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D3-branes dynamics and black holes

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Abstract

Using the D3-brane as the fundamental tool, we adress two aspects of D-branes physics. The first regards the interaction between two electromagnetic dual D-branes in 10 dimensions. In particular, we give a meaning to *both* even and odd spin structure contributions, the latter being non vanishing for non zero relative velocity v (and encoding the Lorentz-like contribution). The second aspect regards the D-brane/black holes correspondence. We show how the 4 dimensional configuration corresponding to a *single* D3-brane wrapped on the orbifold T^6/Z_3 represents a regular Reissner- Nordström solution of $d = 4$ $N = 2$ supergravity

Talk presented by Matteo Bertolini

1 Introduction and Summary

In the last few years the study of D-branes [1] and their dynamics has revealed to be, under many respects, one of the most promising aspects of string theory to be investigated. In this contribution I will adress two different aspects regarding D-brane physics.

Using the D3-brane as the fundamental tool I will first of all consider the interaction of two moving D3-branes in 10 dimensions. The fact that the D3-brane is a dyon (namely charged both electrically and magnetically with respect to the R-R 4-form) implies that the gauge interaction of two of these objects has both a Coulomb-like *and* a Lorentz-like contribution. From string theory point of view the two are encoded respectively in the even and the odd Ramond-Ramond spin structures emerging after GSO projection on the relevant cylinder amplitude. The treatment of the odd spin structure is delicate: a naive computation would give back always a vanishing result because of the presence of at least two fermionic zero-modes coming from transverse directions. This is not peculiar of the D3-brane of course, but it is a general problem affecting the interaction of any couple of electromagnetic dual p and $(6-p)$ D-branes. The first result, therefore, will be to illustrate a way to treat the non-trivial odd spin structure contribution in order to get the correct phase-shift we would expect from a field theory point of view.

In the second part of my contribution I will consider the 4 dimensional configuration corresponding to a D3-brane wrapped on the orbifold T^6/Z_3. Once integrated over the compact coordinates, this configuration corresponds to an exact black hole solution of the effective $N = 2$

supergravity theory in 4 dimensions. More precisely, and this is the most interesting fact, it represents an extremal dyonic Reissner-Nordström (R-N) black hole with non vanishing entropy. This a particular example of what happens in generic Calabi-Yau (CY) compactifications (T^6/Z_3 is an orbifold limit of a CY): as opposite to compactifications on tori, in CY compactifications one can get regular solutions, in 4 dimensions, even with single charged objects. The non-trivial topological structure of CY supersymmetric cycles on which the D-branes are wrapped can "regularize" the solution, as the intersection of different D-branes does on tori. The D-brane/black hole correspondence has been investigated both from a string and supergravity point of view and results are shown to agree. Moreover, it can be shown how the actual values of the 4 dimensional electric and magnetic charges depend explicitly on branes' orientation in the compact space.

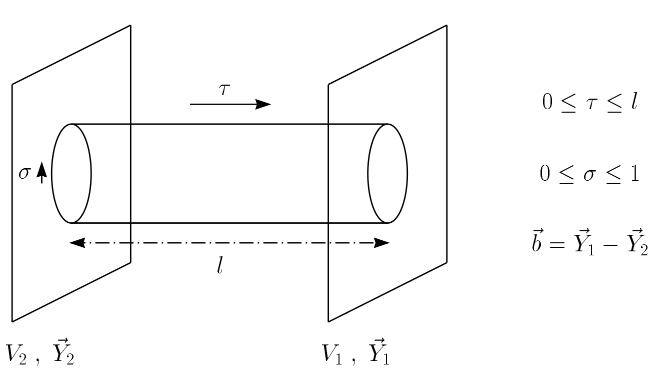
2 The interaction of two D3-branes in 10 dimensions

