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SUPERSYMMETRIC SUPERMEMBRANE VACUA AND SINGLETONS

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ABSTRACT

The eleven-dimensional supermembrane exhibits a local world-volume Siegel symmetry provided the space-time background is a solution of $D = 11$ supergravity. If this background admits Killing spinors, the supermembrane action also admits rigid space-time supersymmetries. We show that in a physical gauge, these symmetries give rise to rigid world-volume supersymmetries and we derive criteria for the existence of surviving supersymmetries of membrane vacua. We then obtain new supermembrane solutions in which space-time is the product of four-dimensional anti-de Sitter space and a compact seven-dimensional internal space M_7 , and the membrane occupies the boundary of AdS at spatial infinity. Applying our criteria for vacuum supersymmetry, we find that for an M_7 with N Killing spinors, the vacuum has N surviving supersymmetries. The physical fluctuations can be interpreted as singletons of $O\text{Sp}(4/N)$.

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The recent discovery of super p-branes^{1),2)}, i.e. theories of p-dimensional extended objects displaying space-time supersymmetry, naturally prompts the question of their quantum consistency. A particularly attractive candidate is the supermembrane ($p = 2$) in eleven space-time dimensions which couples to the background fields of $D = 11$ supergravity. A major difficulty in the analysis of membranes, as opposed to strings, is their intrinsic non-linearity. One approach is therefore to find exact stable solutions of the membrane field equations and to study the quantization of small fluctuations about these backgrounds. In this way, one might hope to determine the mass spectrum and, in particular, the existence or otherwise of massless particles.

It seems that in order to obtain massless particles in the supermembrane theory, a necessary condition is that there be a cancellation of the Casimir energy between the bosons and the fermions as in the Green-Schwarz superstring. The most likely way of achieving this is through a mode-by-mode cancellation due to some Bose-Fermi symmetry. An example of this cancellation phenomenon was discussed in Ref. 3) where the background space-time was chosen to be $R^9 \times S^1 \times S^1$ and the membrane was stabilized by being stretched over the torus. The cancellation would not have come about without both the fermionic Siegel symmetry on the world volume and the space-time supersymmetry.

In fact, the origin of this Bose-Fermi symmetry can be made more precise. Let us suppose that we have chosen a given space-time background but have not yet specified the supermembrane configuration. Provided the space-time background is a solution of $D = 11$ supergravity, the three-dimensional supermembrane action is invariant under a Siegel symmetry local on the world-volume^{1),4)}. If, in addition, this space-time background admits N Killing spinors, then the action is also invariant under N rigid space-time supersymmetries. After going to a physical gauge, a subgroup of these fermionic symmetries survives which is nothing but a rigid $N \leq 8$ world-volume supersymmetry^{*)} which transforms the eight physical boson degrees of freedom and the eight physical fermionic degrees of freedom into one another. Thus the number of world-volume supersymmetries depends on the space-time holonomy group⁵⁾. This is reminiscent of the superstring in the NSR formulation where there are $N = 1, 2$, or 4 world-sheet supersymmetries according as the target space is Riemannian, Kähler or hyper-Kähler.

The vanishing Casimir energy is then a result of choosing a vacuum solution for the membrane in which at least one of these world-volume supersymmetries remains unbroken. The non-vanishing result found recently⁶⁾ for a rotating super-

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membrane in $D = 11$ Minkowski space can be understood in this way, since this configuration breaks all world-volume supersymmetries.

Although the vanishing of the Casimir energy is not guaranteed to yield massless particles, as the torus example of Ref. 3) illustrates, it does seem to be a necessary condition. Accordingly, the purpose of the present paper is to present new examples of supersymmetric membrane vacua. We shall consider the case where the eleven-dimensional space-time takes the form of a direct product of four-dimensional anti-de Sitter space (AdS) and a compact seven-dimensional Einstein space M_7 . As is well known, by taking the four-index field-strength tensor to be proportional to the Levi-Cevita tensor in AdS⁷⁾, such configurations can provide solutions of the $D = 11$ supergravity field equations. An important special case corresponds to the choice $M_7 = S^7$ which yields the maximum of $N = 8$ supersymmetries⁸⁾ and whose vacuum symmetry is $OSp(4/8)$.

