Standard Model Compactifications of IIA $Z_3 \times Z_3$
Orientifolds from Intersecting D6-branes

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Abstract

ABSTRACT

We discuss the construction of chiral four dimensional $T^6/(Z_3 \times Z_3)$ orientifold compactifications of IIA theory, using D6-branes intersecting at angles and not aligned with the orientifold O6 planes. Cancellation of mixed U(1) anomalies requires the presence of a generalized Green-Schwarz mechanism mediated by RR partners of closed string untwisted moduli. In this respect we describe the appearance of three quark and lepton family $SU(3)_C \times SU(2)_L \times U(1)_Y$ non-supersymmetric orientifold models with only the massless spectrum of the SM at low energy that can have either no exotics present and three families of $\nu_R$’s (A’-model class) or the massless fermion spectrum of the N=1 SM with a small number of massive non-chiral colour exotics and in one case with extra families of $\nu_R$’s (B’-model class). Moreover we discuss the construction of SU(5), flipped SU(5) and Pati-Salam $SU(4)_c \times SU(2)_L \times SU(2)_R$ GUTS - the latter also derived from adjoint breaking - with only the SM at low energy. Some phenomenological features of these models are also briefly discussed. All models are constructed with the Weinberg angle to be 3/8 at the string scale.
1 Introduction

Four dimensional perturbative (4D) chiral compactifications (CS) from string theory have a long history. The first perturbative chiral models started with the N=1 4D closed string compactifications of the heterotic string - on either Calabi-Yau manifolds [1], orbifold [2], self-dual lattices [3], Gepner type [4], fermionic constructions [5] or orbifolds with Wilson lines [6]. In these vacua the string scale is of the order of $10^{18}$ GeV, the gauge couplings unify at the same scale and as a consequence - of the high scale - proton is stable. As the string scale is high N=1 susy models are favored phenomenologically, also offering simultaneously a solution to the gauge hierarchy problem. However, these N=1 vacua face severe problems related to the breaking of supersymmetry - creating a non-zero cosmological constant - and also lack satisfactory moduli fixing mechanisms.

On the other hand compact type II orientifold [7] compactifications offer a window into perturbative physics - that has been used in recent years to study supersymmetric and non-supersymmetric D-brane models (See [8] for reviews)- as in type I compactifications the gauge hierarchy problem could be solved by the existence of dimensions transverse to the space that the D6-branes wrap [9]. In particular, stringy D-brane models derived from orientifolds of type II string compactifications that include D6-branes intersecting at angles (IBs) [13-46] and wrapping three cycles in the six dimensional internal space, provide a consistent string framework that determines all the physical quantities in terms of the brane angles and their wrappings. In these constructions chiral fermions get localized in the intersections between branes [12].

Various compact chiral IIA orientifold N=1 supersymmetric constructions with intersecting D6-branes have been produced including orbifolds of $T^6/Z_2 \times Z_2$ [16], $T^6/Z_4$ [27], $T^6/Z_4 \times Z_2$ [31], $T^6/Z_6$ [32]. However, the effect of N=1 supersymmetry seriously affects the spectrum, as in all cases the spectra are semi-realistic with the N=1 SM accompanied by either massless chiral [16, 31] (class $\tilde{A}$) or massless non-chiral exotics (class $\tilde{B}$) [27, 32]. We also note that chiral models could be produced from orientifolds of Gepner models [33] using techniques borrowed from intersecting branes. In this case the models of class $\tilde{B}$ have been produced [34].

On the other hand, the first attempts in a string theory context resulting in non-susy chiral constructions with intersecting D6-branes wrapping 3-cycles | In the T-dual language these models correspond to models with magnetic deformations [10, 26] | have been carried out in compactifications of type IIA on either tori [14] or toroidal orientifolds $IIA/(T^6/\Omega R)$ [13], the latter using the T-dual picture with D9-branes and background fluxes turned on, following the work of [11] on the gauge theory aspects of
magnetic fluxes. Three generation non-supersymmetric models with no extra exotics have been found in toroidal orientifolds (TO)\cite{13} or $Z_3$ orientifolds \cite{15}. Indeed using the constructions \cite{13} it has become possible to derive - using only bifundamental representations - non-supersymmetric vacua that possess *just the observed Standard Model* (SM) spectrum with right handed neutrinos and gauge interactions at low energies \cite{19,20,21} and also Pati-Salam vacua with the same property \cite{22,23}. [In \cite{15}, non-supersymmetric Standard model vacua have been also derived from $Z_3$ orientifolds, but due to existence of antisymmetrics in the spectrum there were no mass terms for the up-quarks.] Important results of these TO constructions \cite{23} include a) *The fixing of all complex structure moduli using N=1 supersymmetry in open string sectors in toroidal orientifolds* \footnote{e.g. see eqn. (4.37) in hep-th/0203187} and the fact that b) *N=1 supersymmetry conditions that introduce supersymmetric sectors - in toroidal orientifolds - solve the condition that the hypercharge survives massless the Green-Schwarz anomaly cancellation* \footnote{The latter mechanism acts as a mass generation mechanism to $U(1)$’s that have a non-zero coupling to RR fields.} mechanism.

In all D-brane models coming from intersecting branes, the SM is accompanied by the simultaneous existence of right handed neutrinos; necessary for RR tadpole cancellation [other proposals in D-brane model building but not based to a particular string construction can be seen in \cite{50,51}].

The purpose of this work, is twofold. Firstly to discuss the main features of the $Z_3 \times Z_3$ orientifolds and to also show that it is not only possible to produce (3-stack) non-supersymmetric models with only the SM at low energy - reproducing the SM fermion spectra of \cite{15} - but also derive new non-supersymmetric vacua which localize the fermion spectrum of the N=1 SM (3- and 5-stack) and extra massive exotics which subsequently break to the SM. These SMs exhibit partial gauge unification of the strong and weak gauge couplings with a Weinberg angle $\sin^2 \theta = 3/8$, as it has been also shown in the split susy scenario \cite{46} [see also section 8 of \cite{47} for different models - distinguished by the different electric and magnetic wrappings that solve the RR tadpoles and intersection numbers - with the same spectrum]. Secondly, to show that GUT models could be constructed. We present explicit examples with two stack flipped $SU(5)$ GUTs and three stack Pati-Salam $SU(4)_c \times SU(2)_L \times SU(2)_R$ GUTS.

The paper is organized as follows. In section two we discuss the formalism of the chiral constructions on the $Z_3 \times Z_3$ orientifolds. We describe explicitly the derivation of the constraints coming from RR tadpole cancellation at the string scale. RR tadpoles constitute an important constraint in string model building as their presence is
equivalent to the constraints coming from the cancellation of cubic gauge anomalies in
the low energy effective theory. We also describe the spectrum rules for the \( Z_3 \times Z_3 \)
orientifold vacua that are based on the AAA tori lattices. The explicit form of the
effective wrappings is seen in appendix A. In addition spectrum rules and RR tadpoles
for different \( Z_3 \times Z_3 \) orientifold vacua - for which we will not present explicit non-
susy or \( N=1 \) supersymmetric models in this work - are presented in appendix B. In
section 3 we discuss the structure of the \( U(1) \) anomaly cancellation necessary for the
cancellation of \( U(1) \)-mixed gauge/gravitational anomalies. The rest of the sections is
dedicated to the study of non-supersymmetric models which in most of the cases exhibit
partial unification of their gauge couplings, that is they unify - at the string scale -
the strong \( SU(3)_c \) and the weak \( SU(2) \) gauge couplings with the Weinberg angle to be
\( \sin^2 \theta = 3/8 \) as in the successful \( SU(5) \) GUT prediction. In section 4 we present three
stack non-susy models which break to only the SM at low energy without any massless
exotics present. These models have been also found before in the \( Z_3 \) orientifolds of [15].
In section 5, we derive other three stack three generation non-supersymmetric models
with the Weinberg angle \( \sin^2 \theta = 3/8 \) which localize the massless fermion spectrum of
the \( N=1 \) supersymmetric SM (with extra generations of \( \nu_R \)'s) in the presence of one
pair of non-chiral exotics. Eventually, the extra beyond the SM massless fermions and
the extra exotics become massive leaving only the SM at low energy [In our companion
paper [47] we presented SMs which generate models with the same chiral spectrum but
for different wrapping numbers].
In section 6 - even though the present models do not exhibit a supersymmetric spec-
trum - we discuss the split susy scenario which was proposed as an alternative signal
for LHC. Hence a comparison with split susy models (SSS) that appear recently in
the literature is also performed as it appears that even though the \( Z3 \times Z3 \) models
are not SSS, they do possess many of their properties. As an application, we discuss
five stack vacua with the spectrum of the \( N=1 \) SM and massive exotics. The models
break to the SM at low energy as the Higgsinos form massive Dirac pairs. In section
7, we discuss the construction of flipped \( SU(5) \) GUTS with the SM at low energy.
Models with the same fermion spectrum have been produced in [43] based on the \( Z_3 
orientifolds [15]. In sections 8, 9 we discuss the construction of \( SU(5) \) and Pati-Salam
\( SU(4)_L \times SU(2)_L \times SU(2)_R \) (with adjoint breaking) non-supersymmetric three family
GUTS respectively. Section 10 contains our conclusions.
2 RR Tadpoles and Spectrum rules for $\mathbb{Z}_3 \times \mathbb{Z}_3$ Orientifolds

In this section we will describe the spectrum rules and the constraints on the parameters (wrappings) that have to be imposed on model building attempts on these constructions. The existence of these rules is independent of the presence of supersymmetry - in the open string sector - where the chiral matter of the Standard model gets localized.

2.1 The background and the RR tadpole cancellation

Our constructions are derived from IIA theory compactified on a six dimensional torus modded out by the orbifold action $(\mathbb{Z}_3 \times \mathbb{Z}_3)$, where the latter symmetry is generated by the twist generators

\begin{align*}
\theta : (z_1, z_2, z_3) &\rightarrow (\alpha z_1, \alpha^{-1} z_2, z_3), \\
\omega : (z_1, z_2, z_3) &\rightarrow (z_1, \alpha z_2, \alpha^{-1} z_3),
\end{align*}

and where $\theta, \omega$ gets associated to the action of the twists $\nu = \frac{1}{3}(1, -1, 0), u = \frac{1}{3}(0, 1, -1)$. Here, $z^i = x^{i+3} + i x^{i+5}, i = 1, 2, 3$ are the complex coordinates on $T^6$, which we consider as being factorizable for simplicity reasons, namely $T^6 \equiv T^2 \otimes T^2 \otimes T^2$. In addition to the orbifold action, IIA is also modded out by the orientifold action $\Omega R$, that combines the worldsheet parity $\Omega$ and the antiholomorphic involution

\[ R : z^i \rightarrow \bar{z}^i. \]  

The orbifold action has to act crystallographically on the lattice. For this reason, we will let the complex structure on all three $T^2$ tori to be fixed at

\[ U_A^I = \frac{1}{2} - i \frac{\sqrt{3}}{2}, \]  

We associate the presence of the above complex structure with the A-torus. The A-torus - will be used in the main body of the paper in which we discuss model building using the AAA tori - is a modified version of the A-torus that appear in [15], as we have chosen the lattice vectors to be $e_1 = (1, 0), e_2 = (-1/2, \sqrt{3}/2)$. The orbifold action (2.1) preserves N=2 SUSY in four dimensions and thus the orbifold describes the singular limit of a Calabi-Yau threefold. Using the cohomology of the internal orbifold space we get from the Hodge numbers describing this threefold $^4$ that in the

\footnote{This class of orbifolds corresponds to models without discrete torsion.}
closed string sector are \( h^{11} = 84, h^{21} = 0 \), where three K"ahler moduli come from the untwisted sector and the rest from the twisted sectors. As the numbers of the independent three cycles is \( b_3 = 2 + 2h^{21} \) all the independent cycles are inherited from the six dimensional toroidal space. There are no exceptional cycles involved and hence we will need only toroidal cycles, as in the orientifolds of \( T^6/Z_3 \) \([15]\), \( T^6/(Z_2 \times Z_2) \) \([16]\).

