N=1 Locally Supersymmetric Standard Models from Intersecting branes

Christos Kokorelis

Dep/to de Física Teórica C-XI and Instituto de Física Teórica C-XVI, Universidad Autónoma de Madrid, Cantoblanco, 28049, Madrid, Spain

ABSTRACT

We construct four dimensional intersecting D6-brane models that have locally the spectrum of the N=1 Supersymmetric Standard Model. All open visible string sectors share the same N=1 supersymmetry. As expected in these supersymmetric classes of models, where the D6-branes wrap a toroidal orientifold of type IIA, the hierarchy may be stabilized if the string scale is low, e.g. below 30 TeV. We analyze the breaking of supersymmetry in the vicinity of the supersymmetric point by turning on complex structure deformations as Fayet-Iliopoulos terms. Positive masses for all squarks and sleptons, to avoid charge/colour breaking minima, may be reached when also two loop contributions may be included. In the ultimate version of the present models N=1 supersymmetry may be broken by gauge mediation. The constructions with four, five and six stacks of D6-branes at $M_s$ are build directly. Next by the use of brane recombination we are able to show that there is a continuous, RR homology flow, between six, five and four stack models. Moreover, we examine the gauge coupling constants of the Standard Model $SU(3)_C$, $SU(2)_L$, $U(1)_Y$ at the string scale in the presence of a non-zero antisymmetric NS B-field.
1 Introduction

It is widely accepted that string theory provides the only consistent theoretical framework for perturbative quantum gravity. In this respect, in the absence of a dynamical principle for selecting a particular string vacuum which may be describing our world, one has to demonstrate beyond doubt that vacua with only the Standard Model (SM) at low energy exist. In this case, string theory may definitely provide us with acceptable evidence for describing the real world. Steps forward in this direction has been taken recently in the context of Intersecting Brane Worlds [1]-[29] (IBW’s). We note that string theory based intersecting brane worlds is the explicit string realization of the idea about the field theory breaking of space-time supersymmetry via the introduction of magnetic fields [2]. See also [3, 30, 31].

In general one expects that string vacua may be such that either they have broken supersymmetry from the start, at the string scale, or they will preserve some amount of supersymmetry \(^1\). In the latter case one might try to build explicit N=1 supersymmetric models using intersecting D6-branes, but one has to solve first serious problems like the existence of extra exotic chiral matter remaining massless to low energies - which cannot become massive from existing stringy couplings - absence of doublet-triplet splitting for supersymmetric GUT constructions, etc. [ See [7, 19, 20, 21] for studies in the context of N=1 SUSY intersecting D6 brane model building] , before discussing realistic N=1 SUSY model building.

In the former case it has been explicitly demonstrated, using intersecting D6-brane language, that non-supersymmetric intersecting D6-brane vacua with only the SM at low energy exist [4, 5, 6] [For some other attempts to construct, using D-branes the (non-susy) SM, but not based on a particular string construction, see [32].]. It has also been shown that non-supersymmetric GUT vacua, either of Pati-Salam type [8] or SU(5)/flipped SU(5) GUTS [11] with only the SM at low energy exist. In these works, it has been evident that by working, in great detail, with four dimensional compactifications of type IIA toroidal orientifolds [1] (or their orbifolds [10] in the case of SU(5)/flipped SU(5) GUTS [11]) - and D6-branes intersecting at angles - one can localize eventually the SM at low energy in all these constructions [see for example [28],[29] for some reviews]. We note that even though models with D6-branes intersecting at angles are related by T-duality to models with magnetic deformations [3, 31], it is much easier to work in the angle picture as one may handle easier the parameters.

\(^1\)The cosmological constant problem remains in both options.
of the theory. We note that the SM at low energy has been shown to exist in another toroidal-like constructions, using D5-branes in [13], [14].

A number of important phenomenological issues have been also examined in the context of model building in IBW’s, including proton decay [9, 11] and the implementation of doublet-triplet splitting mechanism [11].

There is an essential difference between the four D6-brane stack models of [4] and its five, six D6-brane stack model extensions of [5, 6] respectively. The models of [4], which have overall N=0 supersymmetry don’t have open string sectors which preserve some amount of supersymmetry (SUSY). On the contrary, the models of [5, 6] have necessarily some open string sectors which preserve N=1 supersymmetry. This is necessary in order to make massless the previously massive superpartners of $\nu_R$, $s\nu_R$. By giving a vev to $s\nu_R$ one may be able to break the extra U(1)’s, beyond hypercharge, that survive massless the Green-Schwarz anomaly cancellation mechanism. Thus at low energy only the SM survives. Similar considerations apply to the non-supersymmetric Pati-Salam GUTS of [8], where the presence of some open string sectors preserving N=1 supersymmetry with the orientifold plane is necessary in order to create a Majorana mass coupling term for $\nu_R$ and get rid of the massless beyond hypercharge U(1)’s that survive massless the Green-Schwarz mechanism. In those cases, we were able to localize the presence of the SM spectra together with extra matter at the string scale, the latter finally becoming massive by appropriate stringy couplings, leaving only the SM at low energy.

