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New standard model vacua from intersecting branes

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Abstract: We construct new D6-brane model vacua (non-supersymmetric) that have at low energy exactly the standard model spectrum (with right handed neutrinos). The minimal version of these models requires five stacks of branes. and the construction is based on D6-branes intersecting at angles in $D = 4$ type toroidal orientifolds of type I strings. Three U(1)'s become massive through their couplings to RR couplings and from the two surviving anomaly free U(1)'s, one is the standard model hypercharge generator while the extra anomaly free U(1) could be broken from its non-zero couplings to RR fields and also by triggering a vev to previously massive particles. We suggest that extra massless U(1)'s should be broken by requiring some intersection to respect $N = 1$ supersymmetry thus supporting the appearance of massless charged singlets at the supersymmetric intersection. Proton is stable as baryon number is gauged and its anomalies are cancelled through a generalized Green-Schwarz mechanism. Neutrinos are of Dirac type with small masses, as in the four stack standard models of `hep-th/0105155`, as a result of the existence of a similar PQ like-symmetry. The models are unique in the sense that they predict the existence of only one supersymmetric particle, the superpartner of ν_R.

Keywords: D-branes, Standard Model, Compactification and String Models.

Contents

1. Introduction

Over the years, one of the most difficult questions that string theory is facing is the selection of the particular vacuum that includes at low energy all the necessary ingredients of the observable standard model spectrum a low energies. In the absence of an underlying principle of picking up a particular vacuum, one can search for the model with the correct low energy particle content. Such attempts have by far been explored in the context of heterotic string theory as well to branes at singularities [1, 2]. The main characteristics of the models involved include the three generation massless spectrum of the standard model (SM) accompanied by the presence of exotic matter and/or gauge group factors. However, recently there has been some progress in the study of string models as it has been possible in [3] to derive, at low energy, just the SM spectrum together with right handed neutrinos. The models were studied in the context of intersecting branes and have some satisfactory properties including proton stability and small neutrino (of Dirac type) masses.

The purpose of this paper is to extend the four stack construction of type I four dimensional toroidal orientifolded six-torus construction of [3], to a different structure that involves one extra U(1) at the string scale and produces just the standard model (SM) at low energies. Note that in [3] one starts with a $U(3)_a \otimes U(2)_b \otimes U(1)_c \otimes U(1)_d$ open gauge group structure at string scale energies. Note that in our construction, as in [3], right handed neutrinos are present along with the SM particle context at low energies.

Additional non-supersymmetric models along the same Type I backgrounds, which give at low energy the SM structure, have been constructed in [4]. In the latter case one starts with a four dimensional type I background on an orientifolded six dimensional factorized torus, which at the string scale includes[1] a Pati-Salam gauge group. The models incorporate a number of interesting properties, like proton stability and small neutrino masses.

In this work, one starts with five stacks of branes, namely with an $U(3)_a \otimes U(2)_b \otimes U(1)_c \otimes U(1)_d \otimes U(1)_e$ at the string scale and get at low energy only the SM with right handed neutrinos. The models also allow different generations of leptons-neutrinos to be placed at different intersections, that could have interesting implications for the phenomenology of the models. We also note that from the five string scale U(1)'s, four couple to RR fields and one survives massless at low energies. The latter U(1) corresponds to the hypercharge of the SM.

The models are non-supersymmetric and are build on a background of D6-branes intersecting each other at non-trivial angles [5], in an orientifolded factorized six-torus, while O_6 orientifold planes are on top of D6-branes [3]. Note that the latter picture is just the T-dual of type I backgrounds of [6, 7] which make use of D9 branes with background fluxes. Also, we note that the studied backgrounds are T-dual to models with magnetic deformations [8].

Additional models on backgrounds, in the context of intersecting branes, have been discussed in [9]–[15]. For an other proposal for realistic SM D-brane model building, based not on a particular string construction, see [16].

The proposed models have[2] the following distinctive features:

- The model starts with a gauge group at the string scale $U(3) \times U(2) \times U(1) \times U(1) \times U(1)$. The use of Green-Schwarz mechanism in the model renders three of the U(1)'s massive while a combination of the other two remaining anomaly free U(1)'s one makes the standard model hypercharge, while the other one[3] gets broken by its non-zero coupling to RR fields and by turning on $N = 1$ SUSY at an intersection, while keeping the rest of the intersections non-supersymmetric.

- Neutrinos get a Dirac mass, as lepton number L is a gauged symmetry, of the right order[4] in consistency with the LSND neutrino oscillation experiments [17] as a consequence of the existence of a PQ-like symmetry related to chiral symmetry breaking.

- Proton is stable due to the fact that baryon number B is an unbroken gauged global symmetry surviving at low energies whose anomalies cancel through a generalized Green-Schwarz mechanism.

[1]The string scale gauge group structure, that includes four stacks of branes, that is an $U(4)_c \otimes U(2)_L \otimes U(2)_R \otimes U(1)$.

[2]The different classes of models of this work maintain essential characteristics of the classes of models discussed in [3], including proton stability, and sizes of neutrino masses within experimental limits.

[3]In an orthogonal basis, this U(1) will be broken only by the vev of a scalar created by turning $N = 1$ SUSY on an intersection.

[4]The same mechanism was employed in [3] for the four stack D6 orientifold counterpart of the present SM's.

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The paper is organized as follows. In section 2 we describe the general characteristics of the new standard model vacua with particular emphasis on how to calculate the fermionic spectrum from intersecting branes as well providing the classification of multi-parameter solutions to the RR tadpole cancellation conditions. In section 3 we examine the cancellation of U(1) anomalies via a generalized Green-Schwarz (GS) mechanism [18, 19, 20] examining the conditions under which the hypercharge generator remains light at low energies. In section 4 we study the Higgs sector providing for the tachyonic scalars that are used for the electroweak symmetry breaking of the model. We also discuss the breaking of the extra anomaly free, other than hypercharge, U(1), by turning on $N = 1$ supersymmetry at an intersection. In section 5, we examine the Yukawa couplings and the smallness of neutrino masses. Section 6 contains our conclusions.

2. New SM vacua from intersecting branes

In the present work, we are going to describe new type I compactification vacua, that have as their low energy limit just the observable standard model interactions. The proposed three generation non-supersymmetric standard models make use of five stacks of branes at the string scale. They will originate from D6-branes wrapping on 3-cycles of toroidal orientifolds of type IIA in four dimensions. Important characteristic of all vacua coming from these type I constructions is the replication, at each intersection, of the massless fermion spectrum by an equal number of massive particles in the same representations and with the same quantum numbers.

Next, we describe the construction of the standard model. It is based on type I string with D9-branes compactified on a six-dimensional orientifolded torus T^6, where internal background gauge fluxes on the branes are turned on. If we perform a T-duality transformation on the x^4, x^5, x^6, directions the D9-branes with fluxes are translated into D6-branes intersecting at angles. Note that the branes are not parallel to the orientifold planes. Furthermore, we assume that the $D6_a$-branes are wrapping 1-cycles (n_a^i, m_a^i) along each of the ith-T^2 torus of the factorized T^6 torus, namely $T^6 = T^2 \times T^2 \times T^2$. That means that we allow our torus to wrap factorized 3-cycles, that can unwrap into products of three 1-cycles, one for each T^2. We define the homology of the 3-cycles as

$$[\Pi_a] = \prod_{i=1}^{3} \left(n_a^i [a_i] + m_a^i [b_i] \right) \tag{2.1}$$

while we define the 3-cycle for the orientifold images as

$$[\Pi_{a^\star}] = \prod_{i=1}^{3} \left(n_a^i [a_i] - m_a^i [b_i] \right) \tag{2.2}$$

In order to build the SM model structure a low energies, we consider five stacks of D6-branes giving rise to their world-volume to an initial gauge group $U(3) \times U(2) \times U(1) \times U(1) \times U(1)$ or $SU(3) \times SU(2) \times U(1)_a \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e$ at the string

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scale. Also, we consider the addition of NS B-flux [21], such that the tori involved are not orthogonal, thus avoiding an even number of families [6], and leading to effective tilted wrapping numbers as,

$$\left(n^i, m = \tilde{m}^i + \frac{n^i}{2}\right) ; \qquad n,\ \tilde{m}\ \in\ Z\,. \tag{2.3}$$

In this way we allow semi-integer values for the m-wrapping numbers.

Because of the $\Omega\mathcal{R}$ symmetry, where Ω is the worldvolume parity and $\mathcal{R}$ is the reflection on the T-dualized coordinates,

$$T(\Omega\mathcal{R})T^{-1} = \Omega\mathcal{R}\,, \tag{2.4}$$

each $D6_a$-brane 1-cycle, must have its $\Omega\mathcal{R}$ image partner $(n_a^i, -m_a^i)$.

