On gauge-fixed superbrane actions in AdS superbackgrounds

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Abstract

To construct actions for describing superbranes propagating in $AdS \times S$ superbackgrounds we propose a coset space realization of these superbackgrounds which results in a short polynomial fermionic dependence (up to the sixth power in Grassmann coordinates) of target superspace supervielbeins and superconnections. Gauge fixing $\kappa$–symmetry in a way compatible with a static brane solution further reduces the fermionic dependence down to the fourth power. Subtleties of consistent gauge fixing worldvolume diffeomorphisms and $\kappa$–symmetry of the superbrane actions are discussed.

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Worldvolume dynamics of super–p–branes (M-branes and D-branes) propagating in Anti–de–Sitter supergravity backgrounds has been understood to give rise to \((p + 1)\)-dimensional superconformal theories \([1]–[3]\). In connection with this observation an extensive study of superbrane worldvolume actions in these superbackgrounds has been carried out \([4]–[18]\).

To reduce the worldvolume action of the superbrane to that of the superconformal theory one should know an explicit form of the geometrical quantities which describe the AdS background (supervielbeins, superconnections and antisymmetric gauge superfields), to gauge fix worldvolume diffeomorphisms for bosonic target–space coordinates of the superbrane, and to eliminate half the fermionic target–space coordinates by gauge fixing kappa–symmetry of the superbrane worldvolume action.

These problems are mostly technical ones, and to arrive at their simplest solution one, for instance, can try to guess an appropriate realization of the AdS target superspace as a coset superspace \(K = G/H\), where \(G\) is the isometry supergroup of the background and \(H\) is its isotropy (or stability) subgroup. For these purposes the coset–space techniques has been developed in a number of papers \([6]–[9], [13]\). Another possibility is to directly solve supergravity constraints \([15]\). Such realizations produce a complicated form of the target superforms with the dependence on Grassmann coordinates \(\Theta\) up to a 32–nd power. Upon imposing a so called Killing spinor \([10]\) or supersolvable algebra \([9]\) condition for gauge fixing the \(\kappa\)–symmetry, a significant simplification of the structure of the superforms was achieved, so that only the second power in \(\Theta\) remains in the supervielbeins and superconnections. The choice of a gauge fixing condition is a subtle point in the sense that it must be compatible with the classical solution of brane equations of motion chosen as a vacuum over which a perturbative theory is developed. As we shall see the gauge fixing condition mentioned above is not always compatible with a static brane solution embedded into the AdS part of the target superspace. In these cases another \(\kappa\)–symmetry gauge fixing is required.

In this paper we propose a coset space realization of AdS superbackgrounds which results in a short polynomial fermionic dependence (up to the sixth power in \(\Theta\)) of the superforms even before fixing the \(\kappa\)-symmetry. Gauge fixing the \(\kappa\)-symmetry in a way consistent with the static brane solution further reduces \(\Theta\)-dependence down to the fourth power.

The results obtained are generic for the cases of a D3-brane in IIB D=10 supergravity, and an M2–brane and an M5–brane in D=11 supergravity, which are considered simultaneously. They should also be applicable to a superstring propagating in \(AdS^3 \times S^3\).

The general form of the superbrane action is (details can be found in \([19, 20, 21]\))

\[
S = I_0 + I_{WZ},
\]

where

\[
I_0 = -\int_{M_{p+1}} d\xi^{p+1} \sqrt{-detg_{ij}} + ...
\]
is a Nambu–Goto or Dirac–Born–Infeld–type part of the action (the dots stand for the contribution of worldvolume gauge fields, if present),

\[ g_{ij}(\xi) = \partial_i Z^M E^a_M E_{aN} \partial_j Z^N \quad (i, j = 0, 1, \ldots, p) \quad (3) \]

is an induced worldvolume metric and \( E^a_M(Z), E^\alpha_M(Z) \) are supervielbeins of target superspace parametrized by the coordinates \( Z^M = (X^\hat{m}, \Theta^\alpha) \).

