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Deformed intersecting D6-brane GUTS I

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Abstract: By employing D6-branes intersecting at angles in $D = 4$ toroidal orientifolded type IIA strings, we construct *four stack string GUT models* (PS-I class), that contain at low energy *exactly the three generation Standard model* with no extra matter and/or extra gauge group factors. These classes of models are based on the Pati-Salam (PS) gauge group $SU(4)_C \times SU(2)_L \times SU(2)_R$. 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 class) of four stack PS GUTS. Namely, there are no colour triplet couplings to mediate proton decay and proton is stable as baryon number is a gauged symmetry. Neutrinos get masses of the correct sizes. Also the mass relation $m_e = m_d$ at the GUT scale is recovered.
Moreover, we clarify the novel role of *extra* branes, the latter having non-trivial intersection numbers with quarks and leptons and creating scalar singlets, needed for the satisfaction of RR tadpole cancellation conditions. The presence of $N = 1$ supersymmetry in sectors involving the *extra* branes is equivalent to the, model dependent, orthogonality conditions of the U(1)'s surviving massless the generalized Green-Schwarz mechanism. The use of *extra* branes creates mass couplings that predict the appearance of light fermion doublets up to the scale of electroweak scale symmetry breaking.

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

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1. Introduction

Major problems of string theory include among others the hierarchy of scale and particle masses after supersymmetry breaking. These phenomenological issues have by far been explored in the context of construction of semirealistic supersymmetric models of weakly coupled $N = 1$ (orbifold) compactifications of the heterotic string theories [1]. In these theories one of the unsolved problems was the fact that the string scale which is of the order of 10^{18} GeV was in clear disagreement with the observed unification of gauge coupling constants in the MSSM of 10^{16} GeV. The latter problem remains a mystery even though the observed discrepancy between the two high scales was attributed[1] to the presence of the $N = 1$ string threshold corrections to the gauge coupling constants [3].

On the contrary in type I models, the string scale, which is a free parameter, can be lowered in the TeV range [4] thus suggesting that non-SUSY models with a string scale in the TeV region is a viable possibility. In this spirit, recently some new constructions have appeared in a type I string vacuum background which use intersecting branes [2] and give four dimensional non-supersymmetric models.

[1] Among other options.

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In these open string models [2] the use of background fluxes in a D9 brane type I background[2] breaks supersymmetry on the brane and gives chiral fermions with an even number of generations [2]. The fermions on those models get localized in the intersections between branes [5, 6]. The introduction of a quantized background NS-NS B field [7]–[9], that makes the tori tilted gives rise to semirealistic models with three generations [10]. It should be noted that these backgrounds are T-dual to models with magnetic deformations [11]. Additional non-SUSY constructions in the context of intersecting branes, from IIB orientifolds, consisting of getting at low energy the standard model spectrum with extra matter and additional chiral fermions were derived in [12]. Also SUSY constructions in the context of intersecting branes were considered in [13]. In addition, constructions involving intersecting branes in compact Calabi-Yau spaces were discussed in [14], while intersecting brane constructions in the context of non-compact Calabi-Yau spaces were considered in [15]. For some other work in the context of intersecting branes see [16]–[21]. For some recent attempts to construct,[3] non-SUSY GUT models in the context of intersecting branes see [23]–[25].

Furthermore, an important step was taken in [26], by showing how to construct the standard model (SM) spectrum together with right handed neutrinos in a systematic way. The authors considered, as a starting point, IIA theory compactified on T^6 [2] assigned with an orientifold product $\Omega \times R$, where Ω is the worldsheet parity operator and R is the reflection operator with respect to one of the axis of each tori. In this case, the four stack D6-branes contain Minkowski space and each of the three remaining dimensions is wrapped up on a different T^2 torus. In this construction 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 classes of models with just the SM at low energy and having five stacks [27] and six stacks [28] of D6-branes at the string scale. The models of [27, 28] are build as deformations of the QCD intersection numbers, namely they are build around the left and right handed quarks intersection numbers. Also, they hold exactly the same phenomenological properties of [26]. Also, these models — due to the presence of $N = 1$ supersymmetric sectors necessary for the breaking of the U(1)'s surviving massless the Green-Schwarz mechanism — predict the unique existence of one supersymmetric partner of the right neutrino and two supersymmetric partners of the right neutrino in the five and six stack SM's of [27, 28] respectively.

In addition, in [29] we presented the first examples of classes of string derived GUT models (PS-A class) that break completely to the SM at low energies. The models are developed in the same D6-brane backgrounds as the SM's of [26]. They are based in the Pati-Salam (PS) [30] GUT structure G_{422}, $\mathrm{SU}(4)_c \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$. The models predict uniquely the existence of light weak fermion doublets with energy between the range $90 - 246\,\mathrm{GeV}$, that is they can be directly tested at present or future accelerators. We also note another recent construction with D5 branes intersecting at angles on an orientifold of $T^2 \times T^4/\mathbb{Z}_N$ [31]. In this case, four stack models of D5 branes give just the SM at low energy. A

[2] In the T-dual language these backgrounds are represented by D6 branes wrapping 3-cycles on a dual torus and intersecting each other at certain angles.

[3] Non-SUSY GUTS in the context of type IIB with branes on singularities see [22].

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full study of the latter models including an extension to five and six stack SM constructions with just the SM at low energies is performed in [32]. It appears [32] that there is a special class of D5 vacua in four, five and six stacks of SM embeddings that have the same low energy effective theory suggesting that these theories are connected at the infrared.

The purpose of this work is to present further three generation four stack string models (PS-I class) that are based on the PS G_{422} group, and contain at low energy exactly the standard model spectrum, namely $SU(3)_C \times SU(2)_L \times U(1)_Y$, without any extra chiral fermions and/or extra gauge group factors.

We will exhibit the systematics of using *extra* branes with non-trivial intersection numbers with the color and leptonic branes. The use of these *extra* branes will serve as a novel mechanism of scalar singlet generation and breaking of the U(1)'s surviving massless the Green-Schwarz mechanism. The presence of the *extra* brane mechanism will be applied to both PS-A and PS-I classes of PS GUT models. We should note that these *extra* branes are quite different from the use of hidden branes, used in non-GUT based D6-brane model building examined in [26]–[28]. In the latter models the use of additional D6 branes needed to satisfy the RR tadpole cancellation conditions did not charge the SM chiral fermions and thus these extra branes could be characterized as hidden one's. In the present context the use of *extra* D6 branes cannot be characterized as hidden, as there are fields charged under the *extra* branes symmetry group.

The four-dimensional classes of models we study are non-supersymmetric intersecting brane constructions. The basic structure behind the models includes D6-branes intersecting each other at non-trivial angles, in an orientifolded factorized six-torus, where O_6 orientifold planes are on top of D6-branes. The proposed classes of models have some distinctive features:

- the models start, we neglect for the time being the presence of *extra* branes, with a gauge group at the string scale $U(4) \times U(2) \times U(2) \times U(1)$. At the scale of symmetry breaking of the left-right symmetry, M_{GUT}, the initial symmetry group breaks to the the standard model $SU(3)_C \times SU(2)_L \times U(1)_Y$ augmented with an extra anomaly free U(1) symmetry. The additional U(1) symmetry breaks by the vev of charged singlet scalars to the SM itself at a scale set by its vev. The singlets responsible for breaking the U(1) symmetry are obtained by demanding that certain open string sectors of the non-SUSY model respect $N = 1$ supersymmetry.

- neutrinos gets a mass of the right order, consistent with the LSND oscillation experiments, from a see-saw mechanism of the Frogatt-Nielsen type. The structure of Yukawa couplings involved in the see-saw mechanism supports the smallness of neutrino masses thus generating a hierarchy in consistency with neutrino oscillation experiments.

- proton is stable 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. Thus a gauged baryon number provides a natural explanation for proton stability. As in the models of [26]–[29] and [31, 32] the baryon number

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associated U(1) gauge boson becomes massive through its couplings to Green-Schwarz mechanism. That has an an immediate effect that baryon number is surviving as a global symmetry to low energies providing for a natural explanation for proton stability in general brane-world scenarios.

- the model uses small Higgs representations in the adjoint to break the PS symmetry, instead of using large Higgs representations, e.g. 126 like in the standard SO(10) models.

- the bidoublet Higgs fields h responsible for electroweak symmetry breaking do not get charged under the global U(1) and thus lepton number is not broken at the standard model.

- *Extra* branes, with non-trivial intersection numbers with the colour a- and the leptonic d-brane, in addition to the imposition of N=1 SUSY in sectors coming from the intersections of the hidden with the leptonic branes, are being used to engineer the presence of only the SM at low energy. The *extra* branes are added in single pieces, each one being associated with a single U(1), e.g. in the presence of two (2) *extra* U(1) branes the number of extra U(1)'s, which survive massless the Green-Schwarz mechanism, is three, and the theory looks in practical terms like a six-stack model.

The paper is organized as follows. In section 2 we describe the general rules for building chiral models in orientifolded T^6 compactifications and the possible open string sectors. In section 3, we discuss the basic fermion and scalar structure of the PS-I class of models that will mainly focus in this work. In section 4, we make a parenthesis in our study and discuss the role of the *extra* branes in the PS-A class of models of [29]. The methods described in creating singlets scalars fields will serve us as a prototype for an application to the PS-I class of models. In section 5, we discuss the parametric solutions to the RR tadpoles for the PS-I class of models where we will be focusing our attention from now on. In section 6 we discuss the cancellation of U(1) anomalies in the presence of a generalized Green-Schwarz (GS) mechanism and *extra* U(1) D6 branes. In subsection 7.1 we discuss the conditions for the absence of tachyons in the models as well describing the PS breaking Higgs and the electroweak symmetry breaking Higgs fields. In subsection 7.2 we discuss the presence of N=1 supersymmetric sectors and *extra* sector branes. The presence of N=1 SUSY creates scalar singlets which are necessary to make some unwanted fermions massive enough to disappear from the low energy spectrum. In subsection 7.3 we discuss the breaking of the surviving the Green-Schwarz mechanism massless U(1)'s with the use of singlets coming from the non-trivial $N = 1$ SUSY intersections of the *extra* branes and leptonic branes. In section 8 we examine the problem of neutrino masses. We also show that all additional exotic fermions beyond those of SM become massive and disappear from the low energy spectrum. In this section, we describe in detail how the presence of supersymmetry in particular sectors of the theory realizes the particular couplings taking part in the see-saw mechanism. Section 9 contains our conclusions. Finally, appendix, includes the conditions for the absence of tachyonic modes in the spectrum of the PS-I class of models.

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2. Model structure and the rules of computing the spectrum

Next, we describe the construction of the PS classes of models. They are 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. 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. The branes are not paralled to the orientifold planes. We assume that the $D6_a$-branes are wrapping 1-cycles (n^i_a, m^i_a) along each of the T^2 torus of the factorized T^6 torus, that is we assume $T^6 = T^2 \times T^2 \times T^2$.

In order to build a general PS model we consider four stacks of D6-branes giving rise to their world-volume to an initial gauge group $U(4)_c \times U(2)_L \times U(2)_R \times U(1)$ at the string scale. In addition, we consider the addition of NS B-flux, which makes the tori tilted, and leads to the effective tilted wrapping numbers,

$$\left(n^i, m = \tilde{m}^i + \frac{n^i}{2}\right); \qquad n,\, \tilde{m} \,\in\, \mathbb{Z}\,, \tag{2.1}$$

allowing semi-integer values for the m-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,

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

each $D6_a$-brane 1-cycle, must be accompanied by its $\Omega\mathcal{R}$ partner $(n^i_a, -m^i_a)$.

Chiral fermions are obtained by stretched open strings between intersecting D6-branes [6]. The chiral spectrum of the models is obtained after solving simultaneously the intersection constraints coming from the existence of the different sectors together with the RR tadpole cancellation conditions.

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 α, β, by $\alpha^\star, \beta^\star$, 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 ΩR symmetry this sector is mapped to its image, $\alpha^\star\beta^\star + \beta^\star\alpha^\star$ sector. The number, $I_{\alpha\beta}$, of chiral fermions in this sector, transforms in the bifundamental representation $(N_\alpha, \bar{N}_\alpha)$ of $U(N_\alpha) \times U(N_\beta)$, and reads

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

  where $I_{\alpha\beta}$ is the intersection number of the wrapped cycles. Note that the sign of $I_{\alpha\beta}$ denotes the chirality of the fermion and with $I_{\alpha\beta} > 0$ we denote left handed fermions. Negative multiplicity denotes opposite chirality;

- the $\alpha\alpha$ sector: it involves open strings stretching on a single stack of $D6_\alpha$ branes. Under the ΩR 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 appear. This sector is of no importance to us as we are interested in unitary groups;

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- the $\alpha\beta^\star + \beta^\star\alpha$ sector: It involves chiral fermions transforming into the (N_α, N_β) representation with multiplicity given by

$$I_{\alpha\beta^\star} = -(n_\alpha^1 m_\beta^1 + m_\alpha^1 n_\beta^1)(n_\alpha^2 m_\beta^2 + m_\alpha^2 n_\beta^2)(n_\alpha^3 m_\beta^3 + m_\alpha^3 n_\beta^3)\,. \tag{2.4}$$

  Under the ΩR symmetry transforms to itself.

