Standard Model Building from Intersecting D-branes

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ABSTRACT

We provide a general overview of the current state of the art in four dimensional three generation model building proposals - using intersecting D-brane toroidal compactifications [without fluxes] of IIA, IIB string theories - which have only the SM at low energy. In this context, we focus on these model building directions, where non-supersymmetric constructions - based on the existence of the gauge group structure $SU(3)_c \times SU(2)_L \times U(1)_Y$, Pati-Salam $SU(4)_C \times SU(2)_L \times SU(2)_R$, $SU(5)$ and flipped $SU(5)$ GUTS - appear at the string scale $M_s$. These model building attempts are based on four dimensional compactifications that use orientifolds of either IIA theory with D6-branes wrapping on $T^6$, $T^6/Z_3$ and recently on $T^6/Z_3 \times Z_3$ or of IIB theory with D5-branes wrapping on $T^4 \times C/Z_N$. Models with D5-branes are compatible with the large extra dimension scenario and a low string scale that could be at the TeV; thus there is no gauge hierarchy problem in the Higgs sector. In the case of flipped $SU(5)$ GUTS - coming from $T^6/Z_3$ - the special build up structure of the models accommodates naturally a see-saw mechanism and a new solution to the doublet-triplet splitting problem. Baryon number is a gauged symmetry and thus proton is naturally stable only in models with D5 branes or in models with D6-branes wrapping toroidal orientifolds of type IIA. In the rest of the constructions, proton is stable due to the existence of a high $M_s$ that suppress gauge mediating baryon violating couplings. Extra beyond the SM gauge group $U(1)$’s that may be present at $M_s$ and have non-zero couplings to RR fields, become massive by the existence of a generalized Green-Schwarz mechanism and by the existence of either tachyonic singlets receiving a vev or gauge neutral singlets becoming massless at N=1 supersymmetric intersections. Models with intersecting D-branes accommodate naturally right handed neutrinos. Finally, we present new RR tadpole solutions for the 5- and 6- stack toroidal orientifold models of type IIA which have only the Standard Model with right handed neutrinos at low energy.
1 Preliminaries - D6/D5-brane model building with only the SM at Low Energy

At present string theory is our only candidate theory that could unify gravity with the rest of the fundamental interactions in a perturbative framework. In this respect, in the last three years, string constructions based on D-branes intersecting at angles - intersecting brane worlds (IBW's) for short - received a lot of attention [1] - [25], as on these constructions, for the first time in string theory, it became possible \(^1\) to construct four dimensional non-supersymmetric (non-susy) intersecting D6-brane models that have only the Standard model (SM) gauge group and chiral spectrum - with right handed neutrinos \(\nu_R\)'s - at low energies [1, 2, 3]. For N=1 supersymmetric model building attempts in IBW's see for a partial list of references [6, 7]. In IBW's chiral matter get localized as open strings stretching in the intersections between intersecting D-branes [11]. Subsequently, the chiral spectrum of the model may be obtained by solving simultaneously the intersection constraints coming from the existence of the different sectors and the RR tadpole cancellation conditions.

Thus e.g. if a stack of parallel D-branes gets associated to the \(U(3)\) and another one to the \(U(2)\) gauge groups, the chiral fermion localized in their intersection represents a quark; see fig. (1). Hence by considering \(a\) stacks of D-brane configurations with \(N_a, a = 1, \ldots, N\), parallel branes in each stack, we get the gauge group configuration \(U(N_1) \times U(N_2) \times \cdots \times U(N_a)\). Each \(U(N_i)\) factor will also give rise to an \(SU(N_i)\) charged under the associated \(U(1_i)\) gauge group factor that appears in the decomposition \(SU(N_a) \times U(1_a)\).

The initial constructions of IBW's exhibiting the SM at low energy, were based on a background of D6-branes intersecting at angles in 4D toroidal orientifolds [1] of [T-dual to models with magnetic deformations [5]. See also [19.] type IIA theory and exhibit proton stability as baryon number is a gauged symmetry. The primary common phenomenological characteristics of the four stack D6-models of [1] and the five and six SM’s of [2] and [3] respectively - emphasizing that there are no D6-brane models with only the SM at low energy using constructions with more than six stacks of D6-branes at the string scale \(M_s\) - are:

- the prediction of the existence of the SM chiral spectrum together with \(\nu_R\)'s,
- proton stability; conservation of lepton number - the models admit Dirac terms for the neutrinos - their masses appear as a result of the existence of particular Yukawa couplings associated with the breaking of the chiral symmetry.

\(^1\)See also [27] for other reviews on the subject.
Figure 1: Open strings stretching between intersecting branes get identified as fermionic matter. A left handed quark makes its appearance.

Additionally, the SM’s of [2] and [3] exhibit a new phenomenon - not found in the SM’s of [1], namely the prediction of the existence of N=1 supersymmetric (SUSY) partners of $\nu_R$’s, the s$\nu_R$’s. Thus even though the models of [2, 3] are non-susy, they have particular N=1 SUSY partners, the s$\nu_R$’s. There are two ways that vacua with only the SM have been produced in IBW’s: a) with all matter accommodated in bifundamental representations (reps) $[1, 2, 3] (N, \bar{N})$ of $U(N_a) \times U(N_b)$. In this case, we have the clear advantage that all matter Yukawa’s may be realized, b) when some of the matter fields appear in symmetric and antisymmetric reps of $U(N_a)$ [22, 23, 25]. In this case some of the SM matter mass couplings may be missing.

Take for example the general picture with D($3+l_a$)-branes wrapping $l$-cycles with $(n_a^l, m_a^l)$, $l = 1, 2, 3$ the wrapping numbers of a brane along each of the $(T^2_l)$ torus. Thus we allow the six-torus to wrap factorized 2-cycles, so we can unwrap the l-cycle into products of 1-cycles. The definition of the homology of the l-cycles as

$$[\Pi_l] = \prod_{i=1}^{l} (n_a^l[a_i] + m_a^l[a_i]) \quad (1.1)$$

The number of fermions $I_{ab}$ localized in an intersection is described by the intersection number $I_{ab} = [\Pi_a][\Pi_b]$. Major role in these constructions is played by the satisfaction of the so called RR tadpole cancellation conditions (TCC), which in low energies is seen as the cancellation of cubic gauge anomalies. TCC’s can be alternatively be expressed as the cancellation of the RR charge
in homology

\[ \sum_a N_a[\Pi_a] = 0 \]  

(1.2)

We note that while in N=1 susy constructions NSNS automatically cancel, in non-susy constructions uncancelled NS tadpoles remain; even though it remains an open issue whether or not they can be cancelled in higher orders of perturbation theory. The NS tadpoles do not affect the low energy spectrum of the models but they rather imply the instability of the models, in the present order of perturbation theory, in a flat background.