Alternatively, one may try to build non-supersymmetric models where each open string sector preserves some amount of SUSY, and where each brane shares some SUSY with the other branes but not necessarily the same one. Such constructions have been considered in [15, 16] and as the models are not really N=1 supersymmetric they were called quasi-supersymmetric (Q-SUSY). The most important feature of these phenomenological models is the absence of one loop corrections to the scalar masses (CSM’s); necessarily CSM’s appear at two loop order. As a result the lower part of the gauge hierarchy problem is stabilized in these models for a string scale of less than 30 TeV.

\[ \text{Absence of large quadratic corrections to the scalar masses; the higher part of the gauge hierarchy problem we identify as the explanation of the ratio } \frac{M_W}{M_{	ext{Planck}}} . \]
The ‘symmetrical’ choice in the latter constructions is to choose all open string sectors to share the same N=1 supersymmetry. In this case, we may attempt to localize the particle spectrum of the N=1 SUSY SM (N=1 local SUSY models), even though the models overall may have N=0 supersymmetry.

In this work, firstly we consider models with intersecting D6 branes which share the same N=1 SUSY at every intersection. In this way we localize the spectrum of the N=1 Supersymmetric Standard model. Hence we first consider generalizations of the four stack N=1 local SM of [16], in the sense of allowing the most general wrappings in the presence of NS B-field between the intersections of the branes involved. We note that consequences for gauge coupling unification have been examined in [25] based on this model. We also consider further consequences for the gauge couplings for the present models generalizing some of the results of [25], as we consider the additional presence of the NS B-field in the models of [16].

Secondly, we present new solutions with the spectrum of the MSSM, involving a trivial NS B-field across the compact six dimensional space, with D-brane configurations that contain the five and its maximum extended D6-brane six stack constructions of [16]. An interesting feature of these constructions is that these models may have necessarily a low string scale and thus avoid the gauge hierarchy problem related to quadratic Higgs corrections [15]. We also note that non-supersymmetric toroidal orientifold models [15, 5, 6] or their orbifolds, have some non-zero NS tadpoles whose presence acts as an uncancelled cosmological constant, similar (but not equivalent) to the cosmological constant appearing after the breaking of N=1 space-time supersymmetry in supersymmetric models.

The paper is organized as follows. In section 2, we present a brief review of the main features of the constructions of [15, 16]. In particular we present the most general solution to the wrapping numbers of the SM configurations of [16] in the presence of a NS B-field. For clarity reasons we will identify the four, five and six stack models described, as belonging to the Model classes I, II and III respectively. In sections 3, 4 we describe our new five, six stack configurations respectively, which localize the spectrum of the N=1 SUSY SM spectrum, on N=1 supersymmetric intersections, in overall N=0 models. In section 5 we analyze the breaking of the common SUSY preserved in the intersections using FI terms. Similar considerations have appeared in [35, 36, 7, 15]. In section 6 we describe the existence of gauge breaking transitions in the models appearing after brane recombination; the latter effectively corresponding to the existence of adjoint Higgs fields ‘localized’ between the parallel recombined branes. In
In section 8 we discuss gauge mediated N=1 space-time supersymmetry breaking and the issue of the non-factorizable brane needed to cancel RR tadpoles. Finally in section 9 we summarize our concluding remarks.

2 N=1 local SYSY models revisited - Model I

Our aim is to derive models which are of phenomenological interest and which localize the spectrum of the N=1 supersymmetric Standard Model (SSM). The most important property of IBW’s is that open strings stretching in the intersection between branes get associated with chiral fermions localized in their intersections [12]. The multiplicity of the chiral fermions in the intersections is given by the intersection number. In its simplest constructions, the type IIA theory gets compactified to four dimensions (4D) on an orientifolded $T^6$ torus [1]. Moreover, the $D6_a$ branes wrap three-cycles $(n^i_a, m^i_a)$, $i = 1, 2, 3$, across each of the i-th $T^2$ torus of the assumed, for simplicity, factorized $T^6$ torus, $T^6 = T^2 \otimes T^2 \otimes T^2$. In particular, our torus is allowed to wrap factorized products of three 1-cycles, whose homology classes are given exactly by

$$[\Pi^a] = \prod_{i=1}^3 (n^i_a[a_i] + m^i_a[b_i])$$

(2.1)

and their orientifold images as

$$[\Pi^a^\ast] = \prod_{i=1}^3 (n^i_a[a_i] - m^i_a[b_i]).$$

(2.2)

