On D-Branes and Black Holes in Four Dimensions

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Abstract

We find extremal four dimensional black holes with finite area constructed entirely from intersecting D-branes. We argue that the microscopic degeneracy of these configurations agrees with the Bekenstein-Hawking entropy formula. The absence of solitonic objects in these configurations may make them useful for dynamical studies of black holes.

1 Introduction

In recent months there has been considerable progress in accounting for the microscopic degeneracy of black holes using the explicit conformal field theories provided by D-branes [1]. Examples of black holes in four and five dimensions have been constructed for which the degeneracy of microscopic D-brane states matches the Bekenstein-Hawking entropy [2] - [8]. The four dimensional configurations considered so far contain some combination of D-branes and solitons and seem unwieldy for investigation of dynamical phenomena. D–branes are presently understood far better than solitons in general; so it may be an advantage to construct four dimensional black holes entirely from D–branes. In this paper we present some aesthetically pleasing configurations that realize this, and discuss their microscopic entropy.

The paper is organised as follows. In Section 2 we discuss classical black hole solutions to the IIA and IIB supergravities and find their area. Regularity conditions on the classical p-brane solutions lead us to the most general construction of black holes from four intersecting D–branes of Type II on $T^6$. In Section 3 we discuss the microscopic degeneracy of these systems of branes and their relationship to intersecting brane configurations of M-theory.

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2 Brane Surgery

We are interested in configurations of D–branes wrapped on $T^6$ whose four dimensional manifestations are regular extremal black holes. To achieve this, the dilaton and the moduli (the metric of the compactified dimensions) must be finite at the horizon in order to avoid large corrections to the low energy solution. It is also important that the black hole preserves some supersymmetry so that quantum corrections are well controlled. We will consider systems of four intersecting D–branes. This is the smallest number for which all these requirements can be met. The configurations in this paper are dual to the black holes of Cvetiˇc and Youm [9].

The string metric and the dilaton of a single D–brane are [10]:

$$
\begin{align*}
    ds^2 &= F^{-1/2} \left(-dt^2 + dx_1^2 + \ldots + dx_p^2\right) + F^{1/2} \left(dx_{p+1}^2 + \ldots + dx_9^2\right) \\
    e^{-2\phi} &= F^{(p-3)/2}
\end{align*}
$$

The profile function $F$ is a solution of the Laplace equation, so that after compactification to 4 dimensions $F = 1 + \frac{\mathcal{A}}{r}$ for all branes. Here $r^2 = x_4^2 + x_5^2 + x_6^2$ is the radius squared in the uncompactified dimensions. The metric and the dilaton for the intersecting branes of interest here can be found by multiplying the profile functions for each of the branes [11, 12]. The complete solution obtained this way has a horizon at $r = 0$.

The black hole must have a finite dilaton at the horizon. For a single D–brane the dilaton is finite at the horizon only for $p = 3$. For multiple intersecting D–branes the $p$’s must be chosen in such a way that the product of the profile functions tends to a finite value at the horizon. For four intersecting D–branes it is necessary and sufficient that the $p$’s add up to 12. For example, one 6D-brane and three 2D-branes satisfy this criterion.\footnote{We do not consider systems of 2 and 3 intersecting branes because they cannot satisfy all three regularity conditions described in this section.}

Regular black holes must also have finite moduli at the horizon. The profile function $F$ diverges in the same way for all branes, and the dimensions parallel and transverse to a brane are multiplied by $F^{-1/2}$ and $F^{1/2}$ respectively. So each compact dimension must be perpendicular to as many branes as it is parallel to. Four D–branes can satisfy this criterion if each dimension is in the worldvolume of exactly two D–branes. For example, four 3-branes wrapped around the cycles (123), (345), (146) and (256) satisfy both the finite dilaton and finite moduli conditions.

Each D–brane imposes a boundary condition that relates the components of space-time spinors. This identifies the two supersymmetries in 10 dimensions in a way that depends on the orientation of the D–brane. When two D–branes are present we get two boundary conditions. The first identifies the two supersymmetries and the second leads to the projection condition:

$$
\mathcal{Q} = \pm \Gamma \mathcal{Q} \quad ; \quad \Gamma = \Gamma_{a_1} \cdots \Gamma_{a_n}
$$

on the remaining generator in 10 dimensions. The sign in the relation parametrises the distinction between branes and antibranes and is purely conventional. The indices
a_i run over values that are present for one of the branes but not for the other (i.e. (ND) directions). The condition that this is indeed a projection ($\Gamma^2 = 1$) so that there is some surviving supersymmetry, requires that a p-brane intersecting a q-brane along $k$ dimensions satisfies $p + q - 2k = 0 \text{ mod } 4$.