As well known, D-branes can be defined as hypersurfaces on which open strings can end and therefore the potential between two interacting D-branes is expressed by vacuum fluctuations of open strings streched between them. However, thanks to the conformal invariance of the string world sheet, one can study D-branes interactions also from closed string point of view. In this latter case the D-brane can be thought as a source of closed strings and the interaction is mediated by exchange of closed string states. A very appropriate tool in order to describe D-branes and their dynamics from a closed strings point of view is the *boundary state formalism* [2]. A boundary state $|B\rangle$ describing a D-brane can be defined as a coherent state written in terms of closed string oscillators which implement the boundary conditions (N or D) of strings which the brane can emit. The typical structure of a boundary state describing a D-brane is of the following form:

$$|B\rangle \;=\; |B\rangle_{bos}|B\rangle_{ghost}|B\rangle_{fer}|B\rangle_{superghost} \;\simeq\; |B\rangle_0 \;\exp\left[\sum_{n\geq 1}\frac{1}{n}\left(\alpha^{\mu}_{-n}S_{\mu\nu}\tilde{\alpha}^{\nu}_{-n}+...\right)\right]|0\rangle$$

where there are contributions both from physical and ghost oscillators and where $|B\rangle_0$ encodes the zero modes part of the boundary state. The matrix $S_{\mu\nu}$ is a diagonal matrix with $+1$ or -1 entries according to Dirichelet or Neumann nature of the corresponding directions, respectively.

Let us now consider the 10 dimensional interaction of two D3-branes moving with velocities $V_1 = tanhv_1$ and $V_2 = tanhv_2$ along X^1 direction, tilted by angles θ_1^a , θ_2^a on the 3 planes X^a, X^{a+1} ($a = 4, 6, 8$), that will eventually become compact, and with transverse positions $\vec{Y}_1$, $\vec{Y}_2$:



The parameter l is the world sheet distance while $\vec{b}$ is the physical impact parameter.

In the boundary state formalism the interaction is given by the correlation between two boundary states and in our case can then be written in the following way:

$$\mathcal{A} = \frac{\mu^2}{16} \sum_s < v_1, \vec{Y}_1, \theta^1_a | \int_0^\infty dl e^{-lH} | v_2, \vec{Y}_2, \theta^2_a >_s \tag{1}$$

$\mu = \sqrt{2\pi}$ is the electric (and magnetic) density charge of the D3-brane while the sum over s is made on the 4 spin structures emerging after GSO projection on the cylinder. The boundary state $|B, v\rangle$ describing a moving brane is obtained from the static one through a Lorentz transformation with velocity v: $|B, v\rangle = exp(-ivJ^{01})|B\rangle$ [3].

The essential kinematic is encoded in the bosonic zero-modes contribution which can be written as a product of delta functions enforcing the boundary conditions for the center of mass position operator X^μ_o, that is a Fourier superposition of momentum states. Hence, in momentum space, the boundary states for the two branes are:

$$|B, v_1, \theta^1_a, \vec{Y}_1\rangle = \int \frac{d^6\vec{k}}{(2\pi)^6} e^{i\vec{k}_B\cdot\vec{Y}_1} |v_1, \theta^1_a\rangle \otimes |k_B\rangle \ , \ |B, v_2, \theta^2_a, \vec{Y}_2\rangle = \int \frac{d^6\vec{q}}{(2\pi)^6} e^{i\vec{q}_B\cdot\vec{Y}_2} |v_2, \theta^2_a\rangle \otimes |q_B\rangle$$

with $\vec{k}_B$ and $\vec{q}_B$ being the boosted and rotated momenta:

$$\begin{aligned} k^\mu_B &= (\sinh v_1 k^1, \cosh v_1 k^1, k^2, k^3, \cos\theta^1_a k^a, \sin\theta^1_a k^a) \\ q^\mu_B &= (\sinh v_2 q^1, \cosh v_2 q^1, q^2, q^3, \cos\theta^2_a q^a, \sin\theta^2_a q^a) \end{aligned}$$

and k^2, k^3, q^2 and q^3 the momenta along the only two completely transverse directions.

Integrating over momenta and taking into account momentum conservation which for non-vanishing $v \equiv v^{(1)} - v^{(2)}$ and $\theta_a \equiv \theta^{(1)}_a - \theta^{(2)}_a$ forces all the Dirichlet momenta but k^2, k^3 to be zero, the amplitude (1) reads:

$$\mathcal{A} = \frac{\mu^2}{2\sinh|v| \prod_a 2\sin|\theta_a|} \int_0^\infty \frac{dl}{4\pi l} e^{-\frac{b^2}{4l}} \sum_\alpha Z_B Z^\alpha_F \tag{2}$$

where $Z^s_{B,F}$ are the bosonic and fermionic partition functions:

$$Z^s_{B,F} = < v^{(1)}, \theta^{(1)}_a | e^{-lH} | v^{(2)}, \theta^{(2)}_a >^s_{B,F}$$

In the above expression, only the oscillator modes of the string coordinates X^μ appear, since we have already integrated over the center of mass coordinate.

