Supersymmetric D0-Branes in Curved Backgrounds

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Abstract
An action for supersymmetric D0-branes in curved backgrounds is obtained by dimensional reduction of N=1 ten-dimensional supergravity coupled to super Yang-Mills system to 0+1 dimensions. The resultant action exhibits the coset-space symmetry \( \frac{SO(9,9+n)}{SO(9) \times SO(9+n)} \times U(1) \) where \( n = N^2 - 1 \) is the dimension of the SU(N) gauge group.
1 Introduction

Very little is known about D-branes [1], [2], [3], in curved backgrounds [4] and what restrictions, if any, should be imposed on such backgrounds. Supersymmetry imposes constraints on the background metric. In the case of the superstring with world sheet supersymmetry, there is a direct relation between the number of supersymmetries and the background metric [5]. In flat background geometry the D0-brane action has 16 space-time supersymmetries, and it is natural to ask for the type of curved backgrounds compatible with this symmetry. The action with 8 space-time supersymmetries (N = 2 in 4 dimensions) was shown to correspond to Kähler backgrounds [6].

In general it is difficult to construct such actions without having determined the underlying symmetry. There are few possible routes to handle this problem, the most obvious one is to quantize the D-brane action in the presence of a general superspace metric, thus keeping all supersymmetries, and to determine what are the required constraints on such a metric [7], [8]. This is expected to be extremely complicated and only recently some work has been done in this direction [9]. The other possibility is to determine the necessary fields to make a supersymmetric multiplet and to find the corresponding invariant action under such transformations. We shall follow a simpler approach, which is straightforward but which the drawback that the underlying symmetry is not manifest. The idea is based on the observation that the supersymmetric D-0 brane action was obtained by dimensionally reducing the super Yang-Mills action from ten to 0+1 dimensions [8]. The most general supersymmetric interaction in ten dimensions with N = 1 supersymmetry is that of supergravity coupled to super Yang-Mills with an arbitrary gauge group [10]. Dimensional reduction keeps maximal supersymmetry and rearranges the scalar fields to have the action of a non-linear sigma model on a coset space. Reducing to 0+1 dimensions have the peculiarity that there is no gravitational part and the only fields coming from the gravity sector in ten-dimensions yields scalar fields. Similarly the vector fields coming from the Yang-Mills part give scalar fields taking values in the adjoint representation of the gauge group. As with compactification of supergravity theory it is expected that the coset space to be of the form [11]

\[
\frac{SO(9, 9 + n)}{SO(9) \times SO(9 + n) \times U(1)}
\]

where n = N^2 − 1 is the dimension of SU(N) gauge group. The absence of the gravitational sector in 0+1 dimensions also enables us to identify the gravitational coupling in higher dimensions with the string tension to insure that all rescaled scalar fields have the same dimensions and could serve as coordinates of the D-brane.

The aim of this letter is to derive the D0-brane action with the most general curved background compatible with maximal space-time supersymmetry. This is done by compactifying the ten-dimensional theory and grouping all the resultant fields. The result obtained does not have manifest symmetry. Nonetheless, this suggests that a direct derivation in terms of supermultiplets, where some auxiliary fields as well as constrained variables are used, might drastically simplify the answer. This is the situation encountered in the derivation of the four-dimensional N = 4 supersymmetric action where superconformal methods were used to simplify the analysis [12]. This more systematic approach will be
left for the future and our study here will be limited to the action obtained by dimensional reduction. The plan of this paper is as follows. In section 2 we derive the bosonic dimensionally reduced action and in section 3 we give the fermionic part. Section four includes comments on the results.

