Moduli Stabilization for Intersecting Brane Worlds in Type 0’ String Theory

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Abstract

Starting from the non-supersymmetric, tachyon-free orientifold of type 0 string theory, we construct four-dimensional brane world models with D6-branes intersecting at angles on internal tori. They support phenomenologically interesting gauge theories with chiral fermions. Despite the theory being non-supersymmetric the perturbative scalar potential induced at leading order is shown to stabilize geometric moduli, leaving only the dilaton tadpole uncanceled. As an example we present a three generation model with gauge group and fermion spectrum close to a left-right symmetric extension of the Standard Model.

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1 Introduction

Since it was discovered that one can have string excitations as light as a few TeV by allowing for large internal compactification spaces \cite{1,2}, string models without supersymmetry received new attention as phenomenologically plausible scenarios. This may either involve string theories without any supersymmetry already in ten dimensions or compactifications of originally supersymmetric theories with supersymmetry breaking at the string scale. A crucial problem in these constructions is the dynamical stability of the background which is no longer protected by any kind of no-force law. The most radical signal of an inherent instability is the presence of tachyons in the spectrum, scalars with a negative tree level mass squared. By now, the open string tachyons localized on some lower dimensional defect, a D-brane, in the entire ten-dimensional space-time are rather well understood \cite{3,4,5}. They are unstable modes of the gauge theory sector and their condensation translates into a decay of these D-branes into stable configurations but does not affect equally drastically the background space which the branes are wrapped on. The latter is indeed the case for closed string tachyons, unstable gravitational modes, which directly relate to a decay of the space-time. In field theoretic language, open string tachyon condensation and the induced D-brane recombination has its natural interpretation as the Higgs effect, namely the spontaneous breaking of the corresponding gauge group to that subgroup, which is associated to the stable D-brane configuration at the endpoint of the decay. Hence open string tachyons may indeed occur but are not harmful as they can play the role of the Higgs bosons of the effective gauge theory.

While one can often read the statement that a non-supersymmetric configuration is already considered stable if tachyons are absent, this is of course still a very simplistic view. In the absence of supersymmetry, quantum corrections will usually tend to generate potentials for the other massless scalar fields that very often drive the background either to some extreme limit or to a point where tachyons reappear. In particular the dilaton can be pushed to zero or infinite coupling. It will be one of the main objectives of this paper to construct non-supersymmetric D-brane models in which the geometrical moduli fields are stabilized.

There exist two possibilities to construct non-supersymmetric brane world models in the first place: (i) One starts from a supersymmetric theory in ten dimensions and only breaks supersymmetry in the open string, gauge theory sector on the branes. (ii) Supersymmetry is already broken right from the very beginning in the closed string sector. Most of the previous work follows the first pattern, namely one is considering supersymmetric type I and type II string theory where D6-branes intersect at relative angles on an internal torus or orbifold, which has been first pro-

\footnote{See also \cite{6,7} in this context.}
posed in [8, 9] and developed further in [10, 11, 12, 13, 14, 15, 16, 17, 18, 19, 20, 21]. These models display many attractive, generic features such as chiral fermion spectra in the four-dimensional effective theory [22, 23], which are localized at intersections of two D6-branes and therefore mostly transform in the bifundamental representation of the unitary gauge groups living on the respective D6-branes. Scalar fields either decouple or become tachyonic and may serve as Higgs bosons, as their condensation closely resembles standard spontaneous symmetry breaking patterns. In this way intersecting D-brane scenarios can be constructed which come very close to the non-supersymmetric Standard Model. However in the simplest case of D6-branes wrapped around homology 3-cycles of the six-dimensional torus, the background geometry is generically unstable.

In this paper we will discuss the second option (ii) as an alternative, novel approach for intersecting brane world models. To be concrete we will construct a brane world scenario of type 0' string theory [24, 25, 26, 27, 28, 29] with intersecting D6-branes, where we will closely follow the type I and type II D-brane models described before. The low energy description of our type 0' brane worlds has many appealing features close to Standard Model or GUT physics and at the same time stabilizes at least some of the closed string geometric moduli. In fact, while one may expect that these type 0' compactifications could lead to models with worse stability properties compared to the type II and type I D-brane models, the perturbative analysis appears to suggest the opposite: The formerly constructed toroidal intersecting brane worlds of type I string theory suffer from a perturbative instability that drives the complex structure moduli $U^{I}_{2}$ of the torus to a degenerate limit, $U^{I}_{2} \to \infty$, which implies an unacceptable squashing of the internal space. This instability could only be avoided by freezing these fields to particular values by imposing orbifold symmetries [13, 19]. We shall find that the type 0' brane worlds do not display this problem, but, on the contrary, their scalar potential stabilizes the $U^{I}_{2}$ at finite specific values, leaving only the dilaton tadpole uncanceled at leading order. At the same time they share all the generic properties of the former intersecting brane worlds of type I strings and only introduce slight modifications in the explicit construction. Of course, all of these statements can be subject to higher order perturbative or even non-perturbative corrections, which are hoped to stabilize the dilaton somehow, but may also affect the other moduli. Finally, open string tachyon condensation may again occur and play the role of the Higgs effect in the spontaneously broken gauge sector. However now the D-brane configuration after tachyon condensation will still be non-supersymmetric.