In order to compare string and supergravity results, we take in (2) the limit of large impact parameter ($b \to \infty$) . In this limit only world sheets with $l \to \infty$ contribute and the behaviour of partition functions simplifies a lot:

$$\begin{aligned} &Z_B \to 1 \ , \\ &Z^{even}_F \to 2\cosh v \prod_a 2\cos\theta_a - 2(2\cosh 2v + \sum_a 2\cos 2\theta_a) \\ &Z^{odd}_F \to 2\sinh v \prod_a 2\sin\theta_a \cdot 0 \end{aligned} \tag{3}$$

If the two branes are parallel and at rest, i.e. if $\theta_a = 0$ and $v = 0$, the configuration is BPS and one gets $Z^{even}_F = 0$ that expresses the well known no-force condition between two BPS states. If the two branes are parallel but not at rest, i.e. if $\theta_a = 0$ but $v \neq 0$, one gets $Z^{even}_F \sim v^4$ that is infact the phase shift leading order contribution for moving D-branes. Finally, in the more general case, the configuration is not BPS and one has $Z_F^{even} \neq 0$. However, for suitable angles ($\sum_a \theta_a = 2\pi n$) the configuration can still be BPS, although preserving less supersymmetry. This can be seen looking to the low velocity dependence of the partition functions; indeed in this case $Z_F^{even} \sim v^2$.

It seems to be no contribution from Z_F^{odd} because of the 0 always present in its expression. That 0 comes from the fermionic zero-modes of the only two completely transverse directions left, X^2 and X^3, and has to be soaked up if one wants to give a meaning to the odd spin structure. In the Ramond sector, odd and even spin structures are responsible for the gauge interaction between the two branes (the NS sector encodes the gravitational contribution). D3-branes are dyons and therefore, for relative velocity $v \neq 0$ there is both, in general, an electric (Coulomb-like) and a magnetic (Lorentz-like) interaction. The former is encoded in the even spin structure while the latter in the odd one. This means that one would expect a non vanishing contribuiton also from the odd spin structure partition function (and the same must hold for any couple of electromagnetic dual D-branes). One can naively understand this point looking to the tensorial structure of the two (even and odd) RR contributions. From (3) one sees that:

$$< v^1, \theta_1^a | e^{-lH} | v^2, \theta_2^a >_{RR+} \sim \cosh v \prod_a \cos \theta_a \tag{4}$$

while

$$< v^1, \theta_1^a | [...] e^{-lH} | v^2, \theta_2^a >_{RR-} \sim \sinh v \prod_a \sin \theta_a \tag{5}$$

The cosine is a signal of a radial (Coulomb) force while the sine of an orthogonal (Lorentz) one. The simbol [...] in the odd spin structure expression indicates the insertion of suitable supercurrents. In [4] it has been shown infact that with this insertion one can reproduce in string theory the expected field theory result. Namely, togheter with the primary necessity of soaking up fermionic zero-modes, one gets back, as a by-product, the right tensorial structure of the interaction with respect to the transverse non-compact directions, namely the exterior product structure that is charateristic of a Lorentz force. Notice that in order to get a non vanishing contribution from the odd spin structure one also needs the two 3-branes being non parallel (i.e. $\theta_a \neq 0$). This is again quite obvious: the condition for having non zero Lorentz interaction between *extended* objects is to have a complete non parallelism between the corresponding world-volume. This is a generalization of what happens in 4 dimensions to point-like objects where the Lorentz interaction is non vanishing for non zero relative velocity, this conditon being rephrased saying that there is a tilting between the worldlines of the two particles.

An essentially analogous result has been achieved in a T-dual situation, namely a D0-brane moving in the background of a D6-brane [6]. With the insertion of suitable regulator for matter and superghost zero modes it has been found a contribuiton from the odd spin structure as a scale-independent, velocity-dependent potential $V_{06}(r) = -\frac{v}{2r}$ whose *a posteriori* interpretation is of a Lorentz-like potential. In this way the direct result is not that of a phase-shift, however the final essential conclusion does not change, namely that the odd spin structure encodes the magnetic interaction contribution in a scattering amplitude of a pair of electromagnetic dual p and $(6-p)$ D-branes and is in general different from zero.

