New Supergravity Solutions for Branes in $AdS_3 \times S^3$

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Abstract

We find explicit supergravity solutions which describe branes in the $AdS_3 \times S^3$ background. These solutions preserve 8 of the 16 supersymmetries of this background, and are consistent with $\kappa$-symmetry. These represent new $\frac{1}{2}$-BPS states of string theory.

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Introduction

One of the most interesting of the anti-de Sitter solutions of string theory is the $AdS_3 \times S^3 \times T^4$ solution which arises from embedding a stack of fundamental strings within a stack of NS-5-branes. This solution is particularly useful because it is possible to write the action of a string in this background as a Wess-Zumino-Witten model on the group manifold $SL(2, R) \times SU(2)$. While the noncompact nature of $SL(2, R)$ leads to subtleties, great progress in understanding the closed string theory has been made [1]. The open string theory on D-branes in this background has also been studied a great deal recently [2].

It has proven extremely fruitful in the AdS/CFT correspondence to be able to describe the same system in two different ways; as a solution to perturbative string theory and as a solution to classical supergravity. If this can be done, we can generate new dual descriptions of gravitational theories by gauge theories. We will therefore try 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 [3].

There have been several partial results in the previous literature for such classical supergravity solutions. Supergravity solutions for general intersecting branes were considered in [4, 5, 6, 7, 8, 9, 10, 11]. Secondly, in [12], brane probes were considered. 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. Since we want to construct the complete nonlinear solution, we will follow the general approach of [5, 7].

There is an important feature of these branes that simplifies our task. It was shown in [13], using boundary state arguments, that a D-3-brane stretched along an $AdS_2 \times S^2$ submanifold of the background satisfies the equations of motion and is $\frac{1}{2}$ BPS; in other words it preserves 8 supercharges. Note that this is twice the number preserved by a system containing D3-branes, fundamental strings and NS-5 branes. The point is that in the near horizon limit, the supersymmetry is enhanced [14]; $AdS_3 \times S^3 \times T^4$ preserves 16 supercharges. A $\frac{1}{2}$-BPS brane in this background therefore preserves 8 supercharges. We therefore expect considerable simplifications since the branes preserve a greater amount of supersymmetry.
In the following section, we analyze $\kappa$-symmetry for the D-branes in $AdS_3 \times S^3$. This analysis tells us which Killing spinors are preserved in the presence of the brane. Specifically, $\kappa$-symmetry informs us that the preserved Killing spinors are found by applying a particular projector to the Killing spinors.

We then use this to analyze the Killing equations. We require that the Killing equations be satisfied once the projector found above is imposed on the spinors. This then imposes constraints on the metric and field strengths. We can then solve these constraints to find the full solution. Since this procedure is rather tedious, we will simplify by assuming that the axion and dilaton are constant (we will look at the more general case in future work.)

We find that the constraints can indeed be solved, and the explicit solution can be found. The sources are found to be localized at antipodal points on the $S^3$ and wrap an $AdS_2$ in $AdS_3$. This then provides a new $\frac{1}{2}$-BPS solution of supergravity.

2 Imposing $\kappa$-Symmetry

The $AdS_3 \times S^3 \times T^4$ solution of type IIB supergravity can be obtained by taking the near-horizon limit of the solution generated by fundamental strings and NS-5-branes. This system preserves 16 supersymmetries. We shall ignore the $T^4$ directions in the subsequent discussions (it is thereby implicitly assumed that all branes are smeared on the $T^4$).

The metric of $AdS_3 \times S^3$ in global coordinates with unit normalized radius is

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

(1)

It was argued in [13] from boundary state considerations that a D-3-brane could be added in a way which preserved one half of the supersymmetries. In this embedding, the geometry of the D3-brane is $AdS_2 \times S^2$. With a particularly simple choice of parameters\(^3\), the D3-brane stretches along the coordinates $(\tau, \omega, \phi, \chi)$, with the coordinates $\psi$ and $\theta$ appearing as transverse scalars. We can then find a brane solution by solving the Born-Infeld equations of motion.

