Deformed Intersecting D6-Brane GUTS
and N=1 SUSY

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ABSTRACT

We analyze the construction of non-supersymmetric three generation six-stack Pati-Salam (PS) $SU(4)_C \times SU(2)_L \times SU(2)_R$ GUT classes of models, by localizing D6-branes intersecting at angles in four dimensional orientifolded toroidal compactifications of type IIA. Special role in the models is played by the presence of extra branes needed to satisfy the RR tadpole cancellation conditions. The models contain at low energy exactly the Standard model with no extra matter and/or extra gauge group factors. They are build such that they represent deformations around the quark and lepton basic intersection number structure. The models possess the same phenomenological characteristics of some recently discussed examples (PS-A, PS-I; PS-II GUT classes; hep-th/0203187, hep-th/0209202, hep-th/0210004) of four and five stack PS GUTS respectively. Namely, there are no colour triplet couplings to mediate proton decay and proton is stable as baryon number is a gauged symmetry. The mass relation $m_e = m_d$ at the GUT scale is recovered. Even though more complicated, than in lower stack GUTS, the conditions of the non-anomalous U(1)’s to survive massless the generalized Green-Schwarz mechanism are solved consistently by the angle conditions coming from the presence of N=1 supersymmetric sectors involving the presence of extra branes and also required for the existence of a Majorana mass term for the right handed neutrinos.

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1 Introduction

In this work we present four dimensional (4D) three generation six stack Pati-Salam (PS) like type of GUT classes of models that break at low energy exactly to the Standard Model (SM), $SU(3)_C \times SU(2)_L \times U(1)_Y$, without any extra chiral fermions and/or extra gauge group factors (in the form of hidden sectors). These constructions are non-supersymmetric and are centered around the PS $SU(4)_C \times SU(2)_L \times SU(2)_R$ gauge group. These constructions have D6-branes intersecting each other at angles, in orientifolded compactifications of IIA theory in a factorized six-tori, with the O$_6$ orientifold planes on top of D6-branes [1, 2].

The new classes of models have some characteristic features that include:

- The presented PS-III classes of models, are constructed with an initial gauge group $U(4) \times U(2) \times U(2) \times U(1)^3$ at the string scale. At the scale of symmetry breaking of the left-right symmetry $M_{GUT}$, the initial symmetry group breaks to the standard model $SU(3)_C \times SU(2)_L \times U(1)_Y$ augmented with two extra anomaly free $U(1)$ symmetries. The additional $U(1)$'s may break by gauge singlets generated by imposing N=1 SUSY in particular sectors.

- A number of extra $U(1)$'s added to cancel the RR tadpoles breaks by gauge singlets generated in imposing N=1 SUSY on sectors involving the extra $U(1)$'s.

- Neutrinos get a mass of the right order $^1$ from a see-saw mechanism of the Frogatt-Nielsen type.

- Proton stability is guaranteed due to the fact that baryon number is an unbroken gauged global symmetry surviving at low energies and no colour triplet couplings that could mediate proton decay exist. A gauged baryon number is a general feature in D6-brane models [3, 4, 5, 6, 7, 8].

In the major problems of string theory the hierarchy of scale and particle masses after supersymmetry breaking is included. These issues have been explored by explicitly constructing semirealistic supersymmetric models of $^2$ $N=1$ orbifold compactifications of the heterotic string theories. One of the unsolved problems is that the string scale which is of order $10^{18}$ GeV is in clear disagreement with the ‘observed’ unification

$^1$In consistency with LSND oscillation experiments

$^2$weakly coupled
of gauge coupling constants in the MSSM of $10^{16}$ GeV. The latter problem was not
eventually solved even if the discrepancy between the two high scales was attributed
e.g. to the presence of the 1-loop string threshold corrections to the gauge coupling
countants \[11\].

On the opposite side, in type I models, the string scale is a free parameter and thus
may be lowered in the TeV region \[12\] suggesting that non-SUSY models with a TeV
string scale is a possibility. In this content new constructions have appeared in type
I string vacuum background which construct even generation four dimensional non-
supersymmetric models using intersecting branes. The new constructions were made
possible by turning on background fluxes on D9 branes on a type I background \[3\]. Thus
open string models were constructed that break supersymmetry on the brane and give
chiral fermions with an even number of generations \[11\]. In these models the fermions
get localized in the intersections between branes \[13\]. The introduction of a quantized
background NS-NS B field \[14, 15, 16\], that makes the tori tilted, subsequently give rise
to semi-realistic models with three generations \[2\]. The latter backgrounds are T-dual to
models with magnetic deformations \[17\]. In \[8\] the first examples of four D6-brane stack
models which have only the SM at low energy, in the language of D6-branes intersecting
at angles on an orientifolde $T^6$ torus, were constructed. In this construction, as in all
models from the same backgrounds, the proton is stable since the baryon number is a
gauged $U(1)$ global symmetry. A special feature of these models is that the neutrinos
can only get Dirac mass. These models have been generalized to models with five
stacks \[4\] and six stacks of D6-branes at the string scale \[7\]. The models appearing in
\[4\], \[7\] are build as novel deformations of the QCD intersection numbers, namely they
are build around the left and right handed quarks intersection numbers. They hold
exactly the same phenomenological properties of \[3\]. Its is important to stress here that
contrary to their four stack counterparts \[3\] they have unique special features, since
by demanding the presence of $N = 1$ supersymmetric sectors, we are able to break the
extra, beyond the SM gauge group, $U(1)$’s, and thus predicting the unique existence
of one supersymmetric partner of the right neutrino or two supersymmetric partners
of the right neutrinos in the five and the six stack SM’s respectively.

In a parall development, in \[6\] we presented the first examples of GUT models in a
string theory context, and in the context of intersecting branes, that break completely
to the SM at low energies. These models predict the unique existence of light weak

\[3\]In the T-dual backgrounds these constructions are represented by D6 branes wrapping 3-cycles
on a dual six dimensional torus and intersecting each other at certain angles.
fermion doublets with energy between $M_Z - 246$ GeV, thus can be directly tested at present of future accelerators. The constructions of four D6-brane stack of [6] were studied further in [7] in relation to other four stack deformations, and by extending these constructions to five stacks of D6-brane stacks in [8]. In this work we will discuss the six extensions of [6, 7, 8].

We note that apart from D6 models with exactly the SM at low energy just mentioned, there intersecting D5-branes have been studied with only the SM at low energy [9, 10]. In the latter models [10] there are special classes of theories, again appearing as novel deformations of the QCD intersection number structure, which have not only the SM at low energy but exactly the same low energy effective theory including fermion and scalar spectrum.

For additional work, in the context of intersecting branes see [19, 20, 21, 22, 23, 24, 25, 26, 27] 4.

The paper is organized as follows: In section two we describe the general spectra and tadpole rules for building chiral GUT models in orientifolded $T^6$ compactifications. In section 3, we discuss the basic fermion and scalar structure of the present PS-III class of models. In section 4, we discuss the parametric solutions to the RR tadpole cancellation conditions as well giving the general characteristics of the PS-III GUTS. In section 5 we analyze the cancellation of $U(1)$ anomalies in the presence of a generalized Green-Schwarz (GS) mechanism and extra $U(1)$ branes. In section 6, we discuss the conditions for the absence of tachyons, the angle structure between the branes and its role in describing the Higgs sector of the models. In subsection 7.1 we discuss the importance of creating sectors preserving $N=1$ SUSY for the realization of the see-saw mechanism and its relation to the generalized Green-Schwarz mechanism. In subsection 7.2 we discuss the role of extra $U(1)$ branes in creating scalar singlets. In subsection 7.3 we analyze the breaking of the surviving the Green-Schwarz mechanism massless $U(1)$’s. In subsection 8.1 we analyze the structure of GUT Yukawa couplings in intersecting braneworlds and discuss the problem of neutrino masses. In subsection 8.2 we exhibit that all additional exotic fermions beyond those of SM present in the models may become massive and disappear from the low energy $M_Z$ spectrum. We present our conclusions in Section 9. Appendix I, includes the conditions for the absence of tachyonic modes in the spectrum of the PS-III class of models discussed in this work.

4For some other constructions close to the SM but not based on a particular string constructions see [30].
2 Tadpole structure and spectrum rules

Let us now describe the construction of the PS classes of models. 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 \cite{18, 1, 2}. By performing a T-duality transformation on the $x^4$, $x^6$, $x^8$ directions the D9-branes with fluxes are translated into D6-branes intersecting at angles. We assume that the D6-$\alpha$-branes are wrapping 1-cycles $(n^i_\alpha, m^i_\alpha)$ along each of the $T^2$ torus of the factorized $T^6$ torus, namely $T^6 = T^2 \times T^2 \times T^2$.

In order to build a PS model with a minimal Higgs structure we consider six stacks of D6-branes giving rise to their world-volume to an initial gauge group $U(4) \times U(2) \times U(2) \times U(1) \times U(1) \times U(1)$ gauge symmetry at the string scale. In addition, we consider the addition of NS B-flux, such that the tori are not orthogonal, avoiding in this way an even number of families, and leading to effective tilted wrapping numbers,

$$(n^i, m = \tilde{m}^i + n^i/2); \quad n, \tilde{m} \in Z, \quad (2.1)$$

that allows semi-integer values for the m-numbers.

In the presence of $\Omega R$ symmetry, where $\Omega$ is the worldvolume parity and $R$ is the reflection on the T-dualized coordinates,

$$T(\Omega)T^{-1} = \Omega R, \quad (2.2)$$

and thus each D6-$\alpha$-brane 1-cycle, must have its $\Omega R$ partner $(n^i_\alpha, -m^i_\alpha)$.

Chiral fermions are obtained by stretched open strings between intersecting D6-branes \cite{13}. Also the chiral spectrum of the models may be obtained after solving simultaneously the intersection constraints coming from the existence of the different sectors taking into account the RR tadpole cancellation conditions.

There are a number of different sectors, contributing to the chiral spectrum. Denoting the action of $\Omega R$ on a sector $\alpha, \beta$, by $\alpha^*, \beta^*$, respectively the possible sectors are:

- The $\alpha\beta + \beta\alpha$ sector: involves open strings stretching between the D6-$\alpha$ and D6-$\beta$ branes. Under the $\Omega R$ symmetry this sector is mapped to its image, $\alpha^*\beta^* + \beta^*\alpha^*$ sector. The number, $I_{\alpha\beta}$, of chiral fermions in this sector, transforms in the bifundamental representation $(N_\alpha, \bar{N}_\alpha)$ of $U(N_\alpha) \times U(N_\beta)$, being

$$I_{\alpha\beta} = (n^1_\alpha m^1_\beta - m^1_\alpha n^1_\beta)(n^2_\alpha m^2_\beta - m^2_\alpha n^2_\beta)(n^3_\alpha m^3_\beta - m^3_\alpha n^3_\beta), \quad (2.3)$$

•
where $I_{\alpha\beta}$ is the intersection number of the wrapped three cycles. Note that with the sign of $I_{\alpha\beta}$ we denote the chirality of the fermions, $I_{\alpha\beta} > 0$ being those of left handed fermions. Negative multiplicity correspond to opposite chirality.

- The $\alpha\alpha$ sector: it involves open strings stretching on a single stack of D6$_\alpha$ branes. Under the $\Omega R$ symmetry this sector is mapped to its image $\alpha^\star\alpha^\star$. This sector contain $\mathcal{N} = 4$ super Yang-Mills and if it exists SO(N), SP(N) groups in principle may appear. This sector is of no importance to us as we will be dealing with unitary groups.