2 Bosonic action of curved D0-branes.

Our starting point is the \( N = 1 \) supergravity Lagrangian in ten dimensions coupled to super Yang-Mills system. This is given by (up to quartic fermionic terms) \(^ {10} \)

\[
\det \left( e^A_M \right)^{-1} L = -\frac{1}{4\kappa^2} R(\omega) - i\frac{1}{2} \psi_M \Gamma^{MNP} D_N \psi_P + \frac{1}{2\kappa^2} \partial_M \phi \partial_N \phi g^{MN} \\
+ i\frac{1}{2} \Gamma^M D_M \chi + \frac{1}{\sqrt{2}} \partial_M \phi \psi_M \chi + e^{-2\phi} F'_{MNP} F'_{QRS} g^{MQ} g^{NR} g^{PS} \\
+ \frac{i\kappa}{24} e^{-\phi} \psi_M \left( \Gamma^{MNPQR} - 6g^{MP} \Gamma^Q g^{RN} \right) \psi_N F'_{PQR} \\
- \frac{1}{4} \epsilon^{-\phi} \text{Tr} \left( G_{MN} G_{PQ} \right) g^{MP} g^{NQ} + \frac{i}{2} \text{Tr} \left( A^M \Gamma^{NP} \psi_M G_{NP} \right) \\
- \frac{i\kappa}{2\sqrt{2}} \text{Tr} \left( \bar{\chi} \Gamma^M \Gamma^{NP} \psi_M G_{NP} \right)
\]

where \( F'_{MNP} \) is the field strength of the antisymmetric tensor \( B_{MN} \) modified by the gauge Chern-Simons three form.

\[
F'_{MNP} = F_{MNP} + \omega^{(CS)}_{MNP}
\]

and \( F_{MNP} = \frac{3}{2} \partial_{[M} B_{NP]} \), while

\[
\omega^{(CS)}_{MNP} = 6\kappa \text{Tr} \left( A_{[M} \partial_N A_{P]} + \frac{2}{3} A_{[M} A_N A_{P]} \right)
\]

We can reduce this action from 10 to \( d \) dimensions with the following distribution of fields. The metric \( g_{MN} \) gives a metric \( g_{\mu\nu} \), \( m \) vectors and \( \frac{1}{2} m (m + 1) \) scalars in \( d \) dimensions, where \( m = 10 - d \). The antisymmetric tensor \( B_{MN} \) gives \( B_{\mu\nu} \), \( m \) vectors and \( \frac{1}{2} m (m - 1) \) scalars. The gauge fields \( A^M \) give \( A^i_{\mu} \), and \( nm \) scalars \( A^i_m \), where \( n \) is the dimension of the gauge group. All in all we will have \( m (n + m) \) scalars which will span the coset space \(^ {11} \)

\[
\frac{SO(m, n + m)}{SO(m) \times SO(n + m)}
\]

The case when \( d = 4 \) (i.e. \( m = 6 \)) is well established \(^ {10} \). The case we are interested in have \( m = 9 \).

To reduce this action to 0+1 dimensions we decompose:

\[
e^A_M = \begin{pmatrix} e^0_0 & B^a_0 \\ 0 & e^a_m \end{pmatrix}
\]

The inverse metric is

\[
e^M_A = \begin{pmatrix} e^0_0 & e^m_0 \\ 0 & e^m_a \end{pmatrix}
\]
where $e_0^m = \left( e_0^a \right)^{-1}$, $e^m_a e_m = \delta^a_b$, $e_0^m = -e_0^a B_a e_m$.

To evaluate $-\frac{1}{4} \det \left( e_A^a \right) R(\omega)$ we have

$$\frac{1}{4} \det \left( e_M^A \right) R(\omega) = \frac{1}{16} e \left( \Omega_{ABC} \Omega^{ABC} - 2 \Omega_{ABC} \Omega^{CAB} - 4 \Omega_{CA} \Omega^{CB} \right)$$

where $\Omega_{ABC} = -e_M^A e_B^N \left( \partial_M e_N^C - \partial_N e_M^C \right)$. The only non-vanishing $\Omega_{ABC}$ is

$$\Omega_{0bc} = -e_0^a \partial_0 e_{bc}$$

Substituting this gives

$$-\frac{1}{4} \det \left( e_M^A \right) R(\omega) = \frac{e}{8} \left( -\frac{1}{2} \partial_0 g_{mn} \partial_0 g^{mn} - 2 \left( \partial_0 \ln e \right)^2 \right)$$

