in Eq. (2.12) and their frequency can be explained using a tractable model that involves only a small subset of all configurations in the Hilbert space, thus providing a realization of quantum scars in a correlated bosonic system. Our starting point will be the model $\hat{H}_3$, whose graph explicitly separates into disconnected subsets which makes the toy model exact, hence we can analytically calculate the revival frequency. Based on these results, we then introduce an approximation scheme that describes the dynamics from the same initial state under the $\hat{H}_1$ Hamiltonian.

2.3 Quantum scars in $\hat{H}_1$ and $\hat{H}_3$ models

The quench dynamics of fidelity and entanglement entropy in Fig. 2.3 suggest that $\hat{H}_1$ and $\hat{H}_3$ models are candidate hosts for many-body scarred eigenstates that can be probed by initializing the system in product states $|(210)^n\rangle$. We now analyze the structure of these states using our approach called "cluster approximation".

2.3.1 Perfect revivals in the $\hat{H}_3$ model

The dynamics of $\hat{H}_3$ within the sector containing the state $|(210)^n\rangle$ can be solved exactly. We start with a warmup calculation for $\hat{H}_3$ acting on $L = 3$ sites. The connected subspace of 210 contains only two configurations, 120 and 210. The Hamiltonian reduced to this subspace is

$$\hat{H}_3' = -\begin{pmatrix} 0 & 2 \\ 2 & 0 \end{pmatrix}, \tag{2.13}$$

where the basis vectors are

$$\begin{pmatrix} 1 \\ 0 \end{pmatrix} = |210\rangle, \quad \begin{pmatrix} 0 \\ 1 \end{pmatrix} = |120\rangle. \tag{2.14}$$

The eigenvalues of $\hat{H}_3^{'}$ are $E_1 = -2$ and $E_2 = 2$. The initial state $|\psi_1(t=0)\rangle = |210\rangle$ evolves as

$$|\psi_1(t)\rangle = \cos(2t)|210\rangle - i\sin(2t)|120\rangle, \tag{2.15}$$

and the state $|\psi_2(t=0)\rangle = |120\rangle$ evolves as

$$|\psi_2(t)\rangle = -i\sin(2t)|210\rangle + \cos(2t)|120\rangle. \tag{2.16}$$

Previous results can be straightforwardly generalized to larger systems. Let the length of the system be $L = 3n$ for simplicity. The connected component of the state $|(210)^n\rangle$ consists only of combinations of patterns 210 and 120, which means that triplets of sites evolve independently, and dynamically the system behaves as a collection of independent two level systems (spins-1/2). From Eq. (2.15), the initial state $|\psi_n(t=0)\rangle = |(210)^n\rangle$ evolves as

$$|\psi_{L=3n}(t)\rangle = \cos^n(2t)|(210)^n\rangle + (-i)^n\sin^n(2t)|(120)^n\rangle + ... \tag{2.17}$$

where "..." denotes contributions of the basis configurations other than $|(210)^n\rangle$ or $|(120)^n\rangle$. The fidelity is

$$F_{L=3n}(t) = |\langle\psi_n(0)|\psi_n(t)\rangle|^2 = |\cos 2t|^{2n}. \tag{2.18}$$

It follows that the revivals are perfect, with a period $T_3 = \pi/2$. This result is also valid for the translation-invariant initial state $|(210)^n\rangle_{\mathrm{T}}$,

$$|(210)^n\rangle_{\mathrm{T}} \equiv \frac{1}{\sqrt{3}}\left(|(210)^n\rangle + |(021)^n\rangle + |(102)^n\rangle\right), \tag{2.19}$$

as the connected subspaces of 210, 021 and 102 do not overlap and therefore evolve independently.

However, an initial state that is both translation symmetric and inversion symmetric has different dynamics. The inverse of the configuration $|(210)^n\rangle$ is the configuration $|(012)^n\rangle$, which is a translation of the state $|(120)^n\rangle$ that belongs to the connected subspace of $|(210)^n\rangle$. The initial state

$$|\psi_n^{\mathrm{inv}}(t=0)\rangle = \frac{1}{\sqrt{2}}|(210)^n\rangle_{\mathrm{T}} + \frac{1}{\sqrt{2}}|(120)^n\rangle_{\mathrm{T}} \tag{2.20}$$

evolves as

$$|\psi_n^{\mathrm{inv}}(t)\rangle = \left(\cos^n 2t + (-i)^n\sin^n 2t\right)|\psi_n^{\mathrm{inv}}(t=0)\rangle + ... \tag{2.21}$$

and the fidelity is

$$F_n^{\text{inv}}(t) = |\langle\psi_n^{\text{inv}}(0)|\psi_n^{\text{inv}}(t)\rangle|^2 = |\cos^n 2t + (-i)^n \sin^n 2t|^2. \tag{2.22}$$

The frequency of the revivals is now doubled, so the period is $T_3^{\text{inv}} = \pi/4$.

2.3.2 Cluster approximations for the $\hat{H}_1$ model

In contrast to the free dynamics in $\hat{H}_3$, the $\hat{H}_1$ model exhibits decaying revivals and does not admit an exact description. In order to approximate the quench dynamics and scarred eigenstates in $\hat{H}_1$, we project the Hamiltonian to smaller subspaces of the full Hilbert space. These subspaces contain clusters of states which are poorly connected to the rest of the Hilbert space and thereby cause dynamical bottlenecks. The clusters can be progressively expanded to yield an increasingly accurate description of the dynamics from a given initial state.

Here we introduce a scheme for approximating the dynamics from initial states $(210)^n$ in the $\hat{H}_1$ model. As can be observed in Fig. 2.3, the revival periods are approximately the same for different system sizes. We first focus on the non-trivial case $L = 6$. Fig. 2.5 shows part of the graph that contains the initial state, $|210210\rangle$. Configurations labelled inside the ellipses denote representatives of an orbit of translation symmetry, i.e., the configurations are translation-invariant such as the one in Eq. (2.19).

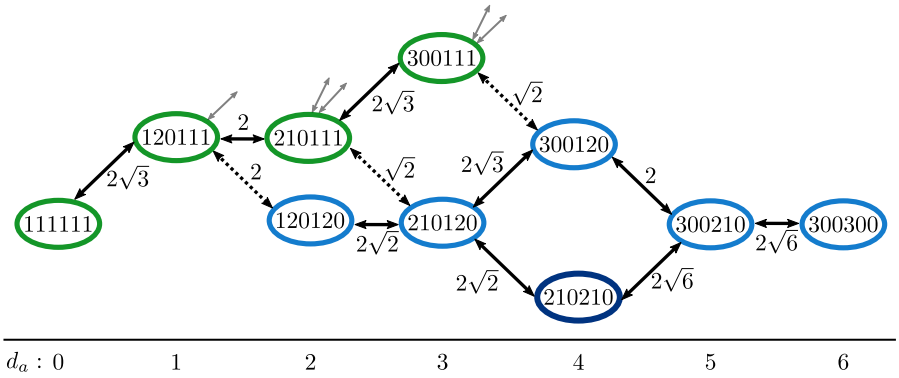

Figure 2.5: Minimal and extended clusters. Hamiltonian $\hat{H}_1$ and system size $L = N_{\text{p}} = 6$. Configurations labelled inside the ellipses are representatives of an orbit of translation symmetry. The minimal cluster is defined by the blue configurations, while green configurations represent the additional components of the extended cluster. Grey arrows connect to configurations outside the extended cluster. The numbers bellow the graph show the distance d_a from the configuration 111111 evaluated using Eq. (2.7).

The minimal subcluster of the graph is highlighted in blue color in Fig. 2.5. This cluster is indeed weakly connected to the rest of the configuration space, as it has only 3 connections that lead outside this cluster (dashed lines) and their hopping coefficients are slightly lower in magnitude than those inside the cluster, meaning that the probability is higher to stay inside the cluster than to leave. The hopping coefficients leading outside are not significantly smaller

than the coefficients staying inside, but in combination with the relatively small number of connections this has significant effects on the dynamics. This effect is even more pronounced when the difference in magnitudes is further increased by squaring the particle number operators, as shown in Appendix A.

The minimal cluster from Fig. 2.5 contains all the states given by tensor products of 210, 120 and 300 configurations. The set of configurations belonging to this cluster could have been chosen differently, but this particular choice has at least two advantages. Firstly, inside this cluster, the evolution of the configuration $|210210\rangle$ can be thought of as two subsystems 210 evolving separately. The evolution of all such configurations at different system sizes can be reduced to the evolution of $L = 3$ subsystems 210, similar to the case of $\hat{H}_3$ in the connected subspace of $(210)^n$. Secondly, this definition allows easy generalization to different system sizes $L = 3n$ with initial states $(210)^n$. The dimension of the reduced Hilbert space grows with the system size as $\dim\mathcal{H}^{\mathrm{c}} = 3^n$. We would like to emphasize that the cluster was not chosen arbitrarily. The calculations of the probability density distribution starting from the initial configuration $|210210\rangle$ and evolving with $\hat{H}_1$ have shown that the probability density stays high in this region of the Hilbert space as long as the revivals in fidelity are visible. The configurations important for the dynamics were then identified by analyzing the structure of the graph around the initial configuration.

As an example, consider system size $L = 3$. The reduced Hilbert space of the cluster $\mathcal{H}^{\mathrm{c}}$ is spanned by the (non-translation-invariant) configurations

$$\begin{pmatrix}1\\0\\0\end{pmatrix} = |300\rangle, \quad \begin{pmatrix}0\\1\\0\end{pmatrix} = |210\rangle, \quad \begin{pmatrix}0\\0\\1\end{pmatrix} = |120\rangle. \tag{2.23}$$

The Hamiltonian reduced to this subspace is

$$\hat{H}_1^{\mathrm{c}} = -\begin{pmatrix}0 & 2\sqrt{3} & 0\\ 2\sqrt{3} & 0 & 2\\ 0 & 2 & 0\end{pmatrix}, \tag{2.24}$$

and its eigenvalues are $E_1 = -4$, $E_2 = 4$, $E_3 = 0$. The initial configuration $|210\rangle$ evolves according to

$$|\psi_1^{\mathrm{c}}(t)\rangle = -\frac{i}{2}\sin(4t)\left(\sqrt{3}|300\rangle + |120\rangle\right) + \cos(4t)|210\rangle. \tag{2.25}$$

By generalizing this result to larger systems, it is easy to prove that the time-evolved state

within the cluster is given by

$$|\psi_n^c(t)\rangle = \cos^n(4t)|(210)^n\rangle + \ldots, \tag{2.26}$$

where the dots denote other configurations, and the fidelity is

$$F_n^c(t) = |\langle\psi_n^c(0)|\psi_n^c(t)\rangle|^2 = |\cos(4t)|^{2n}. \tag{2.27}$$

As in the case of $\hat{H}_3$, this result is also valid for the translation-invariant initial state. We see that the period of revivals is $T_1 = \pi/4$, which is the same as for $\hat{H}_3$ with a translation and inversion symmetric initial state.

The minimal clusters can be expanded by adding several neighboring configurations. For similar reasons as in the case of minimal clusters, the extended clusters are defined as sets of all states which can be obtained using tensor products of the configurations 210, 120, 300 and 111. In the case of $L = 6$, the enlarged cluster can be observed in Fig. 2.5. It contains the minimal cluster studied previously, but it also includes additional configurations shown in green ellipses. Again, the approximation could be improved by including more configurations, but this particular choice is well suited for analytical treatment (see Appendix B) and, as shown above, it gives a good prediction for the first revival peak height.

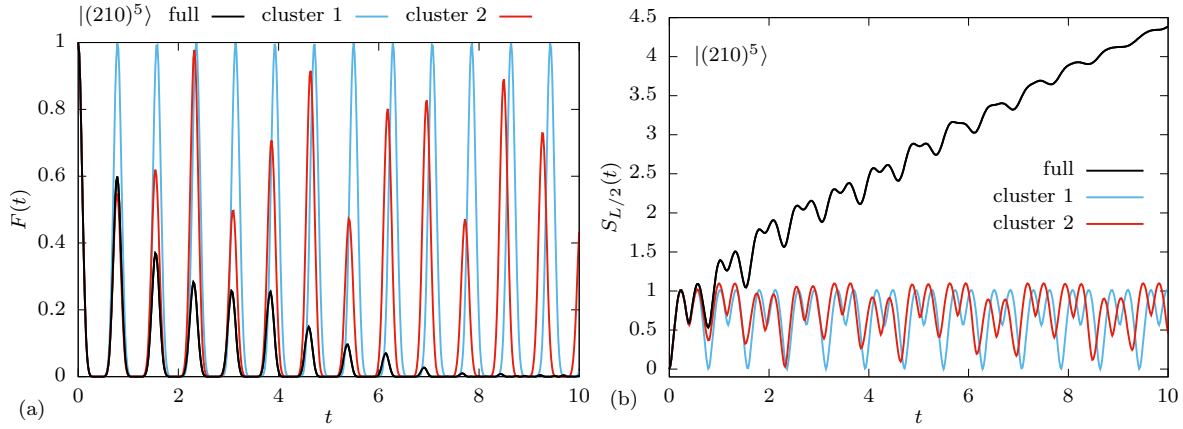

Figure 2.6: Comparison of the full dynamics against the minimal cluster (1) and extended cluster (2) approximation schemes. We consider the system size $L = 15$ with the initial state $|(210)^5\rangle$. (a) Time evolution of the fidelity. The frequency of revivals is approximately the same in both cases, but the results for the extended cluster show better agreement with the results for the full Hilbert space. (b) Time evolution of the entanglement entropy.

The result of the cluster approximation is compared against the exact result for system size $L = 15$ in Fig. 2.6. The frequency of the fidelity revival, shown by the blue line in Fig. 2.6(a), is accurately reproduced in this approximation, however the approximation does not capture the reduction in the magnitude of $F(t)$. Similarly, the dynamics of entanglement entropy, blue line in Fig. 2.6(b), is only captured at very short times. In particular, we observe that the maximum entanglement within the cluster remains bounded even at long times $t \sim 10$, while

the exact entropy continues to increase and reaches values that are several times larger.

To obtain a more accurate approximation, we can expand the minimal cluster with several neighboring configurations on the graph. We define the extended cluster as a set of all states which can be obtained using tensor products of the configurations 210, 120, 300 and 111. The enlarged cluster clearly contains the minimal cluster studied above, but it also includes additional configurations, resulting in a much better prediction for the first revival peak height, while still allowing for analytical treatment. The dimension of the extended cluster grows as $\dim\mathcal{H}^{\tilde{c}} = 4^{L/3}$, and is thus exponentially larger than the minimal cluster approximation. Nevertheless, the extended cluster dimension is still exponentially smaller compared to the full Hilbert space, and within this approximation it is possible to numerically simulate the dynamics of larger systems, $L \lesssim 30$ – see Fig. 2.7(a). The revivals are no longer perfect, while their frequency is independent of the system size and closer to the frequency of revivals for the full Hilbert space compared to to the minimal cluster approximation in Fig. 2.6. The overlap between the eigenstates of the Hamiltonian $\hat{H}_1$ reduced to both the minimal and extended cluster and the state $|(210)^8\rangle$ is given in Fig. 2.7(b). The eigenstates that correspond to the minimal cluster approximately survive in the extended cluster, where they form a band with the highest overlap.

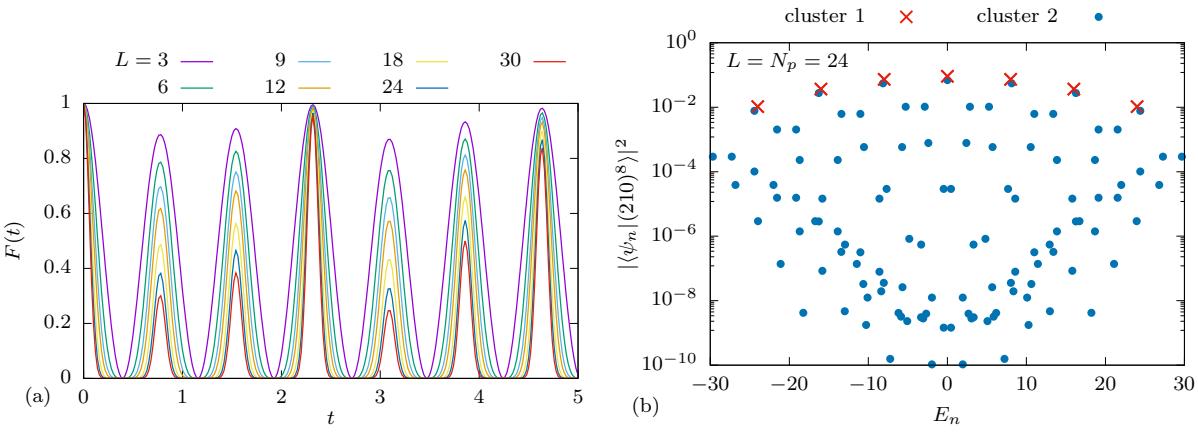

Figure 2.7: Cluster approximations. (a) Fidelity $F(t)$, for the Hamiltonian $\hat{H}_1$ and initial states $|(210)^n\rangle$, in the extended cluster approximation for various system sizes. (b) Eigenstate overlap with the initial state $|(210)^8\rangle$ plotted on a log scale, for both cluster approximations. In the case of degenerate eigenstates the sum of their overlaps is shown.

For the initial product state $(210)^n$, it is possible to analytically obtain the fidelity within the improved approximation for arbitrary system size. Similar to the previous methods, it can be shown (see Appendix B)

$$F^{\tilde{c}}_{L=3n}(t) \quad = \quad 4^n|b^2\cos(\alpha t) + d^2\cos(\beta t)|^{2n}, \tag{2.28}$$

where $\alpha = \sqrt{9+\sqrt{57}} \approx 4.06815$, $\beta = \sqrt{9-\sqrt{57}} \approx 1.20423$, $b \approx 0.694113$ and $d \approx 0.134933$. Eq. (2.28) is in excellent agreement with the numerical results in Fig. 2.7(a). It was also found

to be a very good approximation for the translation-invariant initial state when $L \geq 9$ (data not shown).

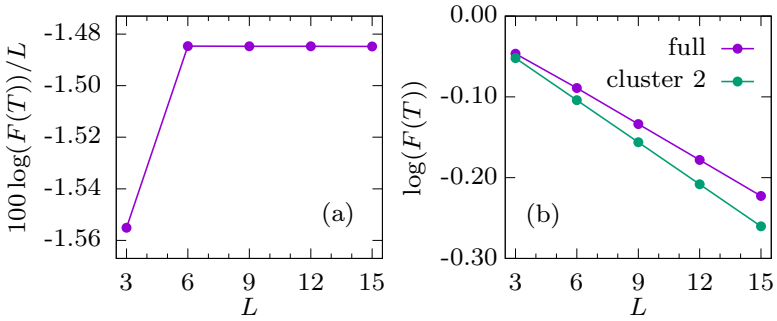

Figure 2.8: First peak height. (a) Logarithm (base 10) of the first revival peak divided by the system size, $\log(F(T))/L$, seems to saturate at a finite value in the thermodynamic limit. (b) Comparison of the logarithm of the first revival peak height for the full dynamics and the improved cluster approximation. The approximation serves as a lower bound.

Fig. 2.8(a) shows that the logarithm of the fidelity per site, $\log(F(T))/L$, at the first peak, saturates at a finite value for large L. In the improved cluster approximation, the first peak height decays as $\mathrm{e}^{-0.04L}$, as shown in Appendix B. For a completely random state, the fidelity would be $F \sim 1/\dim\mathcal{H}$. In the case $\nu = 1$ and large L, the Hilbert space dimension grows with the system size as

$$\dim\mathcal{H} = \binom{2L-1}{L} \approx \binom{2L}{L} \approx \frac{4^L}{\sqrt{\pi L}}. \tag{2.29}$$

This back-of-the-envelope estimate suggests the fidelity of a random state is $F \sim \mathrm{e}^{-1.39L}$, which decays considerably faster than the first peak height in Fig. 2.8. The improved cluster approximation correctly reproduces the short-time dynamics, including the first revival peak, and sets a lower bound for the first peak height – see Figs. 2.6 and 2.8(b).

The evolution of the entanglement entropy for the extended cluster approximation is shown in Fig. 2.6(b). Inside the cluster, entropy remains approximately constant with periodic oscillations that have the same frequency as the wave function revivals. Any further growth of the entanglement entropy can be attributed to the spreading of the wave-function outside the cluster. To illustrate the "leakage" of the wave function outside the cluster, in Fig. 2.9 we compute the time evolution of the overlap with a cluster, i.e., the probability to remain inside a cluster at time t,

$$\mathcal{O}_{\text{Cluster}} = \sum_{a \in \text{Cluster}} |\langle a|\psi(t)\rangle|^2. \tag{2.30}$$

We consider several initial configurations that lie inside or outside the cluster. The configurations initially inside the cluster mostly stay there, and the configuration $|(210)^4\rangle$ that has the highest revivals also has the highest overlap. Similarly, configurations initially outside the cluster continue to have negligible overlaps. The overlap starting from the configuration $|(210)^4\rangle$

approximately predicts the revival peak heights for the full dynamics.

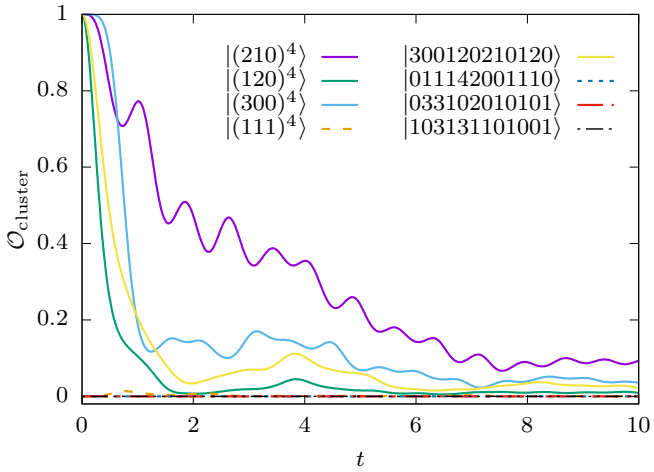

Figure 2.9: Evolution of the probability to remain inside the minimal cluster. $\mathcal{O}_{\text{Cluster}}$, as defined in Eq. (2.30). Initial configurations are indicated in the legend. Solid lines: configurations initially inside the cluster. Dashed lines: configurations initially outside the cluster (all except $|(111)^4\rangle$ are randomly chosen). Similar results are obtained for the extended cluster (not shown). System size $L = 12$.

We now summarize the relation between $\hat{H}_3$ and $\hat{H}_1$ from the point of view of the cluster approximation. For the initial state $|(210)^n\rangle$, the two models yield similar dynamics, compare Eqs. (2.18) and Eq. (2.28). The only difference is that the revival frequency is doubled in the latter case, which can be easily explained by the symmetry of the initial state and that of the Hamiltonian. Hamiltonian $\hat{H}_3$ is inversion-symmetric. If the initial state is also chosen to be inversion-symmetric, the frequency of the revivals doubles. The period is then $T_3^{\text{inv}} = \pi/4$, which is equal to the period of revivals T_1 of $\hat{H}_1$ in the cluster approximation. This was proven analytically in Section 2.3.1, see Eq. (2.22). For comparison, the revival period for the full Hilbert space is approximately 0.77, which is slightly less than $\pi/4 \approx 0.79$. The Hamiltonian $\hat{H}_1$ is not inversion-symmetric, so the frequency does not double for an inversion-symmetric initial state, but the revivals are lower in that case. This shows that it is important for the symmetry of the initial state to match the symmetry of the Hamiltonian.

Finally, the eigenstates of $\hat{H}_1$, projected to the subspace of the minimal cluster approximation, form several degenerate bands whose eigenenergies are equally spaced in integer multiples of 4. Interestingly, some of these eigenstates approximately survive in the full $\hat{H}_1$ model, and they are precisely the eigenstates that have the highest overlap with the configurations $|(210)^n\rangle$, see Fig. 2.10(a). In small system sizes, such as $L = 6$, the surviving eigenstates are also the lowest entropy eigenstates in the middle of the spectrum, which is reminiscent of quantum scars in the PXP model [48]. In larger systems, e.g., $L = 12$, the surviving eigenstates are slightly lower in entropy than their neighbors, but are far from being the lowest entropy eigenstates, as can be seen in Fig. 2.10. The lowest entropy eigenstates have high overlaps with other configurations, such as $|(3100)^3\rangle$, as shown in Figs. 2.10(b) and 2.10(c). In the case of $|(210)^n\rangle$, the

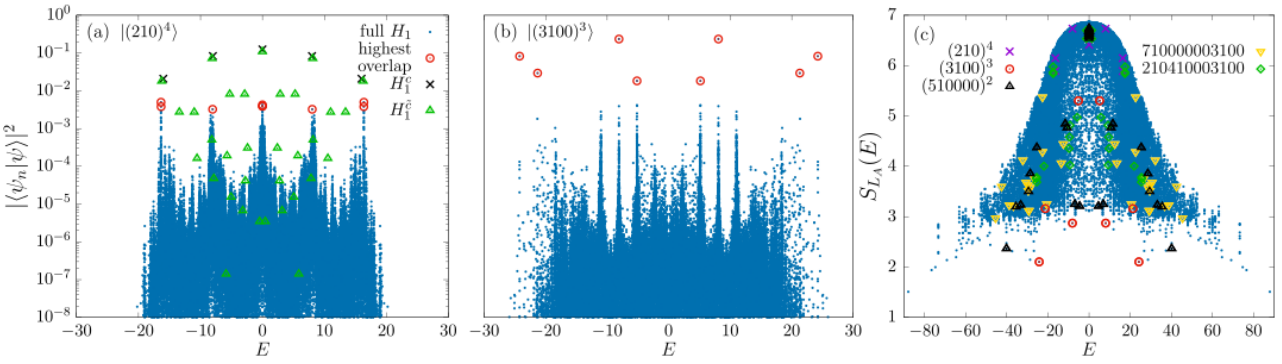

Figure 2.10: Non-ergodic eigenstates. (a) Overlap of the configuration $|(210)^4\rangle$ with all the eigenstates of $\hat{H}_1$, $\hat{H}_1^c$ and $\hat{H}_1^{\bar{c}}$ versus the eigenstate energy for sector $k = 0$ and system size $L = N_p = 12$. (b) Same for $|(3100)^3\rangle$. (c) Entanglement entropy, eigenstates which have the highest overlap with some product states are marked in different colors.

eigenstates surviving in the full system belong to every other band of eigenstates in the cluster approximation and the number of the surviving eigenstates is $n + 1$. For even system sizes this counting includes a zero-energy eigenstate. In Section 2.3.3 we discuss in more details the generalization of the cluster approximations to the states of the form $|(N10...0)^n\rangle$, which were also found to have robust revivals and high overlaps with some low-entropy eigenstates.