- The $\alpha\beta^\star + \beta\alpha^\star$ sector: In this sector which under the $\Omega R$ symmetry transforms to itself, chiral fermions transform into the $(N_\alpha, N_\beta)$ representation with multiplicity given by

$$I_{\alpha\beta^\star} = -(n_1^a m_1^a + m_1^a n_1^a)(n_2^a m_2^a + m_2^a n_2^a)(n_3^a m_3^a + m_3^a n_3^a).$$

(2.4)

• the $\alpha\alpha^\star$ sector: under the $\Omega R$ symmetry is transformed to itself. In this sector the invariant intersections give $8m_1^a m_2^a m_3^a$ fermions in the antisymmetric representation and the non-invariant intersections that come in pairs provide $4m_1^a m_2^a m_3^a(n_1^a n_2^a n_3^a - 1)$ additional fermions in the symmetric and antisymmetric representation of the $U(N_\alpha)$ gauge group.

Additionally any vacuum derived from the previous intersection number constraints of the chiral spectrum is subject to constraints coming from RR tadpole cancellation conditions $[\Pi]$. That is equivalent to cancellation of D6-branes charges$^5$, wrapping on three cycles with homology $[\Pi_a]$ and O6-plane 7-form charges wrapping on 3-cycles with homology $[\Pi_{O6}]$. Explicitly, the RR tadpole cancellation conditions expressed in terms of cancellations of RR charges in homology, obey:

$$\sum_a N_a[\Pi_a] + \sum_{\alpha^\star} N_{\alpha^\star}[\Pi_{\alpha^\star}] - 32[\Pi_{O6}] = 0.$$  

(2.5)

or

$$\sum_a N_a n_1^a n_2^a n_3^a = 16,$$

$$\sum_a N_a m_1^a m_2^a m_3^a = 0,$$

$$\sum_a N_a m_1^a n_2^a m_3^a = 0,$$

$$\sum_a N_a n_1^a m_2^a m_3^a = 0.$$  

(2.6)

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

5
That forces absence of non-abelian gauge anomalies.

A six-stack brane configuration with minimal PS particle content may be obtained by considering six stacks of branes yielding an initial \( U(4)_a \times U(2)_b \times U(2)_c \times U(1)_d \times U(1)_e \times U(1)_f \). In this case the equivalent gauge group is an \( SU(4)_a \times SU(2)_b \times SU(2)_b \times U(1)_a \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e \times U(1)_f \). Thus, in the first instance, we identify, without loss of generality, \( SU(4)_a \) as the \( SU(4)_c \) colour group that its breaking induces the \( SU(3) \) colour group of strong interactions, the \( SU(2)_b \) with \( SU(2)_L \) of weak interactions and \( SU(2)_c \) with the \( SU(2)_R \) of left-right symmetric PS models.

### 3 The basic fermion structure

In this section we will describe the basic characteristics of the GUT models that we will analyze in this work. The models are three family non-supersymmetric GUT model with the left-right symmetric PS model structure \[SU(4)_C \times SU(2)_L \times SU(2)_R\]. The open string background from which the models will be derived are intersecting D6-branes wrapping on 3-cycles of decomposable toroidal \((T^6)\) orientifolds of type IIA in four dimensions \[1, 2\].

The three generations of quark and lepton fields are accommodated into the following representations :

\[
F_L = (4, 2, 1) = q(3, \bar{2}, \frac{1}{6}) + l(1, \bar{2}, -\frac{1}{2}) \equiv (u, d, l),
\]

\[
F_R = (\bar{4}, 1, 2) = u^c(\bar{3}, 1, -\frac{2}{3}) + d^c(\bar{3}, 1, \frac{1}{3}) + e^c(1, 1, 1) + N^c(1, 1, 0) \equiv (u^c, d^c, l^c),
\]

where the quantum numbers on the right hand side of (3.1) are with respect to the decomposition of the \( SU(4)_C \times SU(2)_L \times SU(2)_R \) under the \( SU(3)_C \times SU(2)_L \times U(1)_Y \) gauge group and \( l = (\nu, e) \) is the standard left handed lepton doublet while \( l^c = (N^c, e^c) \) are the right handed leptons. The assignment of the accommodation of the quarks and leptons into the representations \( F_L + F_R \) is the one appearing in the spinorial decomposition of the 16 representation of \( SO(10) \) under the PS gauge group.

Also present are the fermions

\[
\chi_L = (1, \bar{2}, 1), \quad \chi_R = (1, 1, \bar{2}).
\]

These fermions are a general prediction of left-right symmetric theories as their existence follows from RR tadpole cancellation conditions.

The symmetry breaking of the left-right PS symmetry at the \( M_{GUT} \) scale, in principle
as high as the string scale, proceeds through the representations of the set of Higgs fields,

\[ H_1 = (\bar{4}, 1, 2), \quad H_2 = (4, 1, 2), \quad (3.3) \]

where,

\[ H_1 = (\bar{4}, 1, 2) = u_H(3, 1, \frac{2}{3}) + d_H(\bar{3}, 1, -\frac{1}{3}) + e_H(1, 1, -1) + \nu_H(1, 1, 0). \quad (3.4) \]

The electroweak symmetry breaking is achieved through bi-doublet Higgs fields \( h_i \) where \( i = 3, 4 \) in the representations

\[ h_3 = (1, \bar{2}, 2), \quad h_4 = (1, 2, \bar{2}) . \quad (3.5) \]

Because of the imposition of N=1 SUSY on some open string sectors, there are also present the massless scalar superpartners of the quarks, leptons and antiparticles

\[ \tilde{F}_R^H = (\bar{4}, 1, 2) = u_H^c(\bar{3}, 1, -\frac{4}{3}) + d_H^c(3, 1, \frac{1}{3}) + e_H^c(1, 1, 1) + N_H^c(1, 1, 0) \equiv (u_H^c, d_H^c, l_H^c). \quad (3.6) \]

The latter fields \( ^6 \) characterize all vacua coming from these type IIA orientifolded tori constructions is the replication of massless fermion spectrum by an equal number of massive particles in the same representations and with the same quantum numbers.

Also, a number of charged exotic fermion fields, which may receive a string scale mass, appear

\[ 6(6, 1, 1), \quad 6(\bar{10}, 1, 1). \quad (3.7) \]

In addition the following gauge singlet fermion fields appear : \( s_1^R, s_2^R, s_3^R. \)

The complete accommodation of the fermion structure of the PS-III classes of models can be seen in table \([\Pi]\).

\section{4 Tadpole cancellation for PS-III classes of GUTS}

The following comments will be necessary to understand the analysis performed in the following sections of the PS-III classes of GUTS :

a) A proper formulation of a GUT model requires the realization of certain couplings necessary to e.g. for giving masses to right handed neutrinos or make massive non-observed particles. In intersecting brane worlds based on intersecting D6-branes this becomes much easier by demanding that some open string sectors preserve some supersymmetry. Thus some massive fields will be pulled out from the massive spectrum

\( ^6 \)are replicas of the fermion fields appearing in the intersection \( ac \) and they receive a vev
**Table 1:** Fermionic spectrum of the $SU(4)_C \times SU(2)_L \times SU(2)_R$, PS-III class of models together with $U(1)$ charges. We note that at energies of order $M_Z$ only the Standard model survives.

and become massless. Thus a Majorana mass term for the right handed neutrinos is realized if the sector $ac$ preserves $N = 1$ SUSY. As an immediate effect the previously massive $\bar{F}_R^H$ scalar appears.

b) The intersection numbers, in table (1), of the fermions $F_L + \bar{F}_R$ are chosen such that $I_{ac} = -3$, $I_{ab^*} = 3$. Here, $-3$ denotes opposite chirality to that of a left handed fermion. The choice of additional fermion representations $(1, \bar{2}, 1), (1, 1, \bar{2})$ is imposed on us by the RR tadpole cancellation conditions that are equivalent to $SU(N_a)$ gauge anomaly cancellation, in this case of $SU(2)_L, SU(2)_R$ gauge anomalies,

$$\sum_i I_{ia} N_a = 0, \ a = L, R. \quad (4.1)$$

c) In the present classes of models representations of scalar sextets $(6, 1, 1)$ fields, that appear in attempts to construct realistic 4D $N = 1$ PS heterotic models from the fermionic formulation [32], even through heterotic fermionic models where those repre-
sentations are lacking exist \cite{33}, do not appear. These representations were imposed in attempts to produce a realistic PS model in order to save the models from fast proton decay. Thus fast proton decay was avoided by making the mediating $d_H$ triplets of superheavy and of order of the $SU(2)_R$ breaking scale through their couplings to the sextets. In the PS-III models baryon number is a gauged global symmetry, thus proton is stable. Hence there is no need to introduce proton decay saving sextets. Moreover in the PS-III GUTS, there is no problem of having $d_H$ becoming light enough and inducing proton decay, as the only way this could happen, is through the existence of $d_H$ coupling of sextets to quarks and leptons. However, this coupling is forbidden by the symmetries of the models.

d) The mixed anomalies $A_{ij}$ of the six \footnote{We examine for convenience the case of two added extra U(1)'s.} surplus $U(1)$'s with the non-abelian gauge groups $SU(N_a)$ of the theory cancel through a generalized GS mechanism \cite{34} \cite{3}, involving close string modes couplings to worldsheet gauge fields. Crucial for the RR tadpole cancellation is the presence of $N_h$ extra branes. Contrary, of what was found in D6-brane models with exactly the SM at low energy, and a Standard-like structure at the string scale \cite{3} \cite{4} \cite{7} where the extra branes have no intersection with the branes \footnote{A similar phenomenon appears in intersecting D5-brane models on a $T^4 \times C/Z_N$, with exactly the SM at low energy and a Standard-like structure at the string scale \cite{9} \cite{10}.}, in the present PS-III GUT models there is a non-trivial intersection of the extra branes with the branes $a$, $b$, $c$. As a result, this becomes a new singlet generation mechanism after imposing $N = 1$ SUSY between $U(1)$ leptonic (the $d$, $e$, $f$ branes) and the $U(1)$ extra branes. Also, contrary to the SM's of \cite{3} \cite{4} \cite{5} \cite{9} \cite{10} the extra branes do not form a $U(N_h)$ gauge group but rather a $U(1)^{N_1} \times U(1)^{N_2} \cdots U(1)^{N_h}$ one, where $N_1 = N_2 = \cdots = N_h = 1$.

e) We don’t impose the constraint

$$\Pi_{i=1}^3 m^i = 0. \quad (4.2)$$

As a result chiral fermions appear from the $aa^*$, $dd^*$, $ee^*$, $ff^*$, sectors with corresponding fermions $\omega_L$, $y_R$; $s_1^L$; $s_2^R$; $s_3^R$.

f) The PS left-right symmetry is being broken at $M_{GUT}$. As there is no constraint from first principles for $M_{GUT}$ we shall take it equal to the string scale. The surviving gauge symmetry is that of the SM augmented by five anomaly free $U(1)$ symmetries, including the two chosen added extra $U(1)$ branes, surviving the Green-Schwarz mechanism. The breaking of the latter $U(1)$ symmetries will be facilitated by having the $dd^*$, $ee^*$, $ff^*$,
singlets scalars will be localized in these intersections, that are superpartners of the corresponding fermions.

