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A D-brane alternative to the MSSM

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Abstract: The success of SU(5)-like gauge coupling unification boundary conditions $g_3^2 = g_2^2 = 5/3g_1^2$ has biased most attempts to embed the SM interactions into a unified structure. After discussing the limitations of the orthodox approach, we propose an alternative that appears to be quite naturally implied by recent developments based on D-brane physics. In this new alternative: 1) The gauge group, above a scale of order 1 TeV, is the minimal left-right symmetric extension $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ of the SM; 2) Quarks, leptons and Higgs fields come in three generations; 3) Couplings unify at an intermediate string scale $M_s = 9 \times 10^{11}$ GeV with boundary conditions $g_3^2 = g_L^2 = g_R^2 = 32/3\ g_{B-L}^2$. This corresponds to the natural embedding of gauge interactions into D-branes and is different from the standard SO(10) embedding which corresponds to $k_{B-L} = 8/3$. Unification only works in the case of three generations; 4) Proton stability is automatic due to the presence of $\mathbb{Z}_2$ discrete R-parity and lepton parities. A specific type-IIB string orientifold model with the above characteristics is constructed. The existence of three generations is directly related to the existence of three complex extra dimensions. In this model the string scale can be identified with the intermediate scale and SUSY is broken also at that scale due to the presence of anti-branes in the vacuum. We discuss a number of phenomenological issues in this model including Yukawa couplings and a built-in axion solution to the strong-CP problem. The present framework could be tested by future accelerators by finding the left-right symmetric extension of the SM at a scale of order 1 TeV.

Keywords: D-branes, Superstring Vacua, Supersymmetric Standard Model.

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Contents

1. Introduction

The success of coupling unification extrapolations based on the massless spectrum of the MSSM has greatly conditioned the search for a realistic string vacuum. This search has been shaped, to a great extent, by the fact that SM couplings seem to join at a scale $M_X = 2 \times 10^{16}$ GeV, not far from the Planck mass M_p and by the necessity of identifying the string scale M_s essentially with the Planck scale in heterotic model building.

In this situation, it appears natural to look for perturbative heterotic vacua in which, below a scale of order M_X, essentially only the MSSM remains.

Recent p-brane developments have changed our view of the possible ways to embed SM physics into string theory. To start with, it has been realized that the string/M-theory scale may be much below the Planck and unification scales [1]–[12]. This is because in the presence of p-branes (like D-branes in type-II and type-I string theory) gauge interactions can be localized in the world-volume of D-branes (e.g., a 3-brane), whereas gravitational interactions in general live in the full ten (or eleven) dimensions. Then the largeness of the Planck mass may be obtained even if $M_s \ll M_p$ if there are large compact dimensions.

Now, if $M_s \ll M_X$, gauge coupling unification should in principle take place at the string scale M_s and thus the nice unification of MSSM couplings is lost. Of course, this unification problem appearing for string models if $M_s \ll M_X$, could be taken as an argument against them. However, we think that we should first try to answer the following question: Is there any simple alternative framework which is consistent with unification at a string scale $M_s \ll M_X$? After all, the MSSM+big desert orthodoxy is not free of problems. In fact, some unattractive features of the standard scenario are the following:

(i) The quark/lepton generations come in three chiral copies whereas the Higgs fields come only in one copy and are non-chiral. The fact that there is a single Higgs set is crucial to obtain correct unification predictions. This asymmetry among quarks/leptons on one side and Higgs fields on the other looks quite *ad hoc*.

(ii) The MSSM needs to be supplemented by additional symmetries like R-parity in order to ensure proton stability from dimension four operators. It also needs additional symmetries beyond R-parity to get stability against dimension five operators.

(iii) To obtain a viable heterotic string unification, whose massless sector is just the MSSM and includes the above symmetries, turns out to be a very difficult task, if not impossible. All the models studied up to now require a complicated study of possible scalar flat directions and only very particular ones lead to something of that sort [13]. The reason why dynamics should prefer such vacua with only one set of Higgsses and built-in discrete symmetries to suppress too fast proton decay is unclear.

Given the above limitations of the orthodox approach, it seems sensible to look for (if possible, elegant) alternatives. But to be really competitive with the MSSM scenario such alternatives need to 1) improve some of the above problematic aspects of the standard scenario and 2) have a nice and consistent unification of coupling constants. By the latter we mean that couplings unify at the string scale without forcing the structure of the model by adding, for instance, unjustified extra mass scales or ad-hoc extra massless particles.

In the present paper we want to propose such an alternative to the standard MSSM+big desert scenario. We propose that above a scale of order 1 TeV or so, the gauge group is that of the minimal left-right symmetric extension of the supersymmetric standard model, i.e. $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [14, 15, 16]. In addition all quarks, leptons and Higgs fields come in three generations. Interestingly enough, this simple structure leads to very precise unification of gauge coupling constants at an intermediate scale of order 10^{12} GeV as long as the normalization of the

$U(1)_{B-L}$ coupling is the one expected if such gauge group is associated to a collection of D-branes. It is important to remark that this normalization differs from the one predicted by standard GUT left-right symmetric scenarios like SO(10).

We also construct an explicit type-IIB orientifold string compactification leading to the desired massless spectrum and normalization of coupling constants. In this model there is a simple explanation for the family replication: there are three quark-lepton generations because there are three complex compact dimensions and an underlying $\mathbb{Z}_3$ orbifold. These features resemble the first three-generation perturbative heterotic string models built, those of ref. [17]. Furthermore, this model has the interesting property of having natural discrete symmetries including R-parity and Z_2 lepton numbers, thus guaranteeing in a natural way the stability of the proton. Although other baryon number violating processes, such as neutron-antineutron oscillations, are allowed, they are sufficiantly suppressed to be consistent with current experimental bounds. Finally, the structure of the model allows for candidate axion fields with the right couplings to gauge fields needed to solve the strong CP problem.

The structure of this paper is as follows. In section 2 we present this alternative scenario which we call D-brane left-right symmetric model. We also discuss the unification of coupling constants and show how, if this model is correct, new Z' and W' gauge bosons corresponding to left-right symmetry should be found at future or present colliders. In section 3 we present a particular type-IIB orientifold model realizing the above scenario and study the cancellation of U(1) anomalies and generation of Fayet-Iliopoulos terms. In this realization the unification scale is identified with a string scale $M_s \propto 10^{12}$ GeV. The model has also some anti-branes in the bulk which provide for hidden-sector supersymmetry breaking at the same scale of order M_s. We study a number of phenomenological issues of this particular orientifold model in section 4. This includes some aspects of the structure of Yukawa couplings, $SU(2)_R \times U(1)_{B-L}$ symmetry breaking and the presence of natural candidates for invisible axions. In section 5 we present an outlook and some final comments.

2. The D-brane left-right symmetric model and coupling unification

As we discussed above, it has become recently clear that the string scale could well be much below the Planck mass. But, is there any indication or advantage from a lowered string scale? A particularly interesting alternative to the unification at M_X close to the Planck scale is getting unification close to the geometric intermediate scale $M_I = \sqrt{M_W M_p}$. Indeed, if the string scale is of order $M_s \propto M_I$, gauge couplings should unify at that scale. Now, as argued in ref. [10] , if there are non supersymmetric brane configurations, the scale of supersymmetry breaking would also be of order of $M_s = M_I$. This is interesting because hidden sector supersymmetry breaking models also need to have SUSY-breaking at the intermediate scale. Thus in this case the string, unification and SUSY-braking scales would be one and the same.

We would like to argue in what follows that the intermediate scale idea is equally good than the standard one in what concerns coupling unification, at least for a model with the following structure:

(i) The gauge group above a L-R symmetric scale M_R slightly above the weak scale is the minimal left-right symmetric extension of the SM: $\mathrm{SU}(3)\times\mathrm{SU}(2)_L\times \mathrm{SU}(2)_R\times \mathrm{U}(1)_{B-L}$.

(ii) All quarks, leptons and Higgs fields come in three generations. Thus the chiral multiplet content is three copies of $(3,2,1,1/3)+(\bar{3},1,2,-1/3)+(1,2,1,-1)+(1,1,2,+1)+(1,2,2,0)$.

(iii) The boundary conditions at the unification (i.e. string) scale are $g_3^2 = g_L^2 = g_R^2 = 32/3g_{B-L}^2$. This corresponds to a weak angle with $\sin^2\theta(M_s) = 3/14 = 0.215$.

In a model with the above characteristics one finds that gauge couplings naturally unify at a scale of order the intermediate scale $M_s \propto 10^{12}$ GeV as long as the left-right scale M_R is not far from the weak scale M_W. An important point to remark is that the unification boundary conditions are different from those found in GUT schemes. Indeed in SO(10)-like schemes the boundary conditions at unification are $g_3^2 = g_L^2 = g_R^2 = 8/3g_{B-L}^2$ yielding the canonical $\sin^2\theta_W = 3/8$. A remarkable point we find is that *the new boundary conditions we are proposing are precisely the ones which are natural from the point of view of the embedding of the gauge group in a D-brane scheme.*

Let us discuss in some more detail how coupling unification takes place. The above mentioned boundary conditions $g_3^2 = g_L^2 = g_R^2 = 32/3g_{B-L}^2$ have a simple group theoretical interpretation. They correspond to the embedding of the left-right symmetric gauge group into a non-semisimple structure:

$$\mathrm{U}(3)\times\mathrm{U}(2)_L\times\mathrm{U}(2)_R\,, \tag{2.1}$$

with unified coupling constants $g_3^2 = g_L^2 = g_R^2$ at some mass scale (to be identified later on with the string scale M_s). This is in fact the structure one gets in models with gauge groups living on D-branes, as we will discuss in the specific string model below. As we said, the model contain three identical generations under $\mathrm{U}(3)\times\mathrm{U}(2)_L\times\mathrm{U}(2)_R$ with quantum numbers $(3,\bar{2},1)_{(1,-1,0)}+(\bar{3},1,2)_{(-1,0,1)}+(1,\bar{2},1)_{(0,-1,0)}+(1,1,2)_{(0,0,1)}$ $+(1,2,\bar{2})_{(0,1,-1)}$, where the subindices denote the charges with respect to the three U(1)'s. We denote the U(1) generators by Q_3, Q_L and Q_R respectively. It is easy to check that two of them are anomalous and only one of them, the linear combination

$$Q_{B-L} = -\frac{2}{3}\,Q_3 - Q_L - Q_R \tag{2.2}$$

is anomaly free.[1] This is just the familiar $(B-L)$ of left-right symmetric models which is related to weak hypercharge by $Y = -T_R^3 + Q_{B-L}/2$. Now, notice that, if we normalize the original $\mathrm{U}(n)$ generators in the fundamental representation T_a to $\mathrm{Tr}\, T_a^2 = 1/2$, the normalization of the U(1)'s are $\mathrm{Tr}\, Q_3^2 = 3$, $\mathrm{Tr}\, Q_L^2 = \mathrm{Tr}\, Q_R^2 = 2$. Then, the normalization of $\mathrm{U}(1)_{B-L}$ compared to that of the non-abelian generators is $k_{B-L} = 2\,\mathrm{Tr}\, Q_{B-L}^2 = 32/3$, as remarked above. Notice this implies a hypercharge normalization $k_1 = k_R + 1/4 k_{B-L} = 11/3$, and hence a tree level weak angle $\sin^2\theta_W = 3/14 = 0.214$.

Let us study now the one-loop corrections to the couplings. In between the scales M_R and M_s the gauge group is $\mathrm{SU}(3)\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{U}(1)_{B-L}$ and the above chiral field content gives rise to the following one-loop β-function coefficients B_a:

$$B_3 = -3\,, \qquad B_L = +3\,, \qquad B_R = +3\,, \qquad B_{B-L} = +16\,. \tag{2.3}$$

In between the weak scale M_W and M_R the gauge group will be that of the SM with β-function coefficients b_i. Then the one loop running yields:

$$\begin{aligned}
\sin^2\theta_W(M_Z) &= \frac{1}{1+k_1}\left\{1 + k_1\frac{\alpha_e(M_Z)}{2\pi}\left[\left(B_L - \frac{1}{k_1}B_1'\right)\log\left(\frac{M_s}{M_R}\right) + \right.\right.\\
&\qquad\qquad \left.\left. + \left(b_2 - \frac{1}{k_1}b_1\right)\log\left(\frac{M_R}{M_Z}\right)\right]\right\}, \\
\frac{1}{\alpha_e(M_Z)} - \frac{1+k_1}{\alpha_3(M_Z)} &= \frac{1}{2\pi}\left[(b_1 + b_2 - (1+k_1)b_3)\log\left(\frac{M_R}{M_Z}\right) + \right.\\
&\qquad \left. + (B_1' + B_L - (1+k_1)B_3)\log\left(\frac{M_s}{M_R}\right)\right],
\end{aligned} \tag{2.4}$$

where one defines

$$B_1' = B_R + \frac{1}{4}B_{B-L} \tag{2.5}$$

and $k_1 = k_R + k_{B-L}/4$. With the minimal particle content described above one has $B_1' = 7$. Let us suppose for the moment that, below the M_R scale down to M_Z, we were left just with the content of the MSSM. We would then have (for $k_1 = 11/3$, corresponding to $k_{B-L} = 32/3$):

$$\begin{aligned}
\sin^2\theta_W(M_Z) &= \frac{3}{14}\left(1 + \frac{\alpha_e(M_Z)}{2\pi}\left[4\log\left(\frac{M_s}{M_R}\right) - \frac{22}{3}\log\left(\frac{M_R}{M_Z}\right)\right]\right), \\
\frac{1}{\alpha_e(M_Z)} - \frac{14}{3\alpha_3(M_Z)} &= \frac{1}{2\pi}\left[26\log\left(\frac{M_R}{M_Z}\right) + 24\log\left(\frac{M_s}{M_R}\right)\right].
\end{aligned} \tag{2.6}$$

[1] In string theory the other two (anomalous) U(1)'s become massive and decouple due to a generalized Green-Schwarz mechanism. See the discussion in section 3.

