Microscopic Entropy of $N=2$ Extremal Black Holes

David M. Kaplan$^\dagger$, David A. Lowe$^\sharp$, Juan M. Maldacena$^\natural$ and Andrew Strominger$^\ddagger$

$^\dagger$Department of Physics
University of California
Santa Barbara, CA 93106-9530, USA

$^\natural$Department of Physics and Astronomy
Rutgers University
Piscataway, NJ 08855, USA

$^\sharp$California Institute of Technology
Pasadena, CA 91125

Abstract
String theory is used to compute the microscopic entropy for several examples of black holes in compactifications with $N = 2$ supersymmetry. Agreement with the Bekenstein-Hawking entropy and the moduli-independent $N = 2$ area formula is found in all cases.

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1 dmk@cosmic1.physics.ucsb.edu
2 lowe@theory.caltech.edu
3 malda@physics.rutgers.edu
4 andy@denali.physics.ucsb.edu
1. Introduction

More than two decades after its discovery, our understanding of string theory has finally developed to the point where it can be used to provide, in special cases, a precise statistical derivation of the thermodynamic Bekenstein-Hawking area law for the entropy. While some definite relations between the laws of black hole thermodynamics and the statistical mechanics of stringy microstates have been clearly established, much more remains to be understood. For example a universal derivation of the area law for all types of black holes remains elusive. For these reasons it is important to understand as many cases as possible.

In [1] the entropy was microscopically computed for the simplest case of a five-dimensional extremal black hole. Supersymmetry makes it possible to count the microstates at weak coupling and then extrapolate into the strongly-coupled black hole region. This result was extended to include rotation in [2]. The four dimensional case with $N = 8$ supersymmetry was microscopically computed in [3] and [4]. The result agreed with the formulae for the area derived in [5,6]. Given the fact that the area is independent of moduli [7,5], these formulae are fixed up to a few constants by symmetries of $N = 8$ and $N = 4$ supergravity. In this paper we analyze several examples with $N = 2$ supersymmetry, which has not been previously considered. The $N = 2$ area formula was derived for the pure electric case in [7], for the general case with electric and magnetic charges in [8] and elegantly related to central charge minima in [9]. No symmetries are available here, and the formulae have a rather different character involving rational fixed points in the special geometry moduli space. Agreement is again found in the examples considered herein.

2. Type IIA Orientifold Example

First we consider the Type IIA theory on $K3 \times S^1 \times \hat{S}^1$. This theory has $N = 4$ supersymmetry in $d = 4$. In order to construct an $N = 2$ theory, we orientifold this model by a geometric $\mathbb{Z}_2$ symmetry combined with reversal of worldsheet orientation. Black hole solutions of the $N = 4$ theory invariant under this action will also be solutions of the orientifold model. In this manner we construct black hole solutions of an $N = 2$ Type IIA orientifold model and compute their macroscopic and microscopic entropy.

We consider the $\mathbb{Z}_2$ orientifold which combines the Enriques involution on $K3$, reflection on $S^1$, and translation by $\pi$ on $\hat{S}^1$, together with left-right exchange on the worldsheet. In M-theory language, this is a purely geometric orbifold of $K3 \times S^1 \times \hat{S}^1 \times S_{11}$, which acts as Enriques, combined with reflection on $S^1 \times S_{11}$ and translation by $\pi$ on $S^1$. This model was discussed in a different construction in [10] (with the translation on $S_{11}$ rather than $\hat{S}^1$) and with this construction in [11]. It has $N=2$ supersymmetry in $d=4$ with 11 vector multiplets and 12 hypermultiplets.