In the absence of a principle for selecting a particular string vacuum, we select to examine whether or not it is possible to obtain non-susy models with just the SM at low energy by examining different 4D string compactifications. The method that has been used \[1, 2, 3\] follows a bottom-up approach. In particular we embed the localization of SM fermions at particular intersections - as it is seen in tables \[1, 2, 3\] - to different four dimensional (4D) compactifications. As we will see in sections 2 and 3, for constructions that involve a SM-like configuration at \( M_s \) - e.g. \( U(3) \times U(2) \times U(1)^n \), \( n= 2,3,4 \) - these tables were embedded successfully to 4D models with D6-branes wrapping on IIA compactified on a toroidal orientifold \[1, 2, 3\] or to 4D models with D5-branes wrapping \( T^4 \times C/Z_N \) \[8, 9\].

### Matter Fields

| Matter Fields | Intersection | \( Q_a \) | \( Q_b \) | \( Q_c \) | \( Q_d \) | \( Y \) |
|--------------|--------------|-----------|-----------|-----------|-----------|-------|
| \( Q_L \)    | (3, 2)       | \( I_{ab} \) | 1         | -1        | 0         | 0     |
| \( q_L \)    | 2(3, 2)      | \( I_{ab^*} \) | 1         | 1         | 0         | 0     |
| \( U_R \)    | 3(3, 1)      | \( I_{ac^*} \) | -1        | 0         | 1         | 0     |
| \( D_R \)    | 3(3, 1)      | \( I_{bc^*} \) | 1         | 0         | -1        | 0     |
| \( L \)      | 3(1, 2)      | \( I_{bd^*} \) | 0         | -1        | 0         | 1     |
| \( N_R \)    | 3(1, 1)      | \( I_{cd^*} \) | 0         | 0         | -1        | -1   |
| \( E_R \)    | 3(1, 1)      | \( I_{cd^*} \) | 0         | 0         | -1        | 0     |

Table 1: Low energy fermionic spectrum of the four stack string scale \( SU(3)_C \otimes SU(2)_L \otimes U(1)_a \otimes U(1)_b \otimes U(1)_c \otimes U(1)_d \), D5-brane model together with its \( U(1) \) charges.

The hypercharge operator of the SM’s appearing in these tables is defined as a linear combination of the \( U(1)_a, U(1)_c, U(1)_d, U(1)_e, U(1)_f \) gauge groups for the four-, five-, six-stack models respectively:

\[ Y_{Table \ [1]} = \frac{1}{6} U(1)_a - \frac{1}{2} U(1)_c - \frac{1}{2} U(1)_d \]  

(1.3)

\[ Y_{Table \ [2]} = \frac{1}{6} U(1)_a - \frac{1}{2} U(1)_c - \frac{1}{2} U(1)_d - \frac{1}{2} U(1)_e \]  

(1.4)
Table 2: Low energy fermionic spectrum of the five stack string scale $SU(3)_C \otimes SU(2)_L \otimes U(1)_a \otimes U(1)_b \otimes U(1)_c \otimes U(1)_d \otimes U(1)_e$, D5-brane model together with its $U(1)$ charges.

\[
Y_{\text{Table } 3} = \frac{1}{6} U(1)_a - \frac{1}{2} U(1)_c - \frac{1}{2} U(1)_d - \frac{1}{2} U(1)_e - \frac{1}{2} U(1)_f . \tag{1.5}
\]

Moreover the first non-SUSY constructions of string GUTS which have only the SM at low energy \(^2\), were constructed in [10], based on the Pati-Salam structure $SU(4)_C \times SU(2)_L \times SU(2)_R$ at $M_s$. These models may be described in section 3.

## 2 Only the SM at low energy from models compatible with the large extra dimension scenario

In this class of models [8] the general picture involves D5\(_a\)-branes wrapping 1-cycles $(n_a^i, m_a^i)$, $i = 1, 2$ along each of the $i$th-$T^2$ torus of the factorized $T^4$ torus, namely $T^4 = T^2 \times T^2$. Thus we allow the four-torus to wrap factorized 2-cycles, so we can unwrap the 2-cycle into products of two 1-cycles, one for each $T^2$. The definition of the homology of the 2-cycles as

\[
[\Pi_a] = \prod_{i=1}^2 (n_a^i[a_i] + m_a^i[b_i]) \tag{2.1}
\]

defines consequently the 2-cycle of the orientifold images as

\[
[\Pi_{a^*}] = \prod_{i=1}^2 (n_a^i[a_i] - m_a^i[b_i]). \tag{2.2}
\]

\(^2\)These models were also based on the intersecting D6-backgrounds of [4].
Table 3: Low energy fermionic spectrum of the six stack string scale $SU(3)_C \otimes SU(2)_L \otimes U(1)_a \otimes U(1)_b \otimes U(1)_c \otimes U(1)_d \otimes U(1)_e \otimes U(1)_f$, D5-brane model together with its $U(1)$ charges.

We note that because of the $\Omega R$ symmetry each D5$_a$-brane 1-cycle, must be accompanied by its $\Omega R$ orientifold image partner $(n_i^a, -m_i^a)$, $n, m \in \mathbb{Z}$. In addition, because of the presence of discrete NS B-flux \[4\], the tori involved are not orthogonal but tilted. Hence, the wrapping numbers become the effective tilted wrapping numbers,

$$(n^i, m = \tilde{m}^i + b^i \cdot n^i/2); \quad n, \tilde{m} \in \mathbb{Z}, \quad b^i = 0, 1/2,$$

where semi-integer values are allowed for the m-wrappings.