- the $\alpha\alpha^\star$ sector: under the ΩR symmetry is transformed to itself. From this sector the invariant intersections will give $8m_\alpha^1 m_\alpha^2 m_\alpha^3$ fermions in the antisymmetric representation and the non-invariant intersections that come in pairs provide us with $4m_\alpha^1 m_\alpha^2 m_\alpha^3 (n_\alpha^1 n_\alpha^2 n_\alpha^3 - 1)$ additional fermions in the symmetric and antisymmetric representation of the $\mathrm{U}(N_\alpha)$ gauge group.

Any vacuum derived from the previous intersection number constraints of the chiral spectrum is subject to constraints coming from RR tadpole cancellation conditions [2]. That requires cancellation of D6-branes charges,[4] wrapping on three cycles with homology $[\Pi_a]$ and O6-plane 7-form charges wrapping on 3-cycles with homology $[\Pi_{O_6}]$. In formal terms, the RR tadpole cancellation conditions in terms of cancellations of RR charges in homology, read:

$$\sum_a N_a[\Pi_a] + \sum_{\alpha^\star} N_{\alpha^\star}[\Pi_{\alpha^\star}] - 32[\Pi_{O_6}] = 0\,. \tag{2.5}$$

Explicitly, the RR tadpole conditions read:

$$\begin{aligned}
\sum_a N_a n_a^1 n_a^2 n_a^3 &= 16\,,\\
\sum_a N_a m_a^1 m_a^2 n_a^3 &= 0\,,\\
\sum_a N_a m_a^1 n_a^2 m_a^3 &= 0\,,\\
\sum_a N_a n_a^1 m_a^2 m_a^3 &= 0\,.
\end{aligned} \tag{2.6}$$

That ensures absence of non-abelian gauge anomalies but not the reverse. A comment is in order. It is important to notice that the RR tadpole cancellation condition can be understood as a constraint that demands for each gauge group the number of fundamentals to be equal to the number of bifundamentals. As a general rule to D-brane model building, by considering a stacks of D-brane configurations with $N_a, a = 1, \ldots, N$, paralled branes, the gauge group appearing is in the form $\mathrm{U}(N_1) \times \mathrm{U}(N_2) \times \cdots \times \mathrm{U}(N_a)$. Effectively, each $\mathrm{U}(N_i)$ factor will give rise to an $\mathrm{SU}(N_i)$ charged under the associated $\mathrm{U}(1_i)$ gauge group factor that appears in the decomposition $\mathrm{SU}(N_a) \times \mathrm{U}(1_a)$. A brane configuration with the unique minimal PS particle content such that intersection numbers, tadpole conditions and various phenomenological requirements including the absence of exotic representations are accommodated, can be obtained by considering initially four stacks of D6-branes yielding an

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

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initial $\mathrm{U}(4)_a \times \mathrm{U}(2)_b \times \mathrm{U}(2)_c \times \mathrm{U}(1)_d$ gauge group equivalent to an $\mathrm{SU}(4)_a \times \mathrm{SU}(2)_b \times \mathrm{SU}(2)_b \times \mathrm{U}(1)_a \times \mathrm{U}(1)_b \times \mathrm{U}(1)_c \times \mathrm{U}(1)_d$. Thus, in the first instance, we can identify, without loss of generality, $\mathrm{SU}(4)_a$ as the $\mathrm{SU}(4)_c$ colour group that its breaking could induce the usual SU(3) colour group of strong interactions, the $\mathrm{SU}(2)_b$ with $\mathrm{SU}(2)_L$ of weak interactions and $\mathrm{SU}(2)_c$ with $\mathrm{SU}(2)_R$. Note that the condition to satisfy the RR tadpole cancellation conditions will force us to add the presence of *extra* branes.

3. The basic fermion structure

The basic PS-I class of models that we will center our attention in this work, will be a three family non-supersymmetric GUT model with the left-right symmetric Pati-Salam model structure $\mathrm{SU}(4)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$. The open string background on which the models will be build will be intersesting D6-branes wrapping on 3-cycles of decomposable toroidal (T^6) orientifolds of type IIA in four dimensions.

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

$$\begin{aligned} F_L &= (4,2,1) = q\left(3,2,\frac{1}{6}\right) + l\left(1,2,-\frac{1}{2}\right) \equiv (u,\ d,\ l)\,, \\ \bar{F}_R &= (\bar{4},1,2) = u^c\left(\bar{3},1,-\frac{2}{3}\right) + d^c\left(\bar{3},1,\frac{1}{3}\right) + e^c(1,1,1) + N^c(1,1,0) \equiv (u^c, d^c, l^c)\,, \end{aligned} \tag{3.1}$$

where the quantum numbers on the right hand side of (3.1) are with respect to the decomposition of the $\mathrm{SU}(4)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ under the $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge group and $l = (\nu, e)$ is the standard left handed lepton doublet, $l^c = (N^c, e^c)$ are the right handed leptons. Also the assignment of the accommodation of the quarks and leptons into the representations $F_L + \bar{F}_R$ is the one appearing in the spinorial decomposition of the 16 representation of $SO(10)$ under the PS gauge group. A set of useful fermions appear also in the model

$$\begin{aligned} \chi_L^1 &= (1,\bar{2},1)\,, \qquad & \chi_R^1 &= (1,1,\bar{2})\,, \\ \chi_L^2 &= (1,\bar{2},1)\,, \qquad & \chi_R^2 &= (1,1,\bar{2})\,. \end{aligned} \tag{3.2}$$

These fermions are a general prediction of left-right symmetric theories as the existence of these representations follows from RR tadpole cancellation conditions. The symmetry breaking of the left-right PS symmetry at the M_{GUT} scale[5] proceeds through the representations of the set of Higgs fields,

$$H_1 = (\bar{4},1,\bar{2})\,, \qquad H_2 = (4,1,2)\,, \tag{3.3}$$

where,

$$H_1 = (\bar{4},1,\bar{2}) = u_H\left(\bar{3},1,\frac{2}{3}\right) + d_H\left(\bar{3},1,-\frac{1}{3}\right) + e_H(1,1,-1) + \nu_H(1,1,0)\,. \tag{3.4}$$

[5]In principle this scale could be as high as the string scale.

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Fields	Intersection	$\bullet$ $SU(4)_C \times SU(2)_L \times SU(2)_R$ $\bullet$	Q_a	Q_b	Q_c	Q_d
F_L	$I_{ab^*} = 3$	$3 \times (4, 2, 1)$	1	1	0	0
$\bar{F}_R$	$I_{ac} = -3$	$3 \times (\overline{4}, 1, 2)$	-1	0	1	0
χ_L^1	$I_{bd} = -6$	$6 \times (1, \overline{2}, 1)$	0	-1	0	1
χ_R^1	$I_{cd} = -6$	$6 \times (1, 1, \overline{2})$	0	0	-1	1
χ_L^2	$I_{bd^*} = -6$	$6 \times (1, \overline{2}, 1)$	0	-1	0	-1
χ_R^2	$I_{cd^*} = -6$	$6 \times (1, 1, \overline{2})$	0	0	-1	-1
ω_L	I_{aa^*}	$6\beta^2 \times (6, 1, 1)$	2	0	0	0
z_R	I_{aa^*}	$6\beta^2 \times (\overline{10}, 1, 1)$	-2	0	0	0

Table 1: Fermionic spectrum of the $SU(4)_C \times SU(2)_L \times SU(2)_R$, PS-I class of models together with U(1) charges. Note that we have not included fermions coming from the presence of sectors involving the *extra* branes.

The electroweak symmetry breaking is delivered through the bi-doublet Higgs fields h_i $i = 1, 2$, fields in the representations

$$h_1 = (1, 2, 2)\,, \qquad h_2 = (1, \bar{2}, \bar{2})\,. \tag{3.5}$$

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

$$\bar{F}_R^H = (\bar{4}, 1, 2) = u_H^c\left(\bar{3}, 1, -\frac{4}{6}\right) + d_H^c\left(\bar{3}, 1, \frac{1}{3}\right) + e_H^c(1,1,1) + N_H^c(1,1,0) \equiv (u_H^c, d_H^c, l_H^c)\,. \tag{3.6}$$

The latter fields[6] confirm a property shared by all vacua coming from these type IIA constructions. That is the replication of massless fermion spectrum by an equal number of massive particles in the same representations and with the same quantum numbers. This is the basic fermionic structure appearing in the PS models that we have considered in [29] and will be appearing later in this work. Also, a number of charged exotic fermion fields, which receive a string scale mass, appear

$$6(6, 1, 1)\,, \qquad 6(\overline{10}, 1, 1)\,. \tag{3.7}$$

The complete accommodation of the fermion structure of the PS-I classes of models can be seen in table 1.

At this point, before we start discussing the issues of tadpole cancellation for the new PS-I classes of models, we will turn our attention to the PS-A classes of models of [29] as we want to clarify the role of the presence of the *extra* D6 branes in these models. The same methodology will be applied later in the present PS-I classes of models.

4. *Extra* brane engineering — the case of PS-A class of models

The four stack PS-A class of models of [29] are based on the same $U(4) \times U(2)_L \times U(2)_L \times U(1)_d$ gauge structure at the string scale as the PS-I models that we will be discussing extensively in this work. The fermionic field spectrum of PS-A models is the one appearing

[6]Are replicas of the fermion fields appearing in the intersection ac and receive a vev.

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Fields	Intersection	• $SU(4)_C \times SU(2)_L \times SU(2)_R$ •	Q_a	Q_b	Q_c	Q_d
F_L	$I_{ab^*} = 3$	$3 \times (4, 2, 1)$	1	1	0	0
$\bar{F}_R$	$I_{ac} = -3$	$3 \times (\overline{4}, 1, 2)$	-1	0	1	0
χ_L	$I_{bd} = -12$	$12 \times (1, \overline{2}, 1)$	0	-1	0	1
χ_R	$I_{cd} = -12$	$12 \times (1, 1, \overline{2})$	0	0	-1	-1
ω_L	I_{aa^*}	$12\beta^2 \tilde{\epsilon} \times (6, 1, 1)$	$2\tilde{\epsilon}$	0	0	0
z_R	I_{aa^*}	$6\beta^2 \tilde{\epsilon} \times (\overline{10}, 1, 1)$	$-2\tilde{\epsilon}$	0	0	0
s_L	I_{dd^*}	$24\beta^2 \tilde{\epsilon} \times (1, 1, 1)$	0	0	0	$-2\tilde{\epsilon}$

Table 2: Fermionic spectrum of the $SU(4)_C \times SU(2)_L \times SU(2)_R$, PS-A class of models together with U(1) charges.

N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 4$	$(0, \epsilon)$	$(n_a^2, 3\epsilon\beta_2)$	$(\tilde{\epsilon}, \tilde{\epsilon}/2)$
$N_b = 2$	$(-1, \epsilon m_b^1)$	$(1/\beta_2, 0)$	$(\tilde{\epsilon}, \tilde{\epsilon}/2)$
$N_c = 2$	$(1, \epsilon m_c^1)$	$(1/\beta_2, 0)$	$(\tilde{\epsilon}, -\tilde{\epsilon}/2)$
$N_d = 1$	$(0, \epsilon)$	$(n_d^2, 6\epsilon\beta_2)$	$(-2\tilde{\epsilon}, \tilde{\epsilon})$
$\vdots$	$\vdots$	$\vdots$	$\vdots$
N_h	$(1/\beta_1, 0)$	$(1/\beta_2, 0)$	$(2\tilde{\epsilon}, 0)$

Table 3: Tadpole solutions for PS-A type models with D6-branes wrapping numbers 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 two integer parameters, n_a^2, n_d^2, the NS-background β_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_b^1, m_c^1. Note the presence of the N_h *extra* U(1) branes.

in table 2 and the solution to the RR tadpole cancellation conditions is given in table 3. In the PS-A class of models there are present the PS breaking Higgs fields H_1, H_2 from (3.4) as well the electroweak Higgs fields (3.5).

Also we have defined the angles:

$$\theta_1 = \frac{1}{\pi}\cot^{-1}\frac{R_1^{(1)}}{m_b^1 R_2^{(1)}}; \qquad \theta_2 = \frac{1}{\pi}\cot^{-1}\frac{n_a^2 R_1^{(2)}}{3\beta_2 R_2^{(2)}}; \qquad \theta_3 = \frac{1}{\pi}\cot^{-1}\frac{2R_1^{(3)}}{R_2^{(3)}};$$
$$\tilde{\theta_2} = \frac{1}{\pi}\cot^{-1}\frac{n_d^2 R_1^{(1)}}{6\beta_2 R_2^{(1)}}; \qquad \tilde{\theta_1} = \frac{1}{\pi}\cot^{-1}\frac{R_1^{(1)}}{m_c^1 R_2^{(1)}}. \tag{4.1}$$

The presence of $N = 1$ supersymmetry at the sectors ac, $dd^\star$ is compatible with the condition

$$-\tilde{\vartheta}_1 = \tilde{\vartheta}_2 = \vartheta_2 = \vartheta_3 = \frac{\pi}{4}. \tag{4.2}$$

Note that in table 3 we have added the presence of an arbitrary number of *extra* D6 branes which have a non-zero intersection number with the colour a-brane and the leptonic d-brane and thus additional massless fermions seem to be produced.[7] The *extra* N_h branes

[7]The latter issue was not clarified in [29] and we will analyze it here.

