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Matching the BPS spectra of heterotic–type I–type I′ strings [1]

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Abstract

We give a detailed discussion of the matching of the BPS states of heterotic, type I and type I′ theories in $d = 9$ for general backgrounds. This allows us to explicitly identify these (composite) brane states in the type I′ theory that lead to gauge symmetry enhancement at critical points in moduli space. An example is the enhancement of $SO(16) \times SO(16)$ to $E_8 \times E_8$. © 1998 Elsevier Science B.V.

1. Introduction

In nine dimensions the heterotic theories and the orientifold projections of the type II theories are believed to be related by a chain of dualities. Already in ten dimensions it has been conjectured in [1] and further substantiated in [2] that there is a strong-weak coupling duality between the heterotic $Spin(32)/\mathbb{Z}_2$ and the type I theory [6]. The latter, when viewed as an orientifold projection of the IIB

theory, lives on an orientifold nine-plane, and consistency requires the presence of 16 pairs of D9 branes which give rise to an SO(32) gauge group [6]. Upon compactification on a circle this becomes a nine-dimensional duality. The type I theory on S^1 is [7] T-dual to the orientifold projection of the IIA theory known as I′ theory. The compact dimension of the I′ theory is topologically a segment and its two endpoints define two orientifold eight-planes. Consistency requires the presence of 16 mirror pairs of D8 branes parallel to the orientifold planes.

In a background where the unbroken gauge group is SO(16) × SO(16) one gets a closed chain of dualities. For this background the two heterotic theories are related by $R \leftrightarrow 1/R$ duality [8]. Furthermore, one can view both, the nine-dimensional $E_8 \times E_8$ heterotic theory and the type I′ theory as compactifications of M theory on a cylinder $S^1 \times (S^2/\mathbb{Z}_2)$, where in the former case the dilaton is related to the length of the cylinder, whereas in the latter case it is related to its circumference. It has been conjectured

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[6] Various aspects of heterotic-Type I duality have been verified in Refs. [3–5].

in [9] that the two theories are connected by a duality transformation which exchanges $S^1/\mathbb{Z}_2$ with S^1. Attempts to describe the heterotic string as a matrix quantum mechanics of type I′ D particles [10] are based on this conjectured chain of dualities.

In this letter we present a detailed mapping of the type I′ D0 brane states to type I/heterotic BPS states for generic backgrounds. The mapping is based on relating the BPS mass formulas of the corresponding states over the whole moduli space. In this way we can explicitly identify the objects in type I′ theory that are expected to give rise to gauge symmetry enhancement at special points in the moduli space. The generic picture that we find is that the heterotic/type I BPS states at $n = 1$ winding map to bound states of a single D0 brane sitting at one of the orientifold planes with wrapped open and closed type I′ strings. In the $SO(32)$ and other backgrounds where one observes a $U(1) \to SU(2)$ gauge symmetry enhancement at the self-dual heterotic/type I radius, one can identify the $W^{\pm}$ gauge bosons with a D0 and anti-D0 brane sitting at the orientifold plane which has no D8 branes. A D0 brane sitting at the opposite orientifold plane is the image of a spinor weight of the $Spin(32)/\mathbb{Z}_2$ lattice.

2. Review of heterotic–type I–type I′ duality

Before matching the BPS states of the type I′ and the type I theories, or, via duality, of the type I′ and the heterotic theory, we have to review some of the discussion of [2].

The low energy effective action of the type I theory is

$$S_{\mathrm{I}} = \int d^{10}x\sqrt{-g}\ \frac{1}{2\kappa_0^2}\ e^{-2\phi}\big(\mathscr{R} + 4\,\partial_M\phi\,\partial^M\phi\big)$$

$$- T_9 \int d^{10}x\sqrt{-g}\ e^{-\phi}$$

$$\times \left[16 + \frac{(\pi\alpha')^2}{2}\ \mathrm{tr}_v F_{MN} F^{MN} \right]$$

$$- 16\,\mu_9 \int A_{10}\,. \tag{1}$$

The entire space is an orientifold plane which is charged under the ten-form field A_{10}. In addition

there are also 16 D9 branes. F is the $SO(32)$ field strength, the trace being in the vector representation and $\mu_p^2 = 2\kappa_0^2 T_p^2 = 2\pi(4\pi^2\alpha')^{3-p}$, $\kappa_0 = 8\pi^{7/2}\alpha'^2$ [6].