The third tori is permanently tilted. Also b-brane is parallel to the c-brane (and the a D6-brane is parallel to the d, e, f D6-branes). The cancellation of the RR tadpole constraints is solved from multiparameter sets of solutions. They are given in table 2. Their satisfaction needs a number of extra branes, positioned at \((1/\beta_1, 0)(1/\beta_2, 0), (2, 0)\), having non-trivial intersection numbers with the a, d, e, f branes and thus creating extra fermions which finally become massive by arranging for some sectors to have N=1 SUSY.

| \( N_i \) | \((n_1^i, m_1^i)\) | \((n_2^i, m_2^i)\) | \((n_3^i, m_3^i)\) |
|-------------|-----------------|-----------------|-----------------|
| \( N_a = 4 \) | \( (0, \epsilon) \) | \( (n_2^a, 3\epsilon\tilde{\epsilon}\beta_2) \) | \( (1, \tilde{\epsilon}/2) \) |
| \( N_b = 2 \) | \( (-1, \epsilon m_1^b) \) | \( (1/\beta_2, 0) \) | \( (1, \tilde{\epsilon}/2) \) |
| \( N_c = 2 \) | \( (1, \epsilon m_1^c) \) | \( (1/\beta_2, 0) \) | \( (1, -\tilde{\epsilon}/2) \) |
| \( N_d = 1 \) | \( (0, \epsilon) \) | \( (n_2^d, -3\epsilon\tilde{\epsilon}\beta_2) \) | \( (2, -\tilde{\epsilon}) \) |
| \( N_e = 1 \) | \( (0, \epsilon) \) | \( (n_2^e, -2\epsilon\tilde{\epsilon}\beta_2) \) | \( (2, -\tilde{\epsilon}) \) |
| \( N_f = 1 \) | \( (0, \epsilon) \) | \( (n_2^f, -2\epsilon\tilde{\epsilon}\beta_2) \) | \( (1, -\frac{1}{2}\tilde{\epsilon}) \) |
| \( 1 \) | \( (1/\beta_1, 0) \) | \( (1/\beta_2, 0) \) | \( (2, 0) \) |
| \( \vdots \) | \( \vdots \) | \( \vdots \) | \( \vdots \) |
| \( N_h \) | \( (1/\beta_1, 0) \) | \( (1/\beta_2, 0) \) | \( (2, 0) \) |

Table 2: RR tadpole solutions for PS-III classes of GUT models with six stacks of intersecting D6-branes giving rise to the fermionic spectrum and the SM, \( SU(3)_C \times SU(2)_L \times U(1)_Y \), gauge group at low energies. The wrappings depend on four integer parameters, \( n_a^2, n_d^2, n_e^2, n_f^2 \), the NS-background \( \beta_i \) and the phase parameters \( \epsilon = \tilde{\epsilon} = \pm 1 \). Also there is an additional dependence on the two wrapping numbers, integer of half integer, \( m_1^b, m_1^c \). Note also that the presence of the \( N_h \) extra U(1) branes.

The first tadpole condition in (2.6) depends on the number of extra branes

\[
N_h \frac{2}{\beta_1\beta_2} = 16. \tag{4.3}
\]

9We denoted by \( h \) the presence of extra U(1) branes.
Also the third tadpole condition reads:

\[(2n_a^2 - n_d^2 - n_e^2 - \frac{1}{2}n_f^2) + \frac{1}{\beta_2} (m_b^1 - m_c^1) = 0. \quad (4.4)\]

h) the hypercharge operator is defined as usual in this classes of GUT models (see also [6]) as a linear combination of the three diagonal generators of the $SU(4)$, $SU(2)_L$, $SU(2)_R$ groups:

\[ Y = \frac{1}{2} T_{3R} + \frac{1}{2} T_{B-L}, \quad T_{3R} = \text{diag}(1, -1), \quad T_{B-L} = \text{diag}(\frac{1}{3}, \frac{1}{3}, \frac{1}{3}, -1). \quad (4.5)\]

Also,

\[ Q = Y + \frac{1}{2} T_{3L} \quad (4.6)\]

## 5 Cancellation of $U(1)$ Anomalies

In order to saw that the classes of models described by the solutions to the RR tadpoles of table (2) break at low energy to the SM we have first to exhibit that the massless additional $U(1)$’s originally present in the models in the string scale receive a mass and disappear from the low energy spectrum.

The mixed anomalies $A_{ij}$ of the six $U(1)$’s with the non-Abelian gauge groups are given by

\[ A_{ij} = \frac{1}{2} (I_{ij} - I_{ij'}) N_i. \quad (5.1)\]

In the orientifolded type IIA toroidal models the gauge anomaly cancellation [34] is achieved through a generalized GS mechanism [3] that makes use of the 10-dimensional RR gauge fields $C_2$ and $C_6$ and gives at four dimensions \(^{10}\) the couplings to gauge fields

\[ N_a m_a^1 m_a^2 m_a^3 \int_{M_4} B_2^o \wedge F_a; \quad n_b^1 n_b^2 n_b^3 \int_{M_4} C^o \wedge F_b \wedge F_b, \quad (5.2)\]

\[ N_a n^I K m^I \int_{M_4} B_2^I \wedge F_a; \quad n_b^I n_b^J m^K \int_{M_4} C^I \wedge F_b \wedge F_b, \quad (5.3)\]

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

\[ C^o \equiv \int_{(T^2)^1 \times (T^2)^2 \times (T^2)^3} C_6; \quad C^I \equiv \int_{(T^2)^I} C_2, \quad (5.4)\]

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

\(^{10}\)Note that gravitational anomalies cancel since D6-branes never intersect O6-planes.
The triangle anomalies \((5.1)\) cancel from the existence of the string amplitude involved in the GS mechanism \([34]\) in four dimensions \([3]\). The latter amplitude, where the \(U(1)_a\) gauge field couples to one of the propagating \(B_2\) fields, coupled to dual scalars, that couple in turn to two \(SU(N)\) gauge bosons, is proportional \([3]\) to

\[- N_a m_a^1 m_a^m n_b^1 n_b^m - N_a \sum_I n_a^I n_a^J n_b^K m_a^I m_b^J m_b^K , I \neq J, K \quad (5.5)\]

We make the minimal choice

\[\beta_1 = \beta_2 = 1/2 \quad (5.6)\]

that requires two extra D6 branes.

In this case the structure of \(U(1)\) couplings reads:

\[B_2^3 \wedge \left[\frac{- \tilde{\epsilon}_2}{\beta_2}\right] (F^b + F^c),\]

\[B_2^1 \wedge [\epsilon] [4n_a^2 F^a + 2 \frac{m_1^a}{\beta^2} F^b + 2 \frac{m_1^b}{\beta^2} F^c + 2n_a^2 F^d + 2n_b^2 F^e + n_1^2 F^f],\]

\[B_2^3 \wedge (\beta^2) \left(6F^a + 3F^d + 2F^e + F^f\right). \quad (5.7)\]

As can be seen from \((5.7)\), two anomalous combinations of \(U(1)\)'s, e.g. \(6F^a + 3F^d + 2F^c + F^f\), \((F^b + F^c)\) become massive through their couplings to RR fields \(B_2^a, B_2^2\). In addition, there is an anomaly free model dependent \(U(1)\) which is getting massive via its coupling to the RR field \(B_2^1\). In addition, there are three non-anomalous \(U(1)\)'s, that may be broken by the vevs of singlet scalars generated either, by imposing N=1 SUSY on sectors of the form \(dd^*, ee^*, ff^*,\) or by sectors involving the presence of the extra branes (see later discussion). They are:

\[U(1)^{(4)} = (Q^b - Q^e) + (Q^a - Q^d - Q^e - Q^f),\]

\[U(1)^{(5)} = \frac{1}{3}Q^a - \frac{15}{9}Q^d + Q^e + Q^f\]

\[U(1)^{(6)} = 2Q_a + Q^d - 16Q^e + 17Q^f . \quad (5.8)\]

Crucial for the satisfaction of RR tadpoles is the addition of \(N_h\) hidden branes. We note that for simplicity that when \(\beta^1 = \beta^2 = 1/2 , N_h = 2\). In this case, we simply have to add these extra \(U(1)\)'s to the bunch of \((5.8)\),

\[U(1)^{(7)} = \hat{F}^h_1, \quad U(1)^{(8)} = \hat{F}^h_2 . \quad (5.9)\]

We note that the model independent \(U(1)\)'s \((5.8)\), survive massless the presence of the generalized Green-Shcwarz mechanism imposed by the existence of the couplings \((5.7)\),
as long as
\[
12n_a^2 - 30n_d^2 + 18n_e^2 + 9n_f^2 = 0,
8n_a^2 + 2n_d^2 - 32n_e^2 + 17n_f^2 = 0
\] (5.10)

They can be broken by the existence of singlets generated either, by imposing N=1 SUSY on sectors of the form \(dd^*, ee^*, ff^*\), or on sectors involving the presence of the extra branes needed to satisfy the RR tadpole cancellation conditions. The reader should notice that the conditions for demanding that some sectors respect N=1 SUSY, that in turn guarantee the existence of a Majorana coupling for right handed neutrinos as well creating singlets necessary to break the U(1)’s (5.8), solve the condition (5.10). They are analyzed in section (7.1).

Also the couplings of the dual scalars \(C^I\) of \(B^I_2\), required to cancel the mixed anomalies of the \(U(1)\)- non-abelian \(SU(N_a)\) anomalies, appear as:
\[
C^o \wedge \left[ -\frac{1}{\beta^2}([F^b \wedge F^b] - (F^c \wedge F^c) - \frac{2}{\beta^1}(F^{h^1} \wedge F^{h^1} + F^{h^2} \wedge F^{h^2}))\right],
\]
\[
C^2 \wedge \left[ \frac{\epsilon \tilde{\epsilon}}{2}[n_a^2(F^a \wedge F^a) + \frac{m_1^2}{\beta^2}(F^b \wedge F^b) - \frac{m_e^2}{\beta^2}(F^c \wedge F^c) - 2n_d^2(F^d \wedge F^d)
- 2n_e^2(F^e \wedge F^e) - n_f^2(F^e \wedge F^e)\right],
\]
\[
C^3 \wedge [\beta^2 \tilde{\epsilon}][3(F^a \wedge F^a) - 6(F^d \wedge F^d) - 4(F^e \wedge F^e) - 2(F^f \wedge F^f)]
\] (5.11)

6 Higgs sector

6.1 Stability of the configurations and Higgs sector

We have so far seen the appearance in the R-sector of \(I_{ab}\) massless fermions in the D-brane intersections transforming under bifundamental representations \(N_a, \bar{N}_b\). In intersecting brane words, besides the presence of massless fermions at each intersection, we have present of an equal number of massive bosons, in the NS-sector, in the same representations as the massless fermions [19]. Their mass is of order of the string scale. However, some of those massive bosons may become tachyonic\(^{11}\), 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 [35]. Denoting the twist vector by \((\vartheta_1, \vartheta_2, \vartheta_3, 0)\), in the NS open string

\(^{11}\)For consequences when these set of fields may become massless see [24].
sector the lowest lying states are given by

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

Exactly at the point, where one of these masses may become massless we have preservation of \(\mathcal{N} = 1\) SUSY. We note that the angles at the four different intersections can be expressed in terms of the parameters of the tadpole solutions.

- **Angle structure and Higgs fields for PS-III classes of models**

The angles at the different intersections can be expressed in terms of the tadpole solution parameters. We define the angles:

\[
\begin{aligned}
\theta_1 &= \frac{1}{\pi} \cot^{-1} \frac{R_1^{(1)}}{\epsilon m_1 R_2^{(1)}}, \\
\theta_2 &= \frac{1}{\pi} \cot^{-1} \frac{\epsilon R_1^{(2)}}{3 \epsilon \beta_2 R_2^{(2)}}, \\
\theta_3 &= \frac{1}{\pi} \cot^{-1} \frac{2 R_1^{(3)}}{R_2^{(3)}}, \\
\tilde{\theta}_1 &= \frac{1}{\pi} \cot^{-1} \frac{R_1^{(1)}}{\epsilon m_1 R_2^{(1)}}, \\
\tilde{\theta}_2 &= \frac{1}{\pi} \cot^{-1} \frac{\epsilon \tilde{\epsilon} n^2 R_1^{(1)}}{3 \beta_2 R_2^{(1)}}, \\
\tilde{\theta}_3 &= \frac{1}{\pi} \cot^{-1} \frac{\epsilon \tilde{\epsilon} n^2 R_1^{(1)}}{2 \beta_2 R_2^{(1)}}, \\
\theta'_2 &= \frac{1}{\pi} \cot^{-1} \frac{\epsilon \tilde{\epsilon} n^2 R_1^{(1)}}{2 \beta_2 R_2^{(1)}}
\end{aligned}
\]

where we consider \(\epsilon \tilde{\epsilon} > 0, \epsilon m_1 > 0, \epsilon m_2 > 0\) and \(R_i^{(j)}, i = 1, 2\) are the compactification radii for the three \(j = 1, 2, 3\) tori, namely projections of the radii onto the cartesian axis \(X^{(i)}\) directions when the NS flux B field, \(b^k, k = 1, 2\) is turned on.