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are necessary to cancel the first tadpole condition, the latter giving

$$N_h \frac{2\tilde{\epsilon}}{\beta_1 \beta_2} = 16\,. \tag{4.3}$$

Also, the third tadpole condition in (2.6) gives the constraint

$$2n_a^2 + n_d^2 + \frac{1}{\beta_2}(m_b^1 - m_c^1) = 0\,. \tag{4.4}$$

The presence of N_h branes correspond to the presence of an additional $\mathrm{U}(N_h)$ group at the string scale. However, for our purposes instead of adding an $\mathrm{U}(N_h)$ stack of *extra* D6 branes to cancel the RR tadpoles, we will consider adding N_h $\mathrm{U}(1)$ *extra* branes positioned at $(1/\beta_1, 0)(1/\beta_2, 0)(2, 0)$ each. Note that if we had chosen $\tilde{\epsilon} = -1$, we may have added D6 anti-branes at the same position. That is in order to show that the new massless fields, appearing from the non-zero intersections of each $\mathrm{U}(1)$ *extra* D6-brane, get a mass, it is enough to consider only one of the N_h replicas, the latter we call it $N_{\hat{h}_1}$. Thus due to the non-zero intersection numbers of the $N_{\hat{h}_1}$ brane with a, d branes we have also present the sectors ah, $ah^\star$, dh, $dh^\star$. For our convenience we choose the number of *extra* D6-branes to be

$$\beta_1 = \beta_2 = \frac{1}{2}\,, \qquad N_h = 2\,. \tag{4.5}$$

Also, we define[8]

$$\begin{aligned}
&B_2^3 \wedge [2\epsilon\tilde{\epsilon}][(F^b + F^c)]\,,\\
&B_2^1 \wedge \Big(\epsilon\tilde{\epsilon}[4n_a^2\ F^a - 2n_d^2 F^d + 4m_b^1\ F^b + 4m_c^1\ F^c]\Big)\,,\\
&B_2^o \wedge (3\tilde{\epsilon})\,(F^a + F^d)\,.
\end{aligned} \tag{4.6}$$

The extra U(1)'s in the presence of the *extra* D6-branes h^1, h^2 may be defined as follows:

$$\mathrm{U}(1)^{(4)} = (Q^a - Q^d) + (Q^b - Q^c) + (F^{h_1} - F^{h_2})\,, \tag{4.7}$$
$$\mathrm{U}(1)^{(5)} = (F^{h_1} + F^{h_2})\,, \tag{4.8}$$
$$\mathrm{U}(1)^{(6)} = \frac{1}{2}\Big((Q^a - Q^d) + (Q^b - Q^c)\Big) - F^{h_1} + F^{h_2}\,. \tag{4.9}$$

The latter U(1)'s are defined in an orthogonal basis. We don't get additional constraints due to the presence of the *extra* D6-branes in the RR tadpole parameters. The only issue remaining is the breaking of the extra U(1)'s, (4.7), (4.8) and (4.9).

In the following we will describe the additional sectors of the theory due to the presence of the *extra* branes, h^1, h^2. The analysis concerns the intersections of the h^1 brane, but it is valid for the h^2 brane as well. One needs to mimic the procedure for the h^2 brane, obtaining one more copy of *extra* fermions and scalars. In the analysis below, one needs to set $\beta_1 = \beta_2 = 1/2$. Also the net effect of *extra* branes is the creation of enough singlets from the dh, $dh^\star$ sectors that may be used to break the extra U(1)'s (4.7), (4.8) and (4.9).

[8] The explanation of the origin of the structure of U(1) anomalies will be explained later in detail.

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- *ah-sector.* Because $I_{ah} = 3/\beta_1 > 0$ there are present $|I_{ah}| = |3/\beta_1|$ fermions λ_1^f, appearing in the representations

$$(4, 1, 1)_{(1,0,0,0;-1)} \,, \tag{4.10}$$

where the fifth-entry is the h-brane U(1) charge.

- *$ah^\star$-sector.* Since $I_{ah^\star} = -3/\beta_1 < 0$ there are present $|I_{ah^\star}| = |3/\beta_1|$ fermions $\tilde{\lambda}_2^f$, appearing in the representations

$$(\bar{4}, 1, 1)_{(-1,0,0,0;-1)} \,. \tag{4.11}$$

- *dh-sector.* Since $I_{dh} = -12/\beta_1 < 0$ there are present $|I_{dh}| = |12/\beta_1|$ fermions $\tilde{\lambda}_3^f$, appearing in the representations

$$(1, 1, 1)_{(0,0,0,-1;1)} \,. \tag{4.12}$$

We further require that this sector respects $N = 1$ supersymmetry. The condition for $N = 1$ supersymmetry in this sector is exactly

$$-\frac{\pi}{2} + \tilde{\vartheta}_2 + \ \vartheta_3 = 0 \tag{4.13}$$

which is satisfied when $\tilde{\vartheta}_2$, ϑ_3 take the value $\pi/4$ in consistency with (4.2). In this case we have also present the $\tilde{\lambda}_3^B$ massless scalar fields

$$(1, 1, 1)_{(0,0,0,-1;1)} \,. \tag{4.14}$$

The latter scalars receive a vev. The size of the vev, of order M_s, will be induced from the size of a coupling contributing to the mass of the χ_L fermions.

- *$dh^\star$-sector.* The intersection $I_{dh^\star} = -12/\beta_1 < 0$, thus there are present $|I_{ah^\star}| = |12/\beta_1|$ fermions $\tilde{\lambda}_4^f$, appearing in the representations

$$(1, 1, 1)_{(0,0,0,-1;-1)} \,. \tag{4.15}$$

We require that this sector respects $N = 1$ supersymmetry. The condition for $N = 1$ supersymmetry in this sector is exactly the same as in the dh sector, that is (4.13). In this case we have also present $I_{dh^\star}$ massless scalar fields $\tilde{\lambda}_4^B$, appearing as a linear combination of the representations

$$(1, 1, 1)_{(0,0,0,-1;-1)} \,. \tag{4.16}$$

The latter scalars receive a vev with size of order M_s. The latter will be induced from the size of a coupling contributing to the mass of the χ_L fermions.

We will now show that all fermions receive a mass and disappear from the low energy spectrum.

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- The mass term for the λ_1^f fermion reads:

$$\begin{aligned}&(\bar{4},1,1)_{(-1,0,0,0;1)}\ (\bar{4},1,1)_{(-1,0,0,0;1)}\ \langle(4,1,2)_{(1,0,1,0,;0)}\rangle\times\\&\times\langle(4,1,\bar{2})_{(1,0,-1,0,;0)}\rangle\langle(1,1,1)_{(0,0,0,1;-1)}\rangle\ \langle(1,1,1)_{(0,0,0,-1;-1)}\rangle\end{aligned}\tag{4.17}$$

or

$$\bar{\lambda}_1^f\bar{\lambda}_1^f\langle H_2\rangle\langle F_R^H\rangle\langle\lambda_B^3\rangle\langle\tilde{\lambda}_B^4\rangle\sim\bar{\lambda}_1^f\bar{\lambda}_1^f M_s\,.\tag{4.18}$$

- Similarly the mass term for the $\tilde{\lambda}_2^f$ fermion reads:

$$\begin{aligned}&(4,1,1)_{(1,0,0,0;1)}(4,1,1)_{(1,0,0,0;1)}\langle(\bar{4},1,2)_{(-1,0,1,0;0)}\rangle\times\\&\times\langle(\bar{4},1,\bar{2})_{(-1,0,-1,0;0)}\rangle\langle(1,1,1)_{(0,0,0,-1;-1)}\rangle\ \langle(1,1,1)_{(0,0,0,1;-1)}\rangle\end{aligned}\tag{4.19}$$

or

$$\bar{\tilde{\lambda}}_2^f\bar{\tilde{\lambda}}_2^f\langle\bar{F}_R^H\rangle\langle H_1\rangle\langle\tilde{\lambda}_B^4\rangle\langle\lambda_B^3\rangle\sim\bar{\tilde{\lambda}}_2^f\bar{\tilde{\lambda}}_2^f M_s\,.\tag{4.20}$$

- Similarly the mass term for the $\tilde{\lambda}_3^f$ fermion reads:

$$(1,1,1)_{(0,0,0,1;-1)}(1,1,1)_{(0,0,0,1;-1)}\langle(1,1,1)_{(0,0,0,-1;1)}\rangle\langle(1,1,1)_{(0,0,0,-1;1)}\rangle\tag{4.21}$$

or

$$\bar{\tilde{\lambda}}_3^f\bar{\tilde{\lambda}}_3^f\ \langle\tilde{\lambda}_3^B\rangle\langle\tilde{\lambda}_3^B\rangle\sim M_s\bar{\tilde{\lambda}}_3^f\ \bar{\tilde{\lambda}}_3^f\,.\tag{4.22}$$

- Similarly the mass term for the $\tilde{\lambda}_4^f$ fermion reads:

$$(1,1,1)_{(0,0,0,1;1)}(1,1,1)_{(0,0,0,1;1)}\langle(1,1,1)_{(0,0,0,-1;-1)}\rangle\langle(1,1,1)_{(0,0,0,-1;-1)}\rangle\tag{4.23}$$

or

$$\bar{\tilde{\lambda}}_4^f\bar{\tilde{\lambda}}_4^f\langle\tilde{\lambda}_4^B\rangle\langle\tilde{\lambda}_4^B\rangle\sim M_s\bar{\tilde{\lambda}}_4^f\bar{\tilde{\lambda}}_4^f\,.\tag{4.24}$$

Thus all fermions, bosons receive vevs, that appear due to the non-zero intersection numbers of the *extra* U(1) brane with a, d branes receive a string scale mass and disappear from the low energy spectrum. The surviving massless the Green-Schwarz mechanism U(1)'s (4.7)–(4.9) may be broken by vevs of $\tilde{\lambda}_3^B$, $\tilde{\lambda}_4^B$.

A comment is in order. In [29], the mass of the left handed weak fermion doublets χ_L was shown to receive corrections[9] from the coupling

$$(1,2,1)(1,2,1)e^{-A}\frac{\langle h_2\rangle\langle h_2\rangle\langle\bar{F}_R^H\rangle\langle H_1\rangle\langle\bar{s}_L^H\rangle}{M_s^4}\overset{A\to 0}{\sim}\frac{v^2}{M_s}(1,2,1)(1,2,1)\tag{4.25}$$

or

$$\begin{aligned}&(1,2,1)_{(0,1,0,-1)}\ (1,2,1)_{(0,1,0,-1)}\langle(1,\bar{2},\bar{2})_{(0,-1,-1,0)}\rangle\langle(1,\bar{2},\bar{2})_{(0,-1,-1,0)}\rangle\times\\&\times\langle(\bar{4},1,2)_{(-1,0,1,0)}\rangle\langle(4,1,2)_{(1,0,1,0)}\rangle\langle\mathbf{1}_{(0,0,0,2)}\rangle\,.\end{aligned}\tag{4.26}$$

[9] Where we have included the leading contribution of the worksheet area connecting the seven vertices.

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Thus the mass of χ_L was of the order

$$m_{\mathcal{H}_L} \sim \frac{v^2}{M_s} \tag{4.27}$$

which "localized" χ_L in the area between $100 - 246\,\mathrm{GeV}$. Thus, the necessity to push the mass of χ_L over $90\,\mathrm{GeV}$, pushed the string scale to be below $650\,\mathrm{GeV}$.

It turns out that the couplings (4.25) and (4.26) represent only one part of the couplings that contribute the lowest order correction to the χ_L mass. The use of the *extra* singlets guarantees the existence of another mass coupling of the order of the string scale. It involves the *extra* scalars singlets $\tilde{\lambda}_3^B$, λ_4^B and reads:

$$(1,2,1)(1,2,1)e^{-A}\,\frac{\langle h_2\rangle\langle h_2\rangle\langle \bar{F}_R^H\rangle\langle H_1\rangle\langle \bar{\tilde{\lambda}}_3^B\rangle\langle \lambda_4^B\rangle}{M_s^5} \sim \frac{v^2}{M_s}\ e^{-A}\ (1,2,1)(1,2,1) \tag{4.28}$$

explicitly, in representation form, given by

$$\begin{aligned}&(1,2,1)_{(0,1,0,-1)}\ (1,2,1)_{(0,1,0,-1)}\langle(1,\bar{2},\bar{2})_{(0,-1,-1,0)}\rangle\ \langle(1,\bar{2},\bar{2})_{(0,-1,-1,0)}\rangle \times \\ &\times \langle(\bar{4},1,2)_{(-1,0,1,0)}\rangle\ \langle(4,1,2)_{(1,0,1,0)}\rangle\ \langle(1,1,1)_{(0,0,0,1;-1)}\rangle\langle(1,1,1)_{(0,0,0,1;1)}\rangle\,,\end{aligned} \tag{4.29}$$

where we have included the leading contribution of the worksheet area connecting the eight vertices. We should emphasize that the vev of the scalars $\tilde{\lambda}_3^B$, $\tilde{\lambda}_4^B$, should be of order M_s in order for χ_L to be at least of order v^2/M_s. Any other value for $\tilde{\lambda}_3^B$, $\tilde{\lambda}_4^B$ sends the mass of χ_L at smaller than M_z values. That is, assuming that the exponential area factor in front of the coupling is of order one, beyond the SM fermions, there are light fermions surviving between M_Z and the scale of electroweak symmetry breaking $246\,\mathrm{GeV}$. The latter fermions appear to be a general prediction of GUT left-right constructions in the present D6 intersecting brane backgrounds. Exactly the same features will be found later for the PS-I class of GUT models.