3 The black hole configuration in 4 dimensions

Let us now consider the four dimensional configuration corresponding to a D3-brane wrapped on a particular 6-dimensional compact space, the T^6/Z_3 orbifold. What has been acheived in the last few years about the D-brane/black hole correspondence is essentially that various D-branes configurations can in general give a microscopic description of both extremal and non-extremal black holes emerging in 4 dimensions by compactification of p-branes solutions of type II supergravity in 10 dimensions.

Starting from type II theory in 10 dimensons, for toroidal compactifications one obtains an effective $N = 8$ supergravity in 4 dimensions while for compactifications on more generic CY spaces one ends up with $N = 2$ theory. Configurations for which a microscopic description can be given correspond mainly to toroidal compactifications of "intersecting D-branes"[7]. In $N = 2$ compactifications, from microscopic point of view, much less has been said, the essential reason being that for curved D-branes Polchinski's prescription is more difficult to be implemented. The orbifold T^6/Z_3 is actually a limit of a CY space (with Hodge numbers $h_{(1,1)} = 9$ and $h_{(1,2)} = 0$) and so falls in this latter class. It's utility resides in the fact that it has a sufficently simple structure to be treated with usual boundary state techniques but gives, on the other hand, sufficently interesting results. Indeed, what I will show is that in our case, as in many other CY compactifications, one can obtain regular black hole solutions even with less than 4 charged objects, as opposite to the toroidal case. The intuitive reason for that is that the non-trivial topological structure of CY space implies that a single D-brane can "intersect" with itself on a given supersymmetric cycle, therefore mimicking the actual intersection of different D-branes needed on tori in order to get regular solutions.

Strictly speaking a R-N black hole is a non singular spherical solution of Maxwell-Einstein gravity that however can be consistently extended to be $N = 2$ supersymmetric. More in general, in matter coupled $N = 2$ supergravity, one can have generalized regular solutions which turn out to be R-N near the horizon, that is with an $AdS_2 \times S^2$ topology. In particular, one have the so called *double extreme* solution when the vector multiplets scalars are taken to be constant and equal to their fixed values they anyhow must get at the horizon [8]. The hypermultiplet scalars, on the other hand, couple minimally to the gauge fields and can therefore always taken to be neutral in a given solution.

When one compactifies type IIA or IIB supergravity on a CY manifold, the number of vector and hypermultiplets of the relevant $N = 2$ theory is dictated by the two relevant Hodge numbers ($h_{(1,1)}$,$h_{(1,2)}$) charaterizing the C-Y space. In our case, that is type IIB on T^6/Z_3, one has 0 vector multiplets and 10 hypermultiplets. Therefore there are not vector multiplet scalars and the solution is automatically double extreme. From a supergravity point of view the solution is therefore quite straightforward. Let us now analize the configuration both from a macroscopic and a microscopic point of view.

A type II p-brane usually couples to the metric, the dilaton and the corresponding (p+1) gauge potential. The peculiar property of the D3-brane is that it does not couple to the dilaton and therefore the equations of motions for the relevant field in the supergravity effective theory are simply:

$$\begin{aligned} R_{MN} &= T_{MN} \\ \nabla_M F_{(5)}^{MABCD} = 0 & \qquad \left(\leftarrow \; F^{(5)}_{G_1 \ldots G_5} = \frac{1}{5!} \, \epsilon_{G_1 \ldots G_5 H_1 \ldots H_5} \, F_{(5)}^{H_1 \ldots H_5} \right) \end{aligned}$$

One can make a block-diagonal spherically simmetric ansätz for the metric

$$ds^2 = g_{\mu\nu}(x)dx^\mu dx^\nu + g^{(6)}_{ab}(y)dy^a dy^b \quad \mu,\nu = 0,...,3 \quad a,b = 4,...9$$

choosing in particular the compact components g_{ab} not to depend on the non-compact coordinates x^μ. Notice that in general the compact components of the metric should depend on non-compact coordinates, becoming some of the scalar fields in 4 dimensions. In our case, however, we know that in the effective $d = 4$ $N = 2$ supergravity we don't have any vector multiplet scalar and the hypers have consistently been put to 0.