The orientifold projection breaks the bulk supersymmetry to \( N=1 \) and introduces orientifold O6-planes locked at the fixed locus of the antiholomorphic involution \((2.2)\). The O6-planes carry negative RR charge whose presence - introduces an inconsistency into the theory as the partition function diverges - can be cancelled by the introduction of D6-branes intersecting at angles with the O6-planes. The models that are derived are chiral as the D6-branes are not parallel to the orientifold planes [Non-chiral 4D models on \( Z_3 \times Z_3 \) orientifolds of type IIA have been considered in the past, in the presence of intersecting D-branes parallel to the O6-planes in \([29]\) and in a different content in \([30]\)]. The D6-branes are assumed to be wrapping 3-cycles along the toroidal space, with the 1-cycles being described by the (‘electric’-‘magnetic’) numbers respectively \((n^i_a, m^i_a)\), indicating wrapping along \( i \)-th cycle of the \( T^2 \) tori. There are four possible tori choices, allowing for a consistent twist action within the \( Z_3 \times Z_3 \) orientifold that is the AAA, AAB, ABB, BBB one’s. In this work we will exhibit model building attempts which are based on the AAA tori in the main body of this paper. In appendix B, we will derive the spectrum rules and RR tadpoles for the vacua that accommodate the AAB, BBB tori using the A, B torus choices made in \([15]\). Under the orbifold and orientifold action the branes are organized into orbits of length nine. These orbits are characterized by pairs of wrapping numbers \((n^i, m^i)\). Especially for the AAA torus that we will treat in detail in this work, these orbits are described by

\[
\begin{pmatrix}
  n^i \\
  m^i
\end{pmatrix}
\xrightarrow{\Omega R \uparrow}
\begin{pmatrix}
  -m^i \\
  n^i - m^i
\end{pmatrix}
\xrightarrow{\Omega R \uparrow}
\begin{pmatrix}
  -n^i + m^i \\
  -n^i
\end{pmatrix}
\]

\[
\begin{pmatrix}
  n^i - m^i \\
  -m^i
\end{pmatrix}
\xleftarrow{\Omega R \uparrow}
\begin{pmatrix}
  -n^i \\
  m^i - n^i
\end{pmatrix}
\xrightarrow{\Omega R \uparrow}
\begin{pmatrix}
  m^i \\
  n^i
\end{pmatrix}
\]

(2.4)

We also denote the homology class of the \( i \)-th cycle of the \( a \)-th D6-brane as being defined as

\[
[\Pi_a] = \prod_{i=1}^3 (n^i_a \ [a_i] + m^i_a \ [b_i])
\]

(2.5)

We also denote by \([a_i], [b_i]\) the basis for the homology cycles across the corresponding 5
i-th two-tori lattice of the decomposable six-dimensional tori $T^6 = T^2 \times T^2 \times T^2$. The total homology class of the 06 planes is defined as

$$[\Pi_{O6}] = [\Pi_{\Omega R}] + [\Pi_{\Omega R\omega}] + [\Pi_{\Omega R\omega^2}] + [\Pi_{\Omega R\theta}] + [\Pi_{\Omega R\theta^2}] + [\Pi_{\Omega R\theta^2\omega}] + [\Pi_{\Omega R\theta^2\omega^2}] + [\Pi_{\Omega R\theta\omega}]$$

(2.6)

In turn the individual homology classes of the cycles describing the action of spacetime and worldsheet symmetries for the $\Omega R$, $\Omega R\theta$, $\Omega R\omega$, $\Omega R\omega^2$, $\Omega R\theta^2$, $\Omega R\theta^2\omega$, $\Omega R\theta^2\omega^2$, $\Omega R\theta^2\omega^3$, $\Omega R\theta^3\omega$, $\Omega R\theta^3\omega^2$, actions on the 06-planes are given by

$$[\Pi_{\Omega R}] = [a_1] \times [a_2] \times [a_3], \quad [\Pi_{\Omega R\omega}] = -[a_1] \times ([a_2] + [b_2]) \times [b_3],$$

(2.7)

$$[\Pi_{\Omega R\omega^2}] = -[a_1] \times [b_2] \times ([a_3] + [b_3]), \quad [\Pi_{\Omega R\theta}] = -([a_1] + [b_1]) \times [b_2] \times [a_3],$$

(2.8)

$$[\Pi_{\Omega R\theta^2}] = -[b_1] \times ([a_2] + [b_2]) \times [a_3], \quad [\Pi_{\Omega R\theta^2\omega}] = -([a_1] + [b_1]) \times [a_2] \times [b_3],$$

(2.9)

$$[\Pi_{\Omega R\theta^2\omega^2}] = [b_1] \times [b_2] \times [b_3], \quad [\Pi_{\Omega R\theta^2\omega^3}] = -[b_1] \times [a_2] \times ([a_3] + [b_3]),$$

(2.10)

$$[\Pi_{\Omega R\theta^3\omega}] = -([a_1] + [b_1]) \times ([a_2] + [b_2]) \times ([a_3] + [b_3])$$

(2.11)

The O6-planes fixed under the orientifold actions $\Omega R$, $\Omega R\omega$, $\Omega R\omega^2$, $\Omega R\theta$, $\Omega R\theta^2$, $\Omega R\theta\omega$, $\Omega R\theta^2\omega$, $\Omega R\theta^2\omega^2$, $\Omega R\theta^3\omega$ can be seen in figure 11. In $\Omega R$ orientifolds twisted crosscap tadpoles vanish $^{28,29}$, thus the orientifolds of $T^6/Z_3 \times Z_3$ possess only untwisted RR tadpoles. The images of a D6-brane under a $Z_3 \times Z_3$ action may be denoted by $[\Pi_{a'}]$. The RR tadpole cancellation condition is equivalent to the vanishing of the RR charge in homology

$$\sum_{a} N_a \ [\Pi_a] + \sum_{a'} N_{a'} \ [\Pi_{a'}] + \ (-4) \times [\Pi_{O6}] = 0,$$

(2.12)

where $-4$ is the charge of the O6-plane. Summing over the all the different homology classes the tadpole conditions reduce to the general condition

$$\sum_{a} N_a \ Z_{[a]} = 4,$$

(2.13)

where $Z_{[a]}$ is given in appendix A. A comment is in order. As all complex structure moduli are fixed, in N=1 supersymmetric constructions all NSNS tadpoles are absent. In non-supersymmetric constructions the only remaining NSNS disc tadpole is the one associated to the dilaton. However due to the absence of complex structure moduli in the $Z_3 \times Z_3$ constructions the NSNS potential - which is of no interest to us in the present work - may exhibit the typical runaway dilaton potential behaviour [see related comments for the $Z_4$ case $^{15}$].
2.2 Massless spectrum

The closed string sector has N=1 supersymmetry and the massless spectrum of vector and chiral multiplets can be found in [29]; the latter models are non-chiral as the D6-branes are parallel to the O6-planes. The models of the present work are chiral as the D6-branes wrap on general directions and are not parallel to the O6-planes. In the open string sector, we found that the net number of chiral fermions is given by the simple set of rules seen in table (1), where we denote as

\[ \tilde{I}_{ab} = \sum_{\tilde{\Theta}} I_{a(\tilde{\Theta}b)} \]  

the intersection number (IN) from open strings stretching between the D6-brane \( a \) and the D6-branes taking values in the set of orbifold orbit elements

\[ \tilde{\Theta} = \{1, \theta, \theta^2, \omega, \omega^2, \theta \omega, \theta^2 \omega, \theta \omega^2, \theta^2 \omega^2\} \]

Similarly the IN’s coming from open strings stretching between the brane \( a \) and the images of the orbit for brane \( b \) are denoted as

\[ \tilde{I}_{ab'} = \sum_{\tilde{\Theta}} I_{a(\tilde{\Theta}b')} \]  

Figure 1: O6-planes in the orientifold of \( T^6/(Z_3 \times Z_3) \)
| Sector | Multiplicity | Representation |
|--------|--------------|----------------|
| \( a(\tilde{\Theta} a) \) | 3 | \( U(N_a) \) vector multiplet |
| \( a(\tilde{\Theta} b) \) | \( \tilde{I}_{ab} \) | (\( \square \), \( \square \)) fermions |
| \( a(\tilde{\Theta} b') \) | \( \tilde{I}_{ab'} \) | (\( \square \), \( \square \)) fermions |
| \( a(\tilde{\Theta} a') \) | \( \frac{1}{2}(\tilde{I}_{aa'} - \frac{1}{2}I_{a,O6}) \) | \( \square \) fermions |
| | \( \frac{1}{2}(\tilde{I}_{aa'} + \frac{1}{2}I_{a,O6}) \) | \( \square \) fermions |

Table 1: General spectrum on D6-branes at generic angles (namely, not parallel to any O6-plane in all three tori. The models contain additional non-chiral pieces in the \( aa' \), \( ab \), \( ab' \) sectors with zero intersection, if the relevant branes have an overlap.

In particular the intersection numbers between a brane \( a \) and the orientifold images of the D6\(_b\) brane orbits can be expressed in closed from as

\[
\begin{align*}
\tilde{I}_{ab} &= 3(Z_{[a]}Y_{[b]} - Y_{[a]}Z_{[b]}) \\
\tilde{I}_{ab'} &= 3(Z_{[a]}Z_{[b]} - Z_{[a]}Y_{[b]} - Y_{[a]}Z_{[b]}) \\
\#A &= 3(Z_{[a]} - 2Y_{[a]}) \\
\#(A + S) &= \frac{3}{2}(Z_{[a]} - 2Y_{[a]})(Z_{[a]} - 1)
\end{align*}
\]

in terms of the effective wrapping numbers \( Z, Y \) defined in appendices A,B for the AAA; AAB, BBB lattices respectively.

