

# Representations and Particles of Orthosymplectic Supersymmetry Generalization<sup>1</sup>

Igor Salom

Institute of Physics, Belgrade

e-mail: isalom@ipb.ac.rs

**Abstract**—Orthosymplectic  $osp(1|2n)$  supersymmetry (alternative names: Generalized conformal supersymmetry with tensorial central charges, conformal M-algebra, parabose algebra) has been considered as an alternative to  $d$ -dimensional conformal superalgebra. Due to mathematical difficulties, even classification of its unitary irreducible representations (UIR's) have not been entirely accomplished. We give this classification for  $n = 4$  case (corresponding to four dimensional space-time) and then show how the discrete subset of these UIR's can be constructed in a Clifford algebra variation of the Green's ansatz.

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## 1. INTRODUCTION

Orthosymplectic type of space-time symmetry was first analyzed by C. Fronsdal [1], as early as in 1985, and since then interest for this symmetry reappeared, sometimes independently, in many contexts: M-theory [2], BPS particles [3], higher spin fields [4] and others [5].

When considering a (super)group in the context of a space-time symmetry, one of the first and most natural steps to undertake is to find unitary irreducible representations (UIR's) of the group, as these give us basic information on the particle content of the free theory. And in the case of orthosymplectic supersymmetry, physically most important class of UIR's are, so called, positive energy UIR's. The problem of finding these representations have been solved only for  $n = 1$  and  $n = 2$ . We followed the approach of [6] and completed the task for  $n = 4$  by using computer algorithms to analyze Verma module structure. In this way we managed to make a complete list of positive energy  $osp(1|8)$  UIR's, together with explicit forms of the corresponding Verma module singular and subsingular vectors. In this short report we present the main features of the results, leaving the details to be published separately. In particular, we point out that there is a concrete number of discrete UIR families (precisely nine, or ten if the trivial representation is counted as a separate class), that physically should be related to elementary particles of  $osp(1|8)$  models. In addition, we also point out a method to explicitly construct discrete representations, allowing one to easily perform concrete calculations in these spaces and, in that way, give physical interpretation to the states within. The method, directly generalizable to arbitrary  $n$ , is based

on a Clifford algebra variation of the Green's ansatz and is mathematically related to Howe duality.

## 2. POSITIVE ENERGY UIR's OF $osp(1|8)$

Structural relations of  $osp(1|2n)$  superalgebra can be compactly written in the form of trilinear relations of odd algebra operators  $a_\alpha$  and  $a_\alpha^\dagger$ :

$$[\{a_\alpha, a_\beta^\dagger\}, a_\gamma] = -2\delta_{\beta\gamma}a_\alpha, \quad (1)$$

$$[\{a_\alpha^\dagger, a_\beta\}, a_\gamma^\dagger] = 2\delta_{\beta\gamma}a_\alpha^\dagger,$$

$$[\{a_\alpha, a_\beta\}, a_\gamma], [\{a_\alpha^\dagger, a_\beta^\dagger\}, a_\gamma^\dagger] = 0, \quad (2)$$

where operators  $\{a_\alpha, a_\beta^\dagger\}$ ,  $\{a_\alpha, a_\beta\}$  and  $\{a_\alpha^\dagger, a_\beta^\dagger\}$  span the even part of the superalgebra and Greek indices take values 1, 2, ... $n$  (relations obtained from these by use of Jacobi identity are also implied). If we additionally require that the dagger symbol  $\dagger$  above denotes hermitian conjugation in the algebra representation Hilbert space (of positive definite metrics), then we have effectively constrained ourselves to the, so called, positive energy UIR's of  $osp(1|2n)$ .<sup>2</sup> Namely, in such a space, “conformal energy” operator  $E \equiv 1/2 \sum_\alpha \{a_\alpha, a_\alpha^\dagger\}$  must be a positive operator. Operators  $a_\alpha$  reduce the eigenvalue  $\epsilon$  of  $E$ , so the Hilbert space must contain a subspace that these operators annihilate. This subspace is called vacuum subspace:  $V_0 = \{|\nu\rangle, a_\alpha|\nu\rangle = 0\}$ . From the algebra relations follows:  $|\nu\rangle \in V_0 \Rightarrow \{a_\alpha, a_\beta^\dagger\}|\nu\rangle \in V_0$ , with  $\alpha, \beta$  arbitrary.

<sup>2</sup> Omitting a short proof, we note that in such a Hilbert space all superalgebra relations actually follow from one single relation—the first or the second of (1).

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Therefore vacuum subspace carries a representation of  $U(1) \times SU(N)$  group generated by operators  $\{a_\alpha, a_\beta^\dagger\}$  (with  $U(1)$  part generated by  $E$ ). The positive energy UIR's of  $osp(1|2n)$  are entirely labelled by UIR  $\mu$  of  $SU(N)$  (that can be given by a Young diagram) and a positive real number  $\epsilon_0$  (energy of the vacuum subspace) that labels  $U(1)$  representation. However, for a given representation  $\mu$  only certain values of  $\epsilon_0$  are allowed, and this connection is highly important and nontrivial.