\[ I_{WZ} = -\int_{\mathcal{M}_{p+1}} A^{(p+1)} + ... \quad (4) \]

is the Wess-Zumino part of the superbrane action whose leading term is the worldvolume integral over the pullback of a target–space \((p + 1)\)-form gauge superfield. It is convenient to take the following metric as the bosonic part of the \( AdS^{p+2} \times S^{D-p-2} \) superbackground metric

\[ ds^2 = \left( \frac{r}{R} \right)^{2(D-p-3)} \eta_{ij} dx^i dx^j + \left( \frac{R}{r} \right)^2 dr^2 + R^2 d\Omega^2, \quad (5) \]

where \( x^i, r \quad (i = 0, 1, \ldots, p) \) are coordinates of the AdS and \( d\Omega^2 \) stands for a metric of the sphere of a radius \( R \) parametrized by coordinates \( y^{a'} \quad (a' = 1', \ldots, (D-p-2)'). \)

The nonvanishing pure bosonic components of the gauge field \( A^{(p+1)} \) in the \( AdS \times S \) background are

\[ A^{(p+1)} = dx^p \ldots dx^0 \left( \frac{r}{R} \right)^{D-p-3} + O(\Theta). \quad (6) \]

The action (4) is invariant under the fermionic \( \kappa \)–symmetry transformations

\[ \delta_{\kappa} Z^M E^a_M = 0, \quad \delta_{\kappa} Z^M E^\alpha_M = \kappa^\beta(\xi)(1 + \bar{\Gamma})_{\beta}^\alpha, \quad (7) \]

(plus a variation of worldvolume gauge fields) where \( 1 + \bar{\Gamma} \) is a spinor projector. The form of \( \bar{\Gamma} \) is specific for each superbrane [19, 20, 21] but for every brane the leading term of \( \bar{\Gamma} \) is

\[ \bar{\Gamma} = \frac{1}{(p+1)!\sqrt{-g}} \varepsilon^{i_1 \ldots i_{p+1}} \Gamma_{i_1 \ldots i_{p+1}} + ... \quad (8) \]

where \( \Gamma_{i_1 \ldots i_{p+1}} \) is the antisymmetric product of target–space gamma–matrices pulled back on to the worldvolume, i.e. \( \Gamma_i = \partial_i Z^M E^a_M \Gamma_a \).

Note that a classical static “vacuum” solution of the superbrane equations of motion which follow from (1) is

\[ \xi^i = x^i, \quad r = const, \quad y^{a'} = const, \quad \Theta = 0, \quad (9) \]

and the worldvolume gauge fields are zero.

For this static solution the action (1) vanishes, which is called the no–force condition, since there is no a potential which would push the brane to a boundary of AdS [22, 3, 9]. This, in particular, means that the static gauge \( \xi^i = x^i \) for the worldvolume diffeomorphisms is compatible with the vacuum solution.
In the case of the static vacuum solution the $\kappa$–symmetry projector takes a simple form
\[
1 + \tilde{\Gamma} = 1 + \tilde{\gamma},
\]
where $\tilde{\gamma} = \Gamma^{01..p}$ in the case of the M2 and the M5–branes, and in the case of the IIB D3-brane $\tilde{\gamma} = \epsilon^{IJ} \Gamma^{01..3}$ (where $I, J = 1, 2$ number the two $D=10$ Majorana–Weyl spinors).

Note that if we chose another sign in front of the Wess–Zumino term, the sign in the $\kappa$–symmetry projector (7) would also change. In this case, with the same choice of (3) and (6), the no-force condition would be satisfied if in the solution (10) one of the worldvolume coordinates is equal to minus the corresponding AdS coordinate (and the static gauge must also respect this minus sign). At this the vacuum value of the $\kappa$–symmetry projector would remain the same as in (10). This is important for consistent gauge fixing the kappa–symmetry.