We shall show that there are static solutions of the supermembrane equations of motion for which the three-dimensional world-volume of the membrane occupies the boundary of the four-dimensional AdS. Moreover, the fluctuations of the eight bosonic and eight fermionic co-ordinates about this background admit the interpretation of "singleton" representations of $OSp(4/N)$. We recall that singletons⁹⁾ are those curious representations of the AdS group $SO(2,3)$ which have no Poincaré analogue and no obvious field theory interpretation. In the case of $OSp(4/N)$, they form an ultrashort supersymmetry multiplet consisting of spin 0 bosons and spin $\frac{1}{2}$ fermions transforming in the spinor representations of $SO(N)$ ¹⁰⁾. From our point of view, an important property of singletons is that their action is most appropriately written not on AdS itself but on its three-dimensional boundary at spatial infinity⁹⁾. The possible relation between supermembranes and singletons was conjectured in Ref. 11).

The supermembrane action is¹⁾

$$S = \int d^3\xi \left[-\frac{1}{2} \sqrt{-h} h^{ij} E_i^a E_j^b \eta_{ab} + \frac{1}{2} \sqrt{-h} -\frac{1}{6} \varepsilon^{ijk} E_i^A E_j^B E_k^C B_{CBA} \right] \quad (1)$$

where $\xi^i (i = 0, 1, 2) = (\tau, \sigma, \rho)$ are the co-ordinates on the world-volume, with metric $h_{ij}(\xi)$. $E_i^A = \partial_i z^\Lambda E_\Lambda^A$, where E_Λ^A is the supervielbein, $z^\Lambda = (X^M, \theta)$ and $B_{CBA}(X^M, \theta)$ is

the super three-form potential appropriate to the eleven-dimensional superspace. $X^M(\xi)$ ($M = 0, 1, \dots, 10$) are the bosonic co-ordinates and $\Theta(\xi)$ is a 32-component Majorana spinor which constitutes the fermionic co-ordinates of eleven-dimensional superspace.

We shall seek vacuum states for which both the space-time gravitino $\Psi_M(X)$ and the fermionic co-ordinates $\Theta(\xi)$ are zero, in which case the action (1) reduces to

$$S = \int d^3\xi \left[-\frac{1}{2} \sqrt{-h} h^{ij} \partial_i X^M \partial_j X^N g_{MN}(X) + \frac{1}{2} \sqrt{-h} + \frac{1}{6} \varepsilon^{ijk} \partial_i X^M \partial_j X^N \partial_k X^P B_{MNP}(X) \right] \quad (2)$$

Varying with respect to h_{ij} yields the embedding equation

$$h_{ij} = \partial_i X^M \partial_j X^N g_{MN} \quad (3)$$

while varying with respect to X^M yields the equations of motion

$$\begin{aligned} \partial_i (\sqrt{-h} h^{ij} \partial_j X^N g_{MN}) + \frac{1}{2} \sqrt{-h} h^{ij} \partial_i X^N \partial_j X^P \partial_M g_{NP} \\ + \frac{1}{6} \varepsilon^{ijk} \partial_i X^N \partial_j X^P \partial_k X^Q F_{MNPQ} = 0 \end{aligned} \quad (4)$$

where $F_{MNPQ} \equiv 40 [M^B_{NPQ}]$.