The low energy effective action of the type I′ theory, on the other hand, is

$$S_{\mathrm{I'}} = \int d^{10}x\sqrt{-g}\ \frac{1}{2\kappa_0^2}\ e^{-2\phi}\big(\mathscr{R} + 4\,\partial_M\phi\,\partial^M\phi\big)$$

$$- \frac{1}{2}\int G_{10}\ ^*G_{10}$$

$$- \sum_i \left\{ T_8 \int_{x^9 = x_i^9} d^9x\sqrt{-g}\ e^{-\phi} \right.$$

$$\times \left[n_i + \frac{(\pi\alpha')^2}{2}\ \mathrm{tr}_v F_{\mu\nu} F^{\mu\nu} \right]$$

$$\left. + n_i\,\mu_8 \int_{x^9 = x_i^9} A_9 \right\}. \tag{2}$$

The sum is over the D8 branes, with n_i positioned at x_i^9; $\sum n_i = 16$. G_{10} is the field strength of the nine-form potential to which the D8 branes couple; the gauge group depends on the positions of the D8 branes [6].

Polchinski and Witten have found a non-trivial background which solves the type I′ equations of motion:

$$\kappa = \kappa_0 e^{\phi'(y)} = z(y)^{-5/6}, \quad \Omega(y) = Cz(y)^{-1/6},$$

$$z(y) = \frac{3C}{\sqrt{2}}\big[B(y)\mu_8 - \nu(y)y\big], \tag{3}$$

with $g_{MN}^{\mathrm{I'}} = \Omega^2(y)\eta_{MN}$ and $y \in [0,2\pi]$ parameterizing the segment $S^1/\mathbb{Z}_2$. The equation of motion for the 10-form field strength $G_{10}(y) = \nu(y)dx^0 \wedge \ldots \wedge dx^9$ requires $\nu(y)$ to be piecewise constant with discontinuity $\Delta\nu = n_i\mu_8$ at a crossing of a group of n_i D8 branes at $y = y_i$. In addition the type I′ projection implies the boundary conditions $\nu(2\pi) = -\nu(0) = 8\mu_8$. To get continuous metric and dilaton backgrounds, as required by their equations of motion, $B(y)$ is also piecewise constant with discontinuity $\Delta B = n_i y_i$. $B(y)$ is fixed by $B \equiv B(0)$. The type I′ space of classical vacua is thus parameterized by the constants B and C together with the positions y_i of the D8 branes. To have a meaningful back-

ground also requires $z(y) \geq 0$ everywhere. This implies (up to a physically irrelevant choice for the sign of C) that there is a minimum value B_m for B which depends on the configuration of D8 branes.

The type I′ theory is, by definition, the T-dual of the type I theory. In a non-constant background, the relation between the 9-dimensional type I′ and type I metrics is not known a priori. Following [2] we set $g_{\mu\nu}^{I'} = \Omega^2(y)\gamma_{\mu\nu} = \Omega^2(y)Q^2 g_{\mu\nu}^{I}$, with Q a constant factor to be determined. Comparing the 9-dimensional gravitational actions gives

$$2\pi R e^{-2\phi} = 2Q^7 \int_0^{2\pi} dy \, \Omega^8(y) e^{-2\phi'}$$

$$= 2\kappa_0^2 \, Q^7 C^{25/3} \int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{1/3}, \quad (4)$$

while comparison of the gauge actions results in

$$R e^{-\phi} = \kappa_0 \sqrt{\alpha'} \, C^5 Q^5. \quad (5)$$

Note that since the integration on the right-hand side of (4) is restricted to $S^1/\mathbb{Z}_2$, there is an additional factor of 2 [11].

To establish the relation between R, ϕ and B, C we need one additional equation, as we also have the unknown parameter Q. Still following Ref. [2], we will get this equation by comparing masses of BPS states in the two theories.