In this paper we are interested in the  $n = 4$  case.  $SU(4)$  representation  $\mu$  will be explicitly parameterized by three non negative integers  $s_1, s_2, s_3$  in a way that  $\mu$  is determined by a Young diagram with  $s_1 + s_2 + s_3$  boxes in the first row,  $s_1 + s_2$  boxes in the second and  $s_1$  boxes in the third row. In addition to these three numbers, we will use real parameter  $d$  given by  $d = 1/4(\epsilon_0 - s_1 + s_3)$  to label  $osp(1|8)$  representations.

Classification of all positive energy UIR's was done by computer analysis of the Verma module structure, carried out in the following general manner (that we just briefly describe). First, Kac determinant of a sufficiently high level was considered as a function of parameter  $d$  (for each given class of  $SU(4)$  representation  $\mu$ ). In this way it was possible to locate the highest value of  $d$  for which the determinant vanishes and the Verma module becomes reducible. The singular or subsingular vector responsible for the singularity of the Kac matrix was then calculated, effectively by solving an (optimized) system of linear equations. Next we would find the norm of this vector and look for possible additional discrete reduction points at (lower) values of  $d$  for which the norm also vanishes. If new reduction points with new (sub)singular vectors were found it was also necessary to check that, upon removal of the corresponding submodules, no vectors with zero or negative norm remained. For this, it was enough to check that previously found (sub)singular vectors (i.e. those occurring for higher  $d$  values) belonged to the factored-out submodules. Optimized Wolfram Mathematica code was written to perform all these calculations. We now summarize the main results.

Parameter  $d$  can take the following values, depending on the labels  $s_1, s_2, s_3$ :

1.  $s_1 = s_2 = s_3 = 0$ :  $d > 3/2$  and singular points  $d = 0, 1/2, 1, 3/2$ ;
2.  $s_1 = s_2 = 0, s_3 > 0$ :  $d > s_3/2 + 2$  and singular points  $d = s_3/2 + 1, s_3/2 + 3/2, s_3/2 + 2$ ;
3.  $s_1 = 0, s_2 > 0$ :  $d > (s_2 + s_3)/2 + 5/2$  and singular points  $d = (s_2 + s_3)/2 + 2, (s_2 + s_3)/2 + 5/2$ ;
4.  $s_1 > 0$ :  $d > (s_1 + s_2 + s_3)/2 + 3$  and a singular point  $d = (s_1 + s_2 + s_3)/2 + 3$ .

Case 1 corresponds to “unique vacuum” representations, i.e. when the vacuum subspace  $V_0$  is one dimensional and carries trivial representation of the  $SU(4)$  group. Since spatial rotations are a subgroup of

this  $SU(4)$  group, in representations 1 the lowest conformal energy state is invariant to rotations. In this sense, representations 1 correspond to “fundamentally scalar” particles or, more precisely, multiplets (states of other spin values also belong to the multiplet of the full supersymmetry). Of particular physical interest are representations at singular points, and there are exactly three of such scalar representations. Namely, at singular points additional equations of motion appear directly related to the corresponding singular or subsingular vectors. In the case of the simplest nontrivial representation  $d = 1/2, s_1 = s_2 = s_3 = 0$ , singular vector yields the massless condition  $p^2 = 0$  (this is the well known and studied UIR containing tower of massless particles with increasing helicities). However, relating (sub)singular vectors to physical constraints (i.e. equations of motion) is in general complicated, and this problem is effectively solved by the explicit construction of representations that is discussed in the following section.

Case 2 corresponds to representations where  $V$  subspace transforms w.r.t.  $SU(4)$  subgroup as a Young diagram with  $s_3$  boxes in a single row. The simplest representative of the kind is a single box representation—lowest energy state in these representations behaves as a Lorentz  $1/2$  spinor. There are again three singular points in the case 2, corresponding to three physically interesting classes of representations, including three “fundamentally” spinor multiplets.

Cases 3 and 4 correspond to more complex classes of representations. However, it turns out that states from these classes can be naturally seen as composite states built from states belonging to representations of classes 2 or of classes 1.

Overall, there turns out to be 10 singular points corresponding to 9 different classes of nontrivial multiplets ( $d = 0$  point corresponds to the trivial representation).