In (2.2) we have used the fact that wrapping numbers of the orientifold images of the $(n, m)$ wrappings of the $\alpha$-brane are given by $(n, -m)$. The wrappings numbers $(n^i, m^i)$ may take integer values in which case the tori that the D6-branes wrap are orthogonal. If a non-trivial NS B-field is added, the tori becomes tilted and the effective wrapping numbers become $(n^i, m^i) = (n^i, \tilde{m}^i + n^i/2)$, $n^i, \tilde{m}^i, \in Z$, where the magnetic wrappings $m$, take now fractional values. There are several open string sectors, including the $ab$ sector accommodating open strings stretching between the $D6_a$, $D6_b$ branes and localizing fermions transforming in the bifundamental representation $(N_a, \bar{N}_b)$ with multiplicity

$$I_{ab} = [\Pi_a] \cdot [\Pi_b] = \prod_{i=1}^3 (n^i_a m^i_b - m^i_a n^i_b).$$

(2.3)
Also, the $ab^*$ sector is present with chiral fermions transforming in the bifundamental representation $(N_a, N_b)$ and multiplicity given by

$$I_{ab^*} = [\Pi_a] \cdot [\Pi_{b^*}] = -\prod_{i=1}^{3} (n_a^i m_b^i + m_a^i n_b^i). \quad (2.4)$$

The sign of $I_{ab}$, $I_{ab^*}$, denotes the chirality of the corresponding fermion and it is a matter of convention; the positive sign we identify with the left handed fermions. The gauge group for a set of $a, b$ D6-branes on the above background, is in general $U(N_a) \times U(N_b)$. In the case that the brane $a$ is its own orientifold image then the gauge group may be further enhanced from $U(N_a)$ to $SO(2N_a)$ or $USp(2N_a)$. The wrapping numbers are further constrained by the conditions set out by the RR tadpole cancellation conditions \[1\]

$$\sum_a N_a n_a^1 n_a^2 n_a^3 = 16, \sum_a N_a n_a^1 m_a^2 m_a^3 = 0, \sum_a N_a m_a^1 n_a^2 m_a^3 = 0, \sum_a N_a m_a^1 m_a^2 n_a^3 = 0 \quad (2.5)$$

- **N=1 Supersymmetric SM constructions from intersecting D6-branes**

Let us assume the particular D6-brane configuration, seen in table (1), with gauge group $U(3) \times U(1)_b \times U(1)_c \times U(1)_d$. For the special values of the parameters $\epsilon = \bar{\epsilon} = 1, \beta_1 = \beta_2 = 1$, that is when all tori are untilted, we recover the D-brane configuration proposed in \[16\]. It clearly describes the chiral spectrum of the MSSM. The hypercharge is defined as

$$Q_Y = (1/6) Q_a - (1/2) Q_c - (1/2) Q_d \quad (2.6)$$

Also there are exact identifications between the global symmetries of the SM and the U(1) symmetries set out by the D-brane configurations of table (1). Thus baryon number ($B$) is easily identified as $Q_a = 3B$, the lepton number ($L$) is given by $Q_d = L$, and $Q_c$ is twice $I_R$ the third component of the right-handed weak isospin. The intersection numbers localizing the chiral fermions of table (1) are given by

$$I_{ab} = 3, I_{ab^*} = 3, I_{ac} = -3, I_{acs} = -3, I_{bc} = -\frac{1}{\beta_1 \beta_2}, \quad (2.7)$$

The wrapping numbers of the SM chiral fermions can be seen in table (2). The entries of this table represent the most general solution to the wrapping numbers (2.7). For the special values of the parameters $\epsilon = \bar{\epsilon} = 1, \beta_1 = \beta_2 = 1$, that is when all tori are untilted, we recover the wrapping number solution proposed in \[16\]. Note that we
Table 1: Chiral spectrum of the four stack D6-brane N=1 Supersymmetric Standard Model-I with its $U(1)$ charges.

| Matter Fields | Representation | Intersection | $Q_a$ | $Q_c$ | $Q_d$ | $Y$ |
|---------------|----------------|--------------|-------|-------|-------|-----|
| $Q_L$         | 3(3, 2)        | $(ab), (ab^*)$ | 1     | 0     | 0     | 1/6 |
| $U_R$         | 3(3, 1)        | $(ac)$        | -1    | 1     | 0     | -2/3 |
| $D_R$         | 3(3, 1)        | $(ac^*)$      | -1    | -1    | 0     | 1/3 |
| $L$           | 3(1, 2)        | $(db), (db^*)$ | 0     | 0     | 1     | -1/2 |
| $N_R$         | 3(1, 1)        | $(dc)$        | 0     | 1     | -1    | 0   |
| $E_R$         | 3(1, 1)        | $(dc^*)$      | 0     | -1    | -1    | 1   |
| $H_d$         | $\frac{1}{\zeta_1, \zeta_2}$ (1, 2) | $(cb^*)$      | 0     | 1     | 0     | -1/2 |
| $H_u$         | $\frac{1}{\zeta_1, \zeta_2}$ (1, 2) | $(cb)$        | 0     | -1    | 0     | 1/2 |

have allowed for the introduction of a NS B-field, that makes the tori tilted along the second and the third tori. Lets us define [15, 8] the supersymmetry vectors

$$r_1 = \pm \frac{1}{2}(- - - -),$$
$$r_2 = \pm \frac{1}{2}(- + + -),$$
$$r_3 = \pm \frac{1}{2}(+ - + -),$$
$$r_4 = \pm \frac{1}{2}(+ + - -),(2.8)$$