At each of the eight non-trivial intersections we have the presence of four states \(t_i, i = 1, \cdots , 4\), that could become massless, associated to the states \([6.1]\). Hence we have a total of thirty two different scalars in the model. The setup is seen clearly if we look at figure one. These scalars are generally massive but for some values of their angles could become tachyonic (or massless).

Also, if we demand that the scalars associated with \([6.1]\) and PS-III models may not be tachyonic, we obtain a total of eighteen conditions for the PS-III type models with a D6-brane at angles configuration to be stable. They are given in Appendix I. We don’t consider the scalars from the \(aa^*, dd^*, ee^*, ff^*\) intersections. For these sectors we will require later that they preserve \(N = 1\) SUSY. As a result all scalars in these sectors may become massive or receive vevs and becoming eventually massive.

Let us now turn our discussion to the Higgs sector of PS-III models. In general there are two different Higgs fields that may be used to break the PS symmetry. We

\[^{12}\text{we assume } 0 \leq \vartheta_i \leq 1.\]
Figure 1: Assignment of angles between D6-branes on the PS-III class of models based on the initial gauge group $U(4)_C \times U(2)_L \times U(2)_R$. The angles between branes are shown on a product of $T^2 \times T^2 \times T^2$. We have chosen $m_b^1, m_c^1, n_a^2, n_d^2, n_e^2, n_f^2 > 0, \epsilon = \tilde{\epsilon} = 1$. These models break to low energies to exactly the SM.
remind that they were given in (3.3). The question is if $H_1, H_2$ are present in the spectrum of PS-III models. In general, tachyonic scalars stretching between two different branes $\tilde{a}, \tilde{b}$, can be used as Higgs scalars as they can become non-tachyonic by varying the distance between the branes. Looking at the $I_{ac^*}$ intersection we can confirm that the scalar doublets $H^\pm$ get localized. They come from open strings stretching between the $U(4)$ $a$-brane and $U(2)_R$ $c^*$-brane.

$$
\begin{array}{|c|c|c|c|c|c|}
\hline
\text{Intersection} & \text{PS breaking Higgs} & Q_a & Q_b & Q_c & Q_d \\
\hline
ac^* & H_1 = (4,1,2) & 1 & 0 & 1 & 0 \\
ac^* & H_2 = (4,1,2) & -1 & 0 & -1 & 0 \\
\hline
\end{array}
$$

Table 3: Higgs fields responsible for the breaking of $SU(4) \times SU(2)_R$ symmetry of the $SU(4)_C \times SU(2)_L \times SU(2)_R$ with D6-branes intersecting at angles. These Higgs are responsible for giving masses to the right handed neutrinos in a single family.

The $H^\pm$'s come from the NS sector and correspond to the states\footnote{A similar set of states was used in \cite{3} to provide the model with electroweak Higgs scalars.}

\begin{align*}
\text{State} & \quad \text{Mass}^2 \\
(-1 + \vartheta_1, \vartheta_2, 0, 0) & \quad \alpha'(\text{Mass})^2_{H^+} = \frac{Z_3}{4\pi^2} + \frac{1}{2}(\vartheta_2 - \vartheta_1) \\
(\vartheta_1, -1 + \vartheta_2, 0, 0) & \quad \alpha'(\text{Mass})^2_{H^-} = \frac{Z_3}{4\pi^2} + \frac{1}{2}(\vartheta_1 - \vartheta_2)
\end{align*}

(6.3)

where $Z_3$ is the distance squared in transverse space along the third torus, $\vartheta_1, \vartheta_2$ are the (relative) angles between the $a^-, c^*$-branes in the first and second complex planes respectively. The presence of scalar doublets $H^\pm$ can be seen as coming from the field theory mass matrix

\begin{equation}
(H_1^* H_2) \left( M^2 \right) \left( \begin{array}{c} H_1 \\ H_2^* \end{array} \right) + h.c. \quad (6.4)
\end{equation}

where

\begin{equation}
M^2 = M_s^2 \left( \begin{array}{cc} Z_3^{(ac^*)} (4\pi^2)^{-1} & \frac{1}{2} |\vartheta_1^{(ac^*)} - \vartheta_2^{(ac^*)}| \\ \frac{1}{2} |\vartheta_1^{(ac^*)} - \vartheta_2^{(ac^*)}| & Z_3^{(ac^*)} (4\pi^2)^{-1} \end{array} \right) \quad (6.5)
\end{equation}

The fields $H_1$ and $H_2$ are thus defined as

\begin{equation}
H^\pm = \frac{1}{2} (H_1^* \pm H_2) \quad (6.6)
\end{equation}
where their charges are given in table (3). Hence the effective potential which corresponds to the spectrum of the PS symmetry breaking Higgs scalars is given by

\[ V_{\text{Higgs}} = m_H^2(|H_1|^2 + |H_2|^2) + (m_B^2 H_1 H_2 + \text{h.c}) \]  

(6.7)

where

\[ m_H^2 = \frac{Z_3^{(ac)}}{4\pi^2\alpha'} ; \quad m_B^2 = \frac{1}{2\alpha'}|\psi_1^{(ac)} - \psi_2^{(ac)}| \]  

(6.8)

The precise values of \( m_H^2 \), \( m_B^2 \), for PS-III classes of models are given by

\[ m_H^2_{\text{PS-III}} = \frac{(\xi'_a + \xi'_c)^2}{\alpha'}, \quad m_B^2_{\text{PS-III}} = \frac{1}{2\alpha'}|\tilde{\theta}_1 - \tilde{\theta}_2| \]  

(6.9)

where \( \xi'_a(\xi'_c) \) is the distance between the orientifold plane and the \( a(c) \) branes and \( \tilde{\theta}_1, \tilde{\theta}_2 \) were defined in (6.2). Thus

\[ m_B^2_{\text{PS-III}} = \frac{1}{2} |m_{\chi_2}^2(t_2) + m_{\chi_2}^2(t_3) - m_{\chi_b}^2(t_1) - m_{\chi_b}^2(t_3)| \]

(6.10)

For PS-III models the number of Higgs present is equal to the the intersection number product between the \( a-, c^*- \) branes in the first and second complex planes,

\[ n_{H^\pm}^{\text{PS-III}} I_{ac^*} = 3. \]  

(6.11)

A comment is in order. For PS-III models the number of PS Higgs is three. That means that we have three intersections and to each one we have a Higgs particle which is a linear combination of the Higgs \( H_1 \) and \( H_2 \).

The electroweak symmetry breaking could be delivered through the bidoublets Higgs present in the \( bc^* \) intersection (seen in table (4). In principle these can be used to give mass ot quarks and leptons. In the present models their number is given by the intersection number of the \( b, c^* \) branes in the first tori

\[ n_{b,c^*}^{\text{PS-III}} |\epsilon(m_c^1 - m_b^1)| = |\beta^2(2n_a^2 - n_d^2 - n_e^2 - \frac{n_f^2}{2})| \]  

(6.12)

The number of the electroweak bidoublets in the PS-III models depends on the difference \( |m_b^1 - m_c^1| \), taking into account the conditions for N=1 SUSY in some sectors, e.g. (7.18) in section (7.1), we get \( n_{b,c} = 0 \) and thus \( m_b^1 = m_c^1 \). However, this is not a
Intersection | Higgs | $Q_a$ | $Q_b$ | $Q_c$ | $Q_d$
---|---|---|---|---|---
$bc^*$ | $h_1 = (1, 2, 2)$ | 0 | 1 | 1 | 0
$bc^*$ | $h_2 = (1, 2, 2)$ | 0 | $-1$ | $-1$ | 0

Table 4: Higgs fields present in the intersection $bc^*$ of the $SU(4)_C \times SU(2)_L \times SU(2)_R$ classes of models with D6-branes intersecting at angles. These Higgs give masses to the quarks and leptons in a single family and could have been responsible for electroweak symmetry breaking if their net number was not zero.

Table 5: Higgs fields present in the intersection $bc$ of the $SU(4)_C \times SU(2)_L \times SU(2)_R$ classes of models with D6-branes intersecting at angles. These Higgs give masses to the quarks and leptons in a single family and are responsible for electroweak symmetry breaking.

Problem for electroweak symmetry breaking as (see section 8) a different term is used to provide Dirac masses to quarks, leptons and neutrinos \(^{14}\). In the present models it important that

\[
I_{bc} = |m_c^1 + m_b^1| = 2|m_b^1|
\]

may be chosen different from zero. Thus an alternative set of electroweak Higgs may be provided from the the NS sector where the lightest scalar states $h^\pm$ originate from open strings stretching between the $bc$ branes, e.g. named as $h_3, h_4$.

\[
\begin{align*}
\text{State} & & \text{Mass}^2 \\
(-1 + \vartheta_1, 0, \vartheta_3, 0) & & \alpha'(\text{Mass})^2 = \frac{Z_{bc}^2}{4\pi^2} + \frac{1}{2}(\vartheta_3 - \vartheta_1) \\
(\vartheta_1, 0, -1 + \vartheta_3, 0) & & \alpha'(\text{Mass})^2 = \frac{Z_{bc}^2}{4\pi^2} + \frac{1}{2}(\vartheta_1 - \vartheta_3)
\end{align*}
\]

(6.14)

where $Z_{bc}^2$ is the relative distance in transverse space along the second torus from the orientifold plane, $\vartheta_1$, $\vartheta_3$, are the (relative)angle between the $b$-, $c$-branes in the first and third complex planes.

Hence the presence of scalar doublets $h^\pm$ defined as

\[
h^\pm = \frac{1}{2}(h_3^\pm \pm h_4)
\]

(6.15)

\(^{14}\)the same conditions hold for the PS-II models of \cite{8}
The precise values of the angles $\vartheta$ can be seen as coming from the field theory mass matrix

$$(h_3^* h_4) \left( M^2 \right) \left( \begin{array}{c} h_3 \\ h_4^* \end{array} \right) + h.c.$$  \hspace{1cm} (6.16)

where

$$M^2 = M_s^2 \begin{pmatrix} Z_{23}^{(bc)} (4\pi^2)^{-1} & \frac{1}{2} |\vartheta_1^{(bc)} - \vartheta_3^{(bc)}| \\ \frac{1}{2} |\vartheta_1^{(bc)} - \vartheta_3^{(bc)}| & Z_{23}^{(bc)} (4\pi^2)^{-1} \end{pmatrix},$$  \hspace{1cm} (6.17)

The effective potential which corresponds to the spectrum of electroweak Higgs $h_3, h_4$ may be written as

$$V_{Higgs}^{bc} = m_H^2 (|h_3|^2 + |h_4|^2) + (m_B^2 h_3 h_4 + h.c)$$  \hspace{1cm} (6.18)

where

$$m_H^2 = \frac{Z_{23}^{(bc)}}{4\pi^2 \alpha'}; \quad m_B^2 = \frac{1}{2\alpha'} |\vartheta_1^{(bc)} - \vartheta_3^{(bc)}|$$  \hspace{1cm} (6.19)