Since both Ψ_M and Θ are set to zero, the criterion for surviving supersymmetry is that their vacuum expectation values remain zero under some appropriate combination of the Siegel and space-time supersymmetry transformations. These transformation rules, to linear order in fermions (which is as high as we need consider in order to investigate the supersymmetry of a purely bosonic background), are¹⁾

$$\delta \Psi_M = \tilde{D}_M \varepsilon(X) \quad (5)$$

$$\delta \Theta = (1 + \Gamma) \kappa(\xi) + \varepsilon(X) \quad (6)$$

where κ is the Siegel symmetry parameter and ε the space-time supersymmetry parameter. Here $\tilde{D}_M$ is the supercovariant derivative¹²⁾

$$\tilde{D}_M = D_M - \frac{1}{288} \left(\Gamma^{NPQR}{}_M - 8 \Gamma^{NPQ} \delta_M^R \right) F_{NPQR} \quad (7)$$

where $D_M = \partial_M + 1/4 \omega_M^{AB} \Gamma_{AB}$, and Γ is given by

$$\Gamma = \frac{1}{6\sqrt{-h}} \varepsilon^{ijk} \partial_i X^M \partial_j X^N \partial_k X^P \Gamma_{MNP} \quad (8)$$

As a consequence of the embedding equation (3), Γ satisfied $\Gamma^2 = 1$ ¹⁾. Since Γ is trace-free, this implies that $\frac{1}{2}(1 \pm \Gamma)$ are projection operators with 16 zero eigenvalues. Irrespective of the background, the Siegel symmetry may therefore be used to set 16 of the 32 components of Θ to zero. A convenient choice is to set

$$(1 + \Gamma)\Theta = 0 \quad (9)$$

Acting with (6), we see that around a purely bosonic background, (9) is preserved if $(1+\Gamma)\delta\Theta = 0$, i.e., combined κ and ε transformations for which

$$2(1+\Gamma)\kappa + (1+\Gamma)\varepsilon = 0 \quad (10)$$

which implies that $\delta\Theta = \frac{1}{2}(1-\Gamma)\varepsilon$. Thus a bosonic vacuum has residual supersymmetries corresponding to solutions of

$$\tilde{D}_M \varepsilon = 0, \quad \Gamma \varepsilon = \varepsilon \quad (11)$$

We shall now look for a specific solution in which we take the space-time metric to be the direct product of AdS with co-ordinates X^μ ($\mu = 0, 1, 2, 3$) and radius a and a compact Einstein space M_7 with co-ordinates y^m ($m = 1, \dots, 7$) whose metric $g_{mn}(y)$ satisfies $R_{mn} = 3a^2/2 g_{mn}$. We take the AdS line-element in co-ordinates $X^\mu = (t, \theta, \phi, r)$ to be

$$ds^2 = -(1+a^2r^2)dt^2 + r^2(d\theta^2 + \sin^2\theta d\phi^2) + (1+a^2r^2)^{-1}dr^2 \quad (12)$$

In the orthonormal basis $e^0 = (1+a^2r^2)^{\frac{1}{2}}dt$, $e^1 = rd\theta$, $e^2 = r\sin\theta d\phi$, $e^3 = (1+a^2r^2)^{-\frac{1}{2}}dr$, the non-vanishing components of the spin-connection are given by

$$\begin{aligned} \omega^{03} &= a^2 r dt & \omega^{13} &= (1+a^2r^2)^{\frac{1}{2}} d\theta & \omega^{23} &= (1+a^2r^2)^{\frac{1}{2}} \sin\theta d\phi \\ \omega^{12} &= -\cos\theta d\phi \end{aligned} \quad (13)$$

The four-index field strength $F_{\mu\nu\rho\sigma}$ is $3a\sqrt{-g} \varepsilon_{\mu\nu\rho\sigma}$, where g is the determinant of

the AdS metric $g_{\mu\nu}(X)$ and $\varepsilon_{0123} = +1$, which in these co-ordinates becomes $F_{0123} = 3ar^2 \sin\theta$. This can be derived from a three-index potential $A_{012} = -ar^3 \sin\theta$.

