Singular BPS Saturated States and Enhanced Symmetries of Four-Dimensional N=4 Supersymmetric String Vacua

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Abstract

A class of supersymmetric (BPS saturated), static, spherically symmetric solutions of four-dimensional effective $N = 4$ supersymmetric superstring vacua, which become massless at special points of moduli space, is studied in terms of the fields of the effective heterotic string theory compactified on a six-torus. Those are singular four-dimensional solutions corresponding to $O(6,22,Z)$ orbits of dyonic configurations (with zero axion), whose left-moving as well as right-moving electric and magnetic charges are orthogonal (light-like in the $O(6,22,Z)$ sense), while the $O(6,22,Z)$ norms of both the electric and magnetic charges are negative. Purely electric [or purely magnetic] and dyonic configurations preserve $\frac{1}{2}$ and $\frac{1}{4}$ of $N = 4$ supersymmetry, respectively, thus belonging to the vector and the highest spin $\frac{3}{2}$ supermultiplets, respectively. Purely electric [or purely magnetic] solutions (along with an infinite tower of $SL(2,Z)$ transformed states) become massless at a point of the corresponding “one-torus”, thus may contribute to the enhancement of non-Abelian gauge symmetry, while dyonic solutions become simultaneously massless at a point of the corresponding two-torus, and thus may in addition contribute to the enhancement of the local supersymmetry there.

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Recently, it was recognized \[1,2\] that non-trivial configurations, e.g., Bogomol’nyi-Prasad-Sommerfield (BPS) saturated states, play a crucial role in addressing the full non-perturbative dynamics of string theory. When such configurations become light they can affect the low energy dynamics of the string theory in a crucial way \[2,3\].

Within Type IIA string compactified on Calabi-Yau manifolds, it was pointed out by Strominger \[4\] and further studied in Refs. \[3,4\] that massless supersymmetric black holes play a crucial role in the full string theory dynamics at the conifold points of moduli space.

Within four-dimensional $N = 4$ superstring vacua, Hull and Townsend \[3,5\] show that at special points of moduli space massless BPS saturated states can occur, contributing to a phenomenon which is a generalization of the Halpern-Frenkel-Kač (HFK) mechanism; namely, at special points of moduli space along with the perturbative electrically charged massless string states, which enhance the gauge symmetry to the non-Abelian one, there are massless BPS saturated magnetic monopoles and a tower of $SL(2, \mathbb{Z})$ related BPS saturated dyons which contribute to the new phase of the enhanced non-Abelian gauge symmetry.

In somewhat parallel developments, Behrndt \[8\] found four-dimensional, electrically charged supersymmetric black holes, which can become massless at special points of moduli space. They correspond to exact string backgrounds which are obtained by dimensionally reducing supersymmetric plane-wave solutions of effective ten-dimensional heterotic string theory. These solutions were generalized to the corresponding multi-black holes by Kallosh \[9\], while the corresponding exact magnetic solutions and the physical implications of the whole class of these solutions were studied by Kallosh and Linde \[10\].

The purpose of the paper is to study properties of a general class of BPS saturated, spherically symmetric solutions of effective four-dimensional $N = 4$ supersymmetric string vacua, which become massless at special points of moduli space. We parameterize these solutions in terms of fields of the effective heterotic string theory compactified on a six-torus. BPS saturated solutions in the effective Type II string compactified on $K_3 \times T^2$ surface are related to those of the toroidally compactified heterotic string through field redefinitions, since both of the effective actions in four dimensions are the same. A set of BPS states in one theory may turn out to be identified with elementary string states of a dual theory, thus providing further evidence for the conjectured duality between the two theories, whose origin lies in the string-string duality conjecture \[11,2\] of the corresponding six-dimensional theories.