Let us now turn to the description of the AdS superbackground. The isometry subgroup of the target superspace in question is $G = OSp(8|4)$, $SU(2,2|4)$ or $OSp(2,6|4)$ (which corresponds, respectively, to an M2–brane, D3–brane or an M5–brane) whose bosonic subgroup is $SO(2, p+1) \times SO(D-p-1)$ $(p = 2, 3, 5; D = 10, 11)$ and the stability subgroup is $H = SO(1, p+1) \times SO(D-p-2)$. As a homogeneous coset superspace the target superspace is realised as $K = G/H$, and its bosonic subspace parametrized by coordinates $X^{\hat{m}}$ ($\hat{m} = 0, 1, ..., D$) is $AdS^{p+2} \times S^{D-p-2} = \frac{SO(2,p+1) \times SO(D-p-1)}{SO(1,p+1) \times SO(D-p-2)}$. Super-vielbeins $E^{\hat{a}}(X, \Theta)$, $\hat{E}^{\hat{a}}(X, \Theta)$ and a superconnection $\hat{\Omega}^{\hat{a}\hat{b}}$ defining the geometry of $K$ are determined as components of the Cartan one–form
\[
K^{-1} dK = E^{\hat{a}} P_{\hat{a}} + \hat{E}^{\hat{a}} \hat{Q}_{\hat{a}} + \hat{\Omega}^{\hat{a}\hat{b}} J_{\hat{a}\hat{b}},
\]
where $P_{\hat{a}}$, $\hat{Q}_{\hat{a}}$ and $J_{\hat{a}\hat{b}}$ are the generators of $G$, $J_{\hat{a}\hat{b}}$ are the generators of the stability subgroup $H = SO(1, p+1) \times SO(D-p-2)$, and $P_{\hat{a}}$ and $\hat{Q}_{\hat{a}}$ are, respectively, the bosonic and fermionic generators which correspond to the coset $K = G/H$. The explicit form of the algebra of these generators can be found in [5, 6, 7, 8, 13].

With respect to $SO(1, p+1) \times SO(D-p-2)$ the bosonic generators of $G$ split as follows.

$P_{\hat{a}} \equiv (P_{\hat{a}}, P_{a'})$, $J_{\hat{a}\hat{b}} \equiv (J_{\hat{a}\hat{b}}, J_{a'\alpha})$, where $a = 0, 1, ..., p + 1$ and $a' = 1', 2', ..., (D-p-2)'$ are indices of $AdS^{p+2}$ and $S^{D-p-2}$ coordinates, respectively.

$P_{\hat{a}}$ and $P_{a'}$ are generators of the cosets $\frac{SO(2,p+1)}{SO(1,p+1)}$ and $\frac{SO(D-p-1)}{SO(D-p-2)}$, respectively, and $\hat{Q}_{\hat{a}}$ is a 32–component spinor generator which is a Majorana spinor in $D = 11$, or consists of two Majorana–Weyl spinors in $D = 10$.

Our goal is to find the most suitable realization of the coset $K = G/H$ for getting (upon gauge fixing worldvolume diffeomorphisms and $\kappa$–symmetry) the simplest form of the worldvolume pullbacks of the supervielbeins and the superconnection.

A suitable realization turns out to be the one which corresponds to a manifest superconformal structure of the supergroup $G$ in the $p + 1$-dimensional worldvolume of the
corresponding superbrane, namely

\[ K = e^{x^i \Pi_i} e^{\rho D} e^{y^a'} P_{a'} e^{\theta^\alpha} Q_\alpha e^{\theta^\alpha S_\alpha}, \]  \hspace{1cm} (12)

where (as above) \( x^i (i = 0, ..., p) \) and \( \rho = \frac{D-p-3}{p+1} \log \frac{R}{r} \) are coordinates of \( \text{AdS}^{p+2} \), \( y^{a'} (a' = 1', ..., (D-p-2)') \) are coordinates of \( S^{D-p-2} \),

\[ \Pi_i = P_i + J_{ip+1} \]  \hspace{1cm} (13)

is the momentum generator,

\[ D = P_{p+2} \]  \hspace{1cm} (14)

is the dilatation generator of conformal transformations in \( d = p + 1 \) worldvolume and

\[ K_i = P_i - J_{ip+1} \]  \hspace{1cm} (15)

are special conformal transformations.

\( \theta \) and \( \eta \) are the following projections of \( \Theta \)-coordinates

\[ \theta = \Theta(1 - \bar{\gamma}) = (1 - \bar{\gamma})\Theta, \quad \eta = \Theta(1 + \bar{\gamma}) = (1 + \bar{\gamma})\Theta. \]  \hspace{1cm} (16)

And

\[ Q = \frac{1}{2} (1 + \bar{\gamma}) \hat{Q}, \quad S = \frac{1}{2} (1 - \bar{\gamma}) \hat{Q}. \]  \hspace{1cm} (17)

\( \bar{\gamma} \) is the same as in (11), i.e. it coincides with the classical vacuum value of the matrix \( \Gamma \) in the \( \kappa \)-symmetry projector (4), (5) of the superbrane, which corresponds to the static brane configuration (8).