We now substitute (12) and (13) into (3) and (4) and look for solutions of the form

$$t = \tau, \quad \theta = \sigma, \quad \phi = \varrho \quad (14)$$

with r and y^m independent of τ , σ and ρ . The only non-trivial equation that one has to solve is the $M = 3$ component of equation (4) which reads, after multiplying^{*)} by $(1+a^2r^2)^{-\frac{1}{2}}$,

$$a^2 r^3 (1+a^2 r^2)^{-1} + 2r - 3ar^2 (1+a^2 r^2)^{-1/2} = 0 \quad (15)$$

In addition to the trivial solution $r = 0$, corresponding to a totally collapsed membrane, there is a second solution with $r = \infty$. In this case the world-volume of the membrane occupies the boundary of AdS at spatial infinity, which has the topology of either $S^1 \times S^2$ or $R \times S^2$, according as the time co-ordinate t is identified periodically or not (i.e., space-time is AdS or its covering space). It is this solution, the Membrane at the End of the Universe, on which we shall now concentrate. We begin by showing that this vacuum has the required $N < 8$ unbroken supersymmetries, where N is the number of Killing spinors on M_7 .

In order to show that the vacuum configuration (14) does indeed satisfy the requirements of unbroken supersymmetry (11), we first exhibit a spinor ε satisfying $\tilde{D}_M \varepsilon = 0$ everywhere and then show how it satisfies $\Gamma \varepsilon = \varepsilon$ on the boundary of AdS. We first make a four-seven split of the eleven-dimensional Γ^M matrices

$$\Gamma^\mu = \gamma^\mu \otimes \mathbb{1}, \quad \Gamma^m = \gamma_5 \otimes \Sigma^m \quad (16)$$

where γ_5 is the four-dimensional chirality operator given in an orthonormal frame by $\gamma_5 = -i\gamma_{0123}$. Our conventions are

*) Since (4) is a vector equation, our multiplication by $(1+a^2r^2)^{-\frac{1}{2}}$ can be achieved by changing to a new radial variable r' defined by $\partial r' / \partial r = (1+a^2r^2)^{\frac{1}{2}}$. Thus there is no co-ordinate invariant significance to the factor multiplying (4) and so our choice of scaling in (15), which was made so that $r = \infty$ should yield a regular solution, is justified.

$$\{\Gamma^M, \Gamma^N\} = 2 g^{MN} \quad (17)$$

and we record for future reference that $\gamma_5^2 = 1$ and $\varepsilon_{012} = 1$. We work with metric signature $(-, +, \dots, +)$ for both world-volume and space-time. From (8) and (14), we find

$$\Gamma = \gamma \otimes \mathbb{1} \quad , \quad \gamma \equiv -\gamma_{012} \quad (18)$$

and from (7)

$$\tilde{D}_\mu = D_\mu - \frac{ia}{2} \gamma_\mu \gamma_5 \quad (19)$$

$$\tilde{D}_m = D_m + \frac{ia}{4} \Sigma_m \quad (20)$$

where $D_\mu = \partial_\mu + 1/4 \omega_\mu^{\alpha\beta} \gamma_{\alpha\beta}$ and $D_m = \partial_m + 1/4 \omega_m^{ab} \Sigma_{ab}$ are the usual Lorentz covariant derivative on AdS and M_7 . Note that in (18) and in all subsequent equations, explicit numerical indices always denote tangent space indices. To solve $\tilde{D}_M \varepsilon = 0$ we follow Ref. 8) and write

$$\varepsilon(X, y) = \varepsilon(X) \otimes \eta(y) \quad (21)$$

where $\eta(y)$ is a Killing spinor on M_7 satisfying $\tilde{D}_m \eta = 0$ and $\varepsilon(X)$ is a Killing spinor on AdS satisfying $\tilde{D}_\mu \varepsilon = 0$. By methods analogous to Ref. 13), one can show that the general solution for ε is given by

$$\varepsilon = \sqrt{2} (\cos g)^{-1/2} (\cos g/2 + \gamma \sin g/2) \zeta(t, \theta, \phi) \quad (22)$$