5. Tadpole cancellation for the PS-I models

To understand the solution of the RR tadpole cancellation condition, which will be given in parametric form we should make the following comments:

a) The need to realize certain couplings will force us to demand that some intersections will preserve some supersymmetry. Thus some massive fields will be "pulled out" from the massive spectrum and become massless. For example, in order to realize a Majorana mass term for the right handed neutrinos we will demand that the *ac* sector preserves $N = 1$ SUSY. That will have as an immediate effect to "pull out" from the massive mode spectrum the $\bar{F}_R^H$ particles.

b) The intersection numbers, in table 1, of the fermions $F_L + \bar{F}_R$ are chosen such that $I_{ac} = -3$, $I_{ab^\star} = 3$. 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

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to us by the RR tadpole cancellation conditions the latter being equivalent to $\mathrm{SU}(N_a)$ gauge anomaly cancellation, in this case of $\mathrm{SU}(2)_L$, $\mathrm{SU}(2)_R$ gauge anomalies,

$$\sum_i I_{ia} N_a = 0, \quad a = L, R. \tag{5.1}$$

The theory breaks just to the standard model $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)_Y$ at low energies. The complete spectrum of the model appears in table 1. The tadpole solutions of PS-I models are presented in table 4.

c) The mixed anomalies A_{ij} of the four surplus U(1)'s with the non-abelian gauge groups $\mathrm{SU}(N_a)$ of the theory cancel through a generalized GS mechanism [35, 37], involving close string modes couplings to worldsheet gauge fields. Two combinations of the U(1)'s are anomalous and become massive through their couplings to RR fields, their orthogonal non-anomalous combinations survives, combining to a single U(1) that remains massless. Crucial for achieving the RR tadpole cancellation is the presence of *extra* D6-branes. Contrary, of what is happening in models, with exactly the SM at low energy, and a Standard-like structure at the string scale [26]–[28] and [32] where the *extra* branes have no intersection with other branes, in the GUT models there is a non-vanishing intersection. As an immediate consequence, this becomes a singlet generation mechanism after imposing $N = 1$ SUSY between U(1) leptonic and the U(1) *extra* D6-branes. Also, contrary to the SM's of [26]–[28] and [32], in the GUT constructions on the same open string backgrounds the *extra* branes do not form a $\mathrm{U}(N_h)$ gauge group but rather a $\mathrm{U}(1)^{N_1} \times \mathrm{U}(1)^{N_2} \times \cdots \times \mathrm{U}(1)^{N_h}$ one.

d) The constraint

$$\Pi_{i=1}^{3} m^i = 0\,. \tag{5.2}$$

is not imposed and thus leads to the appearance of the non-trivial chiral fermion content from the aa^*, sector with corresponding fermions ω_L, z_R.

e) After breaking the PS left-right symmetry at M_{GUT}, the surviving gauge symmetry is that of the SM augmented by some anomaly free U(1)'s,[10] their number depending on the number of *extra* U(1)'s that have been added to satisfy the RR tadpole conditions. In general the number of the surviving massless the Green-Schwarz mechanism U(1)'s is $1 + n_h$, where n_h is the number of *extra* U(1) branes.

   To break the latter U(1) symmetries we will impose that the dh, $dh^\star$ sector respects $N = 1$ SUSY.

f) Demanding $I_{ab} = 3$, $I_{ac} = -3$, it implies that the third tori should be tilted. By looking at the intersection numbers of table 1, we conclude that the b-brane should be paralled to the c-brane and the a-brane should be paralled to the d-brane as there is an absence of intersection numbers for those branes. Also, the cancellation of the

[10] Surviving the Green-Schwarz mechanism.

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N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 4$	$(0, \epsilon)$	$(n_a^2, 3\epsilon\beta_2)$	$(\tilde{\epsilon}, \tilde{\epsilon}/2)$
$N_b = 2$	$(-1, \epsilon m_b^1)$	$(1/\beta_2, 0)$	$(\tilde{\epsilon}, \tilde{\epsilon}/2)$
$N_c = 2$	$(1, \epsilon m_c^1)$	$(1/\beta_2, 0)$	$(\tilde{\epsilon}, -\tilde{\epsilon}/2)$
$N_d = 1$	$(0, \epsilon)$	$(n_d^2, -6\epsilon\tilde{\epsilon}\beta_2)$	$(2, 0)$
$\vdots$	$\vdots$	$\vdots$	$\vdots$
N_h	$(1/\beta_1, 0)$	$(1/\beta_2, 0)$	$(\tilde{\epsilon}, \tilde{\epsilon}/2)$

Table 4: Tadpole solutions for PS-I type models with D6-branes wrapping numbers 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 two integer parameters, n_a^2, n_d^2, the NS-background β_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_b^1, m_c^1. Note the presence of the N_h *extra* U(1) branes.

RR crosscap tadpole constraints is solved from parametric sets of solutions. They are given in table 4. We note that we have chosen our *extra* D6-branes to be located in

$$\left(\frac{1}{\beta_1}, 0\right) \left(\frac{1}{\beta_2}, 0\right) \left(1, \frac{1}{2}\right). \tag{5.3}$$

With the above choice, all tadpole conditions but the first and the third,[11], are satisfied, the first giving

$$N_h \frac{\tilde{\epsilon}}{\beta_1 \beta_2} = 16\,. \tag{5.4}$$

Thus the number of *extra* branes depends only on the NS background in the three tori and the sign of $\tilde{\epsilon}$. We note that when $\tilde{\epsilon} = 1$ we add D6 branes, while when $\tilde{\epsilon} = -1$ we add D6 anti-branes. For[12] $\beta_1 = \beta_2 = 1$, $N_h = 16$. If $\beta_1 = 1$, $\beta_2 = 1/2$ or $\beta_1 = 1/2$, $\beta_2 = 1$, $N_h = 8$. For

$$\beta_1 = \beta_2 = \frac{1}{2}\,, \qquad N_h = 4\,. \tag{5.5}$$

Also the third tadpole condition gives

$$2n_a^2 + \frac{1}{\beta_2}(m_b^1 - m_c^1) = 0\,. \tag{5.6}$$

To see clearly the cancellation of tadpoles, we have to choose a consistent numerical set of wrapping numbers, e.g.

$$\epsilon = \tilde{\epsilon} = 1\,, \quad n_a^2 = -1\,, \quad m_b^1 = 2\,, \quad m_c^1 = 1\,, \quad n_d^2 = 2\,, \quad \beta_1 = 1\,, \quad \beta_2 = \frac{1}{2}. \tag{5.7}$$

The latter can be satisfied with the addition of eight D6-branes with wrapping numbers $(1,0)(2,0)(1,1/2)$, effectively giving to the models the structure of table 5.

[11]We have added an arbitrary number of N_D U(1) branes which do not contribute to the rest of the tadpoles and intersection numbers. This is always an allowed choice. We chosen not to exhibit the rest of the tadpoles as they involve the identity $0 = 0$.

[12]In the following examples we choose $\tilde{\epsilon} = 1$.

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N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 4$	$(0, 1)$	$(-1, 3/2)$	$(1, 1/2)$
$N_b = 2$	$(-1, 2)$	$(2, 0)$	$(1, 1/2)$
$N_c = 2$	$(1, 1)$	$(2, 0)$	$(1, -1/2)$
$N_d = 1$	$(0, 1)$	$(2, -3)$	$(2, 0)$
$N_{h^1} = 1$	$(1, 0)$	$(2, 0)$	$(1, 1/2)$
$N_{h^2} = 1$	$(1, 0)$	$(2, 0)$	$(1, 1/2)$
$\vdots$	$\vdots$	$\vdots$	$\vdots$
$N_{h^8} = 1$	$(1, 0)$	$(2, 0)$	$(1, 1/2)$

Table 5: Wrapping number set consistent with the tadpole constraints (5.4) and (5.6).

We note that in the model described by table 5 the non-zero intersection numbers of the *extra* branes with a, d branes, will give us just the SM at low energies, in addition to a number of U(1)'s. The breaking of the latter U(1)'s will be facilited by the use of scalar singlets from the dh^i, $i = 1, \ldots, 8$ sectors as we will explain later.

g) The hypercharge operator is defined as a linear combination of the three diagonal generators of the SU(4), SU(2)$_L$, SU(2)$_R$ groups:

$$Y = \frac{1}{2}T_{3R} + \frac{1}{2}T_{B-L}\,, \quad T_{3R} = \text{diag}(1, -1)\,, \quad T_{B-L} = \text{diag}\left(\frac{1}{3}, \frac{1}{3}, \frac{1}{3}, -1\right). \tag{5.8}$$

Also,

$$Q = Y + \frac{1}{2}T_{3L}\,. \tag{5.9}$$

6. Cancellation of U(1) anomalies

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

$$A_{ij} = \frac{1}{2}(I_{ij} - I_{ij^\star})N_i\,. \tag{6.1}$$

Note that gravitational anomalies cancel since D6-branes never intersect O6-planes. In the orientifolded type I torus models gauge anomaly cancellation [37] proceeds through a generalized GS mechanism [26] that makes use of the 10-dimensional RR gauge fields C_2 and C_6 and gives at four dimensions the couplings to gauge fields

$$\begin{aligned} &N_a m_a^1 m_a^2 m_a^3 \int_{M_4} B_2^o \wedge F_a \;;\; n_b^1 n_b^2 n_b^3 \int_{M_4} C^o \wedge F_b \wedge F_b\,, \\ &N_a n^J n^K m^I \int_{M_4} B_2^I \wedge F_a \;;\; n_b^I m_b^J m_b^K \int_{M_4} C^I \wedge F_b \wedge F_b\,, \end{aligned} \tag{6.2}$$

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:

$$C^o \equiv \int_{(T^2)^1 \times (T^2)^2 \times (T^2)^3} C_6\,; \qquad C^I \equiv \int_{(T^2)^I} C_2\,, \tag{6.3}$$

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

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The triangle anomalies (6.1) cancel from the existence of the string amplitude involved in the GS mechanism [35] in four dimensions [37]. 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 [26] to

$$-N_a m_a^1 m_a^2 m_a^3 n_b^1 n_b^2 n_b^3 - N_a \sum_I n_a^I n_a^J n_b^K m_a^I m_b^J m_b^K \,, \qquad I \neq J, K \,. \tag{6.4}$$

The RR couplings B_2^I of (6.2), appear into three terms. In the general case we should have consider the contribution of N_h U(1) branes. The extra U(1)'s could be shown that are broken by appropriate vevs. The latter will be exhibited for the case of four *extra* U(1) branes. This is a minimal choice of *extra* branes given the choices (5.3), (5.5).

$$\begin{aligned}
&B_2^3 \wedge [2\tilde{\epsilon}][-(F^b + F^c) + \tilde{\epsilon} F^{\hat{h}_1} + \tilde{\epsilon} F^{\hat{h}_2} + \tilde{\epsilon} F^{\hat{h}_3} + \tilde{\epsilon} F^{\hat{h}_4}] \,, \\
&B_2^1 \wedge \left(\epsilon \tilde{\epsilon} [4 n_a^2 \ F^a + 2 n_d^2 \tilde{\epsilon} F^d + 4 m_b^1 \ F^b + 4 m_c^1 \ F^c] \right) \,, \\
&B_2^o \wedge (3\tilde{\epsilon})\, F^a \,.
\end{aligned} \tag{6.5}$$

Moreover, analyzing the mixed anomalies of the extra U(1)'s with the non-abelian gauge groups $SU(4)_c$, $SU(2)_R$, $SU(2)_L$ we can see that there are two anomaly free combinations $Q_b - Q_c$, Q_d. As can be seen from (6.5) two anomalous combinations of U(1)'s, e.g. F^a, $-(F^b + F^c) + \tilde{\epsilon}(F^{\hat{h}_1} + F^{\hat{h}_2} + F^{\hat{h}_3} + F^{\hat{h}_4})$ become massive through their couplings to RR fields B_2^o, B_2^3. Also another non-anomalous combination, which is model dependent, becomes massive through its couplings to the RR field B_2^1. Also there are five non-anomalous U(1)'s which are getting broken by vevs of some scalars, as we will discuss later. A comment is on order. We reming the reader that the presence of four U(1) *extra* D6-branes on top of the four stack GUT model, makes PS-I class of models to behave effectively as eight stack models.

They are:

$$\begin{aligned}
U(1)^{(4)} &= Q^b - Q^c + Q^d + \tilde{\epsilon}(F^{\hat{h}_1} + F^{\hat{h}_2} - F^{\hat{h}_3} - F^{\hat{h}_4}) \,, \\
U(1)^{(5)} &= (Q^b + Q^c) + \frac{\tilde{\epsilon}}{2}(F^{\hat{h}_1} + F^{\hat{h}_2} + F^{\hat{h}_3} + F^{\hat{h}_4}) \,, \\
U(1)^{(6)} &= \frac{1}{3}(Q^b - Q^c + Q_d) + \ \frac{\tilde{\epsilon}}{4}(-F^{\hat{h}_1} - F^{\hat{h}_2} + F^{\hat{h}_3} + F^{\hat{h}_4}) \,, \\
U(1)^{(7)} &= \tilde{\epsilon}(F^{\hat{h}_1} - F^{\hat{h}_2} + F^{\hat{h}_3} - F^{\hat{h}_4}) \,, \\
U(1)^{(8)} &= \tilde{\epsilon}(F^{\hat{h}_1} - F^{\hat{h}_2} - F^{\hat{h}_3} + F^{\hat{h}_4}) \,.
\end{aligned} \tag{6.6}$$

The eight U(1)'s are defined in an orthogonal basis. The orthogonality relations between the latter U(1)'s and the model dependent U(1) field coupled to B_2^1 give us the model dependent constraints

$$2 n_a^2 = \tilde{\epsilon} n_d^2 \,, \tag{6.7}$$

$$m_b^1 = -m_c^1 \,. \tag{6.8}$$

At this point we should list the couplings of the dual scalars C^I of B_2^I required to cancel the mixed anomalies of the U(1)'s with the non-abelian gauge groups $SU(N_a)$. We have included the contribution of the four *extra* U(1) branes $\hat{h}_1, \ldots, \hat{h}_4$.