The paper is organized as follows: In section 2 we briefly review the construction of the ten-dimensional tachyon-free type 0' string theory and its basic properties. We then discuss the modifications that come into play by employing the intersecting brane world concept within this theory in section 3. Next we present the spectra of massless fermions, study the requirement of anomaly cancellation as a consistency
check and discuss the scalar potential of the complex structure moduli in section 4. Finally, in section 5, we present an example with the gauge group of a left-right symmetric extension of the Standard Model and corresponding spectrum of chiral fermions, whose complex structure moduli are being frozen at specific values.

2 Type 0′ string theory

The starting point is the non-supersymmetric ten-dimensional type 0B string theory [31, 32], which is known to be perturbatively unstable due its tachyonic ground state.

2.1 Type 0B string theory

Type 0 string theory can either be defined as a quotient of type II string theory by the space-time fermion number \((-1)^F_s\). Alternatively, one can start with the ten-dimensional superstring theory equipped with the modified GSO projection

\[ P_{GSO} = \frac{1 + (-1)^{F_L+F_R}}{2}. \]

leading to the type 0B torus amplitude

\[ T \sim \frac{1}{|\eta|^{16}} \left( |O_8|^2 + |V_8|^2 + |S_8|^2 + |C_2|^2 \right) \]

\[ \sim \frac{1}{2|\eta|^{24}} \left( |\vartheta[0]|^8 + |\vartheta[1/2]|^8 + |\vartheta[1/2]|^8 \right). \]

We have denoted by \(O_8\) etc. the characters of the \(SO(8)\) affine Lie-Algebra at level \(k = 1\). One peculiarity of this model is that compared to type II superstring theory all Ramond-Ramond (RR) forms are doubled. Type 0B string theory therefore contains even RR-forms \(C^\pm\) originating from the \((R+,R+)\) respectively \((R-,R-)\) sector, where the sign indicates the left respectively right moving worldsheet fermion number. Since each RR-form appears twice, there exist also two kinds of \(D(p-1)\)-branes, which couple to the fields \(C^\pm_p\) respectively. These D-branes are most conveniently be described in another basis, namely

\[ C_p = \frac{1}{\sqrt{2}} \left( C_p^+ + C_p^- \right), \quad C'_p = \frac{1}{\sqrt{2}} \left( C_p^+ - C_p^- \right), \]

where the two \(D(p-1)\)-branes are given by the boundary states [33, 34, 35]

\[ |D(p-1), \eta, \eta'\rangle = |D(p-1), \eta\rangle_{\text{NSNS}} + |D(p-1), \eta'\rangle_{\text{RR}}. \]

with \(\eta = \eta' = 1\) for a \(D(p-1)\)-brane and \(\eta = \eta' = -1\) for a \(D(p-1)'\)-brane. The freedom of choice of the sign \(\eta\) is due to the boundary condition of a worldsheet fermion at the position of a brane

\[ \psi^i_r - i\eta\tilde{\psi}^i_r = 0, \quad \psi^\mu_r + i\eta\tilde{\psi}^\mu_r = 0, \]
with indices $i$ referring to directions transverse to the brane and indices $\mu$ along the brane. Note, that in type II string theory only the superposition $D(p-1)+D(p-1)'$ is invariant under the GSO projection, so that only one D-brane of each even dimensionality survives.

### 2.2 Type 0' string theory

In \cite{24, 26} it was noted for the first time that by taking a particular orientifold of type 0 string theory one can get rid of the closed string tachyon. Dividing by the world-sheet parity $\Omega$ alone does not remove the tachyon but the combination $\Omega' = \Omega(-1)^{FR}$ \cite{28} does. Let us review what the effect of this projection on the RR-forms and D-branes is. The RR-forms $C_p^\pm$ transform as

$$
\Omega : \ C_p^\pm \rightarrow (-1)^{(p-2)/2} C_p^\pm ,
$$

$$
(-1)^{FR} : \ C_p^\pm \rightarrow \pm C_p^\pm ,
$$

so that the combined action gives

$$
\Omega' : \ C_p^\pm \rightarrow \pm (-1)^{(p-2)/2} C_p^\pm .
$$

Thus, the following forms survive the orientifold projection

$$
C_{10}^+, C_8^-, C_6^+, C_4^-, C_2^+, C_0^- .
$$

Summarizing, in type 0' string theory there exist only one RR-form of each even degree. The action of $\Omega'$ on the D-branes is

$$
\Omega'|Dp, \eta, \eta'\rangle = |Dp, -\eta, -\eta'\rangle
$$

implying that the symmetry of the brane spectrum under $\Omega'$ requires any respective Dp-brane to be accompanied by a D$p'$-brane.

