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Intersecting brane models of particle physics and the Higgs mechanism

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Abstract: We analyze a recently constructed class of D-brane theories with the fermion spectrum of the SM at the intersection of D6-branes wrapping a compact toroidal space. We show how the SM Higgs mechanism appears as a brane recombination effect in which the branes giving rise to $\mathrm{U}(2)_L \times \mathrm{U}(1)$ recombine into a single brane related to $\mathrm{U}(1)_{em}$. We also show how one can construct D6-brane models which respect some supersymmetry at every intersection. These are *quasi-supersymmetric* models of the type introduced in hep-th/0201205 which may be depicted in terms of *SUSY-quivers* and may stabilize the hierarchy between the weak scale and a fundamental scale of order 10-100 TeV present in low string scale models. Several explicit D6-brane models with three generation of quarks and leptons and different SUSY-quiver structure are presented. One can prove on general grounds that if one wants to build a (factorizable) D6-brane configuration with the SM gauge group and $N = 1$ SUSY (or quasi-SUSY), also a massless $(B - L)$ generator must be initially present in any model. If in addition we insist on left- and right-handed fermions respecting the *same* $N = 1$ SUSY, the brane configurations are forced to have intersections giving rise to Higgs multiplets, providing for a rationale for the very existence of the SM Higgs sector.

Keywords: Spontaneous Symmetry Breaking, D-branes, Supersymmetric Effective Theories, Compactification and String Models.

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Contents

1. Introduction

Chirality is probably the deepest property of the standard model (SM) of particle physics. It thus seems that, in trying to build a string theory description of the SM, the first thing we have to obtain is massless chiral fermions. There is a number of known ways in order to achieve chirality in string theory models. From the point of view of explicit D-brane model building there are however only a few known ways. One of the simplest is to locate D-branes (e.g., D3-branes) at some (e.g., orbifold or conifold) singularity in transverse space [1]. Explicit models with branes at singularities have been built [2, 3] with three generations and massless spectrum close to that of the SM or some simple left-right symmetric extension (see also [4, 5, 6]).

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Another alternative way in order to get massless chiral fermions in D-brane models is intersecting branes [7]. Under certain conditions, D-branes intersecting at angles give rise to massless chiral fermions localized at the intersections. Explicit D-brane models with realistic three-generation particle spectrum lying at intersecting branes have been constructed in the last couple of years [8]–[17] (see also [18]). One of the nice features of this type of constructions is that family replication appears naturally as a consequence of the fact that in a compact space branes may intersect a multiple number of times. Another attractive property of this scenario is that quarks and leptons corresponding to different generations are localized at different points in the transverse compact space. It has been suggested [10, 12] that this may provide a geometrical understanding of the observed hierarchy of quark and lepton masses.

In ref. [12] a particularly interesting class of models was found, with a massless chiral fermion spectrum identical to that of the SM of particle physics. They were obtained from sets of intersecting D6-brane wrapping an (orientifolded) six-torus [8]. These models are non-supersymmetric but it has recently been found [15] that analogous models with one $\mathcal{N} = 1$ supersymmetry at each brane intersection may be obtained by appropriately varying the geometry (complex structure) of the torus. In these models each intersection respects in general a *different* $\mathcal{N} = 1$ supersymmetry so that the model is globally non-supersymmetric but in some sense locally supersymmetric. These type of models are called in [15] *quasi-supersymmetric* (Q-SUSY) and have the property that quantum corrections to scalar masses appear only at two-loops. The different $\mathcal{N} = 1$ structure of each intersection in this class of models may be represented in terms of *SUSY quivers*. This Q-SUSY property may be of phenomenological interest in order to stabilize the hierarchy between the weak scale and a fundamental scale of order 10-100 TeV present in low string scale models.

The present article has two main purposes:

i). We construct explicit D6-brane Q-SUSY models with three quark-lepton generations and study some of their properties. All of these models have four stacks of D6-branes: the *baryonic*, *leptonic*, *left* and *right* stacks and either the SM group or a slight generalization with an extra U(1). We show that there are just four classes of quivers (see figure 6) with four stacks of branes and no SU(3) anomalies yielding realistic models. We call them the triangle, linear, square and rombic quivers. Each of these classes of quivers differ on the $\mathcal{N} = 1$ SUSY preserved by the quark, lepton and Higgs sector. In some models (*triangle quiver*) all of them preserve the same $\mathcal{N} = 1$ SUSY whereas in others e.g., left-handed fermions preserve one supersymmetry and right-handed fermions a different one (*linear quiver*). The (*square quiver*) class corresponds to Q-SUSY versions of the models of [12] and have the SM fermion spectrum. They have the property that left-handed quarks, right-handed quarks, left-handed leptons and right-handed leptons have each a different $\mathcal{N} = 1$ SUSY inside a $\mathcal{N} = 4$ living in the bulk. Finally, models constructed from the (*rombic quiver*) have the property that all quarks and the Higgs set respect the same $\mathcal{N} = 1$ supersymmetry whereas the leptonic sector does not. The example provided has a left-right symmetric gauge group. Whereas in the linear and square quiver models the Higgs sector is non-SUSY, it is so in the other two, providing explicit examples of models with weak scale stabilization.

There are a couple of results which look fairly general:

- Imposing supersymmetry at all the intersections implies necessarily the presence of a very definite extra U(1), the $U(1)_{B-L}$ familiar from left-right symmetric models.
- If in addition we insist on some left- and right-handed fermions to respect the *same* $\mathcal{N} = 1$ SUSY (as happens e.g. in the triangular and rombic quivers), the brane configurations are forced to have intersections giving rise to SM Higgs multiplets,[1] providing for a rationale for the very existence of the SM Higgs sector.

We also show how slight variations of torus moduli give rise to Fayet-Iliopoulos (FI) terms for the U(1)'s in the theory. These FI-terms are only present for the subset of U(1)'s which become massive by combining with antisymmetric B_2 fields of the closed string sector of the theory. Since those U(1)'s are massive, the scalar D-term masses look like explicit soft SUSY-breaking masses for the sparticles charged under those U(1)'s. We finally analyze the structure of gauge coupling constants in this class of models. For these intersecting brane models gauge couplings do not unify at the string scale. Rather they are inversely proportional to the volume wrapped by each brane. In the case of Q-SUSY models those volumes get a particularly simple expression in terms of the torus moduli, and we give explicit formulae for them.

ii). We describe how the SM Higgs mechanism in intersecting brane world models has a nice geometrical interpretation as a brane recombination process, in which the branes giving rise to $U(2)_L \times U(1)$ recombine into a single brane related to $U(1)_{em}$. In the SM Higgs mechanism the rank of the gauge group is reduced. The stringy counterpart of this rank-reduction is brane recombination.[2] The chiral fermion spectrum in these models is determined by the intersection numbers I_{ab} which are topological in character. As branes recombine, $|I_{ab}|$ decreases, signaling that some chiral fermions become massive. This is the stringy version of the fermions getting masses from Yukawa couplings after the Higgs mechanism takes place. We exemplify this brane recombination interpretation of the Higgs mechanism in some of the specific examples introduced in the paper, although the general physics applies to any intersecting brane model.

[1] This is in contrast with other popular embeddings of the SM like SUSY-GUT, CY_3 or Horava-Witten heterotic compactifications etc., in which the presence of light SM Higgs multiplets is somewhat ad-hoc and misterious.

[2] Note that the SM Higgs mechanism cannot be described by the familiar process in which two parallel branes separate. Brane separation does not lower the rank and corresponds to adjoint Higgsing, which is not what the SM Higgs mechanism requires.

As we said, having some $\mathcal{N}=1$ SUSY at all brane intersections implies generically the presence of a $B-L$ massless generator in the spectrum. We discuss how this symmetry may be Higgssed away in terms of a brane recombination process in which the *leptonic* and *right* branes recombine into a single brane. That recombination process gives masses to the right-handed neutrinos and we argue that, after the SM Higgs/recombination mechanism takes place, the left-handed neutrinos get at some level Majorana masses.

The structure of this paper is as follows. In the next two sections we give short introductions to both the D6-brane toroidal models introduced in [12] and to the concept of quasi-supersymmetric models of ref. [15]. In section 4 we construct several three generation models with SM gauge group (or some simple extension) and with some quasi-supersymmetry. We analyze SUSY-breaking effects from FI-terms for some of those models in section 5 whereas we provide formulae for the gauge coupling constants in section 6. In section 7 we interpret the SM Higgs mechanism in terms of recombination of intersecting branes. Some final comments and general conclusions are presented in section 8. A couple of appendices concerning the presence of extra U(1)'s in Q-SUSY models and models with D6-branes wrapping non-factorizable cycles on the tori are provided.

2. The standard model at intersecting branes revisited

In this section we summarize the construction of the Standard Model fermionic spectrum arising from intersecting brane worlds, as presented in [12]. We refer the reader to this paper for details regarding this construction.

We will consider type IIA string theory compactified on a factorized 6-torus $T^2 \times T^2 \times T^2$. In this setup, we introduce sets of D6-branes with their 7-dimensional volume containing four-dimensional Minkowski space and wrapping 3-cycles $[\Pi]$ of T^6 [8, 9, 11] (for related constructions see also [19]). We will further assume that these 3-cycles can be factorized as three 1-cycles, each of them wrapping on a different T^2. We denote by (n_a^i, m_a^i), $i=1,2,3$ the wrapping numbers of each $D6_a$-brane, on the i^{th} torus, n_a^i (m_a^i) being the number of times the brane is wrapping around the basis vector e_1^i (e_2^i) defining the lattice of the i^{th} torus, as depicted in figure 1.

Actually, we will study an orientifolded version of the theory obtained by modding out by $\Omega\mathcal{R}$ [8], where Ω is the worldsheet parity operator and $\mathcal{R}=\mathcal{R}_{(4)}\mathcal{R}_{(6)}\mathcal{R}_{(8)}$ is a reflection operator with respect to the real axis of each internal complex dimension $Z_j = X_{2j+2}+iX_{2j+3}$, $j=1,2,3$ (that is, $\mathcal{R}_{(4)}: Z_1 \mapsto \bar{Z}_1$, etc.). Since we must have $\Omega\mathcal{R}$ symmetry in our models, for each $D6_a$ brane we introduce we must also add its mirror image $\Omega\mathcal{R}D6_a$ or $D6_{a^*}$, whose geometrical locus will be determined by a reflection on the real axes X_4, X_6, X_8. Being identified, both branes will give rise to the same unitary gauge group. This symmetry under $\Omega\mathcal{R}$ does also imply that we are allowed to consider only some specific choices of T^2. We may either have rectangular tori (as the first two tori in figure 1) or tilted tori (as the third torus in figure 1). When we have a tilted torus, we can easily describe our configurations in terms of fractional wrapping numbers, where the m's may take $\mathbb{Z}/2$ values (see [11, 12] and appendix A for more detailed discussions).

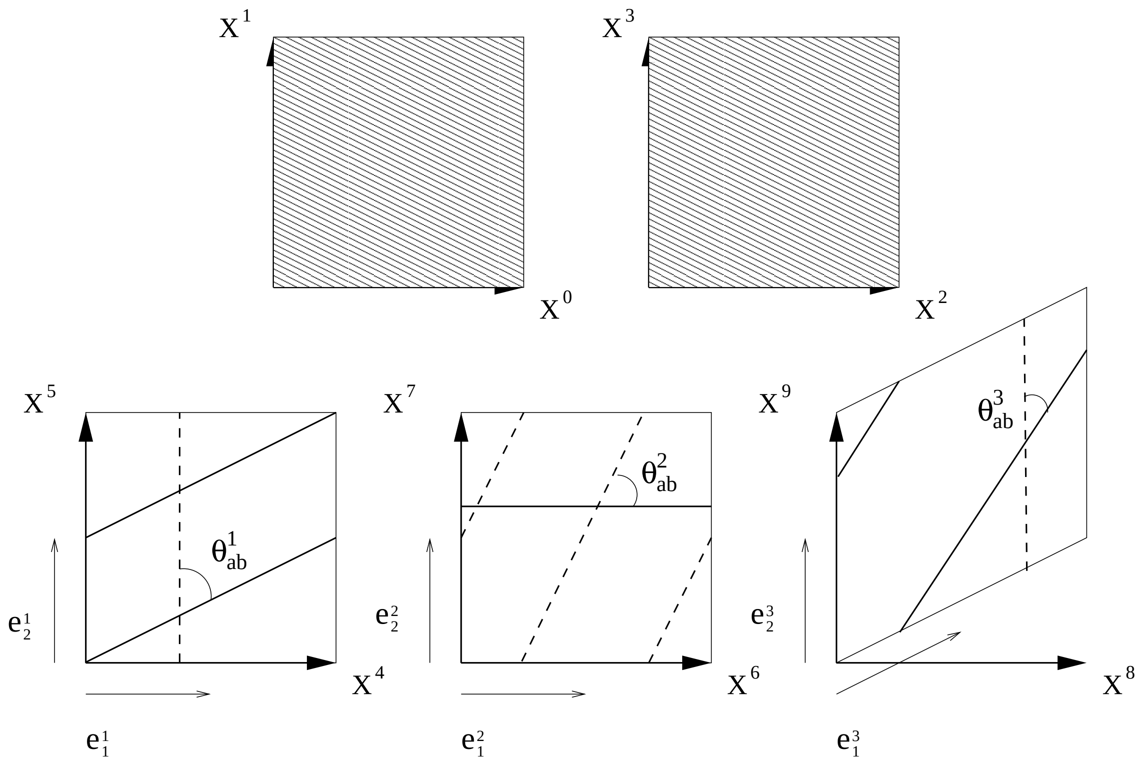

Figure 1: Intersecting brane world setup. We consider two D6-branes containing four non-compact dimensions, to be identified with M_4, and wrapping factorizable 3-cycles of $T^2 \times T^2 \times T^2$. In this specific example the wrapping numbers are $(2,1)(1,0)(1,1)$ (solid line) and $(0,1)(1,2)(0,1)$ (dashed line). Notice that the wrapping numbers of the third torus can be also expressed as fractional cycles, as we have a tilted complex structure (see text). In this fractional language the *dashed* brane remains as before, whereas the *solid* brane wrapping numbers become $(2,1)(1,0)(1,\frac{3}{2})$ (see appendix A). Note that the net intersection number in this example is $2 \times 2 \times 1 = 4$.

The number of times two branes $D6_a$ and $D6_b$ intersect in T^6 is a topological quantity known as the intersection number. When dealing with factorizable branes it can be easily expressed as

$$I_{ab} = \left(n_a^1 m_b^1 - m_a^1 n_b^1\right)\left(n_a^2 m_b^2 - m_a^2 n_b^2\right)\left(n_a^3 m_b^3 - m_a^3 n_b^3\right). \tag{2.1}$$

Open strings stretching around the intersections give rise to chiral fermions in bifundamental representations $(N_a, \bar{N}_b)$ or (N_a, N_b) under the gauge group of the two branes $\mathrm{U}(N_a) \times \mathrm{U}(N_b)$. Thus, these configurations yield I_{ab} copies of the same bifundamental representation. In [12] it was shown how one can build a configuration giving just three generations of a $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)_Y$ gauge group. To achieve this construction one must consider four stacks of D6-branes whose multiplicities and associated gauge group is shown in table 1.

Now, given this brane content, it is relatively easy to achieve the desired SM spectrum as arising from massless fermions living at brane intersections. It is enough to select branes' wrapping numbers (n_a^i, m_a^i) in such a way that the intersection numbers I_{ij}, $i,j = a,b,c,d$

Label	Multiplicity	Gauge Group	Name
stack a	$N_a = 3$	$SU(3) \times U(1)_a$	Baryonic brane
stack b	$N_b = 2$	$SU(2) \times U(1)_b$	Left brane
stack c	$N_c = 1$	$U(1)_c$	Right brane
stack d	$N_d = 1$	$U(1)_d$	Leptonic brane

Table 1: Brane content yielding the SM spectrum.

Intersection	Matter fields		Q_a	Q_b	Q_c	Q_d	Y
(ab)	Q_L	$(3,2)$	1	-1	0	0	1/6
(ab*)	q_L	$2(3,2)$	1	1	0	0	1/6
(ac)	U_R	$3(\bar{3},1)$	-1	0	1	0	$-2/3$
(ac*)	D_R	$3(\bar{3},1)$	-1	0	-1	0	1/3
(bd)	L	$3(1,2)$	0	-1	0	1	$-1/2$
(cd)	N_R	$3(1,1)$	0	0	1	-1	0
(cd*)	E_R	$3(1,1)$	0	0	-1	-1	1

Table 2: Standard model spectrum and U(1) charges. The hypercharge generator is defined as $Q_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d$.

are given by [12]

$$\begin{aligned} I_{ab} &= 1\,, & I_{ab*} &= 2\,, \\ I_{ac} &= -3\,, & I_{ac*} &= -3\,, \\ I_{bd} &= -3\,, & I_{bd*} &= 0\,, \\ I_{cd} &= 3\,, & I_{cd*} &= -3\,, \end{aligned} \tag{2.2}$$

all other intersections vanishing. Here a negative number denotes that the corresponding fermions should have opposite chirality to those with positive intersection number. The massless fermionic spectrum arising from (2.2) is shown in table 2, as well as the charges with respect to the four U(1)'s. Note in this respect that eventually only one U(1) (the hypercharge) remains massless, the other three become massive by swallowing certain closed RR fields, as discussed below.

In [12] a general class of solutions was given for the wrapping numbers (n_a^i, m_a^i) giving rise to a SM spectrum. These are shown in table 3. In this table we have several discrete parameters. First we consider $\beta^i = 1 - b^i$, with $b^i = 0, 1/2$ being the T-dual NS B-background field discussed in [11] (see also [20, 21, 22] for previous related discussions). As shown there, the addition of this background is required in order to get and odd number of quark-lepton generations. From the point of view of branes at angles $\beta^i = 1$ stands for a rectangular lattice for the i^{th} torus, whereas $\beta^i = 1/2$ describes a tilted lattice allowed by the $\Omega\mathcal{R}$ symmetry. Notice that in the third torus one always has $\beta^3 = 1/2$, as in figure 1. We also have two phases $\epsilon, \tilde{\epsilon} = \pm 1$ and the parameter ρ which can only take the values $\rho = 1, 1/3$. Furthermore, each of these families of D6-brane configurations depend on four integers (n_a^2, n_b^1, n_c^1 and n_d^2). Any of these choices lead exactly to the same massless fermion spectrum of table 2.

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

Table 3: D6-brane wrapping numbers giving rise to a SM spectrum. The general solutions are parametrized by two phases $\epsilon, \tilde{\epsilon} = \pm 1$, the NS background on the first two tori $\beta^i = 1 - b^i = 1, 1/2$, four integers $n_a^2, n_b^1, n_c^1, n_d^2$ and a parameter $\rho = 1, 1/3$.

The wrapping numbers of the complete set of D6-branes have to verify the RR tadpole cancellation conditions [8]

$$\begin{aligned} \sum_a N_a n_a^1 n_a^2 n_a^3 &= 16 \\ \sum_a N_a n_a^1 m_a^2 m_a^3 &= 0 \\ \sum_a N_a m_a^1 n_a^2 m_a^3 &= 0 \\ \sum_a N_a m_a^1 m_a^2 n_a^3 &= 0 \end{aligned} \tag{2.3}$$

which just state that the total RR-charge of the configuration has to vanish. The orientifold modding leads to the presence of $8\beta^1\beta^2\beta^3$ orientifold planes wrapping the cycle $(1/\beta^1, 0)(1/\beta^2, 0)(1/\beta^3, 0)$, each with net RR charge -2 (compared to the charge of a pair of mirror D6-branes), as the first condition shows.

These conditions automatically guarantee cancellation of chiral anomalies. In this class of models the last three conditions are automatically verified whereas the first requires:

$$\frac{3n_a^2}{\rho\beta^1} + \frac{2n_b^1}{\beta^2} + \frac{n_d^2}{\beta^1} = 16\,. \tag{2.4}$$

Note however that one can always relax this constraint by adding extra D6-branes with no intersection with the previous ones and not contributing to the rest of the tadpole conditions. In particular, any set of branes with wrapping numbers of the form

$$N_h \left(\frac{1}{\beta_1}, 0\right) \left(\frac{1}{\beta_2}, 0\right) (n_h^3, m_h^3) \tag{2.5}$$

does not intersect the branes from table 1 whereas contributes with $N_h n_h^3/(\beta_1\beta_2)$ to the first condition (without affecting the other three). We will call such D6-brane a *hidden brane*.

As advanced in [15], we can implement this same idea in a more general situation. Let us consider an arbitrary configuration of D6-branes whose low energy chiral spectrum has neither chiral nor mixed anomalies. Since tadpoles imply anomaly cancellation but not the other way round, it could happen that the brane content giving rise to such anomaly-free spectrum did not satisfy tadpoles. We must then 'complete' our model by simply adding the necessary brane content to cancel RR charges. In particular, we can complete our brane configuration by a single brane, let us call it H-brane. Unlike in (2.5), in general this H-brane will not have a simple expression, but will be wrapping a non-factorizable cycle[3] $[\Pi_H]$ of T^6. The important point to notice is that, as long as the low-energy fermion spectrum is anomaly-free, this H-brane will have no net intersection with any brane belonging to our initial configuration, just as in our previous example. Thus, it will constitute a hidden sector, and its presence will not imply the existence of new chiral massless fermions in our low energy spectrum.

Alternatively, and following the recent proposal in [23] (see also [24]), we could equally well turn on an explicit NS-NS background flux H_{NS} in our configuration. This flux will have an associated homology class $[H_{NS}]$, behaving as an RR source just as a D6-brane in this same homology class would do. Thus, we can consider configurations where tadpoles cancel by a combination of D6-brane content and NS-NS flux, the presence of the latter not implying new chiral spectrum. Furthermore, the presence of these fluxes may relax even more the model-building constraints, since the addition of H_{NS} can compensate the anomaly generated by an anomalous $D = 4$ chiral spectrum. The reason for this is that the presence of H_{NS} will induce a Wess-Zumino term in the low energy lagrangian, which will cancel the $\mathrm{SU}(N_a)^3$ and $\mathrm{U}(1)_a - \mathrm{SU}(N_b)^2$ anomaly developed from the naive 'fermion content' point of view (see [23]). Unlike the addition of some H-brane to complete an anomalous configuration, such background flux will not imply new chiral content to our spectrum.[4] We will use the addition of both RR sources in some of the explicit models built in section 4. However, in order to limit the arbitrariness, we will only consider the possible presence of H-flux induced Wess-Zumino terms for anomalous massive U(1)'s so that none of the anomalies of the SM gauge groups have to be canceled via Wess-Zumino terms.

One of the most interesting aspects of this class of theories is the structure of abelian gauge symmetries. The four U(1) symmetries Q_a, Q_b, Q_c and Q_d have clear interpretation in terms of global symmetries of the standard model. Indeed, Q_a is $3B$, B being the baryon number, and Q_d is nothing but lepton number.[5] Concerning Q_c, it is twice I_R, the third component of right-handed weak isospin familiar from left-right symmetric models. Finally Q_b has the properties of a Peccei-Quinn symmetry, having mixed SU(3) anomalies. Two out of the four U(1)'s have triangle anomalies which are canceled by a generalized Green-Schwarz mechanism [12]. The anomalous symmetries correspond to generators Q_b and $3Q_a + Q_d$, the latter being identified with $(9B + L)$. For a general brane configuration

[3] See appendix A for a discussion of non-factorizable cycles.

