Supergravity Solutions For $AdS_3 \times S^3$ Branes

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Abstract

We find a large class of supersymmetric solutions in the $AdS_3 \times S^3$ background with NS fluxes. This two-parameter family of solutions preserves 8 of the 16 supersymmetries of the background.

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1 Introduction

In recent years, a great deal of progress has been made in the study of both
gauge theory and gravity through the AdS/CFT correspondence. This cor-
respondence allows one to relate a theory of quantum gravity in a particular
background to a dual conformal field theory living on the boundary. This
correspondence is best understood when the bulk space-time background is
an anti-de Sitter space. Unfortunately, even in this case, it is quite difficult to
study the bulk gravitational theory beyond the supergravity limit. Recently,
some progress has been made in this area in several different contexts ([1]
[2], among others).

However, the best understood case is the $AdS_3 \times S^3 \times T^4$ solution with
NS-NS gauge fields turned on. This background arises from the embedding of
a stack of fundamental strings within a stack of NS-5-branes (when one then
takes the near-horizon limit). The action of a string in this background is a
Wess-Zumino-Witten model on the group manifold $SL(2,R) \times SU(2)$. Since
$SL(2,R)$ is non-compact, this theory is still somewhat subtle; nevertheless
progress has been made in studying both the open and closed strings in this
background ([3][4]).

It is thus useful to construct supergravity solutions for D-branes in $AdS_3$,
which will then be described in a dual description by a field theory. In the
particular case that we shall analyze, the supergravity solution we find is
believed to be dual to a defect conformal field theory on the boundary [5].

We will be interested in defects which preserve 8 of the 16 supersymme-
tries of the $AdS_3 \times S^3 \times T^4$ background. In [6], it was shown using boundary
state arguments that a D3-brane embedded in an $AdS_2 \times S^2$ submanifold of
this background is $\frac{1}{2}$ BPS. This is twice the number preserved by a system
containing D3-branes, fundamental strings and NS-5 branes, due to the en-
hancement of supersymmetry [7]. General conditions for the preservation of
supersymmetry by branes in $AdS_3$ were also considered by [8].

Supergravity solutions for general intersecting branes were considered in
[9,10,11,12,13,14,15,16]. Secondly, in [17,18], brane probes were consid-
ered. The equations of motion derived from the Born-Infeld action were then
solved to produce stable supersymmetric branes. Unfortunately, the latter
approach only produces solutions to linearized supergravity. Instead, we will
follow the general approach of [10,12]. This generalizes the result of [19].

In the section 2, we study the $\kappa$-symmetry of D3-branes in $AdS_3 \times S^3$.
The gauge-fixing of $\kappa$-symmetry will determine which Killing spinors are
preserved by the supergravity solution. In section 3, we use this to solve the Killing equations. In \cite{19}, a single $\frac{1}{2}$-BPS solution was found. This restriction on the location of the sources was due to certain assumptions made in order to make the Killing equations more tractable. Using much more general assumptions, we are able to find a solution written in terms of one function, with two non-linear partial differential equations as constraints. We argue that a two-parameter family of solutions exists. Generically, the sources for these solutions do not appear to be localized (at least to first order).

2 Obtaining Supersymmetry from $\kappa$-Symmetry

A D-brane typically breaks some of the supersymmetries of the background. The preserved supersymmetries (labelled by Killing spinors) are most easily found by the analysis of $\kappa$-symmetry.

In a consistent supersymmetric brane solution, the preserved Killing spinors (when projected onto the brane) must be invariant under a gauged fermionic symmetry known as $\kappa$-symmetry. The projector which projects onto these invariant spinors can be found very generally, and from this we can deduce the form of the preserved Killing spinors.