These vectors may help us to identify the supersymmetries shared between by the different branes of our models and the orientifold O6 plane. In general a D6-brane, placed in a compactification of type IIA in a toroidal background being a $\frac{1}{2}$ BPS state will preserve N = 4 SUSY. If an orientifold plane is added in our configuration then there are non-trivial consequences for the supersymmetries preserved by the D6-brane, as now the brane may have to share some supersymmetries with its orientifold images and subsequently the orientifold plane. In fact, each brane will now share a N=2 SUSY with its orientifold image. Clearly the O6 plane preserves the SUSY described by the vector $r_1$. On the other hand the intersection between the D6-branes a’, b’ preserves exactly the supersymmetries that are common between them; each brane sharing some supersymmetries with the orientifold. For models based on toroidal orientifolds of type IIA [1], is also possible to allow for non-supersymmetric D6-brane configurations that allow each intersection to preserve either the same or a different supersymmetry with
the rest of the intersections [15, 16]. This is to be contrasted with the constructions of SM’s of [5, 6] where only the intersections where the right handed neutrino was localized were N=1 supersymmetric. In the models we discuss in the present work the former case is realized. Thus the present constructions will allow the same N=1 SUSY to be equally shared by all intersections, achieving a fermion-boson mass degeneracy at every intersection locally, but not globally, as globally the constructions may be non-supersymmetric (N=0).

The Standard Models described by the wrappings (2.7) have all intersections respecting the same N=1 supersymmetry - see table (3) - when the complex structure moduli \( \chi_i = R_i^{(i)} \) satisfy

\[
\beta_i \cdot \chi_2 = \beta_2 \cdot \chi_3 ,
\]

where we have defined

\[
\alpha_1 = \tan^{-1}(3\rho^2 \beta_1 \chi_1)
\]

Let us now focus our attention to the gauge symmetry of the models. After the implementation of the Green-Schwarz anomaly cancellation mechanism, the actual gauge group of the N=1 local SUSY Standard Model I becomes \( SU(3) \times SU(2)_L \times U(1)_Y \times U(1)_{X} \), where

\[
U(1)_{X} = 3Q_a - 9Q_d + 10Q_c
\]

and we have assumed that the brane \( b \) has been brought on top of its orientifold image \( b^* \). The extra \( U(1)_{X} \) generator may be broken by giving a vev to the right handed

| \( N_i \) | \( (n_1^i, m_1^i) \) | \( (n_2^i, m_2^i) \) | \( (n_3^i, m_3^i) \) |
|---|---|---|---|
| \( N_a = 3 \) | \( (1, 0) \) | \( (1/\rho, 3\rho \epsilon \beta_1) \) | \( (1/\rho, -3\rho \epsilon \beta_2) \) |
| \( N_b = 1 \) | \( (0, \epsilon \bar{\epsilon}) \) | \( (1/\beta_1, 0) \) | \( (0, -\epsilon) \) |
| \( N_c = 1 \) | \( (0, \epsilon) \) | \( (0, -\epsilon) \) | \( (\epsilon / \beta_2, 0) \) |
| \( N_d = 1 \) | \( (1, 0) \) | \( (1/\rho, 3\rho \epsilon \beta_1) \) | \( (1/\rho, -3\rho \epsilon \beta_2) \) |

Table 2: General wrapping numbers of Model I, giving rise to the N=1 Standard Model. The wrappings depend on the parameters \( \rho = 1, 1/3; \epsilon = \epsilon = \pm 1 \) and the NS background on the last two tori \( \beta_i = 1 - b^i = 1, 1/2 \). Hence they describe eight different sets of wrappings giving rise to the same intersection numbers.

\[\text{and we have assumed that the brane } b \text{ has been brought on top of its orientifold image } b^* \]. The extra \( U(1)_{X} \) generator may be broken by giving a vev to the right handed
neutrino \( \nu_R \). A comment is in order at this point. We note that the initial gauge symmetry of the models may be further enhanced \(^5\) to the Pati-Salam

\[ SU(4) \times SU(2)_L \times SU(2)_R \]  