Now we wish to construct a four-dimensional black hole solution in this $N = 2$ theory as a collection of intersecting branes [12]. Such a solution is obtained from a set of intersecting branes in the original $N = 4$ theory as follows. First, $Q_5$ symmetric 5-branes wrap $K3 \times \hat{S}^1$. We will consider 5-branes not centered at the fixed points of the reflection on $S^1$. Since we wish to construct a configuration invariant under the $\mathbb{Z}_2$ symmetry each
5-brane is accompanied by its $\mathbb{Z}_2$ image, so that $Q_5$ is even. In addition, a 4-brane wraps the product of a holomorphic 2-cycle, $\Sigma$, of $K3$ with $S^1 \times \hat{S}^1$. It is possible to show that for any $\Sigma \in H_4(K3)$ with even self-intersection number $m = 2Q^2_4$, there is a choice of complex structure of $K3$ such that the cycle may be realized as a holomorphic curve of genus $Q_4^2 + 1$. Since $S^1 \times \hat{S}^1$ is odd under the $\mathbb{Z}_2$, we must impose the additional restriction that the 2-cycle of $K3$ be odd under the Enriques involution. In order that the 4-cycle be supersymmetric in the orientifold theory, the 2-cycle must be holomorphic with respect to the odd self-dual two-form of $K3$. Note also that the symmetric 5-branes cut each of the 4-cycles into $Q_5$ pieces along the $S^1$ direction. Finally, we include $n$ quanta of momentum along the $\hat{S}^1$ direction.

In terms of the original $N = 4$ theory, we are considering a set of $Q_5$ symmetric 5-branes, a 4-brane with self-intersection number $2Q^2_4$ and total momentum $n$. This configuration is related by duality to a configuration of one 6-brane and $Q^2_4 + 1$ 2-branes, $Q_5$ 5-branes and momentum $n$. The statistical entropy of such a configuration was calculated in [3] and found to be

$$S = 2\pi \sqrt{Q^2_4 Q_5 n},$$

in the limit of large charges. This agrees with the Bekenstein-Hawking entropy of the corresponding black hole solution. The entropy is duality invariant, so (2.1) will also hold in the case at hand.  

Now let us consider the Bekenstein-Hawking entropy in the orientifold theory. This may be computed by dividing the horizon area $A_{10}$ of the ten-dimensional solution by the ten-dimensional Newton’s constant $G_{10}$. Newton’s constant is unaffected by the orientifolding but the $\mathbb{Z}_2$ acts freely on the horizon and hence divides the area in half. The entropy $S'$ in the orientifold theory is then

$$S' = \frac{A'_{10}}{4G'_{10}} = \frac{A_{10}}{8G_{10}} = \frac{S}{2},$$

where the prime denotes quantities in the orientifold theory.

It is also easily seen that orientifolding reduces the microscopic entropy by half. The microscopic entropy is carried by $Q^2_4 Q_5$ massless supermultiplets that live on a string wrapping $\hat{S}^1$. The orientifold reduces the length of this string by half. Since the entropy is an extensive quantity it is also reduced by half. The orientifold also introduces twisted spatial boundary conditions for the massless supermultiplets. However this affects mainly the zero mode structure and not the asymptotic form of the entropy for large $n$.

3. Type I Example

In this section we consider black holes in Type I theory on $K3$. These theories have $N = 2$ supersymmetry in four dimensions when we further compactify two more dimensions on a torus. First we describe the classical black hole solutions, then quantize the charges in four and five dimensions and compute the Bekenstein-Hawking entropy. This is found to agree with the number of microscopic configurations obtained using D-brane techniques.
3.1. Classical Solutions

For the five(four)-dimensional black holes we consider Type I on $K3 \times S^1 \times \hat{S}^1$. The four-dimensional classical solution was found in [3] for the case of toroidal compactification. The solution in the $N = 2$ case is the same as the $N = 4$ case as the relevant terms in the low energy supergravity lagrangians involved are identical. In four dimensions the black hole we consider carries charges corresponding to NS solitonic 5-branes wrapping around $K3 \times S^1$, Kaluza-Klein monopoles on $\hat{S}^1$, and fundamental string winding and momentum along $S^1$. The five-dimensional configuration is the same, except for the absence of the Kaluza-Klein monopole. The five dimensional solutions are treated in [13].

It is useful to rewrite the entropy formulas in terms of integer quantized charges. The fundamental strings are winding along $S^1$ so the charge quantization condition will be the same as in $N = 4$. The quantization for momentum will be the same. The NS 5-brane is the Dirac dual to the string winding along $\hat{S}^1$, so it also has the same quantum of charge as in $N = 4$ case, and the Kaluza-Klein monopole is the Dirac dual to momentum along $\hat{S}^1$ so again it is the same as in the $N = 4$ case. Therefore the formula for the entropy is

$$S = 2\pi \sqrt{Q_5 Q_{KK} Q_1 n},$$

where $Q_5, Q_{KK}, Q_1, n$ are the number of NS 5-branes, Kaluza-Klein monopoles, winding strings, and momentum, respectively. The entropy of the five-dimensional black hole is as in (3.1) with the factor $Q_{KK}$ set to one.