The background metric of $AdS_3 \times S^3$ in global coordinates may be written as

$$ds^2 = d\psi^2 + \cosh^2 \psi (d\omega^2 - \cosh^2 \omega d\tau^2) + \sin^2 \theta (d\phi^2 + \sin^2 \phi d\chi^2)$$

(1)

The background $B_{NS}$ fields are

$$\bar{B}_{\phi \chi} = \frac{1}{2} (\theta - \sin \frac{2\theta}{2}) \sin \phi$$
$$\bar{B}_{\tau \phi} = \frac{1}{2} (\psi - \sinh \frac{2\psi}{2}) \cosh \omega$$

(2)

This background preserves a 16-component Killing spinor.

In \cite{6}, it was argued that a D3-brane embedded along an $AdS_2 \times S^2$ submanifold of $AdS_3 \times S^3$ would preserve 8 of the 16 supersymmetries of the background. This embedding may be parameterized as a brane stretching along the $(\tau, \omega, \phi, \chi)$ coordinates of the global coordinate system, with the $\psi$ and $\theta$ coordinates acting as transverse scalars.
One needs to turn on a gauge field strength on the D3-brane, which is of the form

\[ 4\pi\alpha' F_{\phi\chi} = -\pi p \sin \phi \]
\[ 4\pi\alpha' F_{\tau\omega} = -\pi q \cosh \omega \] (3)

The Born-Infeld Lagrangian for this D3-brane is then

\[ L_{DBI} = -T\sqrt{-\det M} \] (4)

where

\[ \sqrt{-\det M} = N(\psi)L(\theta)\cosh \omega \sin \phi \]
\[ L(\theta) = (\sin^4 \theta + (\theta - \frac{\sin 2\theta}{2} - \pi p)^2)^{\frac{1}{2}} \]
\[ N(\psi) = (\cosh^2 \psi - (\psi + \frac{\sinh 2\psi}{2} - \pi q)^2)^{\frac{1}{2}} \] (5)

The resulting equations of motion are solved by

\[ \theta_0 = \pi p \]
\[ \psi_0 = \pi q \] (6)

and we have

\[ L(\theta_0) = \sin \theta_0 \]
\[ N(\psi_0) = \cosh \psi_0 \]
\[ F_{\phi\chi} = -\frac{1}{2} \sin 2\theta_0 \sin \phi \]
\[ F_{\tau\omega} = -\frac{1}{2} \sinh 2\psi_0 \cosh \omega \] (7)

The projection onto \(\kappa\)-invariant spinors is given by \[20\]

\[ d^{p+1} \xi \Gamma_0 = -e^{-\Phi} L_{DBI}^{-1} e^{F_{\text{total}}} \wedge X|_{\text{vol}}, \] (8)

where
\[ X = \bigoplus_n \Gamma_{(2n)} K^n I \]
\[ K \psi = \psi^* \]
\[ I \psi = -\imath \psi \]
\[ \Gamma_{(n)} = \frac{1}{n!} d\xi^{i_n} \wedge ... \wedge d\xi^{i_1} \Gamma_{i_1...i_n} \]  

(9)

We see that

\[ \Gamma_0 = -\imath (\cos \theta_0 \cosh \psi_0 \gamma_{\tau \omega} K + \sin \theta_0 \cosh \psi_0 \gamma_{\phi \chi} K + \sin \theta_0 \sinh \psi_0 \gamma_0 K + \cos \theta_0 \sinh \psi_0) \]  

(10)

where \( \Gamma_0 \) is traceless and \( \Gamma_0^2 = 1. \)

The \( \kappa \)-symmetry projector \( \Gamma_0 \) acts on the background killing spinor at the brane. We also need the projector away from the brane. Let \( \epsilon \) be the Killing spinor preserved in the presence of the brane. We will assume that \( \epsilon \) is of the form \( \epsilon = f \bar{\epsilon} \) where \( f \) is a complex function, and \( \bar{\epsilon} \) is the Killing spinor of the background \( \text{AdS}_3 \times S^3. \)

\[ \bar{\epsilon} = \exp \left( -\frac{\psi}{2} \gamma_{\tau \omega} K \right) \exp \left( -\frac{\theta}{2} \gamma_{\phi \chi} K \right) R_0(\phi, \chi, \omega, \tau) \bar{\epsilon}_0 \]
\[ = \Lambda \bar{\epsilon}(\psi_0, \theta_0) \]  