The precise values of $m_H^2, m_B^2$ are

$$m_H^{PS-II} = \frac{(\chi_b^{(2)} + \chi_c^{(2)})^2}{\alpha'}; \quad m_B^{PS-II} = \frac{1}{2\alpha'} |\theta_1 - \bar{\theta}_1 - 2\bar{\theta}_3|$$  \hspace{1cm} (6.20)

where $\theta_1, \bar{\theta}_1, \bar{\theta}_3$ were defined in (6.2). Also $\chi_b, \chi_c$ are the distances of the $b, c$ branes from the orientifold plane in the second torus. The values of the angles $\vartheta_1, \vartheta_1, \vartheta_2$ can be expressed in terms of the scalar masses at the various intersections

$$\frac{1}{\pi} \vartheta_1 = \frac{1}{2} |m_{F_L}^2 (t_2) + m_{F_L}^2 (t_3) - 1| = \frac{1}{2} |m_{\chi_L}^2 (t_2) + m_{\chi_L}^2 (t_3) + 1|$$

$$= \frac{1}{2} |m_{\chi_L}^2 (t_2) + m_{\chi_L}^2 (t_3) + 1| = \frac{1}{2} |m_{\chi_L}^2 (t_2) + m_{\chi_L}^2 (t_3) + 1|$$  \hspace{1cm} (6.21)

$$\frac{1}{\pi} \bar{\vartheta}_1 = \frac{1}{2} |m_{F_R}^2 (t_2) - m_{F_R}^2 (t_3) - 1| = \frac{1}{2} |m_{\chi_R}^2 (t_2) + m_{\chi_R}^2 (t_3) - 1|$$

$$= \frac{1}{2} |m_{\chi_R}^2 (t_2) + m_{\chi_R}^2 (t_3) - 1| = \frac{1}{2} |m_{\chi_R}^2 (t_2) + m_{\chi_R}^2 (t_3) - 1|$$  \hspace{1cm} (6.22)

$$\frac{1}{\pi} \vartheta_2 = \frac{1}{2} |m_{F_L}^2 (t_1) + m_{F_L}^2 (t_3)| = \frac{1}{2} |m_{F_R}^2 (t_1) + m_{F_R}^2 (t_3)|$$  \hspace{1cm} (6.23)

$$\frac{1}{\pi} \vartheta_3 = \frac{1}{4} |m_{F_L}^2 (t_1) + m_{F_L}^2 (t_2)| = \frac{1}{4} |m_{F_R}^2 (t_1) + m_{F_R}^2 (t_2)|$$

$$= \frac{1}{4} |m_{\chi_L}^2 (t_1) + m_{\chi_L}^2 (t_2)| = \frac{1}{4} |m_{\chi_R}^2 (t_1) + m_{\chi_R}^2 (t_2)|, j = 1, 2, 3$$  \hspace{1cm} (6.24)
7 Singlet scalar generation - $N = 1$ SUSY on Intersections

In this section, we will present a gauge singlet generation mechanism by demanding that certain open string sectors respect $N = 1$ supersymmetry. Similar considerations were useful on the other GUT classes of PS-like models with only the SM at low energy [6, 7, 8] and also in the models from a SM-like configuration at $M_s$ with the number of stacks being only five and six respectively [4, 5]. The singlet scalars will be necessary for giving masses to the $U(1)$’s which they don’t couple to the RR fields. Also they will be used to realize a Majorana mass term for the right handed neutrinos. We note that the spectrum of PS-III classes of models described in table (1) is massless at this point. Supersymmetry will create singlet scalars which receive vevs and generate masses for the otherwise massless fermions $\chi_1^L, \chi_2^L, \chi_3^L, \chi_1^R, \chi_2^R, \chi_3^R, \omega_L, y_R, s_1^R, s_2^R, s_3^R$.

7.1 PS-III models with $N=1$ SUSY

In this part we will show that model dependent conditions, obtained by demanding that the extra $U(1)$’s do not have non-zero couplings to the RR fields, are necessary conditions in order to have scalar singlet generation that could effectively break the extra $U(1)$’s. These conditions will be alternatively obtained by demanding that certain string sectors respect $N = 1$ supersymmetry.

In general, for $N = 1$ supersymmetry to be preserved at some intersection between two branes $L, M$, we need to satisfy $\pm \vartheta_{ab}^{1} \pm \vartheta_{ab}^{2} \pm \vartheta_{ab}^{3}$ for some choice of signs, where $\vartheta_{ab}^{i}, i = 1, 2, 3$ are the relative angles of the branes $L, M$ across the three 2-tori. The latter rule will be our main tool in getting N=1 SUSY on intersections.

- The ac sector respects $N = 1$ supersymmetry.

The condition for $N = 1$ SUSY on the ac-sector is $^{15}$:
\[
\pm \left( \frac{\pi}{2} + \tilde{\vartheta}_1 \right) \pm \vartheta_2 \pm 2\vartheta_3 = 0, \quad (7.1)
\]

This condition can be solved by choosing:
\[
ac \rightarrow \left( \frac{\pi}{2} + \tilde{\vartheta}_1 \right) + \vartheta_2 - 2\vartheta_3 = 0, \quad (7.2)
\]

and thus may be solved by the choice $^{16}$
\[
- \tilde{\vartheta}_1 = \vartheta_2 = \vartheta_3 = \frac{\pi}{4}, \quad (7.3)
\]

$^{15}$We have chosen $m_c^1 < 0$.

$^{16}$We have set $U^{(i)} = \frac{R^{(i)}_L}{R^{(i)}_T}, i = 1, 2, 3$
effectively giving us
\[- \frac{1}{\epsilon m^1_\epsilon} U^{(1)} = \frac{(\epsilon \bar{\epsilon}) n^2_a}{3 \beta^2 U^{(2)}} = \frac{2 \bar{\epsilon}}{U^{(3)}} = \frac{\pi}{4}. \tag{7.4}\]

By imposing \(N = 1\) SUSY on an intersection ac the massless scalar superpartner of \(\bar{F}_R\) appears, the \(\bar{F}_R^\beta\). Note that in (7.4) the imposition of \(N=1\) SUSY connects the complex structure moduli \(U^i\) in the different tori and thus reduces the moduli degeneracy of the theory.

- **The \(dd^*\) sector preserves \(\mathcal{N} = 1\) supersymmetry**

As we noted in the appendix the presence of \(N=1\) supersymmetry in the sectors \(dd^*, \ ee^*\) is equivalent to the absence of tachyons in those sectors.

The general form of the \(\mathcal{N} = 1\) supersymmetry condition on this sector is
\[\pm \pi \pm 2 \bar{\vartheta}_2 \pm 2 \vartheta_3 = 0, \tag{7.5}\]
which may be solved by the choice
\[- \pi + 2 \bar{\vartheta}_2 + 2 \vartheta_3 = 0, \tag{7.6}\]
Hence
\[\bar{\vartheta}_2 = \vartheta_3 = \frac{\pi}{4}, \tag{7.7}\]
that is
\[\frac{\epsilon \bar{\epsilon} n^2_a}{3 \beta^2} U^{(2)} = \frac{2 \bar{\epsilon}}{U^{(3)}} = \frac{\pi}{4}. \tag{7.8}\]

From (7.4) and (7.8) we deduce that
\[n^2_a = n^2_d \tag{7.9}\]

Due to the presence of \(N=1\) SUSY on this sector we have localized the superpartner of the \(s^1_R\), the \(s^1_B\).

- **The \(ee^*\) sector preserves \(\mathcal{N} = 1\) supersymmetry**

The general form of the \(\mathcal{N} = 1\) supersymmetry condition on this sector is
\[\pm \pi \pm 2 \bar{\vartheta}_2 \pm 2 \vartheta_3 = 0, \tag{7.10}\]
which we may recast in the form
\[- \pi + 2 \bar{\vartheta}_2 + 2 \vartheta_3 = 0, \tag{7.11}\]
be solved by the choice
\[ \bar{\vartheta}_2 = \vartheta_3 = \frac{\pi}{4}, \]  
(7.12)
that is
\[ \frac{(\bar{\epsilon}\epsilon)n_e^2}{2\beta_2 U^{(2)}} = \frac{2}{\bar{\epsilon}} U^{(3)} = \frac{\pi}{4}. \]  
(7.13)

Due to the presence of \( N=1 \) SUSY on this sector we have localized the superpartner of the \( s^2_R \), the \( s^2_B \).

- **The \( ff^* \) sector preserves \( \mathcal{N} = 1 \) supersymmetry**

The general form of the \( \mathcal{N} = 1 \) supersymmetry condition on this sector is
\[ \pm \pi \pm 2\vartheta'_2 \pm 2\vartheta_3 = 0, \]  
(7.14)
which we may recast in the form
\[ -\pi + 2\vartheta'_2 + 2\vartheta_3 = 0, \]  
(7.15)
be solved by the choice
\[ \vartheta'_2 = \vartheta_3 = \frac{\pi}{4}, \]  
(7.16)
that is
\[ \frac{(\bar{\epsilon}\epsilon)n_f^2}{2\beta_2 U^{(2)}} = \frac{2}{\bar{\epsilon}} U^{(3)} = \frac{\pi}{4}. \]  
(7.17)

From (7.8), (7.4), (7.13), (7.17), we derive the conditions
\[ 2n^2_a = 3n^2_e = 3n^2_f = 2n^2_d. \]  
(7.18)

These conditions solve exactly the conditions (5.10).

Due to the presence of \( N=1 \) SUSY on this sector we have localized the superpartner of the \( s^3_R \), the \( s^3_B \).

Thus the presence of \( N=1 \) supersymmetry in \( dd^*, ee^*, ff^* \) sectors guarantees the presence of gauge singlets as scalar superpartners of the \( s^1_R, s^2_R, s^3_R \) fermions, e.g. \( s^1_B, s^2_B, s^3_B \) that may may receive vevs of undetermined order.

Also what is is evident by looking at conditions (5.10), (7.18) is that the conditions of orthogonality for the extra \( U(1)'s \) to survive massless the generalized Green-Schwarz mechanism is equivalent to the conditions for \( N=1 \) supersymmetry in the leptonic sectors \( dd^*, ee^* \). The latter condition is equivalent to the absence of tachyons in the sectors \( dd^*, ee^* \).
7.2 Gauge singlet generation from the extra U(1) branes

In this section, we will present an alternative mechanism for generating singlet scalars. We had already seen that in leptonic sectors involving U(1) branes, e.g. \(dd^*, ee^*,\) brane imposing N=1 SUSY creates singlet scalars. This is reflected in the fact that in U(1) \(j\)-branes, sectors in the form \(jj^*\) had localized in their intersection gauge singlet fermions. Thus imposing N=1 SUSY on those sectors help us to get rid of these massless fermions, by making them massive through their couplings to their superpartner gauge singlet scalars.

What we will become clear in this sector is that the presence of supersymmetry in particular sectors involving the extra branes creates singlet scalars that provide the couplings that make massive some non-SM fermions.

In order to show the creation of gauge singlets from sectors involving extra branes we will make our points by using only one of the extra \(N_h U(1)\) branes, e.g. the \(N_{h_1}\) one. The following discussion can be identically repeated for the other extra branes.

Thus due to the non-zero intersection numbers of the \(N_{h_1}\) U(1) brane with \(a,d\) branes the following sectors are present: \(ah, ah^*, dh, dh^*\).

- **ah-sector**

  Because \(I_{ah} = -\frac{3}{\beta_1}\) we have present \(|I_{ah}|\) massless fermions \(\kappa_1^f\) in the representations

  \[
  \kappa_1^f \rightarrow (\bar{4}, 1, 1)_{(-1,0,0,0,0,0,1)}
  \]

  where the subscript last entry denotes the U(1) charge of the of the U(1) extra brane \(^{17}\).