Using these inputs, in \cite{24, 28} the complete string amplitudes at one loop order have been evaluated and the tadpole cancellation conditions deduced. They demand the introduction of $N + N' = 64$ branes, i.e. 32 of each type, and imply a maximal gauge symmetry $U(32)$. In the type I superstring theory $\Omega$ leaves the branes invariant and thus imposes a projection upon the open string excitations that leads to a gauge group $SO(32)$. Here, $\Omega'$ permutes the two kinds of branes and identifies their degrees of freedom, so that one is left with a unitary gauge group $U(32)$.

A property of this theory which will become important later on is that the orientifold planes do only couple to closed string modes from the RR-sector, implying that they have vanishing tension. This can be seen from the tree-channel
Klein-bottle and Möbius strip amplitudes

\[ \tilde{\mathcal{K}} = -2^{10} c \int_0^\infty dl \frac{\vartheta^{[1/2]}_0^4}{\eta^{1/2}}, \] (10)

\[ \tilde{\mathcal{M}} = 2^5 (N + N') c \int_0^\infty dl \frac{\vartheta^{[1/2]}_0^4}{\eta^{1/2}}, \] (11)

with arguments \( \tilde{q} = \exp(-4\pi l) \) and \( \tilde{\bar{q}} = -\exp(-4\pi l) \) respectively and the normalization \( c = \text{Vol}_{10}/(8\pi^2\alpha')^5 \). This leads to a modification of the scalar potential for the moduli fields in the way that the orientifold tension of type I string theory does not appear there.

The interaction between the D-branes at leading order are given by the annulus amplitudes. In the loop channel the amplitude between two D-branes of the same kind looks like

\[ A_{(Dp,Dp)} = \frac{N^2 + N'^2}{8} c \int_0^\infty dt \frac{\vartheta^{[0]}_0^4 - \vartheta^{[1/2]}_0^4}{t^6 \eta^{1/2}}, \] (12)

with argument \( q = \exp(-2\pi t) \). This leads to a repulsive force between the two D-branes, which means that the tension of the type 0 branes is not any longer balanced against their RR-charge. Moreover, it is evident from (12) that between two D-branes of the same type there are only open string excitations which are bosonic in space-time. Contrary, for the loop channel annulus amplitude for two D-branes of opposite type one obtains

\[ A_{(Dp, Dp')} = -\frac{NN'}{4} c \int_0^\infty dt \frac{\vartheta^{[1/2]}_0^4}{t^6 \eta^{1/2}}, \] (13)

leading to space-time fermions. Note, that even though the closed string sector is purely bosonic the open string sector contains space-time fermions. Moreover, the force between two D-branes of opposite type is attractive, as they only interact via exchange of closed string modes from the NSNS sector. This is clear, since the two D-brane are not charged under a common RR-form. However, adding the two annulus amplitudes yields a vanishing force similar to the no-force BPS situation, from which we can deduce that the tension of the type 0 branes is related to the tension of the type II branes via

\[ T_0 = \frac{T_{\Pi}}{\sqrt{2}}. \] (14)

The RR-tadpole cancellation condition is satisfied for \( N = N' = 32 \) leading to a gauge group \( U(32) \) with additional massless Majorana-Weyl fermions in the \( 496 \oplus 496 \) representation of \( U(32) \). Note, that the annulus is also the only amplitude
which receives contributions from the NSNS sector in the tree channel. In ten
dimensions this is attributed to the dilaton tadpole exclusively, which leads to a
run-away behavior for the string coupling. An effective way to deal with it may
consist in a suitable adaption of the Fischler-Susskind mechanism [36, 37, 38, 39].
In lower dimensions the scalar potential will also involve geometric moduli fields.

3 Type 0′ intersecting brane worlds

In [8] magnetic background fields on internal tori were shown to provide effective
means to construct interesting type I string compactifications with chiral fermions
in unitary gauge groups. By an exact perturbative duality transformation, a T-
duality along three of the internal circles, this setting of (non-commutative) D9-
branes with magnetic background fluxes is transformed into a completely equivalent
picture with (commutative) D6-branes intersecting at relative angles on the dual
tori [10, 11]. The later investigations of the prospects and properties of these models
revealed them to be among the most promising candidates to achieve a bottom-up
construction of a Standard Model or GUT field theory out of string theory. While
most effort was spent on studying non-supersymmetric models, some attention was
also devoted to supersymmetric constructions [42, 16, 43] which have close rela-
tions to M-theory vacua with a background space of $G_2$ holonomy.