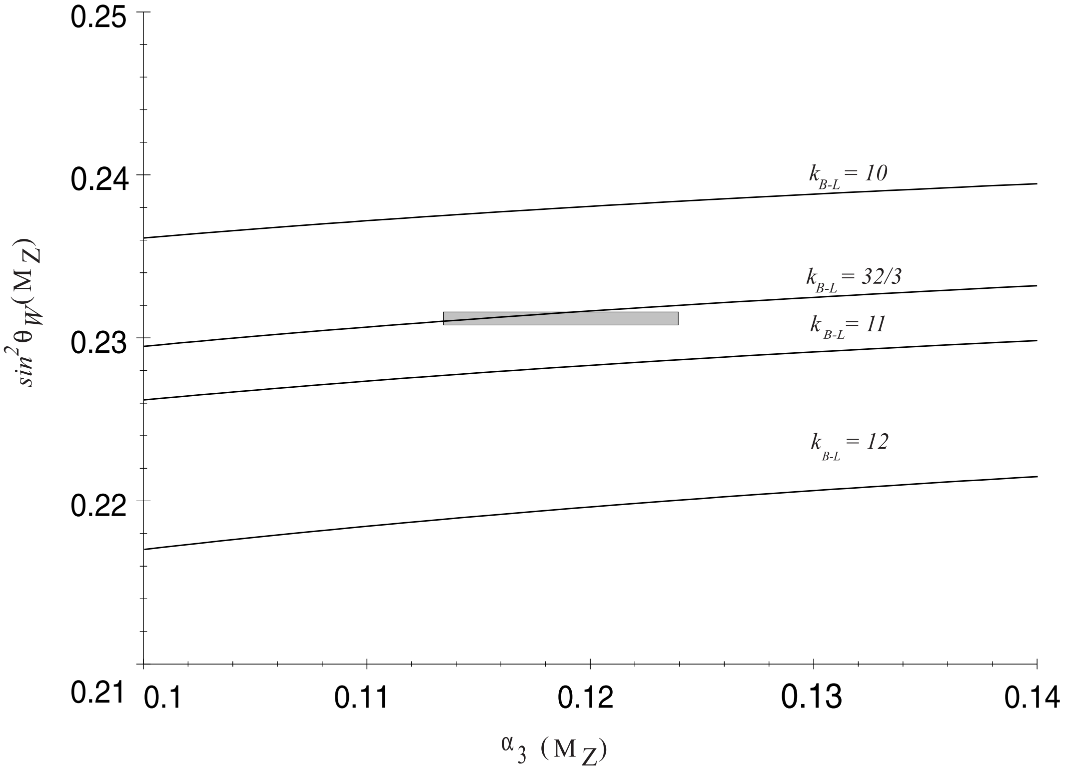

Figure 1: $\sin^2\theta_W(M_Z)$ vs. $\alpha_3(M_Z)$ for different values of k_{B-L}.

Now, using as input $\alpha_e(M_Z)^{-1} = 127.934 \pm 0.027$ in ref. [18] (in particular its central value), one can plot the predicted $\sin^2\theta_W(M_Z)$ versus $\alpha_3(M_Z)$ and compare it to the experimental data for those two quantities. This is done for several values of k_{B-L} in figure 1, where the data plotted correspond to the world-average in ref. [18] with two standard deviation errors. It may be observed that, for the particular normalization corresponding to an embedding of the left-right symmetric interactions into D-branes ($k_{B-L} = 32/3$), a very nice agreement with the data is found. Very slight departures from this value are ruled out. The unification scale corresponding to the successful results is of order $M_s = 9 \times 10^{11}$ GeV. We show for comparison a similar plot for the MSSM obtained from the one-loop formulae in figure 2. Again, for the MSSM standard normalization $k_1 = 5/3$ the prediction nicely goes through the data points, in this case for a unification mass of order $M_X = 2 \times 10^{16}$ GeV. Thus we may conclude that, within the approximations made, the D-brane left-right symmetric model is remarkably successful in obtaining appropriate gauge coupling unification.

It can also be checked that, within this scheme, both the SUSY-breaking scales and the left-right scale M_R cannot be much above the 1 TeV scale. Indeed, let us now use also as input the world average central value for the weak angle, $\sin^2\theta_W(M_Z) = 0.23117 \pm 0.00016$ as well as $\alpha_e(M_Z)$ and let us plot the predicted $\alpha_3(M_Z)$ as a function of the left-right scale M_R for various values of the SUSY-breaking mass M_{sb}.

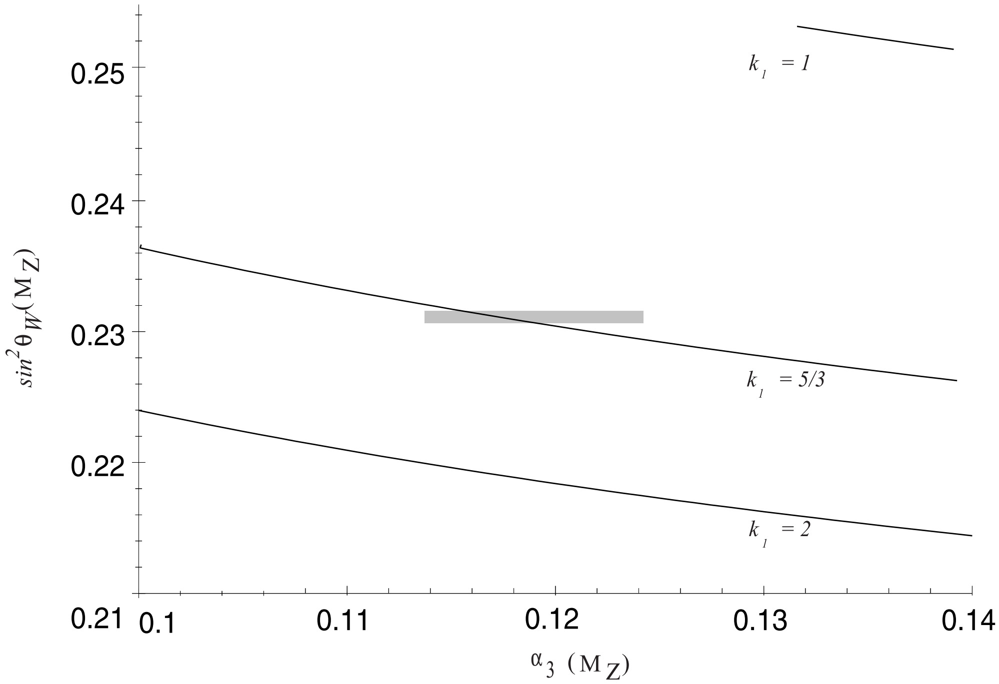

Figure 2: $\sin^2\theta_W(M_Z)$ vs. $\alpha_3(M_Z)$ for different values of k_1 for the MSSM.

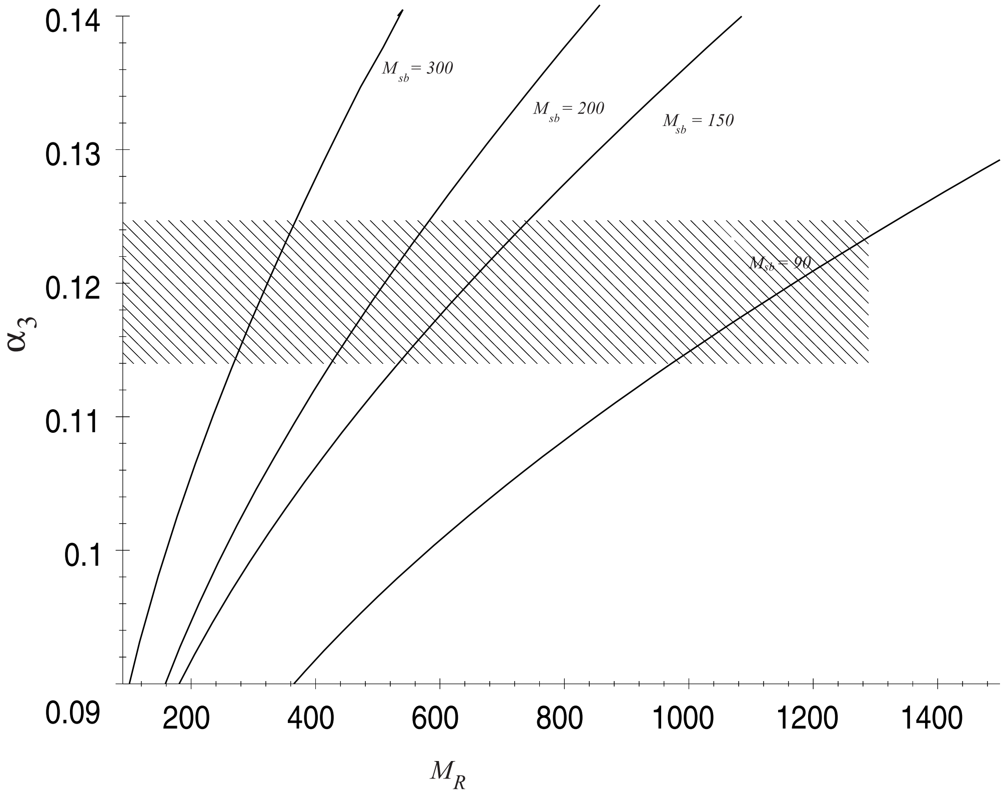

Figure 3: $\alpha_3(M_Z)$ vs. M_R for different values of a universal SUSY-threshold M_{sb}. Here LR symmetry is broken above supersymmetry.

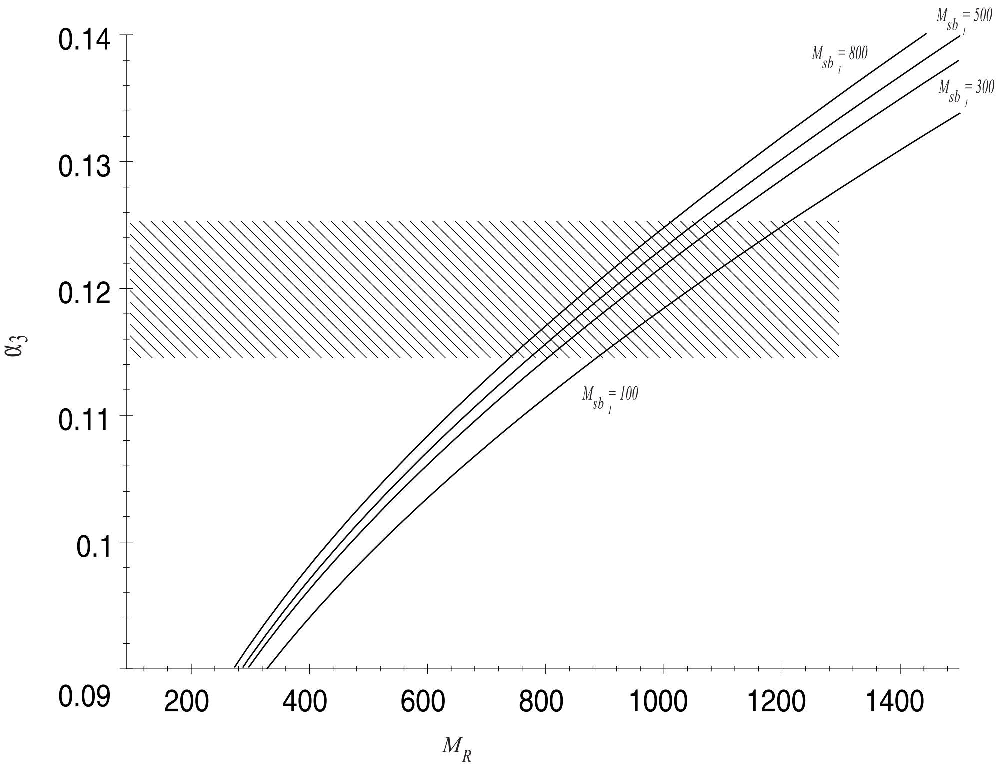

Figure 4: $\alpha_3(M_Z)$ vs. M_R for a non-colored SUSY particle threshold $M_{sb2} = 100$ and different values of a coloured SUSY threshold M_{sb1}.

This is shown in figures 3 and 4, where we have assumed that below the M_R (and above the M_{sb}) scale one is left with the particle content of the MSSM. In figure 3 we consider a universal SUSY-breaking threshold M_{sb} below which the non-SUSY SM is obtained. Figure 4 plots the same quantities but now with two SUSY-breaking thresholds one (M_{sb1}) for the coloured SUSY-particles (squarks and gluinos) and a second lower one M_{sb2} for non-coloured ones. We observe that for large values of M_R, a low value for the SUSY threshold seems to be required to get a consistent value for $\alpha_3(M_Z)$. Thus either the SUSY threshold or the left-right symmetric threshold (or both) should be below the 1 TeV scale. Although we have only used one-loop formulae and have taken step functions for the different thresholds we expect that those refinements will not substantially change the main conclusion that couplings nicely unify as long as the M_R scale is not much above 1 TeV.

Several comments are in order:

(i) Notice that we have not played around with the addition of extra mass scales and/or extra particles beyond the three quark/lepton/Higgs generations in order to get satisfactory unification. So this is not a mere adjustment of the model to get nice coupling unification, it appears naturally as long as M_R is not much higher than 1 TeV.

(ii) In order to obtain the above interesting unification results it is crucial to use the boundary conditions $g_3^2 = g_L^2 = g_R^2 = 32/3g_{B-L}^2$. Thus if we would have used the standard GUT conditions, with $k_{B-L} = 8/3$ instead of 32/3, we would obtain $M_s \propto 10^{13}\,\text{GeV}$ and $\sin^2\theta_W(M_Z) = 0.34$ (for input $\alpha_3(M_Z) = 0.119$ and $M_R \propto 1\,\text{TeV}$). As we have mentioned, the boundary conditions which work correspond to those expected when the gauge group is embedded inside a collection of type-IIB D-branes.

(iii) The tree level result $\sin^2\theta_W = 3/14 = 0.214$ is quite close already to the experimental result 0.231. One loop effects should then be *small and positive* (i.e. of order 8%). This contrasts with the standard MSSM/GUT case, where the tree level result $\sin^2\theta_W = 3/8 = 0.375$ departs from the experimental number. In this case the loop corrections must be *large and negative* (of order 62%). In this connection notice that the $\sin^2\theta_W$ loop corrections have two pieces in our case, one positive and proportional to $\log(M_s/M_R)$ and one negative and proportional to $\log(M_R/M_Z)$. Since a positive correction is needed in order to get agreement for $\sin^2\theta_W$, this is the hidden reason why gauge coupling unification requires in our case large $\log(M_s/M_R)$ but small $\log(M_R/M_Z)$.