- **ah*-sector

  Because \(I_{ah^*} = -\frac{3}{\beta_1}\), there are \(|I_{ah^*}|\) fermions \(\kappa_2^f\) localized in the \(ah^*\) intersection and appearing as

  \[
  \kappa_2^f \rightarrow (\bar{4}, 1, 1)_{(-1,0,0,0,0,0,1)}ldh2
  \]

- **dh-sector**

  Because \(I_{dh} = -\frac{6}{\beta_1}\), there are present \(|I_{dh}|\) fermions \(\kappa_3^f\) transforming in the representations

  \[
  \kappa_3^f \rightarrow (1, 1, 1)_{(0,0,0,-1,0,0,-1)}
  \]

\(^{17}\)We don’t exhibit the beyond the seventh entry of the rest of the extra branes as for the present discussion are identically zero.
We further require that this sector respects $N = 1$ supersymmetry. In this case we have also present the massless scalar fields $\kappa_3^B$,

$$\kappa_3^B \rightarrow (1, 1, 1)_{(0,0,0,-1,0,0,-1)}, \text{laradh5}$$

(7.22)

The latter scalars receive a vev which we assume to be of order of the string scale.

The condition for $N = 1$ supersymmetry in this sector is exactly

$$-\frac{\pi}{2} + \bar{\vartheta}_2 + \vartheta_3 = 0$$

(7.23)

which is satisfied when $\bar{\vartheta}_2, \vartheta_3$ take the value $\pi/4$ in consistency with (7.11) and subsequently (7.11).

- **dh* -sector**

Because $I_{dh*} = -\frac{6}{\beta_1} \neq 0$, there are present $|I_{dh*}|$ fermions $\kappa_4^f$ in the representations

$$\kappa_4^f \rightarrow (1, 1, 1)_{(0,0,0,-1,0,0,-1)} ldh4$$

(7.24)

The condition that this sector respects $N=1$ SUSY is equivalent to the one is the $dh$-sector.

- **eh-sector**

Because $I_{eh} = -\frac{4}{\beta_1}$, there are present $|I_{eh}|$ fermions $\kappa_5^f$ transforming in the representations

$$\kappa_5^f \rightarrow (1, 1, 1)_{(0,0,0,0,0,-1,0,-1)}$$

(7.25)

Also we require that this sector preserves $N=1$ SUSY. Because of the presence of $N=1$ SUSY there is evident the presence of $|I_{eh}|$ bosons $\kappa_5^B$ transforming in the representations

$$\kappa_5^B \rightarrow (1, 1, 1)_{(0,0,0,0,0,0,-1,0,1)}$$

(7.26)

The condition for $N=1$ SUSY is

$$\pm \frac{\pi}{2} \pm \bar{\vartheta}_2 \pm \vartheta_3 = 0$$

(7.27)

which is exactly ‘half’ of the supersymmetry condition (7.11). When it is rearranged into the form

$$\frac{\pi}{2} + \bar{\vartheta}_2 - \vartheta_3 = 0,$$

(7.28)

it is solved by the choice (7.12).
\* eh*-sector In this sector, \( I_{eh^*} = -\frac{2}{\beta_1} \). Thus there are present \(|I_{eh^*}| \) fermions \( \kappa_6^f \) transforming in the representations

\[
\kappa_6^f \rightarrow (1, 1, 1)_{(0,0,0,0,-1,0,-1)}
\] (7.29)

The condition for N=1 SUSY to be preserved by this section is exactly (7.27). Thus we have present \(|I_{eh^*}| \) bosons \( \kappa_6^B \) transforming in the representations

\[
\kappa_6^B \rightarrow (1, 1, 1)_{(0,0,0,0,-1,0,-1)}
\] (7.30)

\* fh-sector

Because \( I_{fh} = -\frac{2}{\beta_1} \), there are present \(|I_{fh}| \) fermions \( \kappa_7^f \) transforming in the representations

\[
\kappa_7^f \rightarrow (1, 1, 1)_{(0,0,0,0,-1;1)}
\] (7.31)

Also we require that this sector preserves N=1 SUSY. Because of N=1 SUSY there are present \(|I_{fh}| \) bosons \( \kappa_7^B \) transforming in the representations

\[
\kappa_7^B \rightarrow (1, 1, 1)_{(0,0,0,0,-1;1)}
\] (7.32)

The condition for N=1 SUSY is

\[
\pm \frac{\pi}{2} \pm \vartheta_2' \pm \vartheta_3 = 0
\] (7.33)

which is ‘half’ of the supersymmetry condition (7.13). When it is rearranged into the form

\[
- \frac{\pi}{2} + \vartheta_2' + \vartheta_3 = 0,
\] (7.34)

it is solved by the choice (7.16).

\* fh*-sector In this sector, \( I_{fh^*} = -\frac{2}{\beta_1} \). Thus there are present \(|I_{fh^*}| \) fermions \( \kappa_8^f \) transforming in the representations

\[
\kappa_8^f \rightarrow (1, 1, 1)_{(0,0,0,0,-1,-1)}
\] (7.35)

The condition for N=1 SUSY to be preserved by this section is exactly (7.34). Thus we have present \(|I_{fh^*}| \) bosons \( \kappa_8^B \) transforming in the representations

\[
\kappa_8^B \rightarrow (1, 1, 1)_{(0,0,0,0,-1,-1)}
\] (7.36)
What we have found\footnote{Similar conclusions could be reached for the 5-stack GUT models of \cite{5}.} we have found that the conditions (5.10) derived as the model dependent conditions of the U(1)'s that survive the generalized Green-Schwarz mechanism, are equivalent:

- to have the leptonic branes, d, e, f, preserve N=1 SUSY on the sectors $dd^*$, $ee^*$, $ff^*$.
- to have the sectors made of a mixture of the extra and leptonic branes preserve N=1 SUSY. The presence of these conditions is independent from the number of extra U(1) branes present.

We will now show that all fermions, appearing from the non-zero intersections of the extra brane $U(N_{h_1})$ with the branes $a$, $d$, $e$, receive string scale mass and disappear from the low energy spectrum (see also a related discussion in the concluding section).

- The mass term for the $\kappa_1^f$ fermion reads:

$$
(4, 1, 1)_{(1,0,0,0,0,0;-1)} \langle (\bar{4}, 1, 1)_{(1,0,0,0,0,0;-1)} \rangle \langle (\bar{4}, 1, 2)_{(-1,0,1,0,0,0)} \rangle 
\times \langle (\bar{4}, 1, 2)_{(-1,0,1,0,0,0)} \rangle \langle (1, 1, 1)_{(0,0,0,0,1;1)} \rangle \langle (1, 1, 1)_{(0,0,0,0,0;1,-1)} \rangle

(7.37)
$$

or

$$
\bar{\kappa}_1^f \bar{\kappa}_1^f \langle H_2 \rangle \langle \bar{F}_H^R \rangle \langle \bar{\kappa}_2^8 \rangle \langle \bar{\kappa}_2^7 \rangle \sim \bar{\kappa}_1^f \bar{\kappa}_1^f M_s

(7.38)
$$

- The mass term for the $\kappa_2^f$ fermion reads:

$$
(4, 1, 1)_{(1,0,0,0,0,0;1)} \langle (\bar{4}, 1, 1)_{(1,0,0,0,0,0;1)} \rangle \langle (\bar{4}, 1, 2)_{(-1,0,1,0,0,0)} \rangle 
\times \langle (\bar{4}, 1, 2)_{(-1,0,1,0,0,0)} \rangle \langle (1, 1, 1)_{(0,0,0,0,0;1,-1)} \rangle \langle (1, 1, 1)_{(0,0,0,0,1;1)} \rangle

(7.39)
$$

or

$$
\bar{\kappa}_2^f \bar{\kappa}_2^f \langle H_2 \rangle \langle \bar{F}_H^R \rangle \langle \bar{\kappa}_2^8 \rangle \langle \bar{\kappa}_2^7 \rangle \sim \bar{\kappa}_2^f \bar{\kappa}_2^f M_s

(7.40)
$$

- The mass term for the $\kappa_3^f$ fermion reads:

$$
(1, 1, 1)_{(0,0,0,1,0;1)} \langle (1, 1, 1)_{(0,0,0,1,0;1)} \rangle \langle (1, 1, 1)_{(0,0,0,0;1;0)} \rangle 
\times \langle (1, 1, 1)_{(0,0,0,0;1;0)} \rangle \langle (1, 1, 1)_{(0,0,0,0;0;1)} \rangle

(7.41)
$$

or

$$
\bar{\kappa}_3^f \bar{\kappa}_3^f \langle \kappa_3^B \rangle \langle \kappa_3^B \rangle \sim M_s \bar{\kappa}_3^f \bar{\kappa}_3^f

(7.42)
$$
• The mass term for the $\kappa^f_4$ fermion reads:

\[
(1, 1, 1)_{(0,0,0,-1,0,-1)} \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\times \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\]

or

\[
\bar{\kappa}^f_4 \kappa^f_4 \langle \kappa^B_4 \rangle \sim M_s \bar{\kappa}^f_4 \kappa^f_4
\]

(7.43)

• The mass term for the $\kappa^f_5$ fermion reads:

\[
(1, 1, 1)_{(0,0,0,1,0,-1)} \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\times \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\]

or

\[
\bar{\kappa}^f_5 \kappa^f_5 \langle \kappa^B_5 \rangle \sim M_s \bar{\kappa}^f_5 \kappa^f_5
\]

(7.45)

• The mass term for the $\kappa^f_6$ fermion reads:

\[
(1, 1, 1)_{(0,0,0,0,1,0;1)} \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\times \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\]

or

\[
\bar{\kappa}^f_6 \kappa^f_6 \langle \kappa^B_6 \rangle \sim M_s \bar{\kappa}^f_6 \kappa^f_6
\]

(7.47)

• The mass term for the $\kappa^f_7$ fermion reads:

\[
(1, 1, 1)_{(0,0,0,0,1,0;1)} \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\times \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\]

or

\[
\bar{\kappa}^f_7 \kappa^f_7 \langle \kappa^B_7 \rangle \sim M_s \bar{\kappa}^f_7 \kappa^f_7
\]

(7.48)

• The mass term for the $\kappa^f_8$ fermion reads:

\[
(1, 1, 1)_{(0,0,0,0,1,0;1)} \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\times \langle (1, 1, 1)(0,0,0,1,0;1) \rangle
\]

or

\[
\bar{\kappa}^f_8 \kappa^f_8 \langle \kappa^B_8 \rangle \sim M_s \bar{\kappa}^f_8 \kappa^f_8
\]

(7.50)

• The mass term for the $\kappa^f_8$ fermion reads:
\[(1, 1, 1)(0, 0, 0, 1, 0; 1) 1, 1, 1)(0, 0, 0, 0, 1) \langle (1, 1, 1)(0, 0, 0, -1, 0; 1) \rangle \times \langle (1, 1, 1)(0, 0, 0, -1, 0; 1) \rangle \] (7.51)

or

\[\bar{\kappa}^f_s \bar{\kappa}^f_s \langle \kappa^B_s \rangle \langle \kappa^B_s \rangle \sim M_s \bar{\kappa}^f_s \bar{\kappa}^f_s \] (7.52)

### 7.3 Breaking the anomaly free massless U(1)'s

After breaking the PS gauge symmetry at \(M_{GUT}\) the initial gauge symmetry \(SU(4)_c \times SU(2)_L \times SU(2)_R \times U(1)_a \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e \times U(1)_f\) breaks to the SM gauge group \(SU(3) \times SU(2) \times U(1)_Y\) augmented by the extra anomaly free U(1)'s \(Q^4, Q^5, Q^6, Q^7, Q^8\). The last two are the hidden U(1)'s needed to satisfy the RR tadpole cancellation conditions. The extra U(1)'s may be broken if appropriate singlets are available. The latter may be created by appropriate choosing the angle parameters between the D6-branes when demanding N=1 supersymmetry to be preserved in particular sectors. In this way, \(U(1)^{(4)}\) may be broken if \(s^4_B\) gets a vev, \(U(1)^{(5)}\) may be broken if \(s^2_B\) gets a vev, \(U(1)^{(6)}\) may be broken if \(s^3_B\) gets a vev. Also \(U(1)^{(7)}\) and \(U(1)^{(8)}\) may be broken if one of the \(\kappa^B_3, \kappa^B_4, \kappa^B_5, \kappa^B_6, \kappa^B_7, \kappa^B_8\) gets a vev. Thus the U(1)'s surviving massless the Green-Schwarz mechanism may be broken easily by using singlets localized on intersections of the extra branes and leptonic branes as well from \(jj^*\) sectors.