It is important to notice that the projector \( (1 - \bar{\gamma}) \) also appears in the form of the Killing spinors on \( \text{AdS} \times S \) \[23, 10\]

\[ \epsilon_{\text{Kill}} = \phi(r, y^{a'}) \left[ 1 - \frac{1}{2} x^i \Gamma_i \Gamma_r (1 - \bar{\gamma}) \right] \epsilon_{\text{const}}, \]  \hspace{1cm} (18)

where \( \phi(r, y^{a'}) \) is a spinor matrix. At the same time notice that it is the Grassmann coordinate \( \theta \) which is constructed by applying to \( \Theta \) the projector \( 1 - \bar{\gamma} \). This means that in the case of the classical static solution (9) \( \eta \) transforms under the \( \kappa \)-symmetry transformations, while \( \theta \) remains invariant. Therefore, \( \eta \) corresponds to the Killing spinor in the bulk which has the property that \( (1 - \bar{\gamma}) \epsilon_{\text{Kill}}^{\text{Bulk}} = 0 \) \[10\].

The generators discussed above plus special conformal boosts \( K_i \), the boosts \( P_{a'} \) on the \( S^{D-p-2} \) sphere and generators \( J_{ij} \) and \( J_{a'b'} \) of \( \text{SO}(1, p) \times \text{SO}(D-p-2) \) form the algebra of superconformal transformations in the \( d = p + 1 \) worldvolume of the superbrane.

The advantages of the realization (12) are the following:

\footnote{Note that there always exists a realization of the gamma-matrices in which \( \bar{\gamma} \) is symmetric.}

\footnote{An analogous nonlinear realization of the \( SU(2,2|N) \) superconformal group in \( D = 4 \) (relevant to the D3–brane) was constructed a long time ago in \[24\]. For the M2–brane case a similar realization has been considered in \[8\], but the choice of the supersolvable algebra gauge caused a brane configuration studied therein to live only on a boundary of \( \text{AdS}^4 \).}
i) At $\theta = \eta = 0$ it directly leads to the $AdS^{p+2}$ metric in the form (2). This choice of the AdS metric allows one to gauge fix worldvolume diffeomorphisms by identifying $x^i$ coordinates of the $AdS^{p+2}$ with coordinates $\xi^i$ of the $d = p + 1$ worldvolume, i.e. to choose the static gauge. As we have already discussed above the choice of the static gauge must be in agreement with the no–force condition.

ii) In (12) used is the splitting of $\hat{Q} = (Q, S)$ defined in (14) with $Q$ and $S$ satisfying simple commutation relations with $\Pi_i$, $K_i$ and $D$

\[
\{Q, Q\} \sim \Gamma^i \Pi_i, \quad [\Pi_i, Q] = 0, \quad \{S, S\} \sim \Gamma^i K_i, \quad [K_i, S] = 0,
\]

\[
[D, Q] = Q, \quad [D, S] = -S, \quad [\Pi_i, S] \sim \Gamma_i Q, \quad [K_i, Q] \sim \Gamma_i S. \tag{19}
\]

As we will see below, this allows one to get in a quite simple way and in a closed (explicit) form the $\theta$– and $\eta$–dependence of the supervielbeins and superconnections.

iii) Conformal supersymmetry generated by $S_\delta$ is nonlinearly realized on the worldvolume fields, this implies that conformal supersymmetry is spontaneously broken and the coordinates $\theta(\xi)$ are Goldstone fermionic fields which are part of the physical spectrum of the $d = p + 1$ superconformal theory in the worldvolume. The fermionic coordinates $\eta(\xi)$ correspond to unbroken supersymmetries generated by $Q_\delta$. They can be gauge fixed to zero by means of $\kappa$–symmetry transformations. As we have discussed the consistency of such a gauge choice is ensured by the presence of the “vacuum” $\kappa$–symmetry projector in the definition of $\eta$ (10).