(iv) Coupling constant unification works only for three generations of quarks, leptons and Higgs fields. Thus there is a connection between generation number and unification (unlike the MSSM which is insensitive, at one loop, to the number of quark-lepton generations). Let us denote the number of quark/lepton/Higgs generations as n_g. Thus, β-function coefficients become now $B_3 = -9 + 2n_g$, $B_L = B_R = -6 + 3n_g$, $B_{B-L} = (16n_g)/3$ and $B_1' = -6 + (13/3)n_g$. The combinations relevant for the running are

$$\left[B_L - \frac{3}{11}B_1'\right] = \frac{4}{11}(5n_g - 12)\,, \qquad \left[B_1' + B_L - \frac{14}{3}B_3\right] = 30 - 2n_g\,. \quad (2.7)$$

Now, notice that in order to have a positive correction for $\sin^2\theta_W$ as required, we need $n_g \geq 3$. On the other hand, for $n_g \geq 4$ either $\sin^2\theta_W$ becomes too large or $\log(M_s/M_R)$ becomes too small to be compatible with unification.

(v) We have assumed that at a scale M_R the gauge group is broken to that of the SM but we have not specified what the fields, giving rise to such breaking, are. As we will show in the specific string construction below, there are simple additions to the model (e.g., from the presence of some D-branes in the bulk, which lead to non-chiral particle content) which contain the required fields for this breaking without modifying the runnings at one loop, thus preserving the interesting results described above.

The above new scheme may probably be obtained in different classes of string models involving D-branes with a gauge group $U(3) \times U(2) \times U(2)$. One of the most interesting points we find is that to construct a specific four-dimensional type-I string model with precisely that massless spectrum and with the necessary gauge coupling boundary conditions is very simple. In fact, such a model was briefly discussed in ref. [19, section 4.1]. We do not claim however that this is the only possible realization of the D-brane left-right symmetric scenario here introduced. Nevertheless, we think it is worth studying such a model, since it may provide clues of more general features of the scheme. We now describe the construction of the mentioned type-IIB orientifold realization.

3. A type-IIB orientifold with left-right symmetry

3.1 The LR orientifold model

The model we are interested in is a $\mathbb{Z}_3$ type-IIB orientifold with both D-branes and anti-D-branes. A general discussion of such kind of models is given in refs. [21, 20, 19].

Here we are interested in the $\mathbb{Z}_3$ type-IIB orientifold models with both D-branes and anti-D-branes as presented in [20, 19] where we refer the reader for notation and details of the construction.

A $\mathbb{Z}_3$ type-IIB orbifold, in four dimensions [22], is obtained by dividing closed type-IIB string theory compactified on a six dimensional torus $\mathbb{T}^6$, by the discrete symmetry group $\mathbb{Z}_3$. The orientifold model is obtained by further dividing the orbifoldized string by world sheet orientation reversal symmetry [23]. The twist eigenvalues, associated to complex coordinates Y_a $a = 0, 1, 2$ are chosen as $v = \frac{1}{3}(1, 1, -2)$ in order to leave $N = 1$ supersymmetry in four dimensions.

The general picture is that the above procedure leads to a Klein-Bottle unoriented world sheet. Amplitudes computed on such a surface contain unphysical tadpole like divergences which can be interpreted as unbalanced charges carried by RR form potentials. Thus, in order to cancel such divergences, D9-branes, carrying opposite charges must be introduced. Moreover, D5-branes and anti-D5-branes can be consistently included. Even if they are not required (in this $\mathbb{Z}_3$ case) for tadpole cancellation, they open the way for achieving interesting supersymmetry breaking patterns and at the same time provide new possibilities for model building.

Let us be more explicit. An open string state is denoted by $|\Psi, ab\rangle \lambda^{pq}_{ab}$ where Ψ refers to world-sheet degrees of freedom whereas a, b are Chan-Paton indices associated to the open string endpoints lying on Dp-branes and Dq-branes respectively [22, 24, 25]. λ^{pq} is the Chan-Paton, hermitean matrix, containing the gauge group structure information. Analogously, $\lambda^{\overline{pq}}$ ($\lambda^{\overline{p}q}$) is introduced for open strings ending at D$\overline{p}$, D$\overline{q}$-antibranes (D$\overline{p}$ antibrane, Dq-brane, etc.).

The $\mathbb{Z}_3$ action (denoted by θ) that twists the internal complex coordinates has a corresponding action on Chan-Paton matrices represented by unitary matrix $\gamma_{\theta,p}$, namely $\theta : \lambda^{pq} \to \gamma_{\theta,p}\lambda^{pq}\gamma_{\theta,q}^{-1}$. Moreover, Wilson lines, wrapping along internal tori directions can also be included and also have a matrix representation when acting on Chan-Paton factors.

Consistency under group algebra operations and the requirement of cancellation of RR tadpoles leads to constraints on the possible twist matrices. Tadpole cancellation, in the $\mathbb{Z}_3$ case we are discussing, imposes the number of nine branes to be 32 and the requirement that the number of D5-branes and anti-D5-branes must be equal. Moreover, cancellation of twisted tadpoles requires

$$\mathrm{Tr}(\mathcal{W})^a\gamma_{\theta,9} + 3(\mathrm{Tr}\,\gamma_{\theta,5,a,i} - \mathrm{Tr}\,\gamma_{\theta,\bar{5},a,i}) = -4\,, \tag{3.1}$$

for $a = 0, 1, 2$.

Here we have allowed for the possibility of having a Wilson line, represented by the matrix $\mathcal{W}$ on Chan-Paton matrices, wrapping along the direction e_1. We denote with a, i with $a,\ i = 0, 1, 2$ the nine orbifold fixed points in the first and second complex planes. For a given $a,\ i = 0, 1, 2$ label the subset of fixed points that feels the twist $(\mathcal{W})^a\gamma_{\theta,9}$.

Generic solutions to these equations are discussed in [19]. Here we consider the specific model characterized by $\gamma_{\theta,9} = (\tilde{\gamma}_{\theta,9}, \tilde{\gamma}^*_{\theta,9})$, and $\mathcal{W} = (\tilde{\mathcal{W}}, \tilde{\mathcal{W}}^*)$ where $*$ denotes the complex conjugate and

$$\begin{aligned}\tilde{\gamma}_{\theta,9} &= \mathrm{diag}(\alpha I_3, \alpha^2 I_2, I_2, I_2, \alpha I_7)\\ \tilde{\mathcal{W}} &= \mathrm{diag}(\alpha I_3, \alpha I_2, \alpha I_2, I_2, I_7)\,.\end{aligned} \tag{3.2}$$

Also,

$$\gamma_{\theta,5,2,i} = \mathrm{diag}(\alpha, \alpha^2)\,, \tag{3.3}$$

where $\alpha = e^{2i\pi/3}$.

It can be easily checked that such a choice satisfies the tadpole cancellation constraint (3.1). Notice that there are two 5-branes ($\mathrm{Tr}\,\gamma_{\theta,5,2,i} = -1$) stuck at each of the three fixed points of the type $(2, i) \equiv (-1, i)$. The effective twist is thus $\mathcal{W}^{-1}\gamma_{\theta,9}$. There are no 5-branes at the other six fixed points ($\mathrm{Tr}(\gamma_{\theta,9}) = \mathrm{Tr}(\mathcal{W})\gamma_{\theta,9} = -4$). Since we must have the same number of branes and antibranes, six anti-5-branes must be present in the bulk.

It is always possible to add an extra Wilson line in the third complex plane in such a way that anti-brane sector is gauge decoupled from the 9, 5-branes sectors. In this way, anti-branes in the bulk, which lead to a non-supersymmetric spectrum, provide a "hidden sector" which will transmit supersymmetry breaking through gravitational interactions to the "observable" brane sector. An alternative description of such a

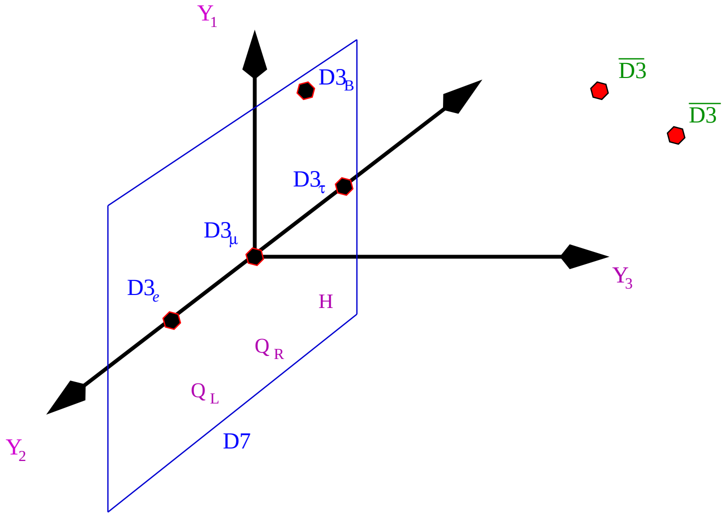

Figure 5: D-brane configuration of the model discussed in the text. The gauge group, Quarks and Higgs fields live in the worldvolume for the D7-brane whereas leptons are located at the three fixed points in the complex Y_2 dimension. Bulk $\mathrm{D3_{bulk}}$ branes, leading to $\mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ breaking also live in the Y_1, Y_2 hyperplane whereas anti-3-branes live away in the bulk of the third (Y_3) complex dimension.

decoupling situation can be achieved by performing a T-duality transformation in the third complex dimension. Thus, 9-branes become 7-branes located at the origin ($Y_3 = 0$) in the third complex plane with their world volume including the first two complex planes. 5-branes (antibranes) turn into 3-branes (antibranes) which can be located anywhere on the compact space. This is pictorically described in figure 5.[2]

Hence, if bulky anti-3-branes are placed away from $Y_3 = 0$ their worldvolume will not overlap that of the supersymmetric 7,3-brane sectors. Thus bulky, non-supersymmetric fields, will decouple from the supersymmetric sector and will only couple to it through the exchange of closed string interactions. This is the T-dual description of the Wilson line accounted for in the preceding paragraph. We adopt this point of view in what follows. Twist matrices in the T-dual description are obtained just by exchanging $7 \leftrightarrow 9$, $3(\bar{3}) \leftrightarrow 5(\bar{5})$ indices. Thus, spectrum and interactions are the same in both descriptions.

Twist matrices can be described in terms of associated shift vectors ([25, 19]). Such a description, which naturally appears when a Cartan-Weyl basis is chosen for the group algebra, is especially adapted for computing the spectrum [25, 19]. Thus,

[2]Notice that in this figure extra 'bulk' branes are added (see below) and only two antibranes are shown, the other antibranes are located at the image of these two under the combination of the $\mathbb{Z}_3$ defining the orbifold and the orientifold Ω (or $(-1)^{F_L}\Omega R_3$ in the 7,3-brane picture) which is a further $\mathbb{Z}_2$ symmetry, making for a total of $2 \times 3 \times 2 = 12$ antibranes which equals the number of 3-branes.

twist matrices above (now for 7,3-branes) correspond to

$$
\begin{aligned}
V_7 &= \frac{1}{3}(1,1,1;-1,-1;0,0;0,0;1,1,1,1,1,1,1) \\
W &= \frac{1}{3}(1,1,1;1,1;1,1;0,0;0,0,0,0,0,0,0)\,,
\end{aligned} \tag{3.4}
$$

and simply $V_{3,(2,i)} = 1/3$ for matrices in (3.3).

We then have for $(\mathcal{W})\gamma_{\theta,9}$ and $\mathcal{W}^{-1}\gamma_{\theta,9}$

$$
V_7 + W = \frac{1}{3}(2,2,2;0,0;1,1;0,0;1,1,1,1,1,1,1) \tag{3.5}
$$

$$
V_7 - W = \frac{1}{3}(0,0,0;1,1;2,2;0,0;1,1,1,1,1,1,1)\,, \tag{3.6}
$$

respectively. The total observable gauge group is $G = G_7 \times G_3$ with the **77**-sector brane group being $G_7 = \mathrm{U}(3) \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R \times \mathrm{SO}(4) \times \mathrm{U}(7)$, while the $\mathbf{33}_{2,i}$ sectors, with branes trapped at $(2,i)$ $(i = 0,1,2)$ fixed points leads to $G_3 = \mathrm{U}(1)_{2,0} \times \mathrm{U}(1)_{2,1} \times \mathrm{U}(1)_{2,2}$. The six anti-3-branes in the bulk give rise to a single Sp(2): two of them give rise to a Sp(2) and the other four are just $\mathbb{Z}_3$ mirrors of the first ones.

Let us analyze the supersymmetric part of the spectrum (see [19]). We find the chiral content

$$
\mathbf{77}\ \text{sector:}\ 3\left[\left(3,2,1,\frac{1}{3}\right) + \left(\bar{3},1,2,-\frac{1}{3}\right) + (1,2,2,0)\right] + 3[(1,21)' + (4,\bar{7})']\,,
$$

where we have indicated the representations under the Left-Right group $G_{LR} = \mathrm{SU}(3) \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ and $G' = \mathrm{SO}(4)' \times \mathrm{U}(7)'$. The abelian factor $\mathrm{U}(1)_{B-L}$ is generated by $Q_{B-L} = -\frac{2}{3}Q_3 - Q_L - Q_R$ introduced in (2.2) where Q_3, Q_L, Q_R are the generators of the abelian factor in the corresponding unitary groups. As mentioned, Q_{B-L} is identified with $B-L$ symmetry generator and it can be shown (see next section) to be non anomalous. The Standard Model hypercharge Y and electromagnetic charge are thus obtained as

$$
Y = \frac{Q_{B-L}}{2} - T_R^3\,, \qquad Q_{em} = Y + T_L^3\,, \tag{3.7}
$$

where $T^3_{R,L}$ are the diagonal generators of $\mathrm{SU}(2)_{R,L}$.