In addition to this condition on each pair, we must consider the complete system of 4-branes. One of the branes identifies the two supersymmetries and the other three impose three projection conditions of the type in Eq. 2. Without loss of generality, the projection operators for four intersecting branes can be taken to be $\pm \Gamma_1 \Gamma_2 \Gamma_3 \Gamma_4$, $\pm \Gamma_1 \Gamma_2 \Gamma_5 \Gamma_6$, and $\mp \Gamma_3 \Gamma_4 \Gamma_5 \Gamma_6$. These matrices commute so that all the different conditions are compatible. Each of the 4 branes imposes one condition and therefore breaks 1/2 of the supersymmetry. However, one of these conditions is redundant - two of the three projection operators multiply to give the third, so the complete state only breaks 1/8 of the supersymmetry.

Note that one of the three projections has a sign that is determined by the others. Therefore the signs in Eq. 2 can not be chosen independently for each of the 4 D–branes. There are $2^4 = 16$ distinct assignments of orientation for the D–branes but only 8 of them lead to 1/8 supersymmetric configurations.

As a final condition recall that a given configuration can only contain even branes in IIA or odd branes in IIB. All the conditions taken together are very restrictive. For example, assume that the configuration includes a 6–brane wrapped around the $T^6$. Then the supersymmetry condition requires that only 2–branes or 6–branes may be intersected with the 6–brane. The dilaton condition dictates that there must be three 2–branes, and the moduli condition requires that the 2–branes should lie in orthogonal dimensions. In this way we get a complete classification of possible configurations of intersecting D–branes that form regular black holes with finite area in 4 dimensions.

Let 1, $\cdots$, 6 denote the dimensions of the $T^6$ and use p–tuples for worldvolume coordinates of a D–brane. For example, $(12)$ is a 2D–brane aligned along the first and second compactified dimension. In this notation, the moduli are stabilised when each of the six digits enters exactly twice. The regularity conditions discussed above yield the configurations:

\[
(123), (345), (146), (256) \\
(1234), (3456), (1256), (1) \\
(1234), (3456), (12), (56) \\
(12345), (126), (346), (5) \\
(123456), (12), (34), (56)
\]

The second configuration was mentioned in [13].

Recall that T–duality along a given direction acts on the branes by adding the corresponding index if it is not there already, and removing it, if it is. Then it is easy to see that the examples given are in fact T–dual to each other. In this sense, there is a unique regular supersymmetric black hole in 4 dimensions that can be made out of four D–branes. It is quite remarkable that such symmetric configurations exist at all. We find it intriguing that this possibility is special to four dimensions.
The configuration consisting of 4 3D–branes is particularly symmetric. The metric is:

\[ ds^2 = \left( \frac{F_1 F_2 F_3 F_4}{F_1^2 - F_2^2 - F_3^2 - F_4^2} \right)^{-1/2} (\frac{F_1 F_2 F_3 F_4}{F_1^2 - F_2^2 - F_3^2 - F_4^2})^{1/2} (dt^2 + dx^2_7 + dx^2_8 + dx^2_9) + \left[ \left( \frac{F_2 F_3}{F_1 F_4} \right)^{1/2} dx^2_1 + \left( \frac{F_1 F_3}{F_2 F_4} \right)^{1/2} dx^2_2 + \left( \frac{F_1 F_4}{F_2 F_3} \right)^{1/2} dx^2_3 + \left( \frac{F_2 F_4}{F_1 F_3} \right)^{1/2} dx^2_4 \right] \] (4)

For 3D–branes the dilaton is constant so there is no need to distinguish between Einstein metric and string metric. The four dimensional area is

\[ A = 4\pi \sqrt{q_1 q_2 q_3 q_4} \] (5)

where the \( q \)'s are the charges of the four different branes.

For this configuration, the volume of the compact space is independent of radial distance from the horizon. Therefore the area of the black hole in the ten dimensional theory is given by Eq. 5, multiplied by the constant internal volume. Also recall that \( G_N \) in the four dimensional theory differs from the one appropriate to the ten dimensional theory by the volume of the compactified space, measured at infinity. Therefore the entropy of this black hole:

\[ S = \frac{A}{4G_N} \] (6)

is the same when \( A \) and \( G_N \) are understood either from the ten dimensional perspective or in the compactified theory.