The effective field theory of massless bosonic fields for the heterotic string on a Narain torus \[13\] at a generic point of moduli space is obtained by compactifying the ten-dimensional $N = 1$ Maxwell/Einstein supergravity theory on a six-torus \[12,13\]. The effective four-dimensional action \[\mathbb{I}\] for massless bosonic fields consists of the graviton $g_{\mu\nu}$, 28 $U(1)$ gauge fields $A^{m}_{\mu} \equiv (A^{(1)}_{\mu} , A^{(2)}_{\mu} , A^{(3)}_{\mu})$, corresponding to the gauge fields of dimensionally reduced ten-dimensional metric (Kaluza-Klein sector), two-form fields, and Yang-Mills fields, respectively, and 134 scalar fields. The scalar fields consist of the dilaton $\phi$, which parameterizes

\[\text{1See Refs. \[14,13\] for notational conventions and the relationship of four-dimensional fields to the corresponding ten-dimensional ones. Also, since we are studying the semiclassical configurations of the effective action we are not addressing $\alpha'$ corrections.}\]
the strength of the string coupling, the axion field $\Psi$, which is obtained from the two-form field $B_{\mu\nu}$ through the duality transformation, and a symmetric $O(6,22)$ matrix $M$ of 132 scalar fields. The matrix $M$ consists of 21 internal metric $g_{mn}$ components, 15 pseudo-scalar fields $B_{mn}$, and 96 scalar fields $a^I_m$, which arise from the dimensionally reduced ten-dimensional metric, two-form field and Yang Mills fields, respectively. Here, $(\mu, \nu) = 0, \cdots, 3$, $(m, n) = 1, \cdots, 6$ and $I = 1, \cdots, 16$.

The classical four-dimensional effective action is invariant under the $O(6,22)$ transformations (T-duality) \cite{14,17}:

$$M \to \Omega M \Omega^T, \quad A^i_{\mu} \to \Omega_{ij} A^j_{\mu}, \quad g_{\mu\nu} \to \Omega_{\mu\rho} \Omega_{\nu\sigma} g_{\rho\sigma}, \quad \phi \to \phi.$$ (1)

Here, $\Omega \in O(6,22)$, i.e., $\Omega^T L \Omega = L$, where $L$ is an $O(6,22)$ invariant matrix. In addition, the corresponding equations of motion and Bianchi identities have the invariance under the $SL(2,R)$ transformations (S-duality) \cite{15,16}:

$$S \to S' = \frac{aS + b}{cS + d}, \quad M \to \hat{M}, \quad g_{\mu\nu} \to \hat{g}_{\mu\nu}, \quad \mathcal{F}^i_{\mu\nu} \to \hat{\mathcal{F}}^i_{\mu\nu} = (c\psi + d)\mathcal{F}^i_{\mu\nu} + ce^{-\phi}(ML)_{ij}\hat{F}^j_{\mu\nu},$$ (2)

where $S \equiv \Psi + ie^{-\phi}$, $\hat{F}^i_{\mu\nu} = \frac{1}{2}(\sqrt{-g})^{-1}\varepsilon^{\mu\rho\sigma\nu}\mathcal{F}^i_{\rho\sigma}$, and $a, b, c, d \in \mathbb{R}$ satisfy $ad - bc = 1$. By embedding an Abelian gauge group in a non-Abelian one, one can see that the instanton effect breaks the $SL(2,R)$ symmetry down to $SL(2,Z)$, referred to as $S$-duality. The allowed discrete electric $\hat{Q}$ and magnetic $\hat{P}$ charges are determined \cite{15} by $T$- and $S$-duality constraints of the toroidally compactified heterotic string \cite{13} and Dirac-Schwinger-Zwanziger-Witten (DSZW) quantization condition \cite{14,18}; both of the “lattice charge vectors” $\bar{\alpha} \equiv L\hat{P}$ and $\bar{\beta} \equiv e^{-\phi}\hat{M}_{\infty}^{-1}\hat{Q} - \Psi_{\infty}\bar{\beta}$, lie in an even, self-dual, Lorenzian lattice $\Lambda$ with signature $(6,22)$. The world-sheet instanton effects break $O(6,22, R)$ invariance of the effective action down to its discrete subgroup $O(6,22, Z)$ referred to as $T$-duality. $T$-duality is an exact string symmetry to all orders in string perturbation and is assumed to survive non-perturbative corrections.