where $e = \det(e_m^a)$. The antisymmetric tensor piece gives

$$\frac{e}{12} \partial_0 e^{-2\phi} \left( 3 F^{'a}_{a0} F^{'a}_{a0} - F^{'a}_{abc} F^{'a}_{abc} \right)$$

where we are raising and lowering tangent space indices with the euclidean metric $\delta_{ab}$, and

$$F^{'a}_{abc} = 4 \kappa e_n^a e_b^m \partial_0 \left( A_m A_n A_p \right)$$

and

$$F^{'a}_{a0} = \frac{1}{\kappa} e_n^a e_b^m \left( \partial_0 B_{mn} - \kappa^2 \partial_0 \left( A_m D_0 A_n - A_n D_0 A_m \right) \right)$$

where $D_0 A_m = \partial_0 A_m + \left[ A', A_m \right]$, and $A' = A - B^m A_m$. The redefinition of $A_0$ will insure that the field $B^m$ will not appear in the action and therefore is irrelevant.

The Yang-Mills part is

$$\frac{1}{4} \det \left( e_M^A \right) e^{-\phi} \partial_0 \left( G_{MNPQ} g^{MN} g^{PQ} \right) = \frac{e}{4} e_0^a e^{-\phi} \partial_0 \left( G_{ab} G_{ab} - 2 G_{a0} G_{a0} \right)$$

where

$$G_{ab} = e_0^m e_n^a [A_m, A_n]$$

$$G_{a0} = -e_0^m e_0^a D_0 A_m$$

Grouping all terms together we obtain the bosonic part of the D0-brane action:

$$e^{-1} L_b = \frac{1}{\kappa^2} \left( e_0^a \left( \frac{1}{16} \partial_0 g_{mn} \partial_0 g^{mn} - \frac{1}{2} \left( \partial_0 \ln e \right)^2 + \frac{1}{2} \partial_0 \phi \partial_0 \phi \right) + \frac{1}{4} e^{-2\phi} g^{mp} g^{nq} D_0 B_{mn} D_0 B_{pq} + \kappa^2 e^{-\phi} g^{mn} \partial_0 \left( D_0 A_m D_0 A_n \right) \right)$$

$$+ e_0^0 \left( -\frac{4\kappa^2}{3} e^{-2\phi} g^{mq} g^{nr} g^{ps} \partial_0 \left( A_{[m} A_{n]} A_{p] \right) \partial_0 \left( A_{q} A_{s} \right) \right)$$

$$- \frac{1}{4} e^{-\phi} g^{mn} g^{pq} \partial_0 \left( [A_{m}, A_{n}] \right) \partial_0 \left( [A_{q}, A_{s}] \right)$$
where $D_0 B_{mn} = \partial_0 B_{mn} - \kappa^2 T r \left( A_m D_0 A_n - A_n D_0 A_m \right)$. To get the correct dimensions we identify the gravitational coupling $\kappa$ with the string tension $\alpha'$ and redefine the gauge fields $A^i_m = \frac{1}{\alpha'} X^i_m$, thus identifying them with the D0-brane coordinates. Multiplying the Lagrangian with an overall factor of $(\alpha')^2$ gives the bosonic part of the D0-brane action. The 82 fields $g_{mn}, B_{mn}$ and $\phi$ are needed, beside the coordinates $X^i_m$, to provide coordinates for a D0-brane action with a curved background. The rescaled bosonic Lagrangian becomes

\[
e^{-1} L_b = \left( e^0_0 \left( -\frac{1}{16} \partial_0 g_{mn} \partial_0 g^{mn} - \frac{1}{4} \left( \partial_0 \ln e \right)^2 + \frac{1}{2} \partial_0 \phi \partial_0 \phi \right. \right.
\]
\[\left. + 4 e^{-2\phi} g^{mp} g^{nq} D_0 B_{mn} D_0 B_{pq} + \frac{\kappa^2}{2} e^{-\phi} g^{mn} T r \left( D_0 X_m D_0 X_n \right) \right) \]
\[\left. + e^0_0 \left( -\frac{4}{3 (\alpha')^2} e^{-2\phi} g^{mr} g^{ns} T r \left( X_{[m} X_{n]} X_{p]} \right) T r \left( X_{[q} X_{r]} X_{s]} \right) \right) \right.
\[\left. - \frac{1}{4 (\alpha')} e^{-\phi} g^{mr} g^{ns} T r \left( X_{[m} X_{n]} \right) T r \left( X_{[q} X_{r]} \right) \right) \right).
\]