The presence of spectrum rules is independent of the presence of any supersymmetry that might be exhibited by the spectrum. N=1 supersymmetry may be preserved by a system of branes if each stack of D6-branes is related to the O6-planes by a rotation in \( SU(3) \), that is the angles \( \tilde{\theta}_i \) of the D6-branes with respect to the horizontal direction in the \( i \)-th two-torus obey the condition \( \tilde{\theta}_1 + \tilde{\theta}_2 + \tilde{\theta}_3 = 0 \). The supersymmetry of the models that is preserved by any pair of branes is determined by the choice of the orbifold and orientifold action. The models that are presented in this work do not have any supersymmetry preserved by a set of branes. In addition to the above chiral matter arising from bifundamentals there is also \(^5\) non-chiral massless matter present in the adjoint (NCMA). NCMA arises from open strings stretching within a D6\(_a\) brane.

\(^5\)and also non-chiral matter coming from bifundamentals for which we comment on the end of section 4
and between a D6-brane and its orbifold images as

$$ (Adj)_L : \prod_{i=1}^{3} ((m_a^I)^2 + (n_a^I)^2 - m_a^I n_a^I)^2, \ A-lattice \tag{2.18} $$

$$ (Adj)_L : \prod_{i=1}^{3} ((m_a^I)^2 + 3(n_a^I)^2 + 3m_a^I n_a^I)^2, \ B-lattice, \tag{2.19} $$

with the non-zero contributions arriving from the $\theta^2 \omega$, $\theta \omega^2$ images. This sector has N=1 supersymmetry as the $Z_3 \times Z_3$ twist preserves it (see comments in the end of section 4).

3 U(1) anomaly cancellation

In any physical theory the existence of chiral fermions induces gauge anomalies which may be absent for the one loop consistency of the theory. In the context of intersecting branes anomalies may cancel by the use of a generalized Green-Schwarz mechanism that couples the U(1) gauge fields $F_a$ to the untwisted RR fields $B_a$. This mechanism differs from the corresponding mechanism in general type IIB orientifolds where the U(1) anomalies are cancelled through exchange of closed string RR twisted moduli [see also the old work in six- [52] or four-dimensions [53], [54], [55]. As it was pointed out in [19] any U(1) gauge field that has a non-zero $B \wedge F$ coupling necessarily gets massive with a mass of the order of the string scale. In the low energy theory the broken symmetry remains as a global symmetry [For related issues in a general orientifold, with “orthogonal branes”, see e.g. [56], [57]]. In fact, for the present $Z_3 \times Z_3$ orientifolds the anomaly cancellation has to take into account the orbits (2.4) that the D6-branes wrap. A sketch of the anomaly argument cancellation proceed as follows. The cubic non-abelian SU($N_a$) gauge anomaly (GGG) is actually the condition [14]

$$ \sum_a N_b I_{ab} = 0 \tag{3.1} $$

which in the present constructions is proportional to

$$ -2Y_a \left( \sum_{b \neq a} N_b Z_b \right) + \sum_b N_b Z_b Z_b + (N_a - 4)(Z_a - 2Y_a) + 2N_a(Z_a - 2Y_a)(Z_a - 1) \tag{3.2} $$

Eqn. (3.2) vanishes by the use of the RR tadpole condition in the first term in (3.2).

The mixed U(1) anomalies should also cancel. In order to cancel the mixed U(1) gravitational $U(1) - g_{\mu\nu}^2$ anomalies that are proportional to

$$ 3N_a(Z_a - 2Y_a) \tag{3.3} $$
and the mixed $U(1)_a - U(1)_b^2$ anomalies that are proportional to

$$N_a N_b (Z_a - 2 Y_a) Z_b$$  \hfill (3.4)

we need to make use of a generalized Green-Schwarz mechanism that makes use of the mediation of the RR partners of the closed string untwisted geometric moduli.

In order to show that the various $U(1)$-anomalies cancel we will not use the usual picture with the D6-branes intersecting at angles but rather their T-dual picture where the D9-branes have magnetic fluxes along the six dimensional orbifolded tori of type I. In ten dimensions there are two RR fields, $C_2$ and $C_6$ with worldvolume couplings

$$\int_{D_{9a}} C_2 \wedge F_a \wedge F_a \wedge F_a \wedge F_a, \quad \int_{D_{9a}} C_6 \wedge F_a \wedge F_a, \quad I = 1, 2, 3. \hfill (3.5)$$

After dimensional reduction and taking into account the orbifold symmetry, (3.5) reduces to the following Chern-Simons terms in the effective action for the D6-branes

$$- 6 N_a (Z_a - 2 Y_a) \int_{M_4} B_2^0 \wedge F_a, \quad 0 \cdot \int_{M_4} C_0^0 \wedge F_a \wedge F_a;$$

$$- 3 N_a (Z_a - 2 Y_a) \int_{M_4} B_2^I \wedge F_a, \quad - 3 Z_b \int_{M_4} C_0^I \wedge F_b \wedge F_b, \hfill (3.6)$$

where the 2-forms, $dC^0 = -*dB_2^0$, $dC^I = -*dB_2^I$ are defined as

$$B_2^0 = C_2, \quad B_2^I = \int_{(T^2)^I \otimes (T^2)^K} C_6$$  \hfill (3.7)

and their duals as

$$C^I = \int_{(T^2)^I} C_2, \quad C^0 = \int_{(T^2)^I \otimes (T^2)^J \otimes (T^2)^K} C_6$$  \hfill (3.8)

These couplings have exactly the form required to cancel the mixed $U(1)$ and gravitational anomalies \hfill (3.4) and \hfill (3.3).

In general $U(1)$ gauge fields that have a non-zero coupling to RR fields get massive while the associated $U(1)$ survives as a global symmetry to low energies. From the form of the RR couplings to the $U(1)$ gauge fields we derive that the only $U(1)$ that becomes massive is the one given by the expression

$$\sum_a N_a (Z_a - 2 Y_a) F_a$$  \hfill (3.9)

All other $U(1)$’s that may be found - in a model building construction - may survive massless below the string scale and unless some Higgs mechanism \hfill 7 is involved that may give masses to them, they will also survive massless to low energies.

\hfill 6\text{for the AAA torus}

\hfill 7\text{Such Higgs mechanisms have been employed in the construction of deformations of the 4-stack intersecting D6-brane toroidal orientifolds SM of [19] in [20] and [21].}
4 Non-supersymmetric D6-brane models with only the SM at low energy

In this section (and the following ones), we will use wrappings that have at least one zero electric or magnetic entry among them. As a result the models we construct may be non-supersymmetric.

Let us make the choice of wrapping numbers

\[(Z_a, Y_a) = (1, 1), \ (Z_b, Y_b) = (1, 1), \ (Z_c, Y_c) = (-1, 0) \quad (4.1)\]

where our three stacks assume the numbers \(N_a = 3, N_b = 2, N_c = 1\). The RR tadpoles are satisfied and the spectrum can be seen in table (2). The initial gauge group is a \(U(3)_a \times U(2)_b \times U(1)_c\). These models belong to the A’-class and have no-exotics present.

| Matter   | \((SU(3) \times SU(2))(Q_a, Q_b, Q_c)\) | \(U(1)^Y\) |
|----------|----------------------------------------|-------------|
| \(\{Q_L\}\) | 3\((3, \bar{2})_{(-1, -1, 0)}\) | 1/6         |
| \(\{u^c_L\}\) | 3\((3, 1)_{(-2, 0, 0)}\) | -2/3        |
| \(\{d^c_L\}\) | 3\((3, 1)_{(1, 0, -1)}\) | 1/3         |
| \(\{L\}\) | 3\((1, 2)_{(0, 1, -1)}\) | -1/2        |
| \(\{e^+_L\}\) | 3\((1, 1)_{(0, -2, 0)}\) | 1           |
| \(\{N_R\}\) | 3\((1, 1)_{(0, 0, 2)}\) | 0           |

Table 2: A three generation SM chiral open string spectrum (A’-model class). The required scalar Higgses may come from bifundamental N=2 hypermultiplets in the N=2 bc, bc* sectors [19, 20, 21] that may trigger brane recombination.

From the three initial U(1)’s one is anomalous and becomes massive by the GS mechanism by having a non-zero couplings to the RR fields, namely the

\[U(1)^{\text{massive}} = 3F_a + 2F_b + F_c \quad (4.2)\]

In addition there are also two anomaly free U(1)’s that correspond to the hypercharge.
and an extra U(1)

\[
U(1)^Y = \frac{1}{3} F_a - \frac{1}{2} F_b, \quad U(1)^{ex} = \frac{3}{2} F_a + F_b - \frac{13}{2} F_c
\]  

(4.3)

Hence we have found - table (2) - exactly the chiral spectrum of the SM as at this point the spectrum for generic angles is non-supersymmetric. The same non-supersymmetric chiral spectrum construction was found in [15] from intersecting D6-branes in $Z_3$ orientifolds. The Higgs fields that participate in the Yukawa couplings give masses to all the chiral fields but the up quarks.

| Brane | $(n^1, m^1) \times (n^2, m^2) \times (n^3, m^3)$ |
|-------|------------------------------------------|
| \{a\} | $(0, 1) \times (1, 0) \times (0, 1)$       |
| \{b\} | $(0, 1) \times (1, 0) \times (0, 1)$       |
| \{c\} | $(0, -1) \times (0, 1) \times (-1, 0)$     |

Table 3: Wrapping numbers responsible for the generation of the three stack D6-brane non-supersymmetric Standard Models of table (2).

Let us now make a deformation of the previous choice of wrapping numbers,

\[
(Z_a, Y_a) = \left(1, 1\right), \quad (Z_b, Y_b) = \left(1, 1\right), \quad (Z_c, Y_c) = \left(-1, -1\right)
\]  

(4.4)

It satisfies the RR tadpoles and corresponds to the spectrum seen in table (2) but with reversed $U(1)_c$ charges. The exchange of the effective wrappings

\[
(Z, Y)_a \leftrightarrow (Z, Y)_b
\]  

(4.5)

is a symmetry of the spectrum as the spectrum do not change.

- **Gauge couplings**

The gauge coupling constants are controlled by the length of the corresponding cycles that the D6-branes wrap

\[
\frac{1}{\alpha_a} = \frac{M_s}{g_s} ||l_i||
\]  

(4.6)

where $||l_i||$ is the length of the corresponding cycle for the i-th set of brane stacks. The canonically normalized U(1)’s as well the normalization of the abelian generators are
given by $\tilde{U}(1)_a = \frac{F_a}{\sqrt{2} N_a}$, $\text{Tr}(T_a T_b) = \frac{1}{2} \delta_{ab}$. As the hypercharge is given $^8$ as a linear combination $Y(1)^Y = \sum_i c_i F_i$, the value of the weak angle becomes

$$\sin^2 \theta_W = \frac{1}{1 + 4 c_2^2 + 6 c_3^2 (\alpha_2/\alpha_3)}$$

(4.7)

Taking into account that in the present models $\alpha_2 = \alpha_3$ we get the successful GUT relation

$$\sin^2 \theta_W \overset{M_s}{=} \frac{3}{8}$$

(4.8)

which means that the strong and the weak couplings unify at the unification scale, the string scale$^9$. Of course the really important issue here is weather or not partial unification can help us to confirm the experimental measured quantities like $\sin^2 \theta_W$ and $\alpha_{EM}$ at low energies. These issues will be examined elsewhere. A comment is in order. In the models of this work, apart from the Standard model chiral matter we have also present non-chiral bifundamental matter (NCBM) and also non-chiral adjoint matter (NCAM). The former arises in sectors that the D6-branes are parallel in at least one torus, the latter from sectors formed from open strings stretching between the D6-branes and their orbifold images [rules (2.18), (2.19)]. In general NCAM, fermions and scalars, is believed to get massive - receiving radiative corrections - once supersymmetry is broken by massive N=1 supermultiplets running in loops by a mechanism that at present has been shown to be at work only $^19$ at the level of the effective theory. If this mechanism is not at work at the level of string perturbation theory then the presence of extra NCAM - may destroy the asymptotic freedom of the gauge groups and the result (1.8) will be useless. Non-chiral matter (NCM) has only been shown that it gets massive in the context of Scherk-Schwarz deformations (SSD) $^44$ in toroidal orientifolds (TO); where only a subset of wrappings from the ones existing in the usual TO’s gives masses to NCM. It remains to be seen if similar results also hold, once SSD is applied to the current orbifolds. We plan to return to this issue elsewhere.