Chiral fermions gets localized at the intersections between branes, by stretched open strings between intersecting D6-branes [5]. Subsequently, the chiral spectrum of the model may be obtained by solving simultaneously the intersection constraints coming from the existence of the different sectors and the RR tadpole cancellation conditions. Note that in the models we examine in this work, there are a number of different sectors, which should be taken into account when computing the chiral spectrum. We denote the action of ΩR on a sector a, b, by $a^\star, b^\star$, respectively. The possible sectors are:

- The $ab+ba$ sector: involves open strings stretching between the $D6_a$ and $D6_b$ branes. Under the ΩR symmetry this sector is mapped to its image, $a^\star b^\star + b^\star a^\star$ sector. The number, I_{ab}, of chiral fermions in this sector, transforms in the bifundamental representation $(\mathbb{N}_a, \bar{N}_a)$ of $U(\mathbb{N}_a) \times U(\mathbb{N}_b)$, and reads

$$I_{ab} = [\Pi_a]\cdot[\Pi_b] = \left(n_a^1 m_b^1 - m_a^1 n_b^1\right)\left(n_a^2 m_b^2 - m_a^2 n_b^2\right)\left(n_a^3 m_b^3 - m_a^3 n_b^3\right), \tag{2.5}$$

  where I_{ab} is the intersection number of the wrapped cycles. Note that with the sign of I_{ab} intersection, we denote the chirality of the fermions, where $I_{ab} > 0$ denotes left handed fermions. Also negative multiplicity denotes opposite chirality.

- The $ab^\star + b^\star a$ sector: It involves chiral fermions transforming into the $(\mathbb{N}_a, N_b)$ representation with multiplicity given by

$$I_{ab^\star} = [\Pi_a]\cdot[\Pi_{b^\star}] = \ -\left(n_a^1 m_b^1 + m_a^1 n_b^1\right)\left(n_a^2 m_b^2 + m_a^2 n_b^2\right)\left(n_a^3 m_b^3 + m_a^3 n_b^3\right). \tag{2.6}$$

  Under the ΩR symmetry it transforms to itself.

- the $aa^\star$ sector: under the ΩR symmetry it transforms to itself. From this sector the invariant intersections will give $8 m_a^1 m_a^2 m_a^3$ fermions in the antisymmetric representation and the non-invariant intersections that come in pairs provide us with $4 m_a^1 m_a^2 m_a^3 (n_a^1 n_a^2 n_a^3 - 1)$ additional fermions in the symmetric and antisymmetric representation of the $U(\mathbb{N}_a)$ gauge group. However as we explain later, these sectors will be absent from our models.

Any vacuum derived from the previous intersection constraints is subject additionally to constraints coming from RR tadpole cancellation conditions [6]. That demands cancellation of D6-branes charges,[5] wrapping on three cycles with homology $[\Pi_a]$ and O6-plane

[5]Taken together with their orientifold images $(n_a^i, -m_a^i)$ wrapping on three cycles of homology class $[\Pi_{\alpha'}]$.

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7-form charges wrapping on 3-cycles with homology $[\Pi_{O_6}]$. Note that the RR tadpole cancellation conditions in terms of cancellations of RR charges in homology are

$$\sum_a N_a[\Pi_a] + \sum_{\alpha'} N_{\alpha'}[\Pi_{\alpha'}] - 32[\Pi_{O_6}] = 0\,. \tag{2.7}$$

In explicit form, the RR tadpole conditions read

$$\begin{aligned}
&\sum_a N_a n_a^1 n_a^2 n_a^3 = 16\,,\\
&\sum_a N_a m_a^1 m_a^2 n_a^3 = 0\,,\\
&\sum_a N_a m_a^1 n_a^2 m_a^3 = 0\,,\\
&\sum_a N_a n_a^1 m_a^2 m_a^3 = 0\,.
\end{aligned} \tag{2.8}$$

That guarantees absence of non-abelian gauge anomalies. In D-brane model building, by considering a stacks of D-brane configurations with $N_a, a = 1, \cdots, N$, parallel branes one gets the gauge group $U(\mathbb{N}_1) \times U(\mathbb{N}_2) \times \cdots \times U(\mathbb{N}_a)$. Each $U(\mathbb{N}_i)$ factor will give rise to an $SU(\mathbb{N}_i)$ charged under the associated $U(1_i)$ gauge group factor that appears in the decomposition $SU(\mathbb{N}_a) \times U(1_a)$. For the five stack model that we examine in this work, the complete accommodation, where all other intersections are vanishing, of the fermion structure can be seen in table (1). We note a number of interesting comments:

(a) There are various gauged low energy symmetries in the models. They are defined in terms of the U(1) symmetries Q_a, Q_b, Q_c, Q_d, Q_e. The baryon number B is equal to $Q_a = 3B$, the lepton number is $L = Q_d + Q_e$ while $Q_a - 3Q_d - 3Q_e = 3(B - L)$. Also note that $Q_c = 2I_{3R}$, I_{3R} being the third component of weak isospin. Also, $3(B-L)$ and Q_c are free of triangle anomalies. The $U(1)_b$ symmetry plays the role of a Peccei-Quinn symmetry in the sence of having mixed SU(3) anomalies. From the study of Green-Schwarz mechanism, see later sections, we deduce that Baryon and Lepton number are unbroken gauged symmetries and thus proton should be stable, while Majorana masses for right handed neutrinos are not allowed. That means that mass terms for neutrinos should be of Dirac type.

The mixed anomalies A_{ij} of the four surplus U(1)'s with the non-abelian gauge groups $SU(\mathbb{N}_a)$ of the theory cancel through a generalized GS mechanism [18, 19], involving close string modes couplings to worldsheet gauge fields. Two combinations of the U(1)'s are anomalous and become massive, their orthogonal non-anomalous combinations survive, combining to a single U(1) that remains massless, the hypercharge.

(b) In order to cancel the appearance of exotic representations in the model appearing from the $aa^\star$ sector, in antisymmetric and symmetric representations of the $U(\mathbb{N}_a)$ group, we will impose the condition

$$\Pi_{i=1}^3 m^i = 0\,. \tag{2.9}$$

Matter Fields		Intersection	Q_a	Q_b	Q_c	Q_d	Q_e	Y
Q_L	$(3,2)$	$I_{ab}=1$	1	-1	0	0	0	1/6
q_L	$2(3,2)$	$I_{ab^*}=2$	1	1	0	0	0	1/6
U_R	$3(\bar{3},1)$	$I_{ac}=-3$	-1	0	1	0	0	$-2/3$
D_R	$3(\bar{3},1)$	$I_{ac^*}=-3$	-1	0	-1	0	0	1/3
L	$2(1,2)$	$I_{bd}=-2$	0	-1	0	1	0	$-1/2$
l_L	$(1,2)$	$I_{be}=-1$	0	-1	0	0	1	$-1/2$
N_R	$2(1,1)$	$I_{cd}=2$	0	0	1	-1	0	0
E_R	$2(1,1)$	$I_{cd^*}=-2$	0	0	-1	-1	0	1
ν_R	$(1,1)$	$I_{ce}=1$	0	0	1	0	-1	0
e_R	$(1,1)$	$I_{ce^*}=-1$	0	0	-1	0	-1	1

Table 1: Low energy fermionic spectrum of the five stack string scale $SU(3)_C \otimes SU(2)_L \otimes U(1)_a \otimes U(1)_b \otimes U(1)_c \otimes U(1)_d \otimes U(1)_e$, type I D6-brane model together with its U(1) charges. Note that at low energies only the SM gauge group $SU(3) \otimes SU(2)_L \otimes U(1)_Y$ survives.

The solutions satisfying simultaneously the intersection constraints and the cancellation of the RR crosscap tadpole constraints are parametric. They are given in table (2). The solutions represent the most general solution of the RR tadpoles and they depend on five integer parameters n_a^2, n_d^2, n_e^2, n_b^1, n_c^1, the phase parameters $\epsilon = \pm 1$, $\tilde{\epsilon} = \pm 1$, and the NS-background parameter $\beta^i = 1 - b^i$, which is associated to the presence of the NS B-field by $b^i = 0,\ 1/2$. Note that the different solutions to the tadpole constraints represent deformations of the D6 brane RR charges within the same homology class. In the rest of the paper we will be discussing, for simplicity the case with $\epsilon = \tilde{\epsilon} = 1$. The multiparameter tadpole solutions of table 2 represent deformations of the D6-brane intersection spectrum of table 1, within the same homology class of the factorizable three-cycles. By using the tadpole solutions of table 2 in (2.8) all tadpole equations but the first are automatically satisfied, the[6] latter yielding:[7]

$$\frac{9n_a^2}{\beta^1} + 2\frac{n_b^1}{\beta^2} + \frac{n_d^2}{\beta^1} + \frac{n_e^2}{\beta^1} + N_D \frac{2}{\beta^1 \beta^2} = 16\,. \tag{2.10}$$

Note that we had added the presence of extra N_D branes (hidden branes). Their contribution to the RR tadpole conditions is best described by placing them in the three-factorizable cycle

$$N_D \left(\frac{1}{\beta^1}, 0\right) \left(\frac{1}{\beta^2}, 0\right) (2, m_D^3) \tag{2.11}$$

[6]We have added an arbitrary number of N_D branes which don't contribute to the rest of the tadpoles and intersection number constraints.