2.3.3 Generalization to other clusters

Building on the previous observation that some of the low-entropy eigenstates have large weight on $|(3100)^3\rangle$ product state, we have investigated periodic revivals from such a larger class of initial states. We find that robust revivals are associated with initial product states of the form

$$|((N-1)1\underbrace{0...0}_{N-2})^n\rangle, \tag{2.31}$$

where N is the length of the unit cell ($L = Nn$). For example, some of these configurations are $|(3100)^n\rangle$, $|(41000)^n\rangle$ and $|(510000)^n\rangle$. Combinations of those patterns such as $|310041000\rangle$ also exhibit similar properties, but we will restrict ourselves to the simpler former cases.

We can construct a generalization of the cluster approximation for configurations of the form in Eq. (2.31). As in the case of $|(210)^n\rangle$, the dynamics inside one unit cell explains the dynamics of the full system. The generalized clusters can be chosen in such a way that their Hilbert spaces are spanned by N configurations

$$|i\rangle = |((N+1-i)(i-1)\underbrace{0...0}_{N-2})^n\rangle, \tag{2.32}$$

where i takes values $1, 2, \ldots N$. If we consider only one unit cell ($n = 1$), the graph that connects these configurations has a linear structure without any loops, i.e., each configuration $|i\rangle$ is solely connected to the configurations $|i \pm 1\rangle$, except the two configurations at the edges,

$|1\rangle$ and $|N\rangle$, which are only connected to $|2\rangle$ and $|N-1\rangle$, respectively.

The projection of the Hamiltonian $\hat{H}_1$ to this cluster, which we denote by $\hat{H}_1^{\mathrm{c}}$, has a very simple structure: it has the form of a tight-binding chain with the only nonzero matrix elements on the upper and lower diagonals:

$$H^{\mathrm{c}}_{1;i,i+1} = H^{\mathrm{c}}_{1;i+1,i} = (N-i)\sqrt{i(N+1-i)}. \tag{2.33}$$

The dynamics within a single unit cell under $\hat{H}_1^{\mathrm{c}}$ corresponds to density fluctuations between the first and the second site. Following the same procedure as previously, we can now diagonalize $\hat{H}_1^{\mathrm{c}}$ and compute the fidelity time series for the initial configuration $|(N-1)10...0\rangle$. This result can be directly generalized to configurations of the form $|((N-1)10...0)^n\rangle$. The derivation is valid for both translation-invariant and non-translation-invariant initial configurations, as the cluster in Eq. (2.32) is disconnected from its translated copies. We stress that this disconnection, namely the absence of a hopping term between $|1(N-1)0...0\rangle$ and $|0N0...0\rangle$, is a consequence of the constraints imposed by $\hat{H}_1$ and it would not hold for $\hat{H}_2$. In this way, we have calculated the time evolution of the fidelity starting from the configurations $|(3100)^n\rangle$ (for $n = 1, 2, 3, 4$), $|(41000)^n\rangle$ ($n = 1, 2, 3$) and $|(510000)^n\rangle$ ($n = 1, 2$), and compared it with the exact numerical results for the full $\hat{H}_1$. The cluster approximation captures both the revival frequency and the height of the first peak. Similar to the $|(210)^n\rangle$ case, the approximation can be improved by adding further configurations to the clusters. Moreover, if we want to consider translation-invariant initial states, we can follow the same recipe for $|(210)^n\rangle$ by summing translated patterns with the required phase factors given in terms of momenta in multiples of $2\pi/N$. We have checked that revivals appear in these momentum sectors, with roughly the same frequency.

We note that the configurations with larger units cells thermalize more quickly on shorter timescales, but slower at long times. Initially, the states starting from configurations with smaller N have lower entanglement entropies than those with larger N. The Hilbert spaces of large N unit cells are larger, so the entanglement entropy starting from these configurations rapidly grows to the maximal value for that unit cell. However, the only way for the wavefunction to spread through the entire Hilbert space is that a unit cell reaches a state close to 111...111, so that particles can hop to the other unit cells. This is less likely for large N, and therefore such configurations need long times to fully thermalize. As a result, smaller N entanglement entropies grow faster and after long enough time they overtake those for larger N. For example, in the case of $L = 12$ and translation-invariant initial states, $(210)^4$ overtakes $(3100)^3$ and $(510000)^2$ around $t \sim 2$, and $(3100)^3$ overtakes $(510000)^2$ around $t \sim 80$, as shown in Fig. 2.11.

Finally, we note that non-thermal behavior reminiscent of the one studied here was previously observed in an unconstrained Bose-Hubbard model, for example in the context of "arrested expansion" [148, 149] and quenches from superfluid to Mott insulator phase [150, 151]. In these

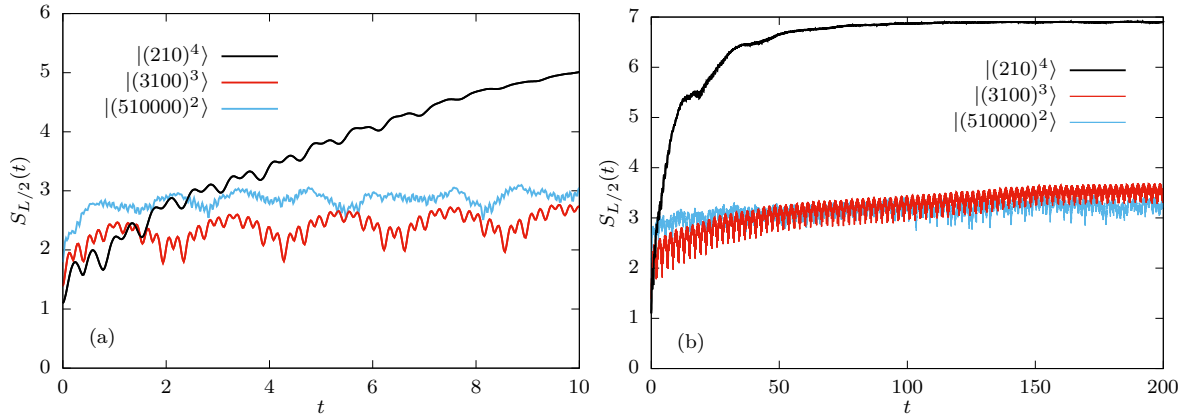

Figure 2.11: Time evolution of entanglement entropy for three different translation-invariant initial states which exhibit slow thermalization. (a) Short timescale. (b) Long timescale. System size $L = 12$. Configurations with larger unit cells (such as $|(510000)^2\rangle$), thermalize more quickly than those with smaller unit cells (such as $|(210)^4\rangle$) on shorter timescales, but slower at longer times.

cases, the main ingredient is the strong on-site interaction, which causes the energy spectrum to split into several bands. Due to the large energy differences between bands, the dynamics of an initial state from a particular band is at first limited only to the eigenstates that belong to the same band. Additionally, these energy bands are approximately equally spaced, which can lead to revivals in fidelity if several bands are populated. In contrast, our models do not feature on-site interaction, and the mechanism which slows down the spread of the wave function is correlated hopping, which suppresses connections between certain configurations and modifies the hopping amplitudes between others, thus creating bottlenecks that separate different clusters of states.

2.4 Zero modes

An interesting feature of $\hat{H}_1$ model is the large number of zero-energy states in the middle of its spectrum. The number of these zero modes, found by brute force diagonalization, is listed in Table 2.1 for different system sizes and momentum sectors. Similar property is found for $\hat{H}_2$ – see Table 2.2, with the notable difference that there are no zero modes when the number of sites L is odd.

The origin of the zero modes is the underlying bipartite structure of the Hamiltonian [143, 144]. As explained in Section 2.1.2, all the basis configurations of the $\hat{H}_1$ model can be separated into two disjoint classes, and the action of the Hamiltonian $\hat{H}_1$ only connects configurations in one class to the configurations in the other class, while $\hat{H}_1$ does not connect configurations within the same class. For example, a graph that shows how the configurations for $L = N_\mathrm{p} = 4$ are connected is displayed in Fig. 2.12. Here we will refer to the two classes as the "green" (even) and the "red" (odd) configurations. Each basis configuration can be uniquely assigned a

Table 2.1: The number of zero-energy states for the Hamiltonian $\hat{H}_1$ and different system sizes. The number of states is resolved per momentum sectors, denoted by their momentum indices i that parametrize the momenta $k_i = \frac{2\pi}{L} i$.

	$\frac{k_i L}{2\pi}$											
$L = N_\mathrm{p}$	0	1	2	3	4	5	6	7	8	9	10	total
2	0	1										1
3	0	1	1									2
4	2	0	1	0								3
5	2	1	1	1	1							6
6	2	3	0	4	0	3						12
7	2	3	3	3	3	3	3					20
8	10	0	8	0	9	0	8	0				35
9	8	8	8	7	8	8	7	8	8			70
10	4	25	2	25	2	26	2	25	2	25		138
11	22	23	23	23	23	23	23	23	23	23	23	252

Table 2.2: The number of zero-energy states for the Hamiltonian $\hat{H}_2$ and different system sizes. The columns are the same as in Table 2.1.

	$\frac{k_i L}{2\pi}$											
$L = N_\mathrm{p}$	0	1	2	3	4	5	6	7	8	9	10	total
2	0	1										1
3	0	0	0									0
4	2	0	1	0								3
5	0	0	0	0	0							0
6	0	3	0	4	0	3						10
7	0	0	0	0	0	0	0					0
8	10	0	8	0	9	0	8	0				35
9	0	0	0	0	0	0	0	0	0			0
10	0	25	0	26	0	26	0	26	0	25		128
11	0	0	0	0	0	0	0	0	0	0	0	0

green or a red label according to the parity of its distance d_a from the configuration $|111...111\rangle$, Eq. (2.7). If this number is even, the configuration is green, and if it is odd, the configuration is red.

This separation into two classes is a consequence of the constraints present in the Hamiltonian $\hat{H}_1$. Hamiltonians without such constraints, for example $\hat{H}_2$ or the standard Bose-Hubbard model, do not exhibit this bipartite structure for odd system sizes, see Section 2.1.2. for more details. In these cases, it is not possible to uniquely determine whether a particular configuration is green or red. The lack of bipartite structure is the reason for the absence of zero-energy eigenstates of $\hat{H}_2$ in odd dimensions, which was observed in Table 2.2. However, the configuration space of $\hat{H}_2$ is still bipartite in even dimensions, allowing for the existence of some zero modes in those cases.

Low-entropy zero-energy states can be constructed as superpositions of either only green or

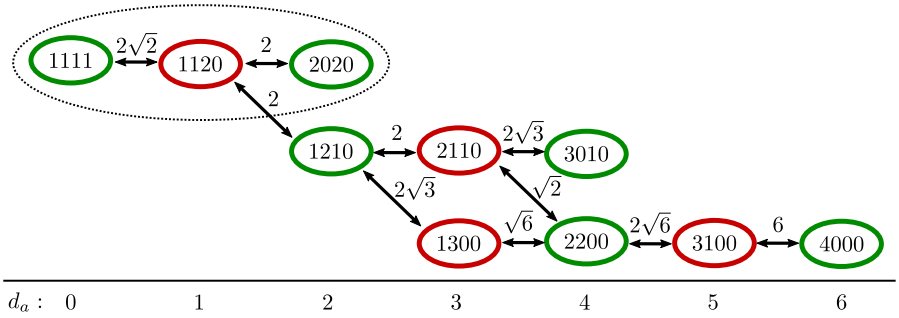

Figure 2.12: Bipartite graph for $\hat{H}_1$ and $L = N_\mathrm{p} = 4$. The two classes of configurations are shown in green (even) and red (odd) ellipses. The configurations are written in the translation-invariant basis. The arrows represent the action of the Hamiltonian $\hat{H}_1$ and the numbers above the arrows are the magnitudes of the corresponding hopping coefficients. The numbers bellow the graph show the distance d_a from the configuration 1111, as defined in Eq. (2.7).

only red configurations. For example, in the case of $L = N_\mathrm{p} = 4$, the simplest and therefore the lowest-entropy zero mode can be constructed using only two green product states (encircled by a dashed line in Fig. 2.12)

$$|\psi_0\rangle = \frac{1}{\sqrt{3}}|1111\rangle - \sqrt{\frac{2}{3}}|2020\rangle_\mathrm{T}, \tag{2.34}$$

where $|\ldots\rangle_T$ was defined in Eq. (2.19). There is another zero mode in this case, and it can be formed by adding more green configurations to the superposition. The number of zero-energy eigenstates is related to the difference between the numbers of green and red configurations [143], as we will now explain.

As the Hamiltonian $\hat{H}_1$ only connects green configurations to red configurations and red to green, we can rewrite it in the following way:

$$\hat{H}_1 = \sum_{i,j} c_{ij}|R_i\rangle\langle G_j| + \sum_{i,j} c_{ij}^\dagger|G_j\rangle\langle R_i|, \tag{2.35}$$

where $|R_i\rangle$ are the red product states and $|G_j\rangle$ are green. Its square, $\hat{H}_1^2$, connects green configurations to green and red to red, and it is therefore block diagonal. The blocks are $\hat{C}\hat{C}^\dagger$ and $\hat{C}^\dagger\hat{C}$, where $\hat{C}$ is a matrix with the elements c_{ij}. The dimensions of $\hat{C}$ are $r \times g$, where r is the number of red configurations and g of green. $\hat{C}$ and $\hat{C}^\dagger$ can be factorized using singular value decomposition. From this structure we can see that the energy spectrum is symmetric around zero and that the minimal number of zero-energy states is $|g - r|$. The zero-energy eigenvectors can also be obtained as the ground states of $\hat{H}_1^2$. Similar analysis and counting of the zero modes in PXP model was performed in Ref. [47, 152].

Table 2.3 shows the difference between the numbers of red and green states $g-r$ for different system sizes and the number of zero-energy states N_0 in those systems. The number of zero-energy states, found by exact diagonalization, in all cases satisfies the anticipated inequality

Table 2.3: The difference between the number of green and red configurations $g - r$ and the number of zero-energy states N_0 (determined by exact diagonalization) for different system sizes. Overall, the derived bound for the number of zero modes is found to be very tight in finite systems where it can be independently confirmed by explicit diagonalization ("NA" denotes cases where this was not possible).

	all sectors		$k = 0$	
$L = N_\mathrm{p}$	$g - r$	N_0	$g - r$	N_0
2	-1	1	0	0
3	-2	2	0	0
4	3	3	2	2
5	6	6	2	2
6	-10	12	0	2
7	-20	20	-2	2
8	35	35	10	10
9	70	70	8	8
10	-126	138	0	4
11	-252	252	-22	22
12	462	NA	80	80
13	924	NA	72	NA
14	-1716	NA	0	NA
15	-3432	NA	-228	NA
16	6435	NA	810	NA

$N_0 \geq |g - r|$. In fact, the bound is almost always saturated, $N_0 = |g - r|$, except when $L = N_\mathrm{p} = 4n + 2$, $n \in \mathbb{Z}$. Interestingly, the minimal number of zero modes $|g - r|$ for $L = N_\mathrm{p} = 4n$ is equal to the Hilbert space dimension for $L = N_\mathrm{p} = 2n$ (both total and for $k = 0$ sector only). This leads to the conclusion that the number of zero-energy states grows exponentially with the system size. It can also be noticed that the total difference $g - r$ for $L = N_\mathrm{p} = 2n + 1$ is always twice the difference for $L = N_\mathrm{p} = 2n$.

2.5 Conclusions

In this Chapter, we have introduced three models of bosons with "soft" kinetic constraints, i.e., density-dependent hopping. We have demonstrated that some of these models exhibit similar phenomenology to other realizations of quantum many-body scars, for example the Rydberg atom system [46]. We have studied quantum dynamics of these systems by performing global quenches from tensor-product initial states. We have shown that both the connectivity of the Hilbert space and the relative magnitude of the hopping coefficients have dramatic effects on the dynamics. For certain initial configurations, the constraints can lead to slow thermalization and revivals in the quantum fidelity. The revival frequency can be predicted by considering an exponentially reduced subset of the Hilbert space. For a family of initial configurations of the form $|(210)^n\rangle$, we have derived analytical expressions for the evolution of quantum fidelity

within this approximation, which accurately capture the revival frequency obtained from exact numerical data. One notable difference between scarred dynamics in the present bosonic models and the PXP model is that the revivals exist in the absence of a hard kinetic constraint, i.e., in the fully connected Hilbert space. Our cluster approximation also explains the structure of some low-entropy eigenstates in the middle of the many-body spectrum. In addition, we have calculated the evolution of two local observables which are experimentally measurable, density correlations between two neighboring sites and density on a single site, and both of them show robust oscillations over a range of system sizes. We have also shown that the introduced models contain additional special properties, like the exponentially large zero-energy degeneracy which is related to the bipartite structure of the model.

We now comment on the possible experimental realizations of the models we studied. The implementation of a correlated hopping term ($\hat{n}_k \hat{b}_i^\dagger \hat{b}_j$) in optical lattices has attracted lot of attention due to a possible onset of quantum phases related to high-Tc superconductivity [153]. An early theoretical proposal exploits asymmetric interactions between the two atomic states in the presence of a state-dependent optical lattice [153]. As a result, the obtained effective model corresponds to the inversion-symmetric form of $\hat{H}_1$. In addition, the same term has been found to feature as a higher-order correction of the standard Bose-Hubbard model [154–157]. Although in this case the term typically represents a modification of the regular hopping term of the order of several percent, its contribution was directly measured [158, 159]. More recently, the set of quantum models accessible in cold-atom experiments has been enriched through the technique of Floquet engineering [86]. As a notable example, a suitable driving scheme can renormalize or fully suppress the bare tunneling rate [160]. On top of that, by modulating local interactions an effective model with the density-dependent tunneling term has been engineered [161]. For the models considered in this Chapter the most promising is a more recent driving scheme exploiting a double modulation of a local potential and on-site interactions [162]. Related sophisticated driving schemes have already enabled a realization of dynamical gauge fields [124, 125, 163] where both the amplitude and the phase of the effective tunneling are density-dependent. Although these experimental proposals explain how to realize some of the correlated hopping terms present in our models using ultracold atoms in optical lattices, finding a scheme that exactly realizes our models requires further study. We emphasize that other models which would exhibit non-ergodic dynamics and scarred eigenstates as a result of the same mechanism that was explained in this work could be built, for example a linear combination of $\hat{H}_1$ and $\hat{H}_2$.

During the completion of this work, we became aware of Ref. [164] which identified non-thermal eigenstates and slow dynamics in the quantum East model. Moreover, a recent study [165] proposed a Floquet scheme for a bosonic model with density-assisted hopping, finding signatures of quantum many-body scars.

Chapter 3

Dynamics of weakly interacting bosons in optical lattices with flux

As already discussed in Section 1.4, a big challenge in the field of ultracold atoms was realization of synthetic magnetic fields, due to the fact that charge-neutral atoms do not feel the Lorentz force. Magnetic field is a key ingredient in various condensed-matter models with nontrivial topological content, such as the the Harper-Hofstadter [68] and the Haldane model [76]. In recent years, the implementation of artificial gauge potentials for neutral atoms [74, 75] has finally enabled the realization of these important models using ultracold atoms in optical lattices [77–80].

Cold-atom realizations of topological models exploit periodic driving, either through laser-assisted tunneling [77, 78] or by lattice shaking [79]. Using Floquet theory [81, 89], a periodically driven system can be related to the time-independent effective Hamiltonian that describes a condensed-matter system of interest. The mapping is known as Floquet engineering and its important features in the context of optical lattices are discussed in Section 1.5 and Refs. [82–88, 166]. Because of important differences of cold-atom setups and their condensed-matter counterparts, new quench protocols for probing topological features were proposed [167–171]. Following up on these studies, a novel experimental protocol was developed which allowed for the first-ever measurement of the Chern number (1.5) in a nonelectronic system by investigating the anomalous deflection of an atomic cloud as a response to external force [70]. The Chern number is a topological invariant which was directly related to the quantization of the Hall conductivity in the integer Hall effect [69].

While Floquet engineering is a highly flexible and powerful technique, it poses several concerns. One of the main open questions is related to the interplay of driving and interactions which causes heating and can quickly lead the system to a featureless, infinite-temperature regime [94, 96]. In particular, it is shown that an initial Bose-Einstein condensate in a periodically driven optical lattice may become unstable due to two-body collisions [172] or through

the mechanism of parametric resonance [96, 173–179]. The preparation protocol, stability and a lifetime of strongly correlated phases, expected in the regime of strong interactions under driving is a highly debated topic at the moment [96, 180, 181].

In order to further explore the role of weak atomic interactions in probing topological features, here we consider the dynamics of weakly interacting incoherent bosons in a driven optical lattice exposed to an external force. The setup that we consider includes all basic ingredients for the Chern-number measurement [70, 168] – the Chern number of the topological band can be extracted from the center-of-mass motion of atomic cloud in the direction transverse to the applied force. We assume an ideal initial state where the lowest topological band of the effective model is almost uniformly populated. The optimal loading sequence necessary to reach this state is considered in Refs. [182, 183]. Following the recent experimental study [70], we assume that atoms are suddenly released from the trap and exposed to a uniform force. We perform numerical simulations for the full time-dependent Hamiltonian and take into account the effects of weak repulsive interactions between atoms within the mean-field approximation. We make a comparison between the dynamics governed by the effective and time-dependent Hamiltonian and delineate the contribution of interactions to the center-of-mass response and to the overall cloud expansion dynamics. Our results show that interactions lead to the undesirable atomic transitions between topological bands [184], but we also find that a weak atomic repulsion can facilitate the Chern-number measurements in several ways.

This Chapter is organized as follows. In Section 3.1 we describe the model and introduce a method that we apply for the description of incoherent bosons. In Section 3.2 we address the dynamics of noninteracting incoherent bosons, and then in Section 3.3 we address the regime of weak repulsive interactions. Finally, we summarize our results in Section 3.4. Appendixes C to F provide further details.

3.1 Model and method

In this section, we first present the driven model introduced in Ref. [70], and then derive the corresponding effective model and discuss its basic characteristics. At the end, we explain our choice of the initial state and outline the method that we use to treat the dynamics of weakly interacting incoherent bosons.

3.1.1 Effective Floquet Hamiltonian

Interacting bosons in a two-dimensional optical lattice can be described by the Bose-Hubbard Hamiltonian

$$\begin{aligned}\hat{H}_{\mathrm{BH}} &= -J_x \sum_{l,m} \left(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}\right) - J_y \sum_{l,m} \left(\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}\right) \\ &+ \frac{U}{2}\sum_{l,m} \hat{n}_{l,m}\left(\hat{n}_{l,m}-1\right), \end{aligned} \tag{3.1}$$

where $\hat{a}^\dagger_{l,m}$ and $\hat{a}_{l,m}$ are creation and annihilation operators that create and annihilate a particle at the lattice site $(l,m) = la\mathbf{e}_x + ma\mathbf{e}_y$ (a is the lattice constant), $\hat{n}_{l,m} = \hat{a}^\dagger_{l,m}\hat{a}_{l,m}$ is the number operator, J_x and J_y are the hopping amplitudes along $\mathbf{e}_x$ and $\mathbf{e}_y$, and U is the on-site interaction. In the derivation of the model (3.1) we use the single-band tight-binding approximation [5]. Although the experimental setup [70] is actually three dimensional, with an additional confinement in the third direction, our study is simplified to a two-dimensional lattice.

In order to engineer artificial gauge field in the experiment [70], hopping along $\mathbf{e}_x$ was at first inhibited by an additional staggered potential

$$\hat{W} = \frac{\Delta}{2}\sum_{l,m}(-1)^l \hat{n}_{l,m}, \tag{3.2}$$

and then restored using resonant laser light. For more details, see Fig. 1.5 in Chapter 1. The experimental setup can be described by a time-dependent Hamiltonian

$$\tilde{H}(t) = \hat{H}_{\mathrm{BH}} + \hat{V}(t) + \hat{W}, \tag{3.3}$$

where $\hat{V}(t)$ is a time-dependent modulation

$$\begin{aligned}\hat{V}(t) = \kappa \sum_{l,m} \hat{n}_{l,m}\Bigg[&\cos\left(\frac{l\pi}{2}-\frac{\pi}{4}\right)\cos\left(\omega t - \frac{m\pi}{2}+\phi_0\right) \\ &+\cos\left(\frac{l\pi}{2}+\frac{\pi}{4}\right)\cos\left(-\omega t - \frac{m\pi}{2}+\frac{\pi}{2}+\phi_0\right)\Bigg], \end{aligned} \tag{3.4}$$

κ is the driving amplitude, and $\omega = \Delta$ is the resonant driving frequency. We set the relative phase ϕ_0 between the optical-lattice potential and the running waves used for laser-assisted tunneling to $\phi_0 = \pi/4$.