Using the methods developed in [6] we get the following general form of the Cartan forms $K^{-1}dK$ (with $K$ defined in (12)):

\[
K^{-1}dK = e^j \Pi_i + e^\hat{\alpha} Q_\delta + f^i K_i + f^\hat{\alpha} S_\alpha + E^A T_A, \tag{20}
\]

where $T_A$ stand for the generators $D$, $J_{ij}$, $P_a$ and $J_{d'\psi}$,

\[
e^j = e^j_0(x, \rho) + \Delta \eta \Gamma^i \eta, \quad \Delta \eta \equiv d\eta + e^0_\eta t_A; \tag{21}
\]

\[
e^\hat{\alpha} = (\Delta \eta + e^0_\eta \Gamma^i)^\hat{\alpha}, \tag{22}
\]

\[
f^\hat{\alpha} = [\Delta \theta + (\Delta \eta h^A \theta) g_A + e^1(\theta \Gamma_i h^A \theta) g_A]^{\hat{\alpha}}, \quad \Delta \theta \equiv d\theta + E^A \theta g_A, \tag{23}
\]

\[
E^A = e^3_0(x, \rho, y) + \Delta \eta h^A \theta + e^1(\theta \Gamma_i h^A \theta), \tag{24}
\]

\[
f^i = \Delta \theta \Gamma^i \theta + (\Delta \eta h^A \theta) (\theta g_A \Gamma^i \theta), \quad E^\hat{\alpha} = e^2_0(x, \rho, y' \alpha') \tag{25}
\]

where $e^j_0(x^i, \rho)$ are bosonic vielbeins of $AdS^{p+2}$, $e^A_0(x^i, \rho, y^a)$ are bosonic vielbeins and connections corresponding to $T_A$, and $t_{A\hat{\alpha}}$, $h^A_{\hat{\alpha}\hat{\beta}}$ and $g_{A\hat{\alpha}\hat{\beta}}$ are the following structure constants of the superalgebra $G$

\[
[T_A, Q_\delta] = t_{A\hat{\alpha}} \hat{\beta} Q_\beta, \quad \{Q_\alpha, S_\beta\} = h^A_{\alpha\hat{\beta}} T_A, \quad [T_A, S_\hat{\alpha}] = g_{A\hat{\alpha}\hat{\beta}} S_\hat{\beta}. \tag{26}
\]

Upon gauge fixing $\kappa$–symmetry by putting $\eta = 0$, the supervielbeins and superconnections take a simpler form

\[
K^{-1}dK|_{\eta=0} = e^0_0 \Pi_i + e^\hat{\alpha}_0 Q_\delta + f^i_0 K_i + f^\hat{\alpha}_0 S_\alpha + E^A_0 T_A, \tag{27}
\]
\[
e^{\hat{\alpha}} = e^i_{\hat{\alpha}}(\theta \Gamma_i)^{\hat{\alpha}}, \quad E^A_0 = e^i_0 + e^i_0(\theta \Gamma_i h^A \theta), \quad (28)
\]
\[
f^\hat{\alpha}_0 = d\theta^{\hat{\alpha}} + E^A_0(\theta g_A)^{\hat{\alpha}} + e^i_0(\theta \Gamma_i h^A \theta)(\theta g_A)^{\hat{\alpha}} = \Delta \theta^{\hat{\alpha}} + e^i_0(\theta \Gamma_i h^A \theta)(\theta g_A)^{\hat{\alpha}},
\]
\[
f^i_0 = \Delta \theta^i \theta + e^i_0(\theta \Gamma_i h^A \theta)(\theta g_A^\Gamma i \theta), \quad (29)
\]

In eqs. (21)–(29) the exact coefficients in the terms of the superforms depend on the superalgebra of \(G\) (and in each case one may expect further simplification of eqs. (22)–(25)). We observe that upon the gauge fixing (\(\eta = 0\)) the fourth power is the maximum power of \(\theta\) in the definition of the supervielbeins and the superconnections.