The form of the entropy that is suitable for comparison with microscopic considerations is the expression in terms of integer quanta of the D–brane charges. The appropriate formula is

\[ S = 2\pi \sqrt{Q_1 Q_2 Q_3 Q_4} \] (7)

Here the notation \( Qi \) denotes the integer quanta of the four branes. This can be shown from saturation of Dirac’s quantization condition. Despite this topological origin the derivation known at present involves explicit consideration of the normalization of charges [14].

The entropy is a pure number independent of both the moduli of the torus and the string coupling. This is a necessary condition for any counting to work [13, 16]. Moreover, the Einstein area is invariant under T–duality so that same entropy formula is relevant to all the configurations considered here. In fact, the entropy formula can be uniquely extended to the full U–duality group \( E(7) \) [14].

3 Microscopic Entropy

The regularity requirements imposed in the previous section allowed us to identify collections of D-branes that act as black holes with finite area. The specific classical
solution written down in Eq. 4 was chosen to have translational symmetry in the compact dimensions. Such choices regarding the details of the configuration in the internal space are not unique - there is a spectrum of degenerate microscopic states consistent with the choice of charges measurable at infinity. In this section we examine the microscopic D–brane derivation of the corresponding entropy.

### 3.1 Intersecting D–branes

Let us briefly recall the main features of some successful D–brane state counts. As an example, consider $N$ 0-branes on $M$ parallel, intersecting 4-branes. The 0-branes are described as instantons of the 4-brane SU(M) world-volume theory. Consequently, each 0-brane comes equipped with $4M$ degrees of freedom associated with orientations of the instantons of SU(M). A given classical configuration breaks $4NM$ symmetries because it defines a point in the moduli space. The corresponding Goldstone modes (and their superpartners) are the relevant elementary excitations leading to the entropy. They are in one-to-one correspondence with the strings on the intersection manifold that run between branes and bind them together. This correspondence is one of the attractive features of D–branes that allows a simple counting of the Goldstone modes.

We are interested in multiple intersecting D-branes that are not parallel. We expect a non-trivial interacting theory on the intersection manifold and therefore it appears to be difficult to count the number of zero modes explicitly. Presumably the intersecting configurations are described by certain classical excitations of the branes can be “oriented” relative to each other, and these “orientations” are described by the zero modes of the string condensate binding the branes together.

### 3.2 4-4-4-0

The Type IIA 4-4-4-0 configuration has $Q_1$, $Q_2$ and $Q_3$ 4-branes wrapped around the (1234), (3456) and (1256) cycles of $T^6$. Let us first consider the special case $Q_1 = Q_2 = Q_3 = 1$. The three 4-branes on $T^6$ intersect at exactly one point. Now bind 0-branes to this mutual intersection point by turning on a background of strings running between the 0-branes and each of the 4-branes. This can be thought of as a condensate of strings that describes orientation degrees of freedom of the 0-brane on the intersecting 4-branes. A 0–brane bound to the intersection point of three 4-branes can carry any integer multiple of the 0-brane RR quantum. These multiply charged 0-brane states arise as fundamental zero-branes bound together by giving an expectation value to the massless strings running between them. (This picture follows naturally from M-theory as we will describe in the next section.) Such 0-brane bound states should be counted as separate sectors of the Hilbert space in counting the degeneracy of a configuration.

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2 For a discussion of some of the spacetime aspects of this point see [17].

3 The twisted sectors of the moduli space of identical 0-branes on 4-branes ([18, 19]) can be interpreted as precisely such bound states, and they are indeed counted as separate states in the Hilbert space.
Let us suppose that a 0-brane attached to the intersecting 4-branes has \( k \) bosonic orientation degrees of freedom. Then the unbroken supersymmetry implies an equal number of fermionic orientations. Define \( a_{\mu}^n \) (\( \mu = 1 \cdots k \)) to be an operator that creates a charge \( n \) 0-brane with bosonic orientation \( \mu \) at the intersection point. Let \( b_{\mu}^n \) be the corresponding fermionic operator. Then the zero-brane charge operator is \( Q = \sum_{n=1}^{\infty} na_{\mu}^n a_{\mu}^n + \sum_{n=1}^{\infty} nb_{\mu}^n b_{\mu}^n \). The generating function of the degeneracy of states with charge \( n \) is then given by the familiar formula:

\[
\sum_n d(n)q^n = \frac{\prod_n (1 + q^n)^k}{\prod_n (1 - q^n)^k}
\]  