(2.12)

if the brane \( d \) is brought on top of brane \( a \) and the brane \( c \) is brought on top of its orientifold image. In this case, breaking the Pati-Salam symmetry \( SU(4) \to SU(3) \times U(1) \) corresponds to moving apart the branes \( a, d \) along different points of the complex plane; giving vev’s to the adjoint multiplets localized in the intersection between the branes. We should also note that when the gauge symmetry of the theory is in the form of a GUT group, like of the Pati-Salam type (2.12), and one necessarily has to have some sectors preserving a SUSY, one should expect some of the branes \( a, b, c, d \) to form angles \( \pi/4 \) with respect to the orientifold plane. This have been firstly observed in all the non-supersymmetric toroidal Pati-Salam orientifold constructions of [8] with only the SM at low energy - where on phenomenological grounds some sectors were N=1 SUSY preserving - and also observed lately in a Pati-Salam N=1 SUSY GUT example in [21].

Assuming the moduli fixing relation (2.9) the model shares in each intersection the same N=1 supersymmetry with the orientifold plane. On the other hand, as have been also discussed in [35, 36] adjusting the complex structure moduli slightly off their SUSY values results in the breaking of supersymmetry [These issues will be discussed in detail in section 6]. The wrapping numbers corresponding to the D6-brane configuration of table (1), do not satisfy RR tadpoles. In general RR tadpoles may be satisfied in a number of ways. In the present models the D6-brane configuration of table (1) is free of any gauge and mixed U(1)-gauge anomalies. Thus RR tadpoles may be cancelled by adding a non-factorizable (NF) D6 brane which has intersections with the SM ‘visible’ branes but where the produced chiral fermions may become massive from existing stringy couplings and simultaneously don’t contribute to anomalies. In the latter case this messenger sector may contain \( N_f \) flavours of chiral superfields \( \Phi_I, \bar{\Phi}_i \) transforming in the representations \( \bar{r} + \bar{\bar{r}} \) of the messenger gauge group. Hence the combined system of the D6-brane/NF-brane may be non-supersymmetric and the messenger sector may be responsible for breaking the supersymmetry through gauge mediation [40]. The same conclusion may be reached for local N=1 SM’s II and III. These issues will be further analyzed in section 8.

- The N=1 Higgs system

\(^5\)as have been already noted in [16]
Table 3: The N=2 supersymmetries shared by the different intersections for the four stack N=1 Supersymmetric Standard Model I.

The number of Higgses present depends on the number of tilted tori in the models. Their charges may be read easily from table (1). At the level of effective theory these Higgses will appear as a mixture of the fields $H^+ = H_u + H_d^*$, $H^- = H_d + H_u^*$. Thus we can have either one pair of Higgs fields with $\beta_1 = \beta_2 = 1$ in which case the chiral fermions of table (1) represent the spectrum of the MSSM, or $\beta_1 = 1/2, \beta_2 = 1/2$ which is doubling the Higgs content of MSSM. Another choice will be when $\beta_1 = \beta_2 = 1/2$ (or $\beta_1 = 1/2, \beta_2 = 1$), where only the one of the Higgs doublet of the MSSM is doubled.

3 N=1 local SM vacua from five stacks of Intersecting D6-branes - Model II

In this section, we will analyze some five stack D6-brane models which are constructed in such a way that the charged lepton and neutrino generations are localized in different intersections. As it will be explained later this has non-trivial consequences for the Yukawa couplings of the models. Our main aim in this section is to exhibit the basic properties of these D6-brane configurations.

3.1 Supersymmetric Standard Model II

Our D6-brane configuration may be seen in table (4). The initial generic gauge symmetry of the models is

$$U(3)_a \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e , \quad (3.1)$$

further enhanced - when the brane b is brought on top of its orientifold image - to

$$U(3)_a \times SP(2)_b \times U(1)_c \times U(1)_d \times U(1)_e , \quad (3.2)$$
The intersection numbers giving rise to the chiral spectrum of the SM in table (4) may be described by

| Matter Fields | Representation | Intersection | $Q_a$ | $Q_c$ | $Q_d$ | $Q_e$ | Y |
|--------------|----------------|--------------|-------|-------|-------|-------|---|
| $Q_L$        | 3(3, 2)        | $(ab), (ab^*)$ | 1     | 0     | 0     | 0     | 1/6 |
| $U_R$        | 3(3, 1)        | $(ac)$        | -1    | 1     | 0     | 0     | -2/3 |
| $D_R$        | 3(3, 1)        | $(ae^*)$      | -1    | -1    | 0     | 0     | 1/3 |
| $L$          | 2(1, 2)        | $(db), (db^*)$ | 0     | 0     | 1     | 0     | -1/2 |
| $l_L$        | (1, 2)         | $(be), (be^*)$ | 0     | 0     | 0     | 1     | -1/2 |
| $N_R$        | 2(1, 1)        | $(dc)$        | 0     | 1     | -1    | 0     | 0   |
| $E_R$        | 2(1, 1)        | $(dc^*)$      | 0     | -1    | -1    | 0     | 1   |
| $\nu_R$      | (1, 1)         | $(ce)$        | 0     | 1     | 0     | -1    | 0   |
| $e_R$        | (1, 1)         | $(ce^*)$      | 0     | -1    | 0     | -1    | 1   |
| $H_d$        | $\frac{1}{\beta_1 \beta_2}$ (1, 2) | $(cb^*)$ | 0     | 1     | 0     | 0     | -1/2 |
| $H_u$        | $\frac{1}{\beta_1 \beta_2}$ (1, 2) | $(cb)$ | 0     | -1    | 0     | 0     | 1/2 |