(11)

where

\[ \Lambda = \exp \left( -\frac{(\psi - \psi_0)}{2} \gamma_{\tau \omega} K \right) \exp \left( -\frac{(\theta - \theta_0)}{2} \gamma_{\phi \chi} K \right) \]  

(12)

To preserve supersymmetry, we must have \((1 - \Gamma_0) \bar{\epsilon} = 0\) at \( \psi = \psi_0 \) and \( \theta = \theta_0 \) (but at arbitrary \( \tau, \omega, \phi \) and \( \chi \)), i.e. \((1 - \Gamma_0) \bar{\epsilon}(\psi_0, \theta_0) = 0.\)

We can write this globally as \((1 - \Gamma) \epsilon = 0\) where

\[ \Gamma = f \Lambda \Gamma_0 \Lambda^{-1} f^{-1} \]  

(13)
After some algebra, we find
\[
\Gamma = -\mathbf{i}(M + N\gamma_\tau K + O\gamma_\phi K + P\gamma_\tau\omega K)
\]
\[
M = \cos \theta \sinh \psi \quad N = \frac{f}{f^*} \cos \theta \cosh \psi
\]
\[
O = \frac{f}{f^*} \sin \theta \sinh \psi \quad P = \sin \theta \cosh \psi
\] (14)

Note that this projector is now independent of $\psi_0$ and $\theta_0$. We can rewrite this projector in the form
\[
\gamma_\tau K \epsilon = (A + B\gamma_\psi \theta) \epsilon
\] (15)

where
\[
A = -\frac{f^*}{f} \frac{\cosh \psi}{\cosh^2 \psi - \sin^2 \theta} (\sinh \psi + \mathbf{i} \cos \theta)
\]
\[
B = \frac{f^*}{f} \frac{\mathbf{i} \sin \theta}{\cosh^2 \psi - \sin^2 \theta} (\sinh \psi + \mathbf{i} \cos \theta)
\] (16)

3 The Killing Equations

We have found the form of the supersymmetry. We now turn to a more detailed analysis of the supersymmetry.

The supersymmetry variations are generally of the form
\[
\delta \psi_\mu = D_\mu \epsilon + ...
\] (17)

Supersymmetry is preserved if the variations all vanish. Every choice of $\epsilon$ for which these variations vanish is called a Killing spinor, and produces a different preserved supersymmetry.

In our case we have already found the form of the Killing spinor. We are thus requiring that the supersymmetry variations vanish for the Killing spinors of the form $\epsilon = f \bar{\epsilon}$, with $(1 - \Gamma)\epsilon = 0$.

We will now explicitly write out the form of the supersymmetry variations, substitute the form of the Killing spinor into them, and find the background solution that preserves this Killing spinor.

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3 We use the conventions of [19], which arise from [21] and [22]
We begin with a few assumptions to constrain the solution. First, we assume that this solution is self-dual in 6 dimension, ie. $G^{\psi \tau \omega} = G^{\phi \chi \theta}$ and $G^\tau \omega \theta = - G^{\psi \phi \chi}$. The killing equations (for the dilatino) then imply that complex axion-dilaton $\Phi$ is constant. Furthermore, the equations of motion and generalized Bianchi identity imply that the 5-form field strength $F$ vanishes.

We also assume that all fields depend only on the coordinates $\psi$ and $\theta$.