We note that up to this point the only issue remaining is how we can give non-zero masses to all exotic fermions of table (1) beyond those that accommodate the quarks and leptons of the SM.

### 8 Yukawa couplings, neutrino and lepton masses

In this section, we discuss the issue of neutrino masses in the \(SU(4) \times S(2)_L \times SU(2)_R\) classes of PS-III GUTS. Also, we discuss in which way the additional fermions of table (1), receive a mass and disappear from the SM spectrum at the \(M_Z\) scale.

#### 8.1 Yukawa couplings and Neutrino masses

Proton decay is the most important problem of grand unified theories. In the usual versions of left-right symmetric PS models this problem is avoided as B-L is a gauged
symmetry but the problem is embedded in baryon number violating operators of sixth order, contributing to proton decay. In the PS-III models there is no-proton decay as baryon number is a gauged symmetry, the corresponding gauge boson becomes massive through its couplings to RR fields, and thus survives as a global symmetry to low energies. That is a plausible explanation for the origin of proton stability in general brane-world scenarios. Baryon B and lepton L numbers are related by \( Q_a = 3B + L \) and are given by

\[
B = \frac{Q_a + Q_{B-L}}{4}.
\]  

(8.1)

For intersecting brane worlds the usual tree level SM fermion mass generating trilinear Yukawa couplings between the fermion states \( F^i_L, \bar{F}^j_R \) and the Higgs fields \( H^k \) depends on the stretching of the worldsheet area between the three D6-branes which cross at those intersections. In the present Pati-Salam GUTS the trilinear Yukawa is

\[
Y^{ijk} = F^i_L \bar{F}^j_R h^k
\]

(8.2)

For a six dimensional torus in the leading order \[19\] we have,

\[
Y^{ijk} = e^{-\tilde{A}_{ijk}}
\]

(8.3)

where \( \tilde{A}_{ijk} \) is the worldsheet area \[19\] connecting the three vertices. The areas of each of the \( T^2 \) tori taking part in this interaction are typically of order one in string units. In \[6\] e.g. we have assumed that the areas of the second and third tori are close to zero. Thus in this case, the area of the full Yukawa coupling \(8.3\) may be given in the form

\[
Y^{ijk} = e^{-\frac{R_1 R_2}{a} A_{ij}}
\]

(8.4)

where \( R_1, R_2 \) the radii and \( A_{ij} \) the area of the two dimensional tori in the first complex plane. Here we exhibit the leading worldsheet correction coming from the first tori \[20\].

For the present class of GUTS we have seen that the electroweak bidoublets \[S.2\] are absent at tree level. However there is another coupling, which is non-renormalizable and of the same order :

\[
F_L \bar{F}_R \langle h_3 \rangle \langle F^H_R \rangle \langle H_2 \rangle \sim \nu F_L \bar{F}_R
\]

(8.5)

\[19\]We note that it is a general property of string theories for their Yukawa couplings to depend exponentially on the worldsheet area.

\[20\]The same hypothesis holds for any PS GUT model constructed so far, as a deformation of the quark and lepton intersection numbers, e.g. the PS-A, PS-I, PS-II classes of \[6, 7, 8\] respectively.
For a dimension five interaction term, like those involved in the Majorana mass term for the right handed neutrinos the interaction term is in the form
\[ Y_{lmni} = e^{-\tilde{A}_{lmni}}, \]  
(8.6)
where \( \tilde{A}_{lmni} \) the worldsheet area connecting the four interaction vertices. Assuming that the areas of the second and third tetragonal are close to zero the four term coupling may be approximated as
\[ Y_{ijk} = e^{-R_1R_2A_2}, \]  
(8.7)
where the area of the \( A_2 \) may be of order one in string units.
A Majorana mass term for right neutrinos appears only once we impose \( N = 1 \) SUSY on an intersection. As a result the massless scalar superpartners of the \( \bar{F}_R \) fermions, the \( \bar{F}_H^R \)'s appears, allowing the dimension five Majorana mass term for \( \nu_R, F_R F_R^H F_R^H \).
Hence the full Yukawa interaction for the fermionic spectrum is
\[ \lambda_1 F_L \bar{F}_R \langle h_3 \rangle \langle F_R^B \rangle \langle H_2 \rangle + \frac{\lambda_2 F_R F_R^H \langle \bar{F}_R^H \rangle}{M_s}, \]  
(8.8)
where
\[ \lambda_1 \equiv e^{-\frac{R_1R_2A_1}{\alpha'}}, \quad \lambda_2 \equiv e^{-\frac{R_1R_2A_2}{\alpha'}}. \]  
(8.9)
and the Majorana coupling involves the massless scalar superpartners \( \bar{F}_R^H \). The \( \bar{F}_R^H \) has a neutral direction that receives the vev \( \langle H \rangle \). There is no restriction on its vev from first principles and its vev can be anywhere between the scale of electroweak symmetry breaking and \( M_s \).
The Yukawa term
\[ F_L \bar{F}_R \langle h_3 \rangle \langle F_R^B \rangle \langle H_2 \rangle \sim \nu F_L \bar{F}_R \]  
(8.10)
is responsible for the electroweak symmetry breaking. This term generates Dirac masses to up quarks and neutrinos. Thus we get
\[ \lambda_1 F_L \bar{F}_R \langle h_3 \rangle \langle F_R^B \rangle \langle H_2 \rangle \rightarrow (\lambda_1 \nu)(u_i u_j^c + \nu_i \nu_j^c) + (\lambda_1 \bar{\nu}) \cdot (d_i d_j^c + e_i e_j^c), \]  
(8.11)
where we have assumed that
\[ \langle h_3 \rangle \langle F_R^B \rangle \langle H_2 \rangle = \begin{pmatrix} \nu & 0 \\ 0 & \bar{\nu} \end{pmatrix} \]  
(8.12)
These mass relations may be retained at tree level only, since as the models are non-supersymmetric, they will receive higher order corrections. Interestingly from (8.12) we derive the GUT relation \[ 36 \]
\[ m_d = m_e. \]  
(8.13)
as well the unnatural
\[ m_u = m_{N^c \nu}. \quad (8.14) \]

In the case of neutrino masses, the unnatural \( \text{(8.14)} \), associated to the \( \nu - N^c \) mixing, is modified due to the presence of the Majorana term \( \text{(8.8)} \) leading to the see-saw neutrino mass matrix, of an extended Frogatt- Nielsen type mixing light with heavy states,

\[
\begin{pmatrix}
\nu \\
N^c
\end{pmatrix} \times
\begin{pmatrix}
0 & m \\
m & M
\end{pmatrix} \times
\begin{pmatrix}
\nu \\
N^c
\end{pmatrix}, \quad (8.15)
\]

where
\[ m = \lambda_1 \nu. \quad (8.16) \]

After diagonalization the neutrino mass matrix gives us two eigenvalues, the heavy eigenvalue
\[ m_{\text{heavy}} \approx M = \frac{\lambda_2 < H >^2}{M_s}, \quad (8.17) \]
corresponding to the right handed neutrino and the light eigenvalue
\[ m_{\text{light}} \approx \frac{m^2}{M} = \frac{\lambda_1^2}{\lambda_2^2} \times \frac{\nu^2 M_s}{< H >^2} \quad (8.18) \]
corresponding to the left handed neutrino. Values of the parameters giving us values for neutrino masses between 0.1-10 eV, consistent with the observed neutrino mixing in neutrino oscillation measurements, have already been considered in [6]. The analysis is identical and it will not repeated here. We note that the hierarchy of neutrino masses has been investigated by examining several scenarios associated with a light \( \nu_L \) mass including the cases \( \langle H \rangle = |M_s|, \langle H \rangle < |M_s| \). In both cases the hierarchy of neutrino masses is easily obtained.

8.2 Exotic fermion couplings

Up to this point we have shown that all the additional U(1)'s originally present at \( M_s \) have received a heavy mass and have disappeared from the low energy spectrum. Thus the SM gauge group is present at low energy. Also, we have shown that all additional particles created from the non-zero intersections of the extra branes with the colour a-brane and the U(1) leptonic d, e, f, branes receive a mass of the order of the string scale and thus are not present to low energies. The SM at low energy is really attainable only if we show that the fermions beyond those incorporated in \( F_L, \bar{F}_R \) of table receive a non-zero mass. We will now show that this is the case, emphasizing that only the
weak fermion doublets $\chi^{1}_{L}, \chi^{2}_{L}, \chi^{3}_{L}$ receive a light mass between $M_{Z}$ and the order of the electroweak symmetry breaking $v$.

Hence the left handed fermions $\chi^{1}_{L}$ receive a mass of order $v^{2}/M_{s}$ from the coupling

$$(1, 2, 1)(1, 2, 1)e^{-A(\bar{h}_{3})\langle h_{3}\rangle}\langle F_{R}^{H}\rangle\langle H_{2}\rangle \sim 0 \frac{v^{2}}{M_{s}} (1, 2, 1)(1, 2, 1) \text{ (8.19)}$$

that is in representation form

$$(1, 2, 1)(0,1,0,−1,0,0) (1, 2, 1)(0,1,0,−1,0,0) \langle(1, 2, 2)(0,−1,1,0,0,0)\rangle \langle(1, 2, 2)(0,−1,1,0,0,0)\rangle \times \langle(4, 1, 2)(1,0,−1,0,0,0)\rangle \langle(4, 1, 2)(1,0,−1,0,0,0)\rangle \text{ (8.20)}$$

In (8.19) we have incorporated the leading contribution of the worksheet area connecting the six vertices. In the following this contribution will be set for simplicity to one ($A \to 0$).

The left handed fermions $\chi^{2}_{L}$ receive a mass of order $v^{2}/M_{s}$ from the coupling

$$(1, 2, 1)(1, 2, 1)\frac{\langle h_{3}\rangle\langle h_{2}\rangle\langle F_{R}^{H}\rangle\langle H_{2}\rangle\langle s_{B}^{3}\rangle}{M_{s}^{4}} \sim 0 \frac{v^{2}}{M_{s}} (1, 2, 1)(1, 2, 1) \text{ (8.21)}$$

that is in representation form :

$$(1, 2, 1)(0,1,0,−1,0,0) (1, 2, 1)(0,1,0,−1,0,0) \langle(1, 2, 2)(0,−1,1,0,0,0)\rangle \langle(1, 2, 2)(0,−1,1,0,0,0)\rangle \times \langle(4, 1, 2)(1,0,−1,0,0,0)\rangle \langle(4, 1, 2)(1,0,−1,0,0,0)\rangle \text{ (8.22)}$$

The $\chi^{3}_{L}$ doublet fermions receive heavy masses of order $v^{2}/M_{s}$ from the realization of the coupling :

$$(1, 1, 2)(1, 1, 2)\frac{\langle h_{3}\rangle\langle h_{3}\rangle\langle F_{R}^{H}\rangle\langle H_{2}\rangle\langle s_{B}^{3}\rangle}{M_{s}^{3}} \text{ (8.23)}$$

In explicit representation form

$$(1, 1, 2)(0,1,0,0,0,−1) (1, 1, 2)(0,1,0,0,0,−1) \langle(4, 1, 2)(1,0,−1,0,0,0)\rangle \langle(4, 1, 2)(1,0,−1,0,0,0)\rangle \times \langle(1, 1, 1)(0,0,0,0,0,2)\rangle \text{ (8.24)}$$

Thus the left handed fermion weak doublets $\chi^{1}_{L}, \chi^{2}_{L}, \chi^{3}_{L}$ receive a low mass of order $v^{2}/M_{s}$. This is a general prediction of all classes of GUT models based on non-supersymmetric toroidally intersecting D6-branes [2]. In (8.19), (8.21), (8.23) we have assumed $21$ vev’s $< H_{2} > \sim < F_{R}^{H} > \sim M_{s}$. For a general string model the issue of determining the size of the vev’s and whether these fields really receive a vev may be made precise only after the calculation of the effective potential.