We observe that the **77** sector contains the standard three quark generations plus a set of three chiral Higgs fields $(1,2,2,0)$. The factor three here is associated to the three compact complex dimensions Y_a $(a = 0,1,2)$ each one of them feeling the same orbifold twist.

$$
\begin{aligned}
\mathbf{37}_{2,i}\ \text{sector:}\ &\left(3,1,1,-\frac{2}{3}\right)_{-1} + \left(\bar{3},1,1,\frac{2}{3}\right)_{-1} + (1,2,1,-1)_1 + \\
&+ (1,1,2,+1)_1 + [(1,7)'_1 + (4,1)'_1]\,.
\end{aligned}
$$

A subindex indicates the charge with respect to the $\mathrm{U}(1)_{3,(2,i)}$ 3-brane group. Since $i = 0,1,2$, we will have three identical copies. Hence, these sectors provide three generations of standard leptons.

77 sector	**37** sector	**33** sector
$3[(3,2,1)+(\bar{3},1,2)+(1,2,2)+$ $+(1,21)'+(4,\bar{7})']$	$3[(3,1,1)+(\bar{3},1,1)+$ $+(1,2,1)+(1,1,2)+$ $+(1,7)'+(4,1)']$	$3[(1)]$
37 sector (bulk)	**33** sector (bulk)	$\bar{\mathbf{3}}\bar{\mathbf{3}}$ sector
$(3,1,1;2)+(1,2,1;2)+(1,1,2;2)+$ $+(1,7;2)'+h.c.+(4,1;2)'$	$2(1)$ $+(3)$	$f_-:(1)\quad f_+:2(3)+(1)$ $s:2(1)+(3)$

Table 1: Spectrum of Left-Right model. We present the quantum numbers under the groups $\mathrm{SU}(3)_c\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{SO}(4)'\times\mathrm{SU}(7)'$ on the 7-branes, $\mathrm{U}(1)^3$ on the trapped 3-branes, $\mathrm{Sp}(2)$ on the bulk 3-branes and $\mathrm{Sp}(2)^2$ from the $\bar{3}$-branes.

We display the summary of the massless spectrum of the model in table 1.[3] Each of the three **33** sectors contain a singlet chiral field $(1)_2$. As we will see in the next section, these three singlets generically get vacuum expectation values of the order of the string scale, giving masses to the extra colour triplets in the **37** sector of the model. In this way, the massless spectrum of the model coupling to the $\mathrm{SU}(3)\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{U}(1)_{B-L}$ group is indeed the one proposed in section 2. Thus, the results for gauge coupling unification obtained there directly apply to the present model.

3.2 Anomalous U(1)'s and Fayet-Iliopoulos terms

The LR model contains seven abelian U(1) factors. Four of these terms, namely Q_3, Q_L, Q_R and Q_7, appear in the unitary groups of the **77** sector, whereas the other three $Q_{(2,i)}$ come from each of the $\mathbf{33}_{2,i}$ sectors. As we know, some of these U(1)'s are anomalous. The spectrum of the model with the corresponding U(1) charges is given in table 2. From there the matrix of mixed U(1)-non-abelian anomalies [19] can be computed to be

$$T_{IJ}^{\alpha\beta}=\begin{pmatrix} 0 & 9 & -9 & 0 & 0 \\ -6 & 0 & 6 & 0 & 0 \\ 6 & -6 & 0 & 0 & 0 \\ 0 & 0 & 0 & 21 & -21 \\ -2 & 1 & 1 & 1 & -1 \end{pmatrix}. \tag{3.8}$$

The rows correspond to the seven factors Q_3, Q_L, Q_R, Q_7 and $Q_{(2,i)}$ (same structure repeats for $i=0,1,2$). The columns correspond to the nonabelian groups: SU(3), $\mathrm{SU}(2)_L$, $\mathrm{SU}(2)_R$, SU(7) and SO(4) respectively.

[3] In tables 1, 2 and 3 we also display the extra massless fields which might appear if, in addition to the branes discussed above, there are further 3-branes living in the bulk (see figure 5) of the first two complex dimensions (but at $Y_3=0$). These 3-branes may be used to break the left-right symmetry down to the SM, as we discuss in next section. They do not modify one-loop coupling unification, though.

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There are two linear independent combinations of above generators which are free of anomalies whereas the other five have triangle anomalies.[4]

A possible choice for the non anomalous generators is[5]

$$
\begin{aligned}
Q_{B-L} &= -\frac{2}{3}Q_3 - Q_L - Q_R \\
Q_X &= Q_R - Q_L - \frac{2}{7}Q_7 + 2\sum_i Q_{n_2^i}\,,
\end{aligned}
$$

while anomalous ones can be chosen as

$$
\begin{aligned}
Q_{A1} &= Q_{n_2^0} - Q_{n_2^1} \\
Q_{A2} &= Q_{n_2^0} + Q_{n_2^1} - 2Q_{n_2^2} \\
Q_{A3} &= Q_3 - Q_L + Q_7 \\
Q_{A4} &= Q_3 - Q_R - Q_7 \\
Q_{A5} &= \sum_i Q_{n_2^i} + 3Q_7\,.
\end{aligned} \tag{3.9}
$$

Matter fields	Q_3	Q_L	Q_R	Q_7	$Q_{(2,i)}$
77 sector					
$(3,2,1)$	1	-1	0	0	0
$(\bar{3},1,2)$	-1	0	1	0	0
$(1,2,2)$	0	1	-1	0	0
$(4,\bar{7})'$	0	0	0	-1	0
$(1,21)'$	0	0	0	2	0
37 sector					
$(3,1,1)$	1	0	0	0	-1
$(\bar{3},1,1)$	-1	0	0	0	-1
$(1,2,1)$	0	1	0	0	1
$(1,1,2)$	0	0	-1	0	1
$(4,1)'$	0	0	0	0	-1
$(1,7)'$	0	0	0	1	1
33 sector					
(1)	0	0	0	0	2
37 bulk					
$(3,1,1;2)$	1	0	0	0	0
$(\bar{3},1,1;2)$	-1	0	0	0	0
$(1,2,1;2)$	0	1	0	0	0
$(1,2,1;2)$	0	-1	0	0	0
$(1,1,2;2)$	0	0	1	0	0
$(1,1,2;2)$	0	0	-1	0	0
$(1,7;2)'$	0	0	0	1	0
$(1,\bar{7};2)'$	0	0	0	-1	0
$(4,1;2)'$	0	0	0	0	0

Table 2: Spectrum of Left-Right model. We present the quantum numbers under the $U(1)^7$ groups. The first 4 U(1)'s come from the 7-brane sector. The next three come from the 3-brane sector, these we have written as a single column with the understanding that for instance in the **37** sector, each of the three copies have that charge under one of the three U(1)'s and zero under the other two.

Notice that even though Q_{A1} and Q_{A2} present no mixed U(1)-non-abelian anomalies, they have cubic and mixed U(1) anomalies. Under these new combinations, the charges of the particles in the spectrum are displayed in table 3.

As usual in type-I theory [26], U(1) anomalies are cancelled by a generalized Green-Schwarz mechanism through the coupling to twisted close string RR fields [27]. Anomalous U(1)s become massive [29]. At the same time, because of supersymmetry, a Fayet-Iliopoulos term, associated to each of the anomalous groups, appears [27]–[31]. The corresponding D-term potential is

$$
V_r = \frac{1}{2}\left(\xi_r + \sum_l q_r^l |\phi_l|^2\right)^2, \tag{3.10}
$$

[4]This is a generic feature of this type of models with 5-branes at just one (a,i) (a fixed) set of fixed points. If branes are stuck at two different a's then there are ten U(1)'s and three of them are non anomalous. If there are branes at the three sets $a=0,1,2$ then four of the thirteen abelian factors are anomaly free.

[5]It is amusing that if one looks at the Q_X charges of quarks, leptons and Higgs fields, they are identical to the ones such fields have under the U(1) contained in the branching $E_6 \to SO(10)\times U(1)$. Here, though, there is no E_6 nor SO(10) symmetry present.

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Matter fields	Q_{B-L}	Q_X	Q_{A1}	Q_{A2}	Q_{A3}	Q_{A4}	Q_{A5}
77 Sector							
$(3,2,1)$	1/3	1	0	0	2	1	0
$(\bar{3},1,2)$	-1/3	1	0	0	-1	-2	0
$(1,2,2)$	0	-2	0	0	-1	1	0
$(4,\bar{7})'$	0	2/7	0	0	-1	1	-3
$(1,21)'$	0	-4/7	0	0	2	-2	6
37 sector							
$(3,1,1)$	-2/3	-2	(-1,1,0)	(-1,-1,2)	1	1	-1
$(\bar{3},1,1)$	2/3	-2	(-1,1,0)	(-1,-1,2)	-1	-1	-1
$(1,2,1)$	-1	1	(1,-1,0)	(1,1,-2)	-1	0	1
$(1,1,2)$	1	1	(1,-1,0)	(1,1,-2)	0	1	1
$(4,1)'$	0	-2	(-1,1,0)	(-1,-1,2)	0	0	-1
$(1,7)'$	0	12/7	(1,-1,0)	(1,1,-2)	1	-1	4
33 sector							
(1)	0	4	(2,-2,0)	(2,2,-4)	0	0	2
37 bulk							
$(3,1,1;2)$	-2/3	0	0	0	1	1	0
$(\bar{3},1,1;2)$	2/3	0	0	0	-1	-1	0
$(1,2,1;2)$	-1	-1	0	0	-1	0	0
$(1,2,1;2)$	1	1	0	0	1	0	0
$(1,1,2;2)$	-1	1	0	0	0	-1	0
$(1,1,2;2)$	1	-1	0	0	0	1	0
$(1,7;2)'$	0	-2/7	0	0	1	-1	3
$(1,\bar{7};2)'$	0	2/7	0	0	-1	1	-3
$(4,1;2)'$	0	0	0	0	0	0	0

Table 3: Spectrum of Left-Right model. We present the quantum numbers under the 4 anomaly free U(1) groups and the three anomalous U(1)'s. Some of the fields in the **37** sector have several entries for Q_Y and Q_Z, the reason being that the fields come in three copies which differ by those charges.

where ϕ_l is the scalar field with charge q_l^r under the anomalous group $\mathrm{U}(1)_{Ar}$. The ξ_r $r=1,\ldots,5$ terms can be explicitly computed (see ref. [19, eq. 3.20]) in terms of the fields $M_{(a,i)}$, the Neveu-Schwarz partners of the RR antisymmetric forms mentioned above. We find

$$
\begin{aligned}
\xi_1 &= \frac{3\sqrt{3}}{2}\left(M_{02}-M_{12}\right), & \xi_2 &= \frac{3\sqrt{3}}{2}\left(M_{02}+M_{12}-2M_{22}\right),\\
\xi_3 &= \frac{\sqrt{3}}{2}\sum_i\left(12M_{0i}+4M_{1i}+5M_{2i}\right), & \xi_4 &= -\frac{\sqrt{3}}{2}\sum_i\left(4M_{0i}+12M_{1i}+5M_{2i}\right),\\
\xi_5 &= \frac{\sqrt{3}}{2}\sum_i(7M_{0i}+7M_{1i}+8M_{2i})\,.
\end{aligned}
\tag{3.11}
$$

Notice that all the FI-terms ξ_r are linearly independent combinations of the $NS-NS$ twisted moduli. This means that, unlike what usually happens in the perturbative heterotic vacua [32] , there is no need to check for D-flatness of the scalar potentials eq. (3.10). This is because for any field direction of the scalars ϕ_l charged under each anomalous $U(1)_r$, there will be vevs for the twisted moduli yielding ξ_r's compensating them. Furthermore, this indicates, in general, a departure of the orbifold limit and also that the anomalous U(1)'s do not remain as effective global symmetries, as it would have happened if all twisted moduli were vanishing [33].

3.3 The structure of mass scales

The unification of coupling constants in this model takes place at a scale of order 9×10^{11} GeV which should then be identified with the string scale M_s. In addition that is also the order of magnitude of the compactification scales M_1, M_2 of the radii of the first two complex dimensions. This is desirable for two reasons: 1) Since the worldvolume of 7-branes includes the first two complex dimensions, if $M_{1,2}$ were much smaller there would be charged Kaluza-Klein fields which might spoil gauge coupling unification; 2) Some phenomenologically interesting non-renormalizable Yukawa couplings involve (see next section) 3-branes living at different locations in the first two compact directions. Such couplings would be considerably suppressed if $M_{1,2}$ were much smaller than M_s. On the other hand, the compactification scale M_3 along the third complex plane (which is transverse to the 7-branes worldvolume) is unconstrained by these considerations. The Planck mass is related to the string scale M_s and the compactification scales M_i by (see e.g. [11]):

$$M_p = \frac{2\sqrt{2} M_s^4}{\lambda M_1 M_2 M_3} = \frac{\sqrt{2}}{\alpha_7} \frac{M_1 M_2}{M_3}\,, \tag{3.12}$$

where $\alpha_7 = \lambda M_1^2 M_2^2 / 2 M_s^4$ is the unified coupling of the group coming from 7-branes, which includes the standard model group. Thus, for $M_{1,2} \propto M_s = 9 \times 10^{11}$ GeV, one can obtain the measured M_p for $M_3 \propto (100)/\alpha_7$ TeV. In this scheme (see figure 5) the size of the Y_3 coordinate would be thus very large compared to $Y_{1,2}$.

The present class of models contain anti-3-branes in the bulk in transverse space. As depicted in figure 5, their worldvolume does not have overlap with that of the "visible world" of 3-branes and 7-branes once the latter are located at the origin in the third compact dimension. Anti-3-branes are instead in the bulk in that dimension. The global configuration of the model is non-supersymmetric, since the supersymmetries preserved by branes are broken by the anti-branes and viceversa.[6] Closed string states living in the bulk of space will generically communicate supersymmetry breaking from the anti-3-brane sector to the visible sector of 3-branes and 7-branes. We will assume that the presence of SUSY-breaking anti-3-branes

[6]It is worth pointing out that even though the presence of the anti-branes explicitly break supersymmetry, the number of massless bosonic degrees of freedom still matches the number of massless fermionic degrees of freedom as can be easily seen in all models of this type, following the general spectrum of reference [19].

in the bulk constitutes a SUSY-breaking hidden sector for this model. Since these anti-3-branes live far away in the bulk of the (very large) third complex dimension, SUSY-breaking effects in the 7-branes and 3-branes where the SM resides will be Planck mass suppressed. Thus one expects SUSY-breaking soft terms of order:

$$M_{\text{soft}} = \epsilon \frac{M_s^2}{M_p} \,, \tag{3.13}$$

where the value of the fudge factor ϵ will depend on the details of how SUSY-breaking effects in the antibranes are transmitted to the branes by the massless closed string fields. Since gauge coupling unification predicts $M_s = 9 \times 10^{11}$ GeV, in order to get soft terms of order, say 1 TeV, we need[7] $\epsilon \propto 10^{-2}$.