[7]We have set for simplicity $\tilde{\epsilon} = 1$.

N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 3$	$(1/\beta^1, 0)$	$(n_a^2, \epsilon\beta^2)$	$(3, \tilde{\epsilon}/2)$
$N_b = 2$	$(n_b^1, -\epsilon\beta^1)$	$(1/\beta^2, 0)$	$(\tilde{\epsilon}, 1/2)$
$N_c = 1$	$(n_c^1, \epsilon\beta^1)$	$(1/\beta^2, 0)$	$(0, 1)$
$N_d = 1$	$(1/\beta^1, 0)$	$(n_d^2, 2\epsilon\beta^2)$	$(1, -\tilde{\epsilon}/2)$
$N_e = 1$	$(1/\beta^1, 0)$	$(n_e^2, \epsilon\beta^2)$	$(1, -\tilde{\epsilon}/2)$

Table 2: Tadpole solutions of D6-branes wrapping numbers giving rise to the standard model gauge group and spectrum at low energies. The solutions depend on five integer parameters, n_a^2, n_d^2, n_e^2, n_b^1, n_c^1, the NS-background β^i and the phase parameters $\epsilon = \pm 1$, $\tilde{\epsilon} = \pm 1$.

and we have set $m_D^3 = 0$. To see clearly the cancellation of tadpoles, we have to choose a consistent numerical set of wrappings, e.g.

$$n_a^2 = 1\,, \quad n_b^1 = 1\,, \quad n_c^1 \in Z\,, \quad n_d^2 = -1\,, \quad n_e = 1\,, \quad \beta^1 = \frac{1}{2}\,, \quad \beta^2 = 1, \epsilon = 1\,. \tag{2.12}$$

With the above choices, all tadpole conditions but the first are satisfied, the latter is satisfied when we add one $\bar{D}6$ brane, e.g. $N_D = -1$. Thus the tadpole structure[8] becomes

$$\begin{aligned}
N_a &= 3 \quad (2,0)(1,1)\left(3, \frac{1}{2}\right) \\
N_b &= 2 \quad \left(1, -\frac{1}{2}\right)(1,0)\left(1, \frac{1}{2}\right) \\
N_c &= 1 \quad \left(n_c^1, \frac{1}{2}\right)(1,0)(0,1) \\
N_d &= 1 \quad (2,0)(-1,2)\left(1, -\frac{1}{2}\right) \\
N_e &= 1 \quad (2,0)(1,1)\left(1, -\frac{1}{2}\right).
\end{aligned} \tag{2.13}$$

Actually, the satisfaction of the tadpole conditions is independent of n_c^1. Thus, when all other parameters are fixed, n_c^1 is a global parameter that can vary. However, finally it will be fixed in terms of the remaining parameters once we specify, the tadpole subclass that corresponds to the massless spectrum with the hypercharge embedding of the standard model.

Note that there are always choices of wrapping numbers of wrapping numbers that satisfy the RR tadpole constraints without the need of adding extra parallel branes, e.g. the following choice satisfies all RR tadpoles

$$n_a^2 = 1\,, \quad n_b^1 = 3\,, \quad n_c^1 \in Z\,, \quad n_d^2 = -3\,, \quad n_e = -1\,, \quad \beta^1 = \frac{1}{2}\,, \quad \beta^2 = 1\,, \quad \epsilon = 1\,. \tag{2.14}$$

[8]Note that the parameter n_c^1 should be defined such that its choice is consistent with a tilted tori, e.g. $n_c^1 = 1$.

with cycle wrapping numbers

$$\begin{aligned}
N_a &= 3 \quad (2,0)(1,1)\left(3, \frac{1}{2}\right) \\
N_b &= 2 \quad \left(3, -\frac{1}{2}\right)(1,0)\left(1, \frac{1}{2}\right) \\
N_c &= 1 \quad \left(n_c^1, \frac{1}{2}\right)(1,0)(0,1) \\
N_d &= 1 \quad (2,0)(-3,2)\left(1, -\frac{1}{2}\right) \\
N_e &= 1 \quad (2,0)(-1,1)\left(1, -\frac{1}{2}\right).
\end{aligned} \tag{2.15}$$

Another alternative choice, satisfied by all RR tadpoles will be

$$n_a^2 = 1\,, \quad n_b^1 = 1\,, \quad n_c^1 \in Z\,, \quad n_d^2 = 2\,, \quad n_e = 1\,, \quad \beta^1 = 1\,, \quad \beta^2 = \frac{1}{2}\,, \quad \epsilon = 1\,. \tag{2.16}$$

with cycle wrapping numbers

$$\begin{aligned}
N_a &= 3 \quad (1,0)\left(1, \frac{1}{2}\right)\left(3, \frac{1}{2}\right) \\
N_b &= 2 \quad (1,-1)(2,0)\left(1, \frac{1}{2}\right) \\
N_c &= 1 \quad \left(n_c^1, \frac{1}{2}\right)(2,0)(0,1) \\
N_d &= 1 \quad (1,0)(2,1)\left(1, -\frac{1}{2}\right) \\
N_e &= 1 \quad (1,0)\left(1, \frac{1}{2}\right)\left(1, -\frac{1}{2}\right).
\end{aligned} \tag{2.17}$$

(c) the hypercharge operator in the model is defined as a linear combination of the three generators of the SU(3), $\mathrm{U}(1)_c$, $\mathrm{U}(1)_d$, $\mathrm{U}(1)_e$ gauge groups:

$$Y = \frac{1}{6}\,\mathrm{U}(1)_a - \frac{1}{2}\,\mathrm{U}(1)_c - \frac{1}{2}\,\mathrm{U}(1)_d - \frac{1}{2}\,\mathrm{U}(1)_e\,. \tag{2.18}$$

3. Cancellation of U(1) anomalies

In general the mixed anomalies A_{ij} of the four U(1)'s with the non-abelian gauge groups are given by

$$A_{ij} = \frac{1}{2}(I_{ij} - I_{ij^\star})N_i\,. \tag{3.1}$$

Moreover, analyzing the mixed anomalies of the extra U(1)'s with the non-abelian gauge groups $\mathrm{SU}(3)_c$, $\mathrm{SU}(2)_b$, we can conclude that there are two anomaly free combinations Q_c, $Q_a - 3Q_d - 3Q_e$. Also, note that the gravitational anomalies cancel since D6-branes never intersect O6-planes. In the orientifolded type I torus models gauge anomaly cancellation [18, 20] is guaranteed through a generalized GS mechanism [3] that uses the 10-dimensional RR gauge fields C_2 and C_6 and gives at four dimensions the following couplings to gauge fields

$$N_a m_a^1 m_a^2 m_a^3 \int_{M_4} B_2^o \wedge F_a \, ; \qquad n_b^1 n_b^2 n_b^3 \int_{M_4} C^o \wedge F_b \wedge F_b \, , \tag{3.2}$$

$$N_a n^J n^K m^I \int_{M_4} B_2^I \wedge F_a \, ; \qquad n_b^I m_b^J m_b^K \int_{M_4} C^I \wedge F_b \wedge F_b \, , \tag{3.3}$$

where $C_2 \equiv B_2^o$ and $B_2^I \equiv \int_{(T^2)^J \times (T^2)^K} C_6$ with $I = 1, 2, 3$ and $I \neq J \neq K$. Notice the four dimensional duals of B_2^o, B_2^I:

$$C^o \equiv \int_{(T^2)^1 \times (T^2)^2 \times (T^2)^3} C_6 \, ; \qquad C^I \equiv \int_{(T^2)^I} C_2 \, , \tag{3.4}$$

where $dC^o = -\star dB_2^o$, $dC^I = -\star dB_2^I$.

The triangle anomalies (3.1) cancel from the existence of the string amplitude involved in the GS mechanism [19] in four dimensions [18]. The latter amplitude, where the $\mathrm{U}(1)_a$ gauge field couples to one of the propagating B_2 fields, coupled to dual scalars, that couple in turn to two SU($\mathbb{N}$) gauge bosons, is proportional [3] to

$$-N_a m_a^1 m_a^2 m_a^3 n_b^1 n_b^2 n_b^3 - N_a \sum_I n_a^I n_a^J n_b^K m_a^I m_b^J m_b^K \, , \qquad I \neq J, K \tag{3.5}$$

Taking into account the phenomenological requirements of eq. (2.9) the RR couplings B_2^I of (3.3) then appear into three terms:[9]

$$\begin{aligned}
&B_2^1 \wedge \left(\frac{-2\epsilon\tilde{\epsilon}\beta^1}{\beta^2} \right) F^b \, , \\
&B_2^2 \wedge \left(\frac{\epsilon\beta^2}{\beta^1} \right) (9F^a + 2F^d + F^e) \, , \\
&B_2^3 \wedge \left(\frac{3\tilde{\epsilon} n_a^2}{2\beta^1} F^a + \frac{n_b^1}{\beta^2} F^b + \frac{n_c^1}{\beta^2} F^c - \frac{\tilde{\epsilon} n_d^2}{2\beta^1} F^d - \frac{\tilde{\epsilon} n_e^2}{2\beta^1} F^e \right) .
\end{aligned} \tag{3.6}$$