3.2. D-Brane Counting

Consider first the five-dimensional case. The Type I configuration on $K3 \times S^1$ consists of D 5-branes wrapping $K3 \times S^1$, D-strings wrapping $S^1$ and momentum flowing along $S^1$. Note that these are D-branes of Type I theory so that the D 5-brane has an $SU(2) = Sp(1)$ gauge field living on it. When $Q_5$ D-5-branes coincide we have an $Sp(Q_5)$ gauge theory on the brane. The D-1-brane charge is carried by instantons of this gauge theory which are self dual gauge connections on the $K3$. For large $Q_1$ and $Q_5$ the number of bosonic degrees of freedom of the instanton moduli space is $4Q_1 Q_5$ and they come mainly from the different ways of orienting the instantons inside the gauge group. We could also, as in [14], count the moduli by considering open strings going between the D 1-branes and the D 5-branes. These open strings are unoriented, so there are 2 bosonic ground states for each string, plus 2 possible Chan Paton factors for each 5-brane. This leads to a total of $4Q_1 Q_5$ bosonic states. (In the corresponding Type II counting there is a factor of 2 arising from the 2 possible orientations of the open string). There are also $4Q_1 Q_5$ fermionic degrees of freedom. The momentum $n$ along the $S_1$ direction will be carried by oscillations in the instanton moduli space or, equivalently, by the $(1,5)$ strings. In either case the counting is that of a 1+1-dimensional gas with $4Q_1 Q_5$ bosonic and fermionic flavors.

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5 The actual state will also have some (1,1) and (5,5) strings excited in such a way that D-terms vanish [14,15].

6 It should be kept in mind that when $n$ is not much bigger than $Q_1 Q_5$ then the effects of multiple windings become important and both the number of flavors and the effective length of the circle increase, giving the same result for the entropy [17].
therefore use the standard asymptotic formula for the entropy which yields

\[ S = 2\pi \sqrt{Q_1 Q_5 n} \, . \]

This agrees with the classical result.

Now let us consider the four-dimensional case. We have the same configuration of branes as in five dimensions plus a Kaluza-Klein monopole. It was shown in [1] for \( N = 4 \) (but the arguments also apply to \( N = 2 \)) for the case of one monopole that (3.2) is unchanged, in agreement with (3.1) for \( Q_{KK} = 1 \). The general formula then follows from duality. It is nevertheless instructive, as well as useful for following sections, to see how this works in detail. We first perform a T-duality along \( \hat{S}^1 \) to a Type I' theory. The D 5-branes become D 6-branes and the D-strings become D 2-branes wrapped around \( S^1 \times \hat{S}^1 \). The momentum remains momentum and the Kaluza-Klein monopole becomes a NS 5-brane wrapping K3×\( S^1 \). This theory also has two orientifold planes perpendicular to \( \hat{S}^1 \). The total number of Type II 5-branes is \( 2Q_{KK} \): a 5-brane and an image 5-brane is the minimum that we can have, therefore it is the “unit” of solitonic 5-brane charge in this theory. However only \( Q_{KK} \) of them are in between a pair of orientifold hyperplanes so that the effective number of 5-branes in the locally Type II theory between two orientifold hyperplanes is \( Q_{II}^5 = Q_{KK} \). The number of D-2-branes in the locally Type II theory is \( Q_{II}^2 = Q_1 \) (the same as the number of original one branes), while the number of locally Type II D-6-branes is \( Q_{II}^6 = 2Q_5 \). The extra factor of two comes from the extra \( Sp(1) \) Chan Paton index carried by Type I 5-branes. We can now use the Type II result [3] to count the number of configurations between two orientifold planes. The counting in [3] relied on the fact that the solitonic 5-branes slice the 2-branes in \( Q_{II}^5 \) pieces. Now we must in addition consider the D-8-branes, which also intersect the 2-branes, further dividing up their worldvolumes. However for large charges they will have a subleading effect since the number of 8-branes is much smaller than \( Q_{II}^5 \) in that limit (\( 16 \ll Q_{II}^5 \)). Similarly we will not worry about possible slicing of the 6-branes by the 8-branes\(^7\). The fact that the 2-branes can end on 5-branes [18,19] implies that different slices in between different solitonic 5-branes can move independently. The momentum is carried mainly by (2,6) strings living between particular 5-branes. Note however that the orientifold projection will correlate what happens on one side of the orientifold hyperplane with what happens on the other side. In particular when we put a unit of momentum between two 5-branes we also have to put a unit of momentum on the two image 5-branes on other side of the orientifold hyperplane. Therefore we have only half of the total momentum available for distributing freely in the locally Type II theory, \( n_{II} = n/2 \). Once we have identified the correct number of degrees of freedom in the locally Type II theory we count as in [3] and obtain