\(^3q = 0\) in the notation of [13].
We need to know the background fields. The relevant part of the $B_{NS}$ background is

\[ \overline{B}_{\phi\chi} = \frac{1}{2} \left( \theta - \frac{\sin 2\theta}{2} \right) \sin \phi \]

\[ \overline{B}_{\tau\omega} = \frac{1}{2} \left( \psi + \frac{\sinh 2\psi}{2} \right) \cosh \omega \]  

(2)

We can also turn on a magnetic flux on the D-3-brane of the form

\[ 4\pi \alpha' F_{\phi\chi} = -\pi p \sin \phi \]  

(3)

where $p$ is constant. The Lagrangian for a D-3-brane embedded in this way is

\[ L_{DBI} = -T_D \sqrt{-\det M} \]  

(4)

where

\[ \sqrt{-\det M} = N(\psi)L(\theta) \cosh \omega \sin \phi \]

\[ L(\theta) = \left( \sin^4 \theta + \left( \theta - \frac{\sin 2\theta}{2} - \pi p \right)^2 \right)^{\frac{1}{2}} \]

\[ N(\psi) = \left( \cosh^4 \psi - \left( \psi + \frac{\sinh 2\psi}{2} \right)^2 \right)^{\frac{1}{2}} \]  

(5)

The solution of the $\theta$ equation of motion is then $\theta_0 = \pi p$. Similarly, we can solve the $\psi$ equation of motion by setting $\psi = 0$.

If we set

\[ a = \frac{1}{2} \sin 2\pi p = \sin \theta_0 \cos \theta_0 \]

\[ b = \sin^2 \pi p = \sin^2 \theta_0 \]  

(6)

then

\[ L(\theta_0) = \sin \theta_0 \]

\[ \mathcal{F}_{\phi\chi} = -\frac{1}{2} \sin 2\theta_0 \sin \phi \]  

(7)

We now discuss the supersymmetries preserved by this brane. For this we need the $\kappa$-symmetry projector $\Gamma$. This is found by [15]

\[ d^{p+1} \xi \Gamma = -e^{-\Phi} L_{DBI}^{-1} e^\mathcal{F} \wedge X|_{vol} \]  

(8)
where

\[ X = \bigoplus_n \Gamma_{(2n)} K^n I, \]

\[ K \psi = \psi^* \]

\[ I \psi = -i \psi \]

\[ \Gamma_n = \frac{1}{n!} d\xi^{i_n} \wedge ... \wedge d\xi^{i_1} \Gamma_{i_1 ... i_n} \]  \hspace{1cm} (9)

We see that

\[ \Gamma = -\frac{1}{\sqrt{a^2 + b^2}} \left( a \gamma_{\tau \omega} K I - b \gamma_{\tau \omega \phi \chi} K^2 I \right) \]

\[ = -i (\cos \theta_0 \gamma_{\tau \omega} K + \sin \theta_0 \gamma_{\tau \omega \phi \chi}) \]  \hspace{1cm} (10)

It may be verified that \( \Gamma \) is traceless and \( \Gamma^2 = 1 \). This implies that 8 of the Killing spinors (pulled back to the worldvolume of the brane) will be invariant under the \( \kappa \)-symmetry projection \( \Gamma \epsilon = \epsilon \).

This is not the end of the story, though. The Killing spinors of \( AdS_3 \times S^3 \) in global coordinates can be written as

\[ \epsilon = \exp \left( \frac{h \psi}{2} \gamma_{\tau \omega} K \right) \exp \left( \frac{h \theta}{2} \gamma_{\phi \chi} K \right) R_0(\phi, \chi, \omega, \tau) \epsilon_0, \]  \hspace{1cm} (11)

where \( \epsilon_0 \) is an arbitrary 16-component constant spinor satisfying \( \gamma^{2345} \epsilon_0 = \epsilon_0 \), \( R_0 \) is invertible and where \( h = -1 \). Since the \( \kappa \)-symmetry projection \( \Gamma \epsilon = \epsilon \) is imposed on the brane, the projection can be written explicitly as

\[ \Gamma \exp \left( \frac{h \theta_0}{2} \gamma_{\phi \chi} K \right) R_0(\phi, \chi, \omega, \tau) \epsilon_0 = \exp \left( \frac{h \theta_0}{2} \gamma_{\phi \chi} K \right) R_0(\phi, \chi, \omega, \tau) \epsilon_0 \]  \hspace{1cm} (12)

Since we want to find the projection on the full spinor \( \epsilon \), we still need to conjugate by the \( (\psi, \theta) \) dependence of the Killing spinor to find the invariant Killing spinor throughout the space.