Table 4: Chiral spectrum of the five stack D6-brane Model-II with its $U(1)$ charges. The hypercharge is defined as $Q_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d - \frac{1}{2}Q_e$.

\[
\begin{align*}
I_{ab} &= 3, & I_{ab^*} &= 3, & I_{ac} &= -3, & I_{ac^*} &= -3, \\
I_{db} &= 2, & I_{db^*} &= 2, & I_{be} &= -1, & I_{be^*} &= 1, \\
I_{dc} &= -2, & I_{dc^*} &= -2, & I_{ce} &= 1, & I_{ce^*} &= -1, \\
I_{bc} &= -\frac{1}{\beta_1 \beta_2}, & I_{bc^*} &= \frac{1}{\beta_1 \beta_2}. 
\end{align*}
\]

(3.3)

The D6-brane wrappings satisfying these intersection numbers are given in table (5).

In the present SM’s II the global symmetries of the SM get identified as

\[
Q_a = 3B, \quad L = Q_d + Q_e, \quad Q_c = 2I_R
\]

(3.4)

One can easily confirm that if the condition (2.9) holds then the brane configuration respects the same N=1 SUSY at every intersection. In fact, the supersymmetries preserved by the brane content of the models may be found in table (6). The angle structure of the branes with respect to the orientifold planes may be seen in table (6), where we have defined

\[
\alpha_2 = tan^{-1}(\beta_1 \chi_2), \quad \gamma_1 = tan^{-1}(2\beta_1 \chi_2)
\]

(3.5)
Table 5: Wrapping numbers of Model II, giving rise to the N=1 Supersymmetric Standard Model. These wrappings depend on the phase parameters $\epsilon = \tilde{\epsilon} = \pm 1$, the parameters $\beta_1 = \beta_2 = 1, 1/2$ parametrizing the NS B-field and the parameter $\rho = 1, 1/3$.

and

$$F_L = (4, 2, 1), \quad \bar{F}_R = (4, 1, \bar{2})$$  \hspace{1cm} (3.6)

the matter multiplets (see [8] for detailed realizations of the Pati-Salam GUT in intersecting brane worlds and classes of models with only the SM at low energy) of the Pati-Salam $SU(4)_C \times SU(2)_L \times SU(2)_R$ GUT.

By implementing the Green-Schwarz anomaly cancellation mechanism [4] and analyzing the U(1) BF couplings to the RR fields

$$N_a \cdot m^1_a \cdot m^2_a \cdot m^3_a \cdot \int_{M_4} B^0_2 \wedge F_a, \quad N_a \cdot m^I_a \cdot m^J_a \cdot m^K_a \cdot \int_{M_4} B^I_2 \wedge F_a,$$

we find, the rest of the couplings having a zero strength with the RR fields,

$$B^1_2 \wedge (\epsilon \beta_1) \cdot (9F^a + 2F^d + F^e),$$
$$B^3_2 \wedge (-\tilde{\epsilon} \beta_2) \cdot (9F^a + 2F^d + F^e);$$  \hspace{1cm} (3.8)

and thus conclude that only one anomalous U(1), $9Q_a + 2Q_d + Q_e$ gets massive by having a non-zero coupling to RR fields, the other three U(1)’s

$$Q^{(1)} = \frac{1}{6}(Q_a - 3Q_d - 3Q_e) - \frac{1}{2}Q_c, \quad Q^{(2)} = \frac{1}{6}(Q_a - 3Q_d - 3Q_e) + \frac{19}{18}Q_c,$$
$$Q^{(3)} = (-\frac{3}{28}Q_a + Q_d - \frac{29}{28}Q_e)$$  \hspace{1cm} (3.9)

remain massless [See [34] for further consequences of for gauge field anomalies in general orientifold models]. The $Q^{(1)}$ generator is identified as the hypercharge in the models.
Figure 1: Gauge group enhancement (and vice-versa) in the five stack quiver to a Pati-Salam
$SU(4) \times SU(2)_L \times SU(2)_R \times U(1)^d$. The brane d on the left is shown to different positions.