Let us discuss the effect of the orbifold action on the open string sectors. The $\mathbb{Z}_N$ orbifold twist in the third complex dimension is generated by the twist vector $v = \frac{1}{N}(0, 0, -2, 0)$, which is fixed by the requirements of modular invariance and for the variety to be spin. Subsequently the $\mathbb{Z}_N$ action is embedded in the $U(N)$ degrees of freedom emanating from the $a$th stack of D5-branes, through the unitary matrix in the form

$$\gamma_{\omega, a} = diag \left( 1_{N_a^0}, \quad \alpha 1_{N_a^1}, \ldots, \quad \alpha^{N-1} 1_{N_a^{N-1}} \right),$$

with $\sum_{i=0}^{N-1} N_i^a = N_a$ and $\alpha \equiv \exp(2\pi i/N)$. In the presence of $\Omega R$ orientifold action, we include sectors where its brane is accompanied by is orientifold image. Lets us denote the $\Omega R$ image of the brane D5$_a$-brane by D5$_{a\ast}$. Hence if the D5$_a$-brane is described by the matrix

$$(n^1, m^1) \otimes (n^2, m^2)$$

$$\gamma_{\omega, a} = diag \left( 1_{N_a^0}, \quad \alpha 1_{N_a^1}, \ldots, \quad \alpha^{N-1} 1_{N_a^{N-1}} \right),$$

with semi-integer values are allowed for the m-wrappings.
Figure 2: Assignment of SM embedding in configurations of four stacks of D5 branes depicted by the ‘reflected’ $Z_3$ quiver diagrams. At low energy we get only the SM. Note that $\tilde{\alpha} = \alpha^{-1}$.

The D5$_{a^*}$ is given by

$$\gamma_{\omega, a^*} = \text{diag} \left( 1_{N_0^a}, \alpha^{N-1} 1_{N_1^a}, \ldots, \alpha 1_{N_{N-1}^a} \right), \quad (2.6)$$

The RR tadpoles for the D5-branes at the $Z_N$ orbifold singularity are given by

$$\begin{align*}
c_k^2 \sum_a n_a^1 n_a^2 & \left( \text{Tr} \gamma_{k,a} + \text{Tr} \gamma_{k,a^*} \right) = 16 \sin \left( \frac{2k}{N} \right) \\
c_k^2 \sum_a m_a^1 m_a^2 & \left( \text{Tr} \gamma_{k,a} + \text{Tr} \gamma_{k,a^*} \right) = 0 \\
c_k^2 \sum_a n_a^1 n_a^2 & \left( \text{Tr} \gamma_{k,a} - \text{Tr} \gamma_{k,a^*} \right) = 0 \\
c_k^2 \sum_a m_a^1 m_a^2 & \left( \text{Tr} \gamma_{k,a} - \text{Tr} \gamma_{k,a^*} \right) = 0
\end{align*} \quad (2.7)$$

with $c_k^2 = \sin((2\pi k)/N)$. The presence of a non-zero term in the first tadpole condition should be interpreted as a negative RR charge induced by the presence of an O5-plane. We note that the first of twisted tadpole conditions can be also written as

$$\sum_a n_a^1 n_a^2 \left( \text{Tr} \gamma_{2k,a} + \text{Tr} \gamma_{2k,a^*} \right) = \frac{16}{\alpha^k + \alpha^{-k}}. \quad (2.8)$$

In this subsection we examine the derivation of exactly the SM at low energies from the embedding of the four stack SM structure of table (1) in a $Z_3$ quiver of $Q1$-type seen in figure (2).
The solutions satisfying simultaneously the intersection constraints and
the cancellation of the RR twisted crosscap tadpole cancellation constraints
are given in parametric form in table (4). The multiparameter RR tadpole
solutions appearing in table (4) represent deformations of the D5-brane
branes, of table (1), intersecting at angles, within the same homology class of
the factorizable two-cycles. The solutions of table (4) satisfy all tadpole equations,
in (2.7), but the first. The latter reads:

\[ 9n_a^1 - \frac{1}{\beta^1} + \bar{\epsilon} n_d^1 + \frac{2\epsilon h N_h}{\beta^1} = -8. \]  

(2.9)

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

\[ N_h \ (\epsilon_h/\beta_1, 0) \ (2, 0) 1_{N_h}. \]  

(2.10)

The presence of an arbitrary number of \( N_h \) D5-branes, which give an extra
\( U(N_h) \) gauge group, does not contribute to the rest of the tadpoles and in-
tersection constraints. Thus in terms of the low energy theory no new chiral
matter is generated and only the SM spectrum appears. The analysis of the
\( U(1) \) anomalies in the models also shows that only the SM hypercharge sur-
vives massless at low energies. We will not present the details of this analysis
as they can be found in detail in [9].

In the present models, there are two dimensions transverse to the configu-
rations space that the D5-branes wrap. As the models are non-susy in order not

| \( N_i \) | \( (n^1, m^1) \) | \( (n^2, m^2) \) | \( (n^3, m^3) \) |
|---|---|---|---|
| \( N_a = 3 \) | \( (n^1_a, \epsilon \beta^1) \) | \( (3, \frac{1}{2} \epsilon \epsilon) \) | \( 1_3 \) |
| \( N_b = 2 \) | \( (1/\beta_1, 0) \) | \( (1, \frac{1}{2} \epsilon \epsilon) \) | \( \alpha^2 1_2 \) |
| \( N_c = 1 \) | \( (1/\beta_1, 0) \) | \( (0, \epsilon \epsilon) \) | \( \alpha \) |
| \( N_d = 1 \) | \( (n^1_d, 3 \epsilon \beta^1) \) | \( (\bar{\epsilon}, -\frac{1}{2} \epsilon) \) | \( 1 \) |
| \( N_h \) | \( (\epsilon_h/\beta^1, 0) \) | \( (2, 0) \) | \( 1_{N_h} \) |

Table 4: General tadpole solutions for the four-stack \( Q_1 \) type quiver of intersecting
D5-branes, giving rise to exactly the standard model gauge group and observable
chiral spectrum at low energies. The solutions depend on two integer parameters,
\( n^1_a, n^1_d \), the NS-background \( \beta^1 = 1 - b_i \), which is associated to the presence of the
NS B-field by \( b_i = 0, 1/2. \) and the phase parameters \( \epsilon = \bar{\epsilon} = \pm 1 \), as well as the CP
phase \( \alpha \).
to face the gauge hierarchy problem (GHP) in the Higgs sector of the models the string scale should be at the order of the TeV. This may be realized by lowering the string scale while keeping the Planck scale large, by increasing the volume of the extra 2-dimensional space being transverse to the D5-branes. In this way, GHP is solved in consistency with the well known scenario of \[12\].

