D-branes on Asymmetric Orbifolds

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Abstract

We construct D-brane states on an asymmetric orbifold of type IIA on a four-torus, which is modded out by T-duality. We find explicit boundary states charged under the twisted sector gauge fields. Unlike other cases, the boundary states involve an explicit dependence on the twist fields. The D-brane spectrum is consistent with the model being equivalent to type IIA on a four-torus.

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1 Introduction

Asymmetric orbifolds are string vacua where the orbifold action acts differently on the left- and right-movers [1]. Since the left- and right-moving bosons on the worldsheet see different target spaces, these vacua usually do not have a simple geometrical interpretation. Such vacua are potentially interesting for phenomenology because they typically have very few moduli. In certain cases, all moduli except the dilaton are projected out [2].

Since these spaces have no geometrical interpretation, it is difficult to ascertain where they lie in the web of string dualities. In particular, it is unknown whether all such orbifolds are smoothly connected to large volume Calabi-Yau spaces. As a consequence, an M-theoretic interpretation of these vacua is missing.

Based on the experience with other string dualities, it is natural to guess that the clue to understanding the non-perturbative structure of these orbifolds is the understanding of their D-brane spectra. D-branes have been shown to probe sub-stringy structure, and in other cases (e.g. symmetric orbifolds), the metric on the D-brane worldvolume reproduces expected properties of the underlying geometry [3]. It is of great interest, therefore, to analyze the moduli space metric in cases where no underlying geometric structure is known, as in the case of asymmetric orbifolds.

Furthermore, D-branes control the strong coupling behaviour of a theory in many cases, so the knowledge of the D-brane spectra constrains possible dual models, and may also suggest guesses for the dual theory.

Asymmetric orbifolds have also arisen in a recent attempt to construct non-supersymmetric vacua with zero cosmological constant [4]. It was shown that in a particular model the cosmological constant vanished to two loops and it was conjectured to vanish to all orders in perturbation theory. Non-perturbative contributions then become important, and in [5] an example was analyzed where duality arguments suggested the existence of a non-perturbative contribution to the cosmological constant. It is of interest to see if direct computations using D-branes in this model support this result. So again we are led to the study of D-brane spectra.

From a technical point of view, constructing D-branes in asymmetric spaces is qualitatively different from the symmetric case. D-branes are defined as endpoints of strings, so at the boundary of the string worldsheet one has to provide a boundary condition relating the left moving fields to the right-moving fields. However, in asymmetric spaces, there is no obvious way
to write such a condition. Heuristically, this is why not all string theories have D-branes.

The question studied in this paper is that of the existence (and construction) of “twisted” D-branes, i.e. D-branes that couple to R-R fields coming from a twisted sector. An untwisted brane can be constructed by summing all images of a boundary state under the orbifold action [6]. However, these states are, at least in the example studied here, not the minimal D-brane states. For example they preserve only a quarter of the supersymmetries.

We shall construct D-branes in a special example, where we can avoid many of the more difficult problems. This is the case of type II string theory on an orbifold of the four-torus, where the orbifold action is that of an overall T-duality [7]. This turns out to have a second description in terms of type II string theory on $T^4$, where the D-brane spectrum is known. In this case, using the equivalent description, we are able to construct the D-branes explicitly.

We find that the D-branes indeed arise from twisted sectors of the asymmetric orbifold. We describe these states explicitly as boundary states [8, 9]. Using this formalism, we show that these D-branes behave exactly as expected of branes on a four-torus. In particular they are BPS states, preserving half of the supersymmetry. Furthermore, the D-branes from the untwisted sector are shown to be combinations of twisted D-branes. It is the twisted branes that play the role of the elementary objects in the equivalent torus description.

This result leaves important questions unanswered; it is not known how to construct the twisted D-branes on general asymmetric spaces. It, however, provides an existence proof of twisted D-branes, and emphasizes their importance. Hopefully, the example here may also hint at a more general construction. We hope to return to these questions in the future.

The paper is organized as follows. We review in section 2 the general formalism of boundary states as applied to toroidal backgrounds of type II string theory. In section 3 we introduce the T-duality orbifold, compute the closed string spectrum and demonstrate the spacetime symmetries of the models. These symmetries are used in section 4 to construct the boundary state for a general BPS saturated D-brane in this model.