### 3. CONSTRUCTION OF REPRESENTATIONS

It turns out that all representations with (half)integer  $d$  values (including all representations at (sub)singular points) can be obtained by representing the odd superalgebra operators  $a$  and  $a^\dagger$  as the following sums:

$$a_\alpha = \sum_{a=1}^p b_\alpha^a e^a, \quad a_\alpha^\dagger = \sum_{a=1}^p b_\alpha^{a\dagger} e^a. \quad (3)$$

In these expressions,  $p$  is integer,  $e^a$  are elements of a real Clifford algebra:

$$\{e^a, e^b\} = 2\delta^{ab}, \quad (a, b) = 1, 2, \dots, p, \quad (4)$$

and operators  $b_\alpha^a$  together with adjoint  $b_\alpha^{a\dagger}$  satisfy ordinary bosonic algebra relations:  $[b_\alpha^a, b_\beta^{b\dagger}] = \delta_{\beta\alpha}\delta^{ab}$ ,

$[b_\alpha^a, b_\beta^b] = 0$ . The (reducible) representation space is spanned by the vectors:

$$\mathcal{H} = l.s. \{ \mathcal{P}(b^\dagger) |0\rangle \otimes \omega \}, \quad (5)$$

where  $\mathcal{P}(b^\dagger)$  are monomials in mutually commutative operators  $b_\alpha^{a\dagger}$ ,  $|0\rangle$  is a bosonic vacuum and  $w \in \mathcal{H}_{CL}$ , where  $\mathcal{H}_{CL}$  is the representation space of real Clifford algebra (4).

Representation ansatz in the form (3) possesses certain intrinsic symmetries. Operators:

$$G^{ab} \equiv \sum_{\alpha=1}^n i(b_\alpha^{a\dagger} b_\alpha^b - b_\alpha^{b\dagger} b_\alpha^a) + \frac{i}{4} [e^a, e^b] \quad (6)$$

commute with entire  $osp(1|8)$  superalgebra. Operators  $G^{ab}$  themselves satisfy commutation relations of  $so(p)$  algebra (the full symmetry is actually slightly larger, given by the orthogonal group). This symmetry we will call the gauge symmetry.

The gauge symmetry actually removes all degeneracy in decomposition of (5) to  $osp(1|8)$  UIR's, i.e. the multiplicity of  $osp(1|8)$  UIR's is fully taken into account by labeling transformation properties of the vector w.r.t. the gauge symmetry group. Furthermore, there is one-to-one correspondence between UIR's of  $osp(1|8)$  and of the gauge group that appear in the decomposition, meaning that transformation properties under the gauge group action automatically fix the  $osp(1|8)$  representation.

The vector  $|\mathcal{V}_{\{d, s_1, s_2, s_3\}}^0\rangle$  that is the lowest weight vector of  $osp(1|8)$  positive energy UIR  $\{d, s_1, s_2, s_3\}$  and the highest weight vector of the gauge group UIR (in a standardly defined root system) takes the following explicit form (up to multiplicative constant):

$$\begin{aligned} |\mathcal{V}_{\{d, s_1, s_2, s_3\}}^0\rangle &= (B_{4+}^{(1)\dagger})^{s_3} (B_{4+}^{(1)\dagger} B_{3+}^{(2)\dagger} - B_{4+}^{(2)\dagger} B_{3+}^{(1)\dagger})^{s_2} \\ &\times \left( \sum_{k_1, k_2, k_3=1}^3 \varepsilon_{k_1 k_2 k_3} B_{4+}^{(k_1)\dagger} B_{3+}^{(k_2)\dagger} B_{2+}^{(k_3)\dagger} \right)^{s_1} \\ &\times \left( \sum_{k_1, k_2, k_3, k_4=1}^4 \varepsilon_{k_1 k_2 k_3 k_4} B_{4+}^{(k_1)\dagger} B_{3+}^{(k_2)\dagger} B_{2+}^{(k_3)\dagger} B_{1+}^{(k_4)\dagger} \right)^{s_0} |0\rangle \otimes \omega_{h.w.}, \end{aligned} \quad (7)$$

with  $d = s_0 + \frac{s_1 + s_2 + s_3 + p}{2}$ , where  $B_{\alpha\pm}^{(k)} =$

$\frac{1}{\sqrt{2}} (b_\alpha^{2k-1} \mp i b_\alpha^{2k})$  and  $(e^{2k-1} + i e^{2k}) \omega_{h.w.} = 0$ ,  $k = 0, 1, \dots [p/2]$ . The form above assumes that  $p$  is large enough that all  $B_{\alpha\pm}^{(k)}$  can be defined, i.e.  $p \geq 8$ :  $s_0$  must be 0 when  $p < 8$ ,  $s_1$  must be 0 when  $p < 6$ ,  $s_2$  must be 0 when  $p < 4$  and all  $s_0, s_1, s_2, s_3$  must be 0 when  $p = 1$  ( $p = 0$  is trivial UIR of  $osp(1|8)$ ).

In this way all positive energy UIR-s of  $osp(1|8)$  classified in the previous section with integer or half-integer values of  $d$  can be constructed using ansatz (3) with  $p \leq 9$ . Physically, corresponding states have natural interpretation as particles composed from  $p$  of the simplest  $osp(1|8)$  massless particles (belonging to  $d = 1/2, s_1 = s_2 = s_3 = 0$  UIR).

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