Using Floquet theory, see Eq. (1.9), the time-evolution operator corresponding to the Hamiltonian (3.3) can be represented as

$$\hat{U}(t,t_0) = \mathrm{e}^{-i\hat{W}t}\mathrm{e}^{-i\hat{K}(t)}\mathrm{e}^{-i(t-t_0)\hat{H}_{\mathrm{eff}}}\mathrm{e}^{i\hat{K}(t_0)}\mathrm{e}^{i\hat{W}t_0}, \tag{3.5}$$

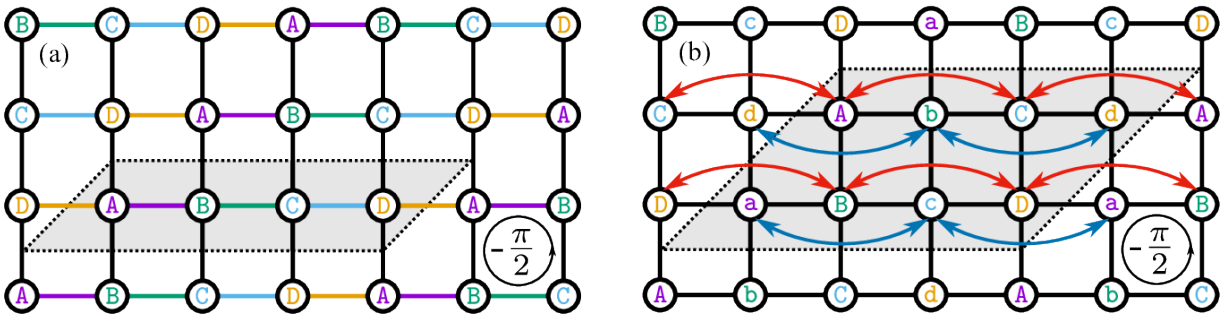

Figure 3.1: Schematic representation of the model. The unit cells are shaded. (a) Effective Hamiltonian without correction, $\hat{H}_{\text{eff},0}$ (3.6). Vertical links correspond to real hopping amplitudes (along $\mathbf{e}_y$ direction), while the horizontal links to the right of lattice sites labeled `A`, `B`, `C`, and `D` correspond to complex hopping amplitudes with phases $\frac{3\pi}{4}$, $\frac{\pi}{4}$, $-\frac{\pi}{4}$, and $-\frac{3\pi}{4}$, respectively (when hopping from left to right). (b) Effective Hamiltonian with correction, $\hat{H}_{\text{eff},1}$ (3.7). Red lines represent positive next-nearest-neighbor hopping amplitudes (connecting uppercase letters), while the blue lines represent negative next-nearest-neighbor hopping amplitudes (connecting lowercase letters). Nearest-neighbor hopping amplitudes are the same as in (a).

where $\hat{H}_{\text{eff}}$ is the full time-independent effective Hamiltonian that describes slow motion and $\hat{K}(t)$ is the time-periodic kick operator that describes micromotion [82, 83].

For the moment, in this subsection we first consider the noninteracting model $U = 0$. We also assume that the driving frequency ω is the highest energy scale, but that it is still low enough that the lowest-band approximation used in deriving Eq. (3.1) is still valid. In the leading order of the high-frequency expansion, the effective Hamiltonian $\hat{H}_{\text{eff}}$ is given by

$$\hat{H}_{\text{eff},0} = J'_x \sum_{l,m} \left[\mathrm{e}^{i\left((m-l-1)\pi/2-\pi/4\right)} \hat{a}^\dagger_{l+1,m} \hat{a}_{l,m} + \text{H.c.} \right] - J'_y \sum_{l,m} \left(\hat{a}^\dagger_{l,m+1} \hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1} \hat{a}_{l,m} \right), \quad (3.6)$$

where the renormalized hopping amplitudes are $J'_x = \frac{J_x \kappa}{\sqrt{2}\omega} = J_y$ and $J'_y = J_y\left(1 - \frac{1}{2}\frac{\kappa^2}{\omega^2}\right)$. A schematic representation of this model is presented in Fig. 3.1(a). The unit cell is shaded and the full lattice is spanned by the vectors $\mathbf{R}_1 = (4, 0)$ and $\mathbf{R}_2 = (1, 1)$. Particle hopping around a plaquette in the counterclockwise direction acquires a complex phase $-\frac{\pi}{2}$ and the model is equivalent to the Harper-Hofstadter Hamiltonian [68] for the case $\alpha = 1/4$ [68]. The explicit form of the kick operator $\hat{K}(t)$ from Eq. (3.3) is given in Appendix C.

Following Refs. [82, 83], we find that additional corrections of the order J_x^2/ω contribute to the system's dynamics and we introduce another approximation for the effective Hamiltonian

$$\hat{H}_{\text{eff},1} = \hat{H}_{\text{eff},0} + \frac{J_x^2}{\omega} \sum_{l,m} (-1)^l \left(2\hat{a}^\dagger_{l,m} \hat{a}_{l,m} + \hat{a}^\dagger_{l+2,m} \hat{a}_{l,m} + \hat{a}^\dagger_{l-2,m} \hat{a}_{l,m} \right). \quad (3.7)$$

The derivation of Hamiltonian (3.7) is given in Appendix C and its schematic representation is given in Fig. 3.1(b). The J_x^2/ω correction introduces next-nearest-neighbor hopping along x direction with opposite signs for lattice sites with either even or odd x-coordinate l. This term

does not change the total complex phase per plaquette, but the unit cell is now doubled and thus the first Brillouin zone is halved. A similar term was engineered on purpose in order to implement the Haldane model [79].

In the next subsection we investigate properties of energy bands of both effective Hamiltonians, $\hat{H}_{\text{eff},0}$ and $\hat{H}_{\text{eff},1}$. We use the units where $\hbar = 1$ and $a = 1$. Unless otherwise stated, we set the parameters to the following values: lattice size 100×100 sites, hopping amplitudes $J'_x = J_y = 1 \equiv J$, and the driving amplitude $\kappa = 0.58\ \omega$. This value of the driving amplitude was chosen to be the same as in the experiment [70]. In order to set the renormalized hopping amplitude along $\mathbf{e}_x$ to $J'_x = 1$, the initial hopping amplitude has to be $J_x = \sqrt{2}\omega/\kappa = 2.44$, and the correction term is therefore proportional to $J_x^2/\omega = 5.95/\omega$, so it cannot be safely neglected unless the driving frequency is very high.

3.1.2 Band structure

Momentum-space representations of the effective Hamiltonians $\hat{H}_{\text{eff},0}$ and $\hat{H}_{\text{eff},1}$, denoted by $\hat{\mathcal{H}}_{\text{eff},0}(\mathbf{k})$ and $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$, respectively, are derived in Appendix C. Band structures for the effective Hamiltonian $\hat{\mathcal{H}}_{\text{eff},0}$ without the J_x^2/ω correction, Eq. (C.20), as well as for the effective Hamiltonian $\hat{\mathcal{H}}_{\text{eff},1}$ including the correction term, Eq. (C.21), are shown in Fig. 3.2 for the two values of driving frequencies $\omega = 20$ and $\omega = 10$.

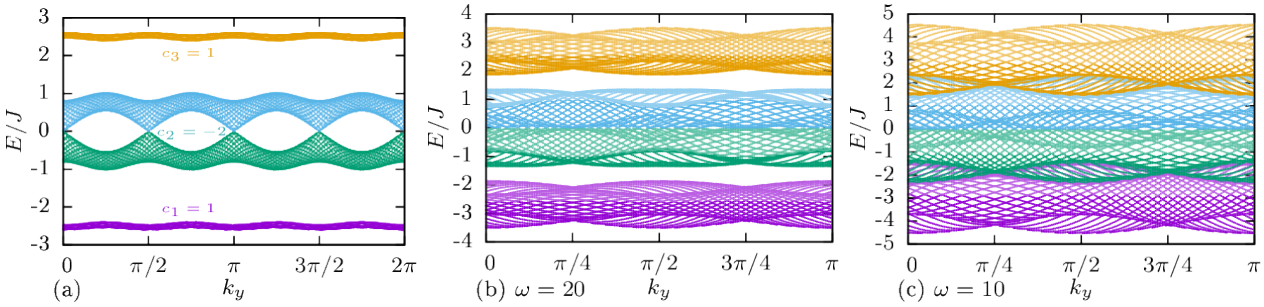

Figure 3.2: Energy bands of the effective Hamiltonians. (a) $\hat{\mathcal{H}}_{\text{eff},0}(\mathbf{k})$ Eq. (C.20), which is without the J_x^2/ω correction term. (b) $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$ Eq. (C.21), which includes the correction term. Driving frequency $\omega = 20$; gaps are open. (c) Same as (b), but with $\omega = 10$. Gaps are closed.

The Hamiltonian $\hat{H}_{\text{eff},0}$ is the Harper-Hofstadter Hamiltonian (1.4) for the flux $\alpha = 1/4$. It has four energy bands, where the middle two bands touch at $E = 0$ and can therefore be regarded as a single band; see Fig. 3.2(a). The topological content of these bands is characterized by the topological index called the Chern number (1.5). The Chern numbers of the three well-separated bands are $c_1 = 1$, $c_2 = -2$, and $c_3 = 1$.

Because the correction from Eq. (3.7) includes next-nearest-neighbor hopping terms, the elementary cell in real space is doubled [see Fig. 3.1(b)] and, as a consequence, the first Brillouin zone for the Hamiltonian $\hat{\mathcal{H}}_{\text{eff},1}$ is reduced by a factor of 2 compared to $\hat{\mathcal{H}}_{\text{eff},0}$. There are now eight lattice sites in the unit cell and eight energy bands, but the number of gaps depends on the

driving frequency. The new bands touch in pairs, in such a way that there are always maximally three well-separated bands. When the driving frequency is high enough, the correction is small and the gaps between the three bands remain open; see Fig. 3.2(b). The original band structure of $\hat{\mathcal{H}}_{\mathrm{eff},0}$ is recovered in the limit $\omega \to \infty$. The Berry curvature and the Chern number can be calculated using the efficient method presented in Ref. [185]. Our calculations confirm that the Chern numbers of $\hat{\mathcal{H}}_{\mathrm{eff},1}$ are equal to those of $\hat{\mathcal{H}}_{\mathrm{eff},0}$ ($c_1 = 1$, $c_2 = -2$, and $c_3 = 1$), as long as the gaps between the energy bands are open. The gaps close when the driving frequency is too low, see Fig. 3.2(c), and the Chern numbers of the subbands can no longer be properly defined.

3.1.3 Dynamics of incoherent bosons

We need to take into account a contribution of weak, repulsive interactions. Full numerical simulations of an interacting many-body problem are computationally demanding, so we need a reasonable, numerically tractable approximation. To this end we will use the classical field method [186], which belongs to a broader class of truncated Wigner approaches [187]. This method is similar to the approach used to treat incoherent light in instantaneous media [188, 189], known in optics as the modal theory.

The underlying idea of the method is to represent the initial state as an incoherent mixture of coherent states $|\psi\rangle$, $\hat{a}_{l,m}|\psi\rangle = \psi_{l,m}|\psi\rangle$ [186]. This is explained in more detail in Appendix D. In our study, we sample initial configurations of these coherent states with

$$|\psi(t=0)\rangle = \sum_{k=1}^{N_m} \mathrm{e}^{i\phi_k}|k\rangle, \tag{3.8}$$

where $\phi_k \in [0, 2\pi)$ are random phases and the states $|k\rangle$ correspond closely to the lowest-band eigenstates of $\hat{H}_{\mathrm{eff}}$. Each of N_{samples} initial states is time evolved and physical variables can be extracted by averaging over an ensemble of different initial conditions.

The time evolution of each of these coherent states is governed by

$$i\frac{d\psi_{l,m}(t)}{dt} = \sum_{ij} H_{lm,ij}(t)\psi_{i,j}(t) - F\,m\,\psi_{l,m}(t) + U|\psi_{l,m}(t)|^2\psi_{l,m}(t), \tag{3.9}$$

where $H_{lm,ij}(t) = \langle l, m|\hat{H}(t)|i, j\rangle$ are matrix elements of $\hat{H}(t)$ from Eq. (3.3), F is the external force, and interactions U contribute with the last, nonlinear term. Formally, Eq. (3.9) takes the form of the Gross-Pitaevskii equation [1, 190, 191]. The performances and limitations of the method are discussed and reviewed in Ref. [192].

For comparison, we also consider the related time evolution governed by the effective Hamiltonian

$$i\frac{d\psi_{l,m}(t)}{dt} = \sum_{ij} h^{\mathrm{eff}}_{lm,ij}\psi_{i,j}(t) - F\,m\,\psi_{l,m}(t) + U|\psi_{l,m}(t)|^2\psi_{l,m}(t), \tag{3.10}$$

where $h^{\text{eff}}_{lm,ij} = \langle l, m|\hat{h}^{\text{eff}}|i, j\rangle$, with $\hat{h}^{\text{eff}}$ being either $\hat{H}_{\text{eff},0}$ from Eq. (3.6), or $\hat{H}_{\text{eff},1}$ from Eq. (3.7). Equation (3.10) should be considered only as a tentative description of the system: the mapping between $\hat{H}(t)$ and $\hat{H}_{\text{eff}}$ is strictly valid only in the noninteracting regime and the interaction term may introduce complex, nonlocal, higher-order corrections [94]. However, we expect their contribution to be small in the limit $U \to 0$, and for time scales which are not too long [97, 98, 193, 194].

In the following we use $N_m = 300$ modes and accommodate $N_p = 300$ particles per mode, so in total in the simulations we have $N = N_m N_p = 90,000$ bosons. Typical densities in real space are up to 100 particles per site and we choose the values of U in the range $U \in [0, 0.05]$. Other parameters: $J'_x = J_y = 1$, $\kappa/\omega = 0.58$, $\omega = 10, 20$, and $F = 0.25J/a$. The correction terms are non-negligible in this frequency range. In practice, we first numerically diagonalize the Hamiltonian (D.2) from Appendix D and set our parameters in such a way that the lowest N_m modes have high overlap with the lowest band of the effective model. In the next step, we sample initial configurations (3.8). For each of $N_{\text{samples}} = 1,000$ sets of initial conditions we then time evolve Eq. (3.9) and extract quantities of interest by averaging over resulting trajectories. This value of N_{samples} is chosen to be high enough, so that the fluctuations are weak. We present and discuss results of our numerical simulations in the following sections.

3.2 Noninteracting case

We start by addressing the dynamics of noninteracting bosons. In this case we set $U = 0$ in Eq. (3.9) and numerically solve the single-particle Schrödinger equation without further approximations. Our aim is to numerically validate and compare the two approximations, Eqs. (3.6) and (3.7), for the effective Hamiltonian. To this purpose, we juxtapose results of the two approximative schemes with the numerically exact results obtained by considering the full time evolution governed by $\hat{H}(t)$. For clarity, the four different time evolutions that we consider in this section are summarized in Table 3.1. We calculate the center-of-mass position $x(t)$ and plot the results in Fig. 3.3. In this way we also find the regime of microscopic parameters where the Chern-number measurement can be optimally performed.

First, we consider the basic Harper-Hofstadter Hamiltonian (3.6) and select the occupied modes $|k\rangle$ of the initial state (D.1) as eigenstates of the model from Eq. (3.8) for $\hat{h}_{\text{eff}} = \hat{H}_{\text{eff},0}$. As explained in the previous section, at the initial moment $t_0 = 0$, the confinement is turned off and the force $\mathbf{F} = -F\mathbf{e}_y$ is turned on. As a consequence of the applied external force and the nonzero Chern number of the lowest band of the model (3.6), the particles exhibit an anomalous velocity in the direction perpendicular to the force [71]. In the ideal case, when the lowest band is fully populated, the theoretical prediction for the center-of-mass position in the

Table 3.1: Four different cases: the same effective Hamiltonian is always used for the initial state and band definitions, either with or without the correction. The evolution is governed either by the time-dependent Hamiltonian or by the same effective Hamiltonian as the one that was used for the initial state and calculation of band populations.

case	initial state	band populations	evolution
1	$\hat{H}_{\text{eff},1}$	$\hat{H}_{\text{eff},1}$	$\hat{H}_{\text{eff},1}$
2	$\hat{H}_{\text{eff},1}$	$\hat{H}_{\text{eff},1}$	$\hat{H}(t)$
3	$\hat{H}_{\text{eff},0}$	$\hat{H}_{\text{eff},0}$	$\hat{H}_{\text{eff},0}$
4	$\hat{H}_{\text{eff},0}$	$\hat{H}_{\text{eff},0}$	$\hat{H}(t)$

$\mathbf{e}_x$ direction is [70]

$$x(t) = x(t_0) + c_1 \frac{2Fa^2}{\pi\hbar} t, \tag{3.11}$$

where $c_1 = 1$ is the Chern number (1.5) of the lowest band. However, even in the ideal case, due to the sudden quench of the linear potential, a fraction of particles is transferred to the higher bands. To take this effect into account, the authors of Ref. [70] introduced a filling factor $\gamma(t)$

$$\gamma(t) = \eta_1(t) - \eta_2(t) + \eta_3(t), \tag{3.12}$$

where $\eta_i(t)$ are populations of different bands of Hamiltionian (3.6) from Eq. (D.4) in Appendix D and the plus and minus signs in Eq. (3.12) are defined according to the Chern numbers $c_1 = 1$, $c_2 = -2$, and $c_3 = 1$. The final theoretical prediction is then [70]

$$x(t) = x(t_0) + c_1 \frac{2Fa^2}{\pi\hbar} \int_0^t \gamma(t')dt'. \tag{3.13}$$

In Fig. 3.3(a) we consider the anomalous drift for a high value of the driving frequency $\omega = 20$, where we expect the expansion in $1/\omega$ to be reliable. We find an excellent agreement between the prediction (3.13) (dotted black line) and numerical calculation based on $\hat{H}_{\text{eff},0}$ (solid green line). However, some deviations between the full numerical results (dashed purple line) and the results of the approximation scheme (solid green line) are clearly visible. These deviations are even more pronounced for $\omega = 10$, Fig. 3.3(b).

Now we turn to the effective model (3.7). In this case we select the modes of the initial state as eigenstates of Eq. (3.8) for $\hat{h}_{\text{eff}} = \hat{H}_{\text{eff},1}$. Moreover, we also consider band populations (D.4) of the same model. In the case when $\omega = 20$, Fig. 3.3(c), the anomalous drift obtained using the effective Hamiltonian (3.7) (solid green line) closely follows the theoretical prediction (3.13). Moreover, from the same figure we can see that the effective Hamiltonian $\hat{H}_{\text{eff},1}$ reproduces the behavior of the time-dependent Hamiltonian very well. All three curves almost overlap for intermediate times ($5-40$ ms); see Fig. 3.3(c). We attribute the long-time (> 45 ms) deviations to the finite-size effects introduced by the next-nearest-neighbor hopping terms, which cause

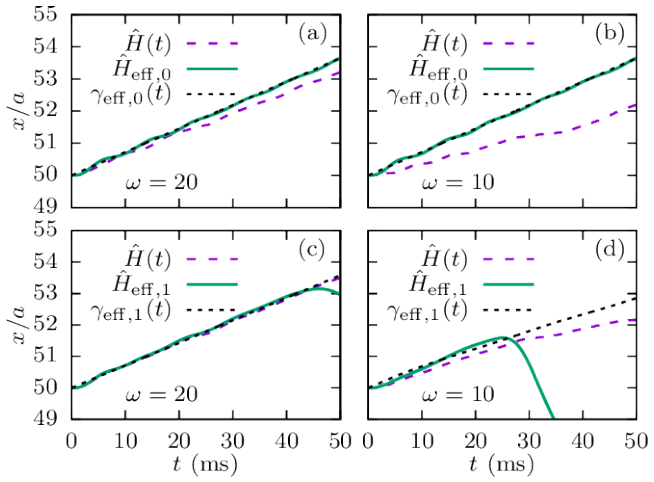

Figure 3.3: Anomalous drift $x(t)$. Dashed purple lines: numerical simulations using the time-dependent Hamiltonian $\hat{H}(t)$ (cases 2 and 4 from Table 3.1). Solid green lines: effective Hamiltonians $\hat{H}_{\mathrm{eff},1}$ (c) and (d) and $\hat{H}_{\mathrm{eff},0}$ (a) and (b) (cases 1 and 3). Dotted black lines: theoretical prediction (3.13) from $\gamma_{\mathrm{eff},1}(t)$ or $\gamma_{\mathrm{eff},0}(t)$. (a) Initial states and band populations obtained using the effective Hamiltonian $\hat{H}_{\mathrm{eff},0}$ without the correction (cases 3 and 4). Driving frequency $\omega = 20$. (b) $\omega = 10$. (c) Hamiltonian $\hat{H}_{\mathrm{eff},1}$ with the J_x^2/ω correction (cases 1 and 2). Driving frequency $\omega = 20$. (d) $\omega = 10$.

the atomic cloud to reach the edge of the lattice faster. This effect is explained in more detail in Section 3.3.2.

For a lower driving frequency $\omega = 10$, the effective and the time-dependent Hamiltonians do not agree so well anymore; see Fig. 3.3(d). The finite-size effects can be observed even earlier in this case (around 25 ms). This happens because the next-nearest-hopping terms are inversely proportional to the driving frequency. It is interesting to note that the prediction (3.13) is close to numerical data for short times even in this case when the gaps of the effective model are closed, see Fig. 3.2(c), and the Chern number of the lowest band is not well defined. In fact, it is surprising that the anomalous drift even exists in this case, as all subbands are now merged into a single band. We attribute this effect to our choice of the initial state. When the gaps are closed, it is hard to set the parameters in such a way that the lowest band is completely filled. The top of this band usually remains empty, and the particles thus do not "see" that the gap is closed.

Time evolution of the filling factor $\gamma(t)$ is plotted in Fig. 3.4 for four different cases from Table 3.1 – evolution using the effective Hamiltonian without correction $\hat{H}_{\mathrm{eff},0}$ [$\gamma_{\mathrm{eff},0}(t)$, case 3, dashed green line in Fig. 3.4(a)], the effective Hamiltonian with correction $\hat{H}_{\mathrm{eff},1}$ [$\gamma_{\mathrm{eff},1}(t)$, case 1, dashed green line in Fig. 3.4(b)], or the time-dependent Hamiltonian $\hat{H}(t)$ [$\gamma(t)$, cases 2 and 4, solid purple lines]. At the initial moment $\gamma(t_0 = 0) < 1$, because the initial state was multiplied by the operator $\mathrm{e}^{-i\hat{K}(0)}$. This introduces a shift between $\gamma(t)$ and $\gamma_{\mathrm{eff},1}(t)$. Apart from the shift, these two curves behave similarly, unlike the $\gamma_{\mathrm{eff},0}(t)$ curve that exhibits completely different behavior. Because of this, we use only $\gamma_{\mathrm{eff},1}(t)$ to estimate the value of the prediction (3.13).

We find that the values of $\gamma_{\mathrm{eff},1}(t)$ for $\omega = 20$ are high: ≥ 0.95; see Fig. 3.4. For this

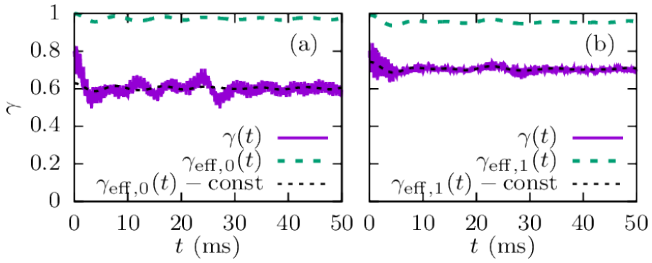

Figure 3.4: Time evolution of the filling factor $\gamma(t)$ for driving frequency $\omega = 20$. Solid purple lines: evolution governed by the time-dependent Hamiltonian $\hat{H}(t)$ (cases 2 and 4 from Table 3.1). Dashed green lines: evolution governed by the effective Hamiltonian $\hat{H}_{\text{eff},1}$ or $\hat{H}_{\text{eff},0}$ (cases 1 and 3). Dotted black lines: green lines shifted in order to compare them with purple lines. Shift is chosen so that the two lines approximately overlap. (a) Initial states and band populations obtained using the effective Hamiltonian $\hat{H}_{\text{eff},0}$, which is without the J_x^2/ω correction term (cases 3 and 4). (b) Hamiltonian $\hat{H}_{\text{eff},1}$ which is with the correction term (cases 1 and 2).

reason, up to 50 ms the center-of-mass position $x(t)$ exhibits roughly linear behavior with some additional oscillations. Interestingly, the anomalous drift $x(t)$ exhibits quadratic behavior on short time scales in all cases from Fig. 3.3. In Appendix E, we explain this feature using the time-dependent perturbation theory and Fermi's golden rule.

3.3 Interacting case

We now investigate the effects of weak repulsive interactions. We work in the high-frequency regime and set $\omega = 20$. As shown in Section 3.1.2, for $U = 0$ the effective Hamiltonian with correction, $\hat{H}_{\text{eff},1}$, is in this case equivalent to the Harper-Hofstadter Hamiltonian with flux $\alpha = 1/4$. Moreover, the same approximative form of the full effective model accurately reproduces the behavior of the time-dependent Hamiltonian up to 50 ms and thus provides a good starting point for the study of weakly interacting particles. We first consider the anomalous drift of the center of mass of the atomic cloud and then we inspect the expansion dynamics more closely in terms of atomic density distributions in real and momentum space.

3.3.1 Anomalous drift and dynamics of band populations

To simulate the dynamics of many incoherent bosons, we use the classical field method presented in Section 3.1.3 and propagate Eq. (3.9) in time. We assume that at $t_0 = 0$ atoms are uniformly distributed over the lowest band of $\hat{H}_{\text{eff},1}$. For this reason, the initial state is the same as the one that we use in the noninteracting regime. In this way, the dynamics is initiated by an effective triple quench: at $t_0 = 0$ the confining potential is turned off, atoms are exposed to the force $\mathbf{F} = -F\mathbf{e}_y$, and also the interactions between particles are introduced. The total number of particles is set to $N = 90,000$, which amounts to approximately 100 particles per lattice site in the central region of the atomic cloud. We consider only weak repulsion $U \leq 0.05$.