In this fashion we now consider a compactification on an internal six-dimensional
complex torus

$$T^6 = T^2_1 \times T^2_2 \times T^2_3$$

(15)

with coordinates $(X^I, Y^I)$ on each $T^2_I$. Each two-dimensional torus is defined by
its complex structure $U^I = U^I_1 + iU^I_2$ and its Kähler structure $T^I = T^I_1 + iT^I_2$. In
order to apply the methods of intersecting brane worlds to the type 0′ theory, one
also needs to switch to a T-dual theory. The world sheet parity projection gets
mapped to $\Omega'\mathcal{R}$ where $\mathcal{R} : Y^I \mapsto -Y^I$ is a reflection along the three circles that
have been dualized. The theory with maximal gauge group now simply contains
32 D6- and 32 D6′-branes located along the fixed circles of $\mathcal{R}$, i.e. along $X^I$, to
cancel the tadpoles. Note, that the discrete parameters $U^I_a = b^I$ can take two
values $b^I \in \{0, 1/2\}$ in order to maintain the $\mathcal{R}$ symmetry of the background, cor-
responding to the two tori shown in figure 1.

The most general D6-brane of type $a$ wrapped on the $T^6$ and of codimension
one on any single $T^2_I$ is characterized by its winding numbers $(n^I_a, m^I_a)$ on the six
elementary circles. It is appropriate to use a basis

$$(n^I_a, m^I_a) = (n^I_a, m^I_a + b^I n^I_a)$$

(16)
denoting the wrapping around the cycles along $X^I$ and $Y^I$ respectively. The parity operation $\Omega'\mathcal{R}$ acts also on the winding numbers

$$\Omega'\mathcal{R} : D6(n_a^I, m_a^I) \mapsto D6(n_a^I, -m_a^I).$$

(17)

This is illustrated in figure 2.

Thus the symmetry of the brane spectrum requires to combine any D6-brane at some angle

$$\varphi_a^I = \arctan \left( \frac{m_a^I}{n_a^I U_2^I} \right)$$

(18)

relative to the $X^I$ axis with another D6'-brane at an opposite angle $-\varphi_a^I$. We henceforth employ the convention to use letters $a = 1...K$ for the branes of type
\( \eta = \eta' = 1 \) and \( a' = 1 \ldots K \) for those of type \( \eta = \eta' = -1 \). Since the two sets of stacks are identified under \( \Omega' \mathcal{R} \), there will be a gauge group \( U(N_a) \) on any individual stack with \( N_a = N_\omega \) being the number of branes. Similar to the type I models we expect chiral fermions at any intersection point of such branes.

In order to get a precise quantitative understanding we need to embark on the computation of the contributions to the tadpole divergencies at one loop order in string perturbation theory. The three relevant diagrams can most easily be found by combining the results of [28] and [8] and have been collected in the appendix. The upshot of this computation is summarized in the following: The Klein bottle does not involve any open strings. The oscillator part is thus given by (10) but momentum integrations need to be replaced by Kaluza-Klein (KK) and winding sums. One then gets

\[
\bar{K} = 2^{11} c_4 \prod_{I=1}^{3} (U_I^2) \int_0^\infty dl + \text{finite} \tag{19}
\]

for the tree channel contribution to the RR tadpole. We have defined \( c_4 = \text{Vol}_4/(8\pi^2\alpha')^2 \). Note again that there is no NSNS contribution whatsoever. The open string amplitudes involve \((D6,D6)\) and \((D6,D6')\) open strings separately. The Möbius strip receives only contributions from open strings invariant under \( \Omega' \mathcal{R} \), i.e. those stretching between branes of opposite type. Its RR tadpole reads

\[
\tilde{\mathcal{M}}_{aa'} = -2^6 c_4 N_a \prod_{l=1}^{3} (n_a^l U_2^l) \int_0^\infty dl + \text{finite}. \tag{20}
\]

Again, there is no NSNS tadpole. Finally, the annulus diagram involves all kinds of open strings, where the contribution to the RR tadpole arises merely from the \((D6,D6)\) and \((D6',D6')\) open strings

\[
\tilde{\mathcal{A}}_{ab}^{(RR)} = c_4 N_a N_b \prod_{l=1}^{3} \left( n_a^l n_b^l U_2^l + \bar{m}_a^l \bar{m}_b^l \frac{1}{U_2^l} \right)^3 \int_0^\infty dl + \text{finite}. \tag{21}
\]

The NSNS tadpole arises from all open string sectors in the annulus and will be dealt with in section 4 when the scalar potential is discussed.