At this point we should list the couplings of the dual scalars C^I of B_2^I that required to cancel the mixed anomalies of the five U(1)'s with the non-abelian gauge groups SU($\mathbb{N}_a$). They are given by

$$\begin{aligned}
&C^1 \wedge \left[\frac{\epsilon\tilde{\epsilon}\beta^2}{2\beta^1} (F^a \wedge F^a) - \frac{\epsilon\tilde{\epsilon}\beta^2}{\beta^1} (F^d \wedge F^d) - \frac{\epsilon\tilde{\epsilon}\beta^2}{2\beta^1} F^e \wedge F^e) \right] , \\
&C^2 \wedge \left[\frac{-\epsilon\beta^1}{2\beta^2} (F^b \wedge F^b) + \frac{\epsilon\beta^1}{\beta^2} (F^c \wedge F^c) \right] , \\
&C^o \wedge \left(\frac{3n_a^2}{\beta^1} (F^a \wedge F^a) + \frac{\tilde{\epsilon} n_b^1}{\beta^2} (F^b \wedge F^b) + \frac{n_d^2}{\beta^1} (F^d \wedge F^d) + \frac{n_e^2}{\beta^1} (F^e \wedge F^e) \right) ,
\end{aligned} \tag{3.7}$$

[9] For convenience we have included the dependence on $\epsilon, \tilde{\epsilon}$ parameters.

Notice that the RR scalar B_2^0 does not couple to any field F^i as we have imposed the condition (2.9) which excludes the appearance of any exotic matter.

Looking at (3.6) we can conclude that there are two anomalous U(1)'s that become massive through their couplings to the RR fields. They are the model independent fields, $U(1)_b$ and the combination $9\,U(1)_a + 2\,U(1)_d + U(1)_e$, which become massive through their couplings to the RR 2-form fields B_2^1, B_2^2 respectively. In addition, there is a model dependent, non-anomalous and massive U(1) field coupled to B_2^3 RR field. That means that the two massless and anomaly free combinations are $U(1)_c$ and $U(1)_a - 3\,U(1)_d - 3\,U(1)_e$. Also, note that the mixed anomalies A_{ij} are cancelled by the GS mechanism set by the couplings (3.6), (3.7).

The question we want to address at this point is how we can, from the general class of models, associated with the generic SM's of tables 1 and 2, pick up the subclass that corresponds to the ones associated with just the observable SM at low energies. Clearly, for this to happen we have to identify the subclass of tadpole solutions of table (2) that corresponds to the hypercharge assignment (2.18) of the standard model[10] spectrum.

In general, the generalized Green-Schwarz mechanism that cancels non-abelian anomalies of the U(1)'s to the non-abelian gauge fields involves couplings of closed string modes to the U(1) field strengths[11] in the form

$$\sum_a f_a^i B_a \wedge \mathrm{tr}(F_i)\,. \tag{3.9}$$

Effectively, the mixture of couplings in the form

$$A_{ik} + \sum_a f_a^i \; g_a^k = 0 \tag{3.10}$$

cancels the all non-abelian U(1) gauge anomalies. That means that if we want to keep some U(1) massless we have to keep it decoupled from some closed string mode couplings that can make it massive, that is

$$\sum_a \left(\frac{1}{6} f_a^{\alpha} - \frac{1}{2} f_a^c - \frac{1}{2} f_a^d - \frac{1}{2} f_a^e \right) = 0\,. \tag{3.11}$$

In conclusion, the combination of the U(1)'s which remains light at low energies, is

$$(3n_a^2 + 3n_d^2 + 3n_e^2) \neq 0\,, \qquad Q^l = n_c^1 (Q_a - 3Q_d - 3Q_e) - \frac{3\tilde{\epsilon}\beta^2 (n_a^2 + n_d^2 + n_e^2)}{2\beta^1} Q_c\,. \tag{3.12}$$

The subclass of tadpole solutions of (3.12) having the SM hypercharge assignment at low energies is exactly the one which is proportional to (2.18). It satisfies the condition,

$$n_c^1 = \frac{\tilde{\epsilon}\beta^2}{2\beta^1} (n_a^2 + n_d^2 + n_e^2)\,. \tag{3.13}$$

[10] At this point, we recall an argument that have appeared in [3].

[11] In addition, to the couplings of the Poincare dual scalars η_a of the fields B_a,

$$\sum_a g_a^k \, \eta_a \;\, \mathrm{tr}(F^k \wedge F^k)\,. \tag{3.8}$$

We note that there is one extra anomaly free, model dependent U(1) beyond the hypercharge combination, and orthogonal to the latter, which is[12]

$$Q^N = \frac{3\tilde{\epsilon}\beta^2}{2\beta^1}(Q_a - 3Q_d - 3Q_e) + 19n_c^1 Q_c\,. \tag{3.14}$$

Lets us summarize what we have found up to now. The tadpole solutions of table 2, taking into account the condition (3.13), give at low energies classes of models that have the low energy spectrum of the SM with the correct hypercharge assignments. At this level the gauge group content of the model includes beyond $SU(3) \otimes SU(2) \otimes U(1)_Y$ the additional $U(1)^N$ generator and all SM particles gets charged under the additional $U(1)^N$ symmetry. However, notice that Q^N has a non-zero coupling to RR field B_2^3. That is it receives a mass of order of the string scale M_s and disappears from the low energy spectrum. Hence at low energy only the SM remains.

In the next sections we will see that in the present five stack constructions it is possible to use an additional mechanism to break this extra U(1) symmetry, complementary to the one associated with the RR fields. In involves the usual mechanism of giving a vev to a scalar field. In this case we will have to require that the intersection where the right handed neutrino is localized, respects $N = 1$ supersymmetry. In the latter case the immediate effect on obtaining just the SM at low energies will be one additional linear condition on the RR tadpole solutions of table 2. We note that when $n_c^1 = 0$, it is possible to have massless in the low energy spectrum both the U(1) generators, Q_c, and the B-L generator $(1/3)(Q_a - 3Q_d - 3Q_e)$ as long as $n_c^1 = 0$, $n_a^2 = -n_d^2 - n_e^2$.

4. Electroweak Higgs and symmetry breaking from open string tachyons

The mechanism of electroweak symmetry breaking is a well understood effect at the level of gauge theories with or without supersymmetry. At the string theory level the mechanism is believed to take place either by using open string tachyonic modes between parallel branes or following a recent suggestion using brane recombination [13]. In the former mechanism, the mass of the Higgs field receives contributions from the distance between the branes b, c (b, $c^\star$) which are parallel across the second tori (see table 2). By varying the distance between the branes, across the 2nd tori, the Higgs mass could become tachyonic signalling electroweak symmetry breaking. In the latter mechanism, one of the two factorizable b-branes, making the SU(2)-stack, recombine with the single U(1) c-brane into a single non-factorizable j-brane. After the recombination the electroweak symmetry is broken, a result better seen from the new intersection numbers produced. Note that the latter procedure is topological and cannot be described using field theoretical methods. During the recombination process instead of us working with our usual wrapping numbers (2.3), one must work with cycles associated with a non-factorizable T^6 torus. Also one has to

[12]We note that alternatively, in an orthogonal basis, the three U(1)'s are coupled to B_2^i's, the 4th U(1) is the hypercharge (3.12) and its condition (3.13), the fifth U(1) is $U(1)^{(5)} = (-\frac{3}{29} + \frac{3}{28})F^a - \frac{1}{29}F^d + \frac{1}{28}F^e$, the latter surviving massless the Green-Schwarz mechanism when $n_a^2 = -(28/9)n_d^2$. In this case, $U(1)^{(5)}$ should be broken by the vev of $s\nu_R$ only (see section 5).

preserve the RR charge in homology before and after the recombination process. In this work, we will follow the former method.

4.1 The angle structure

We have up to now describe the appearance in the R-sector of open strings of I_{ab} massless chiral fermions in the D-brane intersections that transform under bifundamental representations $N_a, \bar{N}_b$. We should note that in backgrounds with intersecting branes, besides the actual presence of massless fermions at each intersection, we have evident the presence of an equal number of massive scalars (MS), in the NS-sector, in exactly the same representations as the massless fermions [10]. The mass of the MS is of order of the string scale. In some cases, it is possible that some of those MS may become tachyonic, triggering a potential that looks like the Higgs potential of the SM, especially when their mass, that depends on the angles between the branes, is such that is decreases the world volume of the 3-cycles involved in the recombination process of joining the two branes into a single one [22].