Defining

\[ \Lambda = \exp \left( \frac{\psi^0}{2} \gamma_{\tau \omega} K \right) \exp \left( \frac{\theta_0}{2} \gamma_{\phi \chi} K \right) \exp \left( -\frac{\psi^0}{2} \gamma_{\tau \omega} K \right) \exp \left( -\frac{\theta}{2} \gamma_{\phi \chi} K \right), \]  \hspace{1cm} (13)
we see we can rewrite the above projection (12) as \( \tilde{\Gamma} \epsilon = \epsilon \) where

\[
\tilde{\Gamma} \equiv \Lambda^{-1} \Gamma \Lambda = -\frac{i}{\sqrt{a^2 + b^2}} (\cos \delta \theta + \sin \delta \theta \gamma_{\phi \chi} K) \times (\cosh \delta \psi + \sinh \delta \psi \gamma_{\tau \omega} K) \gamma_{\tau \omega} (aK + b\gamma_{\phi \chi})
\] (14)

We then have

\[
(1 + \tilde{\Gamma})\epsilon = [1 - i(M + N \gamma_{\tau \omega} K + O \gamma_{\phi \chi} K + P \gamma_{\tau \omega \phi \chi})] \epsilon
\]

\[
M = \frac{1}{\sqrt{a^2 + b^2}} (a \cos \delta \theta \sinh \delta \psi - b \sin \delta \theta \sinh \delta \psi)
\]

\[
N = \frac{1}{\sqrt{a^2 + b^2}} (a \cos \delta \theta \cosh \delta \psi - b \sin \delta \theta \cosh \delta \psi)
\]

\[
O = \frac{1}{\sqrt{a^2 + b^2}} (a \sin \delta \theta \sinh \delta \psi + b \cos \delta \theta \sinh \delta \psi)
\]

\[
P = \frac{1}{\sqrt{a^2 + b^2}} (a \sin \delta \theta \cosh \delta \psi + b \cos \delta \theta \cosh \delta \psi)
\] (15)

which reduces to

\[
M = \sinh \psi \cos \theta \quad N = \cosh \psi \cos \theta
\]

\[
O = \sinh \psi \sin \theta \quad P = \cosh \psi \sin \theta
\] (16)

Note that the \( \theta_0 \) dependence of the space-time projector has dropped out.

For any location \( \theta_0 \) of the brane, the Killing spinors preserved by \( \kappa \)-symmetry are the same.

Using this projection and Hodge duality, one can rewrite the projection as

\[
\gamma_{\tau \omega} K \epsilon = (A + B \gamma_{\phi \theta}) \epsilon
\] (17)

where

\[
A = -\frac{\cosh \psi}{(\cosh^2 \psi - \sin^2 \theta)} (\sinh \psi + i \cos \theta)
\]

\[
B = \frac{i \sin \theta}{(\cosh^2 \psi - \sin^2 \theta)} (\sinh \psi + i \cos \theta)
\] (18)

\[4\delta \theta = \theta - \theta_0. \] To connect to more general embeddings, we also write \( \delta \psi = \psi - \psi_0 \), although in this case \( \psi_0 = 0 \).
The Killing equations

We will impose the projection $\frac{1}{2} \left( 1 + \tilde{\Gamma} \right) \epsilon = \epsilon$ and demand that the spinors which satisfy it are solutions of the Killing equations. This will generate a $\frac{1}{2}$-BPS solution.\(^5\) Note that because the choice of Killing spinor is independent of $\theta_0$ (and presumably $\psi_0$), the solution may correspond to a smeared brane, and not one necessarily localized at the place originally anticipated.

The Killing equations are of the form

\[
\partial_\mu \epsilon - \frac{1}{4} \omega^a_{\mu b} \gamma_{ab} + \frac{i}{192} \Gamma^b_{\mu cde} \gamma_{b cde} \epsilon^* - \frac{i}{48} e^\Phi (G_{\mu a b c} - 9 G_{\tilde{\mu} a b}) \epsilon^* = 0
\] (19)