where

$$\tan g = ar$$

and where ζ satisfies

$$\nabla_\alpha \zeta + \frac{a}{2} \gamma_\alpha \gamma_3 \zeta = 0 \quad , \quad \frac{\partial \zeta}{\partial t} - \frac{ia}{2} \gamma_0 \gamma_5 \zeta = 0 \quad (23)$$

where ∇_α is the covariant derivative on S^2 with metric

$$\tilde{a}^2 (d\theta^2 + \sin^2 \theta d\phi^2) \quad (24)$$

Note that (23) implies the three-covariant equation

$$(1 + \gamma) \left(\nabla_i + \frac{a}{2} \gamma_i \gamma_3 \right) \zeta = 0 \quad (25)$$

where $\nabla_i = (\partial/\partial t, \nabla_\alpha)$ is the covariant derivative on the boundary of AdS with metric $-dt^2 + a^{-2}[d\theta^2 + \sin^2\theta d\phi^2]$. To show that (23) has solutions, we make a three-one split of the γ_μ matrices

$$\gamma_i = \begin{pmatrix} 0 & i\tau_i \\ -i\tau_i & 0 \end{pmatrix}, \quad \gamma_3 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} \quad (26)$$

where $\tau_0 = i\sigma_3$, $\tau_1 = \sigma_1$, $\tau_2 = \sigma_2$ and σ_i are the Pauli matrices. We write

$$\zeta = \left(\exp \frac{ia}{2} \gamma_0 \gamma_5 \tau \right) \begin{pmatrix} u_+(\theta, \phi) \\ u_-(\theta, \phi) \end{pmatrix} \quad (27)$$

and substituting this into (23), we find

$$\nabla_\alpha u_\pm \pm \frac{ia}{2} \tau_\alpha u_\pm = 0 \quad (28)$$

This is just the Killing spinor equation on S^2 for which there are two solutions for each choice of sign. This implies the well-known result that AdS has four Killing spinors.

Since $r = \infty$ corresponds to $\rho = \pi/2$, it follows from (18), (21) and (22) that ϵ becomes an eigenstate of γ and hence $\Gamma\epsilon = \epsilon$ on the boundary of AdS. Thus we have shown that our vacuum is supersymmetric whenever M_7 admits Killing spinors. The number of such spinors, N , depends on the Weyl holonomy of M_7 and lies between 0 and 8⁸⁾.

Having established the criteria for unbroken vacuum supersymmetry and having demonstrated that the solution (14) satisfies these criteria, we now turn to the examination of small fluctuations about this vacuum. For the bosons, we write $X^M = \overset{\circ}{X}^M + \delta X^M$ where $\overset{\circ}{X}^M$ is the vacuum solution. We shall work in the gauge where $\delta X^i = 0$ ($i = 0, 1, 2$). From (4) we find

$$\left[\square - \frac{2}{r^2} \left(1 + \frac{3}{2} a^2 r^2 \right) - 4 a^2 (1 + a^2 r^2)^{-1} \left(1 + \frac{3}{4} a^2 r^2 \right) + \frac{6a}{r} (1 + a^2 r^2)^{1/2} \right] \delta r = 0 \quad (29)$$

and

$$\square \delta y^m = 0 \quad (30)$$

where $\square = (-h)^{-1/2} \partial_i (\sqrt{-h} h^{ij} \partial_j)$.