The models examined in this work, are based on orientifolded six-tori on type I strings. In those configurations the bulk has $\mathcal{N} = 4$ SUSY. Lets us now describe the open string sector of the model. In order to describe the open string spectrum we introduce a four dimensional twist [9, 10] vector υ_θ, whose I-th entry is given by ϑ_{ij}, with ϑ_{ij} the angle between the branes i and j-branes. After GSO projection the states are labeled by a four dimensional twisted vector $r + \upsilon_\theta$, where $\sum_I r^I$ =odd and $r_I \in \mathbf{Z}, \mathbf{Z{+}1/2}$ for NS, R sectors respectively. The Lorentz quantum numbers are denoted by the last entry. The mass operator for the states is provided by:

$$\alpha' M_{ij}^2 = \frac{Y^2}{4\pi^2\alpha'} + N_{bos}(\vartheta) + \frac{(r+\upsilon)^2}{2} - \frac{1}{2} + E_{ij}\,, \tag{4.1}$$

where E_{ij} the contribution to the mass operator from bosonic oscillators, and $N_{osc}(\vartheta)$ their number operator, with

$$E_{ij} = \sum_I \frac{1}{2}|\vartheta_I|(1 - |\vartheta_I|)\,, \tag{4.2}$$

and Y measures the minimum distance between branes for minimum winding states.

If we represent the twisted vector $r + \upsilon$, by $(\vartheta_1, \vartheta_2, \vartheta_3, 0)$, in the NS open string sector, the lowest lying states are given[13] by:

State	**Mass**
$(-1+\vartheta_1, \vartheta_2, \vartheta_3, 0)$	$\alpha' M^2 = \frac{1}{2}(-\vartheta_1 + \vartheta_2 + \vartheta_3)$
$(\vartheta_1, -1+\vartheta_2, \vartheta_3, 0)$	$\alpha' M^2 = \frac{1}{2}(\vartheta_1 - \vartheta_2 + \vartheta_3)$
$(\vartheta_1, \vartheta_2, -1+\vartheta_3, 0)$	$\alpha' M^2 = \frac{1}{2}(\vartheta_1 + \vartheta_2 - \vartheta_3)$
$(-1+\vartheta_1, -1+\vartheta_2, -1+\vartheta_3, 0)$	$\alpha' M^2 = 1 - \frac{1}{2}(\vartheta_1 + \vartheta_2 + \vartheta_3)$

(4.3)

[13] We assumed $0 \le \vartheta_i \le 1$.

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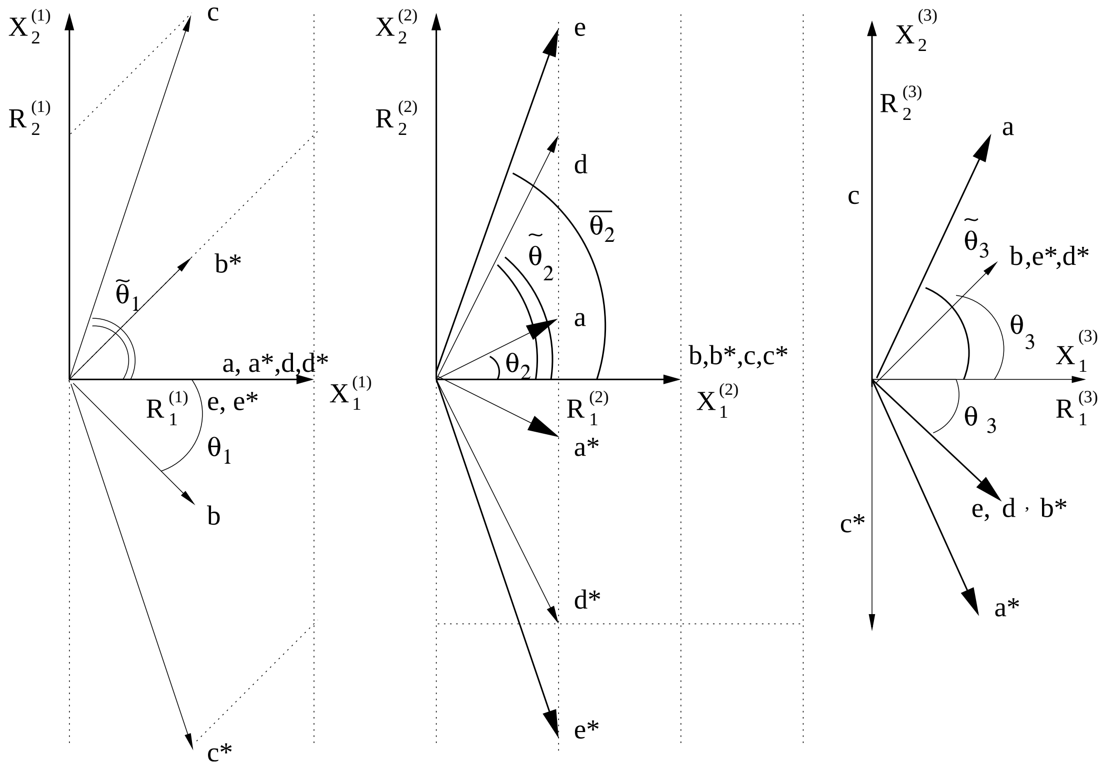

Figure 1: Assignment of angles between D6-branes on the five stack type I model giving rise to the SM at low energies. The angles between branes are shown on a product of $T^2 \times T^2 \times T^2$. We have chosen $\beta^1 = \beta^2 = 1$, $n_b^1, n_c^1, n_a^2, n_d^2 > 0$, $\epsilon = \tilde{\epsilon} = 1$.

The angles at the ten different intersections can be expressed in terms of the tadpole solutions parameters. Let us define the angles:

$$\begin{aligned}
\theta_1 &= \frac{1}{\pi}\tan^{-1}\frac{\beta^1 R_2^{(1)}}{n_b^1 R_1^{(1)}}\,, & \theta_2 &= \frac{1}{\pi}\tan^{-1}\frac{\beta^2 R_2^{(2)}}{n_a^2 R_1^{(2)}}\,, & \theta_3 &= \frac{1}{\pi}\tan^{-1}\frac{R_2^{(3)}}{6R_1^{(3)}}\,, \\
\tilde{\theta_1} &= \frac{1}{\pi}\tan^{-1}\frac{\beta^1 R_2^{(1)}}{n_c^1 R_1^{(1)}}\,, & \tilde{\theta_2} &= \frac{1}{\pi}\tan^{-1}\frac{2\beta^2 R_2^{(2)}}{n_d^2 R_1^{(2)}}\,, & \tilde{\theta_3} &= \frac{1}{\pi}\tan^{-1}\frac{R_2^{(3)}}{2R_1^{(3)}}\,, \\
\bar{\theta_2} &= \frac{1}{\pi}\tan^{-1}\frac{\beta^2 R_2^{(2)}}{n_e^2 R_1^{(2)}}\,, & \bar{\theta_3} &= \frac{1}{\pi}\tan^{-1}\frac{R_2^{(3)}}{2R_1^{(3)}}\,,
\end{aligned} \tag{4.4}$$

where $R_{1,2}^{(i)}$ are the compactification radii for the three $i = 1, 2, 3$ tori, namely projections of the radii onto the $X_{1,2}^{(i)}$ directions when the NS flux B field, b^i, is turned on and we have chosen for convenience $\epsilon = \tilde{\epsilon} = 1$.

At each of the ten non-trivial intersections we have the presence of four states t_i, $i = 1, \ldots, 4$, associated to the states (4.3). Hence we have a total of forty different scalars in the model.[14]

[14]In figure one, we can see the D6 branes angle setup in the present models.

The following mass relations hold between the different intersections of the classes of models:

$$
\begin{aligned}
m^2_{ab}(t_2) + m^2_{ab}(t_3) &= m^2_{ab^\star}(t_2) + m^2_{ab^\star}(t_3) = m^2_{bd}(t_2) + m^2_{bd}(t_3)\,, \\
m^2_{ac}(t_2) + m^2_{ac}(t_3) &= m^2_{ac^\star}(t_2) + m^2_{ac^\star}(t_3) = m^2_{cd}(t_2) + m^2_{cd}(t_3)\,, \\
m^2_{ce^\star}(t_1) + m^2_{ce^\star}(t_2) &= m^2_{cd^\star}(t_1) + m^2_{cd^\star}(t_2)\,,
\end{aligned}
\tag{4.5}
$$

or equivalently

$$
\begin{aligned}
m^2_{Q_L}(t_2) + m^2_{Q_L}(t_3) &= m^2_{q_L}(t_2) + m^2_{q_L}(t_3) = m^2_{L}(t_2) + m^2_{L}(t_3)\,, \\
m^2_{U_R}(t_2) + m^2_{U_R}(t_3) &= m^2_{D_R}(t_2) + m^2_{D_R}(t_3) = m^2_{N_R}(t_2) + m^2_{N_R}(t_3)\,, \\
m^2_{e_R}(t_1) + m^2_{e_R}(t_2) &= m^2_{E_R}(t_1) + m^2_{E_R}(t_2)\,,
\end{aligned}
\tag{4.6}
$$

We note that in this work, we will not discuss the stability conditions for absence of tachyonic scalars such that the D-brane configurations discussed will be stable as this will be discussed elsewhere. Similar conditions have been examined before in [3, 4].