3 Building string vacua with D6-branes and only the SM at low energy

Historically, the first construction of models which have the SM at low energy - with all matter in bifundamental representations - appeared in \[1\], while five and six stack models have been studied in \[2, 3\] respectively. In this section we will describe the derivation of the SM at low energy using five (5) stacks of D6-branes at \(M_s\). These models have been described before in \[2\] but in this section we will present a new solution to the RR tadpoles not found in \[2\]. We will discuss in some detail the 5-stack SM solutions, while by the end of this section we will simply present the alternative RR tadpole cancellation solution to the 6-stack SM’s that appear in \[3\].

Next, we turn our attention to the construction of the standard models. It is based on type I string with D9-branes compactified on a six-dimensional orientifolded torus \(T^6\), where internal background gauge fluxes on the branes are turned on. If we perform a T-duality transformation on the \(x^4, x^5, x^6\), directions the D9-branes with fluxes are translated into D6-branes intersecting at angles. Note that the branes are not parallel to the orientifold planes. Furthermore, we assume that the D6\(_a\)-branes are wrapping 1-cycles \((n^i_a, m^i_a)\) along each of the ith-\(T^2\) torus of the factorized \(T^6\) torus, namely \(T^6 = T^2 \times T^2 \times T^2\). That means that we allow our torus to wrap factorized 3-cycles, that can unwrap into products of three 1-cycles, one for each \(T^2\). We define the homology of the 3-cycles as \([\Pi_a] = \prod_{i=1}^3(n^i_a[a_i] + m^i_a[b_i])\) while we define the 3-cycle as in \((2.1)\) for the orientifold images as \([\Pi_{a^*}] = \prod_{i=1}^3(n^i_a[a_i] - m^i_a[b_i])\)

In order to build the SM model structure a low energies, we consider five stacks of D6-branes giving rise to their world-volume to an initial gauge group \(U(3) \times U(2) \times U(1) \times U(1) \times U(1)\) or \(SU(3) \times SU(2) \times U(1)_a \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e\) at the string scale. Also, we consider the addition of NS B-flux \[4\], such that the tori involved are not orthogonal, and leading to effective tilted wrapping numbers as in \((2.3)\), thus allowing semi-integer values for the m-wrapping numbers.

Because of the \(\Omega R\) symmetry, where \(\Omega\) is the worldvolume parity and \(R\) is the reflection on the T-dualized coordinates, \(T(\Omega R)T^{-1} = \Omega R\) each D6\(_a\)-brane
1-cycle, must have its $\Omega R$ image partner $(n^i_a, -m^i_a)$.

In the toroidal D6-models there are a number of different sectors contributing to the chiral spectrum. To establish notation we denote the action of $\Omega R$ on a sector $a, b$, by $a^*, b^*$, respectively. We recognize the following sectors:

- The sector containing open strings stretching between the D6$_a$ and D6$_b$ branes. Under the $\Omega R$ symmetry this sector is mapped to its image. The number of chiral fermions, namely $I_{ab}$, transforms in the bifundamental representation $(N_a, \bar{N}_a)$ of $U(N_a) \times U(N_b)$, and their multiplicity is given by

$$I_{ab} = [\Pi_a] \cdot [\Pi_b] = (n^1_a m^1_b - m^1_a n^1_b) \left( n^2_a m^2_b - m^2_a n^2_b \right) \left( n^3_a m^3_b - m^3_a n^3_b \right), \quad (3.1)$$

where $I_{ab}$ is the intersection number of the wrapped cycles. Note from the sign of $I_{ab}$ intersection, we get the chirality of the fermions, where $I_{ab} > 0$ by convention denotes left handed fermions.

- The sector containing open string stretching between the brane $a$ and the orientifold image of the $b$-brane. In this sector chiral fermions transforming into the $(N_a, N_b)$ representation with multiplicity given

$$I_{ab^*} = [\Pi_a] \cdot [\Pi_{b^*}] = - (n^1_a m^1_b + m^1_a n^1_b) \left( n^2_a m^2_b + m^2_a n^2_b \right) \left( n^3_a m^3_b + m^3_a n^3_b \right). \quad (3.2)$$

Similar conventions to $I_{ab}$ hold for the sign of $I_{ab^*}$.

- The sector containing open strings stretching between the brane $a$ and its orientifold image, namely $aa^*$. In this sector the invariant intersections contribute $8m^1_a m^2_a m^3_a$ fermions in the antisymmetric representation and the non-invariant intersections add $4m^1_a m^2_a m^3_a (n^1_a n^2_a n^3_a - 1)$ additional fermions in the symmetric and antisymmetric representation of the $U(N_a)$ gauge group.

In general the calculation of the spectrum for a particular string vacuum is subject also to constraints coming from the RR tadpole cancellation conditions \[1\]. That means the cancellation of D6-branes charges,$^3$ wrapping on three cycles with homology $[\Pi_a]$ and O6-plane 7-form charges wrapping on 3-cycles with homology $[\Pi_{O6}]$. For the toroidal orientifolds we deal in this section, the RR tadpole cancellation conditions in terms of cancellations of RR charges in homology become

$$\sum_a N_a [\Pi_a] + \sum_{\alpha'} N_{a'} [\Pi_{\alpha'}] - 32 [\Pi_{O6}] = 0. \quad (3.3)$$

In explicit form, the RR tadpole conditions read

$$\sum_a N_a n^1_a n^2_a n^3_a = 16,$$
\[ \sum_a N_a m_a^1 m_a^2 m_a^3 = 0, \]
\[ \sum_a N_a m_a^1 n_a^2 m_a^3 = 0, \]
\[ \sum_a N_a n_a^1 m_a^2 m_a^3 = 0. \] (3.4)

The complete accommodation of the chiral matter for the five stack SM’s can be seen in table [2]. A number of interesting observations are in order:

• There are various low energy gauged symmetries in the models (as it also happens in the models [1, 3]). They are defined in terms of the $U(1)$ symmetries $Q_a, Q_b, Q_c, Q_d, Q_e$. Hence the baryon number $B$ is defined as $Q_a = 3B$, the lepton number is $L = Q_d + Q_e$ while $Q_a - 3Q_d - 3Q_e = 3(B - L)$; $Q_e = 2I_{3R}$, $I_{3R}$ being the third component of weak isospin. Also, $3(B - L)$ and $Q_e$ are free of triangle anomalies. The $U(1)_b$ symmetry plays the role of a Peccei-Quinn symmetry in the sense of generating mixed $SU(3)$ anomalies. From the study of Green-Schwarz mechanism cancellation of anomalies, we deduce that Baryon and Lepton number are unbroken gauged symmetries and thus proton should be stable. Also Majorana masses for right handed neutrinos are not allowed, meaning that mass terms for neutrinos should be of Dirac type.