The scalar fields can be regrouped into a set $X^R_m$ where $R = 1, \ldots, 9 + n$ plus an additional scalar field as a combination of the fields $g_{mn}, B_{mn}, \phi$ and $X^i_m$. Can one take the limit where $g_{mn} = \delta_{mn}, B_{mn} = 0$ and $\phi = 0$? This gives the flat background D0-brane action plus the order 6 terms in $X^i_m$. This is usually incompatible with supersymmetry as we shall show later. The proper limit to flat backgrounds can be obtained by keeping the couplings $\kappa$ and $\alpha'$ independent, then taking the limit $\kappa \to 0$.

The transformation law for $e^0_0$ with respect to time transformation is given by

\[
\delta e^0_0 = \partial_0 \left( e^0_0 \xi^0 \right)
\]

which would allow us to set $e^0_0 = 1$.

In this action the coset space symmetry is not manifest. The coset space metric is a non-polynomial function of $g_{mn}, B_{mn}, X^i_m$ and $\phi$. To obtain manifest symmetry, one method would be to start with the symmetry $SO(9, 9 + n)$ using supersymmetric multiplets, and then gauge the $SO(9) \times SO(9 + n)$ subgroup. This will be the topic of a forthcoming project where a systematic analysis of all possible background symmetries would be carried out.

### 3 The fermionic action

The Rarita-Schwinger term

\[
-\frac{i}{2} \det(e^M_A) \overline{\psi}_{A} \Gamma^{ABC} D_B \psi_C
\]

where $\psi_A = e^M_A \psi_M$, and $D_M \psi_N = (\partial_M + \frac{1}{2} \omega_M^{AB} \Gamma_{AB}) \psi_N$, gives upon compactification

\[
-\frac{i}{8} e \left( \overline{\psi}_a \Gamma_b \psi_d \left( e^e_b \partial_0 e_{nd} + e^e_d \partial_0 e_{nb} \right) - 2 \overline{\psi}_0 \Gamma^c e^e_c \left( e^e_d \partial_0 e_{nd} \right) \right)
\]
\[
+ \frac{i}{2} \epsilon \left( \overline{\psi}_{a} \Gamma^{ac} \Gamma_{0} \partial_{a} \psi_{c} + \frac{1}{4} \overline{\psi}_{a} \Gamma^{ac} \Gamma_{de} \Gamma_{0} \psi_{c} \left( e_{d}^{n} \partial_{0} e_{ne} \right) \right) \\
+ \frac{i}{4} \epsilon e_{0} \left( \overline{\psi}_{a} \Gamma^{n} \psi_{0} \left( e_{b}^{n} \partial_{0} e_{nb} \right) + \frac{1}{2} \overline{\psi}_{a} \Gamma_{0} \Gamma_{b} \psi_{0} \left( e_{b}^{n} \partial_{0} e_{na} + e_{a}^{n} \partial_{0} e_{nb} \right) \right)
\]

where we have used \( \psi_{0} = e_{0}^{a} \left( \psi - B_{0} \psi_{m} \right) \) and \( \psi_{a} = \epsilon_{a}^{n} \psi_{m} \). Again, the definition of \( \psi_{0} \) insures that \( B_{0} \) does not appear in the action. The nonvanishing components of the spin-connection are

\[
\omega_{0bc} = \frac{1}{2} e_{0}^{a} \left( e_{b}^{n} \partial_{0} e_{nc} - e_{c}^{n} \partial_{0} e_{nb} \right) \\
\omega_{ab0} = -\frac{1}{2} e_{0}^{a} \left( e_{b}^{n} \partial_{0} e_{na} + e_{a}^{n} \partial_{0} e_{nb} \right)
\]