4.2 Tachyon Higgs mechanism in detail

In this section, we will study the electroweak Higgs sector of the models. We note that below the string scale the massless spectrum of the model is that of the SM with all particles having the correct hypercharge assignments but with the gauge symmetry being $SU(3) \otimes SU(2) \otimes U(1) \otimes U(1)^N$. For the time being we will accept that the additional $U(1)^N$ generator breaks to a scale higher than the scale of electroweak symmetry breaking. The latter issue will be discussed in detail in the next section. Thus in the following we will focus our attention to the Higgs sector of the theory.

In general, tachyonic scalars stretching between two different branes can be used as Higgs scalars as they can become non-tachyonic by varying the distance between parallel branes. This happens in the models under discussion as the complex scalars $h^{\pm}$, $H^{\pm}$ get localized between the b, c and between b, c^* branes respectively and can be interpreted from the field theory point of view [3] as Higgs fields which are responsible for the breaking the electroweak symmetry. We note that the intersection numbers of the b, c and $b, c^\star$ branes across the six-dimensional torus vanish as a result of the fact that the b, c and $b, c^\star$ branes are parallel across the second tori. The electroweak Higgs fields, appearing as H_i (resp. h_i), $i = 1, 2$, in table 3, come from the NS sector, from open strings stretching between the parallel b, $c^\star$ (resp. c) branes along the second tori, and from open strings stretching between intersecting branes along the first and third tori.

Initially, the Higgses of table 3, are part of the massive spectrum of fields localized in the intersections $bc, bc^\star$. However, we emphasize that the Higgses H_i, h_i become massless by varying the distance along the second tori between the $b, c^\star$, b, c branes respectively. In fact, a similar set of Higgs fields appear in the four stack models of [3], but obviously with different geometrical data. We should note that the representations of Higgs fields H_i, h_i is the maximum allowed by quantization. Their number is model dependent.

Intersection	EW breaking Higgs	Q_b	Q_c	Y
bc	h_1	1	-1	1/2
bc	h_2	-1	1	-1/2
$bc^\star$	H_1	-1	-1	1/2
$bc^\star$	H_2	1	1	-1/2

Table 3: Higgs fields responsible for electroweak symmetry breaking.

The number of complex scalar doublets present in the models is equal to the non-zero intersection number product between the bc, $bc^\star$ branes in the first and third complex planes. Thus

$$n_{H^\pm} = I_{bc^\star} = |\epsilon\beta_1(n_b^1 - n_c^1)|\,, \qquad n_{h^\pm} = I_{bc} = |\epsilon\beta_1(n_b^1 + n_c^1)|. \tag{4.7}$$

The precise geometrical data for the scalar doublets are

State	**Mass^2**
$(-1+\vartheta_1, 0, \vartheta_3, 0)$	$\alpha'(\text{Mass})^2_Y = \dfrac{Z_2}{4\pi^2} + \dfrac{1}{2}(\vartheta_3 - \vartheta_1)$
$(\vartheta_1, 0, -1+\vartheta_3, , 0)$	$\alpha'(\text{Mass})^2_X = \dfrac{Z_2}{4\pi^2} + \dfrac{1}{2}(\vartheta_1 - \vartheta_3)$

(4.8)

where $X = \{H^+_{bc^\star}, h^+_{bc}\}$, $Y = \{H^-_{bc^\star}, h^-_{bc}\}$ and Z_2 is the distance^2 in transverse space along the second torus, ϑ_1, ϑ_3 are the (relative) angles between the b-, $c^\star$ (for $H^\pm$) (or b, c for $h^\pm$) branes in the first and third complex planes.

Also we note the presence of two "Higgsino masses" at each of the bc or $bc^\star$ intersections, with the same quantum numbers and representations as the Higgs fields and masses corresponding to

State	**Mass^2**
$\left(-\frac{1}{2}+\vartheta_1, \mp\frac{1}{2}, -\frac{1}{2}+\vartheta_3, \pm\frac{1}{2}\right)$	$(\text{Mass})^2 = \dfrac{Z_2}{4\pi^2\alpha'}\,.$

(4.9)

We note that in this picture while the Higgs fields can be made massless by varying the distance between the branes, the Higginos are not massless and are part of the $N = 2$ massive spectrum accompanying the "massless" Higgs fields at the intersections bc, $bc^\star$.

As we noted the presence of scalar doublets $H^\pm, h^\pm$, can be seen as coming from the field theory mass matrix

$$(H_1^* \; H_2)\left(\mathbf{M^2}\right)\begin{pmatrix} H_1 \\ H_2^* \end{pmatrix} + (h_1^* \; h_2)\left(\mathbf{m^2}\right)\begin{pmatrix} h_1 \\ h_2^* \end{pmatrix} + \text{h.c.} \tag{4.10}$$

where

$$\mathbf{M^2} = M_s^2 \begin{pmatrix} Z_2^{(bc^*)} & \frac{1}{2}|\vartheta_1^{(bc^*)} - \vartheta_3^{(bc^*)}| \\ \frac{1}{2}|\vartheta_1^{(bc^*)} - \vartheta_3^{(bc^*)}| & Z_2^{(bc^*)} \end{pmatrix}, \tag{4.11}$$

$$\mathbf{m^2} = M_s^2 \begin{pmatrix} Z_2^{(bc)} & \frac{1}{2}|\vartheta_1^{(bc)} - \vartheta_3^{(bc)}| \\ \frac{1}{2}|\vartheta_1^{(bc)} - \vartheta_3^{(bc)}| & Z_2^{(bc)} \end{pmatrix} \tag{4.12}$$

The fields H_i and h_i are thus defined as

$$H^{\pm} = \frac{1}{2}(H_1^* \pm H_2); h^{\pm} = \frac{1}{2}(h_1^* \pm h_2)\,. \tag{4.13}$$

As a result the effective potential which corresponds to the spectrum of Higgs scalars is given by

$$\begin{aligned} V_{\text{Higgs}} = {} & m_H^2\left(|H_1|^2 + |H_2|^2\right) + m_h^2\left(|h_1|^2 + |h_2|^2\right) + \\ & + m_B^2 H_1 H_2 + m_B^2 h_1 h_2 + \text{h.c.}\,, \end{aligned} \tag{4.14}$$

where

$$\begin{aligned} m_h^2 &= \frac{Z_2^{(bc)}}{4\pi^2\alpha'}\,; & m_H^2 &= \frac{Z_2^{(bc^*)}}{4\pi^2\alpha'} \\ m_b^2 &= \frac{1}{2\alpha'}\left|\vartheta_1^{(bc)} - \vartheta_3^{(bc)}\right|; & m_B^2 &= \frac{1}{2\alpha'}\left|\vartheta_1^{(bc^*)} - \vartheta_3^{(bc^*)}\right| \end{aligned} \tag{4.15}$$

We note that the Z_2 is a free parameter, a moduli, and can become very small in relation to the Planck scale. However, the m_B^2 mass can be expressed in terms of the scalar masses of the particles present at the different intersections. Going one step further, we can express the "angle" part of the Higgs masses in terms of the angles defined in (4.4) and in figure 1. Explicitly, we find:[15]

$$\begin{aligned} m_B^2 &= \frac{1}{2\alpha'}|\tilde{\vartheta}_1 - \vartheta_1|\,; & m_b^2 &= \frac{1}{2\alpha'}\left|\tilde{\vartheta}_1 + \vartheta_1 + \vartheta_3 - \frac{1}{2}\right| \\ m_h^2 &= \frac{1}{2\alpha'}(\chi_b - \chi_c)^2\,; & m_H^2 &= \frac{1}{2\alpha'}(\chi_b + \chi_c)^2, \end{aligned} \tag{4.16}$$

where χ_b, χ_c the distances from the orientifold plane of the branes b, c. Making use of the scalar mass relations at the intersections of the model we can reexpress the mass relations (4.16), in terms of (4.6). The values of m_B^2, m_b^2 are given in appendix A.

5. SUSY at intersections and intermediate scale

In the present classes of models the U(1) symmetry gets broken and the associated gauge boson receives a mass, as there is a non-zero coupling of (3.14) to RR two form field B_2^3. However, for special values of the RR tadpole solution parameters is is possible that the usual mechanism, of giving a vev to a scalar, contributes additionally to the mass of the massive gauge boson associated with (3.14).

Thus our aim in this section is to provide us with a complementary mechanism to break the additional generator $\mathrm{U}(1)^N$. That may happen by demanding that the sector ce preserves $N = 1$ SUSY. That will have as an effect the appearance of I_{ce} massless scalars in the intersection with the same quantum numbers as the massless I_{ce} fermions. Because $I_{ce} = 1$, and the massless fermion localized in the intersection is ν_R, the massive partner of ν_R which become massless will be a $s\nu_R$. Consequently, by giving a vev to $s\nu_R$, the $s\nu_R$ gets

[15] We have chosen a configuration with $\epsilon = \tilde{\epsilon} = 1$, $n_b, n_c, n_d, n_e > 0$.