While this work was in progress, we became aware of [6], where related issues are discussed.
2 Boundary States on a Torus

Let us briefly review the construction of boundary states \([8, 9]\) in toroidal backgrounds. We denote by \(x^6, \ldots, x^9\) four compactified directions, with radii \(R_6, \ldots, R_9\). To avoid introducing (super)ghosts, we will work in the light-cone gauge as in \([10]\). We consider Dirichlet boundary conditions localizing the D-branes at the origin of all non-compact directions.

The problem at hand is to construct the boundary states \(|B\rangle\) satisfying the boundary conditions

\[
(L_n - \bar{L}_{-n})|B\rangle = 0 \tag{1}
\]

This condition ensures that the resulting open string theory is conformally invariant. Further conditions (called Cardy’s conditions \([12]\)) are necessary to make the open string sector a sensible boundary CFT (e.g. degeneracies should be integral).

In practice, the above condition is hard to solve. Instead one imposes (in the free field case) the more restrictive condition \([9]\)

\[
(\alpha_i^n - \bar{\alpha}_{i-n})|B\rangle = 0
\]

\[
(\alpha^\mu_n + \bar{\alpha}^\mu_{-n})|B\rangle = 0, \tag{2}
\]

where \(i = p+1, \ldots, 9\) labels the Dirichlet directions, and \(\mu\) labels the Neumann directions. This manifestly implies (1), since

\[

L_n = \sum_m \left( \alpha_i^m \alpha_{i-n}^m + \alpha^\mu_m \alpha^\mu_{i-n} \right)
\]

\[
\bar{L}_{-n} = \sum_m \left( \bar{\alpha}_{i-m}^{i} \bar{\alpha}_{i-n}^{i} + \bar{\alpha}^\mu_{i-m} \bar{\alpha}^\mu_{i-n} \right) \tag{3}
\]

Note that the condition requires the \(U(1)^4\) invariance of the torus, or generally a chiral symmetry algebra larger than simply the Virasoro algebra \([9]\).

The condition (2) is solved, in the bosonic case, by the coherent state

\[
|Dp; k\rangle = \exp \left\{ \sum_{n>0} \frac{1}{n} (\alpha_{i-n}^i \bar{\alpha}_{i-n}^i - \alpha_{i-n}^\mu \bar{\alpha}_{i-n}^\mu) \right\} |k\rangle, \tag{4}
\]

where \(|k\rangle\) denotes the string ground state of momentum \(k^i\) in the \(i\) directions and winding \(m^\mu\) in the \(\mu\) directions. This state is called an Ishibashi state.
To maintain worldsheet supersymmetry, we need to impose conditions on the fermions as well. In the NS-NS sector, the conditions are

\[
(p^{i}_{r} - i\eta \tilde{p}_{i-r}) |Dp; \eta; k\rangle_{NSNS} = 0 \\
(p^{\mu}_{r} + i\eta \tilde{p}^{\mu}_{r}) |Dp; \eta; k\rangle_{NSNS} = 0
\]

where \(\eta\) is a sign needed for GSO projection.

The corresponding Ishibashi state is

\[
|Dp; \eta; k\rangle_{NSNS} = \exp\left\{ \sum_{n>0} \frac{1}{n} \left( \alpha^{i}_{-n} \tilde{\alpha}^{i}_{-n} - \alpha^{\mu}_{-n} \tilde{\alpha}^{\mu}_{-n} \right) \right\} |k\rangle \\
\otimes \exp\left\{ i\eta \sum_{r>0} (p^{i}_{-r} \tilde{p}^{i}_{-r} - p^{\mu}_{-r} \tilde{p}^{\mu}_{-r}) |F\rangle_{NSNS}\right\}
\]

where \(|F\rangle_{NSNS}\) is the NSNS fermionic Fock vacuum.