[4] If we added such a D6-brane H in order to complete an *anomalous* configuration, its net intersection number would not vanish for some brane a contained in it. Clearly, in this case we would not be allowed to call this extra brane a *hidden brane*.

[5] In all the brane intersection models discussed in this article the baryon number symmetry $\mathrm{U}(1)_a$ is automatically gauged. Although the corresponding generator becomes eventually massive, baryon number remains as an accidental symmetry in perturbation theory [12]. Thus the proton is perturbatively stable.

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the anomaly cancellation mechanism goes as follows. There are four closed string antisymmetric fields B_2^I, $I = 0,1,2,3$ which couple to the four $\mathrm{U}(1)_a$ fields associated to each set of N_a D6-branes in the form:

$$\int_{M_4} \sum_a N_a \left(m_a^1 m_a^2 m_a^3 B_2^0 + \sum_I n_a^J n_a^K m_a^I B_2^I \right) \wedge F_a \,, \qquad I \neq J \neq K \tag{2.6}$$

On the other hand, the Poincare duals of those antisymmetric fields, denoted C^I, $I = 0,1,2,3$ couple to both abelian and non-abelian fields F_b as follows:

$$\int_{M_4} \sum_b \left(n_b^1 n_b^2 n_b^3 C^0 + \sum_I n_b^I m_b^J m_b^K C^I \right) \wedge F_b \wedge F_b \,. \tag{2.7}$$

The combined effect of these two couplings cancel the residual $\mathrm{U}(1)_a \times \mathrm{U}(N_b)^2$ triangle anomalies by the tree level exchange of the four RR fields. At the same time the $B_2^I \wedge F$ couplings in (2.6) give masses to the following four linear combinations of abelian fields

$$\begin{aligned}
B_2^0 &: \quad \sum_a N_a m_a^1 m_a^2 m_a^3 F_a \,, \\
B_2^1 &: \quad \sum_a N_a m_a^1 n_a^2 n_a^3 F_a \,, \\
B_2^2 &: \quad \sum_a N_a n_a^1 m_a^2 n_a^3 F_a \,, \\
B_2^3 &: \quad \sum_a N_a n_a^1 n_a^2 m_a^3 F_a \,.
\end{aligned} \tag{2.8}$$

Since in the specific models we are constructing at least one of the m_a's of each brane vanishes, the first linear combination is trivially vanishing and there is always a massless U(1). The three linear combinations which are massive are :

$$\begin{aligned}
&Q_b \\
&3Q_a + Q_d \\
&Q_m = \left(3\beta^2 n_a^2 Q_a + 6\rho\beta^1 n_b^1 Q_b - 2\tilde{\epsilon}\beta^1 n_c^1 Q_c - 3\rho\beta^2 n_d^2 Q_d - 2N_h \tilde{\epsilon} m_h^3 Q_h \right)
\end{aligned} \tag{2.9}$$

where for the sake of generality we have included the effect of a set of N_h parallel "hidden" branes as discussed above. The first two massive generators are model independent and correspond to the anomalous U(1)'s, whereas Q_m is model-dependent and anomaly-free. The massless U(1)'s are:

i). $m_h^3 = 0$

In this case the presence of a hidden brane does not affect the form of the massless linear combinations which are given by:

$$Q_0 = \frac{1}{6} Q_a + \frac{r}{2} Q_c - \frac{1}{2} Q_d \,, \qquad r = \frac{\tilde{\epsilon}\beta^2}{2n_c^1 \beta^1} \left(n_a^2 + 3\rho n_d^2 \right) \tag{2.10}$$

and the hidden brane charge Q_h. Note that for $r = -1$ the above generator is the standard hypercharge generator, as discussed in [12]. In this case only this standard hypercharge generator remains massless.[6]

ii). $m_h^3 \neq 0$

In this case the hidden brane affects the form of the visible linear combination. The massless linear combinations are:

$$\begin{aligned} Q_1 &= \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d + (1+r)\frac{n_c^1\beta^1}{2N_h m_h^3}Q_h \\ Q_2 &= \frac{1}{3}Q_a - Q_d + \frac{rn_c^1\beta^1}{N_h m_h^3}Q_h \end{aligned} \tag{2.11}$$

Note that the first of these charges couples like standard hypercharge to the SM particles, whereas the second one couples like $B-L$. Thus, in this more general case two massles U(1)'s coupling to quarks and leptons remain in the massless spectrum.

The general solutions yielding the SM spectrum have the following geometrical properties. The Baryonic (a) and Leptonic (d) stacks are parallel in the first complex dimension and that is why they do not intersection Thus, no lepto-quarks fields appear in our massless spectrum. On the other hand, Left (b) and Right (c) stacks are parallel in the second complex dimension, again not intersecting. Something similar happens for each brane i and its mirror i^*, $i = a, b, c, d$. Strings exchanged between stacks b and c have the quantum numbers of SM Higgs (and Higgsino) fields and eventually we would be interested in some of these scalar states to be relatively light, so as to play the role of Higgs fields. So we will asume that the distance between a and d branes is bigger than that between branes b and c, so that the latter can provide us with a Higgs system, as we will describe in more detail in section 7. Note that we will consider the different distances between branes to be fixed quantities. Hence, they will play the role of external parameters of our models, very much like the geometrical moduli.

Notice that the whole of the previous SM construction has been achieved by using factorizable D6-branes, that is, branes whose geometrical locus can be described as product of 1-cycles, each wrapping a different T^2. This kind of construction is preferred from a model-building point of view, since it has a particularly simple associated geometry. This allows us to easily compute some phenomenologically important quantities as, for instance, scalar masses at intersections. However, we could have also constructed a model in a more general set-up, where the building blocks of our configuration were D6-branes wrapping *general* 3-cycles of our six-torus. In this way, we could have equally well realised a D6-brane tadpole-free configuration giving rise to a SM spectrum, the factorizable models just presented being a particular subfamily of the latter.

[6] As noted in [12], if in addition one has $n_c^1 = 0$, then an extra generator $(1/3)Q_a - Q_d$ (which corresponds to B-L) remains also massless.

For simplicity and better visualization of our constructions, we will give as examples models where the relevant physics arises at factorizable D6-branes configurations, so as to extract the phenomenologically interesting quantities more straightforwardly. However, it is important to notice that reached this point it is not possible to ignore non-factorizable D6-branes anymore. In particular, when studying the recombination process of two branes into a third one (corresponding to some Higgs mechanism) such non-factorizable branes will naturally arise. As we will see, this is due to the fact that conservation of RR charge imposes that this third brane should wrap a 3-cycle which is the sum of the two former 3-cycles. In general, this will yield a non-factorizable brane. In order to appropiately describe these less intuitive objects, we have made use of a "q-basis" formalism, which is described in appendix A. For some results regarding the geometry of branes wrapping general cycles see also [25].

Let us also emphasize that in the present paper we will be dealing mostly with the open string sector of the theory, which is the one which may give rise to the SM physics. We will not discuss here the closed string NS potential and its associated NS tadpoles. We will rather take the value of NS geometric moduli as well as brane positions as frozen external parameters defining the geometry of the configuration. Some work on the NS-tadpole structure of this class of theories may be found in, e.g., ref. [13, 15, 17].

3. Intersecting D6-branes and supersymmetry: SUSY & Q-SUSY models

An interesting question is whether one can construct intersecting brane models analogous to those discussed in the previous section and with $\mathcal{N} = 1$ supersymmetry. As discussed in refs. [8, 15], the answer to this question is no, if we restrict ourself to purely toroidal models (no orbifolds) with factorized D6-branes and impose RR tadpole cancellation conditions. Models with $\mathcal{N} = 1$ SUSY may be constructed if an additional $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifolding is perfomed [14], but at the cost of losing the simplicity of the spectrum, as exotic chiral particles beyond the SM appear.

On the other hand it was shown in [15] the possibility of having all brane intersections preserving *some* unbroken $\mathcal{N} = 1$ SUSY, although not necessarily the same one. These configurations were called quasi-supersymmetric, (Q-SUSY). Roughly speaking, a Q-SUSY field theory consists of different subsectors each preserving at least $\mathcal{N} = 1$ supersymmetry but, not being the same supersymmetry for every sector, the system as a whole has $\mathcal{N} = 0$. Thus, we expect corrections that spoil the boson-fermion mass degeneracy appearing in a truly supersymmetric theory, this effects appearing only at two-loops in perturbation theory. There might also be some *massive* sectors not respecting any supersymmetry at all. In any case, the radiative corrections will be fairly supressed with respect to an ordinary $\mathcal{N} = 0$ field theory. In particular, scalars will have a two-loop protection against quadratic divergences. This may be of interest when dealing with scenarios where the fundamental scale of physics is two or three orders of magnitude above the electroweak scale [15]. A two-loop protection would be enough to understand a 2-3 orders of magnitude hierarchy between the weak scale and a fundamental scale of order 10-100 TeV.

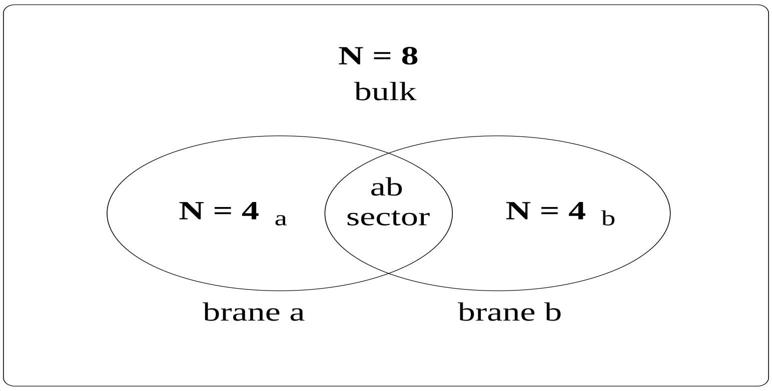

Figure 2: Schematic representation of the supersymmetries preserved by a pair of branes in a toroidal compactification. The closed string sector, living on the bulk, preserves a full $D = 4$ $\mathcal{N} = 8$ superalgebra, whereas each of the branes $D6_a$ and $D6_b$ preserve $\mathcal{N} = 4$ subalgebras, generically different ones. Whether the $D6_aD6_b$ sector is supersymmetric or not depends on the overlap that exists between these two algebras. In the orientifold case the bulk has only $\mathcal{N} = 4$ SUSY.

Let us describe the D-brane setup that allows for an explicit realization of this Q-SUSY structure. Just as in the construction of the Standard Model presented in last section, we will be dealing with type IIA D6-branes wrapping factorized 3-cycles in $T^2 \times T^2 \times T^2$. Any factorized D6-brane constitutes a $\frac{1}{2}$BPS state that preserves half of the supersymmetries coming from a plain toroidal compactification. This will be reflected in the $D6_aD6_a$ sector of the theory (that is, strings beginning and ending on the same $D6_a$-brane), yielding $D = 4$ $\mathcal{N} = 4$ super Yang-Mills as its low energy spectrum. When considering a pair of such factorized branes, say a and b, each of them will preserve some $\mathcal{N} = 4$ subalgebra of the whole $\mathcal{N} = 8$ bulk superalgebra. Whether these two different subalgebras do have some overlap or not will determine if the $D6_aD6_b$ sector is also supersymmetric (see figure 2).

In general, the answer will be positive whenever these two branes are related by a SU(3) rotation in compact dimensions [7]. When dealing with D6-branes at angles, this can be easily computed in terms of the spectrum living at the $D6_aD6_b$ sector. Whereas the Ramond sector always yields a massless chiral fermion, the lightest states coming from the Neveu-Schwarz sector will be four scalars with masses [9]

$$\begin{array}{l} \alpha' \mathbf{Mass^2} \\ \frac{1}{2}(-|\vartheta^1_{ab}| + |\vartheta^2_{ab}| + |\vartheta^3_{ab}|) \\ \frac{1}{2}(|\vartheta^1_{ab}| - |\vartheta^2_{ab}| + |\vartheta^3_{ab}|) \\ \frac{1}{2}(|\vartheta^1_{ab}| + |\vartheta^2_{ab}| - |\vartheta^3_{ab}|) \\ 1 - \frac{1}{2}(|\vartheta^1_{ab}| + |\vartheta^2_{ab}| + |\vartheta^3_{ab}|)\,, \end{array} \tag{3.1}$$

where ϑ^i_{ab} represent the angles between branes a and b in the i^{th} torus (see figure 1). The number of SUSY's shared by sectors $D6_aD6_a$ and $D6_bD6_b$ will correspond to the number of massless scalars in (3.1). Indeed, if for instance $\vartheta^1_{ab} = \vartheta^2_{ab} + \vartheta^3_{ab}$, then the first of such scalars will be massless, and the $D6_aD6_b$ sector will consist of $|I_{ab}|$ $\mathcal{N} = 1$ chiral multiplets.

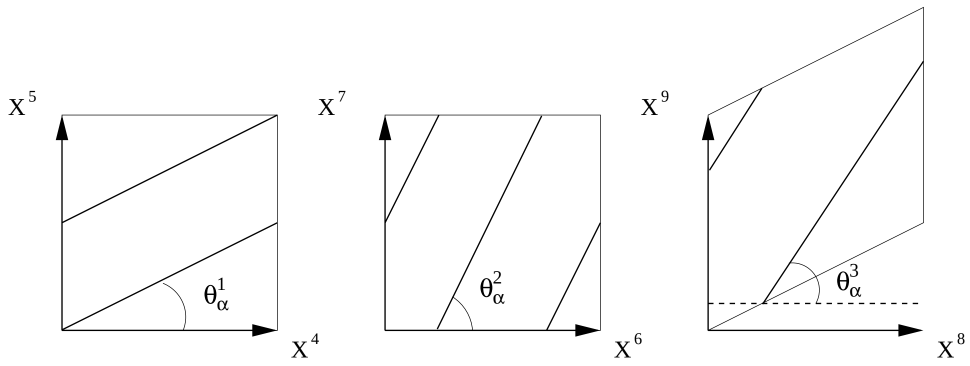

Figure 3: Definition of the twist vector $v_\alpha := (\theta^1_\alpha, \theta^2_\alpha, \theta^3_\alpha)$.

In an orientifold compactification, however, a generic D6-brane a cannot exists on its own, but will be accompanied by its mirror image a^*. In order to achieve our Q-SUSY structure, we will require any brane and its mirror to preserve at least one common SUSY. In fact, we will require them to share a $\mathcal{N} = 2$ subalgebra.[7] As discussed in [15], there exist six types of such branes. In order to describe them, let us define a *twist* vector $v_\alpha = (\theta^1_\alpha, \theta^2_\alpha, \theta^3_\alpha)$ whose three entries contain the relative angles θ^i_α between a D6-brane α and the horizontal axis of the i^{th} torus (see figure 3).

In terms of this twist vector, the six different types of branes can be listed as

type of brane	twist vector	SUSY	
a_1 branes :	$v_{a1} = (0, \theta_{a1}, \theta_{a1})$	r_2, r_3	
a_2 branes :	$v_{a2} = (0, \theta_{a2}, -\theta_{a2})$	r_1, r_4	
b_1 branes :	$v_{b1} = (\theta_{b1}, 0, \theta_{b1})$	r_1, r_3	
b_2 branes :	$v_{b2} = (\theta_{b2}, 0, -\theta_{b2})$	r_2, r_4	
c_1 branes :	$v_{c1} = (\theta_{c1}, \theta_{c1}, 0)$	r_1, r_2	
c_2 branes :	$v_{c2} = (\theta_{c2}, -\theta_{c2}, 0)$	r_3, r_4	(3.2)

where we have suppressed the upper index i. We have named the generators of each $\mathcal{N} = 1$ algebra by r_j, $j = 1, 2, 3, 4$, just in order to keep track of the supersymmetries present on each sector (for a more detailed description see [15]). Each type of brane in (3.2) shares a different $\mathcal{N} = 2$ subalgebra with their mirror brane, which can be represented by a pair of generators.[8] Notice that the twist vector of a mirror brane α^* will be given by $v_{\alpha^*} = -v_\alpha$, whereas the angles between two branes α and β are the components of $v_{\alpha\beta} := v_\beta - v_\alpha$.

When considering a configuration where several of such $\mathcal{N} = 2$ branes appear, we must also consider sectors corresponding to D6-brane intersections. The full Q-SUSY structure can be encoded in a hexagonal quiver-like diagram, shown in figure 4. Each node of this

[7] From the model-building point of view, this is required in order to avoid massless chiral multiplets in the $D6_a D6_{a^*}$ sector, which does only provide us with matter transforming as symmetric and antisymmetric representations. We will forbid such kind of exotic spectrum right from the start.

[8] In a D-brane language, the twist vector v_a contains the relative angles between the D6-brane a and the O6-plane, who lies on the cycle $(1/\beta^1, 0)(1/\beta^2, 0)(1/\beta^3, 0)$. Hence, it encodes the supersymmetries shared by both. If we name the generators of the $\mathcal{N} = 4$ algebra of the O6-plane by r_1, r_2, r_3, r_4, then the six types of branes described above will correspond to different choices of a $\mathcal{N} = 2$ subalgebra.

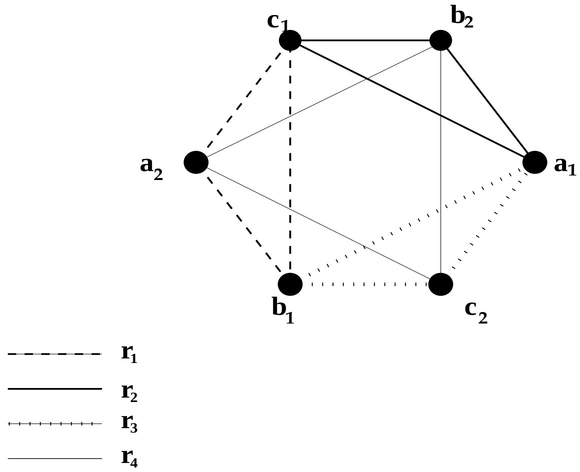

Figure 4: General SUSY quiver diagram. Nodes represent the six different types of branes (plus mirrors) whereas links represent the chiral intersections of those branes. There are four types of links corresponding to the four SUSY's in the bulk.

diagram represents one of the branes in (3.2), and at the same time its mirror image. Notice that gauge groups at intersecting D6-branes models always arise from $D6_\alpha D6_\alpha$ sectors, hence they will be localized at these nodes. Just as in the toroidal case, each gauge group will correspond to a $\mathcal{N}=4$ subsector of the theory, while the matter content living on the $D6_\alpha D6_{\alpha^*}$ sector presents a non-chiral, generically massive $\mathcal{N}=2$ spectrum.

Each of the branes in (3.2) will generically intersect with four other types of brane. Such intersection is represented in our hexagonal quiver by a link joining two nodes. In general, a non-vanishing intersection number $I_{\alpha\beta}$ does signal that some chiral spectrum lives at the $D6_\alpha D6_\beta$ sector. It can easily be seen that any of these sectors will also preserve some $\mathcal{N}=1$ SUSY, yielding chiral multiplets transforming in bifundamentals. Let us take, for instance, the link joining a_2 with b_2. Each node, as a combined system of a brane and its mirror, preserves two supersymmetries, having r_4 in common. Thus, the link between these two nodes will contain $I_{a_2b_2}$ $(N_{a_2}, \bar{N}_{b_2})$ + $I_{a_2b_2^*}$ (N_{a_2}, N_{b_2}) chiral multiplets under the supersymmetry generator r_4. Finally, branes corresponding to opposite nodes (as for instance a_1 and a_2) do never intersect, since they are parallel in one of the tori. This $D_{\alpha_1}D_{\alpha_2}$ sector will contain a non-chiral, generically massive, $\mathcal{N}=0$ spectrum.

We thus see that any low energy effective field theory coming from a D-brane configuration whose building blocks belong to (3.2) will yield a Q-SUSY field theory. In particular, the general hexagonal construction depicted in figure 4 will contain as massless sectors $\mathcal{N}=4$ vector multiplets and $\mathcal{N}=1$ chiral multiplets, whereas some non-chiral $\mathcal{N}=2$ and $\mathcal{N}=0$ massive sectors may also arise. As we mentioned at the beginning of this section, this will protect scalar masses from one-loop corrections, the first non-vanishing contributions coming from two-loops in perturbation theory [15]. There are contributions coming both from light particle exchange (figure 5*a*) and heavy non-SUSY particle exchange (figure 5*b*). All will give corrections to scalar masses of order $\alpha_i/(4\pi)M_s$. For the $\mathcal{N}=4$ gauginos the only source of masses will be the one-loop exchange of massive $\mathcal{N}=0$ sectors of the theory (see figure 5*c*).

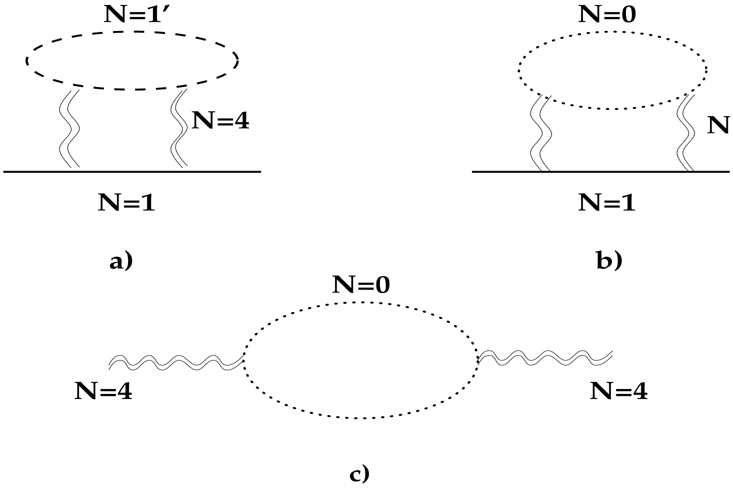

Figure 5: First non-vanishing loop contributions to the masses of scalar fields and gauginos in a Q-SUSY model: *a*) Quadratic divergent contribution present when the upper loop contains *massless* fields respecting different supersymmetries than those of the fields below; *b*) Contribution coming from possible *massive* non-supersymmetric states in the upper loop, *c*) One-loop contribution to the masses of gauginos from heavy non-supersymmetric states.

These general considerations will equally well apply to any subquiver of the hexagonal construction just presented. In particular, we will be interested in models containing four stacks of branes. This is essentially because, just as in the explicit example presented in last section, we need at least four branes in order to arrange the SM chiral spectrum as coming from bifundamental representations. Let us then consider some four-stack subquivers arising from the hexagonal construction. Having four stacks implies that at least there is a pair of them with vanishing intersection, thus yielding either a $\mathcal{N}=2$ or a $\mathcal{N}=0$ sector. Without loss of generality, we will take two a-type branes to be such pair, and one of them to be of one definite kind, say a_2. In order to have a non-trivial chiral spectrum, we need the other two branes to be of type α_i, with $\alpha \neq a$. Our general four-stack configuration will be of the form $(a_2, a'_i, \beta_j, \gamma_k)$. There are some inequivalent choices:

- Either $i = 2$, having a massive $\mathcal{N}=2$ $D6_aD6_{a'}$ sector,
  or $i = 1$, having a massive $\mathcal{N}=0$ $D6_aD6_{a'}$ sector.