5 Non-susy models with the fermion spectrum of the N=1 SM and massive non-chiral exotics

In this section we will generate a three stack non-supersymmetric model (B’-model class) that generates the massless fermionic spectrum of the N=1 Standard Model at the string scale. Eventually, all exotics present become massive due to the existence

$^8$we used the conventions used in $^{50}$

$^9$Similar effects has been observed in $^{35}$. 
of appropriate Yukawa couplings. The only other known examples of models in the context of intersecting branes where all exotics become massive are the Pati-Salam $SU(4)_c \times SU(2)_L \times SU(2)_R$ GUTS of \cite{22}. In the latter case, even though the GUT models are non-supersymmetric they do possess N=1 supersymmetric sectors, the latter being responsible for the generation of gauge singlets.

The spectrum of the models can be seen in table (4). This model possess the SM chiral spectrum as well two massless Higgsinos and a pair of exotic triplets at the string scale. The initial gauge group at the string scale is a $U(3)_c \times U(2)_b \times U(1)_c$ which decomposes to an $SU(3)_c \times SU(2) \times U(1)_a \times U(1)_b \times U(1)_c$. Subsequently, via the Yukawas

$$\lambda_C C_1 C_2 \tilde{N}_R + \lambda_H H_u H_d \tilde{N}_R,$$

the Higgsinos and the exotic colour triplets receive Dirac masses through the vev of the tachyonic superpartner of the neutrino, $\tilde{N}_R$. As it is explained in the next section Higgsinos form a Dirac pair with a mass that receives the exponential suppression of the worldsheet area involved in the Yukawa couplings. Its mass scale can be anywhere below the string scale. The effective wrappings have been chosen to be

$$(Z_a, Y_a) = (1, 0), \quad (Z_b, Y_b) = (1, 0), \quad (Z_c, Y_c) = (-1, 0)$$

(5.2)

At the top of the table (4) we see the massless fermion spectrum of the N=1 SM. The corresponding superpartners are part of the massive spectrum and remain ‘hidden’ in the intersection of each corresponding fermion, unless they become tachyonic by varying appropriately the distances between the branes \cite{47} for some relevant examples.

Via the use of the generalized Green-Schwarz mechanism of section 3, the extra beyond the hypercharge U(1)’s become massive leaving only the hypercharge massless to low energies. The model possess three U(1)’s of which one becomes massive, namely $-3F_a + 2F_b - 3F_c$ - via the use of GS mechanism of section 3. The two remaining $U(1)'s$ are the hypercharge

$$U(1)^Y = -\frac{1}{3}F_a + \frac{1}{2}F_b$$

(5.3)

and the $U(1)^{extra} = 3F_a + 2F_b + (13/3)F_c$ which is broken by $\tilde{N}_R$. Thus at low energy only the SM spectrum remains. The Weinberg angle at $M_s$ in these models is also $\sin^2 \theta = 3/8$, as the volumes of the 3-cycles associated to the SU(3) and SU(2) gauge couplings agree as in the D-brane inspired models of \cite{50}. These models belong to the B’-model class differing in respect of the A’-model class as they have in addition of the SM chiral spectrum extra massive non-chiral exotics. In the next section we will
| Matter       | Intersection | $(SU(3) \times SU(2))_{(Q_a,Q_b,Q_c)}$ | $U(1)^Y$ |
|-------------|--------------|--------------------------------------|-----------|
| $\{Q_L\}$  | $ab^*$       | $3(3,2)_{(1,1,0)}$                   | $1/6$     |
| $\{u_L^*\}$| $A_a$        | $3(3,1)_{(2,0,0)}$                   | $-2/3$    |
| $\{d_L^*\}$| $ac^*$       | $3(3,1)_{(-1,0,-1)}$                 | $1/3$     |
| $\{L\}$    | $bc^*$       | $3(1,2)_{(0,-1,-1)}$                 | $-1/2$    |
| $\{H_d\}$  | $bc^*$       | $3(1,2)_{(0,-1,-1)}$                 | $-1/2$    |
| $\{H_u\}$  | $bc$         | $3(1,2)_{(0,1,-1)}$                  | $1/2$     |
| $\{e_L^+\}$| $A_b$        | $3(1,1)_{(0,2,0)}$                   | $1$       |
| $\{N_R\}$  | $S_c$        | $9(1,1)_{(0,0,2)}$                   | $0$       |
| $\{C_1\}$  | $ac$         | $3(3,1)_{(1,0,-1)}$                  | $1/3$     |
| $\{C_2\}$  | $ac^*$       | $3(3,1)_{(-1,0,-1)}$                 | $-1/3$    |

Table 4: A three generation 4D non-supersymmetric model (B'-model class) with the chiral content of N=1 MSSM on top of the table, in addition to $N_R$'s. There are three pairs of $H_u$, $H_d$ Higgsinos. Note that this model mimics models coming from gauge mediation scenarios and possess $\sin^2(\theta) = 3/8$ at $M_s$.

examine an extended B'-model class that possess extra gauge massive fermions and Higgsinos.

6 Standard Models in the presence of massive exotics and Partial Split SUSY scenario

Next we examine the construction of new non-supersymmetric vacua that break to the SM at low energy by using five stacks of intersecting D6-branes.

The original gauge group is a $U(3)_c \times U(2)_w \times U(1)_c \times U(1)_d \times U(1)_e$. The RR tadpoles are satisfied by the choices of effective wrappings

$$(Z_a, Y_a) = \left( 1, \ 1 \right), \ \ (Z_b, Y_b) = \left( 1, \ 1 \right), \ \ (Z_c, Y_c) = \left( -1, \ -1 \right)$$
(Z_d, Y_d) = \left( -1, \ -1 \right), \ (Z_e, Y_e) = \left( 1, \ 1 \right), \quad (6.1)

where the have chosen \( N_a = 3, \ N_b = 2, \ N_c = 1, \ N_d = 1, \ N_e = 1 \). The massless chiral spectrum can be seen in table (5). The generators \( Q_a, \ Q_b, \ Q_c, \ Q_d \) and \( Q_e \) refer to the \( U(1) \) factors within the \( U(3)_c, \ U(2)_w, \ U(1)_c, \ U(1)_d, \ U(1)_e \), respectively. There is one massive \( U(1) \), namely \( -3F_a - 2F_b + F_c + F_d + F_e \). There are also four \( U(1) \)'s that survive massless the GS mechanism, including the hypercharge,

\[
U(1)^{(1)} = F_c - F_d, \quad U(1)^{(2)} = F_c + F_d - 2F_e, \quad U(1)^{(3)} = 18F_a + 12F_b + 26F_c + 26F_d + 26F_e, \quad U(1)^Y = \frac{1}{3}F^a - \frac{1}{2}F^b.
\]

(6.2)

(6.3)

The \( U(1)^{(i)}, \ i = 1, 2, 3 \) generators could be broken by the vevs of the previously massive scalar superpartners of the \( S^{(1)}, \ S^{(2)}, \ S^{(3)} \) fermions respectively that become tachyonic, leaving only the hypercharge massless to low energies. The fermion singlets \( S^{(I)} \) get masses by their couplings to their scalar tachyonic spartners which can get a vev.

The gauge singlet fermions \( S_I \) could be regarded as extra generations of right handed neutrinos. A good choice of right handed neutrinos is to choose the singlets \( S_4 = N_R \). In this case - \( \langle H_u \rangle = v_u \) - the Yukawa couplings even though they come from dimension five operators

\[
\frac{1}{M_s} L N_R \langle H_u \rangle \langle S_2 \rangle \sim v_u L N_R \quad (6.4)
\]

they deliver a tree level Dirac mass term for neutrinos. Alternative choices of right handed neutrinos are unsatisfactory as they result in suppressed masses. Such a typical mass term is as follows. Let us identify the singlet \( S^I \) with \( N^I_R \). The following Yukawa is allowed

\[
e^{-A} (LN^I_R) \tilde{H}_u \tilde{H}_d \tilde{S}_1 \tilde{S}_2 \tilde{S}_3 \tilde{S}_4 \sim e^{-A} \frac{1}{M^2_s} v^2_u v_d (N_s N^I_R) \quad (6.5)
\]

where we have consider that the scalars vevs \( \langle \tilde{S}^I \rangle = M_s, \ \langle \tilde{H}_u \rangle = v_u, \ \langle \tilde{H}_d \rangle = v_d \). By tilde we denote the tachyonic scalar ‘superpartners’ of the \( S_I, \ H_u, \ H_d \) fermion singlets and Higgsinos respectively. The scalars \( \tilde{H}_u, \ \tilde{H}_d \) take part in electroweak symmetry breaking. We also note of the possibility that the neutrino eigenstate is made of a

\[\text{a alternative - but not phenomenologically interesting - possibility might be that the } S^I (S^I) \text{ are not neutrinos. Thus a candidate mass term for } S^I \text{ fermions (similar terms exist for the rest of } S^I \text{ fermions)} \text{ may be}

\[
(S_1)^2 \tilde{S}_3^2 \tilde{S}_3^2 (\tilde{H}_u \tilde{H}_d)^4 \sim e^{-A'} \frac{1}{M^2_s} v^4_u v^4_d (S_1)^2 \equiv e^{-A'} \frac{1}{M^2_s} \tan^2(\hat{\theta}) v^4_d (S_1)^2; \ \tan(\hat{\theta}) = \frac{v_u}{v_d} \quad (6.6)
\]

which is very suppressed.
Table 5: The three generation SM - from a five stack $SU(3)_C \times SU(2)_L \times U(1)_c \times U(1)_d \times U(1)_e$. On the top of the table (a different $B'$-model class) the fermionic spectrum of the N=1 SM with three pairs of Higgsinos and right handed neutrinos. Either one of the gauge singlets $S^I$ could be identified as the one associated with the right handed neutrino. The exotics triplet $X_1, X_2$ and the $H_u, H_d$ fermion pair receive Dirac masses.