Since in type 0' string theory there exists one RR-form of each even degree, in the original T-dual picture with magnetic fluxes on the D9- and D9'-branes we have to cancel the D9- as well as the effective D7-, D5- and D3-brane charges, leading to eight separate tadpole cancellation conditions. Note, that in type I string theory the D7- and D3-brane charges were automatically canceled due to the symmetry under \( \Omega \). Analogously, here the charges of the second set of RR-forms in type 0B are projected out by \( \Omega' \). Thus, the cancellation conditions which derive from the
above RR tadpole computation are very similar to those of type II string theory, except for a doubled background D9-brane charge:

\[ \sum_{a=1}^{K} N_a \prod_{I=1}^{3} n_a^I = 32, \quad \sum_{a=1}^{K} N_a \prod_{I=1}^{3} \bar{m}_a^I = 0, \]
\[ \sum_{a=1}^{K} N_a n_a^I n_a^J \bar{m}_a^K = 0, \quad \sum_{a=1}^{K} N_a n_a^I \bar{m}_a^J \bar{m}_a^K = 0. \] (22)

It is understood that the conditions are to be satisfied for any combination of indices \( I \neq J \neq K \neq I \). The gauge group

\[ G = \prod_{a=1}^{K} U(N_a) \] (23)

may have been reduced in its rank, where \( \text{rk}(G) < 32 \) is possible for \( n_a^I > 1 \). Note, that in contrast to intersecting brane world models of type I string theory one cannot obtain \( SO(2N_a) \) or \( Sp(N_a) \) gauge groups.

### 4 Chiral fermion spectra and scalar potential

The most important feature for the phenomenological relevance of intersecting branes is the fact that the open string quantization of strings between two D-branes that intersect in a point on the internal space leads to a single chiral fermion in the effective lower dimensional theory [23].

#### 4.1 Chiral fermions

This feature gets slightly modified in type 0' theory. As was pointed out above, only open strings stretched between D6-branes of opposite type give rise to space-time fermions, whereas open string stretched between two brane of the same sort give rise to space-time bosons. If the two intersecting branes are not identified by \( \mathcal{R} \), i.e. at an \((ab')\) or \((a'b')\) intersection, the fermion will simply transform in the bifundamental representation of the respective gauge groups, e.g. the \((N_a, N_{b'})\) of \( U(N_a) \times U(N_{b'}) \). If we are facing an \((aa')\) intersection instead, a further distinction has to be made. For intersection points that are fixed under \( \mathcal{R} \), one needs to regard the projection by \( \Omega' \mathcal{R} \). The Möbius strip amplitude (11) implies the solution

\[ \gamma_{\Omega' \mathcal{R}} = \bigotimes_{a=1}^{K} \begin{pmatrix} 0 & 1_{N_a} \\ 1_{N_{a'}} & 0 \end{pmatrix} \] (24)

for the action of \( \Omega' \mathcal{R} \) on the Chan-Paton labels. This leads to a single fermion in the antisymmetric representation of \( U(N_a) \) at such an invariant intersection of
type \((aa')\). On the contrary, for intersections not invariant under \(R\) no projection applies and they provide a symmetric and an antisymmetric representation for any doublet of intersection points. After defining

\[
I_{ab} = \prod_{I=1}^{3} \left( n_b^I m_a^I - n_a^I m_b^I \right), \quad I_{aa'}^{(R)} = \prod_{I=1}^{3} \left( 2m_a^I \right)
\]

(25)

for the total and \(R\)-invariant intersection numbers of branes \(a\) and \(b\) or \(a'\), the spectrum of chiral fermions is summarized in table 1.

| Representation             | Multiplicity |
|---------------------------|--------------|
| \((\mathbf{A}_a)_L\)      | \(I_{aa'}^{(R)}\) |
| \((\mathbf{A}_a \oplus \mathbf{S}_a)_L\) | \(\frac{1}{2} \left( I_{aa'} - I_{aa'}^{(R)} \right)\) |
| \((\mathbf{N}_a, \mathbf{N}_b)_L\) | \(I_{ab'}\) |

Table 1: Chiral massless fermions

In contrast to the intersecting brane world models of type I strings, where it was necessary to have \(b_i > 0\) for at least one torus in order to achieve odd numbers of generations, one can now also get three generation models with purely imaginary complex structures. It is interesting to note that the charged scalar fields can transform in different representations than the fermions. They arise in the sector of open strings with both ends on the same kind of brane and never experience the projection by \(\Omega'\mathcal{R}\). Thus, space-time bosons and in particular the Higgs fields can never appear in antisymmetric or symmetric representations of the gauge group.