We find in the limit $r \rightarrow \infty$ that (29) reduces to $\square \delta r = 0$ and hence that all eight fluctuating physical degrees of freedom obey the same equation of motion. For the fermions, the equation of motion that follows from (1) is

$$\sqrt{-h} (1-\gamma) h^{ij} E_i^M (\Gamma_M)^\alpha{}_\beta E_j^\beta = 0. \quad (31)$$

To linear order in Θ , the supervielbein components may be found in Ref. 14). We thus obtain the linearized fermion field equation

$$\sqrt{-h} (1-\gamma) e_a^i \gamma^a \tilde{D}_i \Theta = 0 \quad (a = 0, 1, 2) \quad (32)$$

where

$$\tilde{D}_i = \partial_i + \partial_i X^M \left[\frac{1}{4} \omega_M^{AB} \Gamma_{AB} - \frac{1}{288} (\Gamma^{NPQR}{}_M - 8 \Gamma^{NPQ} \delta_M^R) F_{NPQR} \right] \quad (33)$$

In our background, $\tilde{D}_i$ reduces to

$$\begin{aligned} \tilde{D}_t &= \frac{\partial}{\partial t} + \frac{1}{2} \sec \rho \gamma_{03} (\sin \rho - \sigma), \\ \tilde{D}_\theta &= \frac{\partial}{\partial \theta} + \frac{1}{2} \sec \rho \gamma_{13} (1 - \gamma \sin \rho), \\ \tilde{D}_\varphi &= \frac{\partial}{\partial \varphi} - \frac{1}{2} \cos \theta \gamma_{12} + \frac{1}{2} \sin \theta \sec \rho \gamma_{23} (1 - \gamma \sin \rho). \end{aligned} \quad (34)$$

where $\tan \rho = ar$. Since our membrane solution arises in a singular limit of the equations, one has to be careful when interpreting the fluctuation equations on the boundary. We shall discuss the details of this procedure elsewhere. For now, we note that κ and ε symmetry of Eq. (32) on the boundary $r = \infty$ follow from (6) and the Killing spinor equation (11).

On using the Siegel gauge-condition $\Gamma \Theta = -\Theta$, (32) reduces to an equation describing eight physical fermionic degrees of freedom. This, and the equations for the eight bosonic degrees of freedom, are invariant under a certain combination of the original κ and ε symmetries; i.e., (6) together with

$$\delta X^M = i \bar{\Theta} \Gamma^M (1 + \Gamma) \kappa + i \bar{\varepsilon} \Gamma^M \Theta. \quad (35)$$

Specifically, this residual symmetry follows from (10), and taking the appropriate limit as $r \rightarrow \infty$ of

$$\delta\theta = \frac{1}{2}(1-r)\varepsilon, \quad (36)$$

$$\delta\chi^M = 2i\bar{\varepsilon}\Gamma^M\varepsilon. \quad (37)$$

Thus we have a supersymmetric theory with eight spin-0 and eight spin- $\frac{1}{2}$ degrees of freedom, defined on the boundary of AdS. As we showed earlier, the number N of supersymmetries is given by the number of Killing spinors on M_7 , which lies between 0 and 8. The case $N = 8$, which arises when M_7 is the round seven-sphere⁸⁾, is especially interesting; it gives rise to an $N = 8$ supersymmetric singleton theory on the boundary of AdS. (The fact that the $N = 8$ supermultiplet is ultra-short, comprising just spin-0 and spin- $\frac{1}{2}$ fields, demonstrates that it can only be a singleton representation of the AdS algebra.) In the case that M_7 admits $N < 8$ Killing spinors, one obtains multiple copies of the appropriate singleton theory, for which the spin-0 and spin- $\frac{1}{2}$ fields are in spinor representations of $SO(N)$. For instance, if $N = 1$ [which would arise, for example, if M_7 were the squashed seven-sphere⁵⁾], one obtains eight copies of the $N = 1$ singleton.

Finally, we recall that the singleton representations of $OSp(4/8)$ occur in the harmonic expansion of the $D = 11$ supergravity fields on S^7 but can be gauged away¹⁵⁾ everywhere (except on the boundary of AdS?)⁸⁾. This is suggestive¹¹⁾ of a form of bootstrap in which the field theory gives rise to the membrane which in turn yields the field theory. The physical interpretation of this Membrane at the End of the Universe and the issue of massless particles will be discussed elsewhere.

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