- Either $\beta = \gamma$, having a non-chiral spectrum on the $D6_\beta D6_\gamma$ sector,
  or $\beta \neq \gamma$, having a $\mathcal{N}=1$ chiral spectrum.

- Either $j = k$, so that at least one of the a branes preserves the same supersymmetry with both β and γ branes,
  or $j \neq k$, so that both supersymmetries are different.

We thus see that we have eight inequivalent possibilities when considering four-stack subquivers, as the rest of them just amount of some relabeling of the nodes. However, performing a simple anomaly cancellation analysis, we can reject half of them as phenomenologically uninteresting. This comes from the fact that, when performing our model-building,

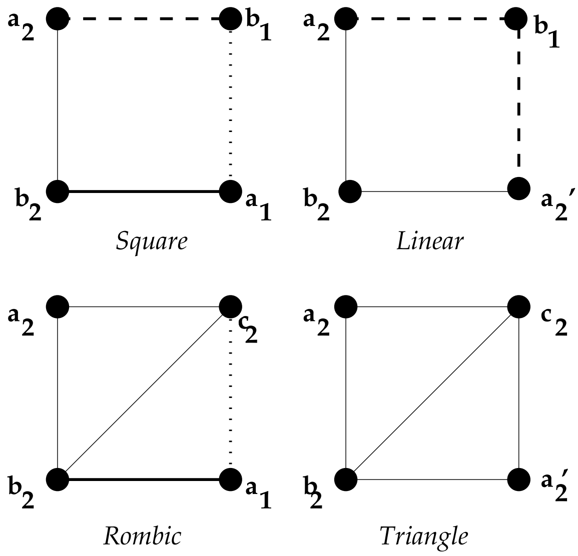

Figure 6: Simplest SUSY-quivers with four nodes. They are obtained from the general hexagonal quiver by deleting two nodes. Other equivalent quivers are obtained by cyclic relabeling $a \to b \to c$. The rightmost examples may be depicted as linear (triangle) quivers respectively if one locates the a_2' branes on top of their parallel relatives a_2.

we will usually choose (a, a') to be the Baryonic and Leptonic branes, respectively. Being a non-intersecting pair of branes, this will avoid massless lepto-quarks in our spectrum. Now, in order to cancel SU(3) anomalies, we need fermions transforming both in **3** and in $\bar{\mathbf{3}}$ representations, same number of them on each. It can be easily shown that, in order for this to be true in our models, we must have either $\beta = \gamma$ and $j \neq k$, or $\beta \neq \gamma$ and $j = k$. So we are finally led to consider only four subquivers, which are depicted in figure 6.

In the next section we will try to combine these four Q-SUSY structures with phenomenologically appealing brane configurations, following the philosophy described in section 2 for getting the SM from general intersecting branes.

4. Model-building: three generation models with SUSY intersections

In this section we are going to consider the construction of explicit intersecting D6-brane models of phenomenological interest. In particular we will focus in the construction of models with quasi-SUSY in the sense described in previous section and ref. [15] with three quark-lepton generations and SM group (or some simple left-right extension). We will concentrate in the class of branes discussed above in which each individual brane preserves an unbroken $\mathcal{N} = 2$ supersymmetry with its mirror brane, having null intersection with it.

As was discussed above, in this class of brane configurations only up to three U(1)'s may become massive by combining with three of the four antisymmetric fields $B^i_{\mu\nu}$ present in the massless spectrum [12]. We will be considering four stacks of branes a, b, c and d each of them containing $N_a = 3$, $N_b = 2$, $N_c = 1, 2$ and $N_d = 1$ parallel branes. Thus to start with our gauge group will be either $\mathrm{U}(3)_a \times \mathrm{U}(2)_b \times \mathrm{U}(1)_c \times \mathrm{U}(1)_d$ or $\mathrm{U}(3)_a \times \mathrm{U}(2)_b \times \mathrm{U}(2)_c \times \mathrm{U}(1)_d$.

The most general phenomenologically interesting Q-SUSY models one can consider involving four stacks of branes may be depicted by the four SUSY-quivers shown in figure 6. Locating the four stacks of branes of the SM at the four corners of those quivers we will get Q-SUSY models with the SM gauge group (and possibly some extra U(1) symmetry as discussed below). Without loss of generality, we will take the Baryonic stack to to be of type a_2, whereas the Left stack can also be chosen to be of some definite type different from a, say b_2. Those two stacks will yield a gauge group $\mathrm{U}(3) \times \mathrm{U}(2)$. Clearly, the leptonic brane must be non-intersecting with the baryonic brane, since we want to avoid lepto-quark particles. It must then be of the same type a as the latter. Depending on what kind of brane one chooses for the right and leptonic stacks, different model building possibilities will be obtained:

i). Square quiver (figure 7). In this case there are four types of intersections respecting four different types of SUSY. They will correspond to Q-SUSY versions of the three generation configurations of ref. [12].

ii). Linear quiver (figure 8). This quiver may be put into a line by putting the leptonic a_2' brane on top of the baryonic stack a_2. It may be considered as a variant of the square quiver in which the left-handed fermions have one supersymmetry and right-handed leptons a different one.

iii). Rombic quiver (figure 9). In this third quiver three different SUSY's are present. One main difference with the previous cases is the presence of Higgs chiral multiplets at the intersections of right and left stacks. Note that the baryonic and Higgs sub-sector will respect the same $\mathcal{N} = 1$ supersymmetry, whereas the leptonic sector respect other ones.

iv). SUSY-triangle (figure 10). In this case baryon and lepton branes are of the same (a_2) type so that the quiver may be deformed into a triangle by locating one on top of each other. Note that now all intersections preserve the same $\mathcal{N} = 1$ SUSY and, as in the rombic quiver, there are massless Higgs multiplets at the intersections.

It turns out that is is easy to find families of examples corresponding to the *square* and *linear* quivers. On the contrary, the *triangle* and *rombic* quivers are more restricted configurations and examples are more scarce. Let us emphasize that all of these brane configurations lead in general to RR-tadpoles. Those may be canceled in a model by model basis by adding extra D-branes (with no intersection with the ones in the quivers, so that the massless chiral spectrum is not modified) and/or the addition of explicit H_{NS}-flux, as described in section 2.

There are some general aspects of these Q-SUSY configurations which can be given without the need of going into the details of the models. One can see that the following properties hold:

1. "In any square, rombic and linear quiver there is, at least, one U(1) generator which remains massless (does not combine with an antisymmetric B-field). In the triangle quiver there are two such massless U(1)'s."

2. "In any SUSY or Q-SUSY (factorizable) D6-brane configuration yielding SM fermion spectrum the $\mathrm{U}(1)_c$ (right) gauge generator is massless."

3. 'If we insist that the left- and right-handed counterparts of the same field respect the *same* $\mathcal{N}=1$ SUSY, then necessarily the brane configuration must have intersections at which multiplets with the quantum numbers of SM Higgs arise'.

The first two general properties are proven in appendix B. The second one follows because it turns out that supersymmetry at all intersections forces the right-stack to have wrapping numbers $m_c^2=0$, $n_c^1=n_c^3=0$. Notice that this massless generator $\mathrm{U}(1)_c$ may or may not coincide with the one(s) mentioned in point 1. Concerning the third one, it is a direct consequence of the quiver classification above: in order for the *Baryonic* stack (of type a_i) to have intersections giving rise to e.g., left- and right-handed quarks respecting the same $\mathcal{N}=1$ SUSY, the *Left* and *Right* branes must be of type b_i and c_i respectively (or viceversa, see figure 6). But branes of type b_i and c_i necessarily intersect, giving rise to massless states with the quantum numbers of SM Higgs multiplets. This is the link drawn between b_2 and c_2 in the lowest quivers in figure 6.

Due to these properties, there is a number of limitations on the possible model building, which we will discuss in a model by model basis below. As a general consequence we will see that in general Q-SUSY forces the hypercharge to come along with an extra $\mathrm{U}(1)_{B-L}$ gauge boson. This general property can only be evaded if Q-SUSY is only an approximate symmetry, as will be discussed in the case of the square SUSY-quiver. Otherwise a gauged $B-L$ must be present, although of course it may be eventually Higgssed. This is an interesting result by itself since it shows an intriguing connection between the presence of supersymmetry and the gauging of $(B-L)$. It is also intriguing the third general property above, since it gives an interesting connection between the presence of SUSY and the existence of Higgs multiplets. We will present specific examples corresponding to the four types of quivers discussed above.

Before presenting the models let us emphasize that these are not the unique models that one can construct starting with the four quiver configurations. Other variations may be constructed depending on how we cancel RR-tadpoles (by explicit factorizable branes and/or non-factorizable branes and/or explicit RR H-flux). One can also easily build models with extended gauge symmetries like models with left-right and/or Pati-Salam symmetries. But we think the examples provide the reader with enough information to look for other interesting variations. A general summary of the models we construct here is given at the end of the paper in table 17.

4.1 The square quiver

In order to describe explicit examples of square quivers let us start from the generic family of models already presented in table 3. As we previously said, this family is parametrized by two phases $\epsilon,\tilde{\epsilon}=\pm1$, the NS background on the first two tori $\beta^i=1-b^i=1,1/2$, four integers n_a^2,n_b^1,n_c^1,n_d^2 and a parameter $\rho=1,1/3$. As we also remarked, without loss of generality one can choose the baryonic(left)-stacks to be of type a_2 (b_2). This implies setting $\epsilon=\tilde{\epsilon}=1$. In order to find out whether some SUSY is unbroken at the intersections, let

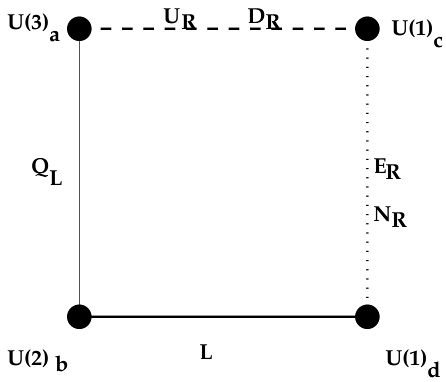

Figure 7: A square SUSY-quiver with SM spectrum. Nodes represent the D6-branes giving rise to the SM gauge group. Links represent the brane intersections, at which quarks and leptons live. The different styles of the links denote that each of the four intersections respects a different $\mathcal{N} = 1$ SUSY.

Brane	θ^1_α	θ^2_α	θ^3_α
a	0	$tg^{-1}\left(\frac{\beta^2 U^2}{n^2_a}\right)$	$tg^{-1}\left(\frac{-\rho U^3}{2}\right)$
b	$tg^{-1}\left(\frac{\beta^1 U^1}{n^1_b}\right)$	0	$tg^{-1}\left(\frac{-3\rho U^3}{2}\right)$
c	$tg^{-1}\left(\frac{3\rho\beta^1 U^1}{n^1_c}\right)$	0	$\frac{\pi}{2}$
d	0	$tg^{-1}\left(\frac{\beta^2 U^2}{\rho n^2_d}\right)$	$tg^{-1}\left(\frac{3\rho U^3}{2}\right)$

Table 4: Components of the twist vector v_α for each D6-brane stack of table 3.

us display the angles that compose the twist vector v_α for each brane $\alpha = a, b, c, d$. Those angles are shown in table 4 in terms of the ratios $U^i = R^i_2/R^i_1$, where R^i_1,R^i_2, $i = 1, 2, 3$ are the radii of the tori.

In order for some SUSY to be preserved at each intersection, one needs to have angles $\vartheta^1_{\alpha\beta} \pm \vartheta^2_{\alpha\beta} \pm \vartheta^3_{\alpha\beta} \in \mathbb{Z}$ for some choice of signs, where $\vartheta^i_{\alpha\beta}$, $i = 1, 2, 3$ are the angles formed by any pair of branes α and β in the three 2-tori [15]. Looking at table 4 one can easily see that for certain choices of the U^i and particular choices of some of the integers involved one can manage so that all intersection preserve one SUSY. Specifically, if the conditions

$$n^1_c = 0 \Rightarrow \rho = \frac{1}{3}, \qquad \beta^1 = 1, \tag{4.1}$$

$$n^1_b > 0, \qquad n^2_a = 3\rho n^2_d = n^2_d > 0, \tag{4.2}$$

$$U^1 = \frac{n^1_b}{2} U^3, \qquad U^2 = \frac{n^2_a}{6\beta^2} U^3, \tag{4.3}$$

are met, the angles of the branes have the general form of table 5, where $\alpha_1 = tg^{-1}(U^3/6)$ and $\alpha_2 = tg^{-1}(U^3/2)$. Note that with these choices the branes a, b, c and d become $\mathcal{N} = 2$ branes of types a_2, b_2, b_1 and a_1 respectively. Thus we recover the square Q-SUSY structure of figure 7. Altogether the wrapping numbers of the resulting models with Q-SUSY properties are displayed in table 6. One can check that there are simple choices

$Brane$	θ_α^1	θ_α^2	θ_α^3
a	0	α_1	$-\alpha_1$
b	α_2	0	$-\alpha_2$
c	$\frac{\pi}{2}$	0	$\frac{\pi}{2}$
d	0	α_2	α_2

Table 5: D6-brane angles when the conditions in (4.3) are met.

N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 3$	$(1, 0)$	(n_a^2, β^2)	$(3, -1/2)$
$N_b = 2$	$(n_b^1, 1)$	$(1/\beta^2, 0)$	$(1, -1/2)$
$N_c = 1$	$(0, 1)$	$(1/\beta^2, 0)$	$(0, 1)$
$N_d = 1$	$(1, 0)$	$(n_a^2, 3\beta^2)$	$(1, 1/2)$
N_h	$(1, 0)$	$(1/\beta^2, 0)$	(n_h^3, m_h^3)

Table 6: D6-brane wrapping numbers giving rise to a Q-SUSY SM spectrum for a square quiver. For the sake of generality we have also considered the possible presence of an extra brane with no intersection with the SM branes.

of the integer parameters and extra branes with no intersection with the SM branes such that all RR tadpoles are canceled. The massless fermionic spectrum is the same as in the models in [12], but now all of them come with a scalar superpartner.

We thus have obtained a family of square Q-SUSY models parametrized by the two integers n_a^2, n_b^1 and the NS parameter β^2, yielding just the spectrum of the standard model. In addition now each chiral fermion will have one massless scalar partner, so there will be the coresponding squarks and sleptons. Thus the initial spectrum (before taking into account the SUSY-breaking effects) will be rather similar to that of the MSSM. It turns out however that the choice of parameters in order to get Q-SUSY (in particular setting $n_c^1 = 0$) changes the structure of the U(1)'s in the model. Indeed, as we already mentioned above, this choice implies that the Q_c generator is massless. In the absence of any extra brane (like that given by the last line in table 6) we would have just Q_c as our only massless generator, rather than hypercharge.

Let us analyze this point in more detail. If we substitute in equation (2.8) the wrapping numbers of the Q-SUSY SM model above, we obtain for the couplings of the antisymmetric fields to the U(1)'s the result

$$\begin{aligned}
&B_2^1 \wedge \frac{2}{\beta^2} F^b \\
&B_2^2 \wedge 3\beta^2 \left(3F^a + F^d\right) \\
&B_2^3 \wedge \left(\frac{-3n_a^2}{2} F^a - \frac{n_b^1}{\beta^2} F^b + \frac{n_a^2}{2} F^d + N_h \frac{m_h^3}{\beta^2} F^h\right)
\end{aligned}$$

Note that the coupling to B_2^0 is zero, because $\prod_i m_a^i = 0 \ \forall a$. Thus we see that, in the absence of additional branes ($N_h = 0$) all U(1)'s but U(1)$_c$ gain masses by combining with antisymmetric tensors. On the other hand, if there are additional h branes present the situation changes. Indeed, recalling eq. (2.11), we can see that we would have two massless U(1)'s: $\mathrm{U}(1)_c$ and the linear combination $\mathrm{U}(1)_a - 3\,\mathrm{U}(1)_d + \frac{3n_a^2\beta^2}{N_h m_h^3}\,\mathrm{U}(1)_h$. In fact, a linear combination of these two gives us a U(1) charge whose coupling to SM fermions is precisely that of the standard hypercharge:

$$Q_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d - \frac{1}{2}\frac{n_a^2\beta^2}{N_h m_h^3}Q_h\,. \tag{4.4}$$

In addition to hypercharge, an extra massless U(1) with analogous couplings to that of a $(B - L)$ generator will remain in the massless spectrum. This is a first realization of our comment above that Q-SUSY requires the gauging not only of hypercharge but of the $B - L$ generator. This latter symmetry may eventually be spontaneously broken, in a way analogous to that discussed in subsection 7.4. An alternative way to guarantee the presence of just standard hypercharge at the massless level will be explained in section 5.1.

An important point is in order concerning the Higgs sector in square quivers of the SM. In these models no Higgs multiplets appear at any of the intersections of the square. As discussed in section 2 and in [12], scalar fields with the quantum numbers of SM Higgs appear if the branes b and c (which are parallel along the second torus) approach to each other. However, since these two branes preserve opposite types of supersymmetries, the combined bc system preserves no SUSY at all. This is reflected in the fact that the scalars in that combined system may become tachyonic when the branes are sufficiently close. Those tachyons signal the recombination of the branes and spontaneous symmetry breaking, as described in section 7. However, since this Higgs system is not supersymmetric, in this case the Q-SUSY property of the rest of the spectrum will not be sufficient to stabilize a hierarchy of masses between the weak scale and the string scale, a (modest) fine-tuning being required to keep apart those two scales. This is one of the main motivations to consider the rombic and triangular quivers which we discuss below. We will describe more aspects of electroweak symmetry breaking for the square quiver in section 7.

4.2 The linear quiver

In some way this may be understood as a variation of the square quiver in which we flip the type of leptonic brane from type a_1 to a_2. Due to this change the left-handed quarks and leptons share now the same SUSY, whereas the right-handed ones respect a different one (see figure 8). One can check that the conditions to get Q-SUSY at the intersections are analogous to those in the square quiver eq. (4.3) and the wrapping numbers giving rise to the SM fermion content are shown in table 7.

One can see that the wrapping numbers are identical to those of the square quiver except for the leptonic brane. This slight change has however an impact in the massless U(1) spectrum. The couplings of the RR fields B_2^i to the U(1)'s is identical to that in the square quiver eq. (4.4) except for a flip in sign of the coefficient of the $B_2^3 \wedge F^d$ term. As a consequence one observes that two U(1)'s will remain necessarily massless: $\mathrm{U}(1)_c$ and $\mathrm{U}(1)_a - 3\,\mathrm{U}(1)_d$. Thus in this model again we will have not only a massless hypercharge $Y = Q_a/6 + Q_c/2 - Q_d/2$ but also an extra $B - L$ generator. One can check that this conclusion does not change even if we add an extra brane as we did in the square quiver case.

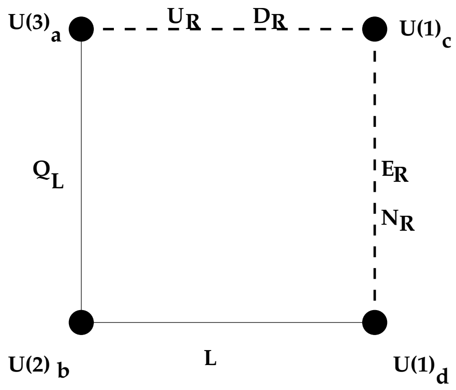

Figure 8: A linear SUSY-quiver with SM spectrum. Note that all left-handed fermions share the same SUSY whereas the right-handed ones respect a different one.

brane type	N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
a_2	$N_a = 3$	$(1, 0)$	(n_a^2, β^2)	$(3, -1/2)$
b_2	$N_b = 2$	$(n_b^1, 1)$	$(1/\beta^2, 0)$	$(1, -1/2)$
b_1	$N_c = 1$	$(0, 1)$	$(1/\beta^2, 0)$	$(0, 1)$
a_2'	$N_d = 1$	$(1, 0)$	$(n_a^2, \beta^2/\rho)$	$(3\rho, -1/2)$

Table 7: D6-brane wrapping numbers giving rise to a Q-SUSY SM spectrum for a linear quiver. Here $\rho = 1, 1/3$, $\beta^2 = 1, 1/2$, $n_a^2, n_b^1 \in \mathbb{Z}$.

Note also that, for $\rho = 1$, the baryonic and leptonic stacks would have the same wrapping numbers and hence one can unify them into a Pati-Salam stack with gauge group $\mathrm{U}(4)_{PS}$ if one locates both stacks on top of each other.

The simple set of branes discussed above are not enough to cancel RR tadpoles in this model. In fact one can check that for the massive anomalous $\mathrm{U}(1)_b$ generator the Green-Schwarz cancellation mechanism will only cancel partially its anomalies and one would in general need to add explicit H_{NS}-flux to cancel tadpoles, as discussed in section 2. This U(1) being anyway massive, decouples from the observable massless spectrum.

Before getting into the rombic class of models let us make a comment concerning the possibility of building left-right symmetric models with gauge group $\mathrm{SU}(3) \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$, i.e., the simplest non-abelian extension of the SM. It is easy to construct such a models starting from the square and linear quivers by just putting a couple of (c)-branes instead of one. We will not present here an analysis of these models, which is straightforward. Let us just make a comment concerning the connection between the property of Q-SUSY and the existence of a gauged $B - L$. In these models, in order to get the desired quantum numbers, we have to ensure that there is a massless $\mathrm{U}(1)_{B-L}$ remaining at the massless level (i.e., no couplings to B_2^i fields) and no other linear combination of U(1)'s. It turns out that in the linear quiver models imposing the Q-SUSY conditions analogous to eq. (4.3) automatically guarantees that there is a $\mathrm{U}(1)_{B-L}$ gauge boson in the massless spectrum.

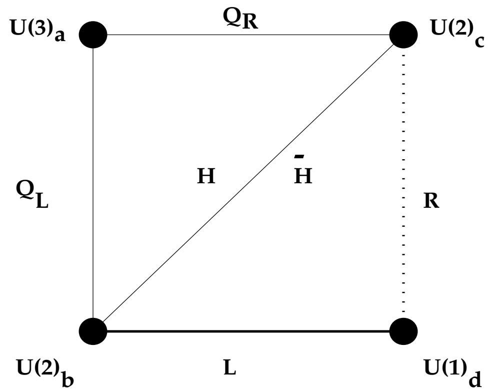

Figure 9: Rombic SUSY-quiver. The quark sector and Higgs fields respect the same $\mathcal{N}=1$ SUSY, while the leptonic sector respects different ones. In the figure a left-right symmetric model has been depicted. Unlike the square case, there are Higgs chiral multiplets at some intersections.

brane type	N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
a_2	$N_a = 3$	$(2, 0)$	$(3, 1/2)$	$(3, -1/2)$
b_2	$N_b = 2$	$(3, -1/2)$	$(2, 0)$	$(3, 1/2)$
c_2	$N_c = 2$	$(3, 1/2)$	$(3, -1/2)$	$(2, 0)$
a_1	$N_d = 1$	$(2, 0)$	$(3, 1/2)$	$(3, 1/2)$

Table 8: D6-brane wrapping numbers giving rise to a Q-SUSY Left-Right symmetric spectrum from rombic quiver.

This connection between Q-SUSY and the presence of $B-L$ is another example of the intriguing connection that we find upon model building in between the presence of Q-SUSY (or SUSY) and the existence of a massless $B-L$ generator.