To GSO project, we note that \((-1)^{F}\) and \((-1)^{\tilde{F}}\) act as

\[
(-1)^{F} |Dp; \eta; k\rangle_{NSNS} = - |Dp; -\eta; k\rangle_{NSNS} \\
(-1)^{\tilde{F}} |Dp; \eta; k\rangle_{NSNS} = - |Dp; -\eta; k\rangle_{NSNS}
\]

and hence the GSO projected state is

\[
|Dp; k\rangle_{NSNS} = \frac{1}{\sqrt{2}} (|Dp; +; k\rangle_{NSNS} - |Dp; -; k\rangle_{NSNS})
\]

We now perform the same analysis in the RR sector. The boundary conditions are then

\[
(p^{i}_{r} - i\eta \tilde{p}^{i}_{r}) |Dp; \eta; k\rangle_{RR} = 0 \\
(p^{\mu}_{r} + i\eta \tilde{p}^{\mu}_{r}) |Dp; \eta; k\rangle_{RR} = 0
\]

with the solution

\[
|Dp; \eta; k\rangle_{RR} = \exp\left\{ \sum_{n>0} \frac{1}{n} \left( \alpha^{i}_{-n} \tilde{\alpha}^{i}_{-n} - \alpha^{\mu}_{-n} \tilde{\alpha}^{\mu}_{-n} \right) \right\} |k\rangle \\
\otimes \exp\left\{ i\eta \sum_{r>0} (p^{i}_{-r} \tilde{p}^{i}_{-r} - p^{\mu}_{-r} \tilde{p}^{\mu}_{-r}) |F_{\eta}\rangle_{RR}\right\}
\]

where \(|F_{\eta}\rangle_{RR}\) is an appropriately chosen R-R vacuum. The Fock vacuum of the R-R sector carries left- and right spinor indices, and as a result of the
condition (9) for the fermionic zero modes, is constrained to satisfy $[15, 18, 16]$: 

$$
(\psi_0^i - i\eta \tilde{\psi}_0^i)|F_\eta\rangle_{R-R} = 0 \\
(\psi_0^\mu + i\eta \tilde{\psi}_0^\mu)|F_\eta\rangle_{R-R} = 0 
$$

The action of $(-1)^F$ and $(-1)^{\tilde{F}}$ is now

$$
(-1)^F|Dp; \eta; k\rangle_{R-R} = (-1)^{7-p}|Dp; -\eta; k\rangle_{R-R} \\
(-1)^{\tilde{F}}|Dp; \eta; k\rangle_{R-R} = |Dp; -\eta; k\rangle_{R-R} 
$$

Therefore the GSO projected state in the R-R sector is

$$
|Dp; k\rangle_{R-R} = \frac{1}{\sqrt{2}} (|Dp; +; k\rangle_{R-R} + |Dp; -; k\rangle_{R-R}) 
$$

To make the state spacetime supersymmetric, we need to combine the NS-NS and R-R sectors appropriately. The result turns out to be $[18, 17]$

$$
|Dp; k\rangle = \frac{1}{\sqrt{2}} (|Dp; k\rangle_{NSNS} \pm 4i|Dp; k\rangle_{R-R}) 
$$

where the $+$, $-$ signs refer to branes and anti-branes respectively.

We still need to satisfy Cardy's conditions. This is achieved by taking a linear combination of the above states. In effect, this is a Fourier transform from momentum basis to position basis. The result is $[17]$ 

$$
|Dp; x\rangle = \int \prod_{l=0}^5 dk_l |Dp; k_l\rangle_D \times \prod_{j=6}^{6+p} \frac{1}{\sqrt{2R_j}} \sum_{n_j \in \mathbb{Z}} e^{-in_j x_j} |Dp; n_j\rangle_D \\
\prod_{\mu=7+p}^9 \sqrt{R_\mu} \sum_{m_\mu \in \mathbb{Z}} e^{-i2R_\mu m_\mu x_\mu} |Dp; m_\mu\rangle_N 
$$

where $n_j, m_\mu$ are the quantized momenta and winding in the compact directions.