Matching the masses of a type I Kaluza-Klein state of mass $1/R$ and a type I′ winding state of mass $m_{I'}$, taking into account the relation between the respective nine-dimensional metrics, i.e. $1/R = Qm_{I'}$, one finds [2]

$$\frac{1}{R} = \frac{2Q}{2\pi\alpha'} \int_0^{2\pi} \Omega^2(y) \, dy$$

$$= \frac{QC^{5/3}}{\pi\alpha'} \int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{-1/3}. \quad (6)$$

Solving (4), (5) and (6) we find

$$Q = C^{-5/6} \frac{\left[\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{1/3} \right]^{1/4}}{\left[\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{-1/3} \right]^{1/4}}, \quad (7)$$

$$R = \pi\alpha' C^{-5/6} \left[\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{-1/3} \right]^{-3/4}$$

$$\times \left[\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{1/3} \right]^{-1/4}, \quad (8)$$

$$e^{\phi} = \frac{1}{8} \pi^{-5/2} \alpha'^{-3/2} C^{-5/3} \left[\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{-1/3} \right]^{1/2}$$

$$\times \left[\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{1/3} \right]^{-3/2}. \quad (9)$$

To complete the map between the vacua of type I′ and type I we need to establish the relation between a Wilson-line in type I and the positions of the D8 branes in type I′. In the absence of a Wilson-line all D8 branes lie at $y = 0$. We then introduce a Wilson line in the type I theory with a single non-vanishing entry A_i; all type I gauge bosons whose roots have a non-vanishing i-th component will acquire a mass $\frac{A_i}{R}$ (we take here for simplicity $A_i > 0$). This picture corresponds in the type I′ description to one in which all branes lie at $y = 0$ except for one which lies at $y = y_i$; the mass of the gauge bosons with a non-vanishing entry in the i-th component of their root vector comes, in this picture, from the stretching of open strings between the branes at 0 and the one at y_i. Comparing the masses one finds

$$A_i = \frac{RQC^{5/3}}{2\pi\alpha'} \int_0^{y_i} dy \left(\frac{z(y)}{C} \right)^{-1/3}$$

$$= \frac{1}{2} \frac{\int_0^{y_i} dy \left(\frac{z(y)}{C} \right)^{-1/3}}{\int_0^{2\pi} dy \left(\frac{z(y)}{C} \right)^{-1/3}}. \quad (10)$$

For $y_i = 0, 2\pi$ we have $A_i = 0, 1/2$, as expected.

3. Matching of the spectra

We will now show how heterotic winding modes, or, via heterotic-type I duality, the D-string winding modes of type I theory, map onto D0 brane states of the type I′ theory.

The mass of a single D0 brane measured in type I units is

$$M_{D0}(y) = Q\Omega(y)\frac{T_0}{2}e^{-\phi'(y)}. \qquad (11)$$

Q and $\Omega(y)$ appear as a result of the non-trivial relation of the nine-dimensional metrics and the non-trivial world-line measure respectively. The factor $1/2$ reflects the fact that a single D0 brane has half the charge and therefore via the BPS condition half the tension of a dynamical object which consists of a brane and its mirror image. Note that the mass of the D0 brane depends on its position in the compact dimension, which for a single D0 brane has to be one of the two fixed points $y = 0$ or $y = 2\pi$.

A state in the $SO(32)$ heterotic string spectrum compactified on a circle S^1 of radius R in the presence of a Wilson line A is characterized by its winding number n, its momentum quantum number m, a weight vector p of the $\text{Spin}(32)/\mathbb{Z}_2$ lattice and oscillator excitations $N_{L,R}$. Its mass is ($a_R = 0(\frac{1}{2})$) for R (NS) sector) (see e.g. [8])

$$\tfrac{1}{2}\alpha'_h M^2 = \left(N_L - 1 + \tfrac{1}{2}p_L^2\right) + \left(N_R - a_R + \tfrac{1}{2}p_R^2\right), \qquad (12)$$

with

$$p_L = \left(p + An, \sqrt{\frac{\alpha'_h}{2}}\left(\frac{m - A\cdot p - \tfrac{1}{2}A^2 n}{R} + \frac{nR}{\alpha'_h}\right)\right),$$

$$p_R = \sqrt{\frac{\alpha'_h}{2}}\left(\frac{m - A\cdot p - \tfrac{1}{2}A^2 n}{R} - \frac{nR}{\alpha'_h}\right). \qquad (13)$$