4.3 The rombic quiver

The general structure of the rombic quiver is presented in figure 9. In order to give an example of such Q-SUSY configuration we are going to present a model with a left-right symmetric $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ gauge group. Getting such gauge group will again imply considering four stacks of branes, following the same general philosophy described in section 2 in order to get the Standard Model. In particular, we will be considering four sets of branes with the wrapping numbers of table 8.

Note that in this brane configuration all the three tori have NS background (i.e., $\beta^i = 1/2$, for all i). From the above configuration one finds for the intersection numbers:

$$\begin{aligned} I_{ab} &= 3\,, & I_{ab*} &= 0\,, \\ I_{ac} &= -3\,, & I_{ac*} &= 0\,, \\ I_{bc} &= 3\,, & I_{bc*} &= 0\,, \\ I_{cd} &= -3\,, & I_{cd*} &= 0\,, \\ I_{bd} &= 0\,, & I_{bd*} &= 3\,, \end{aligned} \tag{4.5}$$

whose associated massless chiral spectrum is presented in table 9.

Intersection	Matter fields		Q_a	Q_b	Q_c	Q_d	B-L
(ab)	Q_L	$3(3,2,2)$	1	-1	0	0	1/3
(ac)	Q_R	$3(\bar{3},1,2)$	-1	0	1	0	-1/3
(bd*)	L	$3(1,2,1)$	0	1	0	1	-1
(cd)	R	$3(1,1,2)$	0	0	-1	1	1
(bc)	H	$3(1,2,2)$	0	1	-1	0	0

Table 9: Spectrum of the Left-Right symmetric model from a rombic quiver.

Unlike the square quiver case, this model has electroweak Higgs fields in the massless chiral spectrum. The property of Q-SUSY is obtained for

$$U^1 = U^2 = U^3 \tag{4.6}$$

in which case each brane respects some $\mathcal{N}=2$ supersymmetry and the branes a, b, c and d are of type a_2, b_2, c_2 and a_1 respectively. Let us analyze the structure of the U(1)'s. In this case one has that the RR fields B_2^i, $i=1,2,3$ couple to the U(1)'s as follows:

$$\begin{aligned} &B_2^1 \wedge 6(-F^b + F^c) \\ &B_2^2 \wedge 3(3F^a - 2F^c + F^d) \\ &B_2^3 \wedge 3(-3F^a + 2F^b + F^d) \end{aligned} \tag{4.7}$$

From here one concludes that there is a unique massless U(1) given by:

$$Q_{B-L} = -\frac{2}{3}Q_a - Q_b - Q_c \tag{4.8}$$

which couples to the fermionic spectrum as the standard $(B-L)$ generator of left-right symmetric models.

An advantage of this model with respect to the square and linear quivers is that the subsector formed by Q_R, Q_L and the electroweak Higgs fields H all respect the same $\mathcal{N}=1$ supersymmetry. Thus the quark subsector of the theory is supersymmetric. Since we know experimentally that the Yukawa couplings of the leptons are in general very small, it is enough to have an approximate SUSY in the quarks sector in order to stabilize the hierarchy between the weak scale and a possible cut-off scale of order 10-100 TeV.

As in the linear model, the set of branes discussed above are not enough to cancel RR tadpoles in this model. In fact one can check that for the massive anomalous $\mathrm{U}(1)_d$ generator the Green-Schwarz cancellation mechanism will only cancel partially its anomalies. As in the linear models, one would in general need to add explicit H_{NS}-flux to cancel tadpoles. Again $\mathrm{U}(1)_d$ is massive and decouples anyway at low energies.

Let us finally mention that the left-right symmetry may eventually break down to the Standard Model by brane recombination, in a way analogous to that explained for the other models in section 7. The reader can verify, following the methods developed in that section, that after the recombination of one of the (c) branes with the (d) brane into a single one, the resulting configuration yields the fermion content and the gauge group of the SM. This would correspond, from the effective action point of view, to giving a vev to the right-handed sneutrinos, what eventually would break $B-L$ into the hypercharge.

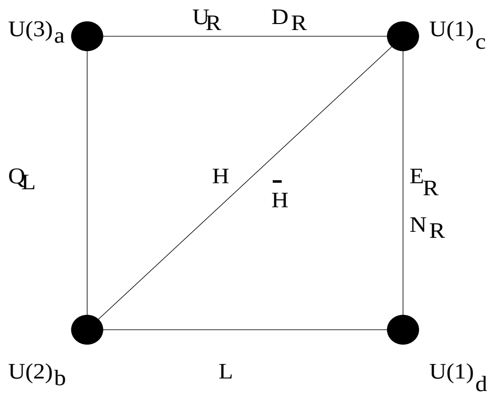

Figure 10: Triangular SUSY-quiver with SM spectrum. The a and d branes are parallel so that actually this may be depicted as a triangle quiver.

4.4 The $\mathcal{N} = 1$ triangle quiver

Let us now show an example based in the triangle quiver in figure 10. Note that, as shown in figure 6, the baryonic (a) and leptonic (d) branes are of the same (a_2) type and could be put on top of each other in the quiver, leading to a triangular shape. In this example all of the intersections will preserve the same $\mathcal{N} = 1$ SUSY so that, at least locally, the models will look fully supersymmetric. As we remarked at the beginning of this section, this quiver predicts the necessary presence of Higgs multiplets appearing at the intersections of *Left* and *Right* brane stacks. Imposing the that left- and right-handed fermion multiplets respect the same $\mathcal{N} = 1$ SUSY forces the existence of these Higgs multiplets.

Consider the wrapping numbers in table 10. It is easy to check that if the conditions

$$2U^1 = 2\beta^2 U^2 = U^3 \tag{4.9}$$

are met, the brane configuration indeed respects the same supersymmetry at all intersections. Here again $U^i = R_2^i/R_1^i$, where R_1^i,R_2^i, $i = 1, 2, 3$ are the radii of the tori. From the brane wrapping numbers above one can obtain the intersection numbers

$$\begin{aligned} I_{ab} &= 1\,, & I_{ab*} &= 2\,, \\ I_{ac} &= -3\,, & I_{ac*} &= -3\,, \\ I_{bd} &= -1\,, & I_{bd*} &= 2\,, \\ I_{cd} &= 3\,, & I_{cd*} &= -3\,, \\ I_{bc} &= \frac{-1}{\beta^2}\,, & I_{bc*} &= \frac{-1}{\beta^2}\,, \end{aligned} \tag{4.10}$$

which gives rise to the massless spectrum in table 11.

Let us now study the structure of the couplings of the U(1) fields to the antisymmetric B-fields, which determine which of them become massive. One finds the couplings:

$$\begin{aligned} &B_2^1 \wedge \frac{2}{\beta^2} F^b \\ &B_2^2 \wedge 3\beta^2 (3F^a + F^d) \\ &B_2^3 \wedge -\frac{1}{2}\left(9F^a + \frac{2}{\beta^2} F^b + 3F^d\right). \end{aligned} \tag{4.11}$$

brane type	N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
a_2	$N_a = 3$	$(1,0)$	$(3,\beta^2)$	$(3,-1/2)$
b_2	$N_b = 2$	$(1,1)$	$(1/\beta^2,0)$	$(1,-1/2)$
c_2	$N_c = 1$	$(0,1)$	$(0,-1)$	$(2,0)$
a_2'	$N_d = 1$	$(1,0)$	$(3,\beta^2)$	$(3,-1/2)$

Table 10: Wrapping numbers of a three generation SUSY-SM with $\mathcal{N} = 1$ SUSY locally.

Intersection	Matter fields		Q_a	Q_b	Q_c	Q_d	Q_Y
ab	Q_L	$(3,2)$	1	-1	0	0	1/6
ab^*	q_L	$2(3,2)$	1	1	0	0	1/6
ac	U_R	$3(\bar{3},1)$	-1	0	1	0	$-2/3$
ac^*	D_R	$3(\bar{3},1)$	-1	0	-1	0	1/3
bd	L	$(1,2)$	0	-1	0	1	$-1/2$
bd^*	l	$2(1,2)$	0	1	0	1	$-1/2$
cd	N_R	$3(1,1)$	0	0	1	-1	0
cd^*	E_R	$3(1,1)$	0	0	-1	-1	1
bc	H	$\frac{1}{\beta^2}(1,2)$	0	-1	1	0	$-1/2$
bc^*	$\bar{H}$	$\frac{1}{\beta^2}(1,2)$	0	-1	-1	0	1/2

Table 11: Chiral spectrum of the SUSY's SM obtained from the triangular quiver. The hypercharge generator is defined as $Q_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d$.

One concludes from here that both $\mathrm{U}(1)_b$ and $3\,\mathrm{U}(1)_a + \mathrm{U}(1)_d$ generators become massive. The massless anomaly-free generators are thus $\mathrm{U}(1)_c$ and $\mathrm{U}(1)_a - 3\,\mathrm{U}(1)_d$, which correspond with the third component of right-handed isospin and B-L, respectively. The hypercharge is given by the linear combination of generators:

$$\mathrm{U}(1)_Y = \frac{\mathrm{U}(1)_a}{6} - \frac{\mathrm{U}(1)_c}{2} - \frac{\mathrm{U}(1)_d}{2}\,. \tag{4.12}$$

Thus, as expected, the massless hypercharge comes along with a massless B-L generator, if supersymmetric intersections are imposed. Note also that, as in the linear quiver case, if one locates the leptonic stack of branes d on top of the baryonic stack one would get an enlarged $\mathrm{U}(4)_{PS}$ Pati-Salam gauge symmetry.

Note that there are two versions of the model with $\beta^2 = 1, 1/2$. Interestingly enough, for the $\beta^2 = 1$ case *the massless chiral spectrum is exactly the same of the MSSM*. For $\beta^2 = 1/2$ the Higgs sector is doubled, although the rest of the spectrum remains the same. As in the rombic example, the simple brane configuration by itself would give rise to RR tadpoles. Those may be canceled by the addition of either some (non-factorizable) brane system with no intersection with the SM branes (for the $\beta^2 = 1/2$ case) or some explicit H_{NS}-flux (for the $\beta^2 = 1$ case).

Before ending this model-building section let us make a final comment. The square quiver models have an attractive point which the others do not have. Square quiver models are just a subclass of those presented in [12], and the latter have the attractive feature

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that the number of generations is related to the number of colours by cancellation of $U(1)_b$ anomalies. In order for this to work there cannot be extra fermion doublets, like the Higgsinos appearing in the triangle and rombic models, which contribute to $U(1)_b$ anomalies. In the case of the linear quiver models one can check that the $U(1)_b$ anomalies do not cancel and Wess-Zumino terms (induced by the presence of H_{NS}-flux) must appear (in addition to the Green-Schwarz mechanism terms) to complete $U(1)_b$ anomaly cancellation. Thus also in this case the number of generations/colours argument is not present. On the other hand, as we mentioned above, the triangle and rombic quivers predict the generic presence of massless Higgs multiplets to give rise to electroweak symmetry breaking.

5. SUSY-breaking and Fayet-Iliopoulos terms

The Q-SUSY examples discussed in the previous section may present three different sources of supersymmetry breaking:

1. Soft SUSY-breaking masses from loop graphs, as discussed in section 3. These will in general be present in the four types of models considered, since all of them will have non-SUSY massive sectors with masses generically of order M_s. They will induce generically one-loop gaugino masses as well as two-loop scalar masses[2]. These effects will be of order $\alpha/(4\pi)M_s$, M_s being the string scale.

2. Fayet-Iliopoulos terms. These may be present in any of the models and have a nice geometrical interpretation, as discussed in [14, 15]. Indeed, in order to obtain SUSY at the brane intersections, we had to appropriately tune the U^i complex structure parameters (see eqs. (4.3), (4.6), (4.9)). A slight departure from these adjustments is seen in the effective lagrangian as the turning on of FI-terms for the U(1) fields of the models. This source of SUSY-breaking may be large or small, depending on the values given to the U^i.

3. Some models have explicit SUSY-breaking in some sector. In particular we already mentioned that in the square and linear quivers the Higgs sector is non-SUSY and so will be the Yukawa couplings. In the case of rombic models the Higgs, and quark sectors respect SUSY at tree level but there will be explicit SUSY-breaking in the leptonic sector (non-SUSY Yukawa couplings). This is expected because, although all intersections preserve some SUSY, Yukawa couplings involve in general different intersections respecting different supersymmetries. As we mentioned, since the leptonic Yukawas are known to be small, the Higgs mass will be sufficiently protected from loop corrections so as to avoid fine-tuning. Finally, in the triangle quiver case all intersections respect the same SUSY and there is a supersymmetric superpotential.

The first source of SUSY-breaking will depend on the particular massive non-SUSY sector of the given model. On the other hand following ref. [15] we can give a relatively model-independent discussion of the FI-terms in this class of theories. Consider a D6-brane which is forming angles with the orientifold plane given by $(\vartheta^1, \vartheta^2, \vartheta^1 + \vartheta^2 + \delta, 0)$, $\vartheta^i > 0$.

Then, for $\delta = 0$ one would recover one unbroken supersymmetry. For small non-vanishing δ one can approximate the effect of this by turning on a Fayet-Iliopoulos term for the U(1) field associated to the brane and given by

$$\xi_a \approx -\frac{\delta}{2} M_s^2 \,. \tag{5.1}$$

In order to compute the mass of a scalar living at the intersection of two branes, let us take two D6-branes $D6_a$ and $D6_b$ whose separation from the same SUSY limit is given by δ_a and δ_b, respectively. Then the mass of this scalar is given by

$$m_{ab}^2 = (-q_a \xi_a - q_b \xi_b) \approx \frac{1}{2}(\delta_a - \delta_b) M_s^2 \,. \tag{5.2}$$

and in fact one can check that this result is in agreement with the masses for the scalars lying at a intersection obtained from the string mass formulae (see ref. [15] for details). Thus for small deviations from a SUSY configuration, the masses of the scalars at an intersection may be understood as coming from Fayet-Iliopoulos terms.

Let us emphasize some points to try to avoid confusion. Normally, when one talks about a FI-term in a SUSY field theory, one is taking about a massless U(1). On the other hand in our models with exact Q-SUSY it is the massive U(1)'s (masses of order M_s) that get FI-terms. This may be understood because a FI-term is the SUSY partner of a $B \wedge F$ coupling [26], and those U(1)'s with a non-vanishing coupling to some B-field are the ones which become massive. On the other hand these massive U(1)'s decouple well below the string scale. This means that, e.g., there are no quartic scalar couplings proportional to the gauge coupling constants of massive U(1)'s. Thus from the effective low-energy point of view the masses of scalars coming from FI-terms of these massive U(1)'s will look rather like standard soft SUSY-breaking masses. It is about these scalar masses that we will be talking in this section.

5.1 FI-terms and sfermion masses in the square and linear quivers

Let us now consider the turning of FI-terms in the square quiver. In fact we are going to consider a slight generalization of the square quiver which has the relative advantage that the only unbroken U(1) is hypercharge.

Indeed, in order to get exact SUSY at all intersections we imposed $n_c^1 = 0$. This was required so that the brane c makes a $\frac{\pi}{2}$ angle with the O-plane in tori 1 and 3, preserving in this way two supersymmetries with the orientifold and one supersymmetry in each intersection with branes a and d. This had as a consequence that the single U(1) massless (in the absence of extra branes) was Q_c, rather than hypercharge. On the other hand, if we are going to break SUSY anyway by adding FI-terms, we can also relax the Q-SUSY condition and allow for approximate SUSY at the intersections involving the c-brane.

Let us start from the model of table 3 and let us impose the conditions

$$\begin{aligned} n_b^1 &> 0 \,, & n_a^2 &= 3\rho n_d^2 > 0 \,, \\ U^1 &= \frac{3\rho n_b^1}{2\beta^1} U^3 \,, & U^2 &= \frac{n_a^2 \rho}{2\beta^2} U^3 \,. \end{aligned} \tag{5.3}$$

N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 3$	$(1/\beta^1, 0)$	(n_a^2, β^2)	$(3, -1/2)$
$N_b = 2$	(n_b^1, β^1)	$(1/\beta^2, 0)$	$(1, -1/2)$
$N_c = 1$	$\left(-\frac{\beta^2}{\beta^1} n_a^2, \beta^1\right)$	$(1/\beta^2, 0)$	$(0, 1)$
$N_d = 1$	$(1/\beta^1, 0)$	$(n_a^2, 3\beta^2)$	$(1, 1/2)$

Table 12: D6-brane wrapping numbers giving rise to an approximate square Q-SUSY Standard Model.

Brane	θ_α^1	θ_α^2	θ_α^3
a	0	$\alpha_1 + \delta_a$	$-\alpha_1$
b	$\alpha_2 + \delta_b$	0	$-\alpha_2$
c	$\frac{\pi}{2} + \Delta$	0	$\frac{\pi}{2}$
d	0	$\alpha_2 + \delta_d$	α_2

Table 13: Angles that the D6-brane stacks form with the orientifold axis in the approximate square configuration in the text.

Note that we have not imposed the condition $n_c^1 = 0$, and thus we do not need to impose any condition on β^1, ρ. The conditions (5.3) will ensure SUSY for all intersections not involving brane c. Now, the condition to get a massless hypercharge generator is:

$$n_c^1 = -\frac{\beta^2}{2\beta^1}(n_a^2 + 3\rho n_d^2) = -\frac{\beta^2}{\beta^1} 3\rho n_d^2 \,, \tag{5.4}$$

since we have chosen $\epsilon = \tilde{\epsilon} = 1$. This implies that in order to avoid multiple wrappings [12] we also need to set $\rho = 1/3$ (and hence $n_a^2 = n_d^2$). We present in table 12 the wrapping numbers of these models.

Consider now a slight departure of the moduli $U^{1,2}$ from their assigned values in (5.3):

$$\begin{aligned} U^1 &= \frac{n_b^1}{2\beta^1} U^3 + \delta^1 \\ U^2 &= \frac{n_a^2}{6\beta^2} U^3 + \delta^2 \end{aligned} \tag{5.5}$$

and let us also define

$$\Delta \equiv \frac{\pi}{2} - \tan^{-1}\left(\frac{(\beta^1)^2}{\beta^2 n_d^2} U^1\right) = \frac{\pi}{2} - \tan^{-1}\left[\frac{(\beta^1)^2}{\beta^2 n_d^2}\left(\frac{n_b^1}{2\beta^1} U^3 + \delta^1\right)\right]. \tag{5.6}$$

This parameter Δ will measure the departure from supersymmetry of the brane c. Under these conditions the angles formed by the branes with the orientifold plane are the ones shown in table 13, where $\alpha_1 = \tan^{-1}(U^3/6)$ and $\alpha_2 = \tan^{-1}(U^3/2)$.

There one has (for small δ^i)

$$\delta_a = \frac{\frac{\beta^2}{3n_a^2}\delta^2}{1 + (U^3/6)^2} \,;$$

Sector	$(\theta^1_{\alpha\beta}, \theta^2_{\alpha\beta}, \theta^3_{\alpha\beta})$	sparticle	mass^2
(ab)	$(\alpha_2 + \delta_b, -\alpha_1 - \delta_a, -\alpha_2 + \alpha_1)$	$1 \times \tilde{Q}_L$	$\frac{1}{2}(\delta_a - \delta_b)$
(ab*)	$(-\alpha_2 - \delta_b, -\alpha_1 - \delta_a, \alpha_2 + \alpha_1)$	$2 \times \tilde{q}_L$	$\frac{1}{2}(\delta_a + \delta_b)$
(ac)	$(\frac{\pi}{2} + \Delta, -\alpha_1 - \delta_a, \frac{\pi}{2} + \alpha_1)$	$3 \times \tilde{U}_R$	$\frac{1}{2}(\delta_a + \Delta)$
(ac*)	$(-\frac{\pi}{2} - \Delta, -\alpha_1 - \delta_a, -\frac{\pi}{2} + \alpha_1)$	$3 \times \tilde{D}_R$	$\frac{1}{2}(\delta_a - \Delta)$
(bd)	$(-\alpha_2 - \delta_b, \alpha_2 + \delta_d, 2\alpha_2)$	$3 \times \tilde{L}$	$\frac{1}{2}(\delta_b + \delta_d)$
(cd)	$(-\frac{\pi}{2} - \Delta, \alpha_2 + \delta_d, \alpha_2 - \frac{\pi}{2})$	$3 \times \tilde{N}_R$	$\frac{1}{2}(\delta_d - \Delta)$
(cd*)	$(-\frac{\pi}{2} - \Delta, -\alpha_2 - \delta_d, -\alpha_2 - \frac{\pi}{2})$	$3 \times \tilde{E}_R$	$\frac{1}{2}(\delta_d + \Delta)$

Table 14: Squark and slepton masses from FI-terms in the square quiver model.

$$\delta_b = \frac{\frac{1}{n_b^1}\delta^1}{1 + (U^3/2)^2}\,;$$
$$\delta_d = \frac{\frac{\beta^2}{n_a^2}\delta^2}{1 + (U^3/2)^2}$$

Note also that (again for $\delta^1 \ll 1$):

$$\tan\Delta \simeq \frac{2n_a^2\beta^2}{n_b^1\beta^1}\frac{1}{U^3}\left(1 - \frac{2\beta^1}{n_b^1}\frac{\delta^1}{U^3}\right) \tag{5.7}$$

so that one can check that Δ can be made quite small for not too large U_3. From table 13 one sees that for all practical purposes the effect of not setting $n_c^1 = 0$ induces SUSY-breaking which (for small Δ) can be parametrized as a FI-term for the $\mathrm{U}(1)_c$ generator.

We can now compute the masses of the sparticles in terms of the FI-terms following the discussion given in ref. [15]. The results are shown in table 14. Notice that we could formally put the masses depending on Δ as coming from a Fayet-Iliopoulos term. However, and unlike the other cases, supersymmetry in these intersections would only be recovered in the $U^3 \to \infty$ limit.

An important point to note is that there is a wide range of parameters for which all squark and slepton (mass)^2 are positive and hence there are no unwanted charge and colour-breaking minima. It is enough to satisfy $\delta_a, \delta_d > \Delta$ and $\delta_a > |\delta_b|$ in order to have all (mass)^2 positive. Note also that (for a given choice of discrete parameters) the squark and slepton masses come determined from just three independent parameters δ_1, δ_2 and U^3. In some limits the expressions for these masses becomes particularly simple. For example, if $\delta_1 = 0$ one gets

$$m^2_{\tilde{Q}_L} = \frac{\delta_a}{2}M_s^2\,; \qquad m^2_{\tilde{L}} = \frac{\delta_d}{2}M_s^2$$
$$m^2_{\tilde{U}_R,\tilde{D}_R} = \left(\frac{\delta_a}{2} \pm \frac{\Delta}{2}\right)M_s^2\,; \qquad m^2_{\tilde{E}_R,\tilde{N}_R} = \left(\frac{\delta_d}{2} \pm \frac{\Delta}{2}\right)M_s^2 \tag{5.8}$$

Note that in this case the ratio of *average* squark masses versus slepton masses is controlled only by the value of U^3:

$$\frac{m^2_{\text{squark}}}{m^2_{\text{slepton}}} = \frac{\delta_a}{\delta_d} = \frac{1}{3}\frac{1+(U^3/2)^2}{1+(U^3/6)^2} = \frac{g_a^2}{g_d^2} \tag{5.9}$$

The last equality follows from eqs. (6.1)and (6.6) in the next subsection. In this connection note that using those equations we can rewrite eqs. (5.7) as

$$\delta_a = g_a^2(\beta^2\delta^2)\,; \qquad \delta_b = g_b^2\left(\frac{\delta^1}{\beta^2}\right); \qquad \delta_d = g_d^2(\beta^2\delta^2) \tag{5.10}$$

Thus the (masses)^2 of sfermions from FI-terms are proportional to the square of the U(1) coupling constants, very much like if they were coming from a one-loop effect.[9] In any event, it is interesting how one can obtain the masses of squarks and sleptons in this scheme in terms of a few geometrical parameters. Notice however that one should add to these masses the contributions coming from the loops involving massive non-SUSY particles, as described above.