The above assumption of a very large Y_3 dimensions is a possible simple explanation for the observed large size of M_p compared to our predicted $M_s = 9 \times 10^{11}$ GeV. Recently an alternative explanation has been proposed [36] to obtain such an effect which may occur (in some simple models) even if the extra dimensions are infinite. This occurs due to the presence of warp factors in the space-time metric exponentially depending on the extra dimensions. Furthermore, it has also been argued [37] that a localized set of $D3$ branes does indeed induce a warped geometry around its location. It would be interesting to explore whether this kind of arguments extend to configurations like the one discussed here which involve intersections of both 3-branes and 7-branes.[8] An exponential warp factor depending on the dimension Y_3 transverse to both 3-branes and 7-branes could in this case be a possible alternative origin for the M_p/M_s hierarchy in a model like the one studied here.

3.4 Yukawa couplings and conservation rules

The general structure of renormalizable couplings in this class of orientifolds was already discussed in ref. [19]. Let us review the couplings involving the supersymmetric sector for the present model, leaving their phenomenological implications for the next section.

(i) $(\mathbf{77})^3$ couplings. These have the form:

$$\phi_i^{77} \phi_j^{77} \phi_k^{77} \,, \qquad i \neq j \neq k \neq i \,, \tag{3.14}$$

where ϕ_i^{77}, $i = 1, 2, 3$ are any of the charged chiral fields in the **77** sector associated to the complex plane i. These type of couplings give rise for example to quark Yukawa couplings, as we discuss below. The coupling is proportional to the gauge coupling constant for the **77** gauge interactions g_7, which is the one associated to the

[7] This seems to suggest a one-loop transmission of SUSY-breaking to the observable D-brane sectors, as occurs for example in moduli dominated [34] and/or anomaly mediated [35] scenarios.

[8] For recent studies of the Randall-Sundrum scenario in the presence of brane intersections see [38].

physical gauge fields. Recall that the latter is related to the string scale M_s and the compactification scales $M_{1,2}$ of the first two complex planes by:

$$\alpha_7 = \frac{g_7^2}{4\pi} = \frac{\lambda M_1^2 M_2^2}{2M_s^4}\,, \tag{3.15}$$

where λ is the type-IIB dilaton coupling. It is this α_7 which provides the boundary conditions for the running of the gauge couplings of the left-right symmetric model.

(ii) (73)(73)(77) couplings These in principle only involve the **77** sector associated to the third complex plane [19]:

$$\psi_i^{73}\psi_i^{73}\phi_3^{77}\,, \tag{3.16}$$

where $i = 0, 1, 2$ labels the fixed points where the 3-brane is localized. Notice that these couplings are diagonal in the i label, i.e. there are no renormalizable couplings involving different fixed points. These Yukawa couplings are also proportional to the **77** gauge coupling constant g_7.

(iii) (73)(73)(33) couplings In a similar manner there are superpotential couplings of the form

$$\psi_i^{73}\psi_i^{73}\phi_{3,i}^{33} \tag{3.17}$$

in which again i labels the fixed point. Again, only the **33** chiral fields in the third complex plane appear in the coupling. For example, we already mentioned that there is a coupling of this type between the singlets $(1)_2$ in the $\mathbf{33_i}$ sectors and the coloured triplets in the $\mathbf{73_i}$ sectors. Notice however that the gauge coupling $\tilde{g}$ is now different, with $\tilde{\alpha} = \lambda/2$.

Several comments concerning the above couplings are in order. From the string point of view, these couplings are obtained from a disk-shaped worldsheet at which boundaries three open string vertex operators are attached. The boundaries of the disk represent the relevant p-branes, 3-branes and 7-branes in our case. Thus an insertion of a vertex operator of a particle in a $\mathbf{73_i}$ sector turns a 7-brane boundary into a 3_i-brane boundary (and viceversa). This implies that, for the disk worldsheet to make sense, $\mathbf{73_i}$ vertex insertions (for each different i) have to come in pairs (see figure 6). Thus there is a $\mathbb{Z}_2 \times \mathbb{Z}_2 \times \mathbb{Z}_2$ symmetry which is respected by all disk couplings. This is obviously respected in the couplings discussed above.

These symmetries will have an important phenomenological role in the present model, as discussed in the next section.

A second question concerns the structure of the couplings $\mathbf{(77)^3}$ above. The reader familiar with heterotic orbifold constructions will realize that the same type of couplings involving necessarily the three different complex planes are present for the untwisted particles in those constructions. In the case of heterotic orbifolds this antisymmetric structure may be understood in terms of the conservation of the so

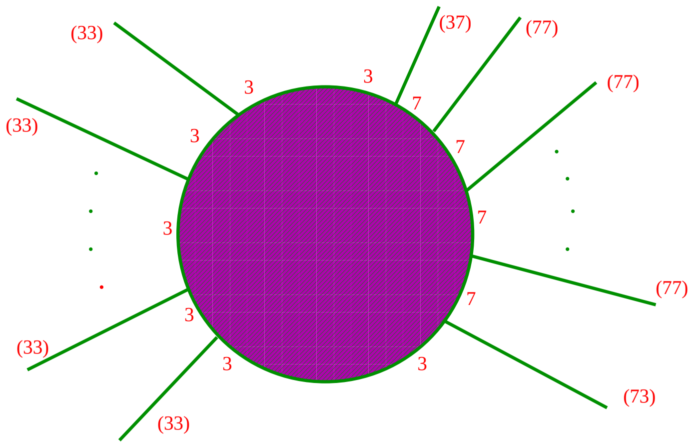

Figure 6: Disk couplings of vertex operators of massless fields from $(\mathbf{33}), (\mathbf{37}), (\mathbf{77})$ sectors. Vertices of $(\mathbf{37}), (\mathbf{73})$ particles must come in pairs in order to get a consistent D-brane boundary on the disk.

called *H-momentum* (see e.g., ref. [39, 40] for a discussion of these symmetries). The right-moving vertex operators have factors proportional to the RNS fermions which, when bosonized, can be written as $\exp(i\mathbf{p}.\mathbf{H})$ where $\mathbf{p}$ is an SO(10) (space-time) weight. H-momentum conservation is the statement that the overall momentum $\mathbf{p}$ in a correlator has to vanish for a coupling to be allowed. Now, in the type-I case something completely analogous may be defined leading to equivalent symmetries.

The above $\mathbb{Z}_2$ symmetries and H-momentum conservation rules are still valid for disk (i.e. tree-level) amplitudes leading to non-renormalizable couplings involving charged open string fields. If a given coupling involves branes living at different points in transverse space, it will be exponentially suppressed by the distance between those branes. Thus, for example, couplings involving the 3_i branes and the $3_{\rm bulk}$ branes will get such a suppression. Notice however that if the compactification scales along the first two complex planes $M_{1,2}$ are of order the string scale M_s, no such a suppression will be present. This is in fact the case considered in the previous subsection in order to understand the hierarchy $M_p \gg M_s$: only the third complex dimension is large and only 3 branes distant in the 3-d complex dimension are exponentially suppressed. This is what we will assume in the phenomenological analysis in the next section.

Another point concerning non-renormalizable couplings is the existence of couplings violating the conservation of *anomalous* U(1) symmetries. Tree level couplings renormalizable or not should respect all non-anomalous U(1) symmetries. However they may violate anomalous U(1) symmetries since those, as we discussed above, are broken by the vevs of twisted NS-NS fields as long as we are away from the orbifold limit (which is the case in the models discussed). As an example of this we show in appendix B how non-renormalizable couplings violating anomalous U(1) symmetries are expected to appear (on the basis of heterotic/type-I duality) in the standard $\mathbb{Z}_3$ orientifold.

4. Phenomenology of the D-brane left-right symmetric orientifold

In this section we discuss several phenomenological aspects of this model. We will not attempt a detailed description of all possible aspects like fermion masses or spontaneous gauge symmetry breaking, rather we will only discuss some possible avenues enabling to address the gross phenomenological issues in this model.

(i) A nearby vacuum. As we mentioned in the previous section, apart from the three left-right symmetric generations and Higgs fields, this particular string model has three copies of vector-like colour triplets form the (37) sectors. However, these states are generically massive. Indeed, looking at Table 3 one sees that all gauge interactions allow for a coupling between the three singlets 1_4 from the (33) sectors to the three pairs $(3, 1, 1, -2/3)_{-2}+(\bar{3}, 1, 1, +2/3)_{-2}$, where the subindices denote the Q_X charge. These Yukawa couplings do indeed exist, as discussed in previous section. Thus, if the three singlets 1_4 get a vev of order M_s, the colour triplets will disappear from the low-energy spectrum and we will be left at low energies with precisely the massless spectrum discussed in section 2, leading to very good predictions for gauge coupling unification.[9]

If we give a vev $\langle 1_4\rangle_a \propto M_s$, $a = 1, 2, 3$ to these (33) singlets, we have to ensure D-flatness and F-flatness for this direction. In fact we should not care too much about the D-terms of the anomalous U(1)'s because they can be easily cancelled for appropriate values of the blowing-up fields M discussed in section 3. On the other hand, the D-term corresponding to the anomaly-free Q_X generator has to cancel, which requires giving a vev to some fields with negative Q_X charge. A natural option seems to be giving vevs to the antisymmetric A^{ij} of SU(7) present in the $(77)_3$ sector as follows:[10]

$$A^{12} = A^{23} = A^{34} = A^{45} = A^{56} = A^{67} = A^{71} = v\,, \tag{4.1}$$

with $v \propto M_s$ (actually it should be somewhat lower than M_s in order to avoid higher-order operators contributions). This direction can be easily seen to be D-flat and F-flat. Below M_s the only U(1) interaction left is now Q_{B-L} since Q_X is broken by the above vevs.

(ii) The breaking of the left-right symmetry and bulk 3-branes. The chiral multiplet content of our model below M_s includes just thee generations of quarks/leptons/Higgs fields. Some additional fields (particularly, $\mathrm{SU}(2)_R$ doublets) are needed if we want to break our theory down to the SM gauge group. Probably

[9]Notice that, on the other hand, equivalent couplings to the Higgs doublets $(1, 2, 2, 0)$ in the (77) sector do not exist. Thus a doublet-triplet splitting mechanism is built-in in the symmetries of the model.

[10]This also turns out to give rise to the required masses for right-handed neutrinos.

there is more than one way to modify the model in such a way that one has additional massless $SU(2)_R$ doublets for symmetry breaking while the good coupling unification predictions are not spoiled.[11] The simplest possibility seems to add some additional 3-branes moving in the bulk in the first two compact directions but at the origin in the third compact direction (so that their worldvolume overlaps with that of 7-branes). The simplest set of 3-branes that one can add in the bulk are 6 of them (one 3-brane and their orbifold and orientifold mirrors). They lead to a SU(2) gauge group in the $(\mathbf{33})_{\rm bulk}$ sector and massless chiral fields in the $(\mathbf{73}_{\rm bulk})$ sector transforming like:

$$\left[\left(3,1,1,-\frac{2}{3};2\right)+(1,2,1,-1;2)+(1,1,2,+1;2)+\text{h.c.}\right]+(1,7;2)'+\text{h.c.}+(4,1;2)'\,. \tag{4.2}$$

In addition there are chiral fields in the $(\mathbf{33})_{\rm bulk}$ sector transforming like $2(1)+(3)$ under the SU(2) group on the 3-branes. We will see later on when we discuss neutrino masses that the SU(2) group coming from this bulky 3-branes should be broken close to the M_s scale by vacuum expectation values of the fields $(1,7;2)'+\text{h.c.}$ above.

The chiral fields in eq.(4.2) include $SU(2)_R$ doublets which can in principle get a vev and break the symmetry. Thus we will assume that some of the fields in $(1,1,2,+1;2)+(1,1,2,-1;2)$ will get vacuum expectation values of order $M_R\propto$ 1 TeV and break the symmetry to that of the SM. We will briefly discuss below how that could take place due to a radiative symmetry breaking mechanism.

One interesting point of the extra particle content provided by the addition of these "bulky 3-branes" is that they give a net vanishing contribution to the combinations $(B_1'+B_L-\frac{14}{3}B_3)$ and $(B_L-\frac{3}{11}B_1')$ which control the joining of coupling constants. Indeed we can easily check that extra contributions to the β-functions are obtained:

$$\Delta B_L=\Delta B_R=\Delta B_3=+2\,,\qquad \Delta B_1'=\Delta B_R+\frac{1}{4}\Delta B_L=\frac{22}{3} \tag{4.3}$$

so that unification of couplings is not modified[12] at one loop.

(iii) Quark and charged lepton masses. In this model renormalizable quark Yukawa couplings of type $(77)^3$ exist with the structure:

$$g\,\epsilon_{ijk}\left(3,2,1,\frac{1}{3}\right)_i\left(\bar{3},1,2,-\frac{1}{3}\right)_j(1,2,2,0)_k\,, \tag{4.4}$$

where $i,j,k=1,2,3$ label the three complex planes and g is the (77) gauge coupling constant. With this simple structure, there would be a massless quark generation

[11] In particular, the variant left-right symmetric model displayed in table 1 of ref. [19] is another possibility. That model has additional matter but one can check that gauge coupling unification along similar lines to those of the present model takes place. A discussion of this variant model is presented in appendix A.