The result is a D-brane state localized in the origin of the noncompact dimensions. The generalization to more general states is straightforward.
3 The T-Duality Orbifold

We now discuss the example of type IIA theory compactified on an asymmetric orbifold of $T^4$. The orbifold group is chosen to be $Z_2$, generated by a reflection of all left moving oscillators:

$$\alpha_n^i \to -\alpha_n^i \quad \tilde{\alpha}_n^i \to \tilde{\alpha}_n^i$$
$$\psi_r^i \to -\psi_r^i \quad \tilde{\psi}_r^i \to \tilde{\psi}_r^i$$

$$|F\rangle_{RR} \to \Gamma^1 \Gamma^2 \Gamma^3 \Gamma^4 |F\rangle_{RR}$$

where $i = 1, 2, 3, 4$.

This action is an overall T-duality, therefore we refer to this model as the T-duality orbifold. This action is only a symmetry of the torus at an $SO(8)$ point. We are interested in computing the massless spectrum of this model, using lightcone quantization of the RNS string.

In order for the model to be consistent (modular invariant), one has to check level matching. For an orbifold element of order $n$, one has to satisfy [11]:

$$E^L - E^R = 0 \mod \frac{1}{n}$$

where $E^L, E^R$ are the ground state energies of the left- and right- movers respectively. This guarantees that the physical states, which have to be invariant under the orbifold action, are level matched [13].

Here the only non-trivial check is for the $Z_2$ generator. The ground state energies for the left movers include contributions from the twisted bosons and fermions, as well as from the non-compact directions, and are found to be:

$$E^L_{R} = 0$$
$$E^L_{NS} = 0$$

$$E^R_{R} = 0$$
$$E^R_{NS} = -\frac{1}{2}$$

The right movers have the standard ground state energies:

$^{1}$To be precise, we must take the torus at the $SO(8)$ point, rather than the $SU(2)^4$ point, so that the zero mode contributions level match [14]. We thank R. Blumenhagen and E. Silverstein for correspondence on this point.
therefore level matching is satisfied in all 4 sectors of the RNS string.

The massless spectrum of the orbifold is identical to that of type II on a 4-torus, making it natural to conjecture that these two models are in fact equivalent, including their D-brane spectrum. To study the mapping between these two vacua we concentrate on the correspondence between space-time supersymmetry generators, as well as vector fields corresponding to the isometries of the 4-torus.

In the untwisted sector of the orbifold, one finds that 2 out of the 8 gravitini are projected out, leaving 6 gravitini (24 real components) in 6 dimensions. Specifically, in the R-NS sector, half of the components of the left-moving spin operator are not invariant under the orbifold action, and are projected out. We denote the surviving left moving spin operator by $S_\alpha$. It is defined by:

$$\Gamma^1\Gamma^2\Gamma^3\Gamma^4 S_\alpha = S_\alpha$$

Here $\alpha = 1, 2$ is an $SU(2)$ index, which is denoted by $SU(2)_I$.

On the other hand, two supersymmetry generators return in the twisted sector. The R-NS sector of the twisted sector contains one gravitino, made entirely from non-compact bosons and fermions. This is because the internal fermions in the twisted Ramond sector have no zero modes. The degeneracy of the twisted sector ground state is computed as in [1], and is found to be 2. Altogether the model has the maximal supersymmetry in six dimensions.

The study of the spacetime vectors is related to the isometries of the 4-torus, and is important in constructing D-brane states.

On the orbifold model, 4 vectors arise in the untwisted sector

$$V_L = \psi^\mu_{-1/2} \bar{\psi}^r_{-1/2} |F\rangle_{NSNS}$$

where the directions of the 4-torus in the orbifold model are denoted by the index $r = 1, ..., 4$.

The additional 4 vectors arise in the twisted sector as follows. The twisted sector vacuum for the left movers is generated by bosonic twist operators acting on the usual NS vacuum. There are two such twist operators, $\tau_\beta$, as the ground state degeneracy of the twisted sector is 2. Here $\beta$ is also an $SU(2)$ index; we denote this $SU(2)$ by $SU(2)_T$.

The bosonic twist operators $\tau_\beta$ have dimension $\frac{1}{4}$. In addition there is a fermionic twist operator, which in the $\mathbb{Z}_2$ case reduces to the usual spin operator. There are 2 spin operators, $S_\alpha$, that are invariant under the orbifold action. They have dimension $\frac{1}{4}$ as well, so the total twist operator carries
dimension $\frac{1}{2}$. In summary we find that the left moving ground state has degeneracy 4, which we denote by an $SO(4) = SU(2)_I \times SU(2)_T$ vector index $r$.