| Matter for $Y^1$ | $Y^1$ | $SU(3) \times SU(2)_L(Q_a, Q_b, Q_c, Q_d, Q_e)$ |
|------------------|-------|--------------------------------------------|
| $\{Q_L\}$       | $\frac{1}{6}$ | $3(\bar{3}, 2)(-1, -1, 0, 0, 0)$ |
| $\{u^c_L\}$     | $-\frac{2}{3}$ | $3(\bar{3}, 1)(-2, 0, 0, 0, 0)$ |
| $\{d^c_L\}$     | $\frac{1}{3}$  | $3(3, 1)(1, 0, 0, 1, 0)$ |
| $\{e^+_L\}$     | $1$           | $3(1, 1)(0, -2, 0, 0, 0)$ |
| $\{L\}$         | $-\frac{1}{2}$ | $3(1, 2)(0, 1, 0, 1, 0)$ |
| $\{H_u\}$       | $\frac{1}{2}$  | $3(1, 2)(0, -1, 0, 0, -1)$ |
| $\{H_d\}$       | $-\frac{1}{2}$ | $3(1, 2)(0, 1, 1, 0, 0)$ |
| $\{S_4\}$       | $0$            | $3(1, 1)(0, 0, 1, 1)$ |
| $\{S_1\}$       | $0$            | $3(1, 1, 1)(0, 0, -2, 0, 0)$ |
| $\{S_2\}$       | $0$            | $3(3, 1, 1)(0, 0, 0, -2, 0)$ |
| $\{S_3\}$       | $0$            | $3(1, 1, 1)(0, 0, 1, 0, 1)$ |
| $\{S_5\}$       | $0$            | $3(1, 1, 1)(0, 0, -1, -1, 0)$ |
| $\{X_1\}$       | $\frac{1}{3}$  | $3(3, 1)(1, 0, 1, 0, 0)$ |
| $\{X_2\}$       | $-\frac{1}{3}$ | $3(\bar{3}, 1)(-1, 0, 0, 0, -1)$ |

linear combination of the $S^I$ singlets; in this case also the dominant contribution to their mass comes from $[6.4]$.

- Split Susy scenario for intersecting branes?

Non-supersymmetric models in type I compactifications could solve the gauge hierarchy problem as the string scale could be lowered to the TeV in consistently with gravity as the Planck scale can become large, as long as there are large extra dimensions transverse to the branes [9]. The only known string non-susy models that realize directly the large extra dimension scenario (LEDS) [9] and break to only the SM at low
energy - making use of D5 branes on $T^4 \times C/Z_N$ - have been considered in [23]. In this case there are two transverse dimensions and the string scale could be lowered to the TeV. In the present $Z_3 \times Z_3$ models the D6-branes wrap the whole of the internal space and thus the LEDS scenario cannot be applied. On the other hand the split susy scenario (SSS) [15] has been conjectured as a signal for high energy susy breaking for LHC [58]. In this scenario the candidate model could solve (potentially) the gauge hierarchy problem in a theory that breaks supersymmetry at a high scale, also accompanied by gauge coupling unification and should also predict simultaneously light Higgsinos and gauginos. As it has been emphasized in [47] (and also suggested in [46] in a different context) in models coming from intersecting branes, the SSS scenario should be modified (partial split susy criteria) with respect of the gauginos and Higgsino mass scales. Thus gauginos which are massless at tree level receive loop corrections from massive $N=1$ supermultiplets running in the loops [19] and thus receive string scale masses 11. Also Higgsinos could form Dirac pairs and their mass range could be anywhere from $M_Z$ to the string scale. Hence the present models even though are not split susy models could also provide signatures for LHC (neglecting stability/cosmological constant problems for non-susy/susy models respectively). The modified criteria of the SSS proposal will be also verified using the models with the wrappings (6.1) in this section.

In general in the context of intersecting branes the existence of split susy models could be examined in either a $N=1$ supersymmetric construction [ For some related work see [48] or in constructions that have local supersymmetry as the SM models of [41] that have been generalized with the inclusion of B-field in [42] (In [42] we have also considered the maximal five and six stack SM generalizations of [41]). The latter models even though they are non-supersymmetric they have explicit the presence of $N=1$ susy locally (see also related work on [49]). On the other hand the present models where $N=1$ supersymmetry does not make its appearance even though are not split susy models can still satisfy some of the SSS existence criteria namely like gauge coupling unification (GCU) and still predict light Higgsinos $H_u, H_d$. In the models of table (5) GCU is achieved for only two of the three gauge couplings; the strong and the weak unify at a string scale where the Weinberg angle is $sin^2\theta = 3/8$. Gauginos becomes massive with a mass of order $M_s$ as usual in intersecting brane models. The Yukawa couplings are given by:

\[ \text{Even though the field theoretical study of [19] was exhibited for non-susy models in toroidal orientifolds the result also hold for N=1 susy models coming from orbifolds of orientifolds, as in the N=1 susy case one has to additionally take into account the different orbifold orbits.} \]
• $100 \text{ GeV} < \text{Light Higgsinos/Exotic triplets} < M_s$

$$Y_{\text{table } 4} = \lambda_u L R H_u \tilde{S}_2 / M_s + \lambda_e L e R H_d \tilde{S}_5 / M_s + \lambda_d Q L d R \tilde{H}_d \tilde{S}_5 / M_s + \lambda_{H_u H_d (\tilde{S}_1 \tilde{S}_3) / M_s} + \lambda_{X_1 X_2 (\tilde{S}_1 \tilde{S}_3) / M_s}$$

(6.7)

where the tachyon scalars $H_u, H_d$ play the role of the Higgs fields needed for electroweak symmetry breaking. The Higgsinos $H_u, H_d$ and the triplets $X_i$ form Dirac mass terms in (6.7). If the area involved (the usual suppression factor in intersecting branes) in the relevant Yukawa (in Planck units)

$$\lambda_{X,S} \approx e^{-A}$$

(6.8)

is vanishing and assuming that the scalar singlets $\tilde{S}_1, \tilde{S}_3$ get - their natural scale - vevs of order $M_s$, then the Higgsino/triplets can reach their maximum value of order of the string scale, since $\lambda_{X,S} \approx e^{-A} = 1$ [The Higgsinos get a tree level mass if no massive colour triplets are present in the models of the previous section]. On the other hand different value of the areas involved can make certain that lower mass scales are reached for the Higgsino and the colour triplets $X_i$ Dirac masses. Hence Higgsino Dirac masses of 100, 500 and 1000 GeV are obtained for area values of $\approx 32.9, 31.3, 30.6$ respectively.

7 Three generation non-supersymmetric flipped SU(5) GUTS

The methodology to construct three generation non-supersymmetric flipped SU(5) GUTS that break to the SM at low energy have been exhibited in [43]; where it was shown how one can properly identify the electroweak and Higgs multiplets in the intersecting brane flipped SU(5) and SU(5) context. These models [43] are considered within the $Z_3$ orientifold constructions of [15].

In the present constructions it is also possible to construct flipped SU(5) and SU(5) models which may break to only the SM at low energy. The minimal case that we consider involves the presence of two stacks of D6-branes at the string scale. In this case, we can choose the effective D6-brane wrappings to be $(Z_a, Y_a), (Z_b, Y_b)$, and solve the RR tadpoles (2.13) by making the choices $Z_a = 1, Z_b = -1$, and also

$$(Z_a, Y_a) \equiv (1, Y_a), \quad (Z_b, Y_b) \equiv (-1, Y_b),$$

(7.1)

that corresponds to a two stack model with an initial gauge group $U(5)_a \times U(1)_b$. Since the nature of the intersection numbers gives chiral fermions on intersections that are
multiples of three, in the general case we can assume that the number of generations is 3n and then solve for the 3G case, \( n = 1 \). Having 3n generations enforces us to choose 3n copies of \( \bar{10} \) representations. Hence an appropriate choice to generate flipped SU(5) models will be

\[
I_{51} = 3n, \quad I_{51^*} = 0, \quad \#(A)_a = -3n, \quad \#(\bar{10})_b = 3n, \quad (7.2)
\]

Explicitly the intersection numbers are given by

\[
I_{51} \#(5,1) = 3(Y_b + Y_a), \quad (A)_a = 3(1 - 2Y_a), \quad (S)_b = 3(2Y_b - 1) \quad (7.3)
\]

with the solution of (7.2), (7.3) to be

\[
(Y_a, Y_b) = \left( \frac{n + 1}{2}, \frac{n - 1}{2} \right) \quad (7.4)
\]

Choosing the model to have 3 generations, we recover the spectrum of table (6). The

| Sector | Multiplicity | Repr   | \( Q_a \) | \( Q_b \) | \( Q^{fl} \) |
|--------|--------------|--------|-----------|-----------|-------------|
| \{51\} | 3            | (5,1)  | 1         | -1        | 3           |
| (A)\_b| 3            | (\bar{10},1) | -2       | 0         | -1          |
| S      | 3            | (1,1)  | 0         | 2         | 5           |

Table 6: Three generation flipped SU(5) GUT models. The last column indicates the flipped U(1) charge.

application of the Green-Schwarz mechanism of section 3, suggests that only the U(1) gauge boson

\[
U(1)^{fl} = \frac{1}{2}U(1)_5 - \frac{5}{2}U(1)_b \quad (7.5)
\]

survives massless the Green-Schwarz mechanism by not having a non-zero coupling to the RR fields. It exactly corresponds to the U(1) generator of the flipped SU(5) x U(1)^{fl}. The generation content is as usual

\[
F = 10_1 = (u, d, d^c, \nu^c), \quad f = \bar{5}_3 = (u^c, \nu, e), \quad l^c = 1_5 = e^c \quad (7.6)
\]

A consistent set of wrappings for the three generation flipped SU(5) model is given in table (7) where \((Z_5, Y_5) = (1, 1)\) and \((Z_1, Y_1) = (-1, 0)\). The SU(3) x SU(2) x U(1)\(Y\) models of section (4) can be reinterpreted as coming from adjoint breaking of the U(5) x U(1) models of this section. During this process the adjoint 24 gets a non-vanishing vev, and thus the U(5) stack splits into two stacks accommodating the U(3) and U(2) gauge groups. The reverse process corresponds to moving the U(3) and U(2) factors on top of each other by tuning the adjoint 24 to a vanishing vev.
Table 7: D6-brane wrapping numbers for the three family flipped SU(5)(and SU(5)) GUT.

| Brane/Gaugegroup | $N_a$ | $a_1^1,m_1^1)(a_2^2,m_2^2)(a_3^3,m_3^3)$ |
|------------------|-------|----------------------------------------|
| U(5)             | 5     | (0, 1)(0, 1)(1, 1)                     |
| U(1)             | 1     | (1, 0)(0, 1)(0, 1)                     |

**7.1 Higgs sector in the flipped SU(5) GUTS**

The flipped SU(5) GUT symmetry breaks to the SM one by the use of the tachyonic Higgs excitations seen in table 8. Electroweak Higgses may come also from open strings stretching between the orbits of branes 5 and 1*. Their quantum numbers may be seen in Table 9.