For an arbitrary asymptotic value $M_{\infty}$ of the moduli fields, one can perform \cite{15} the following simultaneous $O(6,22,R)$ rotations on the lattice $\Lambda$ and the matrix $M_{\infty}$: $M_{\infty} \to \hat{M}_{\infty} = \Omega M_{\infty} \Omega^T$ and $\Lambda \to \hat{\Lambda} = L\Omega L \Lambda$ $(\Omega \in O(6,22,R))$, in such a way that $M_{\infty} = I_{28}$. The charge configuration is now described by a new lattice $\hat{\Lambda}$. The subset of $O(6,22,Z)$ transformations that preserves this new asymptotic value of $\hat{M}_{\infty}$ is $O(6,Z) \times O(22,Z)$. Additionally, one imposes $SL(2,R)$ transformation to bring $S_{\infty} \to \hat{S}_{\infty} = i$. The subset of $SL(2,Z)$ transformations that preserves the asymptotic value $\hat{S}_{\infty} = i$ is $SO(2,Z)$. And the charge configuration is described by a new lattice $\hat{\Lambda}$. After imposing a subset of $O(6,Z) \times O(22,Z)$ and $SO(2,Z)$ transformations to generate a new solution from the generating solution, one has to undo the above $O(6,22,R)$ and $SL(2,R)$ rotations in order to obtain a new configuration with an arbitrary asymptotic value of the moduli field specified by $M_{\infty}$ and an arbitrary asymptotic value $S_{\infty}$ of the complex scalar.

The spectrum of BPS saturated (supersymmetric), static, spherically symmetric configurations at generic points of moduli space of toroidally compactified heterotic string is both $O(6,22,Z)$ and $SL(2,Z)$ invariant. In Ref. \cite{19} we derived the explicit form of a general class of such configurations with 56 charges subject to one constraint. It corresponds to the states which can be obtained by $SL(2,Z)$ transformations on configurations with zero axion
and the most general allowed dyonic charges. The latter set of configurations corresponds to \(O(6, 22, Z)\) orbits of dyonic configurations (with zero axion), whose left-moving as well as right-moving electric and magnetic charges are orthogonal, i.e., light-like in the \(O(6, 22, Z)\) sense.

The BPS states with general charge configurations can be obtained from the generating solutions through the \(O(6, 22, Z)\) and the \(SL(2, Z)\) transformations. It turns out [19] that the generating solution corresponds to the one where from scalar fields only the diagonal components \(g_{mm}\) of the internal metric and the dilaton field \(\phi\) are turned on. The charge lattice can be decomposed into sublattices which are \(O(6, 22, Z)\) orbits, each of which having a fixed value of the \(O(6, 22, Z)\) norm for the charge vectors. All the BPS states related through the \(O(6, 22, Z)\) and \(SL(2, Z)\) transformations have the same thermal and space-time properties.

For the purpose of illustrating the symmetry enhancement at particular points of moduli space, we shall concentrate on the generating solution, and choose \(M_\infty\) and \(S_\infty\) to be diagonal and purely imaginary, respectively. Other solutions with more general charge configurations and different choices of \(M_\infty\) and \(S_\infty\) allow for a larger enhancement of gauge symmetry.

Now, we recapitulate the results of Ref. [19] for the explicit form of the generating solution along with the above choice of \(M_\infty\) and \(S_\infty\). It is parameterized by two magnetic and two electric charges with electric and magnetic charges arising from different \(U(1)\) groups; the two magnetic [electric] charges arise from the Kaluza-Klein sector gauge field \(A^{(1)}_\phi [A^{(1)}_t]\) and the corresponding two-form \(U(1)\) field \(A^{(2)}_{\phi m} [A^{(2)}_t]\). Here, \(m \neq n\). Without loss of generality, we choose the non-zero charges to be \(P^{(1)}_1, P^{(2)}_1, Q^{(1)}_2, Q^{(2)}_2\).

The upper \(\varepsilon_u\) and lower \(\varepsilon_\ell\) two-components of the four-component Killing spinors are subject to the constraints: \(\Gamma^1 \varepsilon_{u,\ell} = i \eta_P \varepsilon_{t,\ell}\) if \(P^{(1)}_1 \neq 0\) and/or \(P^{(2)}_1 \neq 0\), and \(\Gamma^2 \varepsilon_{u,\ell} = \mp \eta_Q \varepsilon_{t,\ell}\) if \(Q^{(1)}_2 \neq 0\) and/or \(Q^{(2)}_2 \neq 0\). Here, \(\eta_P\) and \(\eta_Q\) are \(\pm 1\) and \(\Gamma^{1,2}\) are the gamma matrices of the corresponding \(SO(6)\) Clifford algebra. Non-zero magnetic and electric charges each break \(\frac{1}{2}\) of the remaining supersymmetries. Thus, purely electric [or purely magnetic] configurations preserve \(\frac{1}{2}\), while dyonic solutions preserve \(\frac{1}{4}\), of \(N = 4\) supersymmetry in four dimensions. The first and the second sets of configurations fall into vector- and hyper-supervectormultiplets with highest spins 1 and \(\frac{3}{2}\), respectively.