The term \( \frac{1}{2} \det \left( e_{M}^{A} \right) \Gamma^{A} D_{A} \chi \) reduces to

\[
- \frac{i}{4} \epsilon \left( e_{0}^{a} \Gamma^{ab} \chi \left( e_{a}^{n} \partial_{0} e_{na} \right) - 2 \Gamma^{0} \left( \partial_{0} + \frac{1}{4} e_{b}^{n} \partial_{0} e_{nc} \Gamma_{bc} \right) \chi \right)
\]

Similarly the gaugino kinetic term gives upon reduction

\[
- \frac{i}{4} e \text{Tr} \left( e_{0}^{a} \Gamma^{ab} \chi \left( e_{a}^{n} \partial_{0} e_{na} \right) - 2 \Gamma^{0} \left( \partial_{0} + \frac{1}{4} e_{b}^{n} \partial_{0} e_{nc} \Gamma_{bc} \right) \chi \right)
\]

Next, the fermi-bose interaction \( \overline{\psi} \psi F \) gives

\[
\frac{i \kappa^{2}}{24} e^{-\phi} \epsilon e_{0}^{a} \left( 4 \left( \overline{\psi}_{a} \Gamma^{abcde} \psi_{b} + 6 \overline{\psi}_{a} \Gamma^{0bcde} \psi_{b} - 6 \overline{\psi}_{a} \Gamma_{d} \psi_{c} \right) \text{Tr} A_{a} A_{d} A_{c} \right) \\
+ \left( 3 \overline{\psi}_{a} \Gamma^{abcde} \psi_{b} - 2 \overline{\psi}_{a} \Gamma_{d} \psi_{c} + 2 \overline{\psi}_{a} \Gamma_{0} \psi_{d} - 2 \overline{\psi}_{a} \Gamma_{d} \psi_{b} \right) \\
\times e_{c}^{n} e_{a}^{n} e_{0}^{a} \left( \frac{1}{\kappa} \partial_{0} B_{mn} - \kappa \text{Tr} \left( A_{m} D_{0} A_{a} - A_{n} D_{0} A_{m} \right) \right)
\]

Finally the \( \overline{\psi} \chi \partial \phi \) coupling gives

\[
\frac{1}{2} e \partial_{0} \phi \overline{\psi} \chi.
\]

The supersymmetry transformations in ten dimensions are

\[
\delta e_{M}^{A} = -i \kappa \Gamma^{A} \psi_{M} \\
\delta \phi = - \frac{\kappa}{\sqrt{2}} \chi \\
\delta B_{MN} = \kappa e^{\phi} \left( i \Gamma_{[M} \psi_{N]} + \frac{1}{2} \Gamma_{MN} \chi \right) \\
\delta \psi_{M} = \frac{1}{\kappa} \Gamma M \epsilon + \frac{1}{48} e^{-\phi} \left( \Gamma^{NPQ} M + 9 \delta^{N}_{M} \Gamma^{PQ} \right) \epsilon F_{NPQ} \\
\delta \chi = \frac{i}{2 \kappa} \Gamma^{M} \epsilon \partial_{M} \phi \\
\delta A_{M} = \frac{i}{\sqrt{2}} e^{\phi} \Gamma^{M} \lambda \\
\delta \lambda = e^{-\frac{1}{2} \phi} \Gamma^{MN} \epsilon G_{MN}
\]
The compactified supersymmetry transformations become after rescaling

\[ \delta e_0^0 = -i \kappa \Gamma^0 \psi_0 \]
\[ \delta e_a^0 = -i \kappa \Gamma^a \psi_0 \]
\[ \delta e_m = -i \Gamma^a \psi_m \]
\[ \delta \phi = -\frac{\kappa}{\sqrt{2}} \chi \]

\[ \delta B_{mn} = \kappa \epsilon^\phi \left( \pi \Gamma_{[m} \psi_n] + \frac{1}{2 \sqrt{2}} \Gamma_{mn} \chi \right) \]