Table 8: GUT breaking Higgses for flipped SU(5) classes of GUT models.

| Sector | Field | Repr | $Q_a$ | $Q_b$ | $Q_{fl}$ |
|--------|-------|------|-------|-------|---------|
| $(A)_{1/3}$ | $H_1$ | 10   | 2     | 0     | 1       |
| $(A)_{1/3}$ | $H_2$ | $ar{10}$ | −2   | 0     | −2      |

Table 9: Electroweak Higgses for flipped SU(5) classes of GUT models.

**7.2 Proton decay and Mass generation in the flipped SU(5) GUTS**

- **Proton decay**

  The study of proton decay amplitude, with reference to an SU(5) GUT as those coming from $Z_2 \times Z_2$ orientifolds [16] has been performed to [18]. Further studies of the disk amplitudes that contribute to the gauge mediated proton decay (GMPD) in a general flipped SU(5) GUT and also in general SU(5) GUTS - with emphasis on the orientifolds of $Z_3$ orbifolds has been performed in [13]. As in the present $Z_3 \times Z_3$ orientifold flipped SU(5) GUTS baryon number is not a gauged symmetry GMPD operators do contribute to proton decay - these contributions are identical to those
appearing in [43] - and thus the string scale should be higher than $10^{16}$ GeV in order to safeguard the stability of the proton. The string scale in the present constructions is naturally high as $12^{12}$ the D6-branes wrap the whole of the internal space.

- Quark, lepton masses

\[ Y = \lambda^{\text{quark}} u \cdot \bar{f} \cdot \langle h_2 \rangle + \lambda^{\text{lepton}} c \cdot \bar{l} \cdot \langle h_1 \rangle \]  

(7.7)

- Neutrino masses

\[ \lambda^{(1)} F \cdot \bar{f} \cdot \langle h_2 \rangle + \lambda^{(2)} (F \cdot H_2)(F \cdot H_2)/M_s \]  

(7.8)

where in the first line there are mass terms for the up-quarks and charge lepton masses; in the second line a see-saw mass matrix for the neutrinos. There are no tree mass terms for the down quarks but this is not a severe problem as the magnitudes of their masses is small and it is possible that they could be generated by higher non-renormalizable terms.

8 SU(5) GUT generation

SU(5) models may be produced by using the flipped SU(5) construction of the previous sections. We choose to break the massless U(1) - that survives the GS mechanism - of flipped SU(5) by turning on a singlet tachyon scalar field coming from the open strings stretching between the branes that support the orbit of the $U(1)_b$ brane. Then our model becomes an SU(5) class of GUTS. The breaking to the SM in this case is achieved by the use of the adjoint 24, part of the N=1 Yang-Mills multiplet in the aa-sector that utilizes itself by splitting the U(5) stack into two stacks of U(3) and U(2) branes - with identical wrappings - away from each other. Such a process have been described in [47]. Further details on the identification of GUT and electroweak Higgses can be found in [43]. We note that the Weinberg angle in these SU(5) GUTS also receives the value $\sin^2\theta = 3/8$. This can be proved along the same lines as the ones used in [15] as in the present discussion we have reproduced the spectra of SU(5) GUTS of [15] in the context of $Z_3 \times Z_3$ orientifolds using equal volume cycles for the relevant gauge couplings.

\[^{12}\text{there are no dimensions transverse to the D6-branes that could be made large and lower the string scale}\]
9 Pati-Salam models

A non-supersymmetric Pati-Salam three family model could be constructed [the chiral spectrum may be seen in table 10] using three stacks of branes and the choice

\[(Z_a, Y_a) \equiv (1, 0), (Z_b, Y_b) \equiv (1, 0), (Z_c, Y_c) \equiv (-1, -1)\]

(9.1)
giving a gauge group \(U(4)_c \times U(2)_L \times U(2)_R\). The chiral content of these PS models (e.g. with surplus 3-, 6-plet exotics) is similar to the observable group chiral content of the PS models of tables 3, 4 in the 1st ref. of [17].

| Matter | SU(4)_c × SU(2)_L × SU(2)_R | Q_c | Q_L | Q_R |
|--------|-------------------------------|-----|-----|-----|
| \{F_L\} | 3(4, 2, 1) | 1   | 1   | 0   |
| \{F_R\} | 3(4, 1, 2) | -1  | 0   | 1   |
| \{ω_L\} | 3(6, 1, 1) | 2   | 0   | 0   |
| \{χ_L\} | 3(1, 2, 2) | 0   | -1  | 1   |
| \{ψ_L\} | 3(1, 1, 3) | 0   | 0   | -2  |
| \{P_0\} | 3(1, 1, 1) | 0   | 0   | 2   |
| \{P_1\} | 3(1, 1, 1) | 0   | 0   | -2  |
| \{P_2\} | 3(1, 1, 1) | 0   | 2   | 0   |

Table 10: Chiral spectrum for a three generation PS-model. The U(1) charges belong to the respecting U(n) gauge groups.

The PS models of table (10) can be further subjected to gauge symmetry breaking by adjoint splitting of the D6-branes. Hence the gauge symmetry can be broken directly to the SM by splitting the \(U(4)_c\) stack - into parallel but not overlapping stacks, namely \(a_1\) and \(a_2\), made from 3 and 1 branes - and also by splitting the \(U(2)_R\) stack into two stacks, namely \(c\), \(d\), made from parallel 1 branes. Moving away the D6-branes in the \(U(4)_c\), \(U(2)_R\) stacks corresponds to giving vevs to the appropriate scalars in the adjoints of \(SU(4)_c\), \(U(2)_R\). The application of the Green-Schwarz mechanism makes massive only the \(U(1)^{\text{mass}} = 3F_{a_1} + 2F_b + F_{a_2} + F_c + F_d\). From the rest of the U(1)’s, one is the \(U(1)^Y = (-1/3)F_{a_1} + (1/2)F_b\) SM hypercharge which remains massless while the rest three massless U(1)’s namely, \(U(1)^{(1)} = 2F_{a_2} - F_c - F_d\), \(U(1)^{(2)} = F_c - F_d\),

\[13\] The adjoint breaking is further explained in [16, 17].
$U(1)^{(3)} = -18F_a - 12F_b + 13(F_a + F_c + F_d)$ receive masses if the tachyonic superpartners of the singlets $S_1$, $S_3$, $S_4$ receive a vev respectively. Within this procedure the Pati-Salam gauge group may be broken to the SM. The resulting spectrum is that of table (11). The rest of chiral fermions; namely the colour triplets $X_i$ could receive a Dirac mass term of order $M_s$ by the coupling $^{14}X_1X_2\langle S_1^B \rangle \langle S_3^B \rangle /M_s$, where $S_1^B$, $S_3^B$ tachyon superpartners of the $S_1$, $S_3$ fermions. The singlet fermions $S^I$ also could receive a mass, e.g. the $S^1$ obtain a term $(S_1^1)^2\langle \Psi_1^B \rangle^2 \langle S_3^B \rangle^2$ of order $M_s$, where $S_3^B$, $\Psi_1^B$ the tachyonic superpartners of the $S_3$ fermion and the ones coming from the $a_2c^*$ intersection respectively. Thus at low energies only the SM fermion content remains (with no mass terms for the up-quarks).

10 Conclusions

In this work, we have described the construction of $Z_3 \times Z_3$ orientifolds where the D6-branes intersect at angles and are not parallel with the O6-planes. The presence of O6-planes requires for the cancellation of anomalies in the four dimensional compactification of IIA, intersecting D6-branes that wrap the internal six dimensional toroidal space and also a generalized Green-Schwarz mechanism to cancel the mixed U(1) anomalies. In this context, we have described in full generality the localization of chiral matter that gets consistently localized between the D6-branes. The presence of chiral matter is independent of the presence or not of N=1 supersymmetry in the open string spectrum of the theory.

We focused on the construction of models which break to the SM at low energy, which have equal SU(3), SU(2) gauge couplings so that the Weinberg angle at $M_s$ is equal to $3/8$. We have also described the possibility of constructing a) GUT models and also b) new SM vacua (section 6) with the spectrum of the SM in addition to

$^{14}$Models with identical fermion content as in table (11) which are constructed by direct methods and not by adjoint breaking - where the string scale mass terms for the exotic colour triplets arises at tree level - have been studied in [47]. See model A of the latter work.

$^{15}$singlet fermions in the $a_2c^*$ intersection are non-chiral; in fact $I_{a_2c^*} = 1 - 1 = 0$, the non-zero contributions - where the D6-branes are non-parallel in all tori - are coming from the orbits $\Omega R \omega$, $\Omega R \theta^2 \omega^2$ and we have used the wrapping numbers $(1,0)(0,1)(0,-1)$ and $(1,1)(-1,0)(-1,-1)$ for the $a_2$ and $c$-branes respectively. The non-chiral singlet fermions $(\Psi_1)(\Psi_1)(S_1^B)^2(S_3^B)^2$ and $(\Psi_1)(\Psi_1)(\Psi_2^B)^2$ while $(\Psi_2)$ could get an $M_s$ mass from the coupling $(\Psi_2)(\Psi_2)(\Psi_2^B)^2$.

$^{16}$Problems related to the successful prediction through RG running of low energy phenomenological quantities like $\sin^2(M_Z)$, $a_{EM}$ may be considered elsewhere.
Table 11: The three generation SM from the adjoint splitting of the PS model of table (10).

| Matter | $SU(3) \times SU(2)_L$ | $Q_{e_1}$ | $Q_{e_2}$ | $Q_b$ | $Q_c$ | $Q_d$ | $Y$ |
|--------|-----------------------|----------|----------|-------|-------|-------|-----|
| $\{Q_L\}$ | $(3,2)$ | 1 | 0 | 1 | 0 | 0 | $\frac{1}{6}$ |
| $\{w_L^c\}$ | $(3,1)$ | 2 | 0 | 0 | 0 | 0 | $-\frac{2}{3}$ |
| $\{d_L^c\}$ | $(3,1)$ | -1 | 0 | 0 | 1 | 0 | $\frac{1}{3}$ |
| $\{e_L^+\}$ | $(1,1)$ | 0 | 0 | 2 | 0 | 0 | 1 |
| $\{L\}$ | $(1,2)$ | 0 | 0 | $-1$ | 0 | 1 | $-\frac{1}{2}$ |
| $\{H_u\}$ | $(1,2)$ | 0 | 1 | 1 | 0 | 0 | $\frac{1}{2}$ |
| $\{H_d\}$ | $(3,1)$ | 0 | 0 | $-1$ | 1 | 0 | $-\frac{1}{2}$ |
| $\{S_1\}$ | $(1,1)$ | 0 | $-1$ | 0 | 1 | 0 | 0 |
| $\{S_2\}$ | $(3,1)$ | 0 | $-1$ | 0 | 0 | 1 | 0 |
| $\{S_3\}$ | $(1,1)$ | 0 | 0 | 0 | $-2$ | 0 | 0 |
| $\{S_4\}$ | $(1,1)$ | 0 | 0 | 0 | $-2$ | 0 | 0 |
| $\{S_5\}$ | $(1,1)$ | 0 | 0 | 0 | $-1$ | $-1$ | 0 |
| $\{X_1\}$ | $(3,1)$ | 1 | 1 | 0 | 0 | 0 | $-\frac{1}{4}$ |
| $\{X_2\}$ | $(3,1)$ | -1 | 0 | 0 | 0 | 1 | $\frac{1}{4}$ |