One interesting point concerns the right-handed sneutrino. From the above mass structure one can check that there is a region of parameter space for which one might have a tachyonic mass for the $\tilde{\nu}_R$ but a positive one for the rest of the scalars. It is enough to have $\delta_a > \Delta$ but $\delta_d < \Delta$. One can check that this is possible provided $U^3 > \sqrt{12}$. If this happens lepton number is broken and the $B-L$ gauge boson becomes massive swallowing a linear combination of $\tilde{\nu}_R$'s and antisymmetric B-fields. From the brane point of view the c and d^* branes will recombine into a single one and the obtained structure is similar to the models discussed in subsection 7.4 .

Another point is in order concerning FI-terms. We are used to the fact that in *supersymmetric* theories the masses that scalars get from a FI-term are proportional to the U(1) charges of their fermionic partners. Note that in the case of the Q-SUSY theories that is in general not the case, as observed in ref. [15]. Indeed, consider for example the FI-term associated to the baryonic symmetry $\mathrm{U}(1)_a$. One can see from table 14 that both left- and right-handed squarks get positive masses, although their fermionic partners have opposite baryon number. This is a reflection of the fact that different $\mathcal{N}=1$ SUSY's are preserved at the ab and at the ac intersections.

One can repeat an analogous discussion for the linear quiver case. It is easy to check that one obtains the same results for the masses coming from FI-terms than the ones in the square quiver, table 14. The only difference is that one necessarily has $\Delta = 0$ in the linear case. Indeed, if $\Delta \neq 0$ is taken, then the Q_c generator is massive and the unbroken generator would be $B-L$, rather than hypercharge. This is different to what happens in the square quiver case in which one can chose parameters so that it is hypercharge which remains massless.

[9] In fact in some sense they do, since tree-level closed string couplings are one-loop from the open string channel point of view.

5.2 FI-terms and sfermion masses in the triangle and rombic quivers

One can also easily compute the masses of the scalars coming from turning on FI-terms in the other triangle and rombic type of models. Note that in these cases there are several (or all) intersections respecting the *same* $\mathcal{N}=1$ supersymmetry. In these $\mathcal{N}=1$ subsectors the usual fact will hold and the masses of scalars from FI-terms will be proportional to the U(1) fermionic charges. Consider for example the triangular case. By going slightly away from the supersymmetric configuration

$$U^1=\frac{1}{2}U^3+\delta_1\,;\qquad U^2=\frac{1}{2\beta^2}U^3+\delta_2 \tag{5.11}$$

it is easy to check that one induces FI terms for the U(1)'s as follows:

$$\xi_a=\xi_d=-\frac{\delta_2\beta^2/6}{1+(U^3/6)^2}M_s^2\,;\qquad \xi_b=-\frac{\delta_1/2}{1+(U^3/2)^2}M_s^2 \tag{5.12}$$

and $\xi_c=0$ (recall that $\mathrm{U}(1)_c$ is massless). Note that there are only two independent parameters for the FI-terms, which corresponds to the existence of two massive (anomalous) U(1)'s, $\mathrm{U}(1)_b$ and $3\,\mathrm{U}(1)_a+\mathrm{U}(1)_d$ (more physically $(9B+L)$). Now, unlike the case of the square quiver, the masses for the scalars at the intersections will be just proportional to these charges. Looking at the particle spectrum in the triangle models (table 11) we see that left- and right-handed squarks will have *opposite* masses from the $\mathrm{U}(1)_a$ baryonic FI-term. Thus, FI-terms are unable to give positive masses for all squarks and sleptons at the same time. There is nothing wrong with that, only that the FI-terms contribution cannot be the dominant source of scalar masses if we want to avoid unwanted charge/colour breaking minima, the one-loop contributions should also be important.

A similar general conclusion holds for the rombic model, since the quiver contains a triangle with three intersections respecting the same SUSY.

6. Gauge coupling constants

Unlike the heterotic case, the gauge coupling constants of brane models in which each gauge factor lives in a different stack of branes have no unification. Rather, the size of each coupling constant square is inversely proportional to the volume wrapped by the corresponding brane. The physical gauge couplings will depend on

i). The size of the gauge couplings at the string scale.

ii). The running between the string scale M_s and the weak scale.

For a string scale of order 10-100 TeV there is running for two to three orders of magnitude. The effect of this running may be enhanced due to the fact that in between the string scale, and the weak scale there is not only the spectrum of the SUSY SM but also of the adjoint scalars and fermions corresponding to the $\mathcal{N}=4$ structure of the gauge sectors. Those fields are expected to get masses of order $\alpha/(4\pi)M_s$, which may be well below the string scale. In addition close to the string scale there can be important threshold corrections due to the excitation of KK , winding and "gonion" states [10] some of which may be below the string scale. All these running effects are difficult to evaluate in a model independent way. However one can obtain easy and closed formulae for the size of the gauge couplings at the string scale.

The formulae for the size of gauge couplings get simplified in the case of Q-SUSY and SUSY models compared to the general non-SUSY intersection models. As noted in e.g. [10], the tree-level value of the different gauge coupling constants at the string scale are controlled by the length of the wrapping cycles, i.e.,

$$\frac{4\pi^2}{g_i^2} = \frac{M_s^3}{(2\pi)^2\lambda_{II}}||l_i|| \tag{6.1}$$

where M_s is the string scale, λ_{II} is the type II string coupling, and $||l_i||$ is the length of the cycle of the i-th set of branes

$$\frac{||l_i||^2}{(2\pi)^6} = \left(\left(n_i^1 R_1^{(1)}\right)^2 + \left(m_i^1 R_2^{(1)}\right)^2\right)\left(\left(n_i^2 R_1^{(2)}\right)^2 + \left(m_i^2 R_2^{(2)}\right)^2\right)\left(\left(n_i^3 R_1^{(3)}\right)^2 + \left(m_i^3 R_2^{(3)}\right)^2\right). \tag{6.2}$$

The ratio of the SU(3) and SU$(2)_L$ couplings is thus given by:

$$\frac{\alpha_{QCD}}{\alpha_L} = \frac{||l_b||}{||l_a||} \tag{6.3}$$

For models without a non-abelian extended gauge structure (like the square and triangle quiver examples above) the hypercharge is a linear combination of U$(1)_a$, U$(1)_c$ and U$(1)_d$. One has for the hypercharge coupling:

$$\alpha_Y^{-1} = \frac{1}{6^2}\alpha_a^{-1} + \frac{1}{2^2}\left(\alpha_c^{-1} + \alpha_d^{-1}\right) \tag{6.4}$$

where α_a is the coupling of the U$(1)_a$, which in our normalization verifies at the string scale $\alpha_a = \alpha_{QCD}/6$. We thus have:

$$\begin{aligned} \frac{\alpha_{QCD}}{\alpha_Y} &= \frac{1}{6} + \frac{1}{4}\left(\frac{\alpha_c^{-1}}{\alpha_{QCD}^{-1}} + \frac{\alpha_d^{-1}}{\alpha_{QCD}^{-1}}\right) \\ &= \frac{1}{6} + \frac{1}{4}\left(\frac{||l_c||}{||l_a||} + \frac{||l_d||}{||l_a||}\right) \end{aligned} \tag{6.5}$$

We can now evaluate these ratios for the different models. It turns out that for Q-SUSY models the complicated non-linear expression for the volume of the cycles (6.2) is substantially simplified and the dependence on moduli becomes linear [15].

1. Square and linear quiver.

   In this case we have for the wrapped volume of each brane

$$||l_a|| = 3n_a^2 R_1^{(1)} R_1^{(2)} R_1^{(3)} \left(1 + \left(\frac{U^3}{6}\right)^2\right)$$

$$
\begin{aligned}
||l_b|| &= \frac{n_b^1}{\beta^2} R_1^{(1)} R_1^{(2)} R_1^{(3)} \left(1 + \left(\frac{U^3}{2}\right)^2\right) \\
||l_c|| &= \frac{n_b^1}{2\beta^2} R_1^{(1)} R_1^{(2)} R_1^{(3)} (U^3)^2 \\
||l_d|| &= 3\rho n_a^2 R_1^{(1)} R_1^{(2)} R_1^{(3)} \left(1 + \left(\frac{U^3}{2}\right)^2\right)
\end{aligned}
\tag{6.6}
$$

Here we have $\rho = 1/3$ in the square quiver and both choices $\rho = 1, 1/3$ in the linear quiver case. With these values for the volumes of the cycles one has for the ratio of SU(3) and $\mathrm{SU}(2)_L$ couplings:

$$
\frac{\alpha_{QCD}}{\alpha_L} = \frac{||l_b||}{||l_a||} = \frac{n_b^1}{3n_a^2\beta^2} \frac{1 + (U^3/2)^2}{1 + (U^3/6)^2} \tag{6.7}
$$

and for the hypercharge coupling

$$
\frac{\alpha_{QCD}}{\alpha_Y} = \frac{1}{6} + \frac{1}{12} \left(\frac{3\rho + (3\rho/4 + \frac{n_b^1}{2n_a^2\beta^2})(U^3)^2}{1 + (U^3/6)^2} \right) \tag{6.8}
$$

2. Triangle quiver.

In this case we have for the wrapped volume of each brane

$$
\begin{aligned}
||l_a|| &= 9 R_1^{(1)} R_1^{(2)} R_1^{(3)} \left(1 + \left(\frac{U^3}{6}\right)^2\right) \\
||l_b|| &= \frac{1}{\beta^2} R_1^{(1)} R_1^{(2)} R_1^{(3)} \left(1 + \left(\frac{U^3}{2}\right)^2\right) \\
||l_c|| &= \frac{1}{2\beta^2} R_1^{(1)} R_1^{(2)} R_1^{(3)} (U^3)^2 \\
||l_d|| &= ||l_a||
\end{aligned}
\tag{6.9}
$$

so that we find for the ratio of strong and weak couplings

$$
\frac{\alpha_{QCD}}{\alpha_L} = \frac{||l_b||}{||l_a||} = \frac{1}{9\beta^2} \frac{1 + (U^3/2)^2}{1 + (U^3/6)^2} \tag{6.10}
$$

and for the hypercharge coupling

$$
\frac{\alpha_{QCD}}{\alpha_Y} = \frac{1}{6} + \frac{1}{4} \left(1 + \frac{1}{18\beta^2} \frac{(U^3)^2}{1 + (U^3/6)^2}\right) \tag{6.11}
$$

As we said, this gives us the ratios of the coupling constants at the string scale, and one should compute the running of the couplings in the region in between the weak scale and the string scale. There can be important effects if there are extra particles in that energy region from e.g., KK or winding states, gonions etc. One can in fact check that if one neglects the effect of the running, it is not possible to find values for

U^3 such that one reproduces at the string scale the experimental weak scale ratios $\alpha_{QCD}^{-1} : \alpha_L^{-1} : \alpha_Y^{-1} = 8.3 : 29.6 : 98.4$. In the case of the square quiver it is easy to obtain the weak scale result for $\frac{\alpha_{QCD}}{\alpha_L}$ but the ratio $\frac{\alpha_{QCD}}{\alpha_Y}$ turns out to be too small. In the case of the triangle model presented above it is also difficult to obtain $\frac{\alpha_{QCD}}{\alpha_L}$ large enough. This is due to the fact that this particular example has very large winding numbers n_a^i for the baryonic brane, which then tends to give too small α_{QCD}. It should be interesting to do a systematic search for other triangle models giving rise to larger values for α_{QCD}.

Note that these statements concern these particular models with just 4 stacks of branes discussed above. If for example there is an extra brane h contributing to the hypercharge group as in eq. (4.4) there is an extra piece $\frac{1}{4}\frac{||l_h||}{||l_a||}$ to be added to eq. (6.5) and in the square quiver case one can then easily adjust $\frac{\alpha_{QCD}}{\alpha_Y}$ by varying the extra brane parameters.

7. The SM Higgs mechanism as a brane recombination process

All the above constructions involve the unbroken electroweak gauge symmetry $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ which we know is spontaneously broken by the Higgs mechanism. The different models have scalars in their spectrum with the quantum numbers of SM Higgs fields to do the job. Whereas the mechanism is understood at the field theory level, there should exist a string version of the SM Higgs mechanism in terms of the underlying branes. Specifically, from the string theory point of view we know that the gauge group of the SM originates from open strings starting and ending on D-branes. The gauge group $\mathrm{SU}(2)_L$ originates on two parallel branes whereas hypercharge comes from a linear combination of U(1)'s attached to different branes. In the SM Higgs mechanism the rank of the gauge group is reduced. As we will now discuss, the stringy counterpart of this rank-reduction is brane recombination. Brane recombination is a process in which two intersecting branes fuse into a single one (see figure 11). In particular, it is known that tachyons appearing in string theory signal an instability with respect to the decay of the system into another one with lower energy. Tachyons appearing at a pair of interesting branes show the instability with respect to the recombination of both branes into a single one with less energy, i.e., less wrapped volume and lying in the same homology class than the initial pair.

In a brane recombination process the number of massless chiral fermions decreases. Consider a couple of intersecting branes α and β which recombine into a single final brane γ. Consider an spectator brane ρ which intersects both branes α and β (i.e., $I_{\rho\alpha} \neq 0$, $I_{\rho\beta} \neq 0$) but not participating in the recombination. The net number of chiral fermions before recombination n_i

$$n_i = |I_{\rho\alpha}| + |I_{\rho\beta}| \geq |I_{\rho\alpha} + I_{\rho\beta}| = |I_{\rho\gamma}| = n_f \tag{7.1}$$

is bigger than the net number of chiral fermions after recombination n_f. The effective field theory interpretation is that after a Higgs mechanism some of the chiral fermions may acquire masses from (not-necessarily renormalizable) Yukawa couplings.

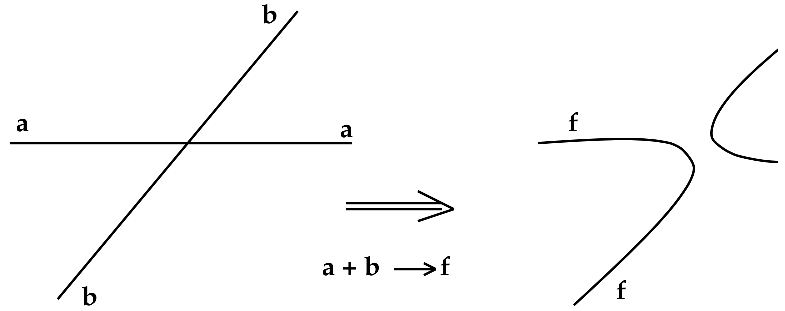

Figure 11: Recombination of two intersecting branes a, b in a compact space. Initially the gauge group is $\mathrm{U}(1)_a \times \mathrm{U}(1)_b$. At the intersection a tachyon scalar triggers the recombination into a single brane f. The final gauge symmetry is $\mathrm{U}(1)_f$, corresponding to the Higgsing $\mathrm{U}(1)_a \times \mathrm{U}(1)_b \to \mathrm{U}(1)_f$ which is induced by the tachyon.

In what follows we will show how general arguments about brane recombination as above yield results which are in agreement with the low-energy field theory expectations. However we will also see that brane recombination has some extra stringy ingredients: it cannot just be described by the lightest tachyon getting a vev, higher excitations of the tachyon must be involved in the full process of brane recombination. This seems to indicate that a full description of the process would require string field theory, rather than an effective field theory lagrangian. Our argumentation on recombination will only make use of the computation of intersection numbers before and after brane recombination, imposing conservation of RR charge. Thus it will be mostly topological rather than geometrical in nature, i.e., we will be able to say what fermions and gauge bosons become massive in a brane recombination process, but not the size of the masses acquired.

The understanding of the Higgs mechanism as a brane recombination process has nothing to do with the presence or not of some SUSY at the brane intersections. So our discussion below also applies to general non-SUSY intersecting models. However for concreteness we will first discuss how the Higgs/recombination process occurs in the square quiver models and will later describe the case of the triangle quiver in which an interesting departure between brane recombination and effective field theory Higgsing occurs.

7.1 The electroweak Higgs system in the square quiver

We will consider for simplicity the version of the square quiver discussed in section 5.1 in which only the hypercharge remains massless, although the discussion in this subsection applies in fact to the larger class of models discussed in [12]. In the square quiver there are not generic intersections giving rise to massless scalars with the quantum numbers of the SM Higgs fields. On the other hand open strings stretched between the 'b' and 'c' stacks of branes do have the quantum numbers of Higgs fields appropriate to yield electroweak symmetry-breaking. These two stacks of branes are parallel on the second torus, and that is why they generically do not intersection However if one approaches those stacks to each other scalar Higgsses with the quantum numbers of table 15 appear

Higgs	Q_b	Q_c	Y
h_1	1	-1	$1/2$
h_2	-1	1	$-1/2$
H_1	-1	-1	$1/2$
H_2	1	1	$-1/2$

Table 15: Electroweak Higgs fields.

in the light spectrum [12]. These states corresponding to open strings stretching between branes b and c (denoted by $h^{\pm}$) and between branes b and c^* (denoted by $H^{\pm}$) have masses (see [12])

$$\begin{aligned} m_{H^{\pm}}{}^2 &= \frac{X_{bc^*}^2}{4\pi} M_s^2 \pm \frac{M_s^2}{2} \left||\theta_{bc^*}^1| - |\theta_{bc^*}^3|\right| = M_s^2 \left(\frac{X_{bc^*}^2}{4\pi} \pm \frac{1}{2}(2\alpha_2 - \Delta + \delta_b) \right), \\ m_{h^{\pm}}{}^2 &= \frac{X_{bc}^2}{4\pi} M_s^2 \pm \frac{M_s^2}{2} \left||\theta_{bc}^1| - |\theta_{bc}^3|\right| = M_s^2 \left(\frac{X_{bc}^2}{4\pi} \pm \frac{1}{2}(2\alpha_2 + \Delta + \delta_b) \right), \end{aligned} \tag{7.2}$$

where X_{bc^*} (X_{bc}) is the distance (in $\alpha'^{\frac{1}{2}}$ units) in transverse space along the second torus. One also has $\alpha_2 = tg^{-1}(U^3/2)$ and Δ, δ_b were defined in (5.6), (5.7). There are also fermionic partners ("Higgsinos") of these two types of complex scalar fields. The above scalar mass spectrum can be interpreted as arising from a field theory mass matrix

$$(H_1^* H_2) \left(\mathbf{M^2}\right) \begin{pmatrix} H_1 \\ H_2^* \end{pmatrix} + (h_1^* \ h_2) \left(\mathbf{m^2}\right) \begin{pmatrix} h_1 \\ h_2^* \end{pmatrix} + \text{h.c.} \tag{7.3}$$

where

$$\mathbf{M^2} = M_s^2 \begin{pmatrix} X_2^{(bc^*)} & \frac{1}{2}(2\alpha_2 - \Delta + \delta_b) \\ \frac{1}{2}(2\alpha_2 - \Delta + \delta_b) & X_2^{(bc^*)} \end{pmatrix}, \qquad \mathbf{m^2} = (bc^*, -\Delta \leftrightarrow bc, \Delta) \tag{7.4}$$

The fields H_i and h_i are thus defined as

$$H^{\pm} = \frac{1}{2}(H_1^* \pm H_2)\,; \qquad h^{\pm} = \frac{1}{2}(h_1^* \pm h_2)\,. \tag{7.5}$$

It is clear from the above formulae that when the distance between each pair of stacks is small enough, some of the scalars $H^{\pm}$ and $h^{\pm}$ become tachyonic, which will be the signal of spontaneous symmetry breaking in the SM. In this process the weak vector bosons Z_0 and $W^{\pm}$ get a mass but also the fermions get masses proportional to their Yukawa couplings to the different Higgs fields. The form of the Yukawa couplings among the SM fields in table 2 and the different Higgs fields are essentially fixed by conservation of gauge symmetries and have the general form [12]:

$$\begin{aligned} & y_j^U Q_L U_R^j h_1 + y_j^D Q_L D_R^j H_2 + \\ + \ & y_{ij}^u q_L^i U_R^j H_1 + y_{ij}^d q_L^i D_R^j h_2 + \\ + \ & y_{ij}^L L^i E_R^j H_2 + y_{ij}^N L^i N_R^j h_1 + \text{h.c.} \end{aligned} \tag{7.6}$$

where $i = 1, 2$ and $j = 1, 2, 3$. With this Yukawa structure one observes that, for example, if only the Higgs fields of type H_i get vevs, two down-like and one up-like quarks still remain massless, as well as all neutrinos. If in addition the Higgs of type h_i get a vev, all fermions get a mass.

As discussed in previous sections, it is clear that this Higgs sector of the square quiver is not supersymmetric. However one can obtain a electroweak scale well below the string scale M_s by "modest tuning" the distances $X_{bc^*}^2/2\pi = 2\alpha_2 - \Delta + \delta_b$, $X_{bc}^2/2\pi = 2\alpha_2 + \Delta + \delta_b$. Let us turn now to a description of the electroweak Higgs mechanism in the brane recombination language.

7.2 The mechanism of brane recombination

We are now going to show how the process of brane recombination yield results consistent with the above field theory description of the Higgs mechanism. For this to happen we have to show that

i. The final gauge group after brane recombination is indeed just $\mathrm{SU}(3) \times \mathrm{U}(1)_{em}$.

ii. The quarks and leptons become massive in the expected way.