• In order to cancel the appearance of exotic representations in the model appearing from the $aa^*$ sector, in antisymmetric and symmetric representations of the $U(N_a)$ group, we will impose the condition

\[ \Pi_{i=1}^3 m^i = 0. \] (3.5)

The solutions satisfying simultaneously the intersection constraints and the cancellation of the RR crosscap tadpole constraints are given in table [5].

By using the tadpole solutions of table [5] in (3.4) all tadpole equations but the first are satisfied, the \(^4\) latter becoming:

\[ \frac{9n_a^2}{\beta_1^1} + 2 \frac{n_b^1}{\beta_2^2} + \frac{n_d^2}{\beta_3^1} + \frac{n_e^2}{\beta_3^1} + N_h \frac{2}{\beta_1^1 \beta_2^2} = 16. \] (3.6)

where we had added the presence of extra $N_h$ branes by setting $m_h^3 = 0$. The cancellation of tadpoles is best described by choosing a consistent numerical set of wrappings, e.g.

\[ n_a^2 = 1, \quad n_b^1 = 1, \quad n_c^1 = -1, \quad n_d^2 = -1, \quad n_e^2 = -1, \quad \beta_1^1 = 1/2, \quad \beta_2^2 = 1, \quad \epsilon \bar{\epsilon} = 1 \] (3.7)

\(^4\)We have added an arbitrary number of $N_h$ branes which don’t contribute to the rest of the tadpoles and intersection number constraints.
Table 5: New general tadpole solutions for the five (5) stack intersecting D6-brane model [not appearing in [2]] coming from \(IIA/T^6/\Omega R\), giving rise to exactly the standard model gauge group and observable chiral spectrum at low energies. These solutions depend on two integer parameters, \(n_1^a, n_1^d\), the NS-background \(\beta^1 = 1 - b_i\), which is associated to the presence of the NS B-field by \(b_i = 0, 1/2\), and the phase parameters \(\epsilon = \tilde{\epsilon} = \pm 1\).

With the above choices, all tadpole conditions but the first are satisfied, the latter is satisfied when we add \(N_h = 0\) extra D6 branes. The wrappings become

\[
\begin{align*}
N_a &= 3 \quad (1/\beta_1, 0) \quad (n_2^a, \tilde{\epsilon}\epsilon\beta_2) \quad (3, -\tilde{\epsilon}/2) \\
N_b &= 2 \quad (n_1^b, \epsilon\beta_1) \quad (1/\beta_2, 0) \quad (1, -\tilde{\epsilon}/2) \\
N_c &= 1 \quad (n_1^c, \epsilon\tilde{\epsilon}\beta_1) \quad (1/\beta_2, 0) \quad (0, 1) \\
N_d &= 1 \quad (1/\beta_1, 0) \quad (n_2^d, 2\epsilon\tilde{\epsilon}\beta_2) \quad (1, \tilde{\epsilon}/2) \\
N_e &= 1 \quad (1/\beta_1, 0) \quad (n_2^e, \epsilon\tilde{\epsilon}\beta_2) \quad (1, \tilde{\epsilon}/2) \\
N_h &= (1/\beta^1, 0) \quad (1/\beta_2, 0) \quad (2, m_3^h) \\
\end{align*}
\]

The mixed anomalies \(A_{ij}\) of the five \(U(1)\)'s with the non-abelian gauge groups \(SU(N_a)\) of the theory cancel through a generalized GS mechanism involving close string modes couplings to worldsheet gauge fields. Three combinations of the \(U(1)\)'s are anomalous and become massive. Two orthogonal
non-anomalous combinations survive massless the Green-Schwarz mechanism, one being the hypercharge seen in (1.4) and also an extra U(1) whose value is model dependent and depends on the choice of parameters seen in table (5). The latter U(1) get broken by demanding that the intersection \( e \bar{e} \) respects N=1 supersymmetry, that is

\[
\pm \tan(-1) \left( \frac{\beta_1 U^1}{n_c^1} \right) \pm \tan(-1) \left( \frac{\beta_2 U^2}{n_c^2} \right) + \frac{\pi}{2} \pm \tan(-1) \left( \frac{U^3}{2} \right) \tag{3.9}
\]

- where we have chosen \( \epsilon \tilde{\epsilon} = 1 \), for some choice of signs. In the special case, \( n_c^2 = 0 \), we get the N=1 susy condition

\[
\frac{\beta_1 U^1}{n_c^1} = \frac{U^3}{2} \tag{3.10}
\]

constraining the complex structure parameters \( U^1, U^3 \). The choice \( n_c^1 = 1 \) determines the fifth U(1) to be \( F_a + (28/3) F_c - F_e \). We note that on this classes of models the string scale cannot be lowered at the TeV region according to the geometric scenario of [12] and thus the solution of gauge hierarchy problem in the Higgs sector is an open question. Whether of not, the supersymmetry present in particular sectors of the theory, or additional superymmetry - that can be implemented in the models by some choice of parameters - could be of any help remains to be seen.

The 6-stack SM configuration that was localized in toroidal orientifold compactifications from intersecting D6-branes in [3] was given in table (3). Here we will present an alternative solution to the RR tadpoles. The new solution can be seen in table (6). We have also added an arbitrary number of extra \( N_h \) branes that might needed in order to satisfy the RR tadpoles. For these solutions the procedures followed in [3] could be also repeated to examine further the SM’s.

4 Building the SU(4)_C × SU(2)_L × SU(2)_R GUTS With Only The SM At Low Energy

Extensions of these GUTS with four, five and six stacks of D6-branes were also considered in [10, 17, 18].

The basic features found in these intersecting D6-brane models can be classified as follows:

- The models even though they have overall N=0 SUSY, possess N=1 SUSY subsectors which are necessary in order to create a Majorana mass term for
Table 6: New general tadpole solutions for the six (6) stack intersecting D6-brane model [not appearing in [3]] coming from $IIA/T\Omega R$, giving rise to exactly the standard model gauge group and observable chiral spectrum at low energies. These solutions depend on two integer parameters, $n^1_{a}, n^1_{d}$, the NS-background $\beta_1 = 1 - b_i$, which is associated to the presence of the NS B-field by $b_i = 0, 1/2$, and the phase parameters $\epsilon = \tilde{\epsilon} = \pm 1$.