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They are given by:

$$C^o \wedge 2\tilde{\epsilon}[-(F^b \wedge F^b) + (F^c \wedge F^c) + 2\tilde{\epsilon}(F^{h_1} \wedge F^{\hat{h}_1} + F^{\hat{h}_2} \wedge F^{\hat{h}_2} + F^{\hat{h}_3} \wedge F^{\hat{h}_3} + F^{\hat{h}_4} \wedge F^{\hat{h}_4})]\,,$$
$$C^3 \wedge \left(\frac{3\tilde{\epsilon}}{2}\right)[(F^a \wedge F^a) - 4(F^d \wedge F^d)]\,,$$
$$C^2 \wedge [\epsilon\tilde{\epsilon}]\left[\frac{n_a^2}{2}(F^a \wedge F^a) + m_b^1(F^b \wedge F^b) - m_c^1(F^c \wedge F^c)\right]. \tag{6.9}$$

As it will be shown later the condition (6.7) will be derived again, when imposing $N = 1$ supersymmetry on some open string sectors.

7. Higgs sector, $N = 1$ SUSY on intersections and extra U(1)'s

7.1 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 actual presence of massless fermions at each intersection, we have evident the presence of an equal number of massive bosons, in the NS-sector, in the same representations as the massless fermions [12]. Their mass is of order of the string scale and it should be taken into account when examining phenomenological applications related to the renormalization group equations. However, it is possible that some of those massive bosons may become tachyonic,[13] 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. Denoting the twist vector by $(\vartheta_1, \vartheta_2, \vartheta_3, 0)$, in the NS open string sector the lowest lying states are given by[14]

State	**Mass**
$(-1+\vartheta_1, \vartheta_2, \vartheta_3, 0)$	$\alpha' M^2 = \frac{1}{2}(-\vartheta_1 + \vartheta_2 + \vartheta_3)$
$(\vartheta_1, -1+\vartheta_2, \vartheta_3, 0)$	$\alpha' M^2 = \frac{1}{2}(\vartheta_1 - \vartheta_2 + \vartheta_3)$
$(\vartheta_1, \vartheta_2, -1+\vartheta_3, 0)$	$\alpha' M^2 = \frac{1}{2}(\vartheta_1 + \vartheta_2 - \vartheta_3)$
$(-1+\vartheta_1, -1+\vartheta_2, -1+\vartheta_3, 0)$	$\alpha' M^2 = 1 - \frac{1}{2}(\vartheta_1 + \vartheta_2 + \vartheta_3)\,.$ (7.1)

Exactly at the point, where one of these masses may become massless we have preservation of $\mathcal{N} = 1$ locally. The angles at the six different intersections can be expressed in terms of the parameters of the tadpole solutions.

[13]For consequences when these set of fields may become massless see [16].

[14]We assume $0 \leq \vartheta_i \leq 1$.

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Angle structure and Higgs fields for PS-I classes of models

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

$$
\begin{aligned}
\theta_1 &= \frac{1}{\pi}\cot^{-1}\frac{R_1^{(1)}}{\epsilon m_b^1 R_2^{(1)}}\,; \quad \theta_2 = \frac{1}{\pi}\cot^{-1}\frac{n_a^2 R_1^{(2)}}{3\epsilon\beta_2 R_2^{(2)}}\,; \quad \theta_3 = \frac{1}{\pi}\cot^{-1}\frac{2R_1^{(3)}}{R_2^{(3)}}\,;\\
\tilde{\theta_2} &= \frac{1}{\pi}\cot^{-1}\frac{n_d^2 R_1^{(1)}}{6\epsilon\tilde{\epsilon}\beta_2 R_2^{(1)}}\,; \quad \tilde{\theta_1} = \frac{1}{\pi}\cot^{-1}\frac{R_1^{(1)}}{m_c^1 R_2^{(1)}}\,,
\end{aligned} \tag{7.2}
$$

where $R_i^{(j)}$, $i = 1,2$ 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 six non-trivial intersections we have the presence of four states $t_i, i = 1,\ldots,4$, associated to the states (7.1). Hence we have a total of twenty four 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 (7.1) and PS-I models may not be tachyonic, we obtain a total of eighteen conditions for the PS-I type models with a D6-brane at angles configuration to be stable. They are given in the appendix. We don't consider the scalars from the dh, $dh^\star$ intersections. For these sectors we will require later that they preserve $N = 1$ SUSY. As a result all scalars, but one, in these sectors may become massive.

Lets us now turn our discussion to the Higgs sector of PS-I models. In general there are two different Higgs fields that may be used to break the PS symmetry. We remind the reader that they were given in (3.3). The question is if H_1, H_2 are present in the spectrum of PS-I 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. By looking at the $I_{ac^\star}$ intersection we can conclude that the scalar doublets $H^\pm$ get localized. They come from open strings stretching between the U(4) a-brane and $U(2)_R$ $c^\star$-brane.

The $H^\pm$'s come from the NS sector and correspond to the states[15]

$$
\begin{array}{cc}
\textbf{State} & \textbf{Mass}^2 \\
(-1+\vartheta_1, \vartheta_2, 0, 0) & \alpha'(\mathrm{Mass})^2_{H^+} = \dfrac{Z_3}{4\pi^2} + \dfrac{1}{2}(\vartheta_2 - \vartheta_1) \\
(\vartheta_1, -1+\vartheta_2, 0, 0) & \alpha'(\mathrm{Mass})^2_{H^-} = \dfrac{Z_3}{4\pi^2} + \dfrac{1}{2}(\vartheta_1 - \vartheta_2)\,,
\end{array} \tag{7.3}
$$

where Z_3 is the distance^2 in transverse space along the third torus, ϑ_1, ϑ_2 are the (relative)angles between the a-, $c^\star$-branes in the first and second complex planes respectively.

[15] A similar set of states was used in [29] to provide the PS-A model with left-right breaking symmetry scalars $H^\pm$.

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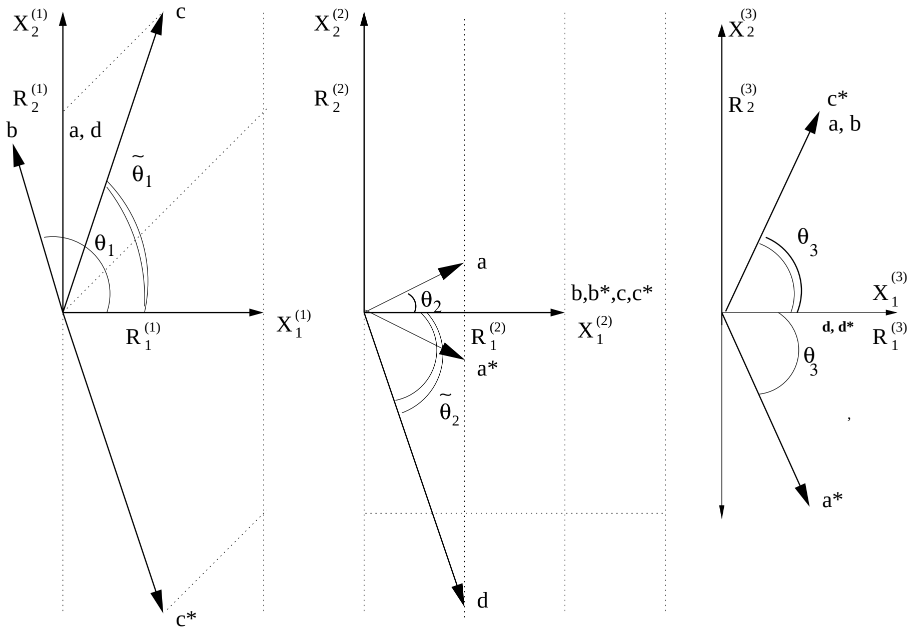

Figure 1: Assignment of angles between D6-branes on a a type I PS-I class of models based on the initial gauge group $\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R$. The angles between branes are shown on a product of $T^2 \times T^2 \times T^2$. We have chosen $\beta_1 = 1$, $m_b^1, m_c^1, n_a^2 > 0$, $\epsilon = \tilde{\epsilon} = 1$. These models break to low energies to exactly the SM.

The presence of scalar doublets $H^{\pm}$ can be seen as coming from the field theory mass matrix

$$(H_1^* \quad H_2) \, (\mathbf{M^2}) \begin{pmatrix} H_1 \\ H_2^* \end{pmatrix} + \text{h.c.} \,, \tag{7.4}$$

where

$$\mathbf{M^2} = M_s^2 \begin{pmatrix} Z_3^{(ac^*)}(4\pi^2)^{-1} & \frac{1}{2}|\vartheta_1^{(ac^*)} - \vartheta_2^{(ac^*)}| \\ \frac{1}{2}|\vartheta_1^{(ac^*)} - \vartheta_2^{(ac^*)}| & Z_3^{(ac^*)}(4\pi^2)^{-1} \end{pmatrix}, \tag{7.5}$$

The fields H_1 and H_2 are thus defined as

$$H^{\pm} = \frac{1}{2}(H_1^* \pm H_2) \,, \tag{7.6}$$

where their charges are given in table 6. Hence the effective potential which corresponds to the spectrum of the PS symmetry breaking Higgs scalars is given by

$$V_{\text{Higgs}} = m_H^2(|H_1|^2 + |H_2|^2) + (m_B^2 H_1 H_2 + \text{h.c}) \,, \tag{7.7}$$

where

$${m_H}^2 = \frac{Z_3^{(ac^*)}}{4\pi^2\alpha'} \,; \qquad m_B^2 = \frac{1}{2\alpha'}|\vartheta_1^{(ac^*)} - \vartheta_2^{(ac^*)}| \,. \tag{7.8}$$

The precise values of m_H^2, m_B^2, for the PS-I classes of models are:

$${m_H}^2 \overset{PS-I}{=} \frac{(\xi_a' + \xi_c')^2}{\alpha'} \,, \qquad m_B^2 \overset{PS-I}{=} \frac{1}{2\alpha'}\left|\frac{1}{2} + \tilde{\theta}_1 - \theta_2\right| , \tag{7.9}$$

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Intersection	PS breaking Higgs	Q_a	Q_b	Q_c	Q_d
$ac^\star$	H_1	1	0	1	0
$ac^\star$	H_2	-1	0	-1	0

Table 6: Higgs fields responsible for the breaking of $\mathrm{SU}(4) \times \mathrm{SU}(2)_R$ symmetry of the $\mathrm{SU}(4)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ type I model with D6-branes intersecting at angles. These Higgs are responsible for giving masses to the right handed neutrinos in a single family.

Intersection	Higgs	Q_a	Q_b	Q_c	Q_d
$bc^\star$	$h_1 = (1,2,2)$	0	1	1	0
$bc^\star$	$h_2 = (1,\bar{2},\bar{2})$	0	-1	-1	0

Table 7: Higgs fields present in the intersection $bc^\star$ of the $\mathrm{SU}(4)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ type I model 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.

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where $\xi'_a(\xi'_c)$ is the distance between the orientifold plane and the a, $c^\star$ branes and $\tilde{\theta}_1$, θ_2 were defined in (7.2). Thus

$$\begin{aligned}
m_B^2 &\overset{PS-I}{=} \frac{1}{2}|m^2_{\bar{F}_R}(t_2) + m^2_{\bar{F}_R}(t_3) - (m^2_{\bar{F}_R}(t_1) + m^2_{\bar{F}_R}(t_3))| \\
&= \frac{1}{2}|m^2_{\bar{F}_R}(t_2) + m^2_{\bar{F}_R}(t_3) - m^2_{F_L}(t_1) - m^2_{F_L}(t_3)| \\
&= \frac{1}{2}|m^2_{\chi^1_R}(t_2) + m^2_{\chi^1_R}(t_3) - (m^2_{\bar{F}_R}(t_1) + m^2_{\bar{F}_R}(t_3))| \\
&= \frac{1}{2}|m^2_{\chi^1_R}(t_2) + m^2_{\chi^1_R}(t_3) - m^2_{F_L}(t_1) - m^2_{F_L}(t_3)| \\
&= 1 - \frac{1}{2}|m^2_{\chi^2_R}(t_2) + m^2_{\chi^2_R}(t_3) - (m^2_{\bar{F}_R}(t_1) + m^2_{\bar{F}_R}(t_3))| \\
&= 1 - \frac{1}{2}|m^2_{\chi^2_R}(t_2) + m^2_{\chi^2_R}(t_3) - m^2_{F_L}(t_1) - m^2_{F_L}(t_3)|\,. \qquad (7.10)
\end{aligned}$$

For PS-I models the number of Higgs present is equal to the the intersection number product between the a-, $c^\star$- branes in the first and second complex planes,

$$n_{H^\pm} \overset{PS-I}{=} |I_{ac^\star}| = 3\,. \qquad (7.11)$$

A comment is in order. For PS-I 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.