The killing equations in Einstein frame can be written as \[23\]
\[
\partial_\mu \epsilon + \frac{1}{4} \omega_{\mu}^{ab} \gamma_{ab} + \frac{l}{192} \Gamma_{\mu}^{\nu \rho} \gamma_{\nu \rho \epsilon} \epsilon^* + \frac{l}{48} e^\Phi (G_{\mu \nu \rho \sigma} \gamma_{\nu \rho \sigma \epsilon} - 9 G_{\mu \nu \rho} \gamma_{\nu \rho \epsilon} \epsilon^*) = 0
\]
where $G$ is a complex field whose real and imaginary parts are the RR and NS-NS three-form field strengths, respectively. The value of this field for the background is $\bar{G}^{\psi \tau \omega} = \bar{G}^{\phi \chi \theta} = -l$.

For example, the $\psi$ Killing equation can be written as
\[
- \partial_\psi f \gamma_{\psi} \epsilon - \frac{1}{2} \omega_{\psi}^{\phi \theta} \gamma_{\theta} \epsilon + \frac{l}{2} G^{\psi \tau \omega} \gamma_{\psi \tau \omega \epsilon} \epsilon^* - \frac{1}{2} G^{\tau \omega \theta} \gamma_{\tau \omega \theta \epsilon} \epsilon^* =
\]
\[
\frac{1}{2} f \bar{e}_{\psi} \gamma_{\psi \tau \omega} \epsilon^*
\]
and the other Killing equations take a similar form.

Since we know that $\epsilon$ satisfies the $\kappa$-symmetry projection, we can use \[15\] to simplify these equations. The above equation, for example becomes:
\[
- \partial_\psi f \gamma_{\psi} \epsilon - \frac{1}{2} \omega_{\psi}^{\phi \theta} \gamma_{\theta} \epsilon + \frac{l}{2} G^{\psi \tau \omega} (A \gamma_{\psi} + B \gamma_{\theta}) \epsilon + \frac{l}{2} G^{\tau \omega \theta} (B \gamma_{\psi} - A \gamma_{\theta}) \epsilon =
\]
\[
\frac{1}{2} f \bar{e}_{\psi} \gamma_{\psi \tau \omega} \epsilon^* (A \gamma_{\psi} + B \gamma_{\theta}) \epsilon (20)
\]

This is only satisfied if the coefficients of $\gamma_{\psi}$ and $\gamma_{\theta}$ are each zero. Therefore we get two complex equations. The other Killing equations similarly yield other complex equations.

By taking linear combinations of these equations, we solve for the various field strengths in terms of the spin connections, and also get equations for the spin connections. The solution for the field strengths in terms of the spin connection is then
\[
- i e^\Phi G^{\tau \omega \theta} (A^2 + B^2) = \frac{1}{2} \omega_{\omega}^{\psi} B - \frac{1}{2} \omega_{\omega}^{\omega \theta} A - \frac{1}{2} \omega_{\phi}^{\psi} B + \frac{1}{2} \omega_{\phi}^{\phi \theta} A
\]
\[ G^{\psi \tau} = i \beta G^{\tau \psi} - i \frac{e^{\psi}}{e^{\phi}} \left( \frac{f^*}{f^2} + \frac{\partial_{\psi} \log f^2}{A} \right) \] (21)

We also have the constraints

\[ Re \log f^2 = \log \frac{e^{\phi}}{e^{\omega}} = \log \frac{e^{\omega}}{e^{\phi}} \]

We now define

\[ X = \frac{e^{\theta}}{e^{\theta}} \quad Y = \frac{e^{\psi}}{e^{\psi}} \quad Z = \frac{e^{\omega}}{e^{\omega}} = \frac{e^{\phi}}{e^{e^{\phi}}} \quad W = \frac{X}{Y} \] (23)