[12] This is analogous to the well known fact that complete SU(5) representations do not modify the one-loop conditions for unification in the MSSM.

and two degenerate generations with masses of order $g\sqrt{\sum_i |\langle H_i\rangle|^2}$. However, this structure is modified by various effects. To start with, the Kahler metric for the (77) matter fields in a model like this one needs not be diagonal. There are Kahler untwisted moduli which mix the different complex planes. Furthermore, other effects mixing different complex planes may come from non-renormalizable D-terms like, e.g. $\langle 21_i 21_j^*\rangle Q_R^i Q_R^j *$. In addition to these, there are mixing terms with the color triplets from the bulky branes. In particular there are renormalizable Yukawa couplings of the form $(77)(73_{\rm bulk})^2$:

$$\begin{aligned} &\left(\bar{3},1,2,-\frac{1}{3}\right)_3 \times \left(3,1,1,-\frac{2}{3};2\right) \times \langle(1,1,2,1;2)\rangle \\ &\left(3,2,1,\frac{1}{3}\right)_3 \times \left(\bar{3},1,1,\frac{2}{3};2\right) \times \langle(1,2,1,-1;2)\rangle\,, \end{aligned} \tag{4.5}$$

and hence the right-handed D-quarks of the (77) sector mix with the colour triplets from the $(73_{\rm bulk})$ sector once the $SU(2)_R$ doublets in that sector get a vev. This means that generically two physical right-handed D-quarks (the three of them in the variant model of the appendices) will be $SU(2)_R$ singlets and the other will be contained in a doublet.[13] The second Yukawa coupling above will give masses to the first two D-quarks and the couplings in (4.4) will give masses to the third.

Concerning the possible Yukawa couplings for the leptons, the following couplings in the disk are allowed by all non-anomalous gauge symmetries:

$$g\,(1,2,2,0)_3 \times (1,2,1,-1)_a \times (1,1,2,+1)_a \times f_a(M_\gamma)\,, \qquad a=1,2,3 \tag{4.6}$$

i.e. the Higgs fields along the third complex plane in the (77) sector couple diagonally to the lepton generations. Looking at tables 2 and 3 we can observe that these couplings are indeed allowed by all non-anomalous symmetries, including the Q_X generator. However they are in principle forbidden by the anomalous U(1)'s which are spontaneously broken at the string scale. The twisted moduli fields discussed in the previous section are charged (non-linearly) under these anomalous U(1)'s so one expects that upon the insertion of coherent sets of twisted vertex M_γ operators the above couplings will be allowed, as discussed in previous section and exemplified in appendix B. This we denote by the addition of the factor $f_a(M_\gamma)$ in the above expression. Notice that this factor does not necessarily mean an exponential suppression, since the physical gauge couplings of the SM gauge interactions are given by the couplings on the (77) sector which are proportional to $\lambda/(M_s^4 R_1^2 R_2^2)$, but not to the dilaton λ itself. Thus λ need not be too small a number (see also the discussion in appendix B). The precise size of the obtained lepton masses depends on the size of the vev for $(1,2,2,0)_3$ and on the value of $f_a(M_\gamma)$ for each a.

[13] This is analogous to the *alternate* left-right models considered in refs. [41, 42]. Those models are interesting from the point of view of supression of FCNC, as we comment below.

(iv) Neutrino masses. In this model there is no right-handed $SU(2)_R$ triplet which might give a large Majorana mass to the right-handed neutrinos. However there are fields N_a which are singlets under the $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ group and can combine with the right-handed neutrinos which then get a Dirac mass of order M_R. Specifically, those singlets are contained in the $(1, 7')$ representations in the three (37) sectors. For the relevant couplings to appear we have to give vevs of order the string scale to the $(1, 7; 2)' + (1, \bar{7}; 2)$ chiral fields in the (73_{bulk}) sector. In particular there is a D-flat and F-flat direction along $\Psi_6^2 = \overline{\Psi}_6^1 = u$ and $\Psi_7^2 = \overline{\Psi}_7^1 = iu$, where in Ψ_r^s $(\overline{\Psi}_r^s)$ the index r runs over SU(7) and s over the SU(2). Then an effective renormalizable Yukawa coupling is induced at low energies of the form:

$$(1, 1, 2, -1)_a \times (1, 1, 2, +1; 2) \times (1, 7)'_a \left\langle (1, 21')^6 \times (1, 7; 2) h(M_\gamma) \right\rangle , \tag{4.7}$$

where the $(1, 1, 2, +1; 2)$ are $SU(2)_R$ doublets from the 3-branes in the bulk. One can check that this coupling is allowed by all anomaly-free gauge interactions of the model. As happened with the masses of charged leptons, insertions of twisted moduli fields will be required, which we parameterize by the factor $h(M_\gamma)$. Notice that this coupling, since it is non-renormalizable, is in principle suppressed by powers of M_s. It is of the general form $(73)^2(73_{\text{bulk}})^2(77)^6$ and hence involves 3-branes located at different points which will also mean exponential suppression in the distance between the location of the 3-branes at the fixed points and those in the bulk. Notice however that, as we discussed in the previous section, we have chosen the first two complex compact directions with sizes of order $1/M_s$ and hence there is not necessarily any extra suppression, only the third compact complex dimension is assumed to be very large.

Once the fields $(1, 1, 2, +1; 2)$ get a vev breaking spontaneously the $SU(2)_R$ symmetry, the right handed neutrinos inside the three $(1, 1, 2, -1)_a$ fields will get a mass of order M_R combining with some singlets inside the $(1, 7)'_a$. Notice in this connection that generically the SU(7) gauge symmetry is broken and those fields behave indeed like singlet partners of the right-handed neutrinos. In this situation the left-handed neutrinos remain massless. However there are mixing terms from analogous couplings involving $SU(2)_L$ doublets of the form:

$$(1, 2, 1, +1)_a \times (1, 2, 1, -1; 2) \times (1, 7)'_a \left\langle (1, 21')^6 \times (1, 7; 2) h(M_\gamma) \right\rangle . \tag{4.8}$$

Then the left handed neutrinos get induced Majorana masses of order $m_{\nu_L} \propto m_l \times (\langle (1, 2, 1, -1; 2) \rangle / \langle (1, 1, 2, +1; 2) \rangle)$. The particular sizes depend on the vev of $\langle (1, 2, 1, -1; 2) \rangle$, since $\langle (1, 1, 2, +1; 2) \rangle$ we know is of order $M_R \propto 1\,\text{TeV}$. One thus gets neutrino masses of order:

$$m_\nu^a \propto m_l^a \times \frac{\langle (1, 2, 1, -1; 2) \rangle}{M_R} . \tag{4.9}$$

For $\langle(1,2,1,-1;2)\rangle \propto m_l$ a seesaw-like formula is obtained but the precise sizes depend on the unknown values of the vevs of the $\mathrm{SU}(2)_L$ doublets $\langle(1,2,1,-1;2)\rangle$. Notice however that these mass contributions are flavour diagonal, there is no mixing between different lepton families. Thus oscillations can only take place into some inert sterile massless neutrino contained in the original $(1,7)'_a$ fields. This is not a generic property of the present scenario. One can check that in the variant model described in the appendices mixing between different neutrino flavors can take place, since there are no $\mathbb{Z}_2$ lepton parities.

(v) Discrete symmetries and proton stability The couplings in this orientifold model respect a number of discrete $\mathbb{Z}_2$ symmetries:

(a) There is a $\mathbb{Z}_2$ symmetry associated to each of the three (37) sectors. Under it all (73) particles are odd and the rest are even. Indeed, if we consider the couplings of (37) particles on the boundary of the disk, they have to appear in multiplets of two (see figure 6). Since in these sectors live the leptons (and some singlets coming from the $(1,7')$'s which, as we saw above behave like neutrino-like fields), this corresponded to a discrete $\mathbb{Z}_2$ lepton number parity. There is one $\mathbb{Z}_2$ symmetry for each of the three flavours.

(b) The flat direction considered gives vevs to the fields $(1,7;2)' + (1,\bar{7};2)'$ and also to some SU(7) antisymmetric fields. Thus this direction respects a $\mathbb{Z}_2$ symmetry under which 7-plets and SU(2) doublets (with respect to the (33_{bulk}) group) are odd. Under this symmetry all quarks and leptons are even but the $(1,7')$ fields in the (73) sectors are odd. The fields in the (73_{bulk}) are odd, since all are SU(2) doublets.

In fact , after breaking of the $\mathrm{SU}(2)_R$ symmetry by the $(1,1,2,\pm1;2)$ fields and of the $\mathrm{SU}(2)_L$ by the $(1,2,2,0)$ (or, in addition, the $(1,2,1,\pm1;2)$ fields), the diagonal $\mathbb{Z}_2$ which is the combination of the original $\mathbb{Z}_2$ and the center of $\mathrm{SU}(2)_R$ and $\mathrm{SU}(2)_L$ remains still unbroken. Thus, even after electroweak breaking a $\mathbb{Z}_2$ symmetry remains under which

- All quarks and leptons are odd.
- Higgs fields breaking $\mathrm{SU}(2)_L$ are even.
- Singlets combining with right-handed neutrinos are odd.
- $\mathrm{SU}(2)_R$ and $\mathrm{SU}(2)_L$ doublets in the (73_{bulk}) sector are even.
- Colour triplets in the (73_{bulk}) sector are odd.

Note that this residual symmetry can be identified with the standard R-parity of supersymmetric models.

In summary, the effective lagrangean has a residual R-parity symmetry[14] and in addition three lepton parities, one per lepton flavour. It is well known that R-parity may be considered as a discrete $\mathbb{Z}_2$ subgroup of the B-L symmetry. Thus combining it with the lepton parities we thus have a $\mathbb{Z}_2$ symmetry associated to baryon number. This means that nucleons are stable, since a $\mathbb{Z}_2$ baryon parity has to be conserved under which baryons are odd and leptons and mesons are even. Thus protons are absolutely stable. On the other hand baryon number can be violated in two units, since there is a $\mathbb{Z}_2$ symmetry. This means that in principle there can be neutron-antineutron transitions allowed. However those transitions violate B-L symmetry and hence are suppressed by high powers of (M_R/M_s) and the rate in this model is negligible. Notice on the contrary that discrete symmetries allow for the neutrino masses discussed in the previous subsection.

(vi) Gauge symmetry breaking and the low energy spectrum. As we discussed in the previous section, due to the presence of anti-3-branes in the bulk, one expects the generation of SUSY-breaking soft terms in the effective action. Once soft SUSY-breaking terms appear, $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ gauge symmetry breaking can occur due to loop corrections. We will not perform a complete analysis of the (quite involved) scalar potential, but will just study what scalar fields are likely to get vevs once loop corrections are included. As usual they will be the SU(3) colour singlets with Yukawa couplings to coloured fields. These include the $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$ doublets in the model, as well as the fields in the (33_{bulk}) sector which are triplets $(1;3)$ under the $\mathrm{SU}(2)_{\text{bulk}}$ gauge group. The following Yukawa couplings appear at the renormalizable level:

$$\begin{aligned}
&\left(3,2,1,\frac{1}{3}\right)_i \times \left(\bar{3},1,2,-\frac{1}{3}\right)_j \times (1,2,2,0)_k\,,\\
&\left(\bar{3},1,2,-\frac{1}{3}\right)_3 \times \left(3,1,1,-\frac{2}{3};2\right) \times (1,1,2,+1;2)\,,\\
&\left(3,2,1,\frac{1}{3}\right)_3 \times \left(\bar{3},1,1,\frac{2}{3};2\right) \times (1,2,1,-1;2)\,,\\
&\left(3,1,1,-\frac{2}{3};2\right) \times \left(\bar{3},1,1,+\frac{2}{3};2\right) \times (1;3)\,.
\end{aligned} \tag{4.10}$$

The first of these couplings is the $(77)^3$ quark Yukawa coupling that we mentioned above. The second and third couplings are of type $(77)_3(73_{\text{bulk}})^2$ and the fourth of type $(33_{\text{bulk}})(73_{\text{bulk}})^2$. All these four Yukawa couplings tend to give negative mass^2 to the above colour singlet scalars from one-loop diagrams in which the colour triplets circulate in the loop.

[14]It is well known that a residual R-parity remains in left-right symmetric models if the $\mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ symmetry is broken by $\mathrm{SU}(2)_R$ triplets $(1,1,3,-2)$. Notice that this is *not* the case here and the origin of the residual R-parity is different.

In addition one also expects generically the presence of trilinear *scalar* couplings ("A-terms") involving the colour singlet scalars. They are proportional to the scalar couplings

$$
\begin{aligned}
&(1;3) \times (1,2,1,-1;2) \times (1,2,1,+1;2) + \text{h.c.} \\
&(1;3) \times (1,1,2,+1;2) \times (1,1,2,-1;2) + \text{h.c.} \\
&(1,2,2,0)_3 \times (1,2,1,-1;2) \times (1,1,2,+1;2) + \text{h.c.} \\
&(1,2,2,0)_3 \times (1,2,1,+1;2) \times (1,1,2,-1;2) + \text{h.c.}\,.
\end{aligned}
\tag{4.11}
$$

The first two have couplings proportional to the $(33_{\rm bulk})$ gauge coupling constant whereas the last two are proportional to the (77) gauge coupling. The corresponding A-terms are proportional to $M_{\rm soft}$ and only involve the corresponding scalars. These contributions to the scalar potential are not positive definite and will favor all $\mathrm{SU}(2)_R$ and $\mathrm{SU}(2)_L$ doublets (and the scalars in $(1,3)$) to get a non-vanishing vev at some level. We will assume that a stable minimum of the scalar potential exists for vevs of the order of magnitude:

$$
\begin{aligned}
&\langle (1,1,2,+1;2)\rangle \propto \langle (1,1,2,-1;2)\rangle \propto \langle (1;3)\rangle \propto M_R \propto 1\,\text{TeV} \\
&\quad \langle (1,2,2,0)_i\rangle \propto M_Z \ll M_R
\end{aligned}
\tag{4.12}
$$

so that the required hierarchy between the left and right gauge symmetries is obtained. Notice that in principle all interactions respect an explicit parity left↔right symmetry and it is the vacuum which will explicitly break parity symmetry and decide who is left-handed and who is right-handed. Whatever SU(2) survives to lower energies we will call $\mathrm{SU}(2)_L$ by definition.

Another relevant question is what is the mass of the extra Higgs and Higgsino fields that this model has both from the (77) and $(73_{\rm bulk})$ sectors. This is a complicate issue which will depend on the detailed structure of vevs. Looking at the first, third and fourth couplings in eqs. (4.10) we see that vevs of order M_R for $(1;3)$ and $(1,1,2,\pm 1;2)$ will make massive some of the $\mathrm{SU}(2)_L$ doublets in the $(73_{\rm bulk})$ sector and also the $(1,2,2,0)_3$ fields in the (77) third complex plane. In this situation we would be left at low energies with the fields $(1,2,2,0)_1$ and $(1,2,2,0)_2$ corresponding to the first two complex planes. However, as we mentioned when we discussed quark Yukawa couplings, there are different effects which will generically mix the particles living in different complex planes in (77) sectors. Thus one also expects that these other doublets could become massive. We will thus assume that at a scale of order M_R only one set of SM doublets remains relatively light, so that they are available for $\mathrm{SU}(2)_L$ spontaneous symmetry breaking.