Taking into account level matching, i.e. $N_L - 1 + \tfrac{1}{2}p_L^2 = N_R - a_R + \tfrac{1}{2}p_R^2$ and restricting to BPS states, i.e. $N_R = 0(\frac{1}{2})$ for R (NS) sectors, one obtains

$$M^2 = \left(\frac{m - A\cdot p - \tfrac{1}{2}A^2 n}{R} - \frac{nR}{\alpha'_h}\right)^2, \qquad (14)$$

$$N_L = 1 - nm - \tfrac{1}{2}p^2. \qquad (15)$$

In these expressions α'_h is related to the heterotic string tension via $\alpha'_h = \frac{1}{2\pi T_h}$. It has been shown in [2] that a single type I closed D string has the worldsheet structure of the $SO(32)$ heterotic string in the fermionic formulation. The tension of this string is $T_D = \tfrac{1}{2}T_1 e^{-\phi} = \frac{1}{4\pi\alpha'}e^{-\phi}$. Here α' is the type I scale and a factor of $1/2$ appears again because the

object we discuss is half a dynamical object. We can then read the spectrum of a single type I D string by substituting in the previous formulas T_D for T_h or $2\alpha'e^{\phi}$ for α'_h.

To be specific, we will consider a Wilson line that breaks the gauge symmetry to $SO(16)\times U(8) \subset SO(32)$ so that at a generic value of R the symmetry is $SO(16)\times U(8)\times U(1)\times U(1)$. The choice of the Wilson line which accomplishes this is

$$A = \left(\frac{\epsilon}{2},\ldots,\frac{\epsilon}{2},0,\ldots,0\right), \qquad (16)$$

with eight entries of each kind and $\epsilon \in [0,1]$. For $\epsilon = 0$ we recover the $SO(32)$ string and for $\epsilon = 1$ the gauge group is $SO(16)\times SO(16)$. In the type I' the 16 physical D8 branes are split into two groups of eight, one positioned at $y = 0$, the other at y_0 which follows from (10) with $A_i = \frac{\epsilon}{2}$.

D string winding modes are expected to map via T-duality to D0 brane states in type I', so in a first step we are trying to determine those states of the D string which are mapped to a D0 brane sitting at $y = 0$ and $y = 2\pi$ respectively. It is straightforward to show that for the choice of quantum numbers $N_L = 0, n = m = 1, p^2 = 0$ one has

$$M^I = \frac{R}{2\alpha'}e^{-\phi} - \frac{1-\epsilon^2}{R} = M_{D0}^{I'}(2\pi). \qquad (17)$$

This in fact is the lightest type I state at $n = 1$. Note also that at the self-dual radius $R_*^2 = 2\alpha'(1-\epsilon^2)e^{\phi}$ this state together with the $n = m = -1$ one (which maps to an anti-D0 brane sitting at $y = 2\pi$) become massless and provide the extra gauge bosons for the enhancement $U(1) \to SU(2)$. The $R = R_*$ locus in the type I/heterotic moduli space corresponds to the $B = B_m$ locus on the type I' side. The type I state that matches $M_{D0}^{I'}(0)$ belongs to the spinor conjugacy class of $\text{Spin}(32)/\mathbb{Z}_2$. Choosing the weight $p = (-\frac{1}{2},\ldots,-\frac{1}{2})$ and quantum numbers $N_L = 0, n = 1, m = -1$ one indeed verifies

$$M^I = \frac{(1-\epsilon)^2}{R} + \frac{R}{2\alpha'}e^{-\phi} = M_{D0}^{I'}(0). \qquad (18)$$

This is again the lightest type I state at $n = 1$ which belongs to the spinor conjugacy class of $\text{Spin}(32)/\mathbb{Z}_2$, although for the specific Wilson line that we chose it is not unique. For the constant

SO(16) × SO(16) background, which corresponds to $\epsilon = 1$, we have $M_{D0}^{I'}(2\pi) = M_{D0}^{I'}(0)$ as expected.