Extra branes are needed to cancel RR tadpoles. The presence of these branes creates extra matter singlets, transforming under both the visible SM gauge group and the extra D6-brane gauge group that may be used to break the extra U(1)'s, beyond hypercharge, surviving massless the presence of the generalized Green-Schwarz mechanism. Their presence is also used to make massive the exotic fermions, seen for example in the bottom part of table (1), taken from [18]. The fermion spectrum of table (1) is consistent with the calculation of RR tadpoles. The RR tadpoles get cancelled with the introduction of extra U(1) branes, $h^i$, that transform under the both the extra U(1) gauge group and the rest of the intersecting D6-branes of table (1). The existence of N=1
The left handed fermions issue in more detail. All fermions of table (1) receive a mass of order $M$.

Table 7: Fermionic spectrum of the $SU(4)_C \times SU(2)_L \times SU(2)_R$, PS-II class of models together with $U(1)$ charges. We note that at energies of order $M_s$ only the Standard model survives.

SUSY at the intersections $dd^*, dh, dh^*, eh, eh^*$, creates the singlets $s^B_1, \kappa^b_3, \kappa^b_4; \kappa^b_5, \kappa^b_6$ respectively, that contribute to the mass of the ‘light’ fermions $\chi^1_L, \chi^2_L$. All fermions of table (1) receive a mass of order $M_s$; the only exception being the light masses of $\chi^1_L, \chi^2_L$, weak fermion doublets. Let us discuss the latter issue in more detail.

The left handed fermions $\chi^1_L$ receive a contribution to their mass from the coupling

$$\begin{align*}
(1, 2, 1)(1, 2, 1)e^{-A}\frac{\langle h_2\rangle\langle h_2\rangle\langle F^H_R\rangle\langle H_1\rangle\langle s^B_1\rangle}{M_s^4} & \sim \frac{\nu^2}{M_s} (1, 2, 1)(1, 2, 1) \quad (4.1)
\end{align*}$$

and from another coupling, of the same order as (4.1), also contributing to the mass of the $\chi^1_L$ fermion as

$$\begin{align*}
(1, 2, 1)(1, 2, 1)\frac{\langle h_2\rangle\langle h_2\rangle\langle F^H_R\rangle\langle H_1\rangle\langle \kappa^B_3\rangle\langle \kappa^B_4\rangle}{M_s^3} & \sim (1, 2, 1)(1, 2, 1) \frac{\nu^2}{M_s} \quad (4.2)
\end{align*}$$

The left handed fermions $\chi^2_L$ receives a non-zero mass from the coupling

$$\begin{align*}
(1, 2, 1)(1, 2, 1)\frac{\langle h_2\rangle\langle h_2\rangle\langle F^H_R\rangle\langle H_1\rangle\langle s^2_B\rangle}{M_s^4} & \sim \frac{\nu^2}{M_s} (1, 2, 1)(1, 2, 1) \quad (4.3)
\end{align*}$$

and the coupling

$$\begin{align*}
(1, 2, 1)(1, 2, 1)\frac{\langle h_2\rangle\langle h_2\rangle\langle F^H_R\rangle\langle H_1\rangle\langle \kappa^B_5\rangle\langle \kappa^B_6\rangle}{M_s^3} & \sim (1, 2, 1)(1, 2, 1) \frac{\nu^2}{M_s}, \quad (4.4)
\end{align*}$$

\[\text{In } (4.1) \text{ we have included the leading contribution of the worksheet area connecting the seven vertices. In the following for simplicity reasons we will set the leading contribution of the different couplings to one e.g. area tends to zero.}\]
Thus assuming that the leading area Yukawa for the couplings is of order $O(1)$, e.g. associated areas going to zero, the masses of $^6$

$$\chi_L^1, \chi_L^2 \sim \frac{2\nu^2}{M_s} \quad (4.5)$$

- As the particles $\chi_L^1, \chi_L^2$ are not observed at present, the fact that their mass may be between

$$100 \text{ GeV} \leq \chi_L^1, \chi_L^1 \leq 2\nu = \max\{\frac{2\nu^2}{M_s}\} = 492 \text{ GeV} \quad (4.6)$$

sends the string scale

$$M_s \leq 1.2 \text{ TeV} \quad (4.7)$$

This is a general feature of all the Pati-Salam models based on toroidal orientifolds; they predict the existence of light weak doublets with masses $^7$ between 100 and $\nu = 492$ GeV. The latter result may be considered as a general prediction of all classes of models based on intersecting D6-brane Pati-Salam GUTS. Another important property of these constructions is that the conditions for some intersections to respect $N=1$ supersymmetry and also needed to guarantee the existence of a Majorana mass term for $s\nu_R$’s:

- solve the orthogonality conditions for the extra - beyond hypercharge - U(1)’s $^8$ to survive massless the presence of a generalized Green-Schwarz mechanism describing the couplings of the U(1)’s to the RR two form fields.

The considerations we have just described $^{10}$, $^{17}$, $^{18}$ are quite generic and the same methodology applies easily to the construction of more general GUT gauge groups in the context of intersecting brane worlds.

We note that at present the only existing string GUT constructions, in the context of Intersecting D6-brane Models, that have only the SM at low energy, with complete cancellation of RR tadpoles, are:

a) the toroidal orientifold II Pati-Salam GUTS of $^{10}$ $^{17}$ $^{18}$ and

b) the constructions of flipped SU(5), and SU(5) GUTS of $^{23}$ described next.

$^6$In this case the masses of $\chi_L^1, \chi_L^2$ are the sum of the contributions of $^{1.1}$ $^{2.2}$ and $^{4.3}$ $^{4.4}$ respectively.

$^7$The reader may convince itself that the maximum value of $2\nu^2/M_s$ is $2\nu$.