More Higgs are present. In the $bc^\star$ intersection we have present some of the most useful Higgs fields of the models. They will be used to give mass to the quarks and leptons of the model as well breaking the electroweak symmetry. They appear in the representations $(1,2,2)$, $(1,\bar{2},\bar{2})$ and have been defined as h_1, h_2 in (3.5).

In the NS sector the lightest scalar states $h^{\pm}$ originate from open strings stretching between the $bc^{\star}$ branes

$$\begin{array}{cc} \textbf{State} & \textbf{Mass}^{\mathbf{2}} \\ (-1+\vartheta_1,0,0,0) & \alpha'(\text{Mass})^2 = \dfrac{\tilde{Z}_{23}^{bc^{\star}}}{4\pi^2} - \dfrac{1}{2}(\vartheta_1) \\ (\vartheta_1,-1,0,0) & \alpha'(\text{Mass})^2 = \dfrac{\tilde{Z}_{23}^{bc^{\star}}}{4\pi^2} + \dfrac{1}{2}(\vartheta_1)\,, \end{array} \tag{7.12}$$

where $\tilde{Z}_{23}^{bc^{\star}}$ is the relative distance in transverse space along the second and third torus from the orientifold plane, θ_1, is the (relative)angle between the b-, $c^{\star}$-branes in the first complex plane.

The presence of scalar doublets $h^{\pm}$ can be seen as coming from the field theory mass matrix

$$(\, h_1^* \quad h_2\,)\,(\,\mathbf{M^2}\,)\begin{pmatrix} h_1 \\ h_2^* \end{pmatrix} + \text{h.c.} \tag{7.13}$$

where

$$\mathbf{M^2} = M_s^2 \begin{pmatrix} Z_{23}^{(bc^*)}(4\pi^2)^{-1} & \dfrac{1}{2}|\vartheta_1^{(bc^*)}| \\ \dfrac{1}{2}|\vartheta_1^{(bc^*)}| & Z_{23}^{(bc^*)}(4\pi^2)^{-1} \end{pmatrix}. \tag{7.14}$$

The fields h_1 and h_2 are thus defined as

$$h^{\pm} = \frac{1}{2}(h_1^* \pm h_2)\,. \tag{7.15}$$

The effective potential which corresponds to the spectrum of electroweak Higgs h_1, h_2 may be written as

$$V_{\text{Higgs}}^{bc^{\star}} = \overline{m}_H^2(|h_1|^2 + |h_2|^2) + (\overline{m}_B^2 h_1 h_2 + \text{h.c})\,, \tag{7.16}$$

where

$$\overline{m}_H^2 \;=\; \frac{\tilde{Z}_{23}^{(bc^*)}}{4\pi^2\alpha'}\,; \qquad \overline{m}_B^2 \;=\; \frac{1}{2\alpha'}|\vartheta_1^{(bc^{\star})}|\,. \tag{7.17}$$

The precise values of for PS-I classes of models $\overline{m}_H^2$, $\overline{m}_B^2$ are

$$\bar{m}_H^2 \stackrel{PS-I}{=} \frac{(\tilde{\chi}_b^{(2)} + \tilde{\chi}_{c^{\star}}^{(2)})^2 + (\tilde{\xi}_b^{(3)} + \tilde{\xi}_{c^{\star}}^{(3)})^2}{\alpha'}\,; \qquad \bar{m}_B^2 \stackrel{PS-I}{=} \frac{1}{2\alpha'}|\tilde{\theta}_1 + \theta_1|\,; \tag{7.18}$$

where θ_1, $\tilde{\theta}_1$ are defined in (7.19), (7.20). Also $\tilde{\chi}_b, \tilde{\chi}_{c^{\star}}$ are the distances of the b, $c^{\star}$ branes from the orientifold plane in the second tori and $\tilde{\xi}_b$, $\tilde{\xi}_{c^{\star}}$ are the distances of the $b, c^{\star}$ branes from the orientifold plane in the third tori. Also, notice that the b, $c^{\star}$ branes are paralled along the second and third tori. The values of the angles θ_1, $\tilde{\theta}_1$ can be expressed in terms of the scalar masses in the various intersections. They are given by

$$\begin{aligned} \frac{1}{\pi}\theta_1 &= \frac{1}{2} + \frac{1}{2}(m_{\chi_R^1}^2(t_2) + m_{\chi_R^1}^2(t_3)) \\ &= \frac{1}{2}(m_{\chi_R^2}^2(t_2) + m_{\chi_R^2}^2(t_3)) - \frac{1}{2} \end{aligned}$$

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N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 4$	$(0, 1)$	$(5, 3/2)$	$(1, 1/2)$
$N_b = 2$	$(-1, -5/2)$	$(2, 0)$	$(1, 1/2)$
$N_c = 2$	$(1, 5/2)$	$(2, 0)$	$(1, -1/2)$
$N_d = 1$	$(0, 1)$	$(10, -3)$	$(2, 0)$
$N_{h^1} = 1$	$(2, 0)$	$(2, 0)$	$(1, 1/2)$
$\vdots$	$\vdots$	$\vdots$	$\vdots$
$N_{h^4} = 1$	$(2, 0)$	$(2, 0)$	$(1, 1/2)$

Table 8: Wrapping number set sonsistent with the constraints (5.6), (6.7), (6.8) and (7.22).

$$= \frac{3}{2} - \frac{1}{2}(m^2_{F_L}(t_2) + m^2_{F_L}(t_3)) \tag{7.19}$$

$$\begin{aligned} \frac{1}{\pi}\tilde{\theta}_1 &= \frac{1}{2}(m^2_{\chi^1_R}(t_2) + m^2_{\chi^1_R}(t_3)) - \frac{1}{2} \\ &= \frac{1}{2}(m^2_{\tilde{F}^2_R}(t_2) + m^2_{\tilde{F}^2_R}(t_3)) - \frac{1}{2} \\ &= \frac{1}{2} - \frac{1}{2}(m^2_{\chi^2_R}(t_2) + m^2_{\chi^2_R}(t_3))\,. \end{aligned} \tag{7.20}$$

The number of h_1, h_2 fields in the $bc^\star$ intersection is given by the intersection number of the b, $c^\star$ branes in the first tori

$$n^{bc^\star}_{h^\pm} \stackrel{PS-I}{=} |\epsilon(m^1_c - m^1_b)|\,. \tag{7.21}$$

A comment is in order. Because the number of the electroweak bidoublets in the PS-I models depends on the difference $|m^1_b - m^1_c|$, by using (5.6), (6.7) and (6.8), we get

$$m^1_b - m^1_c = 2m^1_b = -2\beta_2 n^2_a\,. \tag{7.22}$$

Hence, e.g. by choosing $n^2_a = 5$, $\beta_2 = 1/2$, $m^1_b = -5/2$, we get the constraint

$$n^{bc^\star}_{h^\pm} \stackrel{PS-I}{=} 5\,. \tag{7.23}$$

That is we have effectively choose five electroweak Higgs bidoublet present, each one appearing in each intersection as a linear combination of the h_1, h_2 fields. In this case, a consistent numerical set of wrappings will be, $\epsilon = \tilde{\epsilon} = 1$, $m^1_c = 5/2$, $n^2_d = 10$.

7.2 $N = 1$ SUSY on Intersections

In this section, we will demand that certain open string sectors respect $N = 1$ supersymmetry. The chiral spectrum of PS-I classes of models described in table 1 is massless at this point. The supersymmetry conditions, will create singlet scalars which receive vevs and generate masses for the otherwise massless fermions χ^1_L, χ^2_L, χ^1_R, χ^2_R.

In order that $N = 1$ SUSY is preserved at some intersection between two branes L, M we need to satisfy $\pm\vartheta^1_{ab} \pm \vartheta^2_{ab} \pm \vartheta^3_{ab}$ for some choice of signs, where $\vartheta^i_{\alpha\beta}$, $i = 1, 2, 3$ are the relative angles of the branes L, M across the three 2-tori.

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We note that a Majorana mass term for right neutrinos will appear once we impose $N = 1$ SUSY on an intersection. That will have as an effect the appearance of the massless scalar superpartners of the $\bar{F}_R$ fermions, the $\bar{F}_R^H$'s, allowing a dimension 5 Majorana mass term for ν_R, $F_R F_R \bar{F}_R^H \bar{F}_R^H$ (see section 8).

• *PS-I models with N=1 SUSY*

We demand that the sector ac respects $N = 1$ supersymmetry. The condition for $N = 1$ SUSY on the ac-sector reads:[16]

$$\pm(\frac{\pi}{2} + \tilde{\vartheta}_1) \ \pm \vartheta_2 \ \pm 2\vartheta_3 = 0\,. \tag{7.24}$$

This condition can be solved by choosing:

$$ac \to \left(\frac{\pi}{2} + \tilde{\vartheta}_1\right) + \vartheta_2 - 2\vartheta_3 = 0\,, \tag{7.25}$$

and thus may be solved by the choice[17]

$$-\tilde{\vartheta}_1 = \vartheta_2 = \vartheta_3 = \frac{\pi}{4}\,, \tag{7.26}$$

effectively giving us

$$\epsilon m_c^1\, \mathrm{U}^{(1)} = \frac{3\epsilon\beta_2}{n_a^2}\, \mathrm{U}^{(2)} = \frac{1}{2}\, \mathrm{U}^{(3)} = \frac{\pi}{4}\,. \tag{7.27}$$

By imposing $N = 1$ SUSY on an intersection a massless scalar partner appears in this sector. Thus in the ac-sector it is the massless scalar superpartner of the fermion $\bar{F}_R$, namely the $\bar{F}_R^H$, which is generated. An additional feature of, see (7.27), SUSY on intersections is that the complex structure moduli U^i takes specific values, thus reducing the degeneracy of moduli parameters in the theory.

As in the discussion of section 4, 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 the following discussion we consider only one of the two N_h U(1) branes, e.g. the $\mathrm{U}(1)_{N_{h_1}}$. The discussion may be repeated identically for the other U(1) branes present, e.g. $\mathrm{U}(1)_{N_{h_2}}$, $\mathrm{U}(1)_{N_{h_3}}$, $\mathrm{U}(1)_{N_{h_4}}$.

Due to the non-zero intersection numbers of the N_{h_1} brane with a, d branes the following sectors are present: ah, $ah^\star$, dh, $dh^\star$.

- *ah-sector.* Because $I_{ah} = 0$ there are no fermions[18] present from this sector.

- *$ah^\star$-sector.* Because $I_{ah^\star} = -3/\beta_1 < 0$, there are present $|I_{ah^\star}| = |3/\beta_1|$ fermions ϕ_2^f appearing in the representations

$$(\bar{4}, 1, 1)_{(-1,0,0,0;-1)}\,. \tag{7.28}$$

[16] We have chosen $m_c^1 < 0$.

[17] We have set $U^{(i)} = R_2^{(i)}/R_1^{(i)}$, $i = 1, 2, 3$.

[18] Obviously, there are no massless bosons from this sector.

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- *dh-sector.* Because $I_{dh} = -6/\beta_1 < 0$, there are present $|I_{dh}| = |6/\beta_1|$ fermions ϕ_3^f, appearing in the representations

$$(1,1,1)_{(0,0,0,-1;1)}\,. \tag{7.29}$$

We further require that this sector respects $N=1$ supersymmetry. The condition for $N=1$ supersymmetry in this sector is exactly

$$\frac{\pi}{2} - \tilde{\vartheta}_2 - \vartheta_3 = 0 \tag{7.30}$$

which is satisfied when $\tilde{\vartheta}_2$, ϑ_3 take the value $\pi/4$ in consistency with (4.2). The latter condition (and (4.2)) implies

$$2n_a^2 = \tilde{\epsilon} n_d^2\,, \tag{7.31}$$

which is exactly one of the conditions for the extra U(1)'s $\mathrm{U}(1)^{(4)}$, $\mathrm{U}(1)^{(5)}$, $\mathrm{U}(1)^{(6)}$, to survive massless the Green-Schwarz mechanism! That is, choosing an orthogonal basis for the anomaly free U(1)'s, the requirement on them to survive massless the Green-Schwarz mechanism is equivalent to the creation of singlet scalars from a $N=1$ supersymmetric intersection between a U(1) *extra* brane and the "leptonic" U(1) d-brane. Thus the presence of $N=1$ supersymmetric sectors involving *extra* branes in the non-SUSY PS-I classes of GUT models is equivalent to the presence of the generalized Green-Schwarz anomaly cancellation mechanism. A set of wrapping numbers consistent with this constraint can be seen in (8).

Also present are the $|I_{dh}|$ massless scalar fields ϕ_3^B, appearing in the representations

$$(1,1,1)_{(0,0,0,-1;1)}\,. \tag{7.32}$$

The latter scalars receive a vev which we assume to be of order of the string scale. The size of the vev will be induced once we examine the mass couplings of the χ_L^1 fermion (see also comments on concluding section).