We note that at this point the actual gauge group of the models is $SU(3) \times SU(2) \times U(1)_Y \times Q^{(2)} \times Q^{(3)}$. The additional U(1)’s $Q^{(2), Q^{(3)}}$ that survive massless the Green-Schwarz anomaly cancellation mechanism, can be broken by giving vev’s to the right handed sneutrinos, $sN_R$’s, localized in the $dc$ intersection. Alternatively one may choose to break them by giving vev’s to two linear combinations of $sN_R$’s, $s\nu_R$’s from neutrinos localized in the intersections $dc, ce$ respectively. This mechanism may be described by the presence of the Fayet-Iliopoulos terms for the anomalous U(1)’s and will be described later on.

3.2 Gauge Group Enhancement

The initial gauge symmetry of the models, before implementation of Green-Schwarz mechanism, may be further enhanced in a number of different ways. By choosing for example the values $\rho = 1/3$, $\epsilon = \tilde{\epsilon}$ in table 5, it is clear that as the branes $a, e$ are parallel, by placing brane a on top of brane e, we may have a symmetry enhancement from

$$U(3)_a \times U(1)_e \rightarrow U(4) \quad (3.11)$$

Also, if branes b, c are brought on top of their orientifold images, then the full gauge symmetry of the models may be further enhanced to

$$U(4) \times SU(2)_L^b \times SU(2)_R^c \times U(1)_d \quad (3.12)$$
4 The N=1 Standard Model from Six-stacks of Intersecting D6-branes - Model III

In this section, we will discuss the extension of the four and five stack supersymmetric Standard Models I, II to their closest generalizations with six stacks of D6-branes at the string scale. These models constitute the maximum allowed deformations as it may easily seen from the D-brane configurations of tables (1) and (4). The initial gauge group structure is an

\[ U(3) \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e \times U(1)_f \]  

(4.1)

further enhanced to

\[ U(3) \times SU(2) \times U(1)_c \times U(1)_d \times U(1)_e \times U(1)_f \]  

(4.2)

when brane b is brought on top of its orientifold image.

The chiral spectrum of the SUSY SM gets localized at the various intersections as described in the D-brane configuration of table (7). The observed chiral spectrum gets reproduced by the following intersection numbers

\[ I_{ab} = 3, \quad I_{ab^*} = 3, \quad I_{ac} = -3, \quad I_{ac^*} = -3, \]
\[ I_{db} = 1, \quad I_{db^*} = 1, \quad I_{eb} = 1, \quad I_{eb^*} = 1, \]
\[ I_{fb} = 1, \quad I_{fb^*} = 1, \quad I_{dc} = -1, \quad I_{dc^*} = -1, \]
\[ I_{bc} = -\frac{1}{\beta_1 \beta_2}, \quad I_{bc^*} = \frac{1}{\beta_1 \beta_2}, \]  

(4.3)

where all other intersections are vanishing. The global symmetries of the SM may be expressed in terms of the U(1) symmetries of the models as

\[ Q_a = 3B, \quad L = Q_d + Q_e + Q_f, \quad Q_c = 2I_R \]  

(4.4)
| Matter Fields | Representation | Intersection | $Q_a$ | $Q_c$ | $Q_d$ | $Q_e$ | $Q_f$ | $Y$ |
|---------------|----------------|--------------|-------|-------|-------|-------|-------|-----|
| $Q_L$         | $3(3,2)$       | $(ab), (ab^*)$ | $1$   | $0$   | $0$   | $0$   | $0$   | $1/6$ |
| $U_R$         | $3(3,1)$       | $(ac)$       | $-1$  | $1$   | $0$   | $0$   | $0$   | $-2/3$ |
| $D_R$         | $3(3,1)$       | $(ac^*)$     | $-1$  | $-1$  | $0$   | $0$   | $0$   | $1/3$ |
| $L^1$         | $(1,2)$        | $(db), (db^*)$ | $0$   | $0$   | $1$   | $0$   | $0$   | $-1/2$ |
| $L^2$         | $(1,2)$        | $(be), (be^*)$ | $0$   | $0$   | $0$   | $1$   | $0$   | $-1/2$ |
| $N^1_R$       | $(1,1)$        | $(cd)$       | $0$   | $1$   | $-1$  | $0$   | $0$   | $0$   |
| $E^1_R$       | $(1,1)$        | $(cd^*)$     | $0$   | $-1$  | $-1$  | $0$   | $1$   | $1$   |
| $N^2_R$       | $(1,1)$        | $(ce)$       | $0$   | $1$   | $0$   | $-1$  | $0$   | $0$   |
| $E^2_R$       | $(1,1)$        | $(ce^*)$     | $0$   | $-1$  | $0$   | $-1$  | $0$   | $1$   |
| $N^3_R$       | $(1,1)$        | $(cf)$       | $0$   | $1$   | $0$   | $0$   | $-1$  | $0$   |
| $E^3_R$       | $(1,1)$        | $(cf^*)$     | $0$   | $-1$  | $0$   | $0$   | $-1$  | $1$   |
| $H_u$         | $\frac{1}{\sqrt{3}}(1,2)$ | $(cb^*)$     | $0$   | $1$   | $0$   | $0$   | $0$   | $-1/2$ |
| $H_d$         | $\frac{1}{\sqrt{3}}(1,2)$ | $(cb)$       | $0$   | $-1$  | $0$   | $0$   | $0$   | $1/2$ |