(3) pairs of Higgsinos and exotic triplets (the latter pairs becoming massive by tree level Yukawa couplings) and c) the construction of different (wrapping) solutions to 3- (A’ class) and 5-stack (B’ class) non-supersymmetric models with only the SM at low energy than the ones extensively studied in [47]. At 3-stacks we get only the SM without any chiral exotics present. These 3-stack models exactly reproduce the SM example given in [15] in the context of $Z_3$ orientifolds. At 5-stacks we get chiral models with the chiral fermionic spectrum of the N=1 supersymmetric SM in addition to three identical pairs of massive non-chiral exotics. Even though the models are non-supersymmetric they predict Higgsinos with a light mass that could provide us with a signal for LHC. The construction of GUT models is also possible; in the last section where we have detailed the construction of two stack flipped SU(5) (also SU(5)) models which can break to the SM at low energy and also three stack Pati-Salam GUTS which after adjoint breaking break to the B’-model class.

\[\text{as we have already said in section 4 we neglected through this work the issue of massless non-chiral matter}\]
For the non-supersymmetric SM’s, SU(5) and Pati-Salam $SU(4) \times SU(2)_L \times SU(2)_R$ GUTS - of the present work there is an absence of tree level mass term for the up-quarks. This is a general problem in models involving antisymmetric representations [The same problem persists also in the construction of N=1 SU(5) models from $Z_2 \times Z_2$ orientifolds and in the N=0 models of [15]]. Due to the largeness of the top quark mass, it will be impossible to generate a mass from higher order non-renormalized corrections. On the contrary in flipped SU(5) GUTS, the opposite happens as there are no mass terms for the d-quarks. As the d-masses are generally small, the absence of a tree level mass is not very problematic, as higher order terms may generate in principle the required masses.

As the models we constructed are non-supersymmetric and there are no transverse dimensions that can dilute the strength of gravity, the string/GUT scale may be high [9]. Hence they can safeguard the models from proton decay since baryon number is not a gauged symmetry. However, the high scale is unsatisfactory since we want also to avoid the gauge hierarchy problem (GHP) in the Higgs sector.

In all the models that we have constructed in this work and well as in the preceding [47] one, we have constructed D-brane models using in all cases wrappings $(n^I, m^I)$ which have at least one zero entry for the AAA vacua. Since the models derived in all cases were non-supersymmetric to construct N=1 supersymmetric models we might have to use wrappings with all entries being non-zero for the AAA tori 18. Such an attempt requires an extensive computer search and will be the subject of future work. We also note that we have not investigated the construction of N=1 or N=0 models using the AAB, BBB tori for which we have presented the RR tadpoles, the spectrum rules and the Green-Schwarz anomaly cancellation mechanism in appendix B. We leave this task for future work.

The present constructions even though the orbifold symmetry stabilizes naturally the complex structure moduli do not fix Kähler moduli that are introduced by the orbifold in its twisted sectors. We note that more moduli could be fixed in models with RR and NSNS fluxes e.g. [60], [61] but in this case the exact string description is lost since the results are valid in the low energy supergravity approximation. Alternatively, one could use oblique magnetic fluxes [62], [63] which can fix in principle all moduli but where no gauge group factors of rank large enough to accommodate the SM arise at the moment. On the contrary in the present work we have chosen to build realistic gauge groups, leaving aside problems related to stabilization, fixing moduli and...
gauge hierarchy that finally render our models semirealistic. Some of these moduli in semirealistic models from orbifolded orientifolds could be fixed in principle by the use of multiple gaugino condensates [see for example [16] for examples of this method in N=1 semirealistic models]. We also note that the construction of N=1 supersymmetric models - that have a high scale - may alleviate the hierarchy problem that the present models possess. It will also be interesting to see what will be the effect of discrete torsion in the present models [59].

As the models we have constructed are non-supersymmetric only the NSNS dilaton tadpole remains uncancelled. The presence of the dilaton tadpole does not signal an inconsistency of the theory, but rather that the background is not a solution of the classical equations of motion[see some recent work [64]] and thus may be corrected [65, 66]. In fact it plays the role of an uncancelled effective cosmological constant and is rather connected with the problem of breaking supersymmetry. As even in N=1 supersymmetric models, where no NSNS tadpoles are present initially and after the breaking of supersymmetry we do get a cosmological constant of the wrong size, we will choose to ignore the NSNS issue for the time being as it is connected with whatever mechanism will solve the cosmological constant problem in particle physics.

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11 Appendix A

In this appendix we list the analytic form of the effective wrapping numbers, $Z_{[a]}$, $Y_{[a]}$ in terms of the ‘electric’ and ‘magnetic’ wrapping numbers $(n^i_a, m^i_a)$.

- **AAA torus**

\[
Z^{AAA}_{[a]} = (2m^1_a m^2_a m^3_a + 2n^1_a n^2_a n^3_a - n^1_a m^2_a m^3_a - m^1_a m^2_a n^3_a - n^1_a n^2_a m^3_a - n^1_a m^2_a n^3_a - m^1_a n^2_a m^3_a - n^1_a n^2_a n^3_a, \]
\[
Y^{AAA}_{[a]} = (n^1_a n^2_a n^3_a + m^1_a m^2_a m^3_a - n^1_a n^2_a m^3_a - m^1_a m^2_a n^3_a - n^1_a n^2_a n^3_a - m^1_a n^2_a n^3_a) \quad (11.1)
\]

We can introduce an abbreviation that can help us simplify and shorten the lengthy expressions. Alternatively we can write (11.1) as

\[
Z^{AAA}_{[a]} = (2mmm + 2nnn - nmm - nnn)_a, \]
\[
Y^{AAA}_{[a]} = (nnn + mmm - nnn)_a \quad (11.2)
\]
where each triplet of letters correspond to - reading from left to right - to the wrappings of the first, second and third tori respectively. A single letter inside a triplet of wrappings that may be different than the rest of them is always permuted among the different tori.

12 Appendix B

In order to derive the spectrum, RR tadpoles and apply the Green-Schwarz anomaly cancellation mechanism in those classes of $Z_3 \times Z_3$ four dimensional orientifold string compactifications that the internal space is made of BBB, AAB tori we may use a different notation for the A lattice than the one used in the main body of the paper. The notation for the A-, B- lattices we will follow - in this appendix - may be as in [15]. In particular we will go to a point in moduli space where the complex structure in all three $T^2$'s will be fixed to the values $U_A = \frac{1}{2} + i \frac{\sqrt{3}}{2}$, $U_B = \frac{1}{2} + i \frac{1}{2\sqrt{3}}$.

Hence under the $Z_3$ and $\Omega R$ symmetries the brane orbits are given for the A-torus by

$$
\begin{align*}
\begin{pmatrix}
\tilde{n}^i \\
\tilde{m}^i
\end{pmatrix} &\xrightarrow{Z_3} \begin{pmatrix}
-n^i - m^i \\
n^i
\end{pmatrix} \xrightarrow{\Omega R} \begin{pmatrix}
m^i \\
-n^i - m^i
\end{pmatrix} \\
\begin{pmatrix}
n^i + m^i \\
-m^i
\end{pmatrix} &\xrightarrow{Z_3} \begin{pmatrix}
-m^i \\
n^i
\end{pmatrix} \xrightarrow{\Omega R} \begin{pmatrix}
-n^i \\
n^i + m^i
\end{pmatrix}
\end{align*}
$$

(12.1)

and for the B-torus by

$$
\begin{align*}
\begin{pmatrix}
n^i \\
m^i
\end{pmatrix} &\xrightarrow{Z_3} \begin{pmatrix}
-2n^i - m^i \\
3n^i + m^i
\end{pmatrix} \xrightarrow{\Omega R} \begin{pmatrix}
n^i + m^i \\
-3n^i - 2m^i
\end{pmatrix} \\
\begin{pmatrix}
n^i + m^i \\
-m^i
\end{pmatrix} &\xrightarrow{Z_3} \begin{pmatrix}
n^i \\
-3n^i - m^i
\end{pmatrix} \xrightarrow{\Omega R} \begin{pmatrix}
-2n^i - m^i \\
3n^i + 2m^i
\end{pmatrix}
\end{align*}
$$

(12.2)

• AAB torus

The massless spectrum is given by the rules of eqn’s (2.17). The RR tadpoles are given by

$$
\sum_a N_a Z_{[a]}^{AAB} = 4 ,
$$

(12.3)
where

\[
Z^{AAB}_{[a]} = 2n_a^1 n_a^2 n_a^3 + m_a^1 m_a^2 n_a^3 + m_a^1 n_a^2 n_a^3 + m_a^1 n_a^2 n_a^3 - m_a^1 m_a^2 m_a^3 \\
= \left(2n_a^1 n_a^2 n_a^3 - m_a^1 m_a^2 n_a^3 - m_a^1 m_a^2 m_a^3 + \frac{m_{ab} \cdot m_{bc}}{m_a^1 m_a^2 m_a^3}\right)_a
\]  

(12.4)

and

\[
Y^{AAB}_{[a]} = -m_a^1 m_a^2 m_a^3 - n_a^1 m_a^2 m_a^3 - m_a^1 n_a^2 m_a^3 - 2m_a^1 m_a^2 n_a^3 - n_a^1 m_a^2 n_a^3 - m_a^1 n_a^2 m_a^3 + n_a^1 n_a^2 n_a^3 \\
\]  

(12.5)

- **BBB torus**

The massless spectrum is given by the rules of eqn’s (2.17). The RR tadpoles are given by

\[
\sum_a N_a Z^{BBB}_{[a]} = 12 ,
\]

(12.6)

where

\[
Z^{BBB}_{[a]} = 6n_a^1 n_a^2 n_a^3 + 3m_a^1 m_a^2 n_a^3 + 3n_a^1 m_a^2 m_a^3 + 3n_a^1 n_a^2 m_a^3 + m_a^1 m_a^2 m_a^3 + m_a^1 n_a^2 m_a^3 + n_a^1 m_a^2 m_a^3 \\
= \left(6n_a^1 n_a^2 n_a^3 + 3m_{ab} \cdot m_{bc} + m_{ab} \cdot m_{bc}\right)_a
\]

(12.7)

and

\[
Y^{BBB}_{[a]} = -m_a^1 m_a^2 m_a^3 + 3n_a^1 n_a^2 n_a^3 - m_a^1 m_a^2 m_a^3 - m_a^1 n_a^2 m_a^3 - n_a^1 m_a^2 m_a^3 \\
= \left(-m_a^1 m_a^2 m_a^3 + 3n_a^1 n_a^2 n_a^3 - m_{ab} \cdot m_{bc}\right)_a
\]

(12.8)

- **U(1) anomaly cancellation**

The Green-Schwarz mechanism - for the BBB, AAB tori - makes massive the U(1) that has its couplings given by (3.9). Thus for the BBB vacua the examination of the BF couplings reveals e.g.

\[
-18(Z_a - 2Y_a) \int_{M_4} B_2^a \wedge F_a, \quad -6Z_b \int_{M_4} C^a \wedge F_b \wedge F_b;
\]