Before proceeding let us first make a few comments about the recombination process. In general, even if the two initial intersecting branes α and β wrap factorizable cycles, the resulting recombined brane γ will not be factorizable. In this case the form of the recombined cycle will not have an easy geometrical description. Still, many properties of the recombination process will only depend on the general homology class of the cycles. In the case of non-factorizable cycles one has to work with a $2 \times 2 \times 2 = 8$-dimensional basis for the RR-charges of the cycles. Instead of wrapping numbers n^i, m^i, one has to work with 8-dimensional vectors $\vec{q}_k$, $k = 1 - 8$ (see appendix A). For a given system of branes to be recombinable into another the initial and final configurations must lie in the same homology class, the recombination must preserve RR charge:

$$\sum_{s=\text{initial branes}} N_s (1 + \Omega) \vec{q}_s = \sum_{f=\text{final branes}} N_f (1 + \Omega) \vec{q}_f \tag{7.7}$$

where $\Omega \vec{q}_s$ gives the charge vector of the orientifold mirror of each brane. Consider for example the case of the Higgs fields H_i in our class of models. These states come from the open string exchange between branes of type b and c^*. Thus this Higgs field taking a vev should be related to the recombination of one of the two branes of type b (let us call it b_1) with the brane c^* into a single non-factorizable 'e' brane with charge:

$$2\vec{q}_b + \Omega \vec{q}_c = (\vec{q}_{b_1} + \Omega \vec{q}_c) + \vec{q}_{b_2} \equiv \vec{q}_e + \vec{q}_{b_2} \tag{7.8}$$

Instead of two b-branes (which contained the $\mathrm{SU}(2)_L$ charged current electroweak interactions) we are now left with only one. Thus the electroweak symmetry has been broken. We can now check which chiral fermions, if any, remain in the massless spectrum by computing the new intersection numbers. The latter are easy to compute knowing that $I_{\alpha\beta}$

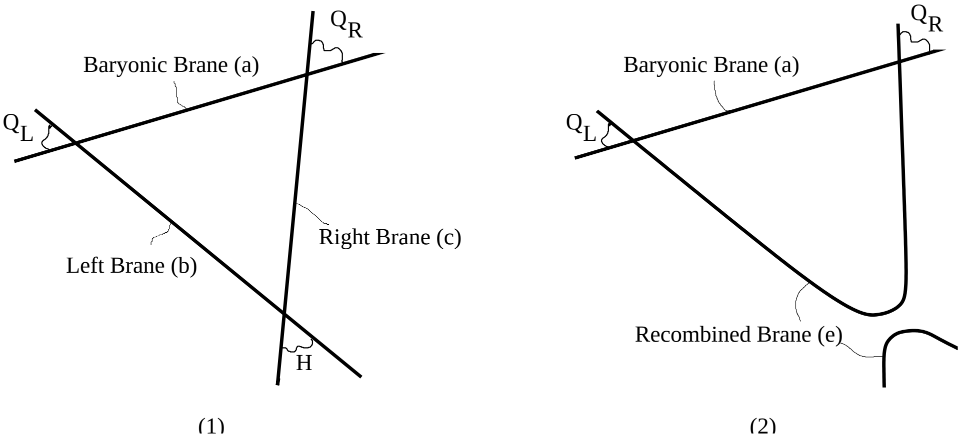

Figure 12: Picture of the recombination/symmetry breaking. (1) Before the recombination, the worldsheets connecting the Q_L, Q_R and Higgs multiplets have a trianglular shape with sides embedded in the branes. (2) After the recombination, branes 'b' and 'c' have recombined into brane 'e', giving a non-factorizable cycle with lower volume. Note that the intersection number between 'a' and 'e' is zero, but these branes are not parallel.

is a bilinear quantity in terms of $\vec{q}_\alpha$ and $\vec{q}_\beta$ (see appendix A). In our case we simply have $I_{ae} = I_{ab_1} + I_{ac^*}$ and $I_{ae^*} = I_{ab_1^*} + I_{ac}$. Thus the new intersection numbers *after* recombination are:

$$\begin{aligned} I_{ab_2} &= 1\,, & I_{ab_2^*} &= 2\,, \\ I_{ae} &= -2\,, & I_{ae^*} &= -1\,, \\ I_{b_2d} &= -3\,, & I_{b_2d^*} &= 0\,, \\ I_{de} &= 0\,, & I_{de^*} &= -3\,, \end{aligned} \tag{7.9}$$

Note that the number of chiral fermions has been reduced, only three quark flavours and some three leptons remain massless. This result is as expected from field theory arguments. Indeed, looking at (7.6) we see that if Higgs fields of type H_i get vevs, two up-like and one down-like quarks become massive and also charged leptons do, due to Yukawa couplings.

In figure 12 we can obtain an intuitive image of what may be happening. The Yukawa coupling amplitudes are obtained from correlators involving vertices of quarks and Higgs fields on the vertices of a triangle in the extra compact dimensions. The different Yukawa couplings are expected to be exponentially suppressed [10, 12] by the area of the corresponding triangle i.e., $Y_{abc^*} = \exp\left(-A_{abc^*}\right)$. After the brane c^* has recombined with one of the branes b, the triangle is smoothed out at the bc^* vertex. This corresponds to the Higgs fields H_i getting a vev. Note that the net contribution to I_{ae} due to the quarks at the other two vertices is zero, these quarks are no-longer chiral, have acquired masses.

At the end of the above recombination process we have four stack of branes a, b_2, e and d, suggesting a gauge group $\mathrm{SU}(3) \times \mathrm{U}(1)_a \times \mathrm{U}(1)_b \times \mathrm{U}(1)_e \times \mathrm{U}(1)_d$. Now, if our interpretation of the Higgs mechanism as a recombination process is correct, we should just obtain electromagnetic charge as our only surviving U(1) boson in the recombined system. Thus as a test we should check that indeed it is the electromagnetic charge with survives the recombination process. The U(1) combination with remains light $\mathrm{U}(1)_{\mathrm{massless}}$ is the one with no coupling to the RR closed string fields B_2^I, $I = 0,1,2,3$. This is the combination

$$\mathrm{U}(1)_{\mathrm{massless}} = \sum_{k=a,b_2,e,d} c^k\, \mathrm{U}(1)_k\,, \tag{7.10}$$

with coefficients c^k such that

$$\sum_k c^k \tilde{d}_k^I = 0\,, \qquad I = 0,1,2,3 \tag{7.11}$$

where $\tilde{d}_k^I$ are the coefficients of the couplings of RR fields to the four recombined U(1)'s, i.e., the coefficients in the couplings $\tilde{d}_k^I\ B_2^I \wedge F_k$. The corresponding coefficients for the starting cycles d_k^I are given in eq. (2.8). Using linearity of the RR charge in the recombination one can compute the new 4 $\tilde{d}_k^I$ coefficients in terms of the ones before recombination d_k^I as follows:

- $\tilde{d}_k^I = d_k^I$ for those brane stacks k not taking part in the recombination.
- $\tilde{d}_k^I = \frac{1}{N_\alpha} d_\alpha^I + \frac{1}{N_\beta} d_\beta^I$, for the new brane resulting from the recombination.

Using this and imposing eq. (7.11) before and after recombination one can compute explicitly the new $\tilde{d}_k^I$'s and obtain which U(1)'s remain massless. In the present case, if before recombination we had as the only surviving U(1) the hypercharge of eq. (2.10) (with $r = -1$), after recombination there is only one remaining massless generator given by

$$Q_{em} = \frac{1}{6} Q_a + \frac{1}{2} Q_e - \frac{1}{2} Q_d - \frac{1}{2} Q_{b_2}\,. \tag{7.12}$$

One can compute the charges of the chiral fermions that remain massless and one finds the residual spectrum (obtained from the intersection numbers (7.9)) after recombination to be:

- *Gauge group:* SU(3)$\times$ $\mathrm{U}(1)_{em}$
- *Matter content:* $2(3)_{-\frac{1}{3}} + 2(\bar{3})_{\frac{1}{3}} + 1(3)_{\frac{2}{3}} + 1(\bar{3})_{-\frac{2}{3}} + 3(1)_0 + 3(1)_0$

Thus we see we have one massless u-quark, two massless d-like quarks (i.e. down and strange) and massless (left and right) neutrinos, as expected from the field theory arguments from the standard Yukawa coupling of the Higgs fields H_i to fermions. This is interesting and in fact not far from the experimental situation in which three quark flavours are relatively light (leading to Gellman's flavour SU(3)) and neutrinos are almost massless.

From the effective lagrangian point of view, we could have in addition a vev for the other Higgs fields of type h_i. If this is the case the rest of the fermions would now become massive, but still a charge generator should remain unbroken. From the brane recombination point of view this corresponds to the recombination of the new brane 'e' with the 'b_2^*' brane which was an spectator in the first recombination. The new recombined brane 'f' will have a RR charge vector $\vec{q}_f \equiv (\vec{q}_e + \Omega\vec{q}_{b_2})$. We are thus only left with three stacks of branes, a,d and f. One can easily check that the intersection numbers between these branes is zero. For example, $I_{af} = I_{ae} + I_{ab_2^*} = -2 + 2 = 0$. Thus there are no massless fermions left after this second recombination, again as expected from field theory arguments. One might worry that, since we have now only three stacks of branes but we still have four RR fields B_2^I, all U(1)'s could become massive and we would be left with no photon in the low energy spectrum. This turns out not to be the case. One can easily check that a U(1) generator

$$Q_{\text{massless}} = \frac{1}{6}Q_a + \frac{1}{2}Q_f - \frac{1}{2}Q_d \tag{7.13}$$

remains unbroken. This massless generator can be identified with electromagnetism by noting that e.g., the massive fields stretched between branes a and d have charges which correspond to the electric charge they had before this last recombination.

Note that from the effective field theory point of view there are other (less interesting) field directions of the scalar potential, particularly if we include in the complete potential additional massive scalars which appear at the intersections of branes b and c with their mirrors b^* and c^*. In particular, open strings stretched between the branes b and b^* contain (massive) scalars transforming like symmetric (triplet) and antisymmetric (singlet) $\mathrm{SU}(2)_L$ representations which couple to the SM Higgs doublets H_i and h_i and modify the scalar potential. These other field directions could in principle lead to unwanted vacua with breaking of electromagnetic U(1). From the field theory point of view this can be controlled by making the unwanted fields heavy, by appropriately separating each brane from its mirror.

These different Higgsing possibilities also exist in the brane recombination language as coming from different choices for brane recombination. The brane recombination discussed above describing the SM Higgs mechanism correspond to the recombination of branes:

$$b_1 + b_2^* + c^* \to e + b_2^* \to f \tag{7.14}$$

Recombining e.g., $b_1 + b_2 + c^*$ or $b_1 + b_2 + c$ would have lead to other less interesting final configurations with e.g., broken electromagnetic charge. Thus from the phenomenological point of view these other possibilities should be somehow energetically disfavoured.

We have seen how the Higgs mechanism in the SM may be understood as a brane recombination process in which the three branes giving rise to the $\mathrm{U}(2)_b \times \mathrm{U}(1)_c$ recombine into a single brane giving rise to a final (in general non-factorizable) brane f. There is only one massless abelian generator which can be identified with the standard electromagnetic charge. Thus the actual string vacuum after recombination involves only three stacks of branes a,d and f and the gauge group is just $\mathrm{SU}(3)_{QCD} \times \mathrm{U}(1)_{em}$.

There is a couple of questions which appear in such an interpretation:

i). The recombination language shows us how the quarks and leptons become massive at each step, but how can one explain the observed existence of large hierarchies of fermion masses? We already mentioned that in the initial picture in which the a, b, c, d branes wrap factorizable cycles, there are Yukawa couplings between Higgs fields and chiral fermions (see figure 12). The worldsheet with one right-handed fermion, one left-handed fermion and one Higgs field have a triangular shape with those particles at the corners [10]. As we mentioned above, the corresponding coupling may be exponentially suppressed. After the branes recombine (i.e., after the Higgs get vevs) the vertex where the Higgs field lies is smoothed out (see figure 12). Still, if the vev of the Higgs is small compared to the string scale, this will amount to a small perturbation, fermion masses would still be exponentially suppressed by the area left between the recombined brane e and the baryonic a-stack (d-stack in the case of leptons), which is the smoothed out triangle. Thus the fermion masses may still have hierarchical ratios.

ii). The actual final vacuum contains only the stacks a, d and f and only $\mathrm{SU}(3)_{QCD} \times \mathrm{U}(1)_{em}$ as gauge group. If we put energy in the system, we should be able to see the electroweak gauge bosons Z_0 and $W^{\pm}$ produced, with masses of order the vev of the Higgs fields. Which are those states in the final recombined system? The Z_0 is neutral and should correspond to open strings beginning and ending on the same recombined brane 'f'. In particular, it should correspond to an open string stretching between the opposite smoothed portions of that brane in figure 12. On the other hand open strings stretching between the recombined brane f and its mirror f^* have charge $= \pm 1$ (see eq. (7.13)) and should give rise to the charged $W^{\pm}$ bosons.

7.3 Higgs mechanism and brane recombination in the triangle quiver

Let us now discuss for comparison the Higgs mechanism and brane recombination in the triangle quiver. Those models have an extra massless $\mathrm{U}(1)_{B-L}$ but we will concentrate first on electroweak symmetry breaking. One important difference is that in these models there are massless chiral multiplets corresponding to Higgs fields at the intersections of the stacks (b) and (c) (see table 11). In the case $\beta^2 = 1$ there is in fact a single Higgs set and the charged chiral massless spectrum is just that of the MSSM (plus right-handed neutrinos). Looking at the charges in table 11 one sees that the only allowed Yukawa superpotential couplings are:

$$\begin{aligned} & y^u_{ij} q^i_L U^j_R \bar{H} + y^d_{ij} q^i_L D^j_R H + \\ & + y^L_{ij} l^i E^j_R H + y^N_{ij} l^i N^j_R \bar{H} + \text{h.c.} \end{aligned} \tag{7.15}$$

with $i = 1, 2$ and $j = 1, 2, 3$. Thus one would say that one of the generations of quarks and leptons does not get masses at this level.

Let us now see how would be the Higgs electroweak symmetry breaking in the brane recombination language. As in the square quiver case, let us assume that one of the two (b) branes, e.g., b_1 recombines with c^*. Looking at table 11 we see that this should correspond to $\bar{H}$ getting a vev. After the recombination $b_1 + c^* \to e$ into a single brane e the electroweak symmetry is broken. We are left with four stacks of branes a, b_2, e and d

and the intersection numbers are given by

$$\begin{aligned}
I_{ab_2} &= 1\,, & I_{ab_2^*} &= 2\,,\\
I_{ae} &= -2\,, & I_{ae^*} &= -1\,,\\
I_{db_2} &= 1\,, & I_{db_2^*} &= 2\,,\\
I_{de} &= -2\,, & I_{de^*} &= -1\,,\\
I_{b_2e} &= -\frac{1}{\beta^2}\,, & I_{b_2e^*} &= -\frac{1}{\beta^2}\,,\\
I_{ee^*} &= \frac{2}{\beta^2}
\end{aligned} \tag{7.16}$$

with $\beta^2 = 1, 1/2$. It is easy to check now that (apart from the $B-L$ symmetry which remains unaffected by this process and will be discussed later) there is an unbroken massless gauge generator which can be identified with electromagnetism

$$\mathrm{U}(1)_{em} = \frac{\mathrm{U}(1)_a}{6} - \frac{\mathrm{U}(1)_d}{2} - \frac{\mathrm{U}(1)_{b_2}}{2} + \frac{\mathrm{U}(1)_e}{2}\,. \tag{7.17}$$

and is not rendered massive by couplings to RR B-fields. The massless chiral spectrum after recombination obtained from the above intersection numbers[10] is (for $\beta^2 = 1$)

- *Gauge Group:* $\mathrm{SU}(3) \times \mathrm{U}(1)_{em} \times [\mathrm{U}(1)_{B-L}]$
- *Matter Content:* $2(3)_{-\frac{1}{3}} + 2(\bar{3})_{\frac{1}{3}} + 1(3)_{\frac{2}{3}} + 1(\bar{3})_{-\frac{2}{3}} + 2(1)_{-1} + 2(1)_1 + (1)_0 + (1)_0 + [(1)_1 + (1)_0 + (1)_{-1}]$

Thus at this level, with only $b_1 + c^*$ recombination, we are left with two D-quarks, one U-quark, two charged leptons and one (Dirac) neutrino. In addition there are charged and neutral Higgsinos (last three states in brackets).

Comparing this massless spectrum with the one expected from the field theory Yukawa couplings with a vev for the Higgs $\bar{H}$, we see that one extra D-quark and one charged lepton became massive after recombination, which was not expected from the effective field theory lagrangian. It seems that this puzzling result can be understood as follows. Consider the brane intersection bc^* giving rise the the Higgs field $\bar{H}$ (see table 11). In addition to the $\bar{H}$ massless chiral multiplet, there are at this intersection other *massive* vector-like pairs with the opposite Q_b and Q_c charges to $\bar{H}$. These are in some way $\mathcal{N} = 4$ partners of the massless Higgs multiplet (see ref. [15, section 4.1]). Thus at the intersection nothere are massive $Y = -1/2$ fields H' with $Q_{b,c}$ charges qual to $(1, 1)$. Note that such H' fields have precisely the gauge quantum numbers required to couple to Q_L and D_R (see table 11), so that if a vev is induced for H' a D-quark (and a charged lepton) will become massive. Thus what it seems is happening here is that when the branes b_1 and c^* recombine at this intersection it is not only the lightest (tachyonic) scalar which is involved but also massive excitations. Recall however that our recombination arguments are purely topological in

[10] In order to compute the chiral fermion spectrum from an intersection of a brane j with its mirror j^* one should use the general formula (A.7) of appendix A. Since in the present case $I_{j,ori} = 0$, the number of chiral fermions is just half the intersection number. This is also true for the spectrum in table (16).

character, they tell us who becomes massive and who remains massless, but it does not tells us how big are the masses. It is reasonable to expect that since massive modes are involved, the masses of these D-quark and charged lepton are small, so that they can perhaps be identified with the d-quark and the electron.

As in the square quiver example, further recombination of the spectator brane b_2^* with the brane e, into a final brane f, $b_2^* + e \to f$ renders massive all of the fermions. Indeed, using bilinearity of the intersection numbers it is easy to check that all intersection numbers vanish. Still an electromagnetic generator

$$\mathrm{U}(1)_{em} = \frac{\mathrm{U}(1)_a}{6} - \frac{\mathrm{U}(1)_d}{2} + \frac{\mathrm{U}(1)_f}{2}\,. \tag{7.18}$$

can be checked to remain in the massless spectrum, i.e., it does not receive any mass from couplings to B-fields. Thus the electroweak symmetry breaking process is completed.

We have not addressed here the question of what triggers electroweak symmetry breaking in the triangle quiver. The Higgs multiplets are massless to start with but the scalars get in general masses from loops and FI-terms, as explained in the previous section. In addition there is the one-loop contribution from top-quark loops which will tend to induce a negative mass^2 to the $\bar{H}$ higgs in the usual way [27]. We are assuming here that those effects combined yield a negative $mass^2$ to the Higgs fields.

7.4 Breaking of B-L and neutrino masses from brane recombination

We have seen that in all of the four classes of models constructed (except for the square quiver with $\Delta \neq 0$) there is an extra gauge boson corresponding to $B-L$ in the massless spectrum. This extra U(1) may be Higgssed in a variety of ways but perhaps the simplest would be to give a vev to some right-handed sneutrino, which are always present in these models. We would like to discuss in this section how this process would be interpreted in the brane recombination language and how it implies the appearance of Majorana neutrino masses.[11]

Let us consider for definiteness the triangle quiver model. Its massless chiral spectrum is the one of the MSSM with right-handed neutrinos. In addition to hypercharge we have the $B-L$ generator at the massless level. The breaking of this extra gauge symmetry could be obtained if some right-handed sneutrino $\tilde{\nu}_R$ gets a vev. The stringy counterpart of this process would be the recombination of branes c and d in the model. So let us assume there is a recombination $c + d \to j$ into a final brane j. We will be left with three stacks of branes now a,b and j. Using the results in (4.10) and the linearity of intersection numbers

[11]Note that in the models of ref. [12] and in the subset of them formed by the models of section 5.1 in the present paper, there is no extra $\mathrm{U}(1)_{B-L}$ which should be spontaneously broken. Thus it is not required to spontaneously break lepton number by giving a vev to a $\tilde{\nu}_R$. In such a case there can only be Dirac neutrino masses, as argued in ref. [12]. In models with an extra $\mathrm{U}(1)_{B-L}$ we will be forced to break lepton number conservation anyhow and at some level Majorana masses will appear, as discussed below.

Intersection	Matter fields		Q_a	Q_b	Q_j	Y
ab	Q_L	$2(3,2)$	1	-1	0	1/6
ab^*	q_L	$(3,2)$	1	1	0	1/6
aj	U_R	$3(\bar{3},1)$	-1	0	1	$-2/3$
aj^*	D_R	$3(\bar{3},1)$	-1	0	-1	1/3
bj	L	$2(1,2)$	0	-1	1	$-1/2$
bj^*	l	$(1,2)$	0	1	1	$-1/2$
$jj*$	E_R	$3(1,1)$	0	0	-2	1

Table 16: Chiral spectrum in the triangle quiver after cd^* brane recombination.

one gets

$$\begin{aligned} I_{ab} &= 1\,, & I_{ab_2^*} &= 2\,, \\ I_{aj} &= -3\,, & I_{aj^*} &= -3\,, \\ I_{bj} &= -1-\frac{1}{\beta^2}\,, & I_{bj^*} &= 2-\frac{1}{\beta^2}\,, \\ I_{jj^*} &= -6\,. \end{aligned} \tag{7.19}$$

One can then easily check that there is only one U(1) with no couplings to RR-fields and hence remains massless. It is given by

$$\mathrm{U}(1)_Y = \frac{\mathrm{U}(1)_a}{6} - \frac{\mathrm{U}(1)_j}{2}\,. \tag{7.20}$$

and corresponds to standard hypercharge. The resulting fermionic spectrum is shown in table 16, for the case $\beta^2 = 1$.

Note that this fermion spectrum is the one of the SM *without right-handed neutrinos*. The breaking of electro-weak symmetry may then proceed as in the previous subsection. Doublet Higgs scalars may be provided both by the scalar partners of the fermionic doublets in the table but also they could be provided by the $H, \bar{H}$ fields of the original model in table 11 if they remained relatively light upon recombination, which is something which will depend on the detailed geometry of the recombined branes.

Note also that in the recombination process not only $\mathrm{U}(1)_{B-L}$ has acquired a mass but also the right-handed neutrinos have dissapeared from the massless spectrum. In addition, compared to the spectrum in table 11, a pair of $\mathrm{SU}(2)_L$ doublets have gained masses. The latter was expected from the point of view of the effective field-theory Higgs mechanism. Looking at the superpotential in (7.15) one observes that a vev for $\tilde{\nu}_R$'s mixes leptons with Higgsinos. On the other hand the fact that the ν_R's get massive is less obvious from the effective field theory point of view, since there are apparently no renormalizable Yukawa couplings giving (Majorana) masses to them. As in the case of the unexpected massive fermions upon EW symmetry breaking discussed in the previous section, an understanding of this fact seems to imply that brane recombination involves in the process also the effects of massive fields. There are massive chiral fields with opposite charges to those of the $\tilde{\nu}_R$'s at the intersections. In general dim=5 couplings of type $(\nu_R\nu_R\tilde{\nu}_R^*\tilde{\nu}_R^*)$ will give masses to ν_R's once the sneutrinos get vevs.[12] Whatever the low-energy field theory interpretation, it is a fact that the right-handed neutrinos get massive upon $c+d$ brane recombination. This is interesting in itself since it has always been difficult to find mechanisms in string-theory giving rise to Majorana masses for neutrinos. It seems that brane recombination gives one possible answer.

One interesting question is what happens now with the left-handed neutrinos. The gauge group of the above model after this recombination is just $\mathrm{SU}(3)\times\mathrm{SU}(2)\times\mathrm{U}(1)_Y$. Once the $\tilde{\nu}_R$'s get vevs, lepton number is violated and there is no longer distinction between sleptons and Higgs fields.[13] Since there are no ν_R's left in the massless spectrum one may argue that there are no possible Dirac fermions for neutrinos and an effective field theory analysis would suggest that the left-handed neutrinos should remain massless after electroweak symmetry breaking. The brane recombination analysis tells us that this will not be the case and after full brane recombination of the branes involved in lepton number violation *and* electroweak symmetry breaking:

$$b_2+b_1^*+c+d\rightarrow f \tag{7.21}$$

there are no massless fermions left. Indeed, using bilinearity of intersection numbers plus eq. (4.10) it is easy to check that $I_{af}=I_{af^*}=I_{ff^*}=0$. So somehow left-handed neutrinos have managed to become massive also after full brane recombination. As in previous cases a possible field theory interpretation of the neutrino masses is the importance of the massive states at the intersections. Dimension 5 operators of the form $(LL\bar{H}\bar{H})$ can give Majorana masses to left-handed neutrinos if the scalars in $\bar{H}$ get a vev.