For the 5-form field strenght the ansätz is again quite simple because the only supersymmetric 3-cycles the 3-brane can be wrapped on are the complete holomorphic or anti-holomorphic ones because $h_{(2,1)} = 0$:

$$F_{(5)}(x,y) = F^0_{(2)}(x) \wedge \left(\Omega^{(3,0)} + \bar{\Omega}^{(0,3)}\right)$$

With these ansätze the equation of motions are satisfied in 10 dimensions and integrating on the 6 compact coordinates the lagrangian one gets in 4 dimensions is of a Maxweel-Einstein type and the solution turns out to be an extreme R-N solution. The electric and magnetic charges has the following values:

$$e = \frac{1}{2}\,\mu\sqrt{V^2_{D3}/V_{CY}}\cos\alpha \;,\; g = \frac{1}{2}\,\mu\sqrt{V^2_{D3}/V_{CY}}\sin\alpha$$

where V_{D3} and V_{CY} are the volumes of the 3-brane and of the full compact space respectively. The extremality conditions reads: $M^2 = \frac{1}{4}(e^2 + g^2)$, with M the mass of the solution. By now α is just an arbitrary parameter. What I will actually show in the following is how α depends explicitely on the D3-brane orientation in the compact space. In particular, the compactification procedure must preserve Dirac quantization conditions which, being satisfied in 10 dimensions with the minimal values of the charges (this result being valid for any for D-brane), must still be valid in 4 dimensions.

Let us now move to the microscopic side of the computation and see what happens to the previously computed scattering amplitude (2) once one compactifies. The first thing one has to do is to project the 10 dimensional boundary state $|B\rangle$ on its Z_3 invariant part in this way:

$$|B_{inv}\rangle = P|B,\theta^a_1\rangle = \frac{1}{3}(1 + g + g^2)|B,\theta^a_1\rangle$$

P is the projection operator and g is the generator of $(2\pi/3)$ rotations in each of the three planes X^a, X^{a+1} $(a = 4,6,8)$. The net effect is that the boundary state is similar to the previous one but with a sum on $\Delta\theta$ angles whose values depend on g (and hence on the particular structure of our manifold). In this way:

$$|B_{inv}\rangle = \frac{1}{3}\sum_{\Delta\theta^a}|B,\theta^a_1 + \Delta\theta^a\rangle \quad , \quad \Delta\theta^a = (\,0 \quad \tfrac{2\pi}{3} \quad \tfrac{4\pi}{3}\,)$$

The full computation gives back the following result for the partition functions, that is similar to (3) but with sums on angles $\Delta\theta$:

$$\begin{aligned} Z_F^{even} &\rightarrow 2\cosh v \sum_{\{\Delta\theta^a\}}\prod_a 2\cos\left(\theta_a + \Delta\theta_a\right) - 2(2\cosh 2v + \sum_{\{\Delta\theta^a\}}\sum_a 2\cos 2\left(\theta_a + \Delta\theta_a\right)) \\ Z_F^{odd} &\rightarrow 2i\sinh v \prod_a 2\sum_{\{\Delta\theta^a\}}\sum_a 2\sin\left(\theta_a + \Delta\theta_a\right) \end{aligned}$$

Using trivial trigonometrical properties of $\Delta\theta_a$ angles which are peculiar, however, just of the orbifold, one obtains a very simple result for Z_F^{even} and the 4 dimensional amplitude reads:

$$\mathcal{A} \sim (\frac{3}{4}\cosh v\ \cos(\sum_a \theta_a)\ -\ \cosh 2v)\ + 0 \tag{6}$$

In 4 dimensional field theory language the three terms represent respectively the exchange of vector fields, of graviton fields and of scalar fields (to which the wrapped D3-brane therefore does not couple): the configuration is then of a R-N type (no scalars excited) as we have already seen from supergravity side. It is easy to see that there is no contribution from the twisted sector because the wrapped D3-brane has mixed boundary conditions in at least one of the three T^2 componing T^6/Z_3 and this is not consistent with twisting.

Notice that in order to find which field the brane couples to in 4 dimensions it would have not been necessary to study any scattering amplitude but rather one point functions of supergravity fields on the disk representing the D-brane. This is precisely what has been done in [5] and the results are the same as here.