(12.9)

that these couplings have the right form to cancel the mixed-U(1) gauge anomalies via the generalized Green-Schwarz mechanism
References

[1] P. Candelas, J. A. Harvey, C. Vafa and E. Witten, Nucl. Phys. B258 (1985) 46
[2] L. Dixon, J.A. Harvey, C. Vafa, E. Witten, Nucl. Phys. B 261 (1985) 678
[3] W. Lerche, D. Lüst and A.N. Schellekens, Nucl. Phys. B287 (1987) 477
[4] D. Gepner, Nucl. Phys. B296 (1988) 757
[5] H. Kawai, D. Lewellen and S.H. Tye, Nucl. Phys. B288 (1987) 1; I. Antoniadis, C. Bachas and C. Kounnas, Nucl. Phys. B289 (1987) 87
[6] L.E. Ibanez, J. Mas, H.P. Nilles, F. Quevedo, Nucl. Phys. B301 (1988) 157
[7] E. Dudas, Clas. Quan. Grav.17 (2000) R41, arXiv:hep-ph/0006190
C. Angelantonj and A. Sagnotti, Phys. Rep. 371 (2002) 1, [Erratum-ibid. 376 (2003) 339], arXiv:hep-th/0204089

[8] C. Kokorelis, “Standard Model Building from Intersecting D-branes”, to appear in the series “New developments in string theory research”, Nova Science Publishers, NY, ISBN:1-59454-488-3, arXiv:hep-th/0410134.
D. Lüst, Class. Quant. Grav. 21:S1399-1424, 2004, arXiv:hep-th/0401156;
A. M. Uranga, Class. Quant. Grav. 20 (2003) S373-S394, [arXiv: hep-th/0301032]

[9] N. Arkani-Hamed, S. Dimopoulos and G.R. Dvali, Phys. Rev. D59:086004, 1999, hep-ph/9807344
I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G.R. Dvali, Phys. Lett. B436 (1998) 257, hep-ph/9804398
[10] C. Angelantonj, I. Antoniadis, E. Dudas and A. Sagnotti, Phys. Lett. B489 (2000) 223, hep-th/0007090
C. Angelantonj and A. Sagnotti, arXiv:hep-th/0010279

[11] C. Bachas, “A way to break supersymmetry”, hep-th/9503030
[12] M. Berkooz, M. R. Douglas and R. G. Leigh, Nucl. Phys. B 480, 265 (1996), arXiv:hep-th/9606139
[13] R. Blumenhagen, B. Körs and D. Lüst, JHEP 0102 (2001) 030, arXiv:hep-th/0012156

[14] G. Aldazabal, S. Franco, L. E. Ibanez, R. Rabadan, A. M. Uranga, J. Math. Phys. 42 (2001) 3103, hep-th/0011073
JHEP 0102 (2001) 047, hep-ph/0011132
[15] R. Blumenhagen, B. Körs, D. Lüst and T. Ott, Nucl. Phys. B616 (2001) 3, arXiv:hep-th/0107138

[16] M. Cvetic, G. Shiu, A. M. Uranga, Nucl. Phys. B615 (2001) 3, arXiv:hep-th/0107166; M. Cvetic, G. Shiu, A. M. Uranga, Phys. Rev. Lett. 87 (2001) 201801, arXiv:hep-th/0107143; M. Cvetic, P. Langacker, G. Shiu, Nucl. Phys. B 642 (2002) 139, arXiv:hep-th/0206115

[17] M. Cvetic, T. Li, T. Liu, Nucl.Phys. B698 (2004) 163, arXiv:hep-th/0403061; M. Cvetic, P. Langacker, T. Li, T. Liu, Nucl.Phys. B709 (2005) 241, arXiv:hep-th/0407178

[18] I. Klebanov and E. Witten, Nucl. Phys. B664 (2003) 3, arXiv:hep-th/0304079

[19] L. E. Ibáñez, F. Marchesano and R. Rabadán, JHEP, 0111 (2001) 002, arXiv:hep-th/0105155

[20] C. Kokorelis, JHEP 09 (2002) 029, arXiv:hep-th/0205147

[21] C. Kokorelis, JHEP 08 (2002) 036, arXiv:hep-th/0206108

[22] Non-susy Pati-Salam GUTS from intersecting branes have been examined in : C.Kokorelis, JHEP 08 (2002) 018, arXiv:hep-th/0203187; arXiv:hep-th/0210004; arXiv:hep-th/0211091; arXiv:hep-th/0212281

[23] C.Kokorelis, JHEP 0211 (2002) 027, arXiv:hep-th/0209202; “Deformed Intersecting D6-Branes II”, arXiv:hep-th/0210200

[24] D. Cremades, L. E. Ibáñez and F. Marchesano, Nucl. Phys. B643 (2002) 93, arXiv:hep-th/0205074; C. Kokorelis, Phys. B677:115, 2004, arXiv:hep-th/0210234

[25] D. Cremades, L. E. Ibáñez, F. Marchesano, JHEP 0207 (2002) 022, arXiv:hep-th/0203160; JHEP 0207 (2002) 009, arXiv:hep-th/0201205;

[26] M. Larosa and G. Pradisi, Nucl. Phys. B667:261, 2003, arXiv:hep-th/0305224

[27] R. Blumenhagen, L. Görlich, T. Ott, JHEP 0301 (2003) 021, arXiv:hep-th/0211059; R. Blumenhagen, J. P. Conlon, K. Suruliz, JHEP 0407:022,2004, arXiv:hep-th/0404254

[28] R. Blumenhagen, L. Görlich, B. Körs, JHEP 0001 (2000) 040, hep-th/9912204
[29] S. Forste, G. Honecker, R. Schreyer, Nucl. Phys. 593 (2001) 127, arXiv:hep-th/0008250

[30] M. Bianchi, J.F. Morales and G. Pradisi, Nucl. Phys. B573 (2000) 314, arXiv:hep-th/9910228

[31] G. Honecker, arXiv:hep-th/0303015 JHEP 0201 (2002) 025, arXiv:hep-th/0201037

[32] G. Honecker and T. Ott, Phys.Rev. D70 (2004) 126010; Erratum-ibid. D71 (2005) 069902, arXiv:hep-th/0404055

[33] R. Blumenhagen and T. Weigand, JHEP 0402 041 (2004), arXiv:hep-th/0401148
R. Blumenhagen, JHEP 09 (2003) 067, arXiv:hep-th/0310244

[34] T.P.T. Dijkstra, L. R. Huiszoon and A. N. Schellekens, arXiv:hep-th/0403196
arXiv:hep-th/0411129

[35] R. Blumenhagen, D. Lust and S. Stieberger, JHEP 0307 (2003) 036, arXiv:hep-th/0305146

[36] D. Lüst and S. Stieberger, arXiv:hep-th/0302221

[37] M. Cvetic, I. Papadimitriou, Phys. Rev. D68 046001 (2003)[Erratum-ibid. D79, 029903 (2004) arXiv:hep-th/0303083];

[38] M. Cvetic, I. Papadimitriou, G. Shiu, Nucl. Phys. B659 (2003) 193, arXiv:hep-th/0212177;

[39] S. A. Abel, A. W. Owen, arXiv:hep-th/0303124; S. A. Abel, A. W. Owen, arXiv:hep-th/0310257 S. A. Abel, M. Masip, J. Santiago, JHEP 0304 (2003) 057, arXiv:hep-ph/0303087; S. Abel, O. Lebedev, J. Santiago, arXiv:hep-ph/0312157;

[40] F. Epple, JHEP 0409 (2004) 021, arXiv:hep-th/0408105

[41] D. Cremades, L. E. Ibáñez, F. Marchesano, arXiv:hep-th/0302105

[42] C. Kokorelis, “N=1 Locally Supersymmetric Standard Models from Intersecting Branes”, arXiv:hep-th/0309070
[43] M. Axenides, E. Floratos and C. Kokorelis, JHEP 0310 (2003) 006, arXiv:hep-th/0307255;
See also the reviews: C. Kokorelis, Proc. SUSY 2003, arXiv:hep-th/0402087

[44] C. Angelantonj, M. Cardella, N. Irges, arXiv:hep-th/0503179

[45] N. Arkani-Hamed and S. Dimopoulos, “Supersymmetric Unification without Low Energy Supersymmetry and Signatures for Fine-Tuning at the LHC”, arXiv:hep-th/0405159

[46] I. Antoniadis and S. Dimopoulos, “Split supersymmetry in string theory”, arXiv:hep-th/0411032;

[47] C. Kokorelis, “Standard Models and Split Supersymmetry from Intersecting Brane Orbifolds”, arXiv:hep-th/0406258

[48] B. Kors and P. Nath, Nucl. Phys. B711 (2005) 112, arXiv:hep-th/0411201

[49] G. L. Kane, P. Kumar, J. D. Lykken, T. T. Wang, [arXiv: hep-ph/0411125

[50] I. Antoniadis, E. Kiritsis and T. Tomaras, Phys. Lett. B486 (2000) 186, arXiv:hep-ph/0004214

[51] D. Berenstein, V. Jejjala and R. G. Leigh, hep-ph/0105042

[52] A. Sagnotti, Phys. Lett. B294 (1992) 196, arXiv:hep-th/9210127

[53] L. E. Ibanez, R. Rabdan, A. M. Uranga, Nucl. Phys. B542 (1999) 112, arXiv:hep-th/9808139

[54] I. Antoniadis, C. Bachas and E. Dudas, Nucl. Phys. B560 (1999) 93, arXiv:hep-th/9906039

[55] C.A. Scrucca, M. Serone, JHEP 9912 (1999) 024, arXiv:hep-th/9912108

[56] E. Kiritsis, Fortsch. Phys. 52 (2004) 200, arXiv:hep-th/0310001

[57] P. Anastasopoulos, JHEP 0308:005,2003, arXiv:hep-th/0306042

[58] N. Arkani-Hamed, S. Dimopoulos, G. Giudice, A. Romanino, Nucl. Phys. B709 (2005) 3, arXiv:hep-ph/0409232

[59] E. Dudas, C. Timirgaziu, Nucl. Phys. B716 (2005) 65, arXiv:hep-th/0502085
[60] R. Blumenhagen, D. Lust, T. Taylor, Nucl. Phys. B663 (2003) 319, arXiv:hep-th/0303016

[61] J. F. G. Cascales, A. M. Uranga, JHEP 0305 (2003) 011, arXiv:hep-th/0303024

[62] I. Antoniadis, T. Maillard, Nucl. Phys. B716 (2005) 3, arXiv:hep-th/0412008

[63] M. Bianchi, E. Trevigne, arXiv:hep-th/0502147

[64] E. Dudas, G. Pradisi, M. Nicolisi and A. Sagnotti, Nucl. Phys. B708 (2005) 3, arXiv:hep-th/0410101

[65] E. Dudas, J. Mourad, Phys. Lett. B486 (2000) 172, arXiv:hep-th/0004165

[66] R. Blumenhagen, A. Font, Nucl. Phys. B599 (2001) 241, arXiv:hep-th/0011269