\[ \delta \psi_0 = \frac{1}{\kappa} \epsilon_0^\phi \left( \partial_0 + \frac{1}{4} e_m^a \partial_0 e_a \Gamma^{ab} \right) \epsilon \]
\[ + \frac{1}{12} e^{-\phi} \epsilon^c \epsilon_d \left( \frac{1}{\kappa} \partial_0 B_{mn} - \frac{\kappa}{(\alpha')^2} \text{Tr} \left( X_m D_n X_n - X_n D_m X_m \right) \right) \epsilon \]
\[ + \frac{\kappa}{12 \alpha')^2} e^{-\phi} \Gamma^{bcd} \epsilon_b \epsilon_c \epsilon_d \left( X_m X_n - X_n X_m \right) \]

\[ \delta \psi_a = \frac{1}{4 \kappa} \epsilon_0^\phi \left( \epsilon_a \partial_0 \epsilon_b + \epsilon_b \partial_0 \epsilon_a \right) \Gamma_{b0} \epsilon \]
\[ - \frac{\kappa}{12 \alpha')^3} e^{-\phi} \left( \Gamma_{abcd} - 4 \delta^b_a \delta^c_d \right) \epsilon_b \epsilon_c \epsilon_d \left( X_m X_n X_p \right) \]
\[ + \frac{1}{16} e^{-\phi} \left( \Gamma_{abcd} - 4 \delta^b_a \delta^c_d \right) \epsilon_b \epsilon_c \epsilon_d \left( X_m X_n X_p \right) \]

\[ \delta \chi = \frac{i}{\kappa \sqrt{2}} \epsilon_0^\phi \partial_0 \epsilon_0 + \frac{i \kappa}{3 \sqrt{2} \alpha')^3} e^{-\phi} \epsilon^a \epsilon_b \epsilon_c \epsilon_d \left( X_m X_n X_p \right) \]
\[ + \frac{i}{4 \sqrt{2} \alpha')^3} \epsilon^a \epsilon_b \epsilon_c \epsilon_d \left( X_m X_n X_p \right) \]

\[ \delta A_0 = \frac{i}{\alpha')^2} e^{-\phi} \Gamma^0 \chi \]
\[ \delta X_m = \frac{i \alpha'}{\sqrt{2}} e^{-\phi} \Gamma m \chi \]
\[ \delta \lambda = \frac{1}{\alpha')^2} e^{-\phi} \Gamma \lambda \]

From these transformations it should be clear that the truncation \( g_{mn} = \delta_{mn}, B_{mn} = 0, \phi = 0 \) is not consistent with supersymmetry because the fields \( X^i, g_{mn}, B_{mn} \) and \( \phi \) are now mixed to form \( 9(9 + n) + 1 \) coordinates for the D-0-brane. A proper way of going to the flat background limit is to keep \( \kappa \) and \( \alpha' \) distinct, and then take the limit \( \kappa \rightarrow 0 \).

## 4 Comments

The D0-brane action with maximal \( N = 16 \) space-time supersymmetry, derived here have the coset symmetry \( SO(9,9+n)/SO(9+1) \) which, however, is not manifest. The 81 fields which are not related to the \( SU(N) \) gauge group are essential to provide curvature for the background. We can say that curved backgrounds are
only possible once $n$ gauge fields are embedded into the above coset structure. One way to improve on this solution is to start with the light-cone formulation of the supermembrane in arbitrary background and find under what conditions the quantized action simplifies in such a manner as not to involve a square root. This is a difficult problem and the recent work of weakly coupling D0-branes to curved backgrounds may help to clarify the situation \cite{Taylor:1999}. The action constructed in \cite{Taylor:1999} contains higher order time derivatives indicating that higher derivative terms in ten-dimensional supergravity should also be included. It would also be very interesting to find out how the 82 scalar fields arise in this formulation.

Another possibility is to study supersymmetry representations in 0+1 dimensions and form multiplets in complete analogy with the one obtained by superconformal methods in four dimensions. This would have the advantage of getting the coset space symmetry in a manifestly invariant way. In addition this would allow to investigate the general problem of finding the relation between the required degree of space-time supersymmetry and the nature of the curved background.

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