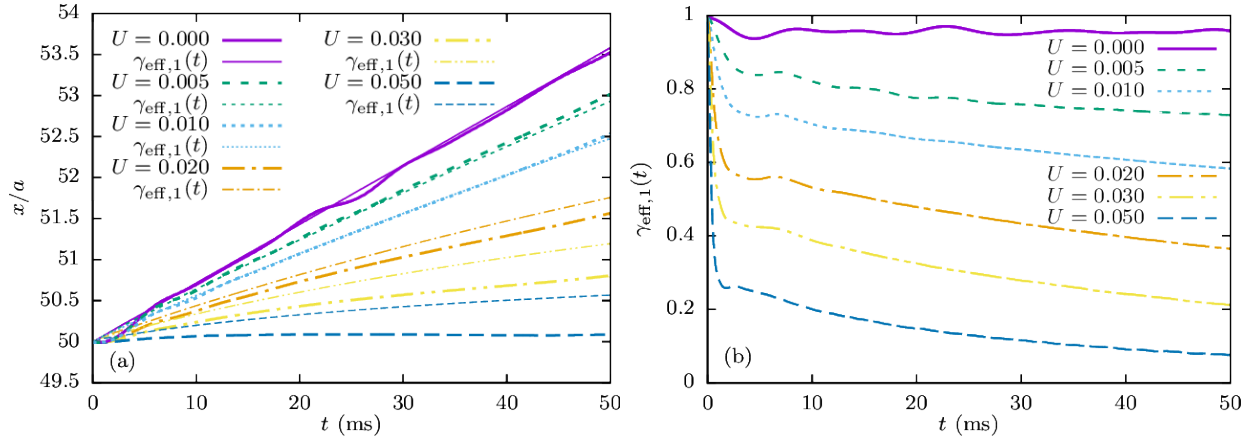

Figure 3.5: Effects of interactions. (a) Anomalous drift $x(t)$ for several different values of the interaction coefficient U. U is given in units where $J = 1$. Thick lines: numerical simulations using the time-dependent Hamiltonian $\hat{H}(t)$. Thin lines: theoretical prediction (3.13) from $\gamma_{\text{eff},1}(t)$. (b) Corresponding $\gamma_{\text{eff},1}(t) = \eta_1(t) - \eta_2(t) + \eta_3(t)$, obtained from simulations using the effective Hamiltonian $\hat{H}_{\text{eff},1}$.

The anomalous drift $x(t)$ obtained using the full time-dependent Hamiltonian is shown in Fig. 3.5(a) for several different values of the interaction strength U. In comparison to the noninteracting regime, we find that the weak repulsive interactions inhibit the response of the center of mass to the external force. In particular, at $t = 50\,\text{ms}$ the drift is reduced by about 15% for $U = 0.005$ and it is further lowered by an increase in U. Finally, at $U = 0.05$, the anomalous drift is barely discernible. Interestingly, for weak $U \in (0.001, 0.01)$ we find that the drift $x(t)$ in the range of $t \in (10, 50)\,\text{ms}$ looks "more linear" as a function of time in comparison to the noninteracting result.

We now analyze the anomalous drift in terms of the filling factor $\gamma(t)$ and compare the results of Eq. (3.9) with the description based on Eq. (3.10). By solving Eq. (3.10) we obtain the filling factor $\gamma_{\text{eff},1}(t)$ following Eq. (D.4) and present our results in Fig. 3.5(b). Whenever the results of Eq. (3.9) reasonably agree with the results obtained from Eq. (3.10), we are close to a steady-state regime with only small fluctuations in the total energy, as Eq. (3.10) preserves the total energy of the system. In this regime, during the expansion dynamics the interaction energy is converted into the kinetic energy and atoms are transferred to higher bands of the effective model. Consequently, the filling factor $\gamma_{\text{eff},1}(t)$ is reduced. Typically, we find three different stages in the decrease of $\gamma_{\text{eff},1}(t)$.

In an early stage, $t \leq t_1 = 5\,\text{ms}$, a fast redistribution of particles over the bands of the effective model sets in due to the sudden quench of U. The factor $\gamma_{\text{eff},1}(t)$ decays quadratically as a function of time down to $\gamma_{\text{eff},1}(t_1) \approx 0.75$ for $U = 0.01$, and $\gamma_{\text{eff},1}(t_1) \approx 0.25$ for $U = 0.05$. In this process the interaction energy of the system is quickly lowered as described in Appendix F. At later times $t > 5\,\text{ms}$, we observe a linear decay of the filling factor $\gamma_{\text{eff},1}(t)$ as a function of time, that finally turns into an exponential decay at even later times ($t > 10$ ms). Similar regimes are observed in other dynamical systems. For example, a decay rate of an initial state

suddenly coupled to a bath of additional degrees of freedom exhibits these three stages [195]. The initial quadratic decay is often denoted as "the Zeno regime." For longer propagation times, Fermi's golden rule predicts the linear decay. At even longer time scales, when the repopulation of the initial state is taken into account, the time-dependent perturbation theory yields the exponential regime, known under the name of the Wigner-Weisskopf theory [195].

We now investigate this last regime in more detail. For the population of the lowest band $\eta_1(t)$, an exponential decay function $f(t) = a + b\mathrm{e}^{-ct}$ provides high quality fits for $t \in (10, 50)$ ms; see Fig. 3.6(a) for an example. Similarly, the populations of two higher bands can also be fitted to exponential functions. The obtained exponential decay coefficients c for the lowest band population are plotted as a function of the interaction strength U in Fig. 3.6(b). The resulting dependence is approximately quadratic: $c(U) = \alpha_0 + \alpha_1\ U + \alpha_2\ U^2$. For small values of U, the exponents $c(U)$ obtained for the dynamics governed by $\hat{H}(t)$ and $\hat{H}_{\mathrm{eff},\,1}$ agree very well and exhibit linear behavior. At stronger interaction strengths $U \geq 0.03$, the approximation of Eq. (3.10) becomes less accurate as it omits the quadratic contribution in $c(U)$ found in the full time evolution. In addition, the values of the exponents c are affected by the force strength F and driving frequency ω.

As we now understand some basic features of $\gamma_{\mathrm{eff},1}(t)$, we make an explicit comparison between the numerical results for the anomalous drift and the expectation (3.13). The dashed lines in Fig. 3.5(a) correspond to the theoretical prediction (3.13) calculated from $\gamma_{\mathrm{eff},1}(t)$. For the intermediate interaction strengths $U \leq 0.01$, we find a very good agreement between the two. From this we conclude that the interaction-induced transitions of atoms to higher bands are the main cause of the reduced anomalous drift $x(t)$ as a function of U. When the interactions become strong enough ($U \sim 0.02$), the numerical results start to deviate from the theoretical prediction (3.13) with $\gamma_{\mathrm{eff},1}(t)$. In this regime, Eq. (3.10) does not provide a reliable description of the dynamics, as higher-order corrections need to be taken into account.

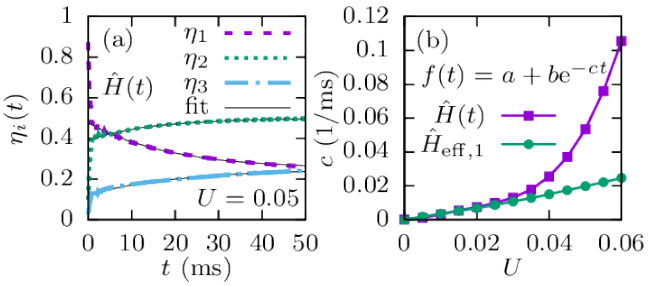

Figure 3.6: (a) Evolution of the band populations $\eta_i(t)$. Dashed lines: numerical results obtained using the time-dependent Hamiltonian $\hat{H}(t)$. Solid black lines: exponential fit using $f(t) = a + b\mathrm{e}^{-ct}$. The coefficient a was fixed to $a_1 = 0.25$, $a_2 = 0.50$ and $a_3 = 0.25$ for the first, second and third band respectively. (b) Dependence of the exponential decay coefficients for the lowest band population $\eta_1(t)$ on the interaction strength. U is given in units where $J = 1$.

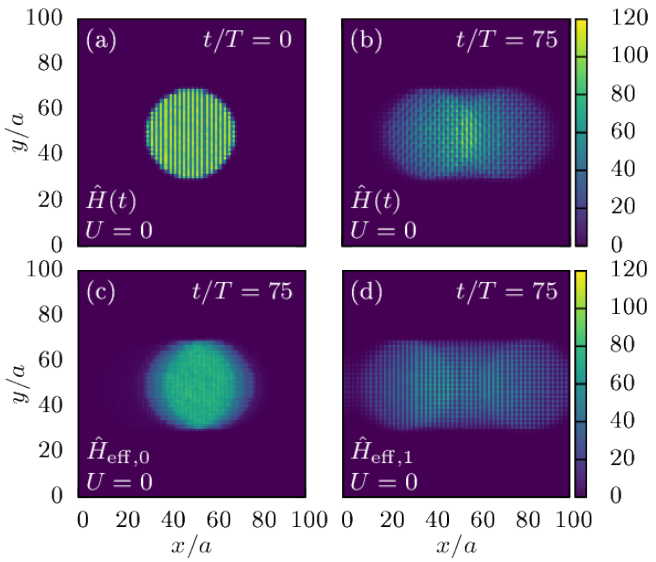

Figure 3.7: Real-space density distribution, noninteracting case $U = 0$. (a) Initial state. (b) After 50 ms (75 driving periods), evolution using the time-dependent Hamiltonian $\hat{H}(t)$. (c) Evolution using effective Hamiltonian without correction $\hat{H}_{\mathrm{eff},0}$. (d) Evolution using effective Hamiltonian with correction $\hat{H}_{\mathrm{eff},1}$.

3.3.2 Real and momentum-space dynamics

So far we have considered the averaged response of the whole atomic cloud. We now inspect the expansion dynamics in a spatially resolved manner. The real-space probability densities at the initial moment and after 50 ms (75 driving periods) are shown in Figs. 3.7 and 3.8, and the corresponding momentum-space probability densities in Appendix F.

At the initial moment, the atomic cloud is localized in the center of the lattice. By setting $r_0 = 20$ in the confining potential of Eq. (D.2) and populating the lowest-lying states, we fix the cloud radius to $r = 20$, Fig. 3.7(a). The cloud density is of the order of 100 atoms per lattice site and a weak density modulation is visible along x direction. After the confining potential is turned off, and the external force in the $-\mathbf{e}_y$ direction is turned on, the cloud starts to expand and move in the $+\mathbf{e}_x$ direction. As shown in the previous subsection, the band populations and therefore the anomalous drift are significantly altered by the interaction strength, and this is also the case with the expansion dynamics; see Figs. 3.7 and 3.8.

In the noninteracting case, Fig. 3.7(b), the atomic cloud nearly separates into two parts moving in opposite directions along x axes (while the center of mass still moves in the $+\mathbf{e}_x$ direction). By comparing Fig. 3.7(c) and Fig. 3.7(d), we conclude that this effect stems from the next-nearest-neighbor hopping along x present in the effective Hamiltonian (3.7), as it does not happen in the effective model without the correction term (3.6). This type of separation was already observed in Ref. [168], where the next-nearest-neighbor hopping terms were also present.

When the interactions between particles are included, this separation is not so prominent [Fig. 3.8(a), $U = 0.01$], and it almost completely disappears when the interactions are strong

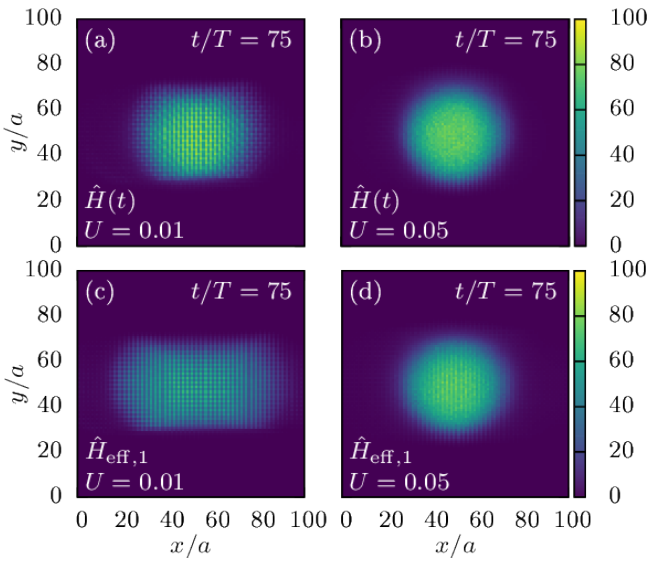

Figure 3.8: Real-space density distribution after 50 ms (75 driving periods), interacting case. U is given in units where $J = 1$. (a) Evolution using the time-dependent Hamiltonian $\hat{H}(t)$, $U = 0.01$. (b) Same with $U = 0.05$. (c) Evolution using the effective Hamiltonian $\hat{H}_{\text{eff},1}$, $U = 0.01$. (d) Same with $U = 0.05$.

enough [Fig. 3.8(b), $U = 0.05$]. This is also the case when the evolution is governed by the effective Hamiltonian $\hat{H}_{\text{eff},1}$; see Figs. 3.8(c) and 3.8(d). Atomic cloud widths $d_x = \sqrt{\langle x^2\rangle - \langle x\rangle^2}$ during the expansion are plotted in Fig. 3.9. We observe a slow expansion of the cloud in y direction, Fig. 3.9(b), and much faster expansion along x direction, Fig. 3.9(a), which comes about as a consequence of the cloud separation. On top of this, we observe that the interactions enhance expansion along y. Surprisingly, the opposite is true for the dynamics along x. This counterintuitive effect is often labeled as self-trapping and its basic realization is known for the double-well potential [196, 197]. In brief, strong repulsive interactions can preserve the density imbalance between the two wells, as the system can not release an excess of the interaction energy. In our case, the situation is slightly more involved as the cloud splitting is inherent (induced by the corrections of the ideal effective Hamiltonian). Apart from this, due to the driving the total energy is not conserved. However, our numerical results indicate that the interaction energy is slowly released in the second expansion stage, Fig. F.1. Effectively, in this

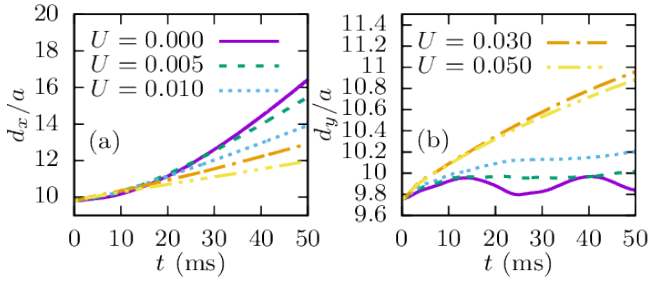

Figure 3.9: Atomic cloud width for different interaction strengths, evolution using the time-dependent Hamiltonian $\hat{H}(t)$. U is given in units where $J = 1$. (a) $d_x = \sqrt{\langle x^2\rangle - \langle x\rangle^2}$. (b) $d_y = \sqrt{\langle y^2\rangle - \langle y\rangle^2}$.

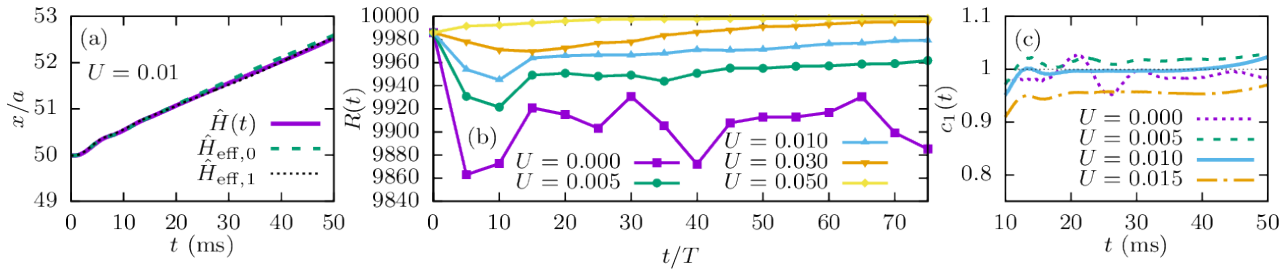

Figure 3.10: (a) Comparison of anomalous drifts obtained from evolution using the time-dependent Hamiltonian $\hat{H}(t)$ (solid purple line), effective Hamiltonian without correction $\hat{H}_{\text{eff},0}$ (dashed green line) and effective Hamiltonian with correction $\hat{H}_{\text{eff},1}$ (dotted black line). Intermediate interaction strength $U = 0.01$. U is given in units where $J = 1$. (b) Time evolution of the inverse participation ratio in momentum space for several different values of U. Evolution is performed using the time-dependent Hamiltonian $\hat{H}(t)$. When the interactions are strong enough, IPR approaches the maximal possible value (10,000 in this case), which is equal to the total number of states and corresponds to the completely delocalized state. U is given in units where $J = 1$. (c) Chern number of the lowest band obtained for different interaction strengths as the ratio of the theoretical prediction for the anomalous drift and numerical results: $c_1(t) = \left(\frac{2Fa^2}{\pi\hbar}\int_0^t \gamma_{\text{eff},1}(t')dt'\right) / \left(x(t) - x(t_0)\right)$.

way the interactions cancel out the contribution of the next-nearest-neighbor hopping and favor the measurement of the properties of the model (3.6). In Fig. 3.10(a) we show that deviations between different approximations based on $\hat{H}(t)$, $\hat{H}_{\text{eff, 1}}$, and $\hat{H}_{\text{eff, 0}}$ in the anomalous drift $x(t)$ nearly vanish at $U = 0.01$.

Another desirable effect might be that the interactions make the momentum-space probability density more homogeneous, see Appendix F, so that the real-space probability density becomes more localized. We can quantify momentum-space homogeneity using the inverse participation ratio $R(t) = \frac{1}{\sum_i P_i^2(t)}$, where $P_i(t) = |\psi_i(t)|^2$ is the probability that the state ψ_i is occupied at time t. Minimal value of the inverse participation ratio (IPR) is 1 and it corresponds to a completely localized state, while the maximal value is equal to the total number of states (in our case 10,000) and corresponds to the completely delocalized state, where the particles have the same probability of being at any quasimomentum $\mathbf{k}$. As stated before, the first Brillouin zone of the lowest band has to be as homogeneously populated as possible in order to properly measure the lowest band Chern number. From Fig. 3.10(b), we see that IPR increases in time when the interaction coefficient U is large enough, so we can conclude that the interactions are actually beneficial for measuring the Chern number, as they can "smooth-out" the momentum-space probability density. In Fig. 3.10(c) we give estimates for the Chern number that can be extracted from our numerical data for different values of U. We find the best estimate $c_1 \sim 0.99$ for the intermediate interaction strength $U \sim 0.01$.

3.3.3 Staggered detuning

Here we briefly consider the effects of staggered detuning that was introduced in the experimental study [70] during the loading and band mapping sequences. This detuning can be described by an additional term

$$\frac{\delta}{2}\sum_{l,m}\left[(-1)^l+(-1)^m\right]\hat{n}_{l,m} \tag{3.14}$$

in the Hamiltonians $\hat{H}(t)$ and $\hat{H}_{\mathrm{eff},1}$. We will ignore the higher-order [at most $\mathcal{O}\left(\frac{1}{\omega^2}\right)$] corrections that this term introduces to the effective Hamiltonian. Staggered detuning does not break the symmetry of the effective Hamiltonian $\hat{H}_{\mathrm{eff},1}$, but if δ is large enough, it can cause a topological phase transition and make all bands topologically trivial. By numerically calculating the Berry curvature and Chern numbers c_i', we find that this transition occurs at $\delta_c \approx 1.38\,J$; see Fig. 3.11. This value is lower than the one for the ordinary Harper-Hofstadter Hamiltonian for $\alpha = 1/4$, which is $\delta_c = 2\,J$ [70], due to the different hopping amplitudes J_x' and J_y', and due to the additional J_x^2/ω correction that we consider.

We now investigate how this topological transition can be probed through the dynamical protocol used in the experiment. We again numerically calculate the anomalous drift and the evolution of the filling factor, but now with staggered detuning (3.14) included in the Hamiltonian $\hat{H}_{\mathrm{initial}}$ (D.2) used to obtain the initial state, in the equations of motion (3.9) and (3.10), and in the definitions of the band populations $\eta_i(t)$ (D.4). Using these results, we repeat the procedure for the extraction of the lowest band Chern number from numerical data that was carried out in the previous section. The Chern number obtained by comparing the anomalous drift to the prediction calculated from the filling factor is then averaged over the time interval $t \in (20, 40)\,\mathrm{ms}$. This interval was chosen in order to avoid the initial quadratic regime and the finite-size effects at later times. The resulting lowest band Chern numbers for several different values of detuning δ in both the noninteracting case and the case of intermediate interaction strength $U = 0.01$ are presented in Fig 3.11.

We can see that the calculated value of the Chern number decreases from $c_1 = 1$ to $c_1 = 0$ with increasing detuning δ. The obtained value of the Chern number is lower than 1 even before the phase transition occurs. This is due to our choice of the initial state, which is not perfectly homogeneous in momentum space. Close to the phase transition, both the energy bands and the Berry curvature have pronounced peaks at the same regions of the first Brillouin zone, and these regions are initially less populated. Because of this, the Berry curvature at these regions contributes less to the anomalous drift, which lowers the measured Chern number. This effect is somewhat reduced by the interactions, as they smooth out the momentum-space probability density, and might also cancel out the detuning term. Similar interplay of interactions and staggering was observed in the fermionic Hofstadter-Hubbard model [198]. The obtained results are in line with experimental measurements [70].

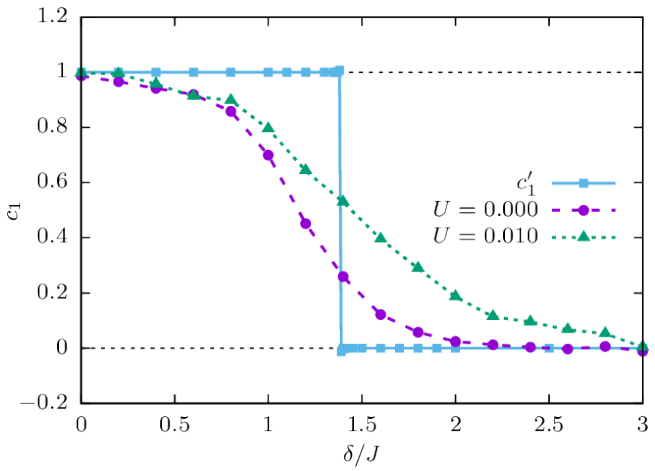

Figure 3.11: Lowest band Chern numbers extracted from numerical data for several different values of detuning δ. Purple circles: noninteracting case, $U = 0$. Green triangles: $U = 0.01$. Blue squares: Theoretical values of the lowest band Chern number c_1'. A topological phase transition is visible at $\delta_c \approx 1.38$. The lines between points are only a guide to the eye.

3.4 Conclusions

Motivated by the recent experimental results reporting the Chern numbers of topological bands in cold-atom setups, we studied numerically bosonic transport in a driven optical lattice. The considered driving scheme and the range of microscopic parameters were chosen to be close to those in a recent experimental study [70]. The driving frequency was set to be high enough in order to avoid strong energy absorption for the relevant time scales. Additionally, the system was restricted to a two-dimensional lattice, even though the actual experimental setup had continuous transverse degrees of freedom. This restriction stabilizes the system [172, 174, 184] and leads to lower heating rates than those in the experiment. It corresponds to the case of strongly confined third dimension.

We investigated bosonic dynamics for the full time-dependent Hamiltonian, the effective Floquet Hamiltonian, and included the effects of weak repulsive interactions between atoms using the mean-field approximation. In the noninteracting case, we found that the effective Hamiltonian and its band structure depend on the frequency of the drive ω through an additional J_x^2/ω correction term. The initial state was set as a mixture of incoherent bosons homogeneously populating the lowest band, but a possible direction of future research could be to simulate the full loading sequence of an initial Bose-Einstein condensate and to try to obtain the incoherent state through driving, as it was done in the experiment.

The main focus of this work is on the effects of weak interactions. For a weak atomic repulsion, atomic transitions to higher effective bands obtained in our simulations mainly occur due to a release of the initial interaction energy during the atomic-cloud expansion. Although the effect is undesirable, it can be properly taken into account in the extraction of the Chern number. At larger interaction strengths, the transitions are more pronounced as the system

absorbs energy from the drive. In this regime the good agreement between the full and effective description is lost and the measurement should become more complicated. In addition to causing redistribution of atoms over bands, our results show that weak interactions can also be beneficial in measuring the Chern number. Their desirable effect comes about due to smoothening the atomic distribution over the topological band and due to canceling out the contribution of some less relevant terms to the bosonic dynamics.

Chapter 4

Bosonic fractional quantum Hall states in driven optical lattices

Since early experiments with quantum gases, there has been a strong interest in the realization of fractional quantum Hall (FQH) states in these setups [199–213]. Despite numerous experimental achievements and a variety of theoretical proposals, FQH physics has still not been reached in cold-atom experiments. At first glance, both key requirements for the emergence of FQH states - atomic interactions and strong synthetic magnetic fields - are now experimentally available. However, there are several specific details in the implementation of strong synthetic magnetic fields for cold atoms that make the realization of FQH states still challenging.

The most advanced recent realizations of artificial gauge potentials exploit periodically driven optical lattices [70, 75, 77–80, 86, 87, 173, 214]. However, general arguments and numerical studies [94, 95, 215] suggest that the interplay of interactions and driving in a thermodynamically large system introduces heating, leading to a featureless infinite-temperature state in the long-time limit. Although this general result might sound discouraging, the heating process can be very slow in some driven systems for specific regime of microscopic parameters. There, the system can be described by a physically interesting "prethermal" Floquet state on experimentally relevant time-scales [96–98, 193, 194, 216, 217]. Moreover, the onset of thermalization in a finite-size interacting system may exhibit unexpected features, not found in the thermodynamic limit [218, 219]. Heating rates and resulting instabilities have been recently investigated both theoretically and experimentally for the driven Bose-Hubbard model in the weakly interacting regime [96, 174, 178, 179]. Moreover, experimental studies of the driven Fermi-Hubbard model in a honeycomb lattice have established a timescale of the order of 100 tunneling times for the regime where the effective-model description applies [124, 220].

In this Chapter, we consider small systems of several interacting bosonic atoms in a periodically driven optical lattice featuring synthetic magnetic flux. The focus of our study is on finding optimal microscopic parameters that would allow to prepare and probe the basic

bosonic Laughlin state in this setup. To this end, we employ exact numerical simulations of the driven Bose-Hubbard model [169] for small system sizes.