An important point would be to know the size of neutrino masses after the full recombination/Higgsing process. Unfortunately the intersection numbers are topological numbers which count the net number of massless fermions but do not give as any information on the size of the masses of the non-chiral fermions. The masses (Yukawa couplings) are geometrical quantities which depend on the precise locations of the wrapping branes. Note in particular that although the initial a,b,c,d stacks of branes of this model are factorizable branes with intersections respecting the same supersymmetry, after any recombination takes place the resulting recombined brane is in general non-factorizable and we do not know the precise shape of the cycle that the brane is wrapping. Thus we cannot compute in detail aspects like Yukawa couplings unless we get that geometrical information. On the other hand it is reasonable to expect that, if the $\tilde{\nu}_R$-vevs are of order the string scale, since the left-handed neutrinos get their masses involving massive modes of order M_s, see-saw-like Majorana masses of order $m_\nu\propto|H|^2/M_s$, with a model dependent coefficient which will depend on the geometry and can perfectly be rather small. So neutrino masses within the experimental indications could be obtained.

[12]In addition there will be mixing of the ν_R's with the gauginos of the $\mathrm{U}(1)_{B-L}$ generator.

[13]This is rather similar to R-parity violating models with lepton number violation.

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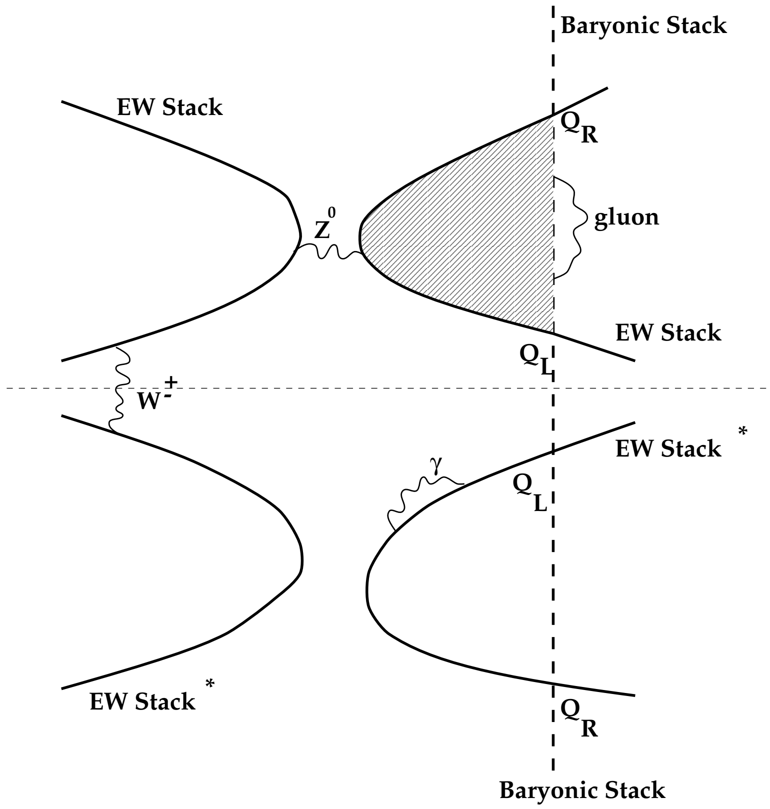

Figure 13: Artist's view of a portion of the *baryonic* (thick dashed line) and *EW* (continuous line) brane configurations after full brane recombination. The mirror of the EW-brane is also depicted below. Massless gluons and photons come from open strings starting and ending on the baryonic and EW branes. The Z^0 and $W^{\pm}$ bosons come from massive open strings states stretching in between different regions of the same EW and/or its mirror EW*. The mass of the quark is exponentially suppressed by the area of shaded world-sheet.

In the scheme we are discussing, after full recombination of the *left*, *right* and *lepton* branes into a single brane, the observed SM particles and interactions would come from open string exchange between a couple of brane stacks (see figure 13 for an artist's view):

- *Baryonic stack.* It contains three parallel branes (and mirrors) and the gauge group $\mathrm{SU}(3) \times \mathrm{U}(1)_a$, with $\mathrm{U}(1)_a$ gauging baryon number. QCD gauge bosons originate in this stack.

- *Electro-Weak stack (f).* It contains only one brane (plus mirror) with a $\mathrm{U}(1)_{EW}$ gauge boson in its worldvolume. It results from the recombination of *left*, *right* and *lepton* branes.

A linear combination of $\mathrm{U}(1)_a$ and $\mathrm{U}(1)_{EW}$ may be identified with the photon and it is massless. The orthogonal combination gets massive by combining with a RR B-field, but its symmetry remains as a global symmetry in perturbation theory, guaranteeing proton stability, as in the examples in previous sections.

Property	Square	Linear	Rombic	Triangle
# different $\mathcal{N}=1$ SUSY's	4	2	3	1
Gauge group	SM (+U(1))	SM+U(1)	LR	SM+U(1)
SUSY Higgs	no	no	yes	yes
Minimal Higgs system	no	no	no	yes
$m^2_{\tilde{q},\tilde{l}} > 0$ from FI	yes	yes	no	no
String Scale	$\sim 10\,\mathrm{TeV}$	$\sim 10\,\mathrm{TeV}$	$\sim 10\,\mathrm{TeV}$	$\sim 10\,\mathrm{TeV}$ -M_{Planck}
Additional RR sources	no	yes	yes	yes

Table 17: Some general aspects of the four classes of SUSY-quiver models with three quark/lepton generations discussed in the text. The linear and triangle quivers lead to an extra $\mathrm{U}(1)_{B-L}$ in the massless spectrum. In the square quiver case the presence of such an extra U(1) is optional. The rombic example provided is a left-right symmetric model.

The $W^{\pm}$,Z^0 bosons correspond to massive string states stretching between different sectors of the electroweak brane (f) or its mirror (f^*). The masses of fermions of the SM will depend on the detailed geography of the configuration. All quarks and leptons are massive but some of their masses may be very small due to the fact that they are located at different distances from the location of the Higgs fields, and some Yukawa couplings are naturally exponentially suppressed, as discussed in ref. [10]. Thus every SM parameter would have a reflection in terms of the detailed geometry of the underlying *baryonic* and *electro-weak* stacks.

8. Final comments and conclusions

Let us make a number of general comments about the explicit intersecting D6-brane models constructed in the previous sections. They consist of type IIA string compactifications on T^6 along with an orientifold operation. We have concentrated on configurations with four stacks of D6-branes : a (*baryonic*), b (*left*), c (*right*), and d (*leptonic*). This is the simplest structure capable of giving rise to all fermions of the SM in bifundamental representations of the underlying group. By varying the geometry of the tori, one can obtain models in which all brane intersections respect *some* $\mathcal{N}=1$ supersymmetry. There are four classes of models with four stacks of branes with SUSY-quiver structures shown in figure 6. They are called respectively square, linear, rombic and triangle quivers depending on the SUSY-quiver structure. We have built explicit models corresponding to these four classes. Some general characteristics of these models are displayed in table 17.

The four classes of models have quarks, leptons and Higgs multiplets respecting in general different SUSY's, as recorded in the table. In general, in addition to the SM gauge group, the property of Q-SUSY (or SUSY) seems to force the presence of an additional U(1) generator corresponding to $\mathrm{U}(1)_{B-L}$, an abelian symmetry well known from left-right symmetric models. This property is quite intriguing since there is in principle no obvious connection between the SUSY properties of a theory and the gauge groups one is gauging. In the case of the square quiver one can give mass to that additional U(1) by departing from the SUSY limit (i.e. $\Delta \neq 0$).

The models from the square and linear quiver have non-SUSY Higgs sectors. If, as suggested in ref. [15], one insists in getting one-loop protection of the Higgs particles to stabilize the weak scale, models with triangle or rombic quiver structure would be required. In particular, in the case of the triangle quiver, specific configurations with the minimal SUSY Higgs sector can be obtained. The latter triangle models have (for $\beta^2 = 1$) the chiral content of the MSSM (with right-handed neutrinos). From this point of view the triangle class of models are particularly attractive. They are also attractive because they predict the presence of a SM Higgs sector. Indeed we showed how, in order for left- and right-handed multiplets to respect the same $\mathcal{N} = 1$ SUSY, brane configurations are forced to have intersections at which Higgs sectors arise. This is an interesting property which is not present in schemes which follow the unification route like SUSY-GUT, CY_3 and/or Horava-Witten heterotic compactifications etc.[14]

On the other hand, we have seen that going slightly away from the Q-SUSY limit, one can parametrize the corresponding SUSY-breaking in terms of FI-terms. As remarked in the table, one can check that in the square and linear quivers all squarks and sleptons can get positive $(mass)^2$ from FI-terms whereas that is not the case for the triangle and rombic quivers. In the latter cases the leading source of SUSY-breaking scalar masses should come from loop effects, as in figure 5.

The general consistency of any of these D6-brane configurations require global cancellation of the total RR charge. In the case of the square quiver it is easy to find simple choices of D6-branes wrapping factorized cycles on the torus and with no intersection with the SM branes, so that all RR-tadpoles cancel. The other three types of quivers are equally consistent from this point of view, although in general non-factorizable extra branes and additional RR H-flux may be required to cancel tadpoles.

In D6-brane models like this, the standard way [28] to lower the string scale M_s compared with the Planck scale M_{Planck} by making large some of the torus radii cannot be performed.[15] Note also that the intersecting brane structure is not necessarily linked to a low string scale (i.e., $M_s \approx$ 10-100 TeV) hypothesis. Consider in particular a triangle quiver configuration, where all intersections preserve the *same* $\mathcal{N} = 1$ SUSY. In principle one can consider this triangle as part of some bigger $\mathcal{N} = 1$ brane configuration somewhat analogous to the class of models in ref. [14]. One could then translate most of our discussion (FI-terms, brane recombination, gauge coupling constants) on triangle quivers to that configuration. This is what we mean by indicating in the table that the string scale for the triangle quiver case is ~ 10 TeV-M_{Planck}.

Independently of the particular models discussed, we have presented a brane interpretation of the SM Higgs mechanism. It is important to realize that the familiar brane interpretation of the Higgs mechanism in terms of the separation of parallel branes is not appropriate for the Higgs mechanism of the SM. Brane separation does not lower the rank of the gauge group and corresponds to adjoint Higgsing. We claim that the appropriate

[14]Also note that in those unification schemes the stability of the proton relies on the asumed presence of symmetries like R-parity or generalizations as well as some doublet-triplet splitting mechanism. In the present case proton stability is a consequence of the gauge character of baryon symmetry.

[15]For a discussion of this point see [8, 9, 15] and references therein.

brane interpretation of the SM Higgs mechanism (and analogous Higgsings lowering the rank) is brane recombination.[16] In our approach the non-abelian weak interactions $SU(2)_L$ live on the worldvolume of two parallel branes i.e., b_1,b_2. If the Higgs field comes from open strings exchanged between branes b and c^* (as in the models constructed), what will happen is that one of the two parallel b-branes (say, b_1) will fuse with c^* giving rise to a single brane $b_1 + c^* \to e$. This is the brane recombination mechanism. Since each brane comes along with its own U(1), it is obvious that the rank has been reduced in the process. At the same time the number of chiral fermions (computed from the intersection numbers of the residual branes left) may be shown to decrease (or vanish), corresponding to the Higgs field giving masses to chiral fermions through Yukawa couplings.

An analogous brane recombination interpretation exists for a process in which a right-handed sneutrino gets a vev. In this case lepton number is broken and at the same time it is shown that both right-handed and left-handed neutrinos get independent (i.e., Majorana masses). If one consider this process in the specific models constructed (like the square or the triangle models) one finds that the distinction between Higgs multiplets and sleptons dissappear, and the properties of the models are somewhat similar to R-parity violating SUSY models with lepton number violation. In this respect we note that it seems quite difficult within the brane intersection scheme to have lepton number violating neutrino masses without having at the same time L-violating dimension-four couplings, both things come along once brane recombination takes place.

At the end of the day, in these schemes the observed SM would have a description in terms of two final recombined brane-stacks : a *baryonic* stack and an *electroweak stack* supporting a $SU(3) \times U(1)_{em}$ gauge group. Every SM parameter would have a reflection in terms of the detailed geometry of the underlying *baryonic* and *electro-weak* stacks.

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A. Branes wrapping general cycles.

As we have mentioned in the text, when dealing with more general intersecting D-brane constructions or considering a generic configuration after some brane recombination has taken place, we are naturally led to consider branes wrapping general cycles.

In our particular setup, a $D6_a$-brane wrapping a general cycle will be located in a 3-submanifold of $T^6 = T^2 \times T^2 \times T^2$, thus corresponding to an element $[\Pi_a]$ of $H_3(T^6, \mathbb{Z})$, which is the group of homology classes of 3-cycles [29]. It turns out that $H_3(T^6, \mathbb{Z})$ is a discrete vector space, so any of its elements can be represented by a vector with integer entries. This vector space has dimension 20. One particular subset of $H_3(T^6, \mathbb{Z})$ is given

[16] An analogous proposal in the context of D4-brane models was put forward in ref. [10].

q comp.	3-cycle	factor. comp.
q_1	$[a_1] \times [a_2] \times [a_3]$	$n^1 n^2 n^3$
q_2	$[b_1] \times [b_2] \times [b_3]$	$m^1 m^2 m^3$
q_3	$[a_1] \times [b_2] \times [b_3]$	$n^1 m^2 m^3$
q_4	$[b_1] \times [a_2] \times [a_3]$	$m^1 n^2 n^3$
q_5	$[b_1] \times [a_2] \times [b_3]$	$m^1 n^2 m^3$
q_6	$[a_1] \times [b_2] \times [a_3]$	$n^1 m^2 n^3$
q_7	$[b_1] \times [b_2] \times [a_3]$	$m^1 m^2 n^3$
q_8	$[a_1] \times [a_2] \times [b_3]$	$n^1 n^2 m^3$

Table 18: Definition of the $\vec{q}$-basis of T^6 in terms of products of 1-cycles.

by what we have called *factorizable cycles*. These are 3-cycles that can be expressed as products of 1-cycles on each T^2 (see figure 1 for an example).[17] Any of those cycles can be expressed by 6 integers as $[\Pi_a] = \prod_{i=1}^3 [(n_a^i, m_a^i)]$, where $n_a^i, m_a^i \in \mathbb{Z}$ describe the 1-cycle the D6$_a$-brane is wrapping on the i^{th} torus. Factorizable cycles can be easily described geometrically, which allows us to compute many phenomenologically interesting quantities. For instance, we can compute the lightest bosonic spectrum living at the intersection of two factorizable D6-branes a and b by simply computing the angles they form on each torus and using (3.1). Thanks to this, we were able to check whether certain $\mathcal{N} = 1$ supersymmetries were preserved at each intersection and define Q-SUSY or SUSY models in this way. Unfortunately, we will be unable to study which supersymmetries are left, if any, after a brane recombination process has taken place. Due to this fact, we have constructed our particle physics models mostly using branes wrapping factorizable cycles, although non-factorized branes will be generically unavoidable after brane recombination.

It turns out that factorizable 3-cycles are not a vector subspace of the homology group $H_3(T^6, \mathbb{Z})$. Indeed, the sum of two factorizable cycles is not, in general, a factorizable cycle. This is an important point, since the homology class $[\Pi_a]$ where a D6$_a$-brane lives determines its RR charges. When two branes a and b fuse into a third one c in a brane recombination process the total RR charge should be conserved, which implies that the final brane will lie in a 3-cycle such that $[\Pi_c] = [\Pi_a] + [\Pi_b]$. Even if we start with a configuration where every brane is factorizable, we are finally led to consider non-factorizable branes as well. To this regard, we will consider the smallest vector subspace of $H_3(T^6, \mathbb{Z})$ that contains factorizable 3-cycles. This is $[H_1(T^2, \mathbb{Z})]^3$, and its dimension is $2^3 = 8$. Following [30], we define a basis on this subspace by table 18, where each element of this basis can be expressed as a product of 1-cycles $[a_i]$, $[b_i]$ of the i^{th} T^2. Each general cycle $[\Pi_a]$ under consideration can then be expressed by a vector $\vec{q}_a$, whose 8 integer components are defined above. In addition, a factorizable 3-cycle will correspond to vector $\vec{q}$ whose components are given in the third column above.[18] See [30] for more details on this construction and some other features involving non-factorizable cycles.

[17] Notice that for this definition to make sense we also need our T^6 to be factorized as $T^2 \times T^2 \times T^2$.

[18] In the construction of the q-basis we are not considering fractional wrapping numbers, so both n and m are integer, see below.

The usefulness of this $\vec{q}$-basis formalism comes from the fact that quantities as the intersection number of two branes a and b can be easily expressed as bilinear products involving an intersection matrix. That is, it can be expressed as

$$I_{ab} \equiv [\Pi_a] \cdot [\Pi_b] = \vec{q}_a^{\,t} I \vec{q}_b \,, \tag{A.1}$$

where the intersection matrix is given by

$$I = \begin{pmatrix} 0 & 1 & & & & & & \\ -1 & 0 & & & & & & \\ & & 0 & 1 & & & & \\ & & -1 & 0 & & & & \\ & & & & 0 & 1 & & \\ & & & & -1 & 0 & & \\ & & & & & & 0 & 1 \\ & & & & & & -1 & 0 \end{pmatrix} \tag{A.2}$$

When dealing with orientifold compactifications each generic brane a must be accompanied by its mirror image a^*. In this q-basis formalism, the corresponding vectors can be related under the action of a linear operator Ω, such that

$$\Omega \vec{q}_a = \vec{q}_{a^*} \,, \quad \Omega^2 = \mathrm{Id} \,. \tag{A.3}$$

As we already mentioned, the geometrical action associated to Ω amounts to a reflection on each complex internal dimension. In terms of a 1-cycle (n^i, m^i) wrapping on the i^{th} T^2, this translates into

$$\Omega : (n^i, m^i) \mapsto (n^i, -m^i - 2b^i n^i) \,, \tag{A.4}$$

where $b^i = 0, \frac{1}{2}$ is the T-dual discrete NS background defined in section 2, and related to the complex structure of the i^{th} torus. From this we can deduce the action of the operator Ω on a general vector $\vec{q}$ describing a D6-brane:

$$\Omega = \begin{pmatrix} 1 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\ -8b^1b^2b^3 & -1 & -2b^1 & -4b^2b^3 & -2b^2 & -4b^1b^3 & -2b^3 & -4b^2b^1 \\ 4b^2b^3 & 0 & 1 & 0 & 0 & 2b^3 & 0 & 2b^2 \\ -2b^1 & 0 & 0 & -1 & 0 & 0 & 0 & 0 \\ 4b^1b^3 & 0 & 0 & 2b^3 & 1 & 0 & 0 & 2b^1 \\ -2b^2 & 0 & 0 & 0 & 0 & -1 & 0 & 0 \\ 4b^1b^2 & 0 & 0 & 2b^2 & 0 & 2b^1 & 1 & 0 \\ -2b^3 & 0 & 0 & 0 & 0 & 0 & 0 & -1 \end{pmatrix} \tag{A.5}$$

Notice that this linear operator should leave fixed the cycle where the O6-plane lies, which is

$$[\Pi_{ori}] = \bigotimes_{i=1}^{3} \left(\frac{1}{1-b^i} [a_i] - 2b^i [b_i] \right) , \tag{A.6}$$

and that can be easily translated into a vector $\vec{q}_{ori}$. The chiral massless spectrum arising at general cycles intersections is given by

$$\sum_{a<b}\left[I_{ab}(N_a,\bar{N}_b)+I_{ab^*}(N_a,N_b)\right]$$
$$\sum_{a}\left[8\beta^1\beta^2\beta^3 I_{a,ori}(\mathbf{A}_a)+\frac{1}{2}\left(I_{aa^*}-8\beta^1\beta^2\beta^3 I_{a,ori}\right)(\mathbf{A}_a+\mathbf{S}_a)\right] \tag{A.7}$$

where $\mathbf{S}_a$ $(\mathbf{A}_a)$ stands for the (anti)symmetric representation of the gauge group $\mathrm{U}(N_a)$. Here $\beta^i=1-b^i$, as defined in the text. All intersections involved in (A.7) can be computed in this q-basis formalism, as for instance

$$I_{ab}=\vec{q}_a^{\,t} I\Omega\vec{q}_b\,, \qquad I_{a,ori}=\vec{q}_a^{\,t} I\vec{q}_{ori}. \tag{A.8}$$

Remarkably enough, RR tadpoles cancellation have an extremely simple expression in this formalism

$$\sum_a N_a\left(\vec{q}_a+\Omega\vec{q}_a\right)=Q_{ori}\vec{q}_{ori} \tag{A.9}$$

where Q_{ori} represents the relative charge between the O6-plane and the D6-branes and is given by $Q_{ori}=32\beta^1\beta^2\beta^3$. Notice that we can define the operators $P_\pm=\frac{1}{2}(1\pm\Omega)$, which satisfy

$$P_\pm^2=P_\pm,\;\; P_\pm\cdot P_\mp=0\,, \tag{A.10}$$

thus being projector operators on the q-basis space. Notice that the projector P_+ is involved in condition (A.9), which means that only some of the components of $\vec{q}_a$ are relevant for tadpoles. On the other hand, the projector P_- is involved on the coupling to branes of the RR B_2 fields that mediate the generalized GS mechanism (see section 2 and [9, 12] for a proper definition of these fields). Indeed, a general D6$_a$-brane whose vector is $\vec{q}_a$ will couple to these antisymmetric four-dimensional fields as $P_-\vec{q}_a$. Notice that since P_- is a projector, it will only couple to four fields. In the same way that was done for factorizable branes in (2.8), we can give an explicit expression for these couplings in terms of the components of the vector $\vec{q}$ of this brane.

$$\begin{aligned}
B_2^0 &: N_a\left(b^1b^2b^3q_1+q_2+b^1q_3+b^2b^3q_4+b^2q_5+b^1b^3q_6+b^3q_7+b^1b^2q_8\right)F_a\,,\\
B_2^1 &: N_a\left(b^1q_1+q_4\right)F_a\,,\\
B_2^2 &: N_a\left(b^2q_1+q_6\right)F_a\,,\\
B_2^3 &: N_a\left(b^3q_1+q_8\right)F_a\,.
\end{aligned} \tag{A.11}$$

In order to relate these expressions with the ones presented in the text, let us recall that in (2.8), we were expressing our D6-brane configurations in terms of *fractional* 1-cycles. These fractional 1-cycles are defined as [11, 15]

$$(n^i,m^i)_{\mathrm{frac}}\equiv(n^i,m^i)+b^i(0,n^i)\,, \tag{A.12}$$

brane type	N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
a	N_α	$(n_\alpha^1, 0)$	(n_α^2, m_α^2)	$(n_\alpha^3, \epsilon_\alpha m_\alpha^3)$
b	N_β	(n_β^1, m_β^1)	$(n_\beta^2, 0)$	$(n_\beta^3, \epsilon_\beta m_\beta^3)$
b or c	N_γ	(n_γ^1, m_γ^1)	$(n_\gamma^2, \epsilon_\gamma^2 m_\gamma^2)$	$(n_\gamma^3, \epsilon_\gamma^3 m_\gamma^3)$
a	N_δ	$(n_\delta^1, 0)$	(n_δ^2, m_δ^2)	$(n_\delta^3, \epsilon_\delta m_\delta^3)$

Table 19: D6-brane wrapping numbers giving rise to a generic Q-SUSY model of those presented in section 3. Here every discrete parameter n_i^j, m_i^j is taken positive, whereas the phases $\epsilon_i^j = \pm 1$ determine the type of brane ($a1$ or $a2$ etc. . .) we are dealing with.

so when computing the coupling of a fractional brane to, say, the field B_2^1, we have the coefficient

$$N_a \left(b^1 q_1 + q_4\right) = N_a \left(b^1 n^1 n^2 n^3 + m^1 n^2 n^3\right) = N_a m^1_{\text{frac}} n^2_{\text{frac}} n^3_{\text{frac}}\,, \tag{A.13}$$

so it reduces to the previous expression. Notice that, apart from this appendix, we have used the fractional notation on the whole text, without any subindex.