In addition there are the extra right-handed chiral fields $(1,1,2,\pm 1;2)$ from the $(73_{\rm bulk})$ sector. Some of these where eaten in the process of $\mathrm{SU}(2)_R$ breaking. The remaining may acquire a mass of order M_R from the second equation in (4.10) , once the scalars $(1;3)$ get a vev. The same applies to the extra colour triplets $(3,1,1,1/3;2)$ + h.c. from the $(73_{\rm bulk})$ sector. We already mentioned that some combination of them mixes with the right-handed quarks from the (77) sector. The orthogonal combination will get a mass of order M_R once the scalars $(1;3)$ get a vev. All in all, the spectrum below the M_R would thus be similar to that of the MSSM: three quark-lepton chiral multiplets and one set of $H_u + H_d$ Higgs fields.

(vii) Ramond-Ramond fields and invisible axions. We already mentioned in section 3 that in this class of orientifold models there are twisted Ramond-Ramond singlet scalars which couple to $F\tilde{F}$. As already discussed in ref. [10] , they are natural candidates to play the role of invisible axions in a model like this. Notice however that, in the absence of other charged scalar vevs, the combinations of twisted Ramond-Ramond fields coupling to the gauge groups get in fact large masses of order the string scale M_s by providing the longitudinal degrees of freedom of anomalous U(1)'s when the latter become massive. This can be easily seen from eq.(3.10).

Now, again from eq. (3.10), since $\xi_r \neq 0$, in the presence of other charged fields from the open string sector acquiring a vev and contributing to the anomalous U(1) breaking, there will be a linear combination of RR-field plus the charged field, which will be swallowed by the U(1) to become massive. The orthogonal combination will remain massless. This massless linear combination will in general couple to the gauge fields (and in particular, to QCD) in the standard axionic fashion with the gauge kinetic function taking the general form $f_\alpha = S + s_\alpha^{(ai)} M_{(ai)}$ with $s_\alpha^{(ai)}$ constant computable coefficients [28, 31] and with a decay constant of order the string scale $M_s \propto 10^{12}\,\mathrm{GeV}$, well within astrophysical limits.

(viii) Experimental signatures. The most obvious experimental implication of the present scheme is the existence of extra W_R, Z'_0 gauge bosons corresponding to the left-right symmetric gauge interactions at a scale of order 1 TeV or below. The phenomenology of $\mathrm{SU}(3)\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{U}(1)_{B-L}$ models has been extensively studied in the past, although most of the studies have tacitly assumed an SO(10) embedding of such gauge symmetry leading to the canonical value for the weak angle [15]. In addition, many studies have concentrated on a scheme in which $\mathrm{SU}(2)_R$ chiral triplets transforming like $(1,1,3,-2)+h.c.$ break the left right symmetry. At the same time these vevs could give rise to large Majorana masses for the right-handed neutrinos, leading to a see-saw structure for neutrino masses. This kind of Higgs fields do not appear in the class of models that we construct, and right-handed neutrinos are expected to become massive by combining with other singlet chiral fields, as explained above. Thus many previous studies do not directly apply to the present model.

There are a number of experimental limits on the masses of the extra gauge bosons [16]. If right-handed neutrinos are lighter than the W_R mass, the channel

$W_R \to l_R \nu_R$ is open leading to clean signatures at the Tevatron. From searches in that channel $D0$ has set [43] the limit $M_{W_R} > 720\,\mathrm{GeV}$ and CDF $M_{W_R} > 650$ [44]. If right-handed neutrinos are heavier than the W_R, this signature disappears and weaker limits coming from dijet production are obtained. $D0$ excludes the range $340 < M_{W_R} < 680\,\mathrm{GeV}$ [45] whereas CDF excludes $300 < M_{W_R} < 420$ [46]. UA2 had excluded the energy range $100 < M_{W_R} < 251$ also from the dijet signature [47]. There are stronger constraints on the W_R mass from the $K_L - K_S$ mass difference but those are much more model dependent [18]. On the other hand, for left-right symmetric models with $g_L = g_R$ like this, one can obtain limits from precision LEP-I measurements and low-energy neutral current data yielding $M_{Z'_0} > 900\,\mathrm{GeV}$, implying $M_{W_R} > 780\,\mathrm{GeV}$ in this class of models [48]. In summary, the extra gauge bosons appearing in a left-right symmetric model like this should weight more than around $800\,\mathrm{GeV}$ or so.[15] Masses of this size or a bit higher are compatible with the gauge coupling unification results in section 2. However those coupling unification results seem to prefer not very high masses for W_R and Z'_0.[16] Thus the extra left-right symmetric degrees of freedom could perhaps soon be discovered.

In addition in this class of models there is a triplication of the number of Higgs fields transforming like $(1, 2, 2, 0)$. Thus one also expects to find at energies of order $1\,\mathrm{TeV}$, charged and neutral Higgs and Higgsino fields. In fact these fields are potentially dangerous. Indeed, it is well known that in generic models with multiple Higgs $\mathrm{SU}(2)_L$ doublets, the unitary transformations which diagonalize the quark mass matrices do not necessarily diagonalize the Yukawa interactions and FCNC can in principle appear. This FCNC problem is generically present in left-right symmetric models with a low M_R scale like this. Thus to suppress sufficiently such kind of transitions, the extra Higgs fields have to be sufficiently heavy and/or the Yukawa couplings will need to have some symmetries. The question of how to evade the problem of FCNC in supersymmetric left-right models with $M_R \propto 1\,\mathrm{TeV}$ has been adressed in ref. [42, 41]. There it is shown that this problem can be avoided if there are present some extra $\mathrm{SU}(2)_R$ singlet D-type quarks mixing with the right-handed doublet quarks in the model in such a way that the physical D-quarks are mostly $\mathrm{SU}(2)_R$ singlets. In addition extra $\mathrm{SU}(2)_R$ and $\mathrm{SU}(2)_L$ doublets with non-vanishing B-L charge are also required [42]. Interestingly enough this type of extra fields and mixings are also present in the D-brane model here discussed. It would be interesting to see whether in a D-brane type of model a similar mechanism as in ref. [42] could be made operative.

[15]If the right-handed D-quarks are mostly $\mathrm{SU}(2)_R$ singlets as discussed above, W_R production is very much supressed and direct limits on the W_R mass are much weakened [41]. However that is not the case for the Z' and hence mass limits of that order are expected to still apply.

[16]It has been recently pointed out [49] that a small amount of missing invisible width in Z decays at LEP I as well as atomic parity violation experiments could perhaps indicate already the existence of some extra Z' with mass of order 500-1000 GeV.

Finally, there are also extra fields from the sector which is in charge of the breaking of the $\mathrm{SU}(2)_R$ symmetry. In order not to spoil gauge coupling unification we have seen that the required $\mathrm{SU}(2)_R$ doublets should come along with the same number of $\mathrm{SU}(2)_L$ doublets and coloured SU(3) triplets. In the D-brane explicit constructions presented in this paper such structure of extra particles is very natural and corresponds to the addition of some extra 3-branes in the bulk of the first two complex dimensions. All these extra particles are expected to have masses also around M_R or the SUSY-breaking scale. Thus new extra heavy lepton-like and quark-like extra fermions and scalars are expected to be produced at future accelerators corresponding to this sector.

In this section we have given an overall discussion of phenomenological aspects of the particular orientifold model introduced in section 3. We believe that the particular $\mathbb{Z}_3$ orientifold here discussed is tantalizingly close to the general new scheme proposed in section 2. Analogous models based on the same orientifold can be constructed which will lead to somewhat different phenomenological properties. For example, one can construct similar models by adding a second discrete Wilson line breaking e.g. the original SU(7) symmetry further. Another possibility is to consider the variant left-right symmetric discussed in the appendices.

5. Final comments and outlook

During more than 15 years the minimal supersymmetric standard model has been thoroughly studied as the best motivated extension of the standard model. During the past 10, this study became more intense due to the realization that gauge coupling unification works very well in the MSSM as compared to the standard model. In this article we have proposed an interesting alternative to the MSSM mostly motivated by the structure of D-brane models. This alternative scenario matches the success of gauge coupling unification of the MSSM, with the interesting feature that unification works only if there are three families of quarks, leptons and Higgs fields. In the present scheme several physical mass scales are unified: the string unification, susy-breaking and axion scales are one and the same and of order $M_s = 9 \times 10^{11}$ GeV. As it has been noted in several occasions, [9, 10], this intermediate scale may have important physical implications regarding neutrino masses, the strong CP problem, ultra high-energy cosmic rays, non thermal dark matter candidates, inflation, etc.

We have also presented a concrete D-brane model satisfying most of the general properties of the proposed scenario. The existence of three quark-lepton generations has an elegant explanation in these constructions: there are three generations because there are three complex compact dimensions and a $\mathbb{Z}_3$ orbifold structure.[17]

[17]This simple explanation, first encountered in the heterotic models in [17], is to be compared

The proton is stable due to a combination of discrete symmetries including lepton and R-parities and Yukawa couplings are obtained for quarks, charged leptons and neutrinos. To our knowledge this is the first example where R-parity appears so naturally in string theoretical models. This guarantees the existence of an LSP and most of the standard searches for supersymmetry. Being a left-right symmetric model, it shares many of the good properties that have been realized over several years about LR models [14, 15]. Recall, however, that the Higgs structure we find differs from the standard treatment in the sense that the fields responsible for breaking the LR group to the standard model are doublets that belong to the spectrum of the corresponding string model instead of SU(2) triplets, as often assumed in the literature [15].

One may ask if going to a left-right symmetric model is unavoidable in this class of theories. Indeed, one can construct explicit three generation $\mathbb{Z}_3$ orientifold models with the gauge group of the standard model and some examples of this type were presented in ref. [19]. However gauge coupling unification does not appear as naturally as in the left-right symmetric models here described. Furthermore, the absence of the $B-L$ gauged symmetry makes difficult to find field directions with a sufficiently stable proton. We believe that this is probably a generic property and both coupling unification and proton stability seem to point towards a left-right symmetric extension of the SM above a TeV scale.

Since the LR scale, M_R, has to be relatively close to the TeV scale in order for gauge unification to work, this scenario can be experimentally tested very soon. In the first place, the existence of new massive gauge fields for which the experimental constraints are becoming very strong can be explored. The scenario can also be tested by looking at the existence of new Higgs particles at the TeV scale, which in principle can give rise to flavour changing neutral currents, and of course with the production of supersymmetric particles.

There are several aspects that remain to be understood in this scenario. First, the Higgs scalar potential needs to be studied in order to understand the possible patterns of gauge symmetry breaking. This is also important in order to address the issue of possible FCNC transitions coming from the multi-Higgs structure. Then, the detailed structure of soft supersymmetry breaking terms needs to be addressed. This is important in order to have more information about the possible experimental signals of supersymmetry. Other issues, such as the origin of baryogenesis, may depend crucially on this knowledge. A through analysis of the general structure of neutrino masses and oscilations in this kind of scenario would also be important.

with that in Calabi-Yau compactifications of perturbative or non-perturbative heterotic vacua in which the net number of generations minus antigenerations is related to the Hodge numbers of the CY. Those tend to be quite large in general and only for very particular manifolds is small.

As for the explicit string model constructed, it would be interesting to study different variations e.g. with further symmetry breaking (Wilson lines) or locations of 3-branes which may lead to different phenomenological details. Also interesting would be to look for similar left-right symmetric models using other constructions like type-IIA orientifolds or non-perturbative heterotic orbifolds [50] leading to similar spectra. In addition the standard issues of moduli stabilization[18] and the cosmological constant[19] remain open.

It is highly remarkable that in this string model, the unification scale coincides with the preferred fundamental string scale determined by supersymmetry breaking, since a priori the two scales did not have to be related. If this turns out to be true, we might say that nature has been misleading us for many years into the belief that the unification scale was much higher. In any case, we believe the present new scenario has many appealing properties and deserves to be considered as a serious alternative to the MSSM.

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A. A variant left-right symmetric orientifold model

In reference [19], we found two LR models with three families in the study of $\mathbb{Z}_3$ orientifold models with a single Wilson line. The model discussed in the text corresponds to the first of such models. Here, for completeness, we will briefly describe the phenomenological aspects of the second model. Following the notation of section 3, the model can be defined by the following shift vector and Wilson lines:

$$\begin{aligned} V_7 &= \frac{1}{3}\ (1,1,1;-1,-1;0,0;0;1,1,1,1,1,1,1,1)\,, \\ W &= \frac{1}{3}(1,1,1;1,1;1,1;0;0,0,0,0,0,0,0,0)\,, \end{aligned} \tag{A.1}$$

which leads to a **77** sector group $\mathrm{U}(3)\times\mathrm{U}(2)_L\times\mathrm{U}(2)_R\times[\mathrm{SO}(2)\times\mathrm{U}(8)]'$. We must also add four 3-branes at each of the six points at (a,i) ($a=0,1$ and $i=0,1,2$). The gauge group on each of the six 3-branes will then be $\mathrm{Sp}(2)\times\mathrm{U}(1)$. Similar to the previous model, the anti-branes live in the bulk and they total 24 (equals to the

[18] For some recent attempts to understand the stabilization of large radii see e.g. ref. [51, 19].

[19] For some recent ideas about the cosmological constant problem in the D-brane context see e.g. refs. [37, 52].