Using the constraints, we can eliminate \( Z, X, Y \) in favor of \( W \), to get the equations:

\[ \frac{\cos^2 \theta - \sinh^2 \psi}{\cosh^2 \psi} (W - \frac{1}{W}) = -\cot \theta \partial_{\psi} W - \tanh \psi \partial_{\psi} W \] (24)

and

\[ \partial_{\psi} \partial_{\theta} W - \frac{1}{W} \partial_{\psi} W \partial_{\theta} W + \frac{1}{1 - \beta^2} \frac{1}{W} \tanh \psi \partial_{\psi} W \]

\[ -\frac{2\beta^2}{(1 - \beta^2)^2} \tanh \psi \cot \theta (1 - W^2) - \frac{\beta^2}{1 - \beta^2} W \cot \theta \partial_{\psi} W = 0 \] (25)
Any solution is determined by the function $W$, which must satisfy the above two non-linear partial differential equations. The constraint equations then allow one to solve for $X, Y$ and $Z$. With $W$, this solves for all of the vierbeins. The other equations then allow one to solve for $f$ and $G$, which gives the full solution.

Note that the constant solution $W = 1$ was the solution found in [19]. Since the equations are non-linear, it is difficult to find the most general consistent solution to both equations. The equations may be expanded around $W = 1$, and we have checked that, at least to third order in $(W - 1)$, there exists a one-parameter family of solutions. This suggests that the full non-linear equations for $W$ in fact exhibits a consistent one-parameter family of solutions, with $W = 1$ being a point on this line.

To linear order, we may write the solution for $W$ as

$$W = 1 + a(\sinh^2 \psi + \cos^2 \theta)$$

where $a$ is a constant of integration. Solving the equations of the vierbein yields another constant of integration, $b$. We then find to linear order in $(W - 1)$

$$X = 1 - \frac{b}{\sin \theta \cosh \psi} - a \sin^2 \theta$$

$$Y = 1 - \frac{b}{\sin \theta \cosh \psi} - a \cosh^2 \psi$$

$$Z = 1 - \frac{1}{2} \frac{b}{\sin \theta \cosh \psi} + \frac{1}{4} a(\sinh^2 \psi + \sin^2 \theta)$$

$$Re f^2 = Z$$

$$Im \log f = \frac{1}{2} a \sinh \psi \cos \theta$$

$$G^{\tau \omega \theta} = -G^{\phi \phi \chi} = \frac{1}{2} \left[ \frac{b}{\sin^2 \theta \cosh^2 \psi} + a \sin \theta \cosh \psi + \frac{b \cot \theta \tanh \psi}{\sin \theta \cosh \psi} \right]$$

$$G^{\psi \tau \omega} = G^{\phi \chi \theta} = -i + \frac{1}{2} a \sinh \psi \cos \theta - \frac{3}{4} a(\sinh^2 \psi + \sin^2 \theta)$$

In conclusion, we have found an apparently consistent order by order expansion for a two-parameter family of $\frac{1}{2}$-BPS supergravity solutions in an asymptotically $AdS_3 \times S^3$ background.

In this solution, the five-form RR field-strength $F$ vanishes, while the 3-form RR field strength $G$ is turned on. This is a bit puzzling, because
we began with the $\kappa$-symmetry projector for a D3-brane wrapped along an $AdS_2 \times S^2$ submanifold, and imposed that projector on the Killing spinors. This would seem to suggest that the source is wrapped on an $S^2$ of vanishing size (ie., a sphere wrapping the $\phi, \chi$ coordinates, but fixed at $\theta = 0, \pi$). For the case $W = 1$, the solution can be found exactly [19] and this is indeed the case. But for the more general case found above, the source does not appear to be localized at a zero-size sphere (at least to first-order).

We believe that the resolution lies in the fact that the location of the source $(\psi_0, \theta_0)$ dropped out of the expression for the $\kappa$-symmetry projector. In other words, supersymmetry alone will not give us a unique solution, but rather will include all $\frac{1}{2}$-BPS solutions that satisfy this projector, even if they arise from a source different from the D3-brane used in our analysis. It would be very interesting to resolve these issues, as well to explicitly construct this solution to all orders.

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