From one point of view, it is expected that a small driven system exhibits low heating rates for a driving frequency set above a finite bandwidth of an effective model [94]. However, driving a system with such a high frequency may lead to undesirable effects, such as coupling of the lowest band to higher bands of the underlying optical lattice, thus making the initial description based on the lowest-band Hubbard model inapplicable. These effects have been addressed in a recent study [88] where an optimal intermediate frequency window for Floquet engineering has been established.

In our study, we go a step further in the search for the optimal regime that might allow for the bosonic Laughlin states under driving. In particular, for a realistic, intermediate value of a driving frequency, the interaction term complicates the effective model by introducing several higher-order terms. Their effect on the topological states has been addressed only recently [221, 222] and it has been found that typically these terms work against the topological state. For this reason, the stability of the Laughlin state at intermediate driving frequency requires a separate study, that we perform here. Moreover, we numerically investigate an experimentally relevant preparation protocol for the Laughlin state in a driven system [183]. For a reference, we note that a simpler but closely related question concerning the static (undriven systems) has gained lot of attention [181, 202, 203, 211].

This Chapter is organized as follows: in Section 4.1 we introduce the model under study and briefly review key features of the particle-entanglement spectra that we will exploit in the identification of the Laughlin-like state. Then, in Section 4.2.1 we investigate general heating effects of interacting bosons exposed to the driving. By extending this approach, in Section 4.2.2 we construct the stroboscopic time-evolution operator and inspect its eigenstates in order to identify possible FQH states. Finally, in Section 4.3 we address the possibility of accessing these states in an experiment through a slow ramp of the driving term.

4.1 Model and method

In this section we first introduce the driven model and explain the basis of Floquet engineering. Then we summarize several key features of the particle-entanglement spectra that we use to characterize the bosonic Laughlin states.

4.1.1 Driven model

Properties of bosonic atoms in a deep optical lattice can be realistically described within the framework of the Bose-Hubbard model given by Eq. (1.1) [5]. We consider a basic driving scheme [169] that introduces a uniform, synthetic magnetic flux into a *square* optical lattice

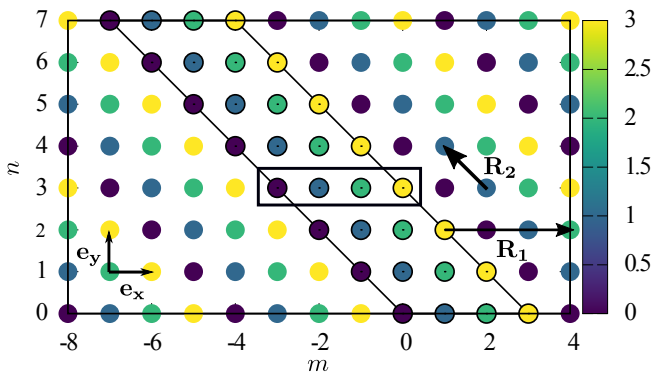

Figure 4.1: Lattice geometry used throughout this Chapter. The parallelogram gives the exemplary lattice size $(L_x, L_y) = (4, 8)$. The color scale is defined by $\mod(m + n, 4)$, in accordance with the driving term from Eq. (4.1). The vectors $\mathbf{R}_1 = 4\,\mathbf{e}_x, \mathbf{R}_2 = -\mathbf{e}_x + \mathbf{e}_y$ are used to implement periodic boundary conditions. The small rectangle gives the magnetic unit cell for the effective model in Eq. (4.2).

here spanned by the two vectors $\mathbf{e_x}$ and $\mathbf{e_y}$. The corresponding Hamiltonian is given by the driven Bose-Hubbard model

$$\begin{aligned}\hat{H}(t) &= -J_x \sum_{m,n} \left(\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \text{H. c.}\right) - J_y \sum_{m,n} \left(e^{i\omega t}\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} + \text{H. c.}\right) \\ &+ \frac{\kappa}{2}\sum_{m,n} \sin\left(\omega\, t - (m + n - 1/2)\,\phi\right)\hat{n}_{m,n} + \frac{U}{2}\sum_{m,n}\hat{n}_{m,n}(\hat{n}_{m,n} - 1), \end{aligned} \tag{4.1}$$

where operators $\hat{a}_{m,n}$ ($\hat{a}^\dagger_{m,n}$) annihilate (create) a boson at lattice position (m, n), and local density operators are $\hat{n}_{m,n} = \hat{a}^\dagger_{m,n}\hat{a}_{m,n}$. J_x and J_y are tunneling amplitudes and U is the on-site local repulsive interaction. We use the units where $\hbar = 1$ and the lattice constant $a = 1$. The driving scheme is defined by the driving frequency ω, the driving amplitude κ and by a phase ϕ. In the following we set $\phi = \pi/2$ and $\kappa/\omega = 0.5$. These values were recently used in an experimental realization of the Harper-Hofstadter model [70]. The derivation of this model is briefly reviewed in Appendix G. We assume periodic boundary conditions implemented using the vectors $\mathbf{R}_1 = 4\,\mathbf{e}_x, \mathbf{R}_2 = -\mathbf{e}_x + \mathbf{e}_y$, as presented in Fig. 4.1. This choice is compatible with the driving term and it allows us to exploit translational symmetry by working in the fixed quasimomentum basis.

Formally, by using the Floquet theory [81–83], it can be shown that the full time-evolution operator corresponding to this model is given by Eq. (1.9). The full-time evolution operator is periodic as well and consequently the (quasi)eigenenergies of the time-independent effective Hamiltonian $\hat{\mathcal{H}}_{\text{eff}}$ are defined up to modulo ω. Eq. (1.9) gives formal mapping of a periodically driven system to an effective model that captures the stroboscopic time evolution of the model. However, according to general analytical arguments and numerical insights, the corresponding effective model of a driven interacting many-body system in the thermodynamic limit exhibits nonphysical features [94, 95]. In particular, the system thermalizes and in the long-time limit its steady state is a featureless, infinite-temperature state, independent of the initial state.

Here we consider small samples of several bosonic atoms. Due to a finite spectrum bandwidth, we expect the high-frequency expansion to be relevant for a finite range of the driving frequency. Within these assumptions, the leading-order (in $1/\omega$) effective Hamiltonian is

$$\begin{aligned}\hat{H}_{\text{eff}} &= -J_x \sum_{m,n} \left(\hat{a}^\dagger_{m+1,n} \hat{a}_{m,n} + \text{H. c.} \right) - J'_y \sum_{m,n} \left(e^{i(m+n)\phi} \hat{a}^\dagger_{m,n+1} \hat{a}_{m,n} + \text{H. c.} \right) \\ &+ \frac{U}{2} \sum_{m,n} \hat{n}_{m,n} \left(\hat{n}_{m,n} - 1 \right). \end{aligned} \tag{4.2}$$

The Hamiltonian (4.2) features complex hopping phases $e^{i(m+n)\phi}$ that result in a uniform synthetic magnetic flux ϕ per lattice plaquette. Due to the driving, the renormalized hopping amplitude along the y direction turns into

$$J'_y \equiv \frac{\kappa}{2\omega} \sin(\phi/2)\, J_y. \tag{4.3}$$

For the values $\phi = 2\pi\alpha$, where the flux density α is set to $\alpha = 1/4$, and $\kappa/\omega = 0.5$, the tunneling amplitude along y direction in the effective model is $J'_y \approx J_y \times 0.1768$.

In a certain regime of microscopic parameters, the ground state of the model defined in Eq. (4.2) is given by the lattice version of the Laughlin state [203, 205, 223–225]. The Laughlin state is stabilized for the filling factor $\nu = N_p/N_\phi = 1/2$, where $N_\phi = \alpha L_x \times L_y$ is the total number of fluxes (N_ϕ being an integer) and N_p is the number of bosons, and for a strong-enough repulsion U. Another important requirement for the Laughlin state is to avoid the strong hopping anisotropy and to keep $J_x \approx J'_y$, so we set $J_x = 0.2 J_y$. We consider system sizes $N_p = 4, 5, 6$ and the respective lattices sizes $(L_x, L_y) = (4, 8), (4, 10)$, and $(4, 12)$, see Fig. 4.1, where we expect the ground state to correspond to the $\nu = 1/2$ Laughlin state. The Hilbert space sizes for $k_x = k_y = 0$ are $\dim \mathcal{H} = 6564, 108604$, and 1913364 respectively. For this choice of microscopic parameters, the model ground state of Eq. (4.2) is approximately twofold degenerate. The two ground-states are found in the sectors $k_x = 0, k_y = 0$ and $k_x = 0, k_y = \pi$. We denote them by $|\psi^{0,0}_{\text{LGH}}\rangle$ and $|\psi^{0,\pi}_{\text{LGH}}\rangle$.

As we are mainly interested in the driven regime, not only the ground state, but the full spectrum of the model from Eq. (4.2) plays a role. A rough argument is that the system does not absorb energy provided that the driving frequency ω is set above the bandwidth of the effective model. Several spectra of the model from Eq. (4.2) for $k_x = 0, k_y = 0$ are presented in Fig. 4.2(a). It can be seen that the ground-state energy is weakly affected by the value of $U \geq J_x$, while the top part of the spectrum with few states is found at $UN_p(N_p - 1)/2$. For higher values of U the spectrum splits into bands where the lowest band corresponds to the hard-core bosons and higher bands include double and higher occupancies.

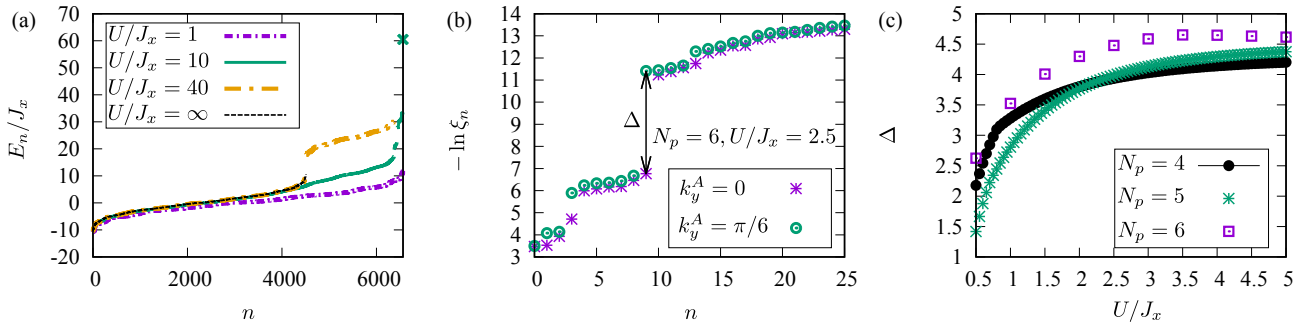

Figure 4.2: (a) The energy spectrum E_n of the model from Eq. (4.2) in the $k_x = 0, k_y = 0$ sector for $N_p = 4$ and different values of interaction $U/J_x = 1, 10, 40$ and $U/J_x = \infty$ (hard-core bosons). The top part of the spectrum is at $\approx (U/J_x)N_p\,(N_p - 1)\,/2$. (Not shown for $U/J_x = 40$.) For a high ratio U/J_x the spectrum splits into bands. The lowest band corresponds to hard-core bosons. (b) The low-lying part of the particle-entanglement spectrum $-\ln \xi_n$ of the ground-state incoherent superposition, Eq. (4.5), in the region A momentum sectors $k_y^A = 0$ and $k_y^A = \pi/6$, and for $N_p = 6, U/J_x = 2.5$. (c) The particle-entanglement gap Δ of the incoherent superposition Eq. (4.5) as a function of interaction strength U for $N_p = 4, 5, 6$.

4.1.2 Particle-entanglement spectra

There are several ways to characterize the ground states of the model from Eq. (4.2) as the Laughlin states. Usually, the starting point in this direction is the identification of the twofold degeneracy expected in the implemented torus geometry for $\nu = 1/2$. Another relevant quantity is the overlap of the numerically obtained state with the Laughlin analytical wave function in the torus geometry [205, 225]. More direct evidence can be obtained through the calculation of the relevant topological index (Chern number) or the quantized Hall conductance. An additional convincing approach, that we pursue here, is based on the analysis of the entanglement spectra of the relevant states.

In the following we will use the particle-entanglement spectrum (PES) [225, 226] to distinguish possible topologically nontrivial states. In order to obtain this type of entanglement spectrum, we partition N_p particles into two sets of N_A and $N_B = N_p - N_A$ particles. For a given mixed state ρ, we construct a reduced density matrix $\rho_A = \mathrm{tr}_B \rho$ by performing a partial trace over N_B particles. The resulting PES is given by $-\ln \xi_n$, where ξ_n are eigenvalues of ρ_A. The related particle-entanglement entropy is given by [227, 228]

$$S_A = -\mathrm{tr}\,(\rho_A \ln \rho_A)\,. \tag{4.4}$$

By partitioning particles, we keep the geometry of the system unchanged. For this reason, we will inspect the PES for the different momentum sectors k_y^A of the remaining N_A particles. An example of a PES is presented in Fig. 4.2(b). As proposed in Refs. [225, 226], we have considered the incoherent superposition of the almost twofold degenerate ground state of Eq. (4.2) as the

density matrix

$$\rho_{\mathrm{GS}} = \frac{1}{2}\left(|\psi^{0,0}_{\mathrm{LGH}}\rangle\langle\psi^{0,0}_{\mathrm{LGH}}| + |\psi^{0,\pi}_{\mathrm{LGH}}\rangle\langle\psi^{0,\pi}_{\mathrm{LGH}}|\right). \tag{4.5}$$

For simplicity, we only present the PES for the two momenta $k_y^A = 0$ and $k_y^A = \pi/6$. We observe a clear particle-entanglement gap Δ. In addition, the counting of low-lying modes below this gap (ten modes for $k_y^A = 0$ and nine modes for $k_y^A = \pi/6$, at $N_A = 3, N_p = 6$) corresponds to the Laughlin state [225, 226]. In this way the PES encodes topological features of the state ρ in the form of well defined number of excitations per momentum sector k_y^A [225, 226]. This type of analysis is useful as it can identify topological features even without model states, as done for the case of fractional Chern insulators [229, 230].

In the following we will consider specific particle partitions $N_A = 2, N_p = 4$; $N_A = 2, N_p = 5$; and $N_A = 3, N_p = 6$. For these cases the counting of excitations $\mathcal{N}_L(k_y^A)$ per momentum sector k_y^A is well established and given in Table 4.1. In Fig. 4.2(c) we show the particle-entanglement gap of the mixtures, Eq. (4.5), obtained at different values of U. Numerical results for the obtained PES indicate that a reasonably large gap is found starting at $U \sim 0.5J_x$ and the characteristic features of the Laughlin state persist with a further increase in U. We note that at lower values of the flux density, $\alpha < 1/4$, the Laughlin state can be found at even lower values of the repulsion U [205, 225].

Table 4.1: Counting of modes $\mathcal{N}_L\left(k_y^A\right)$ in the PES of the Laughlin state for several system sizes and particle partitions. The last column lists the $\mathcal{N}_L(k_y^A)$ values for each momentum sector $k_y^A = 2\pi i/L_y, i = 0, \ldots, L_y - 1$.

N_p	(L_x, L_y)	N_A	PES: $\mathcal{N}_L(k_y^A)$
4	$(4, 8)$	2	$3, 2, 3, 2, 3, 2, 3, 2$
5	$(4, 10)$	2	$4, 3, 4, 3, 4, 3, 4, 3, 4, 3$
6	$(4, 12)$	3	$10, 9, 9, 10, 9, 9, 10, 9, 9, 10, 9, 9$

By analyzing the effective model from Eq. (4.1), we have obtained a guidance for the regime of microscopic parameters and for the geometry of the small system that can give rise to Laughlin states. In the next sections our aim is to go beyond the effective model from Eq. (4.2) and to identify topological states supported by the full driven dynamics as captured by the model given in Eq. (4.1).

4.2 Driven dynamics

In this section we discuss the full driven dynamics as captured by the model given in Eq. (4.1).

4.2.1 Heating

First we address the onset of heating following the standard procedure discussed in Refs. [217, 231]. The initial state of the system is prepared using the ground state of the effective model

$$|\psi(t=0)\rangle = e^{-i\hat{K}(t=0)}|\psi^{0,0}_{\mathrm{LGH}}\rangle \tag{4.6}$$

and we monitor the stroboscopic time-evolution $t = N\,T$, $T \equiv 2\pi/\omega$ governed by the full driven model defined in Eq. (4.1). In our numerical simulations, we approximate the micromotion operator $\hat{K}(t=0)$ using the leading-order high-frequency expansion; see Eq. (G.12). The quantity of interest is the expectation value of the effective Hamiltonian: (4.2)

$$\langle \hat{H}_{\mathrm{eff}}(t=NT)\rangle_K = \langle\psi(t)|e^{-i\hat{K}(t=0)}\hat{H}_{\mathrm{eff}}e^{i\hat{K}(t=0)}|\psi(t)\rangle. \tag{4.7}$$

We expect this quantity to reasonably correspond to the ground-state energy of the effective model E_0 in the regime of very high frequency. On the other hand, for a "low" driving frequency we expect the system to quickly reach the infinite-temperature $\beta \to 0$ regime defined by

$$\lim_{\beta\to 0}\langle \hat{H}_{\mathrm{eff}}\rangle = \frac{1}{\dim\mathcal{H}}\mathrm{tr}\left(\hat{H}_{\mathrm{eff}}\right). \tag{4.8}$$

For this reason we monitor the normalized total energy

$$Q(t=NT) = \frac{\langle \hat{H}_{\mathrm{eff}}(t=NT)\rangle_K - E_0}{\lim_{\beta\to 0}\langle \hat{H}_{\mathrm{eff}}\rangle - E_0} \tag{4.9}$$

and we present it in Fig. 4.3(a), for $U/J_x = 10$. In agreement with the known results [231], we find that the thermalization is quick for both a "high" driving frequency $\omega/J_x \geq 20$ and for a "low" driving frequency $\omega/J_x \leq 10$. For the intermediate values of ω, the heating process is slow [231] and the total energy exhibits a slow exponential growth captured by $Q(t=NT) \approx 1 - b\exp(-c\,t), t \gg 1$. An example of this behavior is given for $\omega/J_x = 15$ in Fig. 4.3(a). The heating process can also be monitored through the particle-entanglement entropy S_A as a function of time. In Fig. 4.3(b) for $N_p = 5$ and low driving frequency we find that this quantity quickly saturates to its maximal value. Indeed, for a thermal state at infinite temperature, S_A is given by

$$S_A^{\max} \approx \ln\binom{L_x\,L_y + N^A - 1}{N^A}, \tag{4.10}$$

marked by the horizontal, dot-dashed line in Fig. 4.3(b). Except for the highest frequency considered ($\omega/J_x = 50$), we find that, in the process of heating, the particle-entanglement gap of the initial state quickly closes (not shown in the plots).

Here we briefly discuss finite-size effects by comparing numerical results for the normalized

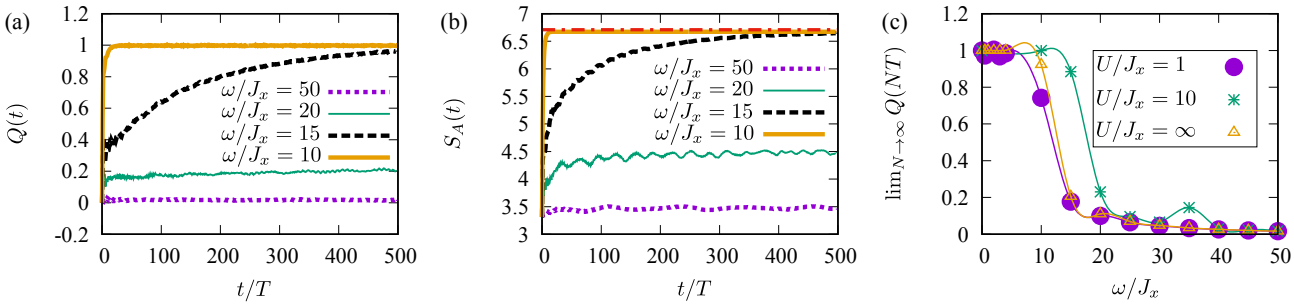

Figure 4.3: (a) The normalized total energy $Q(t = NT)$ from Eq. (4.9), and the (b) particle-entanglement entropy $S_A(t = NT)$, Eq. (4.4), during the time evolution governed by Eq. (4.1) for several driving frequencies $\omega/J_x = 50, 20, 15, 10$. Parameters: $N_p = 5, U/J_x = 10$. Note that the asymptotic value of S_A for $\omega/J_x = 10$ and $\omega/J_x = 15$ matches the one given in Eq. (4.10), as presented by the horizontal, dot-dashed line. (c) The long-time limit $\lim_{N\to\infty} Q(NT)$ for $N_p = 4$ and the on-site interactions $U/J_x = 1, 10$ and $U/J_x = \infty$ (hard-core bosons). The lines are only guides to the eye.

total energy for $N_p = 4$, $N_p = 5$, and $N_p = 6$. In line with the known results [94–96], the "high-frequency" regime with low heating rates moves toward higher ω as the system size increases. However, we find that the estimates obtained in this section ($\omega/J_x \geq 20$ for the high- and $\omega/J_x \leq 10$ for the low-frequency regime, for $U/J_x = 10$) apply to all the three sizes $N_p = 4, 5, 6$, at least for the time-scales that we consider. A comprehensive study of the leading finite-size effects in driven systems can be found in Refs. [94, 217, 231].

4.2.2 The stroboscopic time-evolution operator

In order to better understand the limitations of the effective model, here we time evolve all relevant basis states for a single driving period $T = 2\pi/\omega$ and construct the stroboscopic time-evolution operator

$$\hat{U}_F \equiv \hat{U}(t_0 + T, t_0 = 0) \tag{4.11}$$

such that $\hat{U}(NT+t_0) = \hat{U}_F^N$. In the next step, for a system size $N_p = 4$, $(L_x, L_y) = (4, 8)$ we fully diagonalize this operator and inspect its eigenstates $|n\rangle$. Following the described procedure, we obtain the long-time limit

$$\lim_{N\to\infty} \langle \hat{H}_{\text{eff}}(NT)\rangle_K = \sum_n |\langle n|\psi(t=0)\rangle|^2 \langle n|\hat{H}_{\text{eff}}|n\rangle_K \tag{4.12}$$

where we define

$$\langle n|\hat{H}_{\text{eff}}|n\rangle_K = \langle n|e^{-i\hat{K}(t=0)}\hat{H}_{\text{eff}}e^{i\hat{K}(t=0)}|n\rangle. \tag{4.13}$$

Results for $Q(t = NT)$ from Eq. (4.9) obtained in this way are summarized in Fig. 4.3(c), where we make a comparison between the long-time energies for the case of hard-core bosons ($U \to \infty$) and soft-core bosons (finite values of U). The obtained results indicate that heating rates of hard-core bosons are closer to the case of $U/J_x = 1$ in comparison to $U/J_x = 10$,

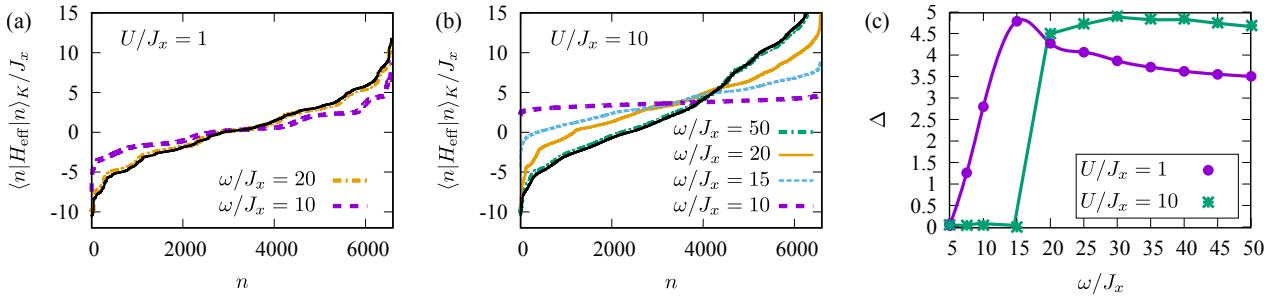

Figure 4.4: Properties of the eigenstates $|n\rangle$ of the stroboscopic time-evolution operator $\hat{U}_F$, Eq. (4.11), in the $k_x = 0, k_y = 0$ sector for $N_p = 4$. Expectation values $\langle n|\hat{H}_{\rm eff}|n\rangle_K$ defined in Eq. (4.13) for (a) $U/J_x = 1$, $\omega/J_x = 10, 20$ and (b) $U/J_x = 10$, $\omega/J_x = 10, 15, 20, 50$. The black solid lines mark eigenenergies of $\hat{H}_{\rm eff}$, Eq. (4.2). Note that in (b) we do not include few states from the top of the spectrum of $\hat{H}_{\rm eff}$, Eq. (4.2), for clarity. (c) The particle-entanglement gap Δ of the incoherent superposition ρ_F, Eq. (4.14), for $U/J_x = 1$ and $U/J_x = 10$, $N_p = 4$. The lines are only guides to the eye.

which is expected from the bandwidths shown in Fig. 4.2(a). Overall we observe that the "high-frequency regime" is wider for lower ratios U/J_x.

In Fig. 4.4, we make a comparison between the exact driven model captured by $\hat{U}_F$ and $\hat{H}_{\rm eff}$. In Figs. 4.4(a) and 4.4(b) we inspect the distribution of expectation values $\langle n|\hat{H}_{\rm eff}|n\rangle_K$. By comparing these values to the eigenenergies of the effective model, Eq. (4.2), we get an insight into the pertinence of the effective description [94, 95]. In particular, for an interacting system in the thermodynamic limit, the distribution is flat and the effective description is useless. We state again that we consider only small atomic samples. For this reason, it is expected that for high values of ω the full stroboscopic description nicely matches to the effective model values. Such an example is given in Fig. 4.4(a) for $U/J_x = 1$ and $\omega/J_x = 20$. As the value of ω gets lower the distribution becomes flatter, as can be seen in Fig. 4.4(b) for $U/J_x = 10$ by comparing results for $\omega/J_x = 50$ and $\omega/J_x = 10$.