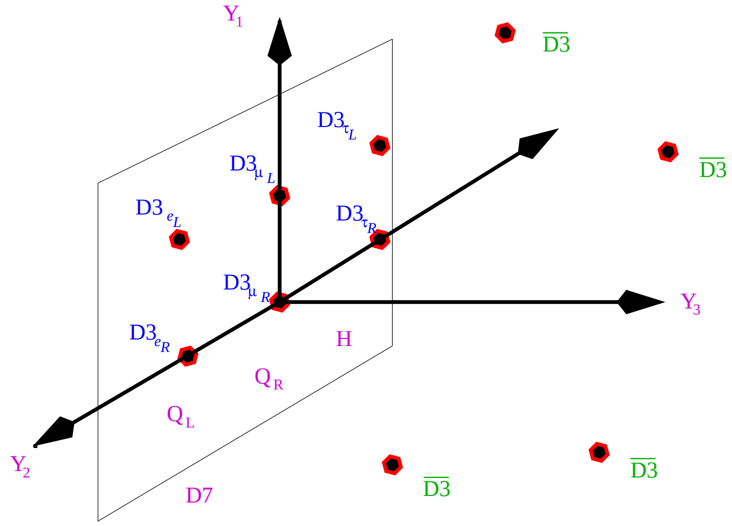

Figure 7: D-brane configuration of the variant left-right model. The gauge group, Quarks and Higgs fields live in the worldvolume for the D7-brane whereas left-handed leptons are located at three fixed points and right-handed leptons in other three fixed points in the complex Y_1, Y_2 dimensions. There is no need of bulk $\mathrm{D3}_{\mathrm{bulk}}$ branes in this case for $\mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ breaking. The anti-3-branes live away in the bulk of the third (Y_3) complex dimension.

total number of 3-branes). Again, in the figure we only show 4 of them, all the other ones are located at the images points under the $\mathbb{Z}_3 \times \mathbb{Z}_2$ orientifold symmetry.

Since there are no branes at the $(2, i)$ points we will have three non-anomalous and seven anomalous U(1)'s.

We choose the linear combinations

$$
\begin{aligned}
Q_{B-L} &= -\frac{2}{3} Q_3 + Q_L - Q_R \\
Q^L_{non} &= \frac{Q_L}{2} + \frac{Q_8}{8} - \frac{2}{3} \sum_i Q_{n_0^i} - \frac{1}{3} \sum_i Q_{n_1^i} \\
Q^R_{non} &= \frac{Q_R}{2} + \frac{Q_8}{8} - \frac{1}{3} \sum_i Q_{n_0^i} - \frac{2}{3} \sum_i Q_{n_1^i} \\
Q^8_A &= Q_8 \\
Q^r_A &= Q_{n_a^i} \,,
\end{aligned}
\tag{A.2}
$$

with $a = 0, 1$ and $i = 0, 1, 2$. The first 3 are non anomalous whereas the last seven are anomalous. The spectrum with the corresponding charges under the U(1)'s is shown in the table 4. We present the quantum numbers under the $\mathrm{U}(1)^{10}$ groups. Q_3, Q_L, Q_R and $Q_8, Q^A_{n_0^i}, Q^A_{n_1^i}$. The last seven can be chosen as independent anomalous combinations. Q_{B-L}, Q_L^{non} and Q_R^{non} are the non anomalous charges.

Mixed anomalies can be computed as we did for the model in the text and they are exactly cancelled. The anomaly matrix is:

$$T_{IJ}^{\alpha\beta} = \begin{pmatrix} 0 & 9 & -9 & 0 & 3 & -3 & -3 & 3 \\ 6 & 0 & -6 & 0 & 2 & 0 & -2 & 0 \\ 6 & -6 & 0 & 0 & 0 & 2 & 0 & -2 \\ 0 & 0 & 0 & 24 & 8 & 8 & -8 & -8 \\ 1 & 1 & -2 & 1 & 3 & 0 & -3 & 0 \\ 1 & -2 & 1 & 1 & 0 & 3 & 0 & -3 \end{pmatrix}, \tag{A.3}$$

where the rows correspond to the 10 U(1)'s ordered as in table 4. The columns correspond to the groups: SU(3), $\mathrm{SU}(2)_L$, $\mathrm{SU}(2)_R$, $\mathrm{SU}(8)'$, $\mathrm{U}(1)^i_0$, $\mathrm{U}(1)^i_1$ and $\mathrm{Sp}(2)^i_0$, $\mathrm{Sp}(2)^i_1$ respectively.

Looking at the spectrum of this model we can immediately see that contrary to the model in the text, there are enough fields on this model to break the group to the standard model one, *without* the need to add extra branes in the bulk. For instance, if we give a nonvanishing vev to the Sp(2) doublets in the **33** sectors we are left with a spectrum in the visible sector consisting only of the three families of quarks, leptons and Higgs fields plus extra matter with *exactly* the same quantum numbers as the 'bulk' matter fields of the model in the text (equation (4.2)), although in three copies.[20] As we have mentioned before, the quantum numbers of those fields are such that they do not modify the analysis of the gauge coupling unification. Therefore we conclude that, quite remarkably, the present model also shares the good properties about gauge coupling unification as the model presented in the text. In order to cancel the D-terms generated by the vevs of the Sp(2) doublets one has to give vevs to other doublets to compensate. The simplest option is to take flat directions for the hidden sector fields such as $(2,\bar{8})'$ of the $\mathbf{37}_{0,1}$ sectors as well as to the $(1,28)'$ of the **77** sector. This allows in addition to generate nonvanishing charged lepton and neutrino masses from couplings such as[21]

$$\begin{aligned} &(1,2,2)(1,2,1;2)^*(1,1,2;2)^* \left\langle (1,\bar{8};2)'(1,\bar{8};2)'(1,28)' \right\rangle \left\langle (1,28)'^4 \right\rangle \\ &(1,1,2;2)^*(1,1,2)(1,1)_{0,i} \left\langle (1,2)_{0,i}(1,\bar{8};2)'(1,\bar{8};2)'(1,28)' \right\rangle , \end{aligned} \tag{A.4}$$

where the * on the Sp(2) doublets stand for the component of the doublet that remains massless after the doublet in the **33** sector got a nonvanishing vev and the

[20] There is an ambiguity about which doublets are assigned to be the leptons and which would be extra matter fields. In fact we can see that the fields $(1,2,1;2)$ in the $\mathbf{37}_0$ sector and one of the $(1,2,1;1)$ of the $\mathbf{37}_1$ sector have the same quantum numbers as the leptons. One of the components of the $(1,2,1;2)$ field will get a mass after the Sp(2) doublet in the $\mathbf{33}_0$ sector gets a vev, the remaining component is what we will identify as the physical leptons. The other choice does not have the appropriate Yukawa couplings.

[21] Yukawa couplings giving rise to quark masses are identical to the model in the text since they only include couplings among fields of the **77** sector.

Matter	Q_3	Q_L	Q_R	Q_8	Q_{B-L}	Q_L^{non}	Q_R^{non}	$Q^A_{n^i_0}$	$Q^A_{n^i_1}$
77 sector									
$(3,2,1)$	1	1	0	0	1/3	1/2	0	0	0
$(\bar{3},1,2)$	-1	0	1	0	-1/3	0	1/2	0	0
$(1,2,2)$	0	-1	-1	0	0	-1/2	-1/2	0	0
$(2,\bar{8})'$	0	0	0	-1	0	-1/8	-1/8	0	0
$(1,28)'$	0	0	0	2	0	1/4	1/4	0	0
$\mathbf{37}_0$ sector									
$(3,1,1;1)$	1	0	0	0	-2/3	-2/3	-1/3	1	0
$(\bar{3},1,1;2)$	-1	0	0	0	2/3	0	0	0	0
$(1,1,2;1)$	0	0	1	0	-1	2/3	1/2+1/3	-1	0
$(1,1,2;1)$	0	0	-1	0	1	2/3	-1/2+1/3	-1	0
$(1,2,1;1)$	0	1	0	0	1	1/2 -2/3	-1/3	1	0
$(1,2,1;2)$	0	-1	0	0	-1	-1/2	0	0	0
$(2,1;1)'$	0	0	0	0	0	2/3	1/3	-1	0
$(1,8;1)'$	0	0	0	1	0	1/8 -2/3	1/8-1/3	1	0
$(1,\bar{8};2)'$	0	0	0	-1	0	-1/8	-1/8	0	0
$\mathbf{37}_1$ sector									
$(3,1,1;2)$	1	0	0	0	-2/3	0	0	0	0
$(\bar{3},1,1;1)$	-1	0	0	0	2/3	-1/3	-2/3	0	1
$(1,1,2;1)$	0	0	1	0	-1	-1/3	1/2-2/3	0	1
$(1,1,2;2)$	0	0	-1	0	1	0	-1/2	0	0
$(1,2,1;1)$	0	1	0	0	1	1/2 +1/3	2/3	0	-1
$(1,2,1;1)$	0	-1	0	0	-1	-1/2+1/3	2/3	0	-1
$(2,1;1)'$	0	0	0	0	0	1/3	2/3	0	-1
$(1,8;1)'$	0	0	0	1	0	1/8 -1/3	1/8-2/3	0	1
$(1,\bar{8};2)'$	0	0	0	-1	0	-1/8	-1/8	0	0
$\mathbf{33}_0$ sector									
$(1,2)$	0	0	0	0	0	2/3	1/3	-1	0
$(1,1)$	0	0	0	0	0	-4/3	-2/3	2	0
$\mathbf{33}_1$ sector									
$(1,2)$	0	0	0	0	0	1/3	2/3	0	-1
$(1,1)$	0	0	0	0	0	-2/3	-4/3	0	2

Table 4: Spectrum of the variant LR model including the charges under all U(1) symmetries. We have written the charges coming from the 3-brane sector as a single column with the understanding that for instance in the **37** sector, each of the three copies have that charge under one of the three U(1)'s and zero under the other two

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subindex a, i indicates that these are fields from the $\mathbf{37}_a$ sector. Notice that in this case the right-handed neutrinos become massive by combining with singlets in the **33** sectors. These couplings are allowed by all the gauge symmetries including the anomalous U(1)'s, therefore, contrary to the model in text, there is no need to introduce insertions of twisted vertex operators in order to generate lepton masses.

Interestingly enough, this flat direction leaves a remaining $\mathbb{Z}_2$ symmetry consisting in changing sign to all SU(8) octets and simultaneously changing sign to all particles in sectors **37**. It is easy to see that this discrete symmetry combined with the two $\mathbb{Z}_2$'s remaining after breaking the LR SU(2)'s imply a residual discrete baryon $\mathbb{Z}_2$ parity making the proton absolutely stable, similar to the model in the text. Unlike that model however, lepton number is not conserved since couplings like QDL are permitted, although suppressed by powers of (M_R/M_s) since they violate $B-L$.

Combining these SU(8) flat directions with those of the Sp(2) doublets does not preserve gauge coupling unification since couplings such as

$$(3,1,1;2)^*(\bar{3},1,1;2)^*\,\langle(1,\bar{8};2)(1,\bar{8};2)(1,28)\rangle \tag{A.5}$$

will in principle give masses to the extra triplets. Thus an alternative which does preserve gauge coupling unification would be combining the SU(8) flat directions with those of singlets in the **33** sectors (not giving vevs to the Sp(2) doublets). This maintains gauge coupling unification with the good properties for lepton masses. In this case giving a mass to all the extra triplets needs the insertion of twisted vertex operators. In summary this model is very similar to the one presented in the text (although probably a bit more complicated) and may deserve further exploration.

B. Anomalous U(1)'s and non-renormalizable couplings in the $\mathbb{Z}_3$ orientifold

In this appendix we will argue that non-renormalizable couplings violating anomalous U(1) charges are in general expected in orientifolds similar to the ones considered in the present article. More specifically heterotic/type-I duality seems to indicate that this is the case. Consider in particular the standard $D=4$, $N=1$, $\mathbb{Z}_3$ orientifold first constructed in ref. [22]. The underlying orbifold is exactly the same than the one considered in the present paper, the only difference being the absence of anti-branes and Wilson lines. The model has 32 9-branes and a gauge group $SU(12)\times SO(8)\times U(1)_A$. In the open string (99) sector there are charged fields transforming like:

$$3(\overline{66},1)_{+2}+3(12,8)_{-1}\,, \tag{B.1}$$

where the subindex denotes the $U(1)_A$ charge. It is easy to check that the $U(1)_A$ interaction is anomalous. In addition there are 27 chiral singlets $(1,1)_0$ coming from the twisted closed string sector. The latter are singlets under $U(1)_A$ but transform non-linearly under that anomalous symmetry, and it is this transformation which cancels the $U(1)_A$ anomalies by means of a generalized Green-Schwarz mechanism. At the same time the $U(1)_A$ becomes massive. Now, SU(12) invariance allows for a non-renormalizable superpotential coupling of the type $(\overline{66},1)^6_{+2}$ by contraction with the 12-index antisymmetric tensor. However, such a coupling would violate $U(1)_A$ conservation in 12 units. Since this $U(1)_A$ symmetry is broken one may suspect that such non-renormalizable coupling may however exist if one moves away from the orbifold limit, i.e. if the twisted NS-NS singlet fields get a vev.

Heterotic/type-I duality seems to indicate that is the case. The above orientifold has a heterotic dual [22, 53, 25] which is a $\mathbb{Z}_3$ orbifold compactification of the SO(32) heterotic string. This model has the same gauge group and the charged particles in the untwisted sector are identical to those in the (99) sector of the type-I model. There are also 27 twisted chiral fields transforming like $(1,1)_{-4}$ under the gauge group, which should be identified as duals of the 27 twisted fields of the type-I dual. Notice however that, unlike their type-I counterparts, the heterotic singlets are charged under the anomalous $\mathrm{U}(1)_A$. Now, one can convince oneself that in this heterotic $\mathbb{Z}_3$ orbifold the following non-renormalizable couplings do in general exist:

$$(1,1)^3_{-4} \times (\overline{66},1)^6_{+2}\,. \tag{B.2}$$

This coupling preserves the $\mathrm{U}(1)_A$ symmetry. One can also check that this coupling preserves the necessary selection rules in order to be present. In particular, it respects the $\mathbb{Z}_3$ point group symmetry (the twisted fields appear to the third power). Also the H-momentum conservation is obeyed. To check this it is enough to consider the vertex operators associated to the $(\overline{66},1)_{+2}$ written in the "0-picture". In addition, as pointed out in ref. [53] , the FI-term present in this heterotic model forces the singlets $(1,1)_{-4}$ to get non-vanishing vevs. Thus away from the orbifold point the heterotic model will present $(\overline{66},1)^6_{+2}$ couplings, effectively violating the anomalous $\mathrm{U}(1)_A$ symmetry. This strongly suggests that similar couplings will also be present in the dual type-I model as long as one stays away from the orbifold point, i.e., as long as the twisted $NS-NS$ closed string singlets have non-vanishing vevs. This is indeed the case considered in the models in the present article.

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