For example, the $\psi$ Killing equation is\(^6\)

\[
\partial_\psi f \left( - \frac{1}{2} \omega_{\psi \theta} \gamma_\theta \epsilon - \frac{i}{8} F^{\tau \omega \phi \chi} \gamma_{\tau \omega \phi \chi} \epsilon + \frac{i}{8} \Gamma^{\psi \tau \omega \phi} \gamma_{\psi \tau \omega \phi} \epsilon \right.
\]

\[
+ 3 \frac{i}{8} e^\Phi G^{\psi \tau \omega} \gamma_{\psi \tau \omega} \epsilon^* - \frac{i}{8} e^\Phi G^{\tau \omega \phi} \gamma_{\tau \omega \phi} \epsilon^* + \frac{3 i}{8} e^\Phi G^{\psi \chi} \gamma_{\psi \chi} \epsilon^*
\]

\[
\left. - \frac{i}{8} e^\Phi G^{\phi \chi} \gamma_{\phi \chi} \epsilon^* \right) = e^\tilde{\psi} \left( - \frac{1}{2} \omega^\psi \gamma_\theta \epsilon + \frac{3 i}{8} e^\Phi G^{\psi \tau \omega} \gamma_{\psi \tau \omega} \epsilon^* - \frac{i}{8} e^\Phi G^{\phi \chi} \gamma_{\phi \chi} \epsilon^* \right) \] (20)

where the bars over expressions indicate that the expression is to be evaluated in the unperturbed background $AdS$ solution and where we have made the ansatz $\epsilon = f(\psi, \theta) \bar{\epsilon}$. The other Killing equations are entirely similar.

The axion-dilaton equation is

\[
\frac{1}{4} \partial_\psi \Phi \gamma_\psi \epsilon + \frac{1}{4} \partial_\theta \Phi \gamma_\theta \epsilon + \frac{i}{8} e^\Phi G^{\omega \psi \tau} \gamma_{\omega \psi \tau} \epsilon^* + \frac{i}{8} e^\Phi G^{\omega \tau \theta} \gamma_{\omega \tau \theta} \epsilon^* + \frac{i}{8} e^\Phi G^{\omega \psi \chi} \gamma_{\omega \psi \chi} \epsilon^* + \frac{i}{8} e^\Phi G^{\omega \chi \theta} \gamma_{\omega \chi \theta} \epsilon^* = 0
\] (21)

\(^5\)We assume that no field strengths have indices along the $T^4$ unless they have indices in all those directions, and can thus be related to a field with no such indices by Hodge duality.

\(^6\)\(\tilde{\mu}\) is a curved space-time index, while $\mu$ is a tangent-space index. We use the notation $\omega^\mu_{\mu \nu} = e^\mu_{\tilde{\mu}} \omega_{\mu \nu}$. 
We shall now assume that $\Phi = 0$. Then (21) is solved by the ansatz
\[
\begin{align*}
\Phi &= 0 \\
G^{\psi \tau \omega} &= G^{\phi \chi \theta} \\
G^{\tau \omega \theta} &= -G^{\psi \phi \chi}
\end{align*}
\tag{22}
\]

We can now return to the other Killing equations. In each of them, we first impose the projector (17). The Killing equation then becomes a matrix equation which is to be satisfied identically. Thus the complex coefficient of each matrix must be zero. In this way, each Killing equation leads to two complex algebraic equations.

For example, the $\psi$ Killing equation, after this procedure, leads to the equations
\[
\frac{\partial \psi}{f} + \frac{i}{8} F^{\tau \omega \phi \chi \theta} + \frac{i}{8} e^\Phi \left(3G^{\psi \tau \omega} A + G^{\tau \omega \theta} B - 3G^{\psi \phi \chi} B + G^{\phi \chi \theta} A\right) = \\
+ \frac{3i}{8} e^\Phi \bar{G}^{\psi \tau \omega} A + \frac{i}{8} e^\Phi \bar{G}^{\phi \chi \theta} A
\tag{23}
\]

and
\[
-\frac{1}{2} \omega^\psi_{\psi} + \frac{i}{8} F^{\psi \tau \omega \phi \chi} + \frac{i}{8} e^\Phi \left(3G^{\psi \tau \omega} B - G^{\tau \omega \theta} A + 3G^{\psi \phi \chi} A + G^{\phi \chi \theta} B\right) = \\
- \frac{1}{2} \bar{\omega}^\psi_{\psi} + \frac{i}{8} e^\Phi \bar{G}^{\psi \tau \omega} B + \frac{i}{8} e^\Phi \bar{G}^{\phi \chi \theta} B
\tag{24}
\]