### 4.2 Anomaly cancellation

It is not difficult to check that the spectrum of table 1 satisfies the cancellation of non-abelian anomalies. By using (22) the sum of all contributions to the triangle anomaly of the factor \(U(N_a)\) vanishes:

\[
\sum_{b' \neq a'} N_{b'} I_{ab'} + (N_a - 4) I_{aa'}^{(R)} + 2N_a \frac{1}{2} \left( I_{aa'} - I_{aa'}^{(R)} \right) = 0.
\]

(26)

This may serve as a consistency check, but is, of course, guaranteed by the tadpole cancellation anyway.
As usual, the mixed $U(1) - SU(N)^2$ and $U(1)^3$ anomalies do not cancel right away but require a suitable Green-Schwarz mechanism. For the $U(1) - SU(N)^2$ anomaly one obtains

\[ A_{aa} = \frac{1}{2} \sum_{b' \neq a'} N_{b'} I_{a b'} + \frac{(N_a - 2)}{2} I_{a a'}^{(R)} + 2N_a \frac{1}{2} \left( I_{a a'}^{(R)} - I_{a a'}^{(R)} \right), \]

\[ A_{ab} = \frac{1}{2} N_a I_{a b'} \quad \text{for } b' \neq a. \]  

(27)

Using the relation (27) this can be written as

\[ A_{ab} = \frac{1}{2} N_a I_{a b'} \quad \text{for all } a, b. \]  

(28)

Analogous to the type I string, we expect that this anomaly is canceled by a generalized Green-Schwarz mechanism invoking the coupling of four-dimensional bulk RR-fields. This issue is more conveniently addressed in the original type $0'$ theory, before any T-duality. The spectrum of RR-forms has been given in (5).

In detail, in type $0'$ string theory one has the following Wess-Zumino terms in the Born-Infeld action

\[
\int_{D_9} C_{0} F^5, \quad \int_{D_9} C_{2} F^4, \quad \int_{D_9} C_{4} F^3, \\
\int_{D_9} C_{6} F^2, \quad \int_{D_9} C_{8} F, \quad \int_{D_9} C_{10},
\]

(29)

These forms give rise via dimensional reduction to the following two-forms and axionic scalars in four dimensions

\[
C^I_2 = \int_{T^2_i} C^I_4, \quad C^I_0 = \int_{T^2_j} C^I_2, \\
B^I_2 = \int_{T^2_j \times T^2_k} C^I_6, \quad B^I_0 = \int_{T^2_j \times T^2_k} C^I_4, \\
B_2 = \int_{T^6} C^+_6, \quad B_0 = \int_{T^6} C^+_2
\]

(30)

and $C_2 = C^+_2, \quad C_0 = C^-_2$ with the following Hodge-duality relations in four dimensions

\[ dC_0 = \ast dB_2, \quad dB^I_0 = \ast dC^I_2, \quad dC^I_0 = \ast dB^I_2, \quad dB_0 = \ast dC_2. \]  

(31)

Thus, by integrating (29) over the internal six-dimensional torus, one obtains the following couplings

\[ N_a m_a m_a \int_{M_4} C_2 F, \quad n_b n_b n_b \int_{M_4} B_0 F_b^2, \]
which combine into the usual tree diagrams to contribute to the respective anomaly. Adding up all terms $F_a F_b^2$ we find that the resulting GS amplitude is proportional to

$$A_{GS} = N_{ab} I_{ab}$$

which has the correct form to cancel the field theory anomaly (28).

### 4.3 The disc level scalar potential

The scalar potential can be read off from the dilaton tadpole divergence of the closed string tree channel amplitude which is essentially the square of the disc expectation value

$$\langle \phi \rangle_{\text{disc}} \sim \frac{\partial V(\phi, U^I_{2})}{\partial \phi}.$$  

Due to the absence of the NSNS contributions in the Klein bottle and Möbius strip diagrams, the resulting potential is identical to the one obtained in type I string theory except for the absence of the orientifold tension,

$$V(\phi, U^I_{2}) = T_6 e^{-\phi_4} \sum_{a=1}^{K} (2N_a) \prod_{I=1}^{3} \sqrt{(n_a^I)^2 U^I_{2} + (m_a^I)^2 U^I_{2}},$$

where $\phi_4$ is the four-dimensional dilaton $\phi_4 = \phi - \ln(\text{Vol}_6)/2$ and $T_6$ the D6-brane tension. The normalization was fixed by comparing to the DBI effective action. Already from the scaling behavior of the potential

$$V(\phi, U^I_{2}) \rightarrow \infty \quad \text{for} \quad U^I_{2} \rightarrow 0, \infty$$

it is clear that there must exist a global minimum at which the $U^I_{2}$ are stabilized at tree level. We will demonstrate this explicitly by presenting an example with stabilized complex structure in the following section.