The intermediate regime of frequencies, e. g., $\omega/J_x = 20$ for $U/J_x = 10$, is of the main experimental relevance [88]. We now investigate whether the driven stroboscopic dynamics supports some Laughlin-like states, by calculating the PES of the mixture

$$\rho_F = \frac{1}{2}\Big(|n_0(0,0)\rangle\langle n_0(0,0)| + |n_0(0,\pi)\rangle\langle n_0(0,\pi)|\Big) \tag{4.14}$$

where $|n_0(k_x, k_y)\rangle$ is the state from the k_x, k_y sector with the lowest expectation value $\langle n|\hat{H}_{\rm eff}|n\rangle_K$. The results are presented in Fig. 4.4(c). We find that the states with a well defined gap and the Laughlin-like PES can be found down to $\omega/J_x \geq 10$ for $U/J_x = 1$, and down to $\omega/J_x \geq 20$ for $U/J_x = 10$. Having established existence of these states for small samples of $N_p = 4$ particles, in the next section we discuss dynamical protocol which can be exploited to prepare these states.

4.3 Slow ramp

The question about an optimal adiabatic protocol that can be used to prepare the Laughlin state in a cold-atom setup has gained lot of attention [181, 202, 203, 211]. The situation becomes even more complex once the full driving process is taken into account. A general wisdom is that, by starting from a topologically trivial state, the topological index of a thermodynamically large system cannot be changed adiabatically. We consider a small atomic sample and follow the proposal of Ref. [211]. Our main contribution is that we extend this protocol to the case of the driven, interacting Bose-Hubbard model.

4.3.1 Model

Following results of Ref. [211], we consider a slow ramp of the tunneling amplitude along y direction, $J_y(t)$, as well as a slow ramp of the driving amplitude $\kappa(t)$. Namely, we start from a series of decoupled wires along the x direction and start coupling them. More precisely, initial states are selected as the ground states of $\hat{H}_{\rm ini}$:

$$\hat{H}_{\rm ini} = -J_x \sum_{m,n} \left(\hat{a}^\dagger_{m+1,n} \hat{a}_{m,n} + \text{H. c.} \right) + \frac{U}{2} \sum_{m,n} \hat{n}_{m,n} (\hat{n}_{m,n} - 1). \tag{4.15}$$

For the filling factors that we consider, the ground states of the $\hat{H}_{\rm ini}$ are simple noninteracting states with the ground state energy $E_{0,\rm ini} = -2J_x N_p$. Out of the several degenerate ground states, we select those where atoms occupy every second wire. There are two such states and we label them as $|\psi_+\rangle$ (even wires occupied) and $|\psi_-\rangle$ (odd wires occupied). These states have finite projections only onto the sectors $k_x = 0, k_y = 0$ and $k_x = 0, k_y = \pi$ of the driven model from Eq. (4.1). Therefore we may expect the two initial states $|\psi_\pm(t=0)\rangle$ to be transformed into the two Laughlin states during the ramp.

Having prepared the initial state, we slowly restore the tunneling amplitude along the y direction, $J_y(t)$, and slowly ramp up the driving amplitude $\kappa(t)$. The time-evolution is governed by

$$\begin{aligned} \hat{H}_{\rm sr}(t) &= -J_x \sum_{m,n} \left(\hat{a}^\dagger_{m+1,n} \hat{a}_{m,n} + \text{H. c.} \right) - J_y(t) \sum_{m,n} \left(e^{i\omega t} \hat{a}^\dagger_{m,n+1} \hat{a}_{m,n} + \text{H. c.} \right) \\ &+ \frac{\kappa(t)}{2} \sum_{m,n} \sin\left(\omega t - (m+n-1/2)\,\phi\right) \hat{n}_{m,n} + \frac{U}{2} \sum_{m,n} \hat{n}_{m,n}(\hat{n}_{m,n} - 1), \end{aligned} \tag{4.16}$$

where $J_y(t) = J_y \tanh(\eta\, t)$, $\kappa(t) = \kappa \tanh(\eta\, t)$, η being the ramping rate. In the long-time limit, we recover the original Hamiltonian from Eq. (4.1). During the ensuing time evolution we

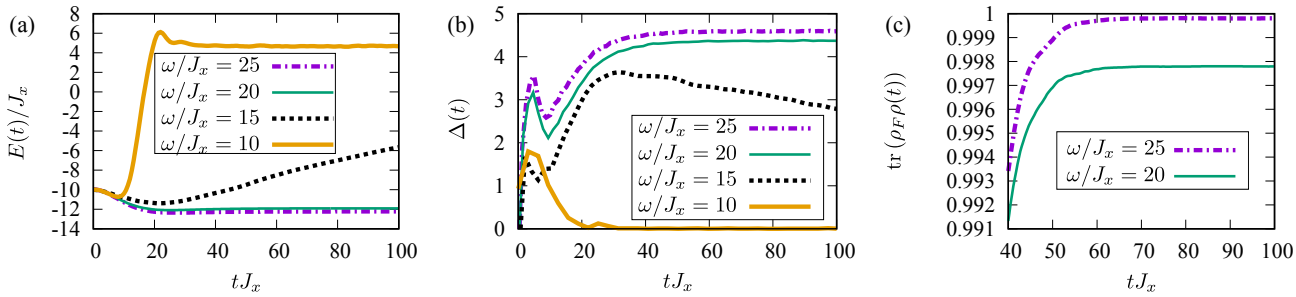

Figure 4.5: (a) The expectation value $E(t)$ defined in Eq. (4.18) and (b) the particle-entanglement gap $\Delta(t)$ of $\rho(t)$, Eq. (4.17), during the time evolution governed by Eq. (4.16) for several driving frequencies $\omega/J_x = 25,\ 20,\ 15,\ 10$. Parameters: $N_p = 5, U/J_x = 10, \eta/J_x = 0.05$. (c) The overlap $\mathrm{tr}\,(\rho(t)\rho_F)$ of the time evolved state with the target eigenstates of $\hat{U}_F$ for $\omega/J_x = 25,\ 20$. Parameters: $N_p = 4, U/J_x = 10, \eta/J_x = 0.05$.

construct the mixture

$$\rho(t) = \frac{1}{2}\Big(|\psi_+(t)\rangle\langle\psi_+(t)| + |\psi_-(t)\rangle\langle\psi_-(t)|\Big). \tag{4.17}$$

We monitor stroboscopically the energy expectation value

$$E(t) = \mathrm{tr}\left(\rho(t)\hat{H}_{\mathrm{eff}}\right) \tag{4.18}$$

and the PES of $\rho(t)$.

4.3.2 Results

In Fig. 4.5(a) we present the energy expectation value from Eq. (4.18) for $U/J_x = 10$ and several driving frequencies $\omega/J_x = 25,\ 20,\ 15,\ 10$. Our numerical results indicate that ramps with the rates up to $\eta/J_x \sim 0.1$ work reasonably well. Slower ramps give better results, but are less practical [211]. By construction, the initial state is a noninteracting state with particles delocalized along the x direction and therefore the initial energy is $E(t = 0) = -2\,N_p\,J_x$. During the ramp with the rate $\eta/J_x = 0.05$, for the regime of high driving frequencies, down to approximately $\omega/J_x = 20$, we find that the energy initially decreases and reaches an almost constant value at around $tJ_x \sim 20$. On the other hand, for $\omega/J_x = 15$, the system slowly heats up during the ramping process, and for $\omega/J_x = 10$ the system quickly reaches the infinite-temperature state.

One of our main results is summarized in Fig. 4.5(b), where we plot the particle-entanglement gap of $\rho(t)$, from Eq. (4.17), as a function of time. In the high-frequency regime $\omega/J_x \geq 20$, starting around $tJ_x \sim 20$ we find a persistent particle-entanglement gap, marking the onset of a topologically nontrivial state. It is even more interesting that, even for $\omega/J_x \sim 15$, the state seems to exhibit a finite gap on intermediate time-scales. This is not the case for $\omega/J_x \leq 10$, where the gap quickly vanishes. In Fig. 4.5(c), we present the value of the overlap $\mathrm{tr}\,(\rho(t)\rho_F)$,

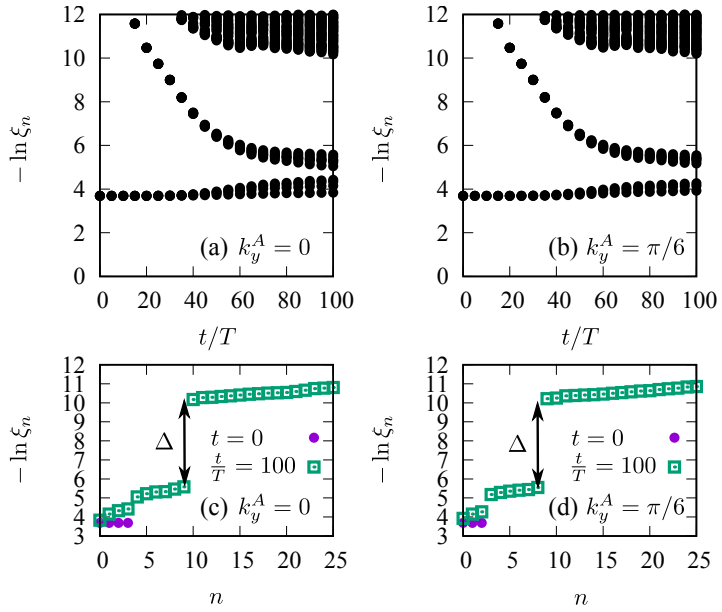

Figure 4.6: The low-lying part of the particle-entanglement spectra $-\ln \xi_n$ of $\rho(t)$, Eq. (4.17), during the time evolution governed by Eq. (4.16) in the (a) $k_y^A = 0$ and (b) $k_y^A = \pi/6$ momentum sectors. The low-lying part of the PES in the sectors (c) $k_y^A = 0$ and (d) $k_y^A = \pi/6$, at two instances of time $t = 0$ and $t/T = 100$. Parameters: $N_p = 6, U/J_x = 5, \omega/J_x = 15, \eta/J_x = 0.05$.

of the time-evolved mixed state with the relevant state from Eq. (4.14) for $N_p = 4$. Clearly, the slow ramp of the type given in Eq. (4.16) allows for the preparation of the relevant eigenstates of $\hat{U}_F$ with high fidelity (better than 1%).

In Figs. 4.6(a) and 4.6(b) we show the time evolution of the PES in the two momentum sectors $k_y^A = 0$ and $k_y^A = \pi/6$ for $N_p = 6$, $U/J_x = 5$, and $\eta/J_x = 0.05$. The PES of the initial state is easy to understand. As the $L_y/2$ wires are occupied by single atoms, the reduced density matrix is proportional to the identity matrix with the proportionality factor yielding $-\ln \xi_n = \ln\left(2\binom{Ly/2}{N_A}\right) \approx 3.69$. During the ramp we find that additional modes in PES are gaining weight and moving down in the spectrum. Finally, the state $\rho(t)$ reached around $t \approx 50T$ exhibits a well defined gap and the correct counting of the low-lying modes: there are ten low-lying modes for $k_y^A = 0$ and nine low-lying modes for $k_y^A = \pi/6$; see Figs. 4.6(c) and 4.6(d) and also Table 4.1.

In Fig. 4.7 we discuss a satisfactory range of ramping rates η for a given interaction strength U and a given driving frequency ω that we fix at $\omega/J_x = 15$. The obtained numerical results suggest that at weaker interaction strengths $U/J_x \leq 2$, slower ramping rates are needed. One way to explain this behavior is by using the effective model and arguing that the gap protecting the Laughlin state is smaller at weaker U. On the other hand, for stronger interaction strengths $U/J_x \geq 8$ the particle-entanglement gap closes at later stages as the heating process becomes dominant. Finally, in the intermediate range $U/J_x \sim 5$, faster ramps with $\eta/J_x = 0.1$ lead to the sought-after state $\rho(t)$ from Eq. (4.17), with persistent features in the PES up to $t = 500T$. These results indicate that, when optimizing the ramping protocol in an actual experiment,

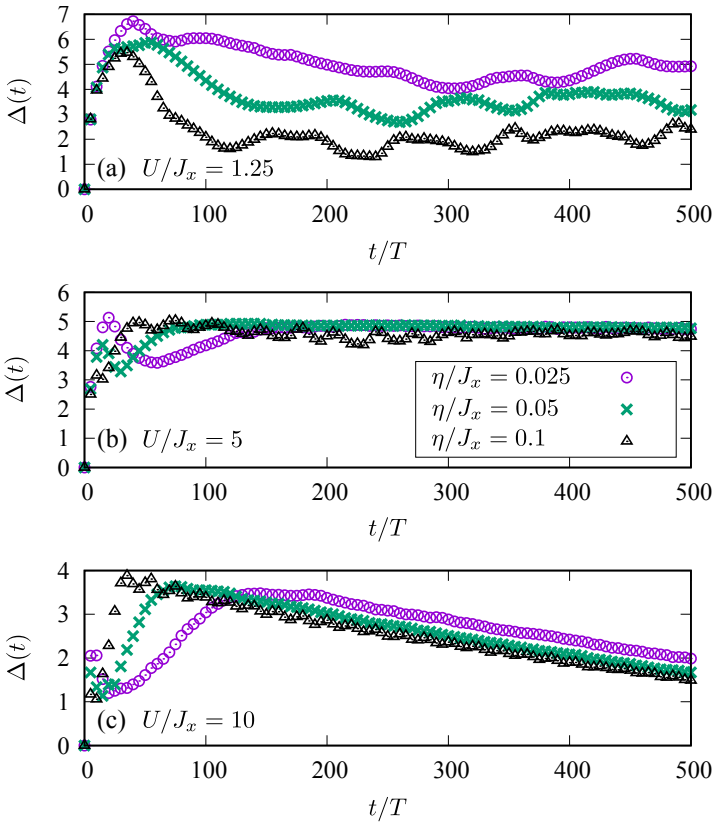

Figure 4.7: The particle-entanglement gap $\Delta(t)$ as a function of time during the time evolution governed by Eq. (4.16), for several interaction strengths (a) $U/J_x = 1.25$ (b) $U/J_x = 5$ and (c) $U/J_x = 10$, and several ramping rates $\eta/J_x = 0.025, 0.05, 0.1$. Other parameters: $N_p = 5, \omega/J_x = 15$.

there will be a tradeoff between the unfavorable heating and a faster ramping into the desired state, as both of these processes are promoted by interactions.

4.4 Conclusions

The technique of Floquet engineering has been successfully exploited for the implementation of synthetic magnetic fields in driven optical lattices. Following up on these achievements and on a long-standing pursuit for the FQH states in cold-atom setups, in this Chapter we have addressed possible realization of the bosonic Laughlin state in a small atomic sample in a periodically driven optical lattice. While a thermodynamically large interacting system generally heats up into an infinite-temperature state under driving, the heating process can be controlled to some extent in a few-particle system.

We have assumed a realistic driving protocol and finite on-site interactions, and we have identified the FQH state based on analysis of its particle-entanglement spectra. Results of our numerical simulations show that the stroboscopic dynamics of $N_p = 4, 5, 6$ particles supports the topological $\nu = 1/2$ Laughlin state down to $\omega/J_x = 20$ for $U/J_x = 10$, and down to $\omega/J_x = 15$ for $U/J_x = 1$, for the driving amplitude $\kappa/\omega = 0.5$. These results are in reasonable

agreement with the recent estimates of the optimal heating times [88] that take into account the contribution of the higher bands of the underlying optical lattice. In addition, we have investigated slow ramping of the driving term and found that it allows for the preparation of the Laughlin state on experimentally realistic time-scales of the order of $20\,\hbar/J_x$, where $\hbar/J_x$ is the tunneling time. Interestingly, we find that some topological features persist during an intermediate stage even in the regime where the system exhibits a slow transition into the infinite-temperature state (e. g., $\omega/J_x = 15$ for $U/J_x = 10$).

In the future, we plan to address the preparation scheme for the relevant correlated states in a driven honeycomb lattice, which exhibits lower heating rates in comparison to a cubic lattice according to the recent experiments [124, 220]. Another highly relevant question, that we have not tackled and that we postpone to future investigation, concerns suitable experimental probes of topological features. The recent progress in the field has led to the development of several detection protocols specially suited for the cold-atom systems [167, 168, 232–235]. For the type of systems considered in this Chapter, the most promising are results of the recent study [235] showing that fractional excitations can be probed even in small systems of several bosons.

Chapter 5

Conclusions

The study of ultracold quantum gases is an important topic in modern physics. The possibility of using ultracold atoms to build versatile quantum simulators is especially promising. These are highly-controllable macroscopic many-particle systems that obey the laws of quantum mechanics and can be used to simulate and study other complex quantum systems, such as those relevant for condensed-matter physics.

Quantum many-body scars have recently been introduced as new paradigm of weak ergodicity breaking in interacting quantum systems. This phenomenon has been first observed in the form of persistent oscillations in the quench dynamics probed in experiments on a Rydberg atom quantum simulator [46]. As a weak form of ergodicity breaking, many-body scars are believed to constitute a new universality class of systems that are distinct from other types of strong ergodicity breaking, such as in integrable models and many-body localized systems. Currently, major efforts are under way to understand the origins of quantum many-body scars. Similar properties have been found in other physical systems, including lattice gauge theories and topological phases of matter.

In Chapter 2 we proposed a realization of quantum many-body scars in a one-dimensional bosonic lattice with kinetically-constrained hopping. This model could be experimentally realized in ultracold atoms under a suitable Floquet engineering scheme. An important open question in this field was the necessity of hard kinetic constraints for the realization of scars. The standard "PXP" model of Rydberg atoms exhibits hard kinetic constraints, while some theoretical works on certain spin models suggested that constraints may not be necessary. Our bosonic models with density-dependent hopping provided a way to tunably control the kinetic constraints and study both limits on an equal footing. Using this approach we demonstrated that scars can occur in the absence of hard kinetic constraints. Another open problem was the relation between scars and integrability, following the observation that certain perturbations can enhance scarring while at the same time making the system non-ergodic. From this point

of view, our bosonic models are important because they demonstrate the presence of scars in a robustly non-integrable regime. Finally, our work points to an experimental platform that could realize quantum many-body scars and potentially allow new probes of this complex phenomenon. To this day there is only a single experiment on a system of Rydberg atoms, in contrast to the wealth of theoretical results. Given the large body of work on bosonic ultracold atoms and the variety of available experimental probes, such systems are the prime candidates for further experimental progress on many-body scars.

The effects of both strong local and long-range interactions have been studied in recent cold-atom experiments, as they can lead to rich phase diagrams. In order to observe topological phases of matter, the effects of magnetic field were also included in some recent experiments. Using periodic driving, synthetic magnetic fields have been experimentally realized in cold-atom systems, which has enabled the realization of seminal condensed-matter models such as the Harper-Hofstedter and Haldane model in cold-atom setups. However the interplay of driving and interactions introduces heating that may preclude the study of relevant topological states. A possibility of finding optimized parameter regimes with slow heating rates is still open and it was the main research topic in the second part of this thesis. These results could contribute to implementing new, even more advanced simulations, which could reveal new quantum phases.

The recent Chern-number measurement [70] was a milestone marking a realization of a topological band in a cold-atom setup. Motivated by the this experiment, in Chapter 3 we investigated the response of incoherent bosons to an external force in driven optical lattices featuring topological bands. The focus of this study was on the role of weak atomic interactions. Using numerical simulations based on a classical-field method, we found that interactions contribute to atomic transitions between different bands, thus complicating the experimental procedure in line with expectations. However, it was also shown that the weak atomic repulsion makes the Chern-number measurement easier in several ways. As this experimental approach is expected to become a routine tool in the near future, a first step in the preparation of more interesting topological phases, the so-far obtained results on the effects of weak interactions are of relevance for the future experiments.

Nowadays, cold-atom setups provide access to both strong synthetic magnetic fields and strong interactions. These ingredients are in principle enough to realize fractional Hall states and address their excitations. The complexity arises when using the driving protocol in the strongly interacting regime, due to heating. A solution for this problem would be to find some optimal parameter regime where the system stays in the prethermalized state for long enough time. At the moment, the possibility of finding fractional quantum Hall states in cold-atom setups is still open and the questions about how to prepare and manipulate these states using cold atoms prompt further theoretical studies. Our work on the stability and lifetime of bosonic fractional quantum Hall states presented in Chapter 4 should provide guidelines in this direction. In our study we took into account important experimental features, such as

a realistic driving scheme and finite on-site interactions. We used the particle-entanglement spectra in order to confirm that the state prepared in our driven model for a high enough driving frequency is indeed a bosonic Laughlin state. By performing numerical simulations, we identified an optimal regime of microscopic parameters for the preparation of these interesting strongly correlated states.

There are many possibilities to further extend the research presented in this thesis. Potential directions of future research are discussed at the ends of Chapters 2-4. It would be particularly exciting to explore whether it is possible to realize quantum many-body scars in topologically-nontrivial driven systems, as this would be a way to slow down the thermalization and increase the lifetimes of interesting topological states.

Appendix A

Relative magnitude of the hopping coefficients

Constraints are not the only factor that slows down the dynamics and leads to weakly-entangled eigenstates in spectrum of the Hamiltonian $\hat{H}_1$ (2.1) from Chapter 2. The relative magnitude of the hopping coefficients between different configurations mapped to each other under the action of the Hamiltonian also has important effects. In order to show this, we introduce two additional Hamiltonians

$$\hat{H}_{1a} = -J \sum_j \left(\hat{b}_j^\dagger \hat{b}_{j+1} \hat{n}_j^2 + \hat{n}_{j-1}^2 \hat{b}_j^\dagger \hat{b}_{j-1} \right), \tag{A.1}$$

$$\hat{H}_{1b} = -J \sum_j \left(\hat{b}_j^\dagger \hat{b}_{j+1} (1 - \delta_{n_j,0}) + (1 - \delta_{n_{j-1},0}) \hat{b}_j^\dagger \hat{b}_{j-1} \right). \tag{A.2}$$

These two Hamiltonians have the same constraints as the Hamiltonian $\hat{H}_1$ and therefore the same graphs as in Fig. 2.1. However, the particle number operators $\hat{n}_j$ are squared in $\hat{H}_{1a}$ and replaced with delta functions $(1 - \delta_{n_j,0})$ in $\hat{H}_{1b}$. This makes the minimal and extended clusters even less connected to the rest of the configurations in the case of $\hat{H}_{1a}$ and more connected in the case of $\hat{H}_{1b}$.

As anticipated, the revivals become more prominent for $\hat{H}_{1a}$, with fidelity peaks reaching more than 95%, while the peaks almost disappear for $\hat{H}_{1b}$, as illustrated in Fig. A.1(a). In addition, the entanglement entropy quickly saturates in the case $\hat{H}_{1b}$, while the growth is significantly suppressed in the case of $\hat{H}_{1a}$, as can be observed in Fig. A.1(b). The distribution of entanglement entropy across all eigenstates is also affected by the change of coefficients (not shown). $\hat{H}_{1a}$ has a spectrum with many low-entropy eigenstates, while the spectrum of $\hat{H}_{1b}$ is almost thermal and resembles that of $\hat{H}_2$. The probability distribution of consecutive gaps in the energy spectrum of $\hat{H}_{1a}$ is close to the Poisson distribution, which implies that $\hat{H}_{1a}$ is almost integrable. On the other hand, the distribution for $\hat{H}_{1b}$ is Wigner-Dyson, like in $\hat{H}_1$.

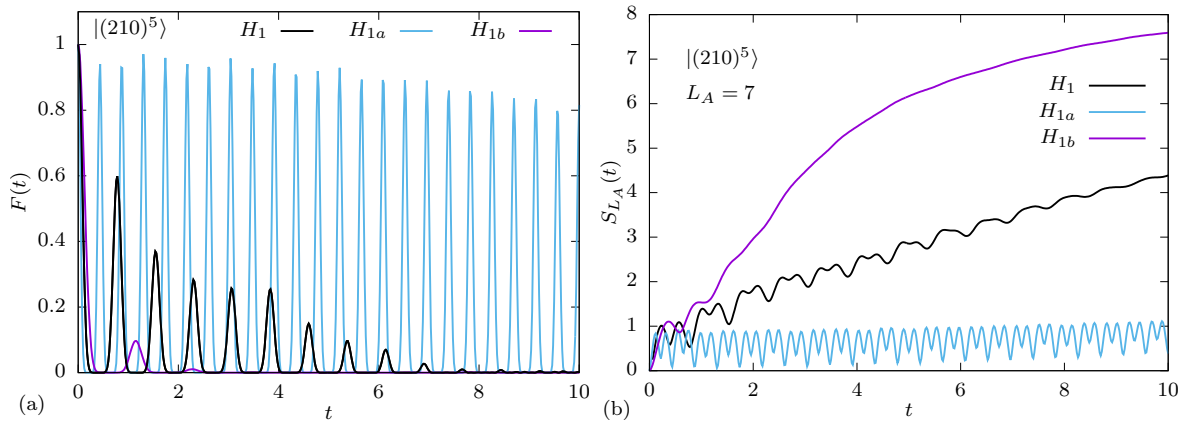

Figure A.1: Time evolution of (a) fidelity $|\langle\psi(t)|\psi(0)\rangle|^2$ and (b) entanglement entropy $S_{L_{\mathrm{A}}=7}(t)$ governed by three different Hamiltonians, $\hat{H}_1$, $\hat{H}_{1\mathrm{a}}$ and $\hat{H}_{1\mathrm{b}}$. System size $L = 15$ and initial state $|(210)^5\rangle$.