With the static, spherically symmetric Ansatz for the four-dimensional space-time metric in the Einstein frame: \(g_{\mu\nu}dx^\mu dx^\nu = \lambda(r)dt^2 - \lambda^{-1}(r)dr^2 - R(r)(d\theta^2 + \sin^2\theta d\phi^2)\), and with the internal diagonal metric \(g_{mm} = \delta_{mn}g_{mm}\) as well as the dilaton \(\phi\) depending only on the radial coordinate \(r\), the explicit form for the solution is given by [19]:

\[
\lambda = r^2/[(r - \eta_P P^{(1)}_1)(r - \eta_P P^{(2)}_1)(r - \eta_Q Q^{(1)}_2)(r - \eta_Q Q^{(2)}_2)]^{\frac{1}{2}},
\]

\[
R = [(r - \eta_P P^{(1)}_1)(r - \eta_P P^{(2)}_1)(r - \eta_Q Q^{(1)}_2)(r - \eta_Q Q^{(2)}_2)]^{\frac{1}{2}},
\]

\[
e^{\phi} = e^{\phi_\infty} \left[\frac{(r - \eta_P P^{(1)}_1)(r - \eta_P P^{(2)}_1)}{(r - \eta_Q Q^{(1)}_2)(r - \eta_Q Q^{(2)}_2)}\right]^{\frac{1}{2}},
\]

\(^2\)Massive highest spin \(\frac{3}{2}\) multiplets were addressed in the context of supergravity in Refs. [20,21] and in the context of massive BPS supervectormultiplets by Kallosh [23].
Here, the radial coordinate is chosen so that the horizon is at \( r = 0 \). The properties of the generating solution depend only on the four screened charges \((P_1, P_2, Q_1, Q_2)\) for which

\[
M_{\text{BPS}} = e^{-\frac{\phi}{2g_2}} \left[ \frac{1}{2 g_{1\infty}} |\beta_1| + g_{1\infty} |\alpha_1| \right] + e^{-\frac{\phi}{2g_2}} \left[ \frac{1}{2 g_{2\infty}} |\alpha_2| + g_{2\infty} |\beta_2| \right].
\]

Note that at generic points of moduli space, in Ref. \cite{19} we studied regular BPS saturated solutions. In this case, the requirement of the absence of naked singularities restricts the relative signs of the two magnetic [and two electric] charges to be the same. For such solutions, \( r = 0 \) corresponds to the Reissner-Nordström-type horizon, null singularity, and naked singularity when four charges, three [or two] charges, and one charge are nonzero, respectively.

On the other hand, Eq. \cite{3} implies when the relative signs for the two magnetic [and/or two electric] charges are opposite the solutions are always singular. When both the magnetic and the electric charges have opposite relative signs, the curvature singularity takes place at \( r = r_{\text{sing}} \equiv \max \{ \min |P_1|, |P_2|, \min |Q_1|, |Q_2| \} > 0 \). In addition, these configurations, in turn, become massless at special points of moduli space; when the magnitudes of the two screened electric [and two magnetic] charges of the generating solution are equal, the corresponding BPS saturated solutions in this class have zero ADM mass for special values of \( M_{\infty} \).

A class of such purely electrically charged configurations is related to the solution, recently found by Behrndt \cite{8,23}, which was obtained by dimensionally reducing supersymmetric gravitational waves of the effective ten-dimensional heterotic string theory. Generalizations to the corresponding multi-black hole solutions and the corresponding exact (in \( \alpha' \) expansion) magnetic solutions were given by Kallosh \cite{8}, and by Kallosh and Linde \cite{10}, respectively. In the latter work, the physical properties of such configurations were further addressed; they repel massive particles.