(8)

The degeneracy of states with three 4-branes and total 0-brane charge \( Q_4 \) is \( d(Q_4) = \exp 2\pi \sqrt{(k/4)} Q_4 \) for \( Q_4 \gg 1 \)

Let us return to the general case where \( Q_1, Q_2 \) and \( Q_3 \) are greater than unity. Then each of the three varieties of 4–branes consists of a number of branes that are parallel but need not coincide. Separating the branes gives \( Q_1 Q_2 Q_3 \) distinct spacetime points where there are three intersecting branes. Attach a total 0-brane charge of \( Q_4 \) to these intersection points. Each of the \( Q_1 Q_2 Q_3 \) intersection points contributes \( k \) bosonic and \( k \) fermionic modes. As before introduce creation operators describing the charge \( n \) 0-brane states, except that the index \( \mu \) on the operators \( a \) and \( b \) now runs between 1 and \( k Q_1 Q_2 Q_3 \). The degeneracy of states becomes

\[
d(Q_4) = \exp 2\pi \sqrt{(k/4)} Q_1 Q_2 Q_3 Q_4 \text{ for large } Q_4 \text{ and the entropy is}
\]

\[
S = \ln d = 2\pi \sqrt{(k/4)} Q_1 Q_2 Q_3 Q_4
\]  

(9)

This agrees with the Bekenstein-Hawking entropy calculated in Section 2 when \( k = 4 \). We therefore expect that a 0-brane bound to the intersection point of three 4-branes has 4 bosonic and 4 fermion zero modes.

It can be understood heuristically why \( k = 4 \): two 4–branes break 3/4 of the 32 real spacetime supersymmetries; so the worldvolume theory on each intersection manifold respects the remaining 8. Half of these are broken when another D–brane is attached and the Goldstone modes associated with this final breaking can condense to bind the 4–branes. There is a separate condensate at each intersection point that comprises 4 fermionic modes and, by virtue of the unbroken worldvolume supersymmetry, 4 bosonic modes. We will motivate in the following section the expectation that excitations of the condensate can carry 0–brane RR charge. As shown in Section 2, the 0–brane RR charge does not break any additional supersymmetry; so it is expected that \( k = 4 \).

This argument is essentially that of [21]. It also works for the five dimensional black holes described by a 1-brane and a 5-brane: the worldvolume theory of the 5-brane has 16 supersymmetries, the 1-brane breaks 1/2 of them, and 8 Goldstone

\footnote{Collective coordinates were attributed to supersymmetries broken by one of the three 4–branes but respected by the others. A more democratic treatment might have suggested a three-fold duplication of these Goldstone modes, but we presume that the correct construction effectively identifies these copies.}
modes appear. Excited states of the binding condensate can carry momentum but, for this to preserve some of the remaining supersymmetry, it is only half of the modes that effectively participate. This leaves $k = 4$ orientations.

### 3.3 M-theory and 4d Black Holes

A number of issues regarding the state-counting for the 4-4-4-0 configuration are clarified from an M-theory perspective. Compactify M-theory on a circle and let 0 denote the compact dimension. Then the three 4-branes arise as compactifications of the 5-branes $(01234)$, $(01256)$ and $(03456)$. The RR 1-form gauge field is the Kaluza-Klein gauge field associated with the 11-dimensional metric. The 0-brane RR charge descends from the quantized momentum along the compact dimension and leads to multiply charged 0-branes in 10 dimensions. These multiply charged states are not just two juxtaposed fundamental 0-branes, because a mode carrying momentum $2n$ in the 0 direction is physically distinct from two modes carrying momentum $n$ each.