Appendix B

Derivation of fidelity in the extended cluster approximation

Here we derive Eq. (2.28) from Chapter 2. For system size $L = 3$, the Hilbert space of the extended cluster is spanned by only four configurations:

$$\begin{pmatrix}1\\0\\0\\0\end{pmatrix} = |300\rangle, \quad \begin{pmatrix}0\\1\\0\\0\end{pmatrix} = |210\rangle, \quad \begin{pmatrix}0\\0\\1\\0\end{pmatrix} = |120\rangle, \quad \begin{pmatrix}0\\0\\0\\1\end{pmatrix} = |111\rangle. \tag{B.1}$$

The Hamiltonian reduced to this subspace is

$$\hat{H}_1^{\check{c}} = -\begin{pmatrix} 0 & 2\sqrt{3} & 0 & 0 \\ 2\sqrt{3} & 0 & 2 & 0 \\ 0 & 2 & 0 & \sqrt{2} \\ 0 & 0 & \sqrt{2} & 0 \end{pmatrix}. \tag{B.2}$$

Its eigenvalues are

$$E_1 = -\alpha, \; E_2 = \alpha, \; E_3 = -\beta, \; E_4 = \beta, \tag{B.3}$$

and its eigenvectors

$$|1\rangle = \begin{pmatrix} a \\ b \\ c \\ d \end{pmatrix}, \; |2\rangle = \begin{pmatrix} -a \\ b \\ -c \\ d \end{pmatrix}, \; |3\rangle = \begin{pmatrix} -c \\ -d \\ a \\ b \end{pmatrix}, \; |4\rangle = \begin{pmatrix} c \\ -d \\ -a \\ b \end{pmatrix}, \tag{B.4}$$

where $\alpha = \sqrt{9+\sqrt{57}} \approx 4.06815$, $\beta = \sqrt{9-\sqrt{57}} \approx 1.20423$, $a \approx 0.591050$, $b \approx 0.694113$, $c \approx 0.388150$ and $d \approx 0.134933$. There are no simple analytical expressions for the coefficients a, b, c and d.

The configuration $|210\rangle$ evolves as

$$\begin{aligned} |\psi_1^{\tilde{c}}(t)\rangle \;=\; & -2i\,(ab\sin\alpha t + cd\sin\beta t)\,|300\rangle \\ & +2\left(b^2\cos\alpha t + d^2\cos\beta t\right)|210\rangle \\ & -2i\,(bc\sin\alpha t - ad\sin\beta t)\,|120\rangle \\ & +2bd\,(\,\cos\alpha t - \,\cos\beta t)\,|111\rangle, \end{aligned} \tag{B.5}$$

which can also be generalized to larger systems

$$|\Psi_n^{\tilde{c}}(t)\rangle = |(210)^n(t)\rangle = 2^n\left(b^2\cos\alpha t + d^2\cos\beta t\right)^n|(210)^n\rangle + ... \tag{B.6}$$

Finally, the fidelity evolves as

$$F_n^{\tilde{c}}(t) = |\langle\Psi_n^{\tilde{c}}(0)|\Psi_n^{\tilde{c}}(t)\rangle|^2 = 4^n|b^2\cos\alpha t + d^2\cos\beta t|^{2n}. \tag{B.7}$$

The period of revivals is approximately $T \approx \pi/\alpha \approx 0.772241$, and the first peak height exponentially decreases as

$$F_{L=3n}^{\tilde{c}}(T) = 4^n|d^2\cos\frac{\pi\beta}{\alpha} - b^2|^{2n} \approx 0.887017^n \approx \mathrm{e}^{-0.119891n} \approx \mathrm{e}^{-0.039964L}. \tag{B.8}$$

Eqs. (B.6) and (B.7) are exact for non-translation-invariant initial states, but just an approximation for the translation symmetric case. This is due to the fact that different translations of 300, 210 and 120 no longer evolve independently in that case, as they are connected to each other through the configuration 111. However, this approximation becomes better with increasing the system size, as the configuration $(111)^n$ becomes further away from the initial state $(210)^n$ and the probability that this configuration will be reached decreases.

Appendix C

The effective model from Chapter 3

In this Appendix we derive the effective Hamiltonian $\hat{H}_{\text{eff},1}$ (3.7) which corresponds to the time-dependent Hamiltonian (3.3) from Chapter 3. We also show the explicit form of its momentum-space representation $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$.

C.1 Real space

After a unitary transformation into the rotating frame $\tilde{\psi} = \mathrm{e}^{-i\hat{W}t}\psi$, where $\tilde{\psi}$ and ψ are the old and the new wave functions, and $\hat{W}$ is the staggered potential, the new time-dependent Hamiltonian that describes the experimental setup is given by [70]

$$\hat{H}(t) = J_y \sum_{l,m} \left(\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m} \right) + \hat{V}^{(+1)}\mathrm{e}^{i\omega t} + \hat{V}^{(-1)}\mathrm{e}^{-i\omega t} + \frac{U}{2}\sum_{l,m} \hat{n}_{l,m}\left(\hat{n}_{l,m} - 1\right), \quad \text{(C.1)}$$

where

$$\hat{V}^{(+1)} = \kappa/2 \sum_{l,m} \hat{n}_{l,m} g(l,m) - J_x \sum_{l_{\text{odd}},m} \left(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m} \right) \quad \text{(C.2)}$$

$$\hat{V}^{(-1)} = \kappa/2 \sum_{l,m} \hat{n}_{l,m} g^*(l,m) - J_x \sum_{l_{\text{even}},m} \left(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m} \right) \quad \text{(C.3)}$$

$$g(l,m) = \cos(l\pi/2 - \pi/4)\mathrm{e}^{i(\phi_0 - m\pi/2)} + \cos(l\pi/2 + \pi/4)\mathrm{e}^{i(m\pi/2 - \phi_0 - \pi/2)}. \quad \text{(C.4)}$$

The kick operator is given by

$$\hat{K}(t) = \frac{1}{i\omega}\left(\hat{V}^{(+1)}\mathrm{e}^{i\omega t} - \hat{V}^{(-1)}\mathrm{e}^{-i\omega t} \right) + \mathcal{O}\left(\frac{1}{\omega^2}\right), \quad \text{(C.5)}$$

and the effective Hamiltonian by

$$\hat{H}_{\text{eff}} = \underbrace{\hat{H}_0}_{\hat{H}^{(0)}_{\text{eff}}} + \underbrace{\frac{1}{\omega}\left[\hat{V}^{(+1)}, \hat{V}^{(-1)}\right]}_{\hat{H}^{(1)}_{\text{eff}}} + \underbrace{\frac{1}{2\omega^2}\left(\left[\left[\hat{V}^{(+1)}, \hat{H}_0\right], \hat{V}^{(-1)}\right] + \left[\left[\hat{V}^{(-1)}, \hat{H}_0\right], \hat{V}^{(+1)}\right]\right)}_{\hat{H}^{(2)}_{\text{eff}}} + \mathcal{O}\left(\frac{1}{\omega^3}\right). \tag{C.6}$$

If we assume that the driving frequency is high and interactions are weak, the interaction term and almost all $\mathcal{O}\left(\frac{1}{\omega^2}\right)$ terms can be neglected. After substituting Eqs. (C.1), (C.2) and (C.3) into Eq. (C.6) we obtain:

$$\hat{H}^{(0)}_{\text{eff}} = -J_y \sum_{l,m}\left(\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m}\right) \tag{C.7}$$

$$\begin{aligned}\hat{H}^{(1)}_{\text{eff}} =& \frac{1}{\omega}\left[\frac{\kappa}{2}\sum_{l,m}\hat{a}^\dagger_{l,m}\hat{a}_{l,m}\; g(l,m) - J_x \sum_{l_{\text{odd}},m}\left(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}\right),\right.\\ &\left.\frac{\kappa}{2}\sum_{l,m}\hat{a}^\dagger_{l,m}\hat{a}_{l,m}\; g^*(l,m) - J_x \sum_{l_{\text{even}},m}\left(\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}\right)\right]\\ =& \hat{H}_1 + \hat{H}_2 + \hat{H}_3 + \hat{H}_4.\end{aligned} \tag{C.8}$$

We will now separately calculate each term:

$$\begin{aligned}\hat{H}_1 =& -\frac{J_x\kappa}{2\omega}\sum_{l_{\text{odd}},m,l',m'} g^*(l',m')\left[\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}, \hat{a}^\dagger_{l',m'}\hat{a}_{l',m'}\right]\\ =& -\frac{J_x\kappa}{2\omega}\sum_{l_{\text{odd}},m}\left[\Big(g^*(l,m) - g^*(l+1,m)\Big)\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \Big(g^*(l,m) - g^*(l-1,m)\Big)\hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}\right]\end{aligned} \tag{C.9}$$

$$\begin{aligned}\hat{H}_2 =& -\frac{J_x\kappa}{2\omega}\sum_{l_{\text{even}},m,l',m'} g(l',m')\left[\hat{a}^\dagger_{l',m'}\hat{a}_{l',m'}, \hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}\right]\\ =& \frac{J_x\kappa}{2\omega}\sum_{l_{\text{even}},m}\left[\Big(g(l,m) - g(l+1,m)\Big)\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \Big(g(l,m) - g(l-1,m)\Big)\hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}\right]\end{aligned} \tag{C.10}$$

$$\begin{aligned}\hat{H}_3 =& \frac{J_x^2}{\omega}\sum_{l_{\text{odd}},m,l'_{\text{even}},m'}\left[\hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m}, \hat{a}^\dagger_{l'+1,m'}\hat{a}_{l',m'} + \hat{a}^\dagger_{l'-1,m'}\hat{a}_{l',m'}\right]\\ =& \frac{J_x^2}{\omega}\sum_{l_{\text{odd}},m}\left(2\hat{a}^\dagger_{l+1,m}\hat{a}_{l+1,m} + \hat{a}^\dagger_{l+3,m}\hat{a}_{l+1,m} + \hat{a}^\dagger_{l-1,m}\hat{a}_{l+1,m} - 2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} - \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} - \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m}\right)\\ =& \frac{J_x^2}{\omega}\sum_{l,m}(-1)^l\left(2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m}\right)\end{aligned} \tag{C.11}$$

$$\hat{H}_4 = \frac{\kappa^2}{4\omega} \sum_{l,m,l',m'} g(l,m)g^*(l',m') \left[\hat{a}^\dagger_{l,m}\hat{a}_{l,m}, \hat{a}^\dagger_{l',m'}\hat{a}_{l',m'}\right] = 0. \tag{C.12}$$

Using trigonometric identities and

$$\begin{aligned} g(l,m) - g(l \pm 1, m) = \pm \sqrt{2}\Big(& \sin((2l \pm 1 - 1)\pi/4) e^{i(\pi/4 - m\pi/2)} \\ & + \sin((2l \pm 1 + 1)\pi/4) e^{i(m\pi/2 - 3\pi/4)} \Big), \end{aligned} \tag{C.13}$$

we can rewrite the sum of terms (C.9) and (C.10) in a more convenient form

$$\hat{H}_1 + \hat{H}_2 = \frac{J_x\kappa}{\sqrt{2}\omega} \sum_{l,m} \left(e^{i\left((m-l)\pi/2 - \pi/4\right)} \hat{a}^\dagger_{l,m}\hat{a}_{l-1,m} + e^{-i\left((m-l-1)\pi/2 - \pi/4\right)} \hat{a}^\dagger_{l,m}\hat{a}_{l+1,m} \right). \tag{C.14}$$

The only $\mathcal{O}\left(\frac{1}{\omega^2}\right)$ ($\hat{H}^{(2)}_{\text{eff}}$) term that cannot be neglected in the parameter range that we use is [70]

$$\frac{J_y}{2}\frac{\kappa^2}{\omega^2} \sum_{l,m} \left(\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m} \right). \tag{C.15}$$

Finally, the effective Hamiltonian becomes

$$\begin{aligned} \hat{H}_{\text{eff},1} = & \frac{J_x\kappa}{\sqrt{2}\omega} \sum_{l,m} \left(e^{i\left((m-l-1)\pi/2 - \pi/4\right)} \hat{a}^\dagger_{l+1,m}\hat{a}_{l,m} + e^{-i\left((m-l)\pi/2 - \pi/4\right)} \hat{a}^\dagger_{l-1,m}\hat{a}_{l,m} \right) \\ & - J_y\left(1 - \frac{1}{2}\frac{\kappa^2}{\omega^2}\right) \sum_{l,m} \left(\hat{a}^\dagger_{l,m+1}\hat{a}_{l,m} + \hat{a}^\dagger_{l,m-1}\hat{a}_{l,m} \right) \qquad \text{(C.16)} \\ & + \frac{J_x^2}{\omega} \sum_{l,m} (-1)^l \left(2\hat{a}^\dagger_{l,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l+2,m}\hat{a}_{l,m} + \hat{a}^\dagger_{l-2,m}\hat{a}_{l,m} \right) \qquad \text{(C.17)} \end{aligned}$$

with the renormalized nearest-neighbor hopping amplitudes

$$J'_x = \frac{J_x\kappa}{\sqrt{2}\omega} = J_y \tag{C.18}$$

and

$$J'_y = J_y\left(1 - \frac{1}{2}\frac{\kappa^2}{\omega^2}\right), \tag{C.19}$$

and a next-nearest-neighbor along $\mathbf{e}_x$ hopping term proportional to $\frac{J_x^2}{\omega}$ in (C.17).

C.2 Momentum space

If we choose the unit cell as in Fig. 3.1(a) (lattice sites $\mathtt{A} = (1,0)$, $\mathtt{B} = (2,0)$, $\mathtt{C} = (3,0)$ and $\mathtt{D} = (4,0)$), the momentum-space representation of the effective Hamiltonian without correction

$\hat{H}_{\text{eff},0}$ (3.6) is given by a 4×4 matrix

$$\hat{\mathcal{H}}_{\text{eff},0}(\mathbf{k}) = \begin{pmatrix} 0 & J'_x e^{-i\frac{3\pi}{4}} - J'_y e^{-i\mathbf{k}\mathbf{R}_2} & 0 & J'_x e^{-i\frac{3\pi}{4}-i\mathbf{k}\mathbf{R}_1} - J'_y e^{i\mathbf{k}(\mathbf{R}_2-\mathbf{R}_1)} \\ J'_x e^{i\frac{3\pi}{4}} - J'_y e^{i\mathbf{k}\mathbf{R}_2} & 0 & J'_x e^{-i\frac{\pi}{4}} - J'_y e^{-i\mathbf{k}\mathbf{R}_2} & 0 \\ 0 & J'_x e^{i\frac{\pi}{4}} - J'_y e^{i\mathbf{k}\mathbf{R}_2} & 0 & J'_x e^{i\frac{\pi}{4}} - J'_y e^{-i\mathbf{k}\mathbf{R}_2} \\ J'_x e^{i\frac{3\pi}{4}+i\mathbf{k}\mathbf{R}_1} - J'_y e^{i\mathbf{k}(\mathbf{R}_1-\mathbf{R}_2)} & 0 & J'_x e^{-i\frac{\pi}{4}} - J'_y e^{i\mathbf{k}\mathbf{R}_2} & 0 \end{pmatrix}, \quad \text{(C.20)}$$

where $\mathbf{R}_1$ and $\mathbf{R}_2$ are the lattice vectors $\mathbf{R}_1 = (4,0)$ and $\mathbf{R}_2 = (1,1)$, and $\mathbf{k}$ is in the first Brillouin zone, which is given by the reciprocal lattice vectors $\mathbf{b}_1 = \frac{\pi}{2}(1,-1)$ and $\mathbf{b}_2 = 2\pi(0,1)$.

When the $\frac{J_x^2}{\omega}$ correction is included in the effective Hamiltonian, $\hat{H}_{\text{eff},1}$ (3.7), the unit cell is doubled, see Fig. 3.1(b), and the first Brillouin zone is therefore halved. If we now choose the lattice sites $\mathtt{a} = (1,0)$, $\mathtt{B} = (2,0)$, $\mathtt{c} = (3,0)$, $\mathtt{D} = (4,0)$, $\mathtt{A} = (2,1)$, $\mathtt{b} = (3,1)$, $\mathtt{C} = (4,1)$ and $\mathtt{d} = (5,1)$ for the unit cell, the momentum-space representation of the effective Hamiltonian will be an 8×8 matrix

$$\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k}) = \begin{pmatrix} -\frac{2J_x^2}{\omega} & J'_x e^{-i\frac{3\pi}{4}} & -\frac{J_x^2}{\omega}(1+e^{i\mathbf{k}\mathbf{R}_1}) & J'_x e^{-i(\frac{3\pi}{4}-\mathbf{k}\mathbf{R}_1)} & 0 & -J'_y e^{i\mathbf{k}\mathbf{R}_2} & 0 & -J'_y e^{i\mathbf{k}\mathbf{R}_1} \\ J'_x e^{i\frac{3\pi}{4}} & \frac{2J_x^2}{\omega} & J'_x e^{-i\frac{\pi}{4}} & \frac{J_x^2}{\omega}(1+e^{i\mathbf{k}\mathbf{R}_1}) & -J'_y & 0 & -J'_y e^{i\mathbf{k}\mathbf{R}_2} & 0 \\ -\frac{J_x^2}{\omega}(1+e^{-i\mathbf{k}\mathbf{R}_1}) & J'_x e^{i\frac{\pi}{4}} & -\frac{2J_x^2}{\omega} & J'_x e^{i\frac{\pi}{4}} & 0 & -J'_y & 0 & -J'_y e^{i\mathbf{k}\mathbf{R}_2} \\ J'_x e^{i(\frac{3\pi}{4}-\mathbf{k}\mathbf{R}_1)} & \frac{J_x^2}{\omega}(1+e^{-i\mathbf{k}\mathbf{R}_1}) & J'_x e^{-i\frac{\pi}{4}} & \frac{2J_x^2}{\omega} & -J'_y e^{-i\mathbf{k}(\mathbf{R}_1-\mathbf{R}_2)} & 0 & -J'_y & 0 \\ 0 & 0 & 0 & -J'_y e^{i\mathbf{k}(\mathbf{R}_1-\mathbf{R}_2)} & \frac{2J_x^2}{\omega} & J'_x e^{-i\frac{3\pi}{4}} & \frac{J_x^2}{\omega}(1+e^{i\mathbf{k}\mathbf{R}_1}) & J'_x e^{-i(\frac{3\pi}{4}-\mathbf{k}\mathbf{R}_1)} \\ -J'_y e^{-i\mathbf{k}\mathbf{R}_2} & 0 & -J'_y & 0 & J'_x e^{i\frac{3\pi}{4}} & -\frac{2J_x^2}{\omega} & J'_x e^{-i\frac{\pi}{4}} & -\frac{J_x^2}{\omega}(1+e^{i\mathbf{k}\mathbf{R}_1}) \\ 0 & -J'_y e^{-i\mathbf{k}\mathbf{R}_2} & 0 & -J'_y & \frac{J_x^2}{\omega}(1+e^{-i\mathbf{k}\mathbf{R}_1}) & J'_x e^{i\frac{\pi}{4}} & \frac{2J_x^2}{\omega} & J'_x e^{i\frac{\pi}{4}} \\ -J'_y e^{-i\mathbf{k}\mathbf{R}_1} & 0 & -J'_y e^{-i\mathbf{k}\mathbf{R}_2} & 0 & J'_x e^{i(\frac{3\pi}{4}-\mathbf{k}\mathbf{R}_1)} & -\frac{J_x^2}{\omega}(1+e^{-i\mathbf{k}\mathbf{R}_1}) & J'_x e^{-i\frac{\pi}{4}} & -\frac{2J_x^2}{\omega} \end{pmatrix}, \quad \text{(C.21)}$$

with the lattice vectors $\mathbf{R}_1 = (4,0)$ and $\mathbf{R}_2 = (2,2)$. The reciprocal lattice vectors are then $\mathbf{b}_1 = \frac{\pi}{2}(1,-1)$ and $\mathbf{b}_2 = \pi(0,1)$.

The energy bands of $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$ are shown in Figs. 3.2 and C.1.

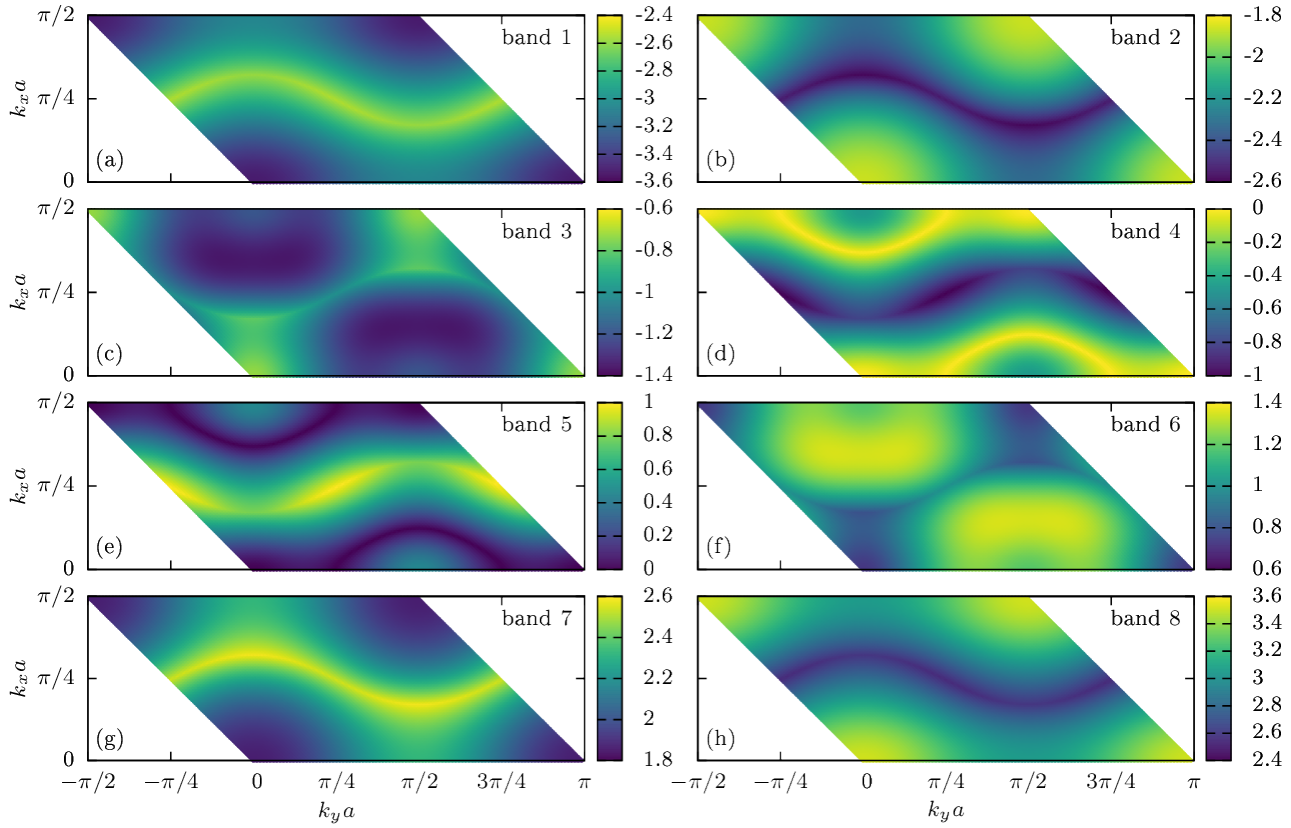

Figure C.1: Eight energy subbands of $\hat{\mathcal{H}}_{\text{eff},1}(\mathbf{k})$ for the driving frequency $\omega = 20$. Subbands 1 and 2 form the lowest band with Chern number $c_1 = 1$, subbands 3, 4, 5, and 6 form the middle band with $c_2 = -2$, and subbands 7 and 8 form the highest band with $c_3 = 1$.

Appendix D

Description of incoherent bosons

Here we explain in more details the method which we use to treat weakly interacting bosons in Chapter 3. In a typical condensed-matter system constituent particles are electrons. Due to their fermionic statistics, at low enough temperatures, and with Fermi energy above the lowest band, that band of the topological model is uniformly occupied, and consequently the transverse Hall conductivity can be expressed in terms of the Chern number (1.5) [69]. In contrast, weakly interacting bosons in equilibrium form a Bose-Einstein condensate in the band minima and only probe the local Berry curvature [167].

Yet in the experiment [70] the Chern number was successfully measured using bosonic atoms of ^{87}Rb. This was possible because in the process of ramping up the drive (3.4), the initial Bose-Einstein condensate was transferred into an incoherent bosonic mixture. Conveniently, it turned out that the bosonic distribution over the states of the lowest band of the effective Floquet Hamiltonian was nearly uniform. Motivated by the experimental procedure, we model the initial bosonic state by a statistical matrix

$$\rho(t=0) = \prod_{k=1}^{N_m} |k, N_p\rangle\langle k, N_p| \tag{D.1}$$

where the states $|k\rangle = a_k^\dagger|0\rangle$ approximately correspond to the lowest-band eigenstates of $\hat{H}_{\text{eff}}$ and each of these N_m states is occupied by N_p atoms $|k, N_p\rangle = \mathcal{N}(a_k^\dagger)^{N_p}|0\rangle$.

A procedure for selecting the states $|k\rangle$ is described in Refs. [70, 168]. In order to probe the Chern number of the lowest band, the states $|k\rangle$ should correspond closely to the lowest-band eigenstates of $\hat{H}_{\text{eff}}$. At the same time, in the experiment in the initial moment the atomic cloud is spatially localized. According to Refs. [70, 168] the optimal approach is to consider a steep confining potential and to use the low-lying eigenstates of

$$\hat{H}_{\text{initial}} = \hat{h}^{\text{eff}} + \left(\frac{r}{r_0}\right)^\zeta, \tag{D.2}$$

where in our calculations $\hat{h}^{\text{eff}}$ is either $\hat{H}_{\text{eff},0}$ from Eq. (3.6) or $\hat{H}_{\text{eff},1}$ from Eq. (3.7) and the parameters of the confining potential are set to $r_0 = 20, \zeta = 20$.

The dynamics of the initial state (D.1) is induced by a double quench: at $t_0 = 0$ the atomic cloud is released from the confining potential and exposed to a uniform force of intensity F along the y direction. During the whole procedure the driving providing the laser-assisted tunneling, defined in Eq. (3.4), is running.

The main observables of interest are the center-of-mass position along x direction

$$x(t) = \Big\langle \sum_{l,m} l |\psi_{l,m}(t)|^2 \Big\rangle, \tag{D.3}$$

and the population of the ith band of the effective model

$$\eta_i(t) = \Big\langle \sum_{|k\rangle \in i\text{-th band}} \Big| \sum_{l,m} \alpha^{k*}_{lm} \psi_{lm}(t) \Big|^2 \Big\rangle, \tag{D.4}$$

where the states $|k\rangle = \sum_{l,m} \alpha^k_{lm} |l,m\rangle$ correspond to the eigenstates of the effective model. Here, angle brackets $\langle\ \rangle$ denote averaging over N_{samples} sets of initial conditions.