By taking linear combinations of these equations such that the barred field strengths cancel, we obtain the equations
\[
\frac{e^0_\theta}{e^0_\tau} \frac{e^\phi_\phi}{e^\phi_\phi} = S(\theta) \quad \frac{e^\psi_\psi}{e^\psi_\psi} \frac{e^\omega_\omega}{e^\omega_\omega} = T(\psi)
\tag{25}
\]

and
\[
2 \frac{\partial \psi}{f} + \omega^\phi_{\phi} = 2 \frac{\partial \theta}{f} + \omega^\theta_{\omega} = 0,
\tag{26}
\]

which tells us that
\[
\bar{e}^\phi_\phi \bar{e}^\phi_\phi = \bar{e}^\phi_\phi \bar{e}^\phi_\phi
\tag{27}
\]
One can substitute all of this back into the Killing equations to derive further relations involving the barred quantities, such as

\[ \frac{i}{8} e^\Phi (3G^\tau\omega + G^\phi\chi) \left( \frac{e^\omega}{e^\tau} - \frac{e^\omega}{e^\phi} \right) A = \frac{\partial_\psi f}{f} + \frac{1}{2} \omega_\psi \omega - \frac{1}{2} \omega_\omega \omega \right) \]

\[ + \frac{i}{4} F^\tau\omega\psi\chi + \frac{i}{2} e^\Phi B (G^\tau\omega - G^\psi\phi\chi). \quad (28) \]

These relations allow us to determine the entire solution, including the fact that \( S(\theta) = T(\psi) = \text{const}. \)

Defining \( \gamma = \sin \theta \cosh \psi = \beta \cosh^2 \psi, \) the complete solution for the metric is found to be

\[ ds^2 = H^\frac{1}{2} d\theta^2 + H^{-\frac{1}{2}} \sin^2 \theta (d\phi^2 + \sin^2 \phi d\chi^2) + H^\frac{1}{2} d\psi^2 + H^{-\frac{1}{2}} \cosh^2 \psi (d\omega^2 - \cosh^2 \omega d\tau^2), \]

\[ H = \left( \frac{c\gamma + 1}{1 + c\gamma} \right)^2. \quad (29) \]

and the field strengths are

\[ F^\tau\omega\psi\chi = F^\phi\psi\omega\chi = 0, \]

\[ G^\tau\omega = -G^\psi\phi\chi = \frac{1}{2} e^{\Phi} \left( \frac{1}{c\gamma(1 + c\gamma)} \left( \frac{\beta}{\sin^2 \theta} + i \cot \theta \tanh \psi \right) \right) \]

\[ G^\psi\tau\omega = G^\phi\phi\chi = -i \sqrt{\frac{1 + c\gamma}{c\gamma}} \left( 1 - \frac{1}{2} \frac{1}{1 + c\gamma} \right). \quad (30) \]

where \( c \) is a constant of integration.

One may verify that the Bianchi identity \( dG = 0 \) is satisfied away from \( \theta = 0. \) We may thus write \( G = dC, \) where

\[ C_{\tau\omega} = \frac{i}{2} \left( \psi + \sinh 2\psi + \frac{\sinh \psi \cosh \omega}{c \sin \theta} \right) + \cosh \omega \cot \theta \]

\[ C_{\phi\chi} = \frac{i}{2} \left( \theta - \frac{\sin 2\theta}{2} + \frac{\cos \theta \sin \phi}{c \cosh \psi} \right) - \frac{\sin \phi}{2c} \tanh \psi \quad (31) \]

The gauge fields are singular at \( \theta = 0, \pi, \) indicating a source there. The net D1-brane charge may then be found up to a normalization constant to
be
\[ Q = Re \int d\psi d\phi d\chi G \tilde{\psi} \tilde{\phi} \tilde{\chi} = \frac{2\pi}{c} \int d\psi \frac{1}{\cosh^2 \psi} = \frac{4\pi}{c} \] (32)

It is also easily seen that the 3-brane charge is zero. The sources can thus be interpreted as string-like objects wrapping the \((\tau, \omega)\) directions in \(AdS_3\), and sitting at antipodal points \(\theta = 0, \pi\) on the \(S^3\).