B. Extra U(1)'s and Q-SUSY structure

One of the most interesting aspects regarding intersecting brane world models involves the massive U(1) structure, arising from couplings with antisymmetric B_2 fields, as shown in section 2 and more extensively in [12]. In particular, we are interested in the abelian gauge symmetries that remain after those couplings have been taken into account. When dealing with Q-SUSY models of factorizable branes there are some general results that can be stated regarding such massless U(1)'s. Indeed, in this second appendix we will try to elucidate the number of massless U(1)'s in terms of the different Q-SUSY structures presented in section 3.

Let us start from a generic brane content consisting of four stacks of factorizable D6-branes, which contains each of the SUSY-quivers considered in section 3. This brane content is presented in table 19.

Notice that we must impose $m_\gamma^2 \cdot m_\gamma^3 = 0$ for the brane γ to belong to the hexagonal structure depicted in figure 4. We should also impose this condition in order to avoid chiral exotic matter appearing in the cc^* sector. This brane content will yield the most general Q-SUSY quiver of four stacks of factorizable branes, modulo renumbering of the tori. However, as we already mentioned in the text, without loss of generality we can take the branes α, β to be of type a_2, b_2, respectively. This amounts to take $\epsilon_\alpha = \epsilon_\beta = -1$ in table 19.

Given a specific brane content, we can easily compute which of the U(1)'s will remain massless in our configuration by looking at the couplings (2.8). For our purposes, it will be useful to encode such information in matrix notation. In general, for a configuration of K stacks of branes we can define the $\mathbf{B}$ as a $4 \times K$ matrix, containing on each column the coupling of the i^{th} brane to the four B_2 fields. When dealing with factorizable branes, such matrix has the form

$$\mathbf{B} = \begin{pmatrix} & N_i m_i^1 m_i^2 m_i^3 & \\ & N_i m_i^1 n_i^2 n_i^3 & \\ \ldots & N_i n_i^1 m_i^2 n_i^3 & \ldots \\ & N_i n_i^1 n_i^2 m_i^3 & \end{pmatrix}. \tag{B.1}$$

Given a general linear combination of U(1) fields, whose generator is

$$Q_X = \sum_{j=1}^{K} c_X^j Q_j \,, \tag{B.2}$$

it will remain as a massless linear combination of the low energy spectrum whenever it does not couple to any of the B_2 fields, that is, if

$$\mathbf{B} \cdot \vec{q}_X = 0 \,, \qquad \vec{q}_X^{\,t} = (\cdots \; c_X^j \; \cdots) \,. \tag{B.3}$$

So we can see that each massless combination of U(1)'s corresponds to a vector $\vec{q}_X$ belonging to the kernel of $\mathbf{B}$, now seen as a linear operator. In particular the number of massive U(1)'s on each configuration equals $Rank(\mathbf{B})$. This simple observation will help us to elucidate how many U(1)'s remain massless on each Q-SUSY configuration arising from table 19. Let us distinguish two different cases

- $m_\gamma^3 = 0$

  This choice contains both rombic and triangular quiver structures. Our $\mathbf{B}$ matrix will have the form

$$\mathbf{B} = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & N_\beta \; m_\beta^1 n_\beta^2 n_\beta^3 & N_\gamma \; m_\gamma^1 n_\gamma^2 n_\gamma^3 & 0 \\ N_\alpha \; n_\alpha^1 m_\alpha^2 n_\alpha^3 & 0 & \epsilon_\gamma^2 N_\gamma \; n_\gamma^1 m_\gamma^2 n_\gamma^3 & N_\delta \; n_\delta^1 m_\delta^2 n_\delta^3 \\ -N_\alpha \; n_\alpha^1 n_\alpha^2 m_\alpha^3 & -N_\beta \; n_\beta^1 n_\beta^2 m_\beta^3 & 0 & \epsilon_\delta N_\delta \; n_\delta^1 n_\delta^2 m_\delta^3 \end{pmatrix}, \tag{B.4}$$

  from which it can easily be seen that the rank of $\mathbf{B}$ will be always lower than 4, so at least one U(1) remains massless. Now, having a Q-SUSY structure of any sort will imply some topological restrictions,[19] such as

$$\frac{m_\alpha^3/n_\alpha^3}{m_\alpha^2/n_\alpha^2} = \frac{m_\delta^3/n_\delta^3}{m_\delta^2/n_\delta^2} \,. \tag{B.5}$$

  Indeed, if we define

$$\lambda_1 = \frac{n_\alpha^2 m_\alpha^3}{m_\alpha^2 n_\alpha^3} \,, \qquad \lambda_2 = \frac{n_\beta^1 m_\beta^3}{m_\beta^1 n_\beta^3} \,, \tag{B.6}$$

[19] Notice that, although supersymmetry between two stacks of branes always implies a geometrical condition given by the angles they form, the *ability* for a full configuration of branes to be Q-supersymmetric does imply some topological restrictions.

then it can easily be seen that having a Q-SUSY structure implies $\mathbf{B}$ taking the form

$$\mathbf{B} = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & N_\beta m_\beta^1 n_\beta^2 n_\beta^3 & N_\gamma m_\gamma^1 n_\gamma^2 n_\gamma^3 & 0 \\ N_\alpha n_\alpha^1 m_\alpha^2 n_\alpha^3 & 0 & \epsilon_\gamma^2 (\lambda_2/\lambda_1) N_\gamma m_\gamma^1 n_\gamma^2 n_\gamma^3 & N_\delta n_\delta^1 m_\delta^2 n_\delta^3 \\ -\lambda_1 N_\alpha n_\alpha^1 m_\alpha^2 n_\alpha^3 & -\lambda_2 N_\beta m_\beta^1 n_\beta^2 n_\beta^3 & 0 & \epsilon_\delta \lambda_1 N_\delta n_\delta^1 m_\delta^2 n_\delta^3 \end{pmatrix}, \tag{B.7}$$

The minor determinants of this matrix will be proportional to

$$\begin{array}{rl} & \det_{(1,2,3)} \propto 1 + \epsilon_\gamma^2 \\ \det_{(1,2,4)}\,, & \det_{(1,3,4)} \propto 1 + \epsilon_\delta \\ & \det_{(2,3,4)} \propto \epsilon_\gamma^2 \epsilon_\delta - 1 \end{array} \tag{B.8}$$

For a triangular quiver we should impose $\epsilon_\gamma^2 = \epsilon_\delta = -1$, so every minor determinant will vanish and the rank of $\mathbf{B}$ will be two. Thus, we will have precisely 2 massless surviving U(1)'s, at least if none of the entries in (B.4) vanishes. In any case, when trying to build either Standard or Left-Right symmetric models, one extra abelian group will arise. When dealing with a rombic quiver, however, we should impose instead $\epsilon_\gamma^2 = -\epsilon_\delta = -1$, and as a result there will be two nonvanishing minor determinants. Thus, we will have just one massless U(1). The other choices of phases correspond to some other SUSY-quivers not considered in this paper.

- $m_\gamma^2 = 0$

  This second choice will contain square and linear Q-SUSY structures. Proceeding in the same manner as done above, we find that in order to have a Q-SUSY structure our $\mathbf{B}$ matrix should take the form

$$\mathbf{B} = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & N_\beta m_\beta^1 n_\beta^2 n_\beta^3 & N_\gamma m_\gamma^1 n_\gamma^2 n_\gamma^3 & 0 \\ N_\alpha n_\alpha^1 m_\alpha^2 n_\alpha^3 & 0 & 0 & N_\delta n_\delta^1 m_\delta^2 n_\delta^3 \\ -\lambda_1 N_\alpha n_\alpha^1 m_\alpha^2 n_\alpha^3 & -\lambda_2 N_\beta m_\beta^1 n_\beta^2 n_\beta^3 & \epsilon_\gamma^3 \lambda_2 N_\gamma m_\gamma^1 n_\gamma^2 n_\gamma^3 & \epsilon_\delta \lambda_1 N_\delta n_\delta^1 m_\delta^2 n_\delta^3 \end{pmatrix}, \tag{B.9}$$

  the minor determinants now being proportional to

$$\begin{array}{l} \det_{(1,2,3)}\,, \quad \det_{(2,3,4)} \propto 1 + \epsilon_\gamma^3 \\ \det_{(1,2,4)}\,, \quad \det_{(1,3,4)} \propto 1 + \epsilon_\delta \end{array} \tag{B.10}$$

  The square quiver case amounts to take $\epsilon_\gamma^3 = \epsilon_\delta = 1$, which implies $Rank(\mathbf{B}) = 3$ and just one massless U(1). For the Linear quiver, in turn, we must take $\epsilon_\gamma^3 = -\epsilon_\delta = 1$, again with the same result.

Apart from these general considerations, let us notice that, when trying to get the Standard Model from a general orientifold configuration whith the brane content of table 19, the hypercharge generator will have an associated vector of the form

$$\vec{q}_Y = \begin{pmatrix} 1/6 \\ 0 \\ \epsilon/2 \\ \tilde{\epsilon}/2 \end{pmatrix}, \tag{B.11}$$

where ϵ, $\tilde{\epsilon}$ are model-dependent phases. For this vector to belong to the kernel of the general matrix in (B.4), we must impose, among others, the condition $m_\gamma^1 n_\gamma^2 n_\gamma^3 = 0$. Now, notice that $m_\gamma^1 \neq 0$, or else there will be no intersection with branes α and δ, so we are finally led to consider $n_\gamma^2 = 0$ or $n_\gamma^3 = 0$. Since this will imply that the brane γ has a twist vector

$$v_\gamma = \left(\theta_\gamma, \pm\frac{\pi}{2}, 0\right) \quad \text{or} \quad v_\gamma = \left(\theta_\gamma, 0, \pm\frac{\pi}{2}\right), \tag{B.12}$$

then in order to belong to one of the types of branes in (3.2), we must impose $\theta_\gamma = \frac{\pi}{2}$, which in turn implies that $n_c^1 = 0$. As a result, the brane γ will not couple to any of the B_2 fields in (2.8), as can be seen by direct substitution in (B.1). The **B** matrix will be effectively reduced to the 3×3 minor $(1, 2, 4)$, and its rank will depend exclusively on ϵ_δ. Indeed, when $\epsilon_\delta = -1$ we find that there are two massless U(1)'s whose generators are

$$\begin{aligned} &Q_\gamma \,, \\ &(N_\delta \; n_\delta^1 m_\delta^2 n_\delta^3) \; Q_\alpha - (N_\alpha \; n_\alpha^1 m_\alpha^2 n_\alpha^3) \; Q_\delta \,, \end{aligned} \tag{B.13}$$

whereas in the case $\epsilon_\delta = 1$ only Q_γ remains massless.

References

[1] M.R. Douglas and G.W. Moore, *D-branes, quivers and ALE instantons*, hep-th/9603167; M.R. Douglas, B.R. Greene and D.R. Morrison, *Orbifold resolution by D-branes*, *Nucl. Phys.* **B 506** (1997) 84 [hep-th/9704151].

[2] G. Aldazábal, L.E. Ibáñez, F. Quevedo and A.M. Uranga, *D-branes at singularities: a bottom-up approach to the string embedding of the standard model*, *J. High Energy Phys.* **08** (2000) 002 [hep-th/0005067].

[3] G. Aldazábal, L.E. Ibáñez and F. Quevedo, *Standard-like models with broken supersymmetry from type I string vacua*, *J. High Energy Phys.* **01** (2000) 031 [hep-th/9909172]; *A D-brane alternative to the MSSM*, *J. High Energy Phys.* **02** (2000) 015 [hep-ph/0001083]; M. Cvetič, A.M. Uranga and J. Wang, *Discrete wilson lines in $N = 1$ $D = 4$ type IIB orientifolds: a systematic exploration for z(6) orientifold*, *Nucl. Phys.* **B 595** (2001) 63 [hep-th/0010091].

[4] D. Berenstein, V. Jejjala and R.G. Leigh, *The standard model on a D-brane*, *Phys. Rev. Lett.* **88** (2002) 071602 [hep-ph/0105042].

[5] D. Bailin, G.V. Kraniotis and A. Love, *Searching for string theories of the standard model*, hep-th/0108127; L.L. Everett, G.L. Kane, S.F. King, S. Rigolin and L.-T. Wang, *Supersymmetric Pati-Salam models from intersecting D-branes*, *Phys. Lett.* **B 531** (2002) 263 [hep-ph/0202100].

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[6] L.F. Alday and G. Aldazábal, *In quest of 'just' the standard model on D-branes at a singularity*, *J. High Energy Phys.* **05** (2002) 022 [hep-th/0203129].

[7] M. Berkooz, M.R. Douglas and R.G. Leigh, *Branes intersecting at angles*, *Nucl. Phys.* **B 480** (1996) 265 [hep-th/9606139];
H. Arfaei and M.M. Sheikh Jabbari, *Different D-brane interactions*, *Phys. Lett.* **B 394** (1997) 288 [hep-th/9608167];
V. Balasubramanian and R.G. Leigh, *D-branes, moduli and supersymmetry*, *Phys. Rev.* **D 55** (1997) 6415 [hep-th/9611165];
M.M. Sheikh Jabbari, *Classification of different branes at angles*, *Phys. Lett.* **B 420** (1998) 279 [hep-th/9710121].

[8] R. Blumenhagen, L. Görlich, B. Körs and D. Lüst, *Noncommutative compactifications of type I strings on tori with magnetic background flux*, *J. High Energy Phys.* **10** (2000) 006 [hep-th/0007024]; *Magnetic flux in toroidal type I compactification*, *Fortsch. Phys.* **49** (2001) 591 [hep-th/0010198].

[9] G. Aldazábal, S. Franco, L.E. Ibáñez, R. Rabadán and A.M. Uranga, *$D = 4$ chiral string compactifications from intersecting branes*, *J. Math. Phys.* **42** (2001) 3103 [hep-th/0011073].

[10] G. Aldazábal, S. Franco, L.E. Ibáñez, R. Rabadán and A.M. Uranga, *Intersecting brane worlds*, *J. High Energy Phys.* **02** (2001) 047 [hep-ph/0011132].

[11] R. Blumenhagen, B. Körs and D. Lüst, *Type I strings with F- and B-flux*, *J. High Energy Phys.* **02** (2001) 030 [hep-th/0012156].

[12] L.E. Ibáñez, F. Marchesano and R. Rabadán, *Getting just the standard model at intersecting branes*, *J. High Energy Phys.* **11** (2001) 002 [hep-th/0105155];
L.E. Ibáñez, *Standard model engineering with intersecting branes*, hep-ph/0109082.

[13] R. Blumenhagen, B. Körs, D. Lüst and T. Ott, *The standard model from stable intersecting brane world orbifolds*, *Nucl. Phys.* **B 616** (2001) 3 [hep-th/0107138]; *Intersecting brane worlds on tori and orbifolds*, hep-th/0112015.

[14] M. Cvetič, G. Shiu and A.M. Uranga, *Three-family supersymmetric standard like models from intersecting brane worlds*, *Phys. Rev. Lett.* **87** (2001) 201801 [hep-th/0107143]; *Chiral four-dimensional $N = 1$ supersymmetric type IIA orientifolds from intersecting D6-branes*, *Nucl. Phys.* **B 615** (2001) 3 [hep-th/0107166]; *Chiral type II orientifold constructions as M theory on G_2 holonomy spaces*, hep-th/0111179.

[15] D. Cremades, L.E. Ibáñez and F. Marchesano, *SUSY quivers, intersecting branes and the modest hierarchy problem*, *J. High Energy Phys.* **07** (2002) 009 [hep-th/0201205].

[16] D. Bailin, G.V. Kraniotis and A. Love, *Standard-like models from intersecting D4-branes*, *Phys. Lett.* **B 530** (2002) 202 [hep-th/0108131];
H. Kataoka and M. Shimojo, $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ *chiral models from intersecting D4/D5 branes*, hep-th/0112247;
G. Honecker, *Intersecting brane world models from D8-branes on $(T^2 \times T^4/Z_3)/\Omega\mathcal{R}_1$ type IIA orientifolds*, *J. High Energy Phys.* **01** (2002) 025 [hep-th/0201037].

[17] R. Blumenhagen, B. Körs and D. Lüst, *Moduli stabilization for intersecting brane worlds in type 0' string theory*, *Phys. Lett.* **B 532** (2002) 141 [hep-th/0202024].

[18] J. García-Bellido, R. Rabadán and F. Zamora, *Inflationary scenarios from branes at angles*, *J. High Energy Phys.* **01** (2002) 036 [hep-th/0112147];
R. Blumenhagen, B. Körs, D. Lüst and T. Ott, *Hybrid inflation in intersecting brane worlds*, hep-th/0202124.

[19] R. Blumenhagen, L. Görlich and B. Körs, *Supersymmetric 4D orientifolds of type IIA with D6-branes at angles*, *J. High Energy Phys.* **01** (2000) 040 [hep-th/9912204]; *A new class of supersymmetric orientifolds with D-branes at angles*, hep-th/0002146;
S. Förste, G. Honecker and R. Schreyer, *Supersymmetric $Z_N \times Z_M$ orientifolds in 4D with D-branes at angles*, *Nucl. Phys.* **B 593** (2001) 127 [hep-th/0008250]; *Orientifolds with branes at angles*, *J. High Energy Phys.* **06** (2001) 004 [hep-th/0105208];
G. Honecker, *Non-supersymmetric orientifolds with D-branes at angles*, hep-th/0112174.

[20] C. Bachas, *A way to break supersymmetry*, hep-th/9503030.

[21] M. Bianchi, G. Pradisi and A. Sagnotti, *Toroidal compactification and symmetry breaking in open string theories*, *Nucl. Phys.* **B 376** (1992) 365;
M. Bianchi, *A note on toroidal compactifications of the type I superstring and other superstring vacuum configurations with 16 supercharges*, *Nucl. Phys.* **B 528** (1998) 73 [hep-th/9711201];
E. Witten, *Toroidal compactification without vector structure*, *J. High Energy Phys.* **02** (1998) 006 [hep-th/9712028];
C. Angelantonj, *Comments on open-string orbifolds with a non-vanishing B_{ab}*, *Nucl. Phys.* **B 566** (2000) 126 [hep-th/9908064];
C. Angelantonj and R. Blumenhagen, *Discrete deformations in type I vacua*, *Phys. Lett.* **B 473** (2000) 86 [hep-th/9911190];
Z. Kakushadze, *Geometry of orientifolds with NS-NS B-flux*, *Int. J. Mod. Phys.* **A 15** (2000) 3113 [hep-th/0001212].

[22] R. Blumenhagen, L. Görlich, B. Körs and D. Lüst, *Asymmetric orbifolds, noncommutative geometry and type I string vacua*, *Nucl. Phys.* **B 582** (2000) 44 [hep-th/0003024];
C. Angelantonj, I. Antoniadis, E. Dudas and A. Sagnotti, *Type I strings on magnetised orbifolds and brane transmutation*, *Phys. Lett.* **B 489** (2000) 223 [hep-th/0007090];
C. Angelantonj and A. Sagnotti, *Type I vacua and brane transmutation*, hep-th/0010279.

[23] A.M. Uranga, *D-brane, fluxes and chirality*, *J. High Energy Phys.* **04** (2002) 016 [hep-th/0201221].

[24] K. Dasgupta, G. Rajesh and S. Sethi, *M theory, orientifolds and G-flux*, *J. High Energy Phys.* **08** (1999) 023 [hep-th/9908088];
S.B. Giddings, S. Kachru and J. Polchinski, *Hierarchies from fluxes in string compactifications*, hep-th/0105097;
S. Kachru, M.B. Schulz and S. Trivedi, *Moduli stabilization from fluxes in a simple IIB orientifold*, hep-th/0201028.

[25] S. Kachru and J. McGreevy, *Supersymmetric three-cycles and (super)symmetry breaking*, *Phys. Rev.* **D 61** (2000) 026001 [hep-th/9908135];
M. Mihailescu, I.Y. Park and T.A. Tran, *D-branes as solitons of an $N = 1$, $D = 10$ non-commutative gauge theory*, *Phys. Rev.* **D 64** (2001) 046006 [hep-th/0011079];
E. Witten, *BPS bound states of D0-D6 and D0-D8 systems in a B-field*, *J. High Energy Phys.* **04** (2002) 012 [hep-th/0012054];
R. Blumenhagen, V. Braun and R. Helling, *Bound states of D(2p)-D0 systems and supersymmetric p-cycles*, *Phys. Lett.* **B 510** (2001) 311 [hep-th/0012157].

[26] M. Dine, N. Seiberg and E. Witten, *Fayet-Iliopoulos terms in string theory*, *Nucl. Phys.* **B 289** (1987) 589;
J.J. Atick, L.J. Dixon and A. Sen, *String calculation of Fayet-Iliopoulos D terms in arbitrary supersymmetric compactifications*, *Nucl. Phys.* **B 292** (1987) 109;
M. Dine, I. Ichinose and N. Seiberg, *F terms and D terms in string theory*, *Nucl. Phys.* **B 293** (1987) 253.

[27] L.E. Ibáñez and G.G. Ross, $\mathrm{SU}(2) \times \mathrm{U}(1)$ *symmetry breaking as a radiative effect of supersymmetry breaking in GUT's*, *Phys. Lett.* **B 110** (1982) 215.

[28] J.D. Lykken, *Weak scale superstrings*, *Phys. Rev.* **D 54** (1996) 3693 [hep-th/9603133];
N. Arkani-Hamed, S. Dimopoulos and G.R. Dvali, *The hierarchy problem and new dimensions at a millimeter*, *Phys. Lett.* **B 429** (1998) 263 [hep-ph/9803315];
I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G.R. Dvali, *New dimensions at a millimeter to a Fermi and superstrings at a TeV*, *Phys. Lett.* **B 436** (1998) 257 [hep-ph/9804398];
G. Shiu and S.H.H. Tye, *TeV scale superstring and extra dimensions*, *Phys. Rev.* **D 58** (1998) 106007 [hep-th/9805157].

[29] M. Nakahara, *Geometry, topology and physics*, Graduate student series in physics, Hilger, Bristol, UK 1990.

[30] R. Rabadán, *Branes at angles, torons, stability and supersymmetry*, *Nucl. Phys.* **B 620** (2002) 152 [hep-th/0107036].

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