The potential (35) does not depend on the Kähler moduli $T^I$. It was explained in [15] that it is to be expected that this feature remains true perturbatively if the D6-branes intersect in points. Only if there are also branes that are parallel on circles of the internal space the propagation of KK and winding modes will introduce a dependence on the $T^I$. It was estimated from their respective proportionality to $1/T^I$ and $T^I$ that at least to the next to leading order, in the annulus diagram, the Kähler moduli may stabilize as well.
5 A left-right symmetrically unified model

In this section we present a concrete model which contains the gauge group of a left-right symmetric extension of the Standard Model with some additional charged chiral matter. Since, in contrast to the type II orientifold models, the type 0' orientifold models only have chiral fermions for $ab'$ intersections, the anomaly cancellation conditions for the $U(N_a)$ gauge factors (including $N_a = 1, 2$) prevent the realization of a pure three generation Standard Model. Recall that for the type II orientifold models two of the left-handed quark doublets transformed in the $(3, 2)$ representation and one in the $(\overline{3}, 2)$ representation. In this way, the $(ab)$ and $(ab')$ intersections could combine to provide three generations, which is now impossible, as only $(ab')$ will provide fermions at all.

We now choose five stacks of D6-D6' branes, $b^I = 1/2$ on all three two-dimensional tori $T^2_I$ and the wrapping numbers shown in table 2.

| $(n^1_a, \bar{m}^1_a)$ | $(n^2_a, \bar{m}^2_a)$ | $(n^3_a, \bar{m}^3_a)$ | $N_a$ |
|------------------|------------------|------------------|-----|
| $(2, 0)$         | $(1, -\frac{1}{2})$ | $(4, 1)$         | 3   |
| $(1, \frac{1}{2})$ | $(1, \frac{3}{2})$ | $(1, \frac{1}{2})$ | 2   |
| $(1, -\frac{1}{2})$ | $(1, \frac{3}{2})$ | $(1, \frac{1}{2})$ | 2   |
| $(2, 0)$         | $(1, \frac{7}{2})$ | $(1, -\frac{1}{2})$ | 1   |
| $(2, 0)$         | $(1, -\frac{1}{2})$ | $(1, -\frac{7}{2})$ | 1   |

Table 2: Wrapping numbers

It can easily be checked that these wrapping numbers satisfy the RR-tadpole cancellation conditions. The resulting massless chiral spectrum is presented in table 3.

This is the matter content of the minimal left-right symmetric extension of the Standard Model plus three generations of a charged $(A, 1) \oplus (1, \overline{A})$ of $SU(2)_L \times SU(2)_R$. Note, that the fourth $U(1)$ factor decouples completely from the chiral spectrum and therefore can be considered as resulting from a spectator brane which is only there to satisfy the RR cancellation conditions. Two of the five $U(1)$ factors are anomalous and besides $U(1)_4$ the anomaly free ones are

\[
U(1)_{B-L} = \frac{1}{3} U(1)_1 + U(1)_5, \\
U(1)_K = U(1)_2 + U(1)_3 + 2 U(1)_5. \tag{37}
\]

It is beyond the scope of this paper to follow the phenomenological properties of this model further. Instead, we investigate the disc level scalar potential for this
SU(3) × SU(2)_L × SU(2)_R × U(1)^5 \quad \text{Multiplicity}

| (3, 2, 1)_{(1,1,0,0,0)} | 3 |
| \overline{(3, 1, 2)}_{(-1,0,-1,0,0)} | 3 |
| (1, 2, 1)_{(0,-1,0,0,-1)} | 3 |
| (1, 1, 2)_{(0,0,1,0,1)} | 3 |
| (1, A, 1)_{(0,2,0,0,0)} | 3 |
| (1, 1, 
 \overline{A})_{(0,0,-2,0,0)} | 3 |

Table 3: Chiral massless spectrum

By performing a numerical analysis of the resulting potential (35) we find that there exist a unique global minimum for the values of the complex structures

\[ U_2^1 = 0.086, \quad U_2^2 = 1.045, \quad U_2^3 = 0.486. \] (38)

We have also determined the ground state energies \( E_0 \) in the various bosonic open string sectors at this minimum where depending on the values of the complex structures tachyons can potentially arise. It is another nice feature of these type 0′ models that only \((ab)\) type intersections need to be considered, thus the number of dangerous sectors is essentially halved. If two D6-branes are parallel on at least one \( T^2 \), then we can move the two D-branes apart on this torus and avoid the tachyonic instability classically. For the remaining open string sectors we find the following ground state energies:

| Sector | \( E_0 \) |
|--------|--------|
| (12)   | 0.05   |
| (13)   | 0.05   |
| (24)   | 0.02   |
| (34)   | 0.02   |
| (25)   | 0.09   |
| (35)   | 0.09   |

Table 4: Ground state energies
All ground state energies are positive implying that this model, in the approximation we used, is perturbatively stable. Note, that the possible tachyon in the (23) open string sector transforms in the (2, 2) representation and is needed for breaking the left-right symmetric down to the Standard Model. However, as we mentioned already there can be no tachyons and therefore no Higgs particles in symmetric or anti-symmetric representations of the gauge group.