To conclude, we have found a new solution of supergravity, representing a brane in an \(AdS_3 \times S^3\) background, which preserves 8 supersymmetries. The preserved killing spinor is independent of the embedding parameters \(\theta_0, \psi_0\), demonstrating that the sources can be smeared or superposed without breaking additional supersymmetry.

The specific case we have analyzed has no 3-brane charge. It is likely that generalizing our ansatz to more general cases with nonconstant scalar fields will produce solutions involving three-branes as well. We will explore this possibility in future work [16]. It will also be interesting to analyze the implications for the dual defect conformal field theory.

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4 Appendix: Notation

In [17] \(\epsilon\) is a 32-component complex spinor in mostly - signature, whereas [18] uses \(\epsilon_{L,R}\), each of which is a 16-component Majorana spinor in mostly + signature.

Considering the case of Type IIB, if we choose the \(\gamma\)'s to be real and define
\[ \epsilon_R = Re \left( \frac{1 \pm \gamma^{11}}{2} \epsilon \right) \]
\[ \epsilon_L = Im \left( \frac{1 \pm \gamma^{11}}{2} \epsilon \right) \]
\[ \gamma^\mu = i \Gamma^\mu \] (33)

then the conventions are consistent.

We use the notation of [17], but will switch to mostly + signature. In the brane action, we use the same sign convention as [19].
References

[1] J. M. Maldacena and H. Ooguri, Phys. Rev. D 65, 106006 (2002) [arXiv:hep-th/0111180].
J. M. Maldacena, H. Ooguri and J. Son, J. Math. Phys. 42, 2961 (2001) [arXiv:hep-th/0005183].
J. M. Maldacena and H. Ooguri, J. Math. Phys. 42, 2929 (2001) [arXiv:hep-th/0001053].

[2] A. Rajaraman and M. Rozali, “Boundary states for D-branes in AdS(3),” Phys. Rev. D 66, 026006 (2002) [arXiv:hep-th/0108001].
B. Ponsot, V. Schomerus and J. Teschner, “Branes in the Euclidean AdS(3),” JHEP 0202, 016 (2002) [arXiv:hep-th/0112198].
P. Lee, H. Ooguri and J. w. Park, “Boundary states for AdS(2) branes in AdS(3),” Nucl. Phys. B 632, 283 (2002) [arXiv:hep-th/0112188].
P. Lee, H. Ooguri, J. w. Park and J. Tannenhauser, “Open strings on AdS(2) branes,” Nucl. Phys. B 610, 3 (2001) [arXiv:hep-th/0106129].
A. Rajaraman, “New AdS(3) branes and boundary states,” arXiv:hep-th/0109200
A. Giveon, D. Kutasov and A. Schwimmer, “Comments on D-branes in AdS(3),” Nucl. Phys. B 615, 133 (2001) [arXiv:hep-th/0106005].
T. Quella, “On the hierarchy of symmetry breaking D-branes in group manifolds,” JHEP 0212, 009 (2002) [arXiv:hep-th/0209157].
S. Ribault, “Two AdS(2) branes in the Euclidean AdS(3),” arXiv:hep-th/0210248

[3] A. Karch and L. Randall, “Locally localized gravity,” JHEP 0105, 008 (2001) [arXiv:hep-th/0011156].
A. Karch and L. Randall, “Open and closed string interpretation of SUSY CFT’s on branes with boundaries,” JHEP 0106, 063 (2001) [arXiv:hep-th/0105132].
O. DeWolfe, D. Z. Freedman and H. Ooguri, “Holography and defect conformal field theories,” Phys. Rev. D 66, 025009 (2002) [arXiv:hep-th/0111135].
C. Bachas, J. de Boer, R. Dijkgraaf and H. Ooguri, “Permeable conformal walls and holography,” JHEP 0206, 027 (2002) [arXiv:hep-th/0111210].
J. Erdmenger, Z. Guralnik and I. Kirsch, “Four-dimensional superconformal theories with interacting boundaries or defects,” Phys. Rev. D
N. R. Constable, J. Erdmenger, Z. Guralnik and I. Kirsch, “Intersecting D3-branes and holography,” arXiv:hep-th/0211222.

N. R. Constable, J. Erdmenger, Z. Guralnik and I. Kirsch, “(De)constructing intersecting M5-branes,” arXiv:hep-th/0212136.