The fact that the two conjugacy classes of Spin(32)/$\mathbb{Z}_2$ map to D0 branes sitting at different ends of the type I′ compact dimension follows from T-duality; cf. also [12]. The type I Wilson-line A determines in type I′ the positions of the D8 branes via the dualization of 9-9 to 8-8 open strings (9-9 momenta in the compact direction become 8-8 windings). The position of a D0 brane in type I′ is determined relative to the position of the D8 branes via the dualization of 1-9 to 0-8 open strings. To be concrete let us consider the case with unbroken SO(32), when all the type I′ D8 branes sit at $y = 0$. In type I the 1-9 strings can feel the presence of a gauge holonomy around a closed D string. The possible gauge holonomies belong to the $\mathbb{Z}_2$ group $\{1, -1\}$ [2]. The massless excitations of the 1-9 strings live in the world-volume of the D string and are fermionic degrees of freedom with respect to the SO(1,1) Lorentz group of this world-volume theory [2]. The transportation of a massless 1-9 state around the closed D string detects the presence or absence of a non-trivial holonomy and induces a -1 or $+1$ phase respectively. This means that in the presence of a non-trivial holonomy the massless 1-9 excitations give rise to anti-periodic fermions on the D string world-volume, while the trivial holonomy results in periodic ones. When the D string is wrapped once around the compact dimension, and in the absence of a gauge holonomy, the momentum mode of the 1-9 strings along this direction is integer moded (in quanta of $1/R$), while in the presence of the non-trivial -1 holonomy it is half-integer moded. This maps via T-duality to integer and half-integer windings of the 0-8 strings in the type I′ theory respectively, which in turn implies that the wrapped D string with a trivial holonomy maps to a D particle at $y = 0$, that is on top of the D8 branes, while the non-trivial holonomy puts the D particle at $y = 2\pi$ making the 0-8 windings half-integral. Given then the fact that the trivial holonomy corresponds to the periodic sector of the D string current algebra, which provides the spinor conjugacy class of Spin(32)/$\mathbb{Z}_2$, one expects that this conjugacy class be represented in type I′ theory by a D0 brane sitting at $y = 0$. Accordingly the adjoint conjugacy class which arises in the anti-periodic sector of the D string current

algebra should be represented by a D0 brane sitting at $y = 2\pi$. The same conclusions about the position of a single D0 brane and its expected relation to weights in the spinor or adjoint conjugacy class of Spin(32)/$\mathbb{Z}_2$ arise in the more complicated case when a type I Wilson line is present moving some or all of the D8 branes away from $y = 0$.

Beyond the two heterotic states that match the mass of a D0 brane placed at $y = 2\pi$ and at $y = 0$ respectively, the heterotic string has a whole host of additional BPS states at $n = 1$. Let us first however recollect what happens at $n = 0$. For $n = 0$ the level-matching condition (15) implies that the BPS spectrum is exhausted by the two possibilities $p^2 = 0$, $N_L = 1$ and $p^2 = 2$, $N_L = 0$ with arbitrary values of the Kaluza-Klein number m. This is nothing but the tower of states resulting from the Kaluza-Klein reduction of the ten-dimensional supergravity and super-Yang-Mills multiplets to nine dimensions. The mass of these states is always proportional to $\frac{1}{R}$. At the lowest level ($m = 0$) it is zero for the gravity and the unbroken gauge multiplets, and $\frac{1}{R}A \cdot p$ for the SO(32) gauge bosons with $A \cdot p \neq 0$. All of these states are mapped in the type I′ theory to winding modes of closed and open fundamental strings. The heterotic Kaluza-Klein momentum maps to type I′ winding number. The class of heterotic states with $p^2 = 0$, $N_L = 1$ map on the I′ side to the lowest level of stretched closed strings if they belong to the supergravity multiplet, or to the lowest level of stretched open strings with endpoints on the same D8 brane if they are in the Cartan subalgebra. The second class of heterotic BPS states ($p^2 = 2$, $N_L = 0$) correspond to non-vanishing roots of the SO(32) lattice and are mapped in type I′ to the lowest level of wound open strings with ends on different D8 branes.