Table 7: Standard model chiral spectrum of the six stack string scale D6-brane N=1 SUSY Model III together with its $U(1)$ charges.

Analyzing the $U(1)$ BF couplings to the RR fields we conclude that the anomalous $U(1)$, $9Q_a + Q_d + Q_e + Q_f$ gets massive by having a non-zero coupling to RR fields, and other four $U(1)$’s remain massless, the following

$$
\tilde{Q}^{(Y)} = \frac{1}{6}(Q_a - 3Q_d - 3Q_e - 3Q_f) - \frac{1}{2}Q_c, \quad (4.5)
$$

$$
\tilde{Q}^{(1)} = \frac{1}{6}(Q_a - 3Q_d - 3Q_e - 3Q_f) + \frac{28}{18}Q_c, \quad (4.6)
$$

$$
\tilde{Q}^{(2)} = ( -2Q_d + Q_e + Q_f), \quad (4.7)
$$

$$
\tilde{Q}^{(3)} = (Q_e - Q_f). \quad (4.8)
$$

The $\tilde{Q}^{(Y)}$ is identified as the hypercharge in the models. The breaking of the remaining $U(1)$’s proceeds once the sneutrino’s, $sN_R$’s get a vev. An obvious choice will be for $\tilde{Q}^{(1)}$ to be broken by a vev from $sN^1_R$; $\tilde{Q}^{(2)}$ to be broken by a vev from $sN^2_R$, and $\tilde{Q}^{(3)}$ to be broken by a vev from $sN^3_R$. Appropriate combinations of vev’s by $sN_R$’s can also have the same effect on the Higgsing of the $U(1)$’s. The SUSY’s shared by each intersection with the orientifold plane may be seen in table (9).

The gauge symmetry may be enhanced to $SU(5)$ when (we choose $\rho = 1/3$) the branes d, e are brought on top of brane a; with $SU(3) \times U(1)_d \times U(1)_e \rightarrow SU(5)$. In
\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|}
\hline
$N_i$ & $(n_1^i, m_1^i)$ & $(n_2^i, m_2^i)$ & $(n_3^i, m_3^i)$ \\
\hline
$N_a = 3$ & $(1, 0)$ & $(1/\rho, 3\rho\epsilon\beta_1)$ & $(1/\rho, -3\rho\epsilon\beta_2)$ \\
$N_b = 2$ & $(0, \epsilon\tilde{\epsilon})$ & $(1/\beta_1, 0)$ & $(0, -\tilde{\epsilon})$ \\
$N_c = 1$ & $(0, \epsilon)$ & $(0, -\epsilon)$ & $(\epsilon/\beta_2, 0)$ \\
$N_d = 1$ & $(1, 0)$ & $(1, \epsilon\beta_1)$ & $(1, -\epsilon\beta_2)$ \\
$N_e = 1$ & $(1, 0)$ & $(1, \epsilon\beta_1)$ & $(1, -\epsilon\beta_2)$ \\
$N_f = 1$ & $(1, 0)$ & $(1, \epsilon\beta_1)$ & $(1, -\epsilon\beta_2)$ \\
\hline
\end{tabular}
\caption{Wrapping numbers of Model III, giving rise to the N=1 standard model. The parameter $\rho = 1, 1/3$ describes two different sets of wrappings giving rise to the same intersection numbers.}
\end{table}

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|}
\hline
$Brane$ & $\theta_a^1$ & $\theta_a^2$ & $\theta_a^3$ & SUSY preserved \\
\hline
$N_a = 3$ & 0 & $\alpha_1$ & $-\alpha_1$ & $r_2, r_1$ \\
$N_b = 2$ & $\frac{\pi}{2}$ & 0 & $-\frac{\pi}{2}$ & $r_3, r_1$ \\
$N_c = 1$ & $\frac{\pi}{2}$ & $-\frac{\pi}{2}$ & 0 & $r_3, r_4$ \\
$N_d = 1$ & 0 & $\alpha_2$ & $-\alpha_2$ & $r_2, r_1$ \\
$N_e = 1$ & 0 & $\alpha_2$ & $-\alpha_2$ & $r_2, r_1$ \\
$N_f = 1$