These twisted ground states are generated from the NS vacuum by 4 operators of dimension $\frac{1}{2}$, which are denoted for later convenience as $J^r$. It is easy to show that $J^r$ are GSO odd. Using the mapping of $SU(2)_I \times SU(2)_T$ into $SO(4)$ we can write:

$$J^r = (\Gamma^r)_{\alpha\beta} S^\alpha \tau^\beta$$  (22)

The additional 4 massless vectors are then:

$$V_R = \tilde{\psi}_{-1/2} J^r |F\rangle_{NSNS}$$  (23)

Note that they are GSO even as required.

We wish to compare these states with their counterparts in the torus description. Type II string theory on the 4-torus has 8 massless vectors in the NS-NS sector, related to momentum and winding modes on the 4-torus. The corresponding states are:

$$V_L = \psi_{-1/2} \tilde{\psi}_{-1/2} |F\rangle_{NSNS}$$

$$V_R = \tilde{\psi}_{-1/2} \psi_{-1/2} |F\rangle_{NSNS}$$  (24)

where $\psi^a, \tilde{\psi}^a$ are fermions in the 4-torus directions ($a = 1, ..., 4$), and $\psi^\mu, \tilde{\psi}^\mu$ are fermions in the noncompact directions.

It is then natural to identify the new torus fields as:

$$\tilde{\psi}^r \rightarrow \psi^a$$

$$J^r \rightarrow \psi^a$$  (25)

Similarly in the left and right-moving Ramond sectors, the ground states carry spinor indices of $SO(4)$. The spinor indices of the non-compact directions are suppressed. Denote, as above, the spinor index which is invariant under reflection by $\alpha$, and the other spinor index by $\tilde{\alpha}$. The ground state degeneracy in the twisted sector is parametrized by the index $\beta$.

With these notations, the right moving spinors of the new torus can be identified as $\tilde{S}_\alpha$ and $\tilde{S}_{\tilde{\alpha}}$. The left-moving spinors are $S_\alpha$ and $\tau_{\tilde{\beta}}$.

Having identified the torus symmetries in the asymmetric orbifold language, we are ready to use those symmetries to write the D-brane boundary states.
4 Boundary States

In the T-duality orbifold, it is possible to construct “untwisted” branes by starting with any brane on the torus, and adding its images under the $Z_2$ action. For example a 0-brane with any location on the torus is T-dual to a 4-brane wrapped on the torus, with some Wilson lines on its worldvolume. The $Z_2$ invariant boundary state is simply the sum of these two states. It is charged only with respect to untwisted sector fields, hence the name “untwisted”.

The interpretation of this state in the new torus description is not trivial. The locations and Wilson lines are separate moduli of the two components in the boundary state. As the state is required to be $Z_2$ invariant, its moduli space is $T^4$. It is therefore natural to interpret this untwisted state as a single D-brane state on the new torus. However, this is incorrect since the untwisted brane preserves only a quarter of the supersymmetry generators. In the equivalent torus description such states are made from two types of D-branes, and have a larger moduli space.

Therefore, similar to the closed string spectrum, the untwisted branes give only a subset of the expected D-brane states on the new torus. In particular, the naive moduli space of an untwisted brane is incorrect unless one allows more general twisted branes.

In order to construct the general D-brane state we are therefore forced to consider twisted branes, that is, branes that carry some twisted R-R charge. For a general asymmetric orbifold this is a difficult task. Constructing boundary states is usually done by imposing not just conformal invariance, but a more restrictive invariance under some chiral algebra. By construction, the chiral algebras for the left or right movers are different in most twist sectors. It is then impossible to use the usual ansatz to construct boundary states in each twist sector separately.

However, in the case considered in this paper, one can use the symmetries identified in the last section to construct boundary states. We start by constructing the NS-NS component of the boundary states.