Just as in the case of purely electric [or purely magnetic] singular solutions \cite{10}, a general class of such dyonic configurations, when the magnetic and/or the electric charges have opposite relative signs, again repel massive particles. In particular, the traversal time of

\[ g_{11} = g_{1\infty} \left( \frac{r - \eta_P P_1(2)}{r - \eta_P P_1(1)} \right), \quad g_{22} = g_{2\infty} \left( \frac{r - \eta_Q Q_2(1)}{r - \eta_Q Q_2(2)} \right), \quad g_{mm} = g_{mm\infty} \quad (m \neq 1, 2). \]
the geodesic motion for a test particle with energy \(E\), mass \(m\) and zero angular momentum along the radial coordinate \(r\), as measured by an asymptotic observer, can be written as

\[
t(r) = \int_{r_{\infty}}^{r} \frac{E \, dr}{\lambda(r) \sqrt{E^2 - m^2 \lambda(r)}}. \tag{5}
\]

The minimum radius that can be reached by a test particle corresponds to \(r_{\text{min}} > r_{\text{sing}}\) for which \(\lambda(r = r_{\text{min}}) = E^2 / m^2\). Note that \(r_{\text{sing}}\) corresponds to the space-time singularity, \(i.e., \lambda(r = r_{\text{sing}}) = \infty\) and \(R(r = r_{\text{sing}}) = 0\). On the other hand, classical massless particles with zero angular momentum do not feel the repulsive gravitational potential due to increasing \(\lambda(r)\), and they reach the space-time singularity in a finite time.

Purely electrically charged solutions have quantum numbers of elementary string states, thus indicating that they should be identified with elementary string states \[26\]. On the other hand, purely magnetically charged solutions are not in the string spectrum, and should therefore be viewed as solitonic configurations. The nature of their charges indicates that they correspond to a hybrid ("bound state") of a Kaluza-Klein monopole and an H-monopole \[27\], whose relative signs of the corresponding magnetic charges are opposite. Note, that these configurations are different in nature from the BPS H-monopoles \[28\].

Dyonic solutions cannot be identified with elementary string excitations either, since there are no dyonic elementary string states. We believe that these configurations should be included in the spectrum of states contributing to the full string dynamics. It would be important to find out whether or not such configurations could arise from higher dimensional soliton configurations as discussed extensively in Ref. \[5\].

In the weak coupling limit, \(i.e., \phi_{\infty} \to -\infty\), and at a generic point of moduli space purely magnetically charged and dyonic configurations decouple, while purely electrically

\[4\text{Note that for the general class of singular solutions studied here, } \lambda \geq 1 \text{ for } r \text{ small enough, while for the regular solutions, studied in Ref. } \[19\] \text{ (with the same relative signs for the corresponding charges), } \lambda \leq 1. \text{ Thus, for the latter type of configurations the particles are always attracted toward the singularity. In the case of only one non-zero charge, the regular solution has a naked singularity at } r = 0, \text{ i.e., } t(r = 0) \text{ is finite.}

\[5\text{Another example of gravitational backgrounds, which repel massive particles, are Type III supergravity walls } \[24\], \text{ which are static, planar configurations interpolating between specific } N = 1 \text{ supergravity vacua with negative cosmological constant (two anti-deSitter vacua). In this gravitational background, a massive particle cannot reach the boundary of the anti-deSitter space-time on one side of such a wall. On the other hand, a classical massless test particle with zero transverse momentum can reach this boundary. This is the reason that on the boundary of the anti-deSitter vacuum one has to specify the boundary conditions } \[25\]. \text{ Note, however, that the boundary of the anti-deSitter space is not singular.}

\[6\text{Dyonic configurations may also have higher dimensional embedding into generalised supersymmetric wave solutions. Note that a subset of regular dyonic solutions which correspond to configurations with dyonic Kaluza-Klein charges } (P_{1}^{(1)} \neq 0, \text{ and } Q_{2}^{(1)} \neq 0) \[29\] \text{ have an embedding into six-dimensional supersymmetric Brinkman plane wave solutions } \[30\].}
charged ones are light. Note also that for purely magnetic [electric] singular configurations the coupling becomes weak, \(i.e., \phi \to -\infty, [\text{strong}, i.e., \phi \to +\infty]\) near the naked singularity at \(r_{\text{sing}}\). On the other hand, for dyonic singular configurations, both of the corresponding charges opposite, the coupling at the naked singularity is either weak when \(\min[|P^{(1)}_{1\infty}|, |P^{(2)}_{1\infty}|] > \min[|Q^{(1)}_{2\infty}|, |Q^{(2)}_{2\infty}|]\) or strong when \(\min[|P^{(1)}_{1\infty}|, |P^{(2)}_{1\infty}|] < \min[|Q^{(1)}_{2\infty}|, |Q^{(2)}_{2\infty}|]\).