The bosonic \(AdS^{p+2}\) vielbeins \(e^i_0\) which appear in the realization under consideration correspond to the choice (5) of the \(AdS \times S\) metric. Note also that the following combinations of supervielbeins contribute into the definition of the induced worldvolume metric (3)

\[
E^{\hat{a}}_M = (e^i_M + f^i_M, E^\rho_M, E^{\hat{a}}_M). \quad (30)
\]

This is because the indices \(\hat{a}\) in (3) correspond to the \(AdS\) boost generators \(P^i\) (eq. (12)), and not to \(\Pi^i\) (eq. (13)).

Substituting the gauge fixed values of the supervielbeins into the superbrane action one gets an action for an interacting superconformal field theory in \(d = p + 1\), whose structure remains rather complicated and contains interacting terms up to an 8–th power in \(\theta\).

**Discussion**

We have considered a coset realization of the \(AdS \times S\) target superspaces which, upon gauge fixing the \(\kappa\)–symmetry, allows one to simplify to a certain extent the form of superbrane actions in these backgrounds.

The \(\kappa\)–symmetry gauge fixing considered in this paper should reduce the \(\Theta\)–dependence of target superforms to the fourth power also in the realizations of the \(AdS\) superspaces considered in [4, 7, 13, 15].

It would be of interest to establish the relationship between different realizations of the supervielbeins and superconnections before fixing \(\kappa\)–symmetry. It may happen that in the realizations of [4, 5, 13, 15], the power of \(\Theta\) is in fact less than 32 due to some matrix identities. Or the superforms of different realizations can be related by transformations (such as super–Weyl transformations) under which the supergravity constraints are invariant.

Much simpler form of the supervielbeins and superconnections (up to the second power in \(\theta\)) might be obtained if in (21)–(23) instead of the \(\eta\)–coordinate it might be possible to put to zero the \(\theta\)–coordinate. Such a gauge choice would correspond to the Killing spinor gauge of [11] or to a supersolvable algebra of [3]. In this gauge the static “vacuum” value of the \(\kappa\)–symmetry projector would be \((1 - \gamma)\) and not \((1 + \gamma)\) as in eq. (11). This would correspond to an “anti”–static gauge with respect to the worldvolume diffeomorphisms, i.e. when one of the worldvolume coordinates is identified with minus the corresponding coordinate of \(AdS\) (for example, identify the time coordinates as \(\xi^0 = -x^0\)). However,
there is no a classical static solution of the superbrane equations of motion obtained from the action (1) in the background (5), (6) which would be compatible with the “anti”–static gauge unless the $AdS$ radial coordinate $r$ is zero.

Alternatively, one might change the sign of the Wess–Zumino term keeping the form (5), (6) of the background. This would correspond to an anti-brane of an opposite charge propagating in the AdS background generated by a large number of branes put on the top of each other, and it is known that there is a force between branes and anti-branes.

Thus, the no–force condition would not hold and the Killing spinor (or supersolvable algebra) gauge is incompatible with the static solution in these cases. It was admissible in the case of a IIB superstring in $AdS^5 \times S^5$ [11, 12, 14] since the superstring $\kappa$–symmetry projector differs from $\tilde{\Gamma}$ in eqs. (7) and (8) for the D3–brane projector, and the no–force problem and the consistency of the gauge choice should be studied there from a different point of view (as discussed, for example, in [14]).

A group–theoretical and geometrical reason why in the cases considered in this paper the $\kappa$–symmetry allows one to eliminate $\eta(\xi)$ and not $\theta(\xi)$ is that $\theta(\xi)$ are worldvolume Goldstone fields of spontaneously broken special conformal supersymmetry in the bulk, while $\eta$ correspond to unbroken worldvolume supersymmetry which in the superembedding approach to describing superbranes has been known to be an irreducible realization of the $\kappa$–symmetry [23, 26, 27] (for recent reviews see [28]). In this approach $\eta$ can be identified with Grassmann coordinates which parametrize worldvolume supersurface embedded into target superspace. Putting the worldvolume Grassmann coordinates to zero one reduces the superembedding models to the Green–Schwarz formulation of the superbranes.

All this does not exclude that there can exist an interesting class of (non)–static (anti)–brane solutions for which $\theta = 0$ is an admissible gauge. In general such solutions will break worldvolume superconformal symmetry.

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