$^8$The latter becoming massive from the use of extra singlets created by the presence of extra branes; the latter needed to satisfy the RR tadpoles.
The Construction of SU(5), Flipped SU(5) GUTS with only the SM at Low Energy

Let us review the intersecting D6-branes constructions of the $Z_3$ orientifolds of [22]. The D6-branes involved satisfy the following RR tadpole conditions where

$$\sum_a N_a Z_a = 2.$$  \hspace{1cm} (5.3)

As it was noticed in [22] the simplest realization of an SU(5) GUT involves two stacks of D6-branes at the string scale $M_s$, the first one corresponding to a $U(5)$ gauge group while the second one to a $U(1)$ gauge group. Its effective wrapping numbers are given by

$$(Y_a, Z_a) = (3, \frac{1}{2}), \quad (Y_b, Z_b) = (3, -\frac{1}{2}),$$  \hspace{1cm} (5.4)

Under the decomposition $U(5) \subset SU(5) \times U(1)$, the models become effectively an $SU(5) \times U(1) \otimes U(1)$ GUT. One combination of $U(1)$’s become massive due to its coupling to a RR field, another one remains massless to low energies. The spectrum of this SU(5) GUT may be seen in the first seven columns (reading from the left) of table (2). At this stage the SU(5) models - have the correct chiral fermion content of an SU(5) GUT - and the extra $U(1)$ surviving the

| Field | Sector | Multip. | SU(5) | U(1)$_a$ | U(1)$_b$ | U(1)$_{mass}$ | U(1)$_{flg}$ = $\frac{5}{2}$ × U(1)$_{mass}$ |
|-------|--------|---------|-------|----------|----------|-------------|------------------|
| $f$   | {51}   | 3       | 5     | -1       | 1        | $-\frac{5}{2}$ | -3               |
| $F$   | $A_a$  | 3       | 10    | 2        | 0        | $\frac{2}{5}$  | 1                |
| $l^c$ | $S_b$  | 3       | 1     | 0        | -2       | 2           | 5                |

Table 8: Chiral Spectrum of a two intersecting D6-brane stacks in a three generation flipped $SU(5) \otimes U(1)_{mass}$ model. Note that the charges under the $U(1)_{flg}$ gauge symmetry, when rescaled appropriately (and $U(1)_{flg}$ gets broken) ‘converts’ the flipped SU(5) model to a three generation (3G) SU(5).

\footnote{The net number of bifundamental massless chiral fermions in the models is defined as}

$$(\tilde{N}_a, N_b)_L : \quad I_{ab} = Z_a Y_b - Y_a Z_b$$  \hspace{1cm} (5.1)

$$(N_a, \tilde{N}_b)_L : \quad I_{ab^\ast} = Z_a Y_b + Y_a Z_b$$  \hspace{1cm} (5.2)
presence of the Green-Schwarz mechanism, breaks by the use of a singlet field present. However, the electroweak 5-plets needed for electroweak symmetry breaking of the models are absent. Later on, attempts to construct a fully N=1 supersymmetric SU(5) models at $M_s$ in \[21\], produced 3G models that were not free of remaining massless exotic 15-plets. Also, later on in \[20\] it was noticed that if one leaves unbroken, and rescales, the U(1) surviving massless the Green-Schwarz mechanism of the SU(5) GUT of \[22\], the rescaled U(1) becomes the flipped U(1) generator. However, the proposed 3G models lack the presence of GUT Higgses or electroweak pentaplets and were accompanied by extra exotic massless matter to low energies.

In \[23\] we have shown that it is possible to construct the first examples of string SU(5) and flipped SU(5) GUTS - where we identified the appropriate GUT and electroweak Higgses - which break to the SM at low energy. E.g. in the flipped SU(5) GUT, the fifteen fermions of the SM plus the right handed neutrino $\nu^c$ belong to the chiral multiplets. The GUT breaking Higgses may come from the ‘massive’ spectrum of the sector localizing the 10-plet $(10^B_1 = (u_H, d_H, d^c_H, \nu^c_H))$ fermions seen in table (8). The lowest order Higgs in this sector, let us call them $H_1$, $H_2$, have quantum numbers as those given in table (9). By looking at the last column of table (9), we realize that the Higgs $H_1$, $H_2$ are the GUT symmetry breaking Higgses of a standard flipped SU(5) GUT. By duplicating the analysis of section (3.1), one may conclude that what it appears in the effective theory as GUT symmetry breaking scalars, is the combination $H^G = H_1 + H_2^*$. In a similar way the correct identification of the electroweak content \[23\] of the flipped SU(5) $5^B_{-2} = (D, h^-, h^0)$-plet (and SU(5))GUTS made possible the existence of the see-saw mechanism which is generated by the interaction

$$\mathcal{L} = \tilde{Y}_{\nu_R} \cdot 10 \cdot 5 \cdot \tilde{h}_4 + \tilde{Y}_{\nu_R} \cdot \frac{1}{M_s} \cdot (10 \cdot \overline{10}^B)(10 \cdot \overline{10}^B).$$

(5.6)

Its standard version can be generated by choosing

$$\langle h_4 \rangle = v, \quad \langle 10^B_i \rangle = M_s$$

(5.7)
and generates small neutrino masses. In these constructions the baryon number is not a gauged symmetry, thus a high GUT scale of the order of \(10^{16}\) GeV helps the theory to avoid gauge mediated proton decay modes like the

\[
\sim \frac{1}{M_s^2} (\bar{u}_L^c u_L) (\bar{e}_R^c) (d_R), \quad \sim \frac{1}{M_s^2} (\bar{d}_R^c u_R) (\bar{d}_L^c) (\nu_L). \tag{5.8}
\]

[In IBW’s proton decay by direct calculation of string amplitudes for SUSY SU(5) D-brane models was examined in [13].] Scalar mediated proton decay modes get suppressed by the existence of a new solution to the doublet-triplet splitting problem

\[
\frac{r}{M_s^2} (HHh)(\bar{F}\bar{F}\bar{h}) + m(\bar{hh})(ar{H}H) + \kappa(\bar{H}H)(\bar{H}H), \tag{5.9}
\]

that stabilizes the vev’s of the triplet scalars \(d_h^H\), \(D\) [23]. This is the first example of a doublet-triplet splitting realization in IBW’s. The full solution of the gauge hierarchy problem, that is avoiding the existence of quadratic corrections to the electroweak Higgses remains an open issue in the present GUTS.

Recently the interest of model building in IBW’s has been focused in the construction of non-susy intersecting D6-brane models which localize the spectrum of MSSM at low energies [24], [25].

6 The SM at low energy from \(Z_3\), \(Z_3 \times Z_3\) orientifolds

The SM but with some of the matter not in bifundamental reps have been derived initially in [22] in the context of \(Z_3\) orientifolds. Quite recently [25], we have shown that using \(Z_3 \times Z_3\) orientifolds, we can reproduce these SM vacua for various choices of wrappings. In this section, we present some novel features of the new constructions, including the derivation of the SM.