Brane	θ_a^1	θ_a^2	θ_a^3
c	$tan^{-1}\alpha$	0	$\frac{\pi}{2}$
e	0	$\frac{\pi}{2}$	$\tan^{-1}(\alpha)$

Table 4: Angle content for branes participating in the gauge symmetry breaking to the SM. Imposing $\mathcal{N}=1$ SUSY on the open sector ce breaks the surplus $\mathrm{U}(1)^N$ by a vev of $s\nu_R$.

charged and thus breaks $\mathrm{U}(1)^N$, leaving only the SM gauge group $\mathrm{SU}(3)\otimes\mathrm{SU}(2)\otimes\mathrm{U}(1)_Y$ at low energies. Lets us describe the procedure in more detail.

We want the particles localized on the intersection ce to respect some amount of SUSY, in our case $N=1$. That means that the relative angle between branes c, e, should obey the SUSY preserving condition

$$\pm\tilde{\theta}_1 \pm \bar{\theta}_2 \pm \left(\frac{\pi}{2} + \theta_3\right) = 0 \tag{5.1}$$

In this case, a massless scalar field appear in the intersection ce, the superpartner of the ν_R, the $s\nu_R$ field. It is charged under the additional $\mathrm{U}(1)^N$ symmetry, thus breaks $\mathrm{U}(1)^N$ by receiving a vev. In this case the surviving gauge symmetry is of SM. The scale of the additional breaking, M_N, is set from the vev of $s\nu_R$ and in principle, can be anywhere between M_Z and the string scale.

The following choice:

$$\tan^{-1}\frac{\beta^1 U^{(1)}}{n_c^1} + \tan^{-1}\frac{\beta^2 U^{(2)}}{n_e^2} - \tan^{-1}\left(\frac{U^{(3)}}{6}\right) - \frac{\pi}{2} = 0\,, \tag{5.2}$$

with

$$n_e^2 = 0\,, \qquad \frac{\beta^1 U^{(1)}}{n_c^1} = \frac{U^{(3)}}{6}\,, \qquad \alpha = \frac{\beta^1 U^{(1)}}{n_c^1}\,, \qquad U^{(i)} = \frac{R_2^{(i)}}{R_1^{(i)}}\,. \tag{5.3}$$

solves (5.1). In particular,

$$n_e^2 = 0 \Rightarrow \beta^2 = 1 \tag{5.4}$$

thus the second tori is not tilted. The angle content of the branes, c, e, when the gauge symmetry breaks to the SM is given in table 4.

Summarizing the SM exists between M_Z and the mass scale M_N of the additional $\mathrm{U}(1)^N$. A set of SM wrappings exists only if we consider the *hypercharge* (3.13) and the *gauge symmetry breaking* condition (5.3) when defining numerically the tadpole solutions of table 2. Taking into account both conditions a consistent set, for the observable SM to exist, wrapping numbers is given by

$$n_e^2 = 0\,, \qquad \beta^2 = 1\,, \qquad \beta^1 = \frac{1}{2}\,, \qquad n_b = -1\,, \qquad n_d^2 = -1\,, \qquad n_a^2 = 2\,, \qquad n_c^1 = 1 \tag{5.5}$$

or

$$N_a = 3 \quad (2,0)(2,1)\left(3,\frac{1}{2}\right)$$

Higgs fields	β_1	β_2	n_b	n_c	n_a^2	N_D
$n_H = 1,\ n_h = 0$	$1/2$	1	1	-1	$-1 - n_d^2$	$8 + 4n_d^2$
$n_H = 1,\ n_h = 0$	$1/2$	1	-1	1	$1 - n_d^2$	$4n_d^2$
$n_H = 0,\ n_h = 1$	$1/2$	1	1	1	$1 - n_d^2$	$-1 + 4n_d^2$
$n_h = 0,\ n_h = 1$	$1/2$	1	-1	-1	$-1 - n_d^2$	$9 + 4n_d^2$
$n_H = 1,\ n_h = 1$	1	1	1	0	$-n_d^2$	$7 + 4n_d^2$
$n_H = 1,\ n_h = 1$	1	1	-1	0	$-n_d^2$	$9 + 4n_d^2$
$n_H = 1,\ n_h = 1$	1	1	0	1	$2 - n_d^2$	$-1 + 4n_d^2$
$n_H = 1,\ n_h = 1$	1	1	0	-1	$-2 - n_d^2$	$16 + 4n_d^2$

Table 5: Families of models with with minimal Higgs structure. They depend on a single integer, n_d. The surplus *gauge symmetry breaking* condition (5.4) has been taken into account.

$$\begin{aligned}
N_b &= 2 \quad \left(-1, -\frac{1}{2}\right)(1,0)\left(1, \frac{1}{2}\right) \\
N_c &= 1 \quad \left(1, \frac{1}{2}\right)(1,0)(0,1) \\
N_d &= 1 \quad (2,0)(-1,2)\left(1, -\frac{1}{2}\right) \\
N_e &= 1 \quad (2,0)(0,1)\left(1, -\frac{1}{2}\right).
\end{aligned} \tag{5.6}$$

It satisfies all tadpole conditions but the first, the latter is satisfied with the addition of four $\bar{D}_6$ located at $(2,0)(1,0)(2,0)$.

The number of electroweak Higgs present in the model can be investigated further. The most interesting cases that have a minimal Higgs content follow:

- *The Higgs system of MSSM*

  For (4.7), we can see that the minimal set of Higgs in the models is obtained for either $n_H = 0$, $n_h = 1$, or $n_h = 1$, $n_H = 0$. The two cases, are studied in table 5. We found two families of models that depend on a single integer n_d^2. We also list the number of necessary N_D branes required to cancel the first tadpole condition. We have taken into account the conditions (3.13), (5.4) necessary to obtain the observable SM at low energies.

  The case $n_H = 1$, $n_h = 0$ appears to be the most interesting as this appears to give a plausible explanation for the existence of small and different neutrino masses to the different generations. These issues are examined in more detail in the next section.

- *Double Higgs system*

  The next to minimal set of Higgses is obtained when $n_H = 1,\ n_h = 1$. In this case, quarks and leptons get their mass from the start.

6. Neutrino couplings and masses

The Yukawa couplings in this model follow the usual pattern that appears in intersecting branes [10]. The couplings between the two fermion states F_L^i, $\bar{F}_R^j$ and the Higgs fields H^k, arise from the stretching of the worldsheet between the three D6-branes which cross at those intersections. For a six dimensional torus they can take the following form in the leading order [10],

$$Y^{klm} = e^{-\tilde{A}_{klm}} , \tag{6.1}$$

where $\tilde{A}_{klm}$ is the worldsheet area connecting the three vertices in the six dimensional space. The areas of each of the two dimensional torus involved in this interaction is typically of order one in string units. For the models discussed in table 1, the Yukawa interactions for the chiral spectrum of the SM's yield:

$$\begin{aligned}
&Y_j^U Q_L U_R^j h_1 + Y_j^D Q_L D_R^j H_2 + \\
&+ Y_{ij}^u q_L^i U_R^j H_1 + Y_{ij}^d q_L^i D_R^j h_2 + \\
&+ Y_h^l \, l_L^h \, \nu_R^h \, h_1 + Y_h^e l_L^h e_R^h H_2 + \\
&+ Y_{ij}^N L^i N_R^j h_1 + Y_{ij}^E L_i E_R^j H_2 + \text{h.c.}
\end{aligned} \tag{6.2}$$

where $i = 1, 2$, $j = 1, 2, 3$, $h = 1$.

The nature of Yukawa couplings is such that the lepton and neutrino sector of the models distinguish between different generations, e.g. the "first" from the other two generations, as one generation of neutrinos (resp. leptons) is placed on a different intersection from the other two one's. For example looking at the charged leptons of table 1 we see that one generation of charged leptons l_L gets localized on *be*-intersection, while the other two generations L get localized in the *bd*-intersection.

There are a number of scalar doublets in the model present that are interpreted in terms of the low energy theory as Higgs doublets. The most interesting ones were mentioned briefly in the last section. There are two possibilities to be discussed. The minimal case when $n_H = 1$, $n_h = 0$ and the next to minimal case, $n_H = 1$, $n_h = 1$.

- *Minimal Higgs presence*

  Without loss of generality we choose $n_H = 1$, $n_h = 0$, only the H_1, H_2 fields are present. The mechanism that will give masses to the charged lepton/quark sector is similar to what happended in the four stack models of [3]. At tree level two U-quarks and one D-quark as well the charge leptons gain masses. Identifying the massive quarks with t, c, b, the rest of the quarks as well as neutrinos remain massless at tree level. The rest of the quarks are expected to receive masses from strong interaction effects, that create effective couplings of the form $Q_L U_R^j H_1$, $q_L^i D_R^j H_2$ creating masses for the u-, d-, s-quarks of less than equal of Λ_{QCD}. Note that the latter couplings are not allowed at tree level since otherwise the $\mathrm{U}(1)_b$ global symmetry would have been violated.

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The neutrino sector is slightly different from the four stack counterparts [3] of the SM's discussed in the present work. Contrary to the models of [3] where all neutrinos come from the same intersection, in this work the SM's are build such that one generation of neutrinos (and leptons) are placed at a different intersection from the remaining two generations.[16] Thus the structure of Yukawa couplings in the neutrino sector suggests that they is a distinction between the neutrinos of the one (e.g. first) generation and the other two one's. That could be used in principle to discuss neutrino mass textures in the context of recent results[17] of SuperKamiokande [24], which suggest that at least two generations of neutrinos may be massive.