We would now like to turn to the role that such states may play in enhancing the symmetries of vacua at special points of moduli space. For the example of the singular solutions, presented above, the allowed quantized magnetic and electric charge lattice vectors, as determined \([3]\) by the symmetries of the toroidally compactified heterotic string \([3]\) and the DSZW quantization rule \([17,18]\), that could give rise to massless BH’s correspond to \(\beta^{(1)} = -\beta^{(2)} = \pm 1\) and \(\alpha^{(1)} = -\alpha^{(2)} = \pm 1\), respectively.

The singular solutions which preserve \(\frac{1}{2}\) of supersymmetries correspond to purely, say, electrically charged solutions with \(two\) choices for charge lattice vectors: \(\alpha^{(1)} = -\alpha^{(2)} = \pm 1\). These two sets of solutions become massless at the self-dual point of the “one-torus”, \(i.e., \) when \(g_{22\infty} = 1\). Each of them belongs to a massless vector super-multiplet, identified with the perturbative string states, and form, together with the \(U(1)^{(1)+(2)}\) gauge field \((A^{(1)}_{\mu 2} + A^{(2)}_{\mu 2})/\sqrt{2}\), the adjoint representation of the non-Abelian \(SU(2)^{(1)+(2)}\) gauge group, thus enhancing the gauge symmetry from \(U(1)^{(1)}_2 \times U(1)^{(2)}_2\) to \(U(1)^{(1)+(2)}_2 \times SU(2)^{(1)+(2)}_2\) at the self-dual point of a one-torus. In addition, there are purely magnetically charged monopole configurations with \(\beta^{(2)} = -\beta^{(2)} = \pm 1\), discussed above, which along with a tower of the corresponding \(SL(2, \mathbb{Z})\) related dyonic states become massless at this point as well, thus providing an explicit realization of the phenomenon, studied by Hull and Townsend \([3]\).

On the other hand, the singular solutions which preserve \(\frac{1}{4}\) of supersymmetries correspond to dyonic solutions with electric and magnetic charges associated with two different one-tori with four charge assignments \(\beta^{(1)} = -\beta^{(2)} = \pm 1\) and \(\alpha^{(1)} = -\alpha^{(2)} = \pm 1\). They become massless at the self-dual point of the corresponding two-torus, \(i.e., \) when \(g_{11\infty} = g_{22\infty} = 1\). Each of these configurations belongs to a massless highest spin \(\frac{3}{2}\) supermultiplet, which contains not only massless vector fields (which combine with the \(U(1)^{(1)+(2)}\) gauge field \((A^{(1)}_{\mu 1} + A^{(2)}_{\mu 1})/\sqrt{2}\) and the \(U(1)^{(1)+(2)}\) gauge field \((A^{(1)}_{\mu 2} + A^{(2)}_{\mu 2})/\sqrt{2}\) to form the adjoint representation of the non-Abelian \(SU(2)^{(1)+(2)}_1 \times SU(2)^{(1)+(2)}_2\) gauge group), but also four massless spin \(\frac{3}{2}\) states, which contribute to the enhancement of the local supersymmetry. Along with this set of dyonic states, there is a tower of \(SL(2, \mathbb{Z})\) related dyonic states, which also become massless; at the self-dual point of the corresponding two-torus one encounters a possibility of a new phase of the theory where, along with a tower of massless vectors fields, a tower of massless spin \(\frac{3}{2}\) fields appears as well.

We would, however, like to caution that at the points of moduli space with the enhanced

For purely electrically charged solutions, such an embedding has been given in Ref. \([23]\). See also Ref. \([31]\) and references therein.
symmetries the description of the effective field theory breaks down due to the appearance of an infinite tower of new massless fields, and thus the full string dynamics at such points is not understood, yet. As for configurations belonging to the massless highest spin $\frac{3}{2}$ supermultiplets, it may turn out that they correspond to unstable states which decay into configurations belonging to the vector supermultiplets. All of these issues await further investigation.

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