In the 4-4-4-0 counting the enormous degeneracy arose because of the many mutual intersection points. This simplification is only suitable for the limit where $Q_4$ is much bigger than $Q_1$, $Q_2$ and $Q_3$. Indeed, the 4-branes may bind, just like the 0-branes; so we may expect states with multiple 4-brane charge that would reduce the number of intersection points and threaten the state counting. The problem is resolved by noting that there are two ways to obtain a bound state of 4-branes from M-theory. The naive way is to bind $Q_1$ 5-branes together and dimensionally reduce to get a bound state of $Q_1$ 4-branes. Alternatively, we can wind a 5-brane $Q_1$ times around the 0 direction and dimensionally reduce to get a bound state of 4-branes. In this case, the effective length of the circle on which the Kaluza-Klein modes live is multiplied by $Q_1$ leading to 0-brane charges quantized in units of $1/Q_1$. Admitting these fractionally charged 0-branes gives the correct state count [22, 23]. In our case, this mechanism has to employed for each of the three intersecting varieties of 5-branes, leading to Kaluza-Klein modes quantized in units of $1/(Q_1Q_2Q_3)$ (for relatively prime $Q_1$, $Q_2$ and $Q_3$). When all four charges are big [23] it is essential to consider these additional states in order to match the Bekenstein-Hawking formula. The “fractional charge” required here is the D-brane analogue of the “tension renormalization” of [13].

M-theory also clarifies the nature of the string condensate binding the 4-4-4-0 configuration together. M-theory 5-branes interact via exchange of membranes with boundaries on the 5-branes [25, 26, 27], and quantization of the boundary states yields the 5-brane low-energy effective theory [21, 28]. This suggests that collapsed membranes live on the intersection manifold of our intersecting 5-branes and bind them together [24]. Dimensional reduction of these collapsed membranes would give the desired string condensate and the momentum of the membranes turns into 0-brane charge. It is suggested in [24] that the collapsed membranes act like a single self-dual string in 6 dimensions with 4 bosonic and fermionic zero modes. This would give $k = 4$ in the 4-4-4-0 counting, but the rules for state counting in M-theory are

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5 Similar arguments have also appeared recently in [24].
not yet firmly established and more work is necessary on this point.

4 Comments on Other Configurations

4.1 3-3-3-3

The 4-4-4-0 configuration is T-dual to the more symmetric IIB 3-3-3-3 configuration. There is a heuristic counting of states for this configuration that is exactly parallel to the IIA setup. Once again, take Q1, Q2 and Q3 3-branes wrapped around the (123), (345) and (146) cycles of $T^6$ respectively. Then there are $Q1 Q2 Q3$ distinct spacetime points where three branes intersect. We attach 3-branes wrapped around the (256) cycle to these points. By T-duality from the 4-4-4-0 configuration we expect that these (256) branes come in multiply charged varieties. We suppose as before that there are $k$ bosonic and $k$ fermionic “orientations” of the fourth species of 3-brane that arise as collective excitations of the strings running between the branes. Then the combinatorics of the previous section goes through unchanged giving an entropy of

$$S = 2\pi \sqrt{\left(\frac{k}{4}\right) Q1 Q2 Q3 Q4}.$$ 

The existence of the 3-3-3-3 configuration makes it manifest that the entropy must be symmetric in all charges. It is disappointing that it appears very difficult to extract any additional insight from this remarkable configuration.

4.2 2-2-2-6

The 2-2-2-6 black hole in Type IIA is particularly suggestive. In this case the Q4 6-branes are completely wrapped around $T^6$ so it is not possible to separate them as we did in the 3-3-3-3 and 4-4-4-0 cases. Now the $\int C^{(3)} \wedge F \wedge F$ coupling between the RR 3-form and the world-brane gauge field strength implies that each 2-brane separately looks like a an instanton of the SU(Q4) 6-brane gauge theory. Therefore the intersecting 2-branes on a 6-brane should correspond to a classical solution of 6 dimensional Euclidean gauge theory with the property that projection onto the (1234), (3456) or (1256) cycles yields a standard four-dimensional instanton. It would very interesting to count the degeneracy of this intersecting brane system by counting the orientation degrees of freedom of these six-dimensional “instantons”.

5 Conclusion

We have presented four dimensional extremal black hole configurations with finite area that are composed entirely out of D-branes. We argued that the entropy of these configurations naturally scaled with the charges to match the Bekenstein-Hawking entropy formula. We gave heuristic arguments that fixed the overall coefficient in the entropy, but much more needs to be done to establish the rules for this counting. The Type IIB configuration of four intersecting 3-branes is particularly interesting because the complete symmetry between the various branes explicitly reflects the symmetry of the entropy formula. A possible advantage of the black hole configurations presented
here is that they do not contain any solitonic objects. Consequently it should be easier to conduct scattering experiments from them than from the previously constructed black hole configurations.

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