Because Lepton number is a gauged symmetry, there are only Dirac masses allowed. Thus a see-saw mechanism is excluded. The origin of small neutrino masses originates from dimension 6 operators in the form, breaking the $U(1)_b$ PQ like symmetry through chiral symmetry breaking,

$$\alpha'(LN_R)\ (Q_L U_R)^{\star}\,, \qquad \alpha'(l\nu_R)(q_L U_R)\,. \tag{6.3}$$

Hence, the smallness of neutrino masses is related to the existence of the dominant u-quark chiral condensate, $\langle u_R u_R \rangle$. For values[18] of the u-chiral condensate, and assuming that all generations of neutrino species receive the same value for u-condensate, of order $(240\,\mathrm{MeV})^3$, the neutrinos get a mass of order

$$\frac{< u_R u_L >}{M_s^2} = \frac{(240\,\mathrm{MeV})^3}{M_s^2} \tag{6.4}$$

Hence neutrino masses of order 0.1–10 eV are easily achieved in consistency with LSND oscillation experiments [17]. In this case, a generation mixing among neutrino species is generated from the leading order coupling behavior (6.1).

Alternatively, if we assume that the chiral condensate generates the generation mixing, receiving different values for different neutrino species the neutrino masses will depend weakly on the precise form of the couplings (6.1).

- *Next to minimal Higgs presence*

  In this case, all couplings to quarks and leptons are realized from the start. All particles get a mass. The hierarchy of masses depend both on the Higgs fields and the leading order Yukawa behavior (6.1).

7. Conclusions and future directions

In this work, we have presented the first examples of string models, not based on a GUT group at the string scale, that have at low energies only the standard model and are derived from five stacks of (D6) branes at the string scale.[19] We note the following:

[16] The latter though are placed both at the same intersection.

[17] See for example ref. [23].

[18] See for example ref. [25].

[19] The first examples of models giving rise to the SM at low energies have been considered in [3], using four stacks of branes and the same type I orientifolded T^6 constructions.

- Baryon number is a gauged symmetry and proton is stable. If proton was not stable in the models then we should have push the string scale higher than 10^{16} GeV to suppress dimension six operators that could potentially contribute to proton decay. However, in the present class of models, this is not necessary as proton is stable.

- The additional $U(1)^N$ symmetry may break from its non-zero coupling to the RR two form fields.

  However, as we have noted an additional mechanism is available for certain values of the tadpole and complex structure parameters, and thus complementary to the abelian gauge field mass term generation induced by their non-zero couplings to the RR two form fields B_2^i, $i = 1, 2, 3$. Crucial for this new mechanism, thus contributing to the breaking of the additional anomaly free U(1) generator and getting only the SM at low energies, was the novel partial imposition of $\mathcal{N} = 1$ supersymmetry at only one intersection, e.g. *ce*. That had as an immediate effect to pull out from the massive modes the superpartner of ν_R. It would be interesting to investigate in detail the symmetry breaking patterns that follow from having a SUSY intersection, at an open string sector, of the non-SUSY SM's examined in this work.

- We emphasize that the breaking of the $U(1)^N$ symmetry implies the existence of an extra Z^o boson above the electroweak scale. Bounds on additional gauge bosons exist [26] placing them in the range between 500-800 GeV. Thus we conclude that the string scale should be at least equal to M_N or higher. Improved bounds of the string scale for the present models would require a generalization of the four stack D6 model [3] analysis of [27], for the masses of the extra U(1) gauge bosons made massive by the Green-Schwarz mechanism. It would be interesting to extend the analysis of [27] to the present SM's that predict an additional intermediate scale between M_Z and M_S.

- A natural extension of the five-stack D6-brane SM's of this work is to examine how we can construct D6-brane models that respect some supersymmetry at every intersection as was detailed[20] recently.

- The models have vanishing RR tadpoles but some NSNS tadpoles remain, leaving a open issue the full stability of the configurations. It is then an open question if the backgrounds can be cured using Fischler-Susskind mechanism [28] in redefining the background [29] as in [30].

Concluding this work, it is very interesting that the present class of models predicts not only the existence of a non-supersymmetric standard model at low energies but in addition other classes of models predicting the unique existence of a SUSY partner of the right handed neutrino, the $s\nu_R$.

[20]See the first two references of [13] for a similar construction for the 4-stack SM's of [3].

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A. Higgs mass parameters

In this appendix, we list the values of the mass parameters, of section 4, involved in the mass of the set of four Higgses taking part in the process of electroweak symmetry breaking. As we remark in the main body of the paper, the quadratic parts of the Higgs mixing mass terms in the effective field theory potential are exactly calculable at tree level in the D6-brane models.

$$
\begin{aligned}
m_B^2 &= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)-m^2_{U_R}(t_2)-m^2_{U_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)-m^2_{D_R}(t_2)-m^2_{D_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)-m^2_{N_R}(t_2)-m^2_{N_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)-m^2_{U_R}(t_2)-m^2_{U_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)-m^2_{D_R}(t_2)-m^2_{D_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)-m^2_{N_R}(t_2)-m^2_{N_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{L}(t_2)+m^2_{L}(t_3)-m^2_{U_R}(t_2)-m^2_{U_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{L}(t_2)+m^2_{L}(t_3)-m^2_{D_R}(t_2)-m^2_{D_R}(t_3)\Big| \\
&= \frac{1}{2}\Big|m^2_{L}(t_2)+m^2_{L}(t_3)-m^2_{N_R}(t_2)-m^2_{N_R}(t_3)\Big| \qquad \text{(A.1)} \\
m_b^2 &= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)+m^2_{U_R}(t_2)+m^2_{U_R}(t_3)-m^2_{e_R}(t_1)-m^2_{e_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)+m^2_{U_R}(t_2)+m^2_{U_R}(t_3)-m^2_{E_R}(t_1)-m^2_{E_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)+m^2_{N_R}(t_2)+m^2_{N_R}(t_3)-m^2_{e_R}(t_1)-m^2_{e_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{Q_L}(t_2)+m^2_{Q_L}(t_3)+m^2_{N_R}(t_2)+m^2_{N_R}(t_3)-m^2_{E_R}(t_1)-m^2_{E_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)+m^2_{U_R}(t_2)+m^2_{U_R}(t_3)-m^2_{e_R}(t_1)-m^2_{e_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)+m^2_{U_R}(t_2)+m^2_{U_R}(t_3)-m^2_{E_R}(t_1)-m^2_{E_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)+m^2_{D_R}(t_2)+m^2_{D_R}(t_3)-m^2_{e_R}(t_1)-m^2_{e_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)+m^2_{D_R}(t_2)+m^2_{D_R}(t_3)-m^2_{E_R}(t_1)-m^2_{E_R}(t_2)\Big| \\
&= \frac{1}{2}\Big|m^2_{q_L}(t_2)+m^2_{q_L}(t_3)+m^2_{N_R}(t_2)+m^2_{N_R}(t_3)-m^2_{e_R}(t_1)-m^2_{e_R}(t_2)\Big|
\end{aligned}
$$

$$
\begin{aligned}
&= \frac{1}{2}\Big| m_{q_L}^2(t_2) + m_{q_L}^2(t_3) + m_{N_R}^2(t_2) + m_{N_R}^2(t_3) - m_{E_R}^2(t_1) - m_{E_R}^2(t_2)\Big| \\
&= \frac{1}{2}\Big| m_L^2(t_2) + m_L^2(t_3) + m_{U_R}^2(t_2) + m_{U_R}^2(t_3) - m_{e_R}^2(t_1) - m_{e_R}^2(t_2)\Big| \\
&= \frac{1}{2}\Big| m_L^2(t_2) + m_L^2(t_3) + m_{U_R}^2(t_2) + m_{U_R}^2(t_3) - m_{E_R}^2(t_1) - m_{E_R}^2(t_2)\Big| \\
&= \frac{1}{2}\Big| m_L^2(t_2) + m_L^2(t_3) + m_{D_R}^2(t_2) + m_{D_R}^2(t_3) - m_{e_R}^2(t_1) - m_{e_R}^2(t_2)\Big| \\
&= \frac{1}{2}\Big| m_L^2(t_2) + m_L^2(t_3) + m_{D_R}^2(t_2) + m_{D_R}^2(t_3) - m_{E_R}^2(t_1) - m_{E_R}^2(t_2)\Big| \\
&= \frac{1}{2}\Big| m_L^2(t_2) + m_L^2(t_3) + m_{N_R}^2(t_2) + m_{N_R}^2(t_3) - m_{e_R}^2(t_1) - m_{e_R}^2(t_2)\Big| \\
&= \frac{1}{2}\Big| m_L^2(t_2) + m_L^2(t_3) + m_{N_R}^2(t_2) + m_{N_R}^2(t_3) - m_{E_R}^2(t_1) - m_{E_R}^2(t_2)\Big| \qquad \text{(A.2)}
\end{aligned}
$$

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