The right movers of the new torus were identified as the fields $\tilde{\psi}^r$ of dimension $(0, \frac{1}{2})$ and their partners $\tilde{\partial}X^r$ of dimension $(0, 1)$. The left movers are the twist fields $J^r$ of dimension $(\frac{1}{2}, 0)$, and their partners denoted as $T^r$, of dimension $(1, 0)$. We denote the modes of any operator $O$ by $O_n$, except the modes of $\partial X$ and $\bar{\partial} X$, which are denoted by $\alpha_n$ and $\bar{\alpha}_n$ respectively.

It is now straightforward to write the Ishibashi states for an arbitrary
D-brane state. For example any D-brane with Dirichlet boundary condition in the torus directions (i.e. an unwrapped brane) has the factor:

\[ |B_\eta\rangle_{\text{NSNS}} = \exp \left( \sum_n T^r_{-n} \tilde{\alpha}^r_{-n} + i \eta \sum_n J^r_{-n} \tilde{\psi}^r_{-n} \right) |F\rangle_{\text{NSNS}} \] (26)

This is multiplied by the part of the boundary state coming from the non-compact directions, which is given in section 2. One has to sum over the states \( |B_\eta\rangle_{\text{NSNS}} \) in order to project into GSO invariant states, as one does in flat space.

As explained in section 2, in order to preserve supersymmetry, one has to consider also R-R boundary states. To find the correct combination of sectors to generate a BPS state we are guided by the explicit form of the spacetime supersymmetry operators. Intuitively, for the complete boundary state to be BPS saturated, the R-R boundary states are to be related to the NS-NS boundary state by spacetime supersymmetry operators (spectral flow).

In the present case the situation is similar. The supersymmetry generators from the left-moving sector are generated by the operators \( S_\alpha \) and \( \tau_\beta \) (the spin operators from the non-compact directions are suppressed). Those two operators create square root cuts in \( J^r \), defined in (22). Therefore we have to consider the sectors defined by the insertion of the supersymmetry generators at the origin. Those sectors define unusual, mixed, periodicity conditions for the original bosons and fermions, but act simply on the new variables \( J^r, T^r \). We refer to these sectors loosely as the Ramond sectors for the left movers.

In each such sector the fermionic currents \( J^r \) are now integer moded. The construction of the R-R boundary states is identical to the one in section 2. One obtains:

\[ |B_\eta\rangle_{\text{RR}} = \exp \left( \sum_{n>0} \frac{1}{n} (T^i_{-n} \tilde{\alpha}^i_{-n}) + i \eta \sum_{r>0} J^i_{-r} \tilde{\psi}^i_{-r} \right) |F_\eta\rangle_{\text{RR}} \] (27)

We have suppressed the dependence on the noncompact directions. The vacua \( |F_\eta\rangle_{\text{RR}} \) are the vacua of the R-R sectors, where the left-mover sector is a Ramond sector in the sense defined above. Those vacua have to be chosen now to satisfy:

\[ (J^i_0 - i \eta \tilde{\psi}^i_0) |F_\eta\rangle_{\text{RR}} = 0 \]
\[(J_0^\mu + i\eta \tilde{\psi}_0^\mu)|F_\eta\rangle_{RR} = 0\]  \hspace{1cm} (28)

The complete boundary state can be written as:

\[|B\rangle = \frac{1}{\sqrt{2}}(|B\rangle_{NSNS} \pm 4i|B\rangle_{RR})\]  \hspace{1cm} (29)

Furthermore, in order to obtain Neumann boundary conditions in any of torus directions one can modify the boundary state as described in section 2. The GSO projection acts, as in type IIA, asymmetrically. Hence, only even dimensional branes survive. We thus obtain the complete spectrum of D-branes in type IIA on \(T^4\).

As a final note we comment on the untwisted D-brane. Since some of the left-moving supersymmetry generators come from the twisted sector, it is clear that one cannot obtain a BPS state that breaks only half of the supersymmetry when one uses only the untwisted fields. Given the general D-brane state above, equation (29), one can recover any untwisted brane by an appropriate linear combination of those states. One then verifies the claim that the untwisted branes are two-object states which leave only a quarter of the supersymmetries unbroken.

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