$^{21}$Also assume in the following discussion.
The $\chi_R^1$ right handed doublet fermions receive heavy masses of order $M_s$ in the following way:

$$ (1, 1, 2)(1, 1, 2) \left\langle H_2 \right\rangle \left\langle F_R^H \right\rangle \left\langle s_B^1 \right\rangle \over M_s^2 $$

(8.25) In explicit representation form

$$ (1, 1, 2)_{(0,0,1,1,0,0)} (1, 1, 2)_{(0,0,1,1,0,0)} \left\langle (4, 1, 2)_{(-1,0,-1,0,0,0)} \right\rangle \left\langle (4, 1, 2)_{(1,0,-1,0,0,0)} \right\rangle \times \left\langle (1, 1, 1)_{(0,0,0,-2,0,0)} \right\rangle \quad (8.26) $$

The $\chi_R^2$ right handed doublet fermions receive heavy masses of order $M_s$ from the following coupling

$$ (1, 1, 2)(1, 1, 2) \left\langle H_2 \right\rangle \left\langle F_R^H \right\rangle \left\langle s_B^2 \right\rangle \over M_s^2 $$

(8.27) In explicit representation form

$$ (1, 1, 2)_{(0,0,1,0,1,0)} (1, 1, 2)_{(0,0,1,0,1,0)} \left\langle (4, 1, 2)_{(-1,0,-1,0,0,0)} \right\rangle \left\langle (4, 1, 2)_{(1,0,-1,0,0,0)} \right\rangle \times \left\langle (1, 1, 1)_{(0,0,0,0,-2,0)} \right\rangle \quad (8.28) $$

The $\chi_R^3$ right handed doublet fermions receive heavy masses of order $M_s$ in the following way:

$$ (1, 1, 2)(1, 1, 2) \left\langle H_2 \right\rangle \left\langle F_R^H \right\rangle \left\langle s_B^3 \right\rangle \over M_s^2 $$

(8.29) In explicit representation form

$$ (1, 1, 2)_{(0,0,1,0,0,1)} (1, 1, 2)_{(0,0,1,0,0,1)} \left\langle (4, 1, 2)_{(-1,0,-1,0,0,0)} \right\rangle \left\langle (4, 1, 2)_{(1,0,-1,0,0,0)} \right\rangle \times \left\langle (1, 1, 1)_{(0,0,0,0,0,-2)} \right\rangle \quad (8.30) $$

The 6-plet fermions, $\omega_L$, receive a mass term of order $M_s$ from the coupling,

$$ (\bar{6}, 1, 1)(\bar{6}, 1, 1) \left\langle H_1 \right\rangle \left\langle F_R^H \right\rangle \left\langle H_1 \right\rangle \left\langle F_R^H \right\rangle \over M_s^3 $$

(8.31) where we have made use of the $SU(4)$ tensor products $6 \otimes 6 = 1 + 15 + 20$, $4 \otimes 4 = 6 + 10$ and have defined

$$ (\bar{6}, 1, 1)_{(-2,0,0,0,0,0)} (\bar{6}, 1, 1)_{(-2,0,0,0,0,0)} \left\langle (4, 1, 2)_{(1,0,1,0,0,0)} \right\rangle \left\langle (4, 1, 2)_{(1,0,1,0,0,0)} \right\rangle \times \left\langle (4, 1, 2)_{(1,0,1,0,0,0)} \right\rangle \left\langle (4, 1, 2)_{(1,0,1,0,0,0)} \right\rangle \quad (8.32) $$

The 10-plet fermions $z_R$ receive a heavy mass of order $M_s$ from the coupling

$$ (10, 1, 1)(10, 1, 1) \left\langle F_R^H \right\rangle \left\langle F_R^H \right\rangle \left\langle H_2 \right\rangle \left\langle H_2 \right\rangle \over M_s^3 $$

(8.33)
and we have used the tensor product representations of $SU(4)$, $10 \otimes 10 = 20 + 35 + 45$, $20 \otimes 4 = 15 + 20$, $\bar{20} \otimes 4 = 6 + 10$, $10 \otimes 4 = 4 + 36$, $4 \otimes 4 = 1 + 15$. Explicitly, in representation form,

\[(10, 1, 1) (2, 0, 0, 0, 0, 0) (10, 1, 1) (2, 0, 0, 0, 0, 0) \langle (\bar{4}, 1, 2)_{(-1,0,1,0,0,0)} \rangle \langle (\bar{4}, 1, 2)_{(-1,0,1,0,0,0)} \rangle \times \langle (\bar{4}, 1, \bar{2})_{(-1,0,-1,0,0,0)} \rangle \langle (\bar{4}, 1, \bar{2})_{(-1,0,-1,0,0,0)} \rangle \quad (8.34)\]

Thus at low energies of order $M_Z$ only the SM remains.

9 Conclusions

Based on the intersecting D6-brane constructions of [2], in [6] we constructed the first examples of string GUT models which break at exactly the SM at low energy without any additional group factors and/or exotic massless matter. The classes of GUT models we have been considering recently [6, 7, 8] and at the present work, have as their low energy theory in energies of order $M_Z$ the Standard model. Their common characteristic is that they represent deformations around the basic intersection structure of the Quark and Lepton structure,

\[I_{ab} = 3, \quad I_{ac^*} = -3 . \quad (9.1)\]

Thus they all share the same intersection numbers along the ‘baryonic’ $a$ and the left and right ‘weak’ $b$ and $c$, D6 branes.

Interestingly the GUT constructions have a number of features independently of the number of D6-stacks that they are defined originally. These general characteristics include:

- The prediction of low mass ($\sim \frac{\sqrt{s}}{M_Z}$) weak left handed doublets ($\chi^i_L$, $i = 1, 2, 3$) with mass between $M_Z$ and 246 GeV. This is a universal feature and appears at GUT constructions with various number of stacks [6, 7, 8].

This result makes automatic the existence of a low scale in the models, e.g. below 650 GeV which makes intersecting D6-brane GUTs directly testable at present or near feature accelerators.

- The conditions for some of the U(1)’s to survive massless the Green-Schwarz mechanism are equivalent to the conditions originating from the existence of N=1 supersymmetry in some sectors. The preservation of N=1 SUSY is necessary in some open string sectors in order to allow the generation of gauge singlets making
massive the unwanted exotic fermions, e.g. $\chi_{1}^{R}$, $\chi_{2}^{R}$, $\chi_{3}^{R}$, $s_{1}^{R}$, $s_{2}^{R}$, $s_{3}^{R}$ of table (1), and to generate the Majorana mass term.

- The presence of extra branes needed to cancel the RR tadpoles, may be arranged so that it has non-zero intersection numbers with the colour brane and the rest, but b, c, of the branes. This feature is to be contrasted with models, with only the SM at low energy build from a Standard like configuration at the string scale $M_{s}$ contained between four- [3] five- [4] and six- [5] stacks of D6-branes (there is no configuration of D6-branes able to give only the SM at low energy beyond 6-stacks), where the extra branes have no intersection number with the rest of the branes. The additional fermions created are made always massive by allowing the presence of N=1 SUSY in sectors in the form $jj^{*}$ and also in sectors having extra and U(1) leptonic branes.

- Even though the models are overall non-supersymmetric they contain N=1 SUSY preserving sectors. These sectors are necessary to be implemented into the theory as without them it will not be possible to generate gauge singlets breaking the surviving massless the Green-Schwarz mechanism U(1)’s. Equally important, in their absence we could not been able to generate a Majorana mass term for right handed neutrinos.

The present non-supersymmetric constructions, if the angle stabilization conditions of appendix A hold are free of tachyons, however NSNS tadpoles remain, thus leaving the full question of stability in these models an open question. We should nevertheless remember that even supersymmetric constructions are free of NSNS tadpoles, supersymmetry breaking may create an unexpected cosmological constant. These tadpoles could be removed in principle by background redefinition in terms of wrapped metrics [38] or by orbifolding as was suggested in [39] and freezing the moduli to discrete values.

However, it is intriguing that we have found intersecting brane constructions, in the absence of a dynamical mechanism which can select a particular string vacuum, that offer the possibility to obtain vacua with just the observable Standard Model spectrum and gauge interactions at low energy.

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10 Appendix I

In this appendix we list the conditions under which the PS-III model classes of intersecting D6-branes discussed in this work are tachyon free. Note that the conditions are expressed in terms of the angles defined in (6.2). We have included the contributions from the sectors $ab^\ast$, $ac$, $bd$, $cd^\ast$, $be$, $ce^\ast$, $bf$, $cf^\ast$. We have not included the tachyon free conditions from the sectors $dd^\ast$, $ee^\ast$, $ff^\ast$, as the latter conditions will be shown to be equivalent to the presence of $N=1$ supersymmetry in these sectors.

\[
\begin{align}
-(\frac{\pi}{2} + \theta_1) &+ \theta_2 + 2\theta_3 \geq 0 \\
-(\frac{\pi}{2} - \tilde{\theta}_1) &+ \theta_2 + 2\theta_3 \geq 0 \\
-(\theta_1 - \frac{\pi}{2}) &+ \tilde{\theta}_2 + 2\theta_3 \geq 0 \\
-(\frac{\pi}{2} + \tilde{\theta}_1) &+ \tilde{\theta}_2 + 2\theta_3 \geq 0 \\
-(\theta_1 - \frac{\pi}{2}) &+ \theta'_2 + 2\theta_3 \geq 0 \\
-(\frac{\pi}{2} + \tilde{\theta}_1) &+ \theta'_2 + 2\theta_3 \geq 0 \\
-(\frac{\pi}{2} + \theta_1) &+ \theta'_2 + 2\theta_3 \geq 0
\end{align}
\]

\[(10.1)\]

\[
\begin{align}
(\frac{\pi}{2} + \theta_1) &- \theta_2 + 2\theta_3 \geq 0 \\
(\frac{\pi}{2} - \tilde{\theta}_1) &- \theta_2 + 2\theta_3 \geq 0 \\
(\theta_1 - \frac{\pi}{2}) &- \tilde{\theta}_2 + 2\theta_3 \geq 0 \\
(\frac{\pi}{2} + \tilde{\theta}_1) &- \tilde{\theta}_2 + 2\theta_3 \geq 0 \\
(\theta_1 - \frac{\pi}{2}) &- \theta'_2 + 2\theta_3 \geq 0 \\
(\frac{\pi}{2} + \tilde{\theta}_1) &- \theta'_2 + 2\theta_3 \geq 0 \\
(\frac{\pi}{2} + \theta_1) &- \theta'_2 + 2\theta_3 \geq 0
\end{align}
\]

\[(10.2)\]

\[
\begin{align}
(\frac{\pi}{2} + \theta_1) &+ \theta_2 - 2\theta_3 \geq 0 \\
(\frac{\pi}{2} - \tilde{\theta}_1) &+ \theta_2 - 2\theta_3 \geq 0 \\
(\theta_1 - \frac{\pi}{2}) &+ \tilde{\theta}_2 - 2\theta_3 \geq 0 \\
(\frac{\pi}{2} + \tilde{\theta}_1) &+ \tilde{\theta}_2 - 2\theta_3 \geq 0 \\
(\theta_1 - \frac{\pi}{2}) &+ \theta'_2 - 2\theta_3 \geq 0 \\
(\frac{\pi}{2} + \tilde{\theta}_1) &+ \theta'_2 - 2\theta_3 \geq 0
\end{align}
\]
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