\[ S = 2\pi \sqrt{Q_{II}^5 Q_{II}^6 Q_{II}^{11} n_{II}} = 2\pi \sqrt{Q_1 Q_5 Q_{KK} n} \, , \]

\( \text{(3.3)} \)

\( ^7\) We also do not worry about (2,8) or (6,8) strings because, for large charges, there are not as many flavors of these as there are for (2,6) strings.
which is the same formula as in the $N = 4, 8$ cases \[3\]. Again, this formula agrees with the classical result of equation \(3.1\).

4. More General $N=4$ Examples

In the previous two $N = 2$ examples the counting is similar to the maximally supersymmetric $N = 8$ cases. This is partly because the black holes we chose to analyze were present in $N = 8$ as well as $N = 2$ theories. In this section we analyze some new features that appear only when there is less than the maximal supersymmetry. In the $N = 8$ case all gauge charges are part of the supergravity multiplet, while in the $N = 4$ case we can have extra gauge multiplets. In this latter, more general case, the entropy formula in four dimensions can be written in terms of an $O(6, 22)$ vector of magnetic charge $P$ and an $O(6, 22)$ vector of electric charges $Q$. We consider Type I/heterotic theory on $T^6$, in which case the electric charges are carried by the Type I D-1-brane or heterotic fundamental string. Define $Q = (\frac{1}{2}p_R, \frac{1}{2}p_L, \frac{1}{\sqrt{2}}q)$ where $p_{R,L}$ are the right and left-moving momenta of a heterotic string on $T^6$ and $q$ are the 16 $U(1)$ charges of a generic compactification. In terms of D-branes these charges are carried by (1,9) strings. The black holes we considered in the previous section had $p_{R,L}^5 = (\frac{2}{n} \pm Q_1 R)$ (with other components of $p_{R,L}$ set to zero) and $q = 0$; now let us consider $q$ different from zero. The magnetic charges are still carried by the D-5-brane and the Kaluza-Klein monopole. As in the preceding section, we go to the Type $I'$ theory with $Q_1$ D-2-branes, $Q_5$ D-6-branes and $Q_{KK}$ solitonic 5-branes.\[8\] Now the open strings that carry momentum $n$ will also have to carry some charge. The charge is carried by (2,8) strings, the D-8-branes appeared when we did the T-duality transformation to the Type $I'$ theory. These (2,8) strings are left-moving fermions on the intersection onebrane. The (6,8) strings can also carry some charge but they are massive when the (2,6) strings are excited.\[10\] As in the case of rotating black holes\[2\] we conclude that the effective momentum that is left to distribute in (2,6) strings, after we have put enough (2,8) strings to account for the charge, is $n_{\text{eff}} = n - q^2/2Q_1$, where the factor of $Q_1$ arises as in \[2\] from the different flavors of (2,8) strings among which the charge is distributed.

The entropy formula becomes

\[
S = 2\pi \sqrt{Q_1 n_{\text{eff}} Q_{KK} Q_5} = 2\pi \sqrt{(Q_1 n - \frac{1}{2}q^2) Q_{KK} Q_5} = 2\pi \sqrt{Q^2 P^2 - (Q \cdot P)^2}, \quad (4.1)
\]

since $Q^2 = Q_1 n - \frac{1}{2}q^2$ and in this case $Q \cdot P = 0$. This is the classical formula \[3\]. Here $q^2$ is an even integer, each left-moving (2,8) fermion carries one unit of charge and the total current carrying fermion number is restricted to even values\[20\].

It is also of interest to consider a black hole that is extremal (in the sense that the mass is such that the solution is on the threshold of developing a naked singularity) but

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8 The subindex indicates what the object was in the original Type I theory, hopefully this will not cause confusion.