At $n = 1$ the picture is more complicated. Level-matching now permits arbitrarily large values of both N_L and p^2, since their contribution can be canceled by negative values of m. Starting with $p^2 = 0$, $N_L = 0$ we get $m = 1$ and this is the largest possible value of m at $n = 1$, corresponding to the lightest state at this winding, as can be seen from (14). We identified this state with a D0 brane sitting at $y = 2\pi$ in the type I′ theory. Further inspection of equations (14) and (15) shows that increasing N_L by one unit has to be compensated for level-matching by a one-unit

decrease in m. This results in an increase [7] in the mass of the state by $\frac{1}{R}$. Similarly increasing p^2 by two has to be compensated by decreasing m by one and the corresponding increase in the mass of the state is now equal to $\frac{1 + A \cdot \Delta p}{R}$, where Δp is the difference between the new and the original weight. It is then evident that all mass increases above the lightest state are proportional to $\frac{1}{R}$, a scale which maps on the type I′ side to the scale of fundamental string windings. These observations suggest the following type I′ picture: Starting with the lightest object at $n = 1$, which is the D0 brane at $y = 2\pi$, one could envisage constructing the states which match the heterotic BPS states with $N_L > 0$, as bound states at threshold of the given D0 brane with closed and open (ends on the same D8 brane) strings at positive winding [8]. These stretched strings have in their ground state a mass which equals their winding number times $\frac{1}{R}$ in heterotic units, as suggested by equation (6), and therefore reproduce exactly the increase in the heterotic BPS mass brought about by non-vanishing values of N_L. The heterotic states with $p^2 > 0$ could then be matched by bound states at threshold of the D0 brane with wound open strings whose ends are on different D8 branes. Here and below we assume the existence of these threshold bound states. For example the heterotic BPS states with $p^2 = 2$, $N_L = 0$ have $m = 1$ and exceed the mass of the D0 brane at $y = 2\pi$ by an amount by $\frac{1 + A \cdot p}{R}$ (in heterotic units). This excess of mass is exactly the stretching energy of an open string which winds once around the compact dimension (i.e. from 0 to 2π and back), and has its ends fixed on the pair of D8 branes which correspond to the considered root $p = \pm(e_i \pm e_j), i, j = 1, \dots, 16$. There is however a subtlety here if we want to adopt the interpretation that the non-trivial gauge quantum numbers of these type I′ states are due to the Chan-Paton indices of the open string part in the bound state: one can never represent a spinor weight of $\mathrm{Spin}(32)/\mathbb{Z}_2$ with open strings stretched between D8 branes. This is the T-dual of the observation that one does not get the

spinor of $\mathrm{Spin}(32)/\mathbb{Z}_2$ in the open string sector of the type I theory. What this suggests is that it is not possible to start with the D0 brane at $y = 2\pi$ and construct the whole set of heterotic BPS states at $n = 1$ by binding this D0 brane with stretched open and closed strings. This however is just as well since we saw that a D0 brane placed at $y = 0$ has exactly the right mass to represent the lightest heterotic states in the spinor of $\mathrm{Spin}(32)/\mathbb{Z}_2$. Given then the fact that the difference of two spinorial weights is always a weight in the 0 conjugacy class, one expects that bound states of the D0 brane at $y = 0$ with stretched open strings will represent the whole set of heterotic $n = 1$ states in the spinor of $\mathrm{Spin}(32)/\mathbb{Z}_2$. The picture just described reproduces in the I′ theory the heterotic $n = 1$ BPS mass formula. At the same time it accounts for the gauge quantum numbers of these states in I′.

The $SU(2)$ symmetry enhancement which occurs in the heterotic theory at the self-dual radius for generic backgrounds comes about by means of the two states $n = m = 1$ and $n = m = -1$ at $p = N_L = 0$. We have argued that these states map to a D0 and an anti-D0 brane respectively placed at $y = 2\pi$. It is noteworthy that when attempting to understand the $SU(2)$ enhancement in the type I′ both a D0 and an anti-D0 brane seem to be essential ingredients, whereas in the proposed matrix description of the strong coupling limit of the IIA theory in the infinite momentum frame only D0 branes are required [13]. For the special choice of background for which the unbroken gauge group is $SO(16) \times SO(16)$ the self-dual radius is zero and the symmetry enhancement which occurs there is more interesting: a large number of states become massless at every value of n, thus enhancing the unbroken gauge group to $E_8 \times E_8$ and at the same time unfolding a new large dimension [9]. At $n = 1$ the states that become massless are in the $(128,1) + (1,128)$ of $SO(16) \times SO(16)$ [14], and the state with $n = m = 1$ at $p = N_L = 0$ is one of them. It is interesting that for this choice of background we can consistently consider only positive values of n (this implies that we restrict ourselves to

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[7] Note that we work with radii larger than the self-dual radius for the type I′ background to be meaningful, i.e. to have $z(y) \geq 0$.