- *$dh^\star$-sector.* Because $I_{dh^\star} = 6/\beta_1 > 0$, there are present $|I_{ah^\star}| = |6/\beta_1|$ fermions κ_4^f, appearing in the representations

$$(1,1,1)_{(0,0,0,1;1)}\,. \tag{7.33}$$

We require that this sector respects $N=1$ supersymmetry. The condition for $N=1$ supersymmetry in this sector is exactly the same as in the dh sector, that is (4.13). In this case we have also present $|I_{dh^\star}| = |6/\beta_1|$ massless scalar fields κ_4^B appearing in the representations

$$(1,1,1)_{(0,0,0,1;1)}\,. \tag{7.34}$$

The latter scalars receive a vev which we assume to be of order of the string scale. The size of the vev will be induced once we examine the mass couplings of χ_L^2 fermions.

We will now show that all fermions, appearing from the non-zero intersections of the *extra* brane $\mathrm{U}(1)_{N_{h_1}}$ with the branes a, d receive string scale mass.

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- The mass term for the ϕ_2^f fermion reads:

$$\begin{aligned}&(4,1,1)_{(1,0,0,0;1)}(4,1,1)_{(1,0,0,0;1)}\ \langle(\bar{4},1,2)_{(-1,0,1,0;0)}\rangle\times\\ &\times\langle(\bar{4},1,\bar{2})_{(-1,0,-1,0;0)}\rangle\langle(1,1,1)_{(0,0,0,-1;-1)}\rangle\langle(1,1,1)_{(0,0,0,1;-1)}\rangle\end{aligned} \tag{7.35}$$

or

$$\bar{\phi}_2^f\bar{\phi}_2^f\langle H_1\rangle\langle\bar{F}_R^H\rangle\langle\phi_B^4\rangle\langle\kappa_B^3\rangle\sim\bar{\phi}_2^f\bar{\phi}_2^f M_s\,. \tag{7.36}$$

- The mass term for the ϕ_3^f fermion reads:

$$(1,1,1)_{(0,0,0,1;-1)}(1,1,1)_{(0,0,0,1;-1)}\langle(1,1,1)_{(0,0,0,-1;1)}\rangle\langle(1,1,1)_{(0,0,0,-1;1)}\rangle \tag{7.37}$$

or

$$\bar{\phi}_3^f\bar{\phi}_3^f\langle\phi_3^B\rangle\langle\phi_3^B\rangle\sim M_s\bar{\phi}_3^f\bar{\phi}_3^f\,. \tag{7.38}$$

- The mass term for the κ_4^f fermion reads:

$$(1,1,1)_{(0,0,0,-1;-1)}(1,1,1)_{(0,0,0,-1;-1)}\langle(1,1,1)_{(0,0,0,1;1)}\rangle\langle(1,1,1)_{(0,0,0,1;1)}\rangle \tag{7.39}$$

or

$$\bar{\kappa}_4^f\bar{\kappa}_4^f\langle\kappa_4^B\rangle\langle\kappa_4^B\rangle\sim M_s\bar{\kappa}_4^f\bar{\kappa}_4^f\,. \tag{7.40}$$

7.3 Breaking the extra U(1)'s

In the standard version of a non-SUSY left-right Pati-Salam model if the neutral component of H_1 (resp. H_2), ν_H, acquires a vev, e.g. $\langle\nu^H\rangle$, then the initial gauge symmetry, $U(4)\times U(2)_L\times U(2)_R\times U(1)_d$, can break to the standard model gauge group $SU(3)\times U(2)\times U(1)_Y$ augmented by the extra, non-anomalous, U(1). In the PS-I models the initial gauge symmetry is $U(4)\times U(2)_L\times U(2)_R\times U(1)_d\times U(1)_{N_{h_1}}\times U(1)_{N_{h_2}}\times U(1)_{N_{h_3}}\times U(1)_{N_{h_4}}$. After the use of the Green-Schwarz mechanism and the use of the PS breaking Higgs scalars the gauge symmetry is $SU(3)\times U(2)\times U(1)_Y\times U(1)^{(4)}\times U(1)^{(5)}\times U(1)^{(6)}\times U(1)^{(7)}\times U(1)^{(8)}$. By appropriate Higgsing we may break the extra, beyond the SM, U(1)'s.

In the PS-A models, by imposing SUSY on sectors $dd^\star$, dh, $dh^\star$ we made it possible to generate the appearance of the scalar superpartners of s_L, $\tilde{s}_L$, ϕ_3^f, κ_4^f. In this case, different singlets $\tilde{s}_L$, are charged under the U(1)'s, $U(1)^{(4)}$, $U(1)^{(6)}$, and thus break them. Also, ϕ_3^B, κ_4^B, get charged under the anomaly free U(1) symmetries $U(1)^{(5)}$ and thus some of them may be used to break it. Thus finally, the PS-A models break to exactly the much wanted SM gauge group structure, $SU(3)\otimes SU(2)\otimes U(1)_Y$.

In the case of the non-anomalous U(1)'s (6.6) of PS-I class of models the available singlets that would break the extra U(1)'s come from the dh_i, $dh_i^\star$, $i=1,\ldots,4$ sectors. Note that in both PS-A and PS-I classes of Pati-Salam type GUTS the extra non-anomalous U(1)'s have some important phenomenological properties. In particular they do not charge the PS symmetry breaking Higgs scalars H_1, H_2 thus avoiding the appearance of axions.

We emphasize that up to this point the only issue remaining is how we can give non-zero masses to all fermions of table 1 beyond those of SM fermions.

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8. Neutrino couplings and fermion masses

Proton decay is one of the most important problems of grand unifies theories. In the standard versions of left-right symmetric PS models this problem is avoided as B-L is a gauged symmetry but the problem persists in baryon number violating operators of sixth order, contributing to proton decay. In the PS-I models proton decay is absent as baryon number survives as a global symmetry to low energies. That provides for an explanation for the origin of proton stability in general brane-world scenarios. Clearly $Q_a = 3B + L$ and the baryon B is given by

$$B = \frac{Q_a + Q_{B-L}}{4} . \tag{8.1}$$

In intersecting brane worlds the usual tree level SM fermion mass generating trilinear Yukawa couplings between the 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. Its general form for a six dimensional torus is in the leading order [12],

$$Y^{ijk} = e^{-\tilde{A}_{ijk}} , \qquad \tilde{A}_{ijk} \equiv A_1 , \tag{8.2}$$

where $\tilde{A}_{ijk}$ is the worldsheet area connecting the three vertices. The areas of each of the two dimensional torus involved in this interaction is typically of order one in string units. As in [29], we can without loss of generality assume that the areas of the second and third tori are close to zero. In this case, the area of the full Yukawa coupling (8.2) takes the form

$$Y^{ijk} = e^{-\frac{R_1 R_2}{a'} A_{ijk}} , \tag{8.3}$$

where R_1, R_2 the radii and A_{ijk} the area of the two dimensional tori in the first complex plane. For the dimension five interaction term, like those involved in the Majorana mass term for the right handed neutrinos the interaction term is scaled in the form

$$Y^{lmni} = e^{-\tilde{A}_{lmni}} , \qquad \tilde{A}_{lmni} \equiv A_2 , \tag{8.4}$$

where $\tilde{A}_{lmni}$ the worldsheet area connecting the four interaction vertices. Assuming that the areas of the second and third tori are close to zero, the four term coupling can be approximated as

$$Y^{ijk} = e^{-\frac{R_1 R_2}{a'} \tilde{A}_{lmni}} , \tag{8.5}$$

where the area of the $\tilde{A}_{lmni}$ may be of order one in string units.

Thus the full Yukawa interaction for the chiral spectrum of the PS-I models reads:

$$\lambda_1 F_L \ \bar{F}_R \ h + \lambda_2 \frac{F_R F_R \bar{F}_R^H \bar{F}_R^H}{M_s} , \tag{8.6}$$

where

$$\lambda_1 \equiv e^{-(R_1 R_2 A_1)/\alpha'} , \qquad \lambda_2 \equiv e^{-(R_1 R_2 A_2)/\alpha'} , \tag{8.7}$$

and the Majorana coupling involves the massless scalar superpartners $\bar{F}_R^H$ of the antiparticles $\bar{F}_R$. This coupling is unconventional, in the sense that the $\bar{F}_R^H$ is generated by imposing

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SUSY on a sector of a non-SUSY model. We note the presence of $N = 1$ SUSY at the sector ac. As can be seen by comparison with (3.6) the $\bar{F}_R^H$ has a neutral direction that receives the vev $\langle H \rangle$. There is no restriction on the vev of F_R^H from first principles and its vev can be anywhere between the scale of electroweak symmetry breaking and M_s.

The Yukawa term

$$F_L \bar{F}_R h \,, \qquad h = \{h_1, h_2\} \,, \tag{8.8}$$

is responsible for the electroweak symmetry breaking. This term generates Dirac masses to up quarks and neutrinos. Thus, we get

$$\lambda_1 F_L \bar{F}_R h \to (\lambda_1 \ \upsilon)(u_i u_j^c + \nu_i N_j^c) + (\lambda_1 \ \tilde{\upsilon}) \cdot (d_i d_j^c + e_i e_j^c) \,, \tag{8.9}$$

where we have assumed that

$$\langle h \rangle = \begin{pmatrix} \upsilon & 0 \\ 0 & \bar{\upsilon} \end{pmatrix} . \tag{8.10}$$

We observe that the model gives non-zero tree level masses to the fields present. These mass relations may be retained at tree level only, since as the model has a non-supersymmetric fermion spectrum, it will receive higher order corrections. It is interesting that from (8.10) we derive the GUT relation [36]

$$m_d = m_e \,, \tag{8.11}$$

as well the unwanted

$$m_u = m_{N^c \nu} \,. \tag{8.12}$$

In the case of neutrino masses, the "unwanted" (8.12), associated to the $\nu - N^c$ mixing, is modified due to the presence of the Majorana term in (8.6) leading to the see-saw mixing type neutrino mass matrix

$$\begin{pmatrix} \nu & N^c \end{pmatrix} \begin{pmatrix} 0 & m \\ m & M \end{pmatrix} \times \begin{pmatrix} \nu \\ N^c \end{pmatrix} , \tag{8.13}$$

where

$$m = \lambda_1 \upsilon \,. \tag{8.14}$$

After diagonalization the neutrino mass matrix gives us two eigenvalues, the heavy eigenvalue

$$m_{\text{heavy}} \approx M = \lambda_2 \frac{\langle H \rangle^2}{M_s} \,, \tag{8.15}$$

corresponding to the right handed neutrino and the light eigenvalue

$$m_{\text{light}} \approx \frac{m^2}{M} = \frac{\lambda_1^2}{\lambda_2} \times \frac{\upsilon^2 \ M_s}{\langle H \rangle^2} \tag{8.16}$$

corresponding to the left handed neutrino.[19]

Values of the parameters giving us values for neutrino masses between $0.1 - 10\,\text{eV}$, consistent with the observed neutrino mixing in neutrino oscillation measurements, will

[19]The neutrino mass matrix is of the type of an extended Frogatt-Nielsen mechanism [33] mixing light with heavy states.

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not be presented here, as they have already been discussed in [29]. The analysis remain the same, as the mass scales as well the Yukawa coupling parametrization of the theory do not change. We note that the hierarchy of neutrino masses has been investigated by examining several different scenaria associated with a light ν_L mass including the cases $\langle H\rangle = |M_s|$, $\langle H\rangle < |M_s|$. In both cases a hierarchy of neutrino masses in the area of $0.1 - 10\,\mathrm{V}$ in consistency with neutrino oscillation experiments can be easily accommodated for a wide choice of parameters.

Our main focus in this part is to show that all additional particles, appearing in table 1, beyond those of $F_L + \bar{F}_R$, get a heavy mass and disappear from the low energy spectrum. The only exception will be the light masses of χ_L^1, χ_L^2, weak fermion doublets which are of order of the electroweak symmetry breaking scale, e.g. 246 GeV. Lets us discuss the latter issue in more detail. The left handed fermions χ_L^1 receive a mass from the coupling

$$(1,2,1)(1,2,1)e^{-A}\frac{\langle h_2\rangle\langle h_2\rangle\langle \bar{F}_R^H\rangle\langle H_1\rangle\langle \phi_3^B\rangle\langle \kappa_4^B\rangle}{M_s^5} \overset{A\to 0}{\sim} \frac{\upsilon^2}{M_s}\,(1,2,1)(1,2,1) \tag{8.17}$$

explicitly, in representation form, given by

$$\begin{aligned}&(1,2,1)_{(0,1,0,-1;0)}\ (1,2,1)_{(0,1,0,-1;0)}\langle(1,\bar{2},\bar{2})_{(0,-1,-1,0;0)}\rangle\ \langle(1,\bar{2},\bar{2})_{(0,-1,-1,0;0)}\rangle\times\\ &\times\langle(\bar{4},1,2)_{(-1,0,1,0;0)}\rangle\langle(4,1,2)_{(1,0,1,0;0)}\rangle\langle(1,1,1)_{(0,0,0,1,-1)}\rangle\ \langle(1,1,1)_{(0,0,0,1,1)}\rangle\,,\end{aligned} \tag{8.18}$$

where we have included the leading contribution of the worksheet area connecting the seven vertices. Also we have assumed that $\langle\bar{\phi}_3^B\rangle = \langle\kappa_4^B\rangle = \langle H_1\rangle = M_s$. Any other value for these scalars will lower the mass of χ_L^1 below M_z something unacceptable.

In the following for simplicity reasons we will set the leading contribution of the different couplings to one (e.g. area tends to zero).