In the case of non-interacting particles, these and other quantities can be numerically accessed by solving the single-particle time-dependent Schrödinger equation for N_m different initial states $|k\rangle$. This is equivalent to sampling the initial state according to Eq. (3.8).

In the end, we give two technical remarks. First, all our calculations are done in the rotating frame; see Eq. (C.1) in Appendix C. The staggered potential (3.2) is removed in this way. Second, in the case when the evolution is governed by the time-dependent Hamiltonian (3.9), the initial state is multiplied by the operator $\mathrm{e}^{-i\hat{K}(0)}$ in order to properly compare these results to the ones obtained from the evolution governed by the effective Hamiltonian (3.10); see Eq. (3.5).

Appendix

E

Initial quadratic regime

In this Appendix we explain the initial quadratic behavior exhibited by several observables in Chapter 3. For simplicity, we will consider only the case without the confining potential and with very weak force $F = 0.01$. The initial state is a Bose-Einstein condensate in one of the eigenstates of the effective Hamiltonian. The results are later averaged over all first band eigenstates.

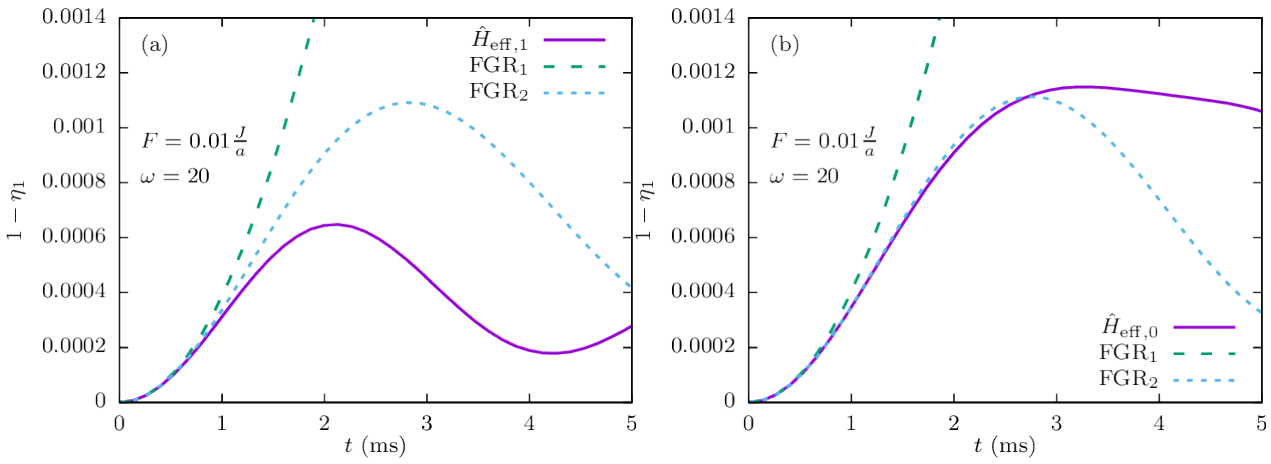

Figure E.1: Population in higher bands, comparison of numerical results (solid line) with the Fermi's golden rule in the first and second approximation (dashed lines). Band populations are calculated for an initial BEC in an eigenstate of the effective Hamiltonian and then averaged over (approximately) all states in the first band. (a) Initial state and evolution from the effective Hamiltonian with correction $\hat{H}_{\text{eff},1}$, Eq. (3.7). (b) Without the correction, $\hat{H}_{\text{eff},0}$, Eq. (3.6).

Fermi's golden rule predicts that the probability for transition from an initial state ψ_i to a final state ψ_f, induced by a perturbation $\Delta\hat{H}$, is proportional to the square of matrix elements $|\langle\psi_i|\Delta\hat{H}|\psi_f\rangle|^2$. In this case, the perturbation is $\Delta\hat{H} = F\hat{y}$. If we assume that the probability of a particle being in the initial state is always $P_i(t) = |\psi_i(t)|^2 \approx 1$, Fermi's golden rule predicts [236]

$$P_{i\to f}^{\text{FGR}_1}(t) = \frac{1}{\hbar^2}|\langle\psi_i|\Delta\hat{H}|\psi_f\rangle|^2 t^2. \tag{E.1}$$

If we now also consider transitions from the other states to the initial state, but keep the assumption that the populations in other states are small $P_{j\neq i}(t) = |\psi_{j\neq i}(t)|^2 \ll 1$, the time-dependent perturbation theory then predicts [236]

$$P_{i\to f}^{\mathrm{FGR}_2}(t) = |\langle i|\Delta\hat{H}|f\rangle|^2 \frac{1 - 2\mathrm{e}^{-\frac{\Gamma}{2\hbar}t}\cos\left(\frac{E_f - E_i}{\hbar}t\right) + \mathrm{e}^{-\frac{\Gamma}{\hbar}t}}{\left(E_f - E_i\right)^2 + \frac{\Gamma^2}{4}}, \tag{E.2}$$

where $\Gamma = \frac{2\pi}{\hbar}|\langle i|\Delta\hat{H}|f\rangle|^2$ and E_i (E_f) is the energy of the initial (final) state.

We plot the numerical results and both theoretical predictions from Fermi's golden rule in Fig. E.1. Here we can see that all three curves agree well for short times, the second approximation longer remains close to the numerical results, and that the initial quadratic regime is reproduced by theory. This is the so-called quantum Zeno regime [195].

Appendix F

Effects of interactions

Here we provide further details about the effects of interactions on the dynamics of weakly interacting incoherent bosons described in Chapter 3. In particular, we consider the evolution of the kinetic and interaction energy, as well as the probability density distribution in momentum space.

F.1 Energy

Time evolution of kinetic and interaction energy per particle for different interaction strengths is plotted in Fig. F.1. Here we define the kinetic energy per particle as the expectation value of the time-dependent Hamiltonian (C.1) divided by the total number of particles

$$E_{\mathrm{kin}}(t) = \frac{1}{N}\Big\langle \sum_{l,m,i,j} \psi^*_{l,m}(t) H_{lm,ij}(t) \psi_{i,j}(t) \Big\rangle, \tag{F.1}$$

while the interaction energy per particle is

$$E_{\mathrm{int}}(t) = \frac{1}{N}\frac{U}{2}\Big\langle \sum_{l,m} |\psi_{l,m}(t)|^2 \left(|\psi_{l,m}(t)|^2 - 1 \right) \Big\rangle. \tag{F.2}$$

Both energies grow with increasing interaction coefficient U.

When the interactions are strong enough and after long enough time, the atoms become equally distributed between the eigenstates of the Hamiltonian $\hat{H}(t)$. As the energy spectrum of $\hat{H}(t)$ is symmetric around zero, the expectation value of $\hat{H}(t)$ (kinetic energy) should be zero when all bands are equally populated. We can see this in Fig. F.1(a), where the kinetic energy approaches zero at $t \approx 50$ ms for the case $U = 0.05$.

The interaction energy at first rapidly decreases, as the cloud rapidly expands after turning off the confinement potential $\hat{V}_{\mathrm{conf}}$, and after that continues to slowly decrease as the cloud

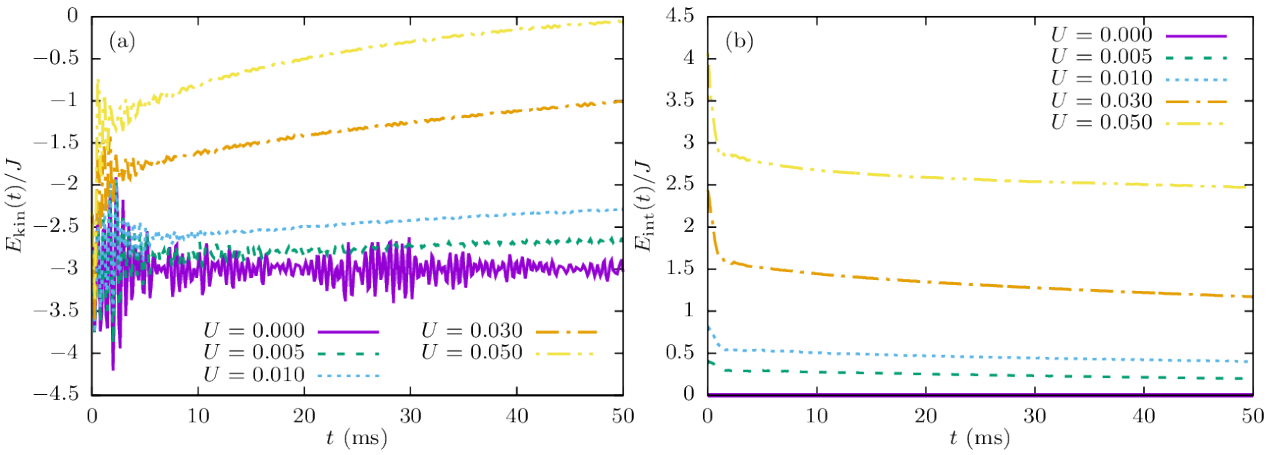

Figure F.1: (a) Kinetic energy per particle (expectation value of the time-dependent Hamiltonian $E_{\mathrm{kin}}(t) = \frac{1}{N}\left\langle \sum_{l,m,i,j} \psi^*_{l,m}(t) H_{lm,ij}(t) \psi_{i,j}(t) \right\rangle$ divided by the total number of particles N) for several different interaction strengths. (b) Interaction energy per particle $E_{\mathrm{int}}(t) = \frac{1}{N}\frac{U}{2}\left\langle \sum_{l,m} |\psi_{l,m}(t)|^2 \left(|\psi_{l,m}(t)|^2 - 1\right) \right\rangle$. U is given in units where $J = 1$.

slowly expands; see Fig. F.1(b).

These considerations also provide a possibility to discuss the applicability of the approximative method introduced in Section 3.3. As we work in the regime of high frequency $\omega = 20$, we find that for weak interaction, at short enough times of propagation, the energy is approximately conserved. At stronger values of $U \geq 0.01$ we observe a slow increase in the total energy on the considered time scales. In both cases we do not find the onset of parametric instabilities [174]. If present, these instabilities are signaled by an order of magnitude increase in energy on a short time scale, that we do not find.

In addition, the two-body interaction can deplete the occupancies of initial coherent modes [172, 184] and limit the validity of our approach. In principle, these types of processes can be addressed by including quantum fluctuations along the lines of the full truncated Wigner approach [187]. Yet, we set our parameters in such a way that these additional contributions are small.

F.2 Momentum-space density distribution

The momentum-space probability densities at the initial moment and after 75 driving periods (50 ms) are shown in Fig. F.2. The interactions deplete the lowest band, but also smooth out the density distribution.

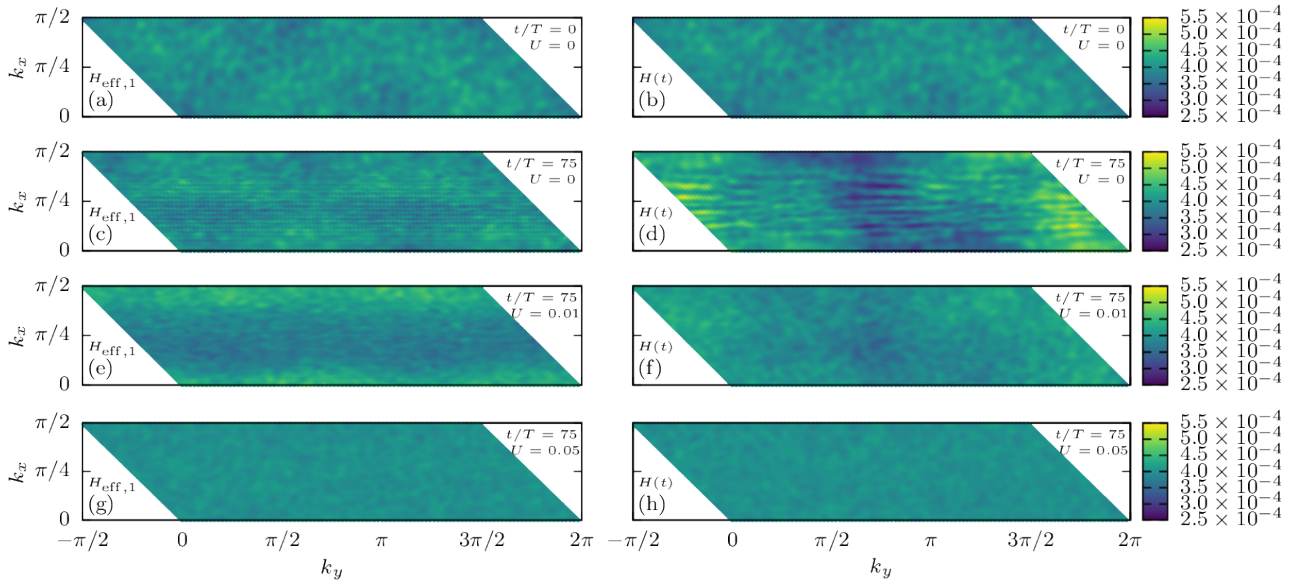

Figure F.2: Momentum-space density distribution in all bands, $\eta_1(\mathbf{k}) + \eta_2(\mathbf{k}) + \eta_3(\mathbf{k})$. U is given in units where $J = 1$. Left: evolution using the time-dependent Hamiltonian $\hat{H}_{\mathrm{eff},1}$. Right: evolution using the time-dependent Hamiltonian $\hat{H}(t)$. (a), (b) Initial state. (c), (d) Final state after 50 ms (75 driving periods), noninteracting case $U = 0$. (e), (f) $U = 0.01$. (g), (h) $U = 0.05$.

Appendix G

The effective model from Chapter 4

In this appendix we review the derivation of the model given in Eq. (4.1) from Chapter 4. The system is described by

$$\hat{H}_{\text{lab}}(t) = \hat{H}_{BH} + \hat{H}_{\text{drive}}(t) + \omega\,\hat{V}, \tag{G.1}$$

where we start with the Bose-Hubbard model

$$\begin{aligned}\hat{H}_{BH} = &- J_x \sum_{m,n} \left(\hat{a}^\dagger_{m+1,n}\hat{a}_{m,n} + \text{H. c.}\right) - J_y \sum_{m,n} \left(\hat{a}^\dagger_{m,n+1}\hat{a}_{m,n} + \text{H. c.}\right) \\ &+ \frac{U}{2}\sum_{m,n} \hat{n}_{m,n}(\hat{n}_{m,n} - 1),\end{aligned} \tag{G.2}$$

and we introduce an offset $\omega\hat{V}$:

$$\hat{V} = \sum_{m,n} n\,\hat{n}_{m,n}. \tag{G.3}$$

This shifted Bose-Hubbard model is exposed to a suitable resonant driving scheme:

$$\hat{H}_{\text{drive}}(t) = \frac{\kappa}{2}\sum_{m,n} \sin\left(\omega t - \phi_{m,n} + \frac{\phi}{2}\right)\hat{n}_{m,n}, \quad \phi_{m,n} = (m+n)\,\phi. \tag{G.4}$$

We assume periodic boundary conditions compatible with the driving term (G.4) in the laboratory frame. To this purpose we use vectors $\mathbf{R}_1 = 4\,\mathbf{e}_x$ and $\mathbf{R}_2 = -\mathbf{e}_x + \mathbf{e}_y$ as presented in Fig. 4.1. For simplicity, we work in the rotating frame

$$|\psi_{\text{rot}}(t)\rangle = e^{i\omega t\hat{V}}|\psi_{\text{lab}}(t)\rangle \tag{G.5}$$

and derive the Schrödinger equation

$$i\frac{d|\psi_{\text{rot}}(t)\rangle}{dt} = \hat{H}_{\text{rot}}(t)|\psi_{\text{rot}}(t)\rangle, \tag{G.6}$$

where

$$\hat{H}_{\mathrm{rot}}(t) = \left(e^{i\omega t\hat{V}}\hat{H}_{\mathrm{lab}}(t)e^{-i\omega t\hat{V}} - \omega\hat{V}\right). \tag{G.7}$$

Now we calculate $\hat{H}_{\mathrm{rot}}(t)$ explicitly. The only nontrivial action of this rotation on $\hat{H}_{\mathrm{lab}}$ comes from the nearest-neighbor hopping along y direction. Indeed, we have

$$e^{i\omega t\hat{V}}\hat{a}^{\dagger}_{m,n}\hat{a}_{m,n'}e^{-i\omega t\hat{V}} = e^{i\omega t(n-n')}\hat{a}^{\dagger}_{m,n}\hat{a}_{m,n'}. \tag{G.8}$$

In total we obtain

$$\begin{aligned}\hat{H}_{\mathrm{rot}}(t) &= -J_x\sum_{m,n}\left(\hat{a}^{\dagger}_{m+1,n}\hat{a}_{m,n} + \text{H. c. }\right) + \frac{U}{2}\sum_{m,n}\hat{n}_{m,n}(\hat{n}_{m,n}-1)\\ &+ e^{i\omega t}\hat{H}_1 + e^{-i\omega t}\hat{H}_{-1} + e^{-i\omega t(L_y-1)}\hat{H}_{L_y-1} + e^{i\omega t(L_y-1)}\hat{H}_{-L_y+1},\end{aligned} \tag{G.9}$$

with

$$\hat{H}_1 = -J_y\sum_{m,n}^{\mathrm{OBC}}\left(\hat{a}^{\dagger}_{m,n+1}\hat{a}_{m,n} - \frac{i}{4}\kappa e^{i(-\phi_{m,n}+\frac{\phi}{2})}\hat{n}_{m,n}\right), \quad \hat{H}_{-1} = \hat{H}_1^{\dagger}, \tag{G.10}$$

$$\hat{H}_{-L_y+1} = -J_y\sum_{m}\hat{a}^{\dagger}_{m,0}\hat{a}_{m-L_y,L_y-1}, \quad \hat{H}_{L_y-1} = \hat{H}^{\dagger}_{-L_y+1}. \tag{G.11}$$

In the terms $\hat{H}_{-L_y+1}$ and $\hat{H}_{L_y-1}$ we take into account periodic boundary conditions along the direction parallel to $\mathbf{R_2}$ as imposed in the laboratory frame. In order to limit the complexity of the numerical calculation, we keep translational invariance and impose the periodic boundary conditions in both directions in the rotating frame. This implies that we will neglect "phasors" $e^{-i\omega t(L_y-1)}$ and $e^{i\omega t(L_y-1)}$. Under these assumptions, we can recast Eq. (G.9) into the time-dependent Hamiltonian given in Eq. (4.1). In practice, this would require engineering additional non-trivial terms in the laboratory frame.

The leading order of the kick operator is given by

$$\hat{K}(t=0) \approx -\frac{\kappa}{2\omega}\sum_{m,n}\cos(\phi_{m,n}-\phi/2)\hat{n}_{m,n}. \tag{G.12}$$

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Biography of the author

Ana Hudomal was born on March 8th, 1991 in Belgrade, Republic of Serbia. She completed elementary and high-school education in Belgrade. In 2010 she started BSc studies in Theoretical and Experimental Physics at the Faculty of Physics, University of Belgrade and graduated in July 2014 (GPA 10.0/10.0). She then continued with MSc studies, also at the Faculty of Physics, University of Belgrade, and defended her MSc thesis in October 2015 (GPA 10.0/10.0). The title of her MSc thesis was *"New Periodic Solutions to the Three-Body Problem and Gravitational Waves"* and it was supervised by Dr. Veljko Dmitrašinović, Research Professor at the Institute of Physics Belgrade. During her studies, she received scholarships from the National Foundation for Development of Science and Art Youth (2007-2011), the City of Belgrade (2011-2013), and the Fund for Young Talents of the Republic of Serbia (2013-2015).

In November 2015 Ana Hudomal started PhD studies in the field of condensed matter physics at the Faculty of Physics, University of Belgrade, under the supervision of Dr. Ivana Vasić, Associate Research Professor at the Institute of Physics Belgrade. Since March 2016, Ana Hudomal is employed at the Institute of Physics Belgrade as a Research Assistant at the Scientific Computing Laboratory of the National Center of Excellence for the Study of Complex Systems. During this time she participated in several research projects, including the national research project ON171017 *"Modeling and Numerical Simulations of Complex Many-Particle Systems"* funded by the Ministry of Education, Science, and Technological Development of the Republic of Serbia, and multiple bilateral projects.

Ana Hudomal has so far published five papers in peer-reviewed international journals:

1. **A. Hudomal**, I. Vasić, N. Regnault, and Z. Papić, *"Quantum scars of bosons with correlated hopping"*, Commun. Phys. **3**, 99 (2020).
2. **A. Hudomal**, N. Regnault, and I. Vasić, *"Bosonic fractional quantum Hall states in driven optical lattices"*, Phys. Rev. A **100**, 053624 (2019).
3. **A. Hudomal**, I. Vasić, H. Buljan, W. Hofstetter, and A. Balaž, *"Dynamics of weakly interacting bosons in optical lattices with flux"*, Phys. Rev. A **98**, 053625 (2018).
4. V. Dmitrašinović, **A. Hudomal**, M. Shibayama, and A. Sugita,*"Linear stability of periodic three-body orbits with zero angular momentum and topological dependence of Kepler's third law: a numerical test"*, J. Phys. A: Math. Theor. **51**, 315101 (2018).
5. V. Dmitrašinović, M. Šuvakov, and **A. Hudomal**, *"Gravitational waves from periodic three-body systems"*, Phys. Rev. Lett. **113**, 101102 (2014).

Three of these papers (1-3) are directly related to the research presented in this thesis. The other two papers (4-5) are the result of the research conducted during her MSc studies. Ana Hudomal has also presented her research through one talk and several poster presentations at international schools, workshops and conferences.

Изјава о ауторству

Име и презиме аутора – **Ана Худомал**

Број индекса – **2015/8006**

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да је докторска дисертација под насловом

Numerical study of quantum gases in optical lattices and in synthetic magnetic fields (Нумеричко проучавање квантних гасова у оптичким решеткама и у синтетичким магнетним пољима)

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Изјављујем да је штампана верзија мог докторског рада истоветна електронској верзији коју сам предао/ла ради похрањивања у **Дигиталном репозиторијуму Универзитета у Београду.**

Дозвољавам да се објаве моји лични подаци везани за добијање академског назива доктора наука, као што су име и презиме, година и место рођења и датум одбране рада.

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Изјава о коришћењу

Овлашћујем Универзитетску библиотеку „Светозар Марковић“ да у Дигитални репозиторијум Универзитета у Београду унесе моју докторску дисертацију под насловом:

Numerical study of quantum gases in optical lattices and in synthetic magnetic fields (Нумеричко проучавање квантних гасова у оптичким решеткама и у синтетичким магнетним пољима)

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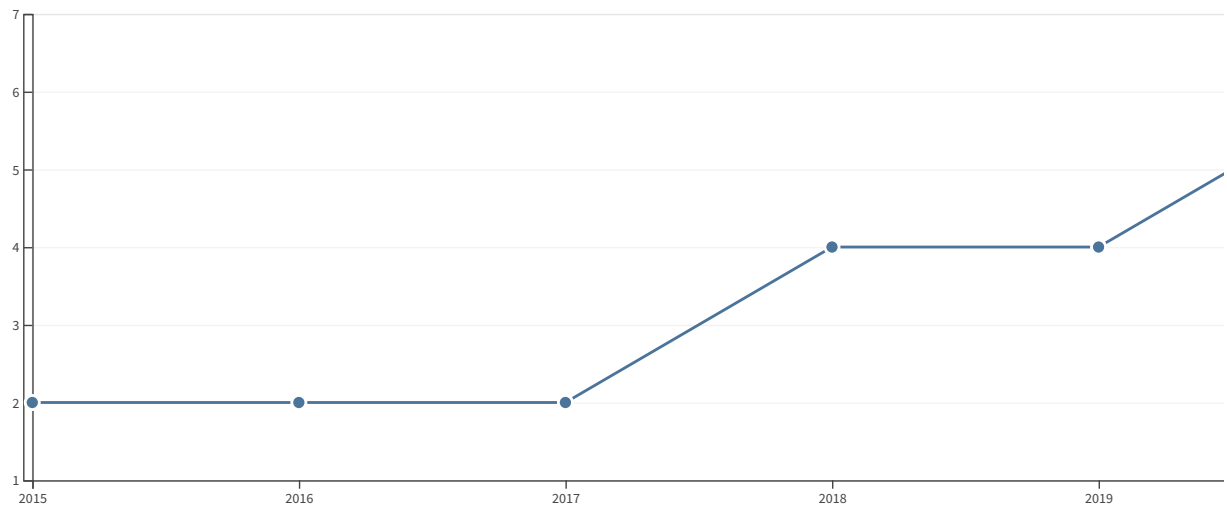

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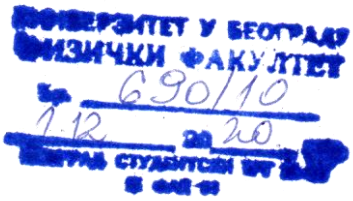

На основу члана 29 Закона о општем управном поступку («Службени гласник РС» број 18/2016 и 95/2018), и члана 149 Статута Универзитета у Београду - Физичког факултета, по захтеву АНЕ ХУДОМАЛ, мастер физичара, издаје се следеће

У В Е Р Е Њ Е

АНА ХУДОМАЛ, мастер физичар, дана 1. децембра 2020. године, одбранила је докторску дисертацију под називом

"NUMERICAL STUDY OF QUANTUM GASES IN OPTICAL LATTICES AND IN SYNTHETIC MAGNETIC FIELDS“ (Нумеричко проучавање квантних гасова у оптичким решеткама и у синтетичким магнетним пољима)

пред Комисијом Универзитета у Београду - Физичког факултета и тиме испунила све услове за промоцију у ДОКТОРА НАУКА – ФИЗИЧКЕ НАУКЕ.

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Проф. др Иван Белча