6 Conclusions

In this paper we have constructed a class of non-supersymmetric string models whose geometric complex structure moduli are stabilized at finite values by the leading order perturbative potential. One may further argue, that the stabilization can be extended to the Kähler moduli as well, hence leaving only the dilaton unstable perturbatively. Our novel construction is intrinsically non-supersymmetric, as it starts from the non-supersymmetric but tachyon-free ten-dimensional type 0’ string theory compactified on a torus. To this theory we have applied the techniques of intersecting brane worlds finding a set of solutions for the effective four-dimensional field theory which displays perspectives for obtaining semi-realistic models that are comparative to the earlier studied type I and type II compactifications. One may consider various extensions and generalizations of the present program. It may for instance be interesting to construct orbifolds of the purely toroidal models along the lines of [44, 45].

Acknowledgements

We are supported in part by the EEC RTN programme HPRN-CT-2000-00131. We would like to thank A. Uranga for discussion.

A Amplitudes of type 0’ with branes at angles

In this appendix we summarize the more technical results for the three amplitudes that contribute to the tadpole divergence at the \(\chi = 0\) level of perturbation theory. They are obtained most directly by combining the results of [28] and [8]. The Klein bottle amplitude of the original ten-dimensional type 0’ theory was given in (10). The modification due to the toroidal compactification is standard and does not take reference to the novel issue of intersecting branes,

\[
\tilde{\mathcal{K}} = -2^7 c_4 \prod_{I=1}^{3} (U_2^{I}) \int_0^\infty dl \frac{\eta^{1/2}}{\eta^{1/2}} \prod_{I=1}^{3} \left( \sum_{r,s \in \mathbb{Z}} e^{-4\pi l (r^2 R^2 + s^2 / R^2)} \right).
\] (39)
For the open string amplitude the peculiarities of type 0′ as well as the modifications due to the relative angles of the D6-branes need to be taken into account. The open string states involve only the NS sector but one needs to include the D6- and D6′-branes now. The relative angles lead to shifts in the oscillator spectrum of string excitations by

$$\epsilon_{ab}^I = \frac{\varphi_{b}^I - \varphi_{a}^I}{\pi} \quad (40)$$

and change the KK and winding zero mode quantization [40]. The annulus diagram for strings with both ends on the same brane a then reads

$$\tilde{A}_{aa} = 2^{-4} c_4 N_a^2 \prod_{I=1}^{3} \left( (m_{a}^I)^2 U_2 + \frac{(m_{a}^I)^2}{U_2^2} \right) \int_0^\infty dl \frac{\vartheta[0]_I^4 - \vartheta[1/2]_I^4}{\eta^{12}} \prod_{I=1}^{3} \left( \sum_{r,s \in \mathbb{Z}} e^{-\pi l \tilde{M}_I^2} \right)$$

which is to be combined with an identical contribution from the a′ branes. The annulus diagram for strings between two different branes a and b (or a′ and b′) but both of the same type is

$$\tilde{A}_{ab} = 2^{-1} c_4 N_a N_b I_{ab} \int_0^\infty dl \frac{\vartheta[0]_I^4 \prod_{I} \vartheta[0]_{ab_I} - \vartheta[1/2]_I^4 \prod_{I} \vartheta[1/2]_{ab_I}}{\eta^3 \prod_{I} \vartheta[1/2+\epsilon_{ab}^I]} \quad (41)$$

while for open strings between branes a and b′ (or a′ and b) one finds

$$\tilde{A}_{ab′} = -2^{-1} c_4 N_a N_{b′} I_{ab′} \int_0^\infty dl \frac{\vartheta[1/2]_I^4 \prod_{I} \vartheta[0]_{ab′_I} - \vartheta[1/2]_I^4 \prod_{I} \vartheta[1/2]_{ab′_I}}{\eta^3 \prod_{I} \vartheta[1/2+\epsilon_{ab′}^I]} \quad (42)$$

Together this combines into the annulus diagram of type II strings, which in particular guarantees the absence of any tachyon contribution to the tadpole divergence. Finally, the Möbius strip diagram gives the contribution

$$\tilde{M}_{aa′} = 2^2 c_4 N_a I_{aa′}^{(R)} \int_0^\infty dl \frac{\vartheta[1/2]_I^4 \prod_{I} \vartheta[1/2]_{a'/\pi}}{\eta^3 \prod_{I} \vartheta[1/2+\varphi_{a}/\pi]} \quad (43)$$

We have used the standard definitions

$$\frac{\vartheta[\alpha \beta]_I(q)}{\eta(q)} = e^{2\pi i \alpha \beta} q^{\frac{\alpha^2}{2} - 1/24} \prod_{n=1}^{\infty} \left( 1 + q^{n-1/2+\alpha e^{2\pi i \beta}} \right) \left( 1 + q^{n-1/2-\alpha e^{-2\pi i \beta}} \right),$$

$$\eta(q) = q^{1/24} \prod_{n=1}^{\infty} (1 - q^n). \quad (44)$$
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