9 Similar observations have been made by C. Vafa (private communication).
not BPS, as for example a black hole with $Q^2 < 0$. It can be seen from the general black hole solutions in [5] that the classical entropy formula is just

$$S_{\text{ext}} = 2\pi \sqrt{|Q^2| P^2}.$$ (4.2)

To be definite consider a black hole with zero momentum $n = 0$, but with some gauge charge $q$, so that $Q = (\frac{1}{2} Q_1 R, -\frac{1}{2} Q_1 R, \sqrt{2} q), \; Q^2 = -\frac{1}{2} q^2$. In this case, we have to have enough $(2,8)$ strings to carry the charge, but since the $(2,8)$ strings can move only in one direction [21] they will carry some net momentum along the direction of the original string. This implies that there must be an equal amount of open strings moving in the opposite direction. These will be $(2,6)$ strings since they carry the most entropy. The black hole, and the D-brane system, are not BPS but they are extreme in the sense that they carry the minimum amount of mass consistent with the given charges. Note that there is no BPS bound for the charges $q$ under which this black hole is charged. In fact there can be light particles charged under this gauge group near an enhanced symmetry point. A real-world electrically charged extremal black hole will be of this type, in the sense that the electron is nearly massless compared to the string scale. Such black holes are not stable and will decay quickly by emitting charged particles.

5. Entropy for more general $N = 2$ cases

Now we construct a black hole with charges that can exist only in the $N = 2$ case and not in the $N = 4$ case. In order to do this we use the Gimon-Polchinski model [21] which is connected to the Type I on K3 considered above. This model contains 9-branes and 5-branes. The 5-branes are oriented along the directions $(012345)$. We compactify the directions $(45)$ and T-dualize along the direction 4. Now we have 2 orientifold hyperplanes; the 9-branes (5-branes) transform into 16 8-branes (4-branes) between the orientifold hyperplanes. The 4-branes are oriented along $(01235)$ and one can choose a point in moduli space where there are not any coinciding branes. There are also orbifold fixed points on the internal torus $(6789)$. These branes are “background” branes in the sense that they are completely extended along the macroscopic four dimensional space and are part of the vacuum state.

Now we include the same configuration that we had before: $Q_5$ solitonic 5-branes along $(056789), \; Q_6$ 6-branes along $(0456789), \; Q_2$ 2-branes along $(045)$ and momentum $n$ along the direction 5. If these are all the charges we have, the state counting for this case is the same as in the previous case, since all “background” branes give contributions that are subleading in the limit of large charges.

The new feature, relative to the $N=4$ case, is that we can have extra charges associated to 4-branes. These charges will be carried by $(4,6)$ strings which are left-moving fermions, they are related by T-duality to the $(2,8)$ or $(1,9)$ strings of the previous section. The $(2,4)$ strings also carry charge, but become massive when a condensate of $(2,6)$ strings form. The $(4,8)$ strings are a purely subleading contribution since they involve only the background branes, and, in any case, we can sit at a point in moduli space where they are massive.

If the black hole also carries charge $p$ under the 4-brane $U(1)s$ and charge $q$ under the 8-branes $U(1)s$, then we are forced to have some left-moving $(2,8)$ and $(4,6)$ strings thus
reducing the available momentum that we can distribute among the highly entropic (2,6) modes by $n_{\text{eff}} = n - \frac{q^2}{2Q_2} - \frac{p^2}{2Q_6}$. The formula for the entropy then becomes

$$S = 2\pi \sqrt{\left(Q_2Q_6n - \frac{1}{2}Q_2p^2 - \frac{1}{2}Q_6q^2\right)Q_5}. \tag{5.1}$$