As anticipated, it is interesting to see how the D3-brane orientation in the compact space affects the values of the electric and magnetic 4 dimensional charges and how Dirac quantization condition changes from 10 to 4 dimensions. In dimensions $d = 2(q+1)$ Dirac quantization condition (Dqc) is: $eg' + (-1)^q e'g = 2\pi n$ hence it has a $+$ sign in 10d and a $-$ sign in 4d. As it is weel known, D-branes realize Dqc with the minimal amount of charge and for any couple of p and $(6-p)$ branes it is always true that $\mu_p\mu_{6-p} + \mu_{6-p}\mu_p = 4\pi$. Once one compactifies the numerical values of the charges is affected by the compact part of the bosonic zero-mode part of the boundary state. Indeed the net effect of the compactification is that compact space zero-modes aquire a discrete structure. Before doing any orbifold projection their whole contribution to the scattering amplitude turns out to be the following:

$$< \theta_1^a, \vec{Y}_1 | e^{-lH} | \theta_2^a, \vec{Y}_2 >_B = \frac{V(B_1)V(B_2)}{\mathrm{Vol}(T_6)} = \prod_a \frac{|\bar{n}_a^{(1)}\bar{n}_{a+1}^{(2)} - \bar{n}_{a+1}^{(1)}\bar{n}_a^{(2)}|}{\sin|\theta_a|} \tag{7}$$

where $V(B_1)$ and $V(B_2)$ are the volumes of the two 3-branes and where various $\bar{n}_a^i$ are integers depending explicetly to the geometric configuration of the branes in the compact space. The 4 dimensional amplitude therefore becomes the following one:

$$\mathcal{A} = \frac{\hat{\mu}^2}{\sinh|v|}\int_0^\infty \frac{dl}{4\pi l} e^{-\frac{b^2}{4l}} \frac{1}{16}\sum_s Z_B Z_F^s \qquad \text{with} \qquad \hat{\mu}^2 \equiv \mu^2 \frac{V(B_1)V(B_2)}{\mathrm{Vol}(T_6)}$$

One can now see that Dqc is still valid in 4 dimensions with an integer that depends on branes' orientation in the compact space. Moreover, using again trivial trigonometrical properties, one can rewrite it as a combinations of 8 different charges (4 electric and 4 magnetic) whose exact values depend on suitable angles combinations. Using (5) and (7):

$$\hat{\mu}^2 \prod_a \sin\theta_a \ = \sum_{i=1}^4 \left(e_i^{(1)} g_i^{(2)} - g_i^{(1)} e_i^{(2)}\right) \ = \ 2\pi \prod_a |\bar{n}_a^{(1)}\bar{n}_{a+1}^{(2)} - \bar{n}_{a+1}^{(1)}\bar{n}_a^{(2)}|$$

$$e_i^{(1,2)} = \frac{\hat{\mu}}{2}\cos\Phi_i^{(1,2)}\ ,\ g_i^{(1,2)} = \frac{\hat{\mu}}{2}\sin\Phi_i^{(1,2)}\ \ ;\ \Phi_i^{(1,2)} = \theta^4_{(1,2)} \pm \theta^6_{(1,2)} \pm \theta^8_{(1,2)} \tag{8}$$

One can see that there are always at least 4 charges with non vanishing values, independently from different values one can choose for the various angles.

By the orbifold projection only one pair of electro-magnetic charges survives, because only one angle combination ($\Phi_1 = \theta^4_{(1,2)} + \theta^6_{(1,2)} + \theta^8_{(1,2)}$) is Z_3 invariant and the final result is the following:

$$e^{(1,2)} = \frac{\hat{\mu}}{2} \cos \sum_a \theta^a_{1,2} \ \ , \ \ g^{(1,2)} = \frac{\hat{\mu}}{2} \sin \sum_a \theta^a_{1,2}$$

still satisfying Dqc. As expected, each 3-brane is a point-like 4 dimensional dyon and couples to the unique gauge field present, the graviphoton. Moreover, the previously introduced parameter α has now a precise value, being the sum of the 3 angles each brane is tilted with respect to the compact space referring directions. For parallel branes one sees that $e^{(1)} = e^{(2)}$, $g^{(1)} = g^{(2)}$ and the Lorentz contribution cancels (see (8)) as expected for identical dyons in 4 dimensions.

It would be interesting to carry on the orbifold compactification on the type IIA side where the number of vector multiplets would be 9 rather than 0 and so there would be non-trivial vector multiplet scalars. A very interesting fact is that in this case the $N = 2$ supergravity effective theory turns out to be a consistent truncation of $N = 8$ supergravity and in particular the solution one obtains is the most general 1/8 supersymmetry preserving $N = 8$ black hole modulo U-duality transformations. The construction of this solution (metric, scalar and vector fields) is the main result of a forthcoming paper I'm doing in collaboration with M. Trigiante and P. Frè.

Of course, it would also be interesting to consider compactifications on more complicated Calabi-Yau spaces on which however a much more involved technique is needed in order to deal with D-branes and boundary states in an efficient way.

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