Also the left handed fermions χ_L^2 receive an M_s mass from the coupling

$$(1,2,1)(1,2,1)\frac{\langle h_2\rangle\langle h_2\rangle\langle \bar{F}_R^H\rangle\langle H_1\rangle\langle \phi_3^B\rangle\langle \bar{\kappa}_4^B\rangle}{M_s^5}\sim\frac{\upsilon^2}{M_s}\,(1,2,1)(1,2,1) \tag{8.19}$$

explicitly, in representation form, given by

$$\begin{aligned}&(1,2,1)_{(0,1,0,-1;0)}\ (1,2,1)_{(0,1,0,-1;0)}\langle(1,\bar{2},\bar{2})_{(0,-1,-1,0;0)}\rangle\ \langle(1,\bar{2},\bar{2})_{(0,-1,-1,0;0)}\rangle\times\\ &\times\langle(\bar{4},1,2)_{(-1,0,1,0;0)}\rangle\ \langle(4,1,2)_{(1,0,1,0;0)}\rangle\ \langle(1,1,1)_{(0,0,0,-1,1)}\rangle\ \langle(1,1,1)_{(0,0,0,-1,-1)}\rangle\,.\end{aligned} \tag{8.20}$$

Altogether, χ_L^1, χ_L^2, receive a mass of order υ^2/M_s and thus are expected to be found between M_Z and the scale of electroweak symmetry breaking.

The χ_R^1 doublet fermions receive heavy masses of order M_s in the following way:

$$(1,1,2)(1,1,2)\frac{\langle H_2\rangle\langle F_R^H\rangle\langle\bar{\phi}_3^B\rangle\langle\kappa_4^B\rangle}{M_s^3}\,. \tag{8.21}$$

In explicit representation form

$$\begin{aligned}&(1,1,2)_{(0,0,1,-1;0)}\ (1,1,2)_{(0,0,1,-1;0)}\ \langle(\bar{4},1,\bar{2})_{(-1,0,-1,0;0)}\rangle\ \langle(4,1,\bar{2})_{(1,0,-1,0;0)}\rangle\times\\ &\times\langle(1,1,1)_{(0,0,0,1,-1;0)}\rangle\ \langle(1,1,1)_{(0,0,0,1,1;0)}\rangle\end{aligned} \tag{8.22}$$

with vevs $\langle H_2\rangle\sim\langle F_R^H\rangle\sim M_s$, the mass of χ_R^1 is of order M_s.

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We note that in principle the vevs of ϕ_3^B, κ_4^B setting the scale of breaking of the extra anomaly free U(1) could be anywhere between $\langle v\rangle$ and M_s.

The χ_R^2 doublet fermions receive heavy masses of order M_s in the following way:

$$(1,1,2)(1,1,2)\frac{\langle H_2\rangle\langle F_R^H\rangle\langle\phi_3^B\rangle\langle\phi_3^B\rangle}{M_s^3}\,. \tag{8.23}$$

In explicit representation form

$$\begin{aligned}&(1,1,2)_{(0,0,1,1;0)}\ (1,1,2)_{(0,0,1,1;0)}\ \langle(\bar{4},1,\bar{2})_{(-1,0,-1,0;0)}\rangle\ \langle(4,1,\bar{2})_{(1,0,-1,0;0)}\rangle\times\\&\times\langle(1,1,1)_{(0,0,0,-1,1;0)}\rangle\ \langle(1,1,1)_{(0,0,0,-1,-1;0)}\rangle\end{aligned} \tag{8.24}$$

with vevs $\langle H_2\rangle\sim\langle F_R^H\rangle\sim M_s$, the mass of χ_R^2 is of order M_s.

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

$$(\bar{6},1,1)(\bar{6},1,1)\frac{\langle H_1\rangle\langle F_R^H\rangle\langle H_1\rangle\langle F_R^H\rangle}{M_s^3}\,, \tag{8.25}$$

where we have made use of the SU(4) tensor products $6\otimes 6=1+15+20$, $4\otimes 4=6+10$. Explicitly, in representation form,

$$\begin{aligned}&(\bar{6},1,1)_{(-2,0,0,0;0)}\ (\bar{6},1,1)_{(-2,0,0,0;0)}\langle(4,1,2)_{(1,0,1,0;0)}\rangle\ \langle(4,1,2)_{(1,0,1,0;0)}\rangle\times\\&\times\ \langle(4,1,\bar{2})_{(1,0,-1,0;0)})\rangle\ \langle(4,1,\bar{2})_{(1,0,-1,0;0)})\rangle\,.\end{aligned} \tag{8.26}$$

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

$$(10,1,1)(10,1,1)\frac{\langle\bar{F}_R^H\rangle\langle\bar{F}_R^H\rangle\langle H_2\rangle\langle H_2\rangle}{M_s^3}\,, \tag{8.27}$$

where we have used the tensor product representations for SU(4), $10\otimes 10=20+35+45$, $20\otimes\bar{4}=\bar{1}5+\bar{2}0$, $\bar{2}0\otimes\bar{4}=\bar{6}+10$, $10\otimes\bar{4}=4+36$, $4\otimes\bar{4}=1+15$. Explicitly, in representation form,

$$\begin{aligned}&(10,1,1)_{(2,0,0,0;0)}(10,1,1)_{(2,0,0,0;0)}\langle(\bar{4},1,2)_{(-1,0,1,0;0)}\rangle\ \langle(\bar{4},1,2)_{(-1,0,1,0;0)}\rangle\times\\&\times\langle(\bar{4},1,\bar{2})_{(-1,0,-1,0;0)}\rangle\ \langle(\bar{4},1,\bar{2})_{(-1,0,-1,0;0)}\rangle\,.\end{aligned} \tag{8.28}$$

Thus only the chiral fermion content of the SM fermions remains at low energy.

9. Conclusions

In this work, we have continue our discussion in [29] of constructing left-right symmetric G_{422} Pati-Salam GUT models in the context of D6 branes intersecting on compactifications of type IIA on an orientifolded factorizable T^6 tori. The GUT models based on the latter open string backgrounds have the unique future of breaking exactly to the SM at low energy. They are constructed as intersecting number deformations, around the basic intersection number structure in which the quarks and leptons of the G_{422} GUT structure $SU(4)_C\times SU(2)_L\times SU(2)_R$ are accommodated.

Most important, we presented a new mechanism of generating singlet scalars in the context of intersecting branes. It amounts in the use of *extra* U(1) branes needed in the satisfaction of the RR tadpole cancellation conditions,[20] the latter having non-trivial intersection numbers with the colour a-brane and the leptonic d-brane. The presence of *extra* branes creates singlets scalars that may be used to break the additional *extra* U(1)'s that survive massless the Green-Schwarz mechanism. Equally important, as we showed in subsection 7.2, the existence of $N = 1$ supersymmetry conditions in open string sectors involving the *extra* branes,[21] is equivalent to the existence of the orthogonality conditions for the U(1)'s surviving massless the presence of the generalized Green-Schwarz mechanism.

The special form of the solutions to the RR tadpole cancellation conditions allows exotic, antisymmetric and symmetric, fermionic representations of the colour degrees of freedom, arising from brane-orientifold image brane, $\alpha\alpha^{\star}$, sectors. Interestingly the models have the capacity to accommodate couplings that give a mass of order M_s to all these exotic fermions.

The models have some important phenomenological features, namely they can easily accommodate small values of neutrino masses of order $0.1 - 10\,\mathrm{eV}$ in consistency with neutrino oscillation experiments and a stable proton. The stability of the proton is guaranteed as baryon number is a gauged symmetry and survives as global symmetry to low energies. Moreover, colour triplet Higgs couplings that could couple to quarks and leptons and cause a problem to proton decay are absent in all classes of models.

Despite the fact, that the non-supersymmetric models we examined are free of RR tadpoles and, if the angle stabilization conditions of appendix hold, free of tachyons, they will always have closed string NSNS tadpoles that cannot all be removed. Some ways that this might be possible have been suggested in [23], by freezing the complex moduli to discrete values, or by background redefinition in terms of wrapped metrics [38]. However, it appears that a dilaton tadpole will always remain that could in principle reintroduce tadpoles in the next leading order. We note that in NS tadpoles are not existent in supersymmetric models but the backgrounds that we examined in this work, are non-supersymmetric.

Also, we note that the complex structure moduli[22] can be fixed to discrete values using the supersymmetry conditions, e.g. see (7.27), and in this way it is possible that some if not all, of the NS tadpoles can be removed. We leave this task for a future investigation.

One point that we want to emphasize is that until recently, in orientifolded six-torus compactifications there was any obvious explanation for keeping the string scale low [4], e.g. to the $1 - 100\,\mathrm{TeV}$ region. Thus controlling the hierarchy by making the Planck scale large, while keeping the string scale low, by varying the radii of the transverse directions [4]

[20]This is to be contrasted with models with just the SM at low energy from a Standard -like structure at the string scale [26]–[28] and [32], where the presence of *extra* branes has no intersection with the rest of the branes.

[21]Needed to satisfy the RR tadpole cancellation conditions.

[22]As was noted in [29] the Kähler moduli could be fixed from its value at the string scale, using relations involving the product radii (see [29]) but in this way we could use a large fine tuning which seems unnatural in a string theory context, where moduli should be assigned values dynamically.

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could not be applied, as there are no simultaneously transverse torus directions to all D6-branes [2]. However, as was noted in [29] and that is also the case for the classes of PS-I models examined in this work, there is an alternative mechanism that keeps the string scale M_s low. In particular the existence of the light weak doublets χ_L^1, χ_L^2 with mass of order up to 246 GeV, makes a definite prediction for a low string scale in the energy range less than 650 GeV. That effectively, makes the PS-I class of D6-brane models (also the PS-A class) directly testable to present or feature accelerators.

We should emphasize that crucial in showing that the GUT classes of models presented break exactly to the SM at low energies was our passive acceptence that there are couplings allowed by charge and gauge invariance selection rules that may give the beyond the SM fermions masses, if some scalars get a vev. However, it should be pointed out that whether of not in the present models these scalars get a vev is a highly non-trivial dynamical problem which in order to be solved precicely, we have to calculate at the string theory level the effective potential for these moduli scalars. However, with our present level of understranding of non-SUSY intersecting braneworld models this is a non-trivial question, as first of all we have to solve the stability problem of the configurations, as we have already commented about. We also note that in the context of intersecting branes it is not clear at all, that at the point in the moduli space that these scalars receive a vev, whether or not the system of recombined branes, signalling the presence of tachyon at the minimum of the scalar potential that would break the gauge symmetry, has a lower energy than the rest of the scalars and thus standard electroweak symmetry breaking will be preferred.[23] However, it is absolute amazing that at the present level of understanding the intersecting brane worlds, that we can find models that have all the necessary couplings in building classes of models with only the Standard model at the low energy limit.

Also, it will be interesting to extend the methods employed in this article, to GUT groups of the same type in higher stacks [40]. Summarizing, in the present work, we have shown that it is possible to consider further four stack classes of GUT models with exactly the SM at low energy, the geometry of which depends on deforming the basic PS quark-lepton intersection structure $I_{ab^\star} = 3$, $I_{ac} = -3$ and the presence of *extra* branes.

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A. Tachyon free conditions for classes of PS-I GUTS

In this appendix we list the conditions, mentioned in section 5, under which the PS-I model D6-brane configurations of tadpole solutions of table 4, are tachyon free. Note that the conditions are expressed in terms of the angles defined in (7.2).

$$-\left(\frac{3\pi}{2} - \vartheta_1\right) + \vartheta_2 + 2\vartheta_3 \geq 0$$

[23] See also the 2nd reference of [12] and [39] for some discussions relevant to this problem.

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$$
\begin{aligned}
-\left(\frac{\pi}{2}+\tilde{\vartheta}_1\right)+\vartheta_2+2\vartheta_3 &\geq 0\\
-\left(-\frac{\pi}{2}+\vartheta_1\right)+\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
-\left(\frac{\pi}{2}+\tilde{\vartheta}_1\right)+\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
-\left(\frac{\pi}{2}+\vartheta_1\right)+\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
-\left(\frac{\pi}{2}-\tilde{\vartheta}_1\right)+\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
\left(\frac{3\pi}{2}-\vartheta_1\right)-\vartheta_2+2\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}+\tilde{\vartheta}_1\right)-\vartheta_2+2\vartheta_3 &\geq 0\\
\left(-\frac{\pi}{2}+\vartheta_1\right)-\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}+\tilde{\vartheta}_1\right)-\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}+\vartheta_1\right)-\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}-\tilde{\vartheta}_1\right)-\tilde{\vartheta}_2+\vartheta_3 &\geq 0\\
\left(\frac{3\pi}{2}-\vartheta_1\right)+\vartheta_2-2\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}+\tilde{\vartheta}_1\right)+\vartheta_2-2\vartheta_3 &\geq 0\\
\left(-\frac{\pi}{2}+\vartheta_1\right)+\tilde{\vartheta}_2-\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}+\tilde{\vartheta}_1\right)+\tilde{\vartheta}_2-\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}+\vartheta_1\right)+\tilde{\vartheta}_2-\vartheta_3 &\geq 0\\
\left(\frac{\pi}{2}-\tilde{\vartheta}_1\right)+\tilde{\vartheta}_2-\vartheta_3 &\geq 0\,.
\end{aligned}
\tag{A.1}
$$

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