[8] Negative winding strings would then bind with the anti-D0 brane to provide the corresponding heterotic states at $n = -1$.

[9] An alternative approach to recovering the $E_8 \times E_8$ symmetry in the presence of a $SO(16) \times SO(16)$ background can be found in Ref. [5].

positive momenta along the new dimension which emerges). For generic backgrounds on the other hand, one needs states at $n = \pm 1$ to get the $W^{\pm}$ bosons for the gauge symmetry enhancement $U(1) \rightarrow SU(2)$.

Heterotic BPS states at windings $n > 1$ map to type I′ bound states of n D0 branes at a fixed point with wrapped closed and open fundamental strings. An interesting observation is that the spectrum of the heterotic theory does not contain BPS states at $n > 1$ whose mass can match bound states at threshold of n D0 branes only, i.e. without wrapped fundamental strings.

Acknowledgements

References

[1] E. Witten, String Theory Dynamics in Various Dimensions, Nucl. Phys. B 443 (1995) 85, hep-th/9503124

[2] J. Polchinski and E. Witten, Evidence for Heterotic-Type I String Duality'', Nucl. Phys. B 460 (1996) 525, hep-th/9710169

[3] C. Bachas and E. Kiritsis, F^4 Terms in N = 4 String Vacua, Nucl. Phys. Proc. Suppl. 55B (1997), hep-th/9611205

[4] C. Bachas, C. Fabre, E. Kiritsis, N.A. Obers and P. Vanhove, Heterotic/Type I Duality and D-brane Instantons, hep-th/9707126

[5] J. Blum and K. Dienes, From the Type I String to M Theory: a Continuous Connection, hep-th/9708016

[6] J. Polchinski, TASI Lectures on D-Branes, hep-th/9611050

[7] J. Dai, R.G. Leigh and J. Polchinski, New Connections Between String Theories, Mod. Phys. Lett. A, Vol. 4, 21 (1989) 2073

[8] P. Ginsparg, Comment on Toroidal Compactification of Heterotic Superstrings, Phys. Rev. D 35 (1986) 648

[9] P. Horava and E. Witten, Heterotic and Type I String Dynamics from Eleven Dimensions, Nucl. Phys. B 460 (1996) 506, hep-th/9510209

[10] U. Danielsson and G. Ferretti, The Heterotic Life of the D Particle, Int. J. Mod. Phys. A 12 (1997) 4581, hep-th/9610082; S. Kachru and E. Silverstein, On Gauge Bosons in the Matrix Model Approach to M Theory, Phys. Lett. 396B (1997) 70, hep-th/9612162; T. Banks, N. Seiberg and E. Silverstein, Zero and One-Dimensional Probes With N = 8 Supersymmetry, Phys. Lett. B 401 (1997) 30, hep-th/9703052; T. Banks and L. Motl, Heterotic Strings from Matrices, hep-th/9703218; D. Lowe, Bound States of Type I′ D Particles and Enhanced Gauge Symmetry, Nucl. Phys. B 501 (1997) 134, hep-th/9702006; Heterotic Matrix String Theory, Phys. Lett. B 403 (1997) 243, hep-th/9704041; S.J. Rey, Heterotic M(atrix) Strings and their Interactions, hep-th/9704158; P. Horava, Matrix Theory and Heterotic Strings on Tori, hep-th/9705055; D. Kabat and S.-J. Rey, Wilson Lines and T-Duality in Heterotic Matrix Theory, hep-th/9707099

[11] E. Witten, Strong Coupling Expansion of Calabi-Yau Compactifications, Nucl. Phys. B 471 (1996) 135, hep-th/9602070

[12] J. Polchinski, S. Chaudhuri and C. Johnson, Notes on D-Branes, hep-th/9602052

[13] T. Banks, W. Fischler, S.H. Shenker and L. Susskind, M Theory as a Matrix Model: a Conjecture, Phys. Rev. D 55 (1997) 5112, hep-th/9610043

[14] S. Kachru and E. Silverstein, as in [10]