In \(Z_3\), \(Z_3 \times Z_3\) orientifolds all complex structure moduli are fixed from the beginning. In particular the \(Z_3 \times Z_3\) orientifold models are subject to the cancellation of untwisted RR tadpole conditions [25] given by

\[
\sum_a N_a Z_a = 4, \tag{6.1}
\]

where

\[
Z_a = 2m_a^2 m_a^1 m_a^0 + 2n_a^2 n_a^1 n_a^0 - n_a^2 m_a^1 m_a^0 - n_a^1 m_a^2 n_a^0 - m_a^1 n_a^2 n_a^0 - m_a^1 m_a^2 n_a^0 - m_a^1 n_a^2 m_a^0 - n_a^1 m_a^2 m_a^0. \tag{6.2}
\]
The gauge group $U(N_a)$ supported by $N_a$ coincident D6$_a$-branes comes from the $a(Oa)$ sector, the sector made from open strings stretched between the $a$-brane and its images under the orbifold action. In addition, we get three adjoint N=1 chiral multiplets. The $a(Ob)$ sector, strings stretched between the brane $a$ and the orbit images of brane $b$, will give $I_{ab}$ fermions in the bifundamental ($N_a, \bar{N}_b$) where

$$I_{ab} = 3(Z_a Y_b - Z_b Y_a),$$

and ($Z, Y$) are the effective wrapping numbers with $Y$ given by

$$Y_a = m_a^1 m_a^2 m_a^3 + n_a^1 n_a^2 n_a^3 - n_a^1 n_a^2 m_a^3 - n_a^1 m_a^2 n_a^3 - m_a^1 n_a^2 n_a^3$$

The sign of $I_{ab}$ denotes the chirality of the associated fermion, where we choose positive intersection numbers for left handed fermions. In the sector $ab'$ - strings stretching between the brane $a$ and the orbit images of brane $b$, there are $I_{ab'}$ chiral fermions in the bifundamental ($N_a, N_b$), with

$$I_{ab'} = 3(Z_a Z_b - Z_a Y_b - Z_b Y_a),$$

The following numbers of chiral fermions in symmetric (S) and antisymmetric (A) representations of $U(N_a)$, open strings stretching between the brane $a$ and its orbit images ($Oa$), are also included

$$(A_a) = 3(Z_a - 2Y_a),$$

$$(A_a + S_a) = \frac{3}{2}(Z_a - 2Y_a)(Z_a - 1)$$

Also, from open strings stretched between the brane $a$ and its orbifold images we accommodate non-chiral massless fermions in the adjoint representation,

$$(\text{Adj})_L : \prod_{i=1}^{3} (L^I_{[a]})^2,$$

where

$$L^I_{[a]} = \sqrt{(m^I_a)^2 + (n^I_a)^2 - (m^I_a)(n^I_a)}$$

Adjoint massless matter, including fermions and gauginos, is expected to become massive from loops once supersymmetry is broken leaving only the gauge bosons massless. Supersymmetry may be preserved by a system of branes if each stack of D6-branes is related to the O6-planes by a rotation in $SU(3)$, that is the angles $\bar{\theta}_i$ of the D6-branes with respect to the horizontal direction in the i-th two-torus obeys the condition $\bar{\theta}_1 + \bar{\theta}_2 + \bar{\theta}_3 = 0$. In the low energy
theory, cubic gauge anomalies automatically cancel, due to the RR tadpole conditions \[6.1\]. Mixed U(1)-gauge anomalies also cancel due to the existence of a generalized Green-Schwarz mechanism that makes massive only one U(1) gauge field given by

\[
\sum_a N_a (Z_a - 2Y_a) F_a
\]  

(6.10)

By choosing to work with three stacks of intersecting D6-branes and making the choice of wrapping numbers

\[
(Z_a, Y_a) = \left( 1, \ 0 \right), \ (Z_b, Y_b) = \left( 1, \ 1 \right), \ (Z_c, Y_c) = \left( -1, \ -1 \right)
\]  

(6.11)

we find exactly the (non-supersymmetric) SM spectrum that may be seen in table (10).

| Matter       | \((SU(3) \times SU(2))(Q_a, Q_b, Q_c)\) | \(U(1)Y\) |
|--------------|----------------------------------------|------------|
| \{QL\}      | 3(3, 2)\((1, -1, 0)\)                 | 1/6        |
| \{u_L\}     | 3(3, 1)\((2, 0, 0)\)                  | -2/3       |
| \{d_L\}     | 3(3, 1)\((-1, 0, 1)\)                | 1/3        |
| \{L\}       | 3(1, 2)\((0, 1, 1)\)                  | -1/2       |
| \{e_L\}     | 3(1, 1)\((0, -2, 0)\)                | 1          |
| \{N_R\}     | 3(1, 1)\((0, 0, -2)\)                | 0          |

Table 10: A three generation chiral (open string) spectrum accommodating the SM. The required Higgs may come from bifundamental \(N=2\) hypermultiplets in the \(N=2\) bc, bc* sectors \([1, 2, 3]\) that may trigger brane recombination.

This is exactly the SM found in \([22]\). It can be reproduced \([25]\) with many choices of effective wrappings \((Z, Y)\). As we can see, many of the matter fields present are not in bifundamental representations. In this case, we find that there is no mass term for the up-quarks allowed by charge conservation, which is not the case for the SM’s of section 3 and \([1, 2, 3]\).

7 Conclusions

IBW’s offer a lot of promising research directions. We note that one of the major problems in string model building, and in string theory in general, is
the determination of the values of the moduli. As moduli are free parameters, the determination of the single vacuum in string theory requires that there is a dynamical determination of their value. IBW’s offer a solution towards this direction, as we have seen in section 3, as the presence of N=1 supersymmetry in particular sectors of the theory can fix the complex structure moduli. On the other hand SM’s coming from $Z_3 \times Z_3$ orientifolds have from the start all complex structure moduli fixed. In this respect, $Z_3 \times Z_3$ orientifolds solve in principle the problem of determining complex structure moduli. Thus model building attempts coming from the latter constructions offer a promising avenue for further studies as it only remains to show that it is possible that also the Kähler moduli can be also fixed (possibly from non-perturbative effects). We also note that the phenomenology of models with intersecting D5-branes on $IIB/\Omega R/(T^4 \times C/Z_N)$ (of section 2) need to be further studied as these Standard models suffer from no gauge hierarchy problem while they are also consistent with the extra dimension scenario [12].

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