[4] A. Rajaraman, “Comments on D-branes in flux backgrounds,” arXiv:hep-th/0208085.

[5] A. Rajaraman, “Supergravity duals for N = 2 gauge theories,” JHEP 0210, 009 (2002) arXiv:hep-th/0011279.

[6] A. Rajaraman, “Supergravity solutions for localised brane intersections,” JHEP 0109, 018 (2001) arXiv:hep-th/0007241.

[7] A. Fayyazuddin and D. J. Smith, “Warped AdS near-horizon geometry of completely localized intersections of M5-branes,” JHEP 0010, 023 (2000) arXiv:hep-th/0006060.

[8] B. Brinne, A. Fayyazuddin, T. Z. Husain and D. J. Smith, “N = 1 M5-brane geometries,” JHEP 0103, 052 (2001) arXiv:hep-th/0012194.

[9] B. Brinne, A. Fayyazuddin, S. Mukhopadhyay and D. J. Smith, “Supergravity M5-branes wrapped on Riemann surfaces and their QFT duals,” JHEP 0012, 013 (2000) arXiv:hep-th/0009047.

[10] S. A. Cherkis and A. Hashimoto, “Supergravity solution of intersecting branes and AdS/CFT with flavor,” arXiv:hep-th/0210105.

[11] O. Lunin and S. D. Mathur, “AdS/CFT duality and the black hole information paradox,” Nucl. Phys. B 623, 342 (2002) arXiv:hep-th/0109154.

O. Lunin, S. D. Mathur and A. Saxena, “What is the gravity dual of a chiral primary?,” arXiv:hep-th/0211292.

[12] C. Bachas, M. R. Douglas and C. Schweigert, “Flux stabilization of D-branes,” JHEP 0005, 048 (2000) arXiv:hep-th/0003037.

A. Alekseev and V. Schomerus, “RR charges of D2-branes in the WZW model,” arXiv:hep-th/0007096.

A. Y. Alekseev, A. Recknagel and V. Schomerus, “Brane dynamics in background fluxes and non-commutative geometry,” JHEP 0005,
C. Bachas and M. Petropoulos, “Anti-de-Sitter D-branes,” JHEP 0102, 025 (2001) [arXiv:hep-th/0012234].

R. Kallosh and J. Kumar, “Supersymmetry enhancement of D-p-branes and M-branes,” Phys. Rev. D 56, 4934 (1997) [arXiv:hep-th/9704189].

M. Cederwall, A. von Gussich, B. E. Nilsson and A. Westerberg, “The Dirichlet super-three-brane in ten-dimensional type IIB supergravity,” Nucl. Phys. B 490, 163 (1997) [arXiv:hep-th/9610148].

M. Aganagic, C. Popescu and J. H. Schwarz, “D-brane actions with local kappa symmetry,” Phys. Lett. B 393, 311 (1997) [arXiv:hep-th/9610249].

M. Cederwall, A. von Gussich, B. E. Nilsson, P. Sundell and A. Westerberg, “The Dirichlet super-p-branes in ten-dimensional type IIA and IIB supergravity,” Nucl. Phys. B 490, 179 (1997) [arXiv:hep-th/9611159].

E. Bergshoeff and P. K. Townsend, “Super D-branes,” Nucl. Phys. B 490, 145 (1997) [arXiv:hep-th/9611173].

M. Aganagic, C. Popescu and J. H. Schwarz, “Gauge-invariant and gauge-fixed D-brane actions,” Nucl. Phys. B 495, 99 (1997) [arXiv:hep-th/9612080].

E. Bergshoeff, R. Kallosh, T. Ortin and G. Papadopoulos, “Kappa-symmetry, supersymmetry and intersecting branes,” Nucl. Phys. B 502, 149 (1997) [arXiv:hep-th/9705040].

K. Skenderis and M. Taylor, JHEP 0206, 025 (2002) [arXiv:hep-th/0204054].

J. Kumar and A. Rajaraman, work in progress.

E. Bergshoeff, “p-branes, D-branes and M-branes,” [arXiv:hep-th/9611099]

J. M. Maldacena, “Black holes in string theory,” [arXiv:hep-th/9607235]

R. Kallosh, J. Kumar and A. Rajaraman, “Special conformal symmetry of worldvolume actions,” Phys. Rev. D 57, 6452 (1998) [arXiv:hep-th/9712073].