5.1. Classical Solution

The Gimon-Polchinski model is U-dual to a heterotic theory on $K^3$ with a instanton numbers (12,12) embedded in the two $E_8$ factors $[22,23]$. The six-dimensional low-energy lagrangian for this heterotic theory has been considered in $[24]$. This is equivalent to the Type I action. The relevant terms in this action, in heterotic variables, are:

$$S = \frac{(2\pi)^3}{\alpha'^2} \int d^6x \sqrt{-g} \left( e^{-\phi} \left[ R + (\partial \phi)^2 - \frac{1}{12} H^2 - \frac{\alpha'}{8} F^2 \right] - \frac{\alpha'}{8} \tilde{F}^2 \right) + \frac{(2\pi)^3}{\alpha'^2} \int_{M_6} -\frac{\alpha'}{4} B \wedge \tilde{F} \wedge \tilde{F} - \frac{\alpha'^2}{8} \omega_3 \wedge \tilde{\omega}_3, \tag{5.2}$$

where $F$ ($\tilde{F}$) denotes the field strength of the gauge fields arising from the 9-branes (5-branes), $\phi$ is the six-dimensional dilaton, and $\omega_3$ is the Chern-Simons form defined by $d\omega_3 = -F \wedge F/2$, and likewise for $\tilde{\omega}_3$. The field strength for the antisymmetric tensor field is defined in the usual way as $H = dB + \frac{\alpha'}{2} \omega_3$. Note we have dropped the higher derivative terms that appear in the action of $[24]$ which will only be relevant for large curvatures.

The equations of motion that follow from this action are invariant under a $\mathbb{Z}_2$ duality transformation which acts as

$$\begin{align*}
\phi &\rightarrow -\phi \\
g_{mn} &\rightarrow e^{-\phi} g_{mn} \\
H &\rightarrow e^{-\phi} * H \\
A &\rightarrow \tilde{A} \\
\tilde{A} &\rightarrow A.
\end{align*} \tag{5.3}$$

This symmetry is actually just a T-duality symmetry on the Type I side which inverts the size of the $K^3$.

Now let us compactify on a torus down to four dimensions and consider the classical black hole solution which carries the charges mentioned above. It follows from $[6,8]$ that there exists a solution with constant scalar fields, provided the asymptotic values of these scalars are adjusted to special values. Because the entropy does not depend on the asymptotic values of the moduli $[6,8]$ there is no loss of generality in restricting our considerations to this case. For this solution, it may be seen from the equations of motion that $H = *H$, where the Hodge dual $*$ is defined with respect to string metric. Taking into account the fact that $F \wedge F$ and $\tilde{F} \wedge \tilde{F}$ vanish for the solution at hand, the equations of motion take the same form as the usual $N = 4$ Type I (or heterotic) equations $[25]$. The extremal BPS
black hole solutions of the $N = 4$ equations have been classified in \cite{5}. Using these results it may be shown that the field $\phi$ satisfies

\[ e^\phi = \frac{Q_6}{Q_2}, \quad (5.4) \]

and the entropy is

\[ S = 2\pi \sqrt{\left( Q_2 Q_6 n - \frac{1}{2} Q_6 (q^2 + s^2) \right) Q_5}, \quad (5.5) \]

where all charges are as defined above, and $s = \gamma p$ with $\gamma$ a constant to be determined.

The difference between the $F$ and $\tilde{F}$ fields arises when one considers the relationship between the integer-valued quantized charges and the physical charges $Q^i$ (defined by $F^i = Q^i/r^2$, with $F^i$ the four-dimensional field strengths, as defined in \cite{25}). Since the gauge kinetic terms for the $\tilde{F}$ fields do not have the usual $e^{-\phi}$ factor in front, we find the relation $s = e^{-\phi/2} p$. Substituting this into (5.4) and (5.5) we find the Bekenstein-Hawking entropy of the black hole agrees with the microscopic counting (5.1).

6. Conclusions

We have found agreement between the macroscopic Bekenstein-Hawking entropies for BPS black holes in $N = 2$ supergravities and the microscopic entropy in string theory. In the first two examples the counting is very similar to the counting for the $N = 4,8$ cases. The only real difference is that the various branes are on a less supersymmetric background. The physical mechanism that gives rise to the large degeneracy is basically the same as in the more symmetric cases. We explored more intrinsically $N = 4,2$ cases by considering black holes which carry gauge charges that exist in these less supersymmetric theories.

It would be interesting to present a general argument testing the full $N = 2$ spectrum of charged black holes. In particular, D-brane counting for Type II string theories compactified on generic Calabi-Yau 3-folds that are not orbifolds of more symmetric cases is an unexplored problem. The results of \cite{8,9} describe a universal geometric formula for the entropy which must be somehow reproduced by D-branes. In particular the simple relationship \cite{9} of the entropy formula to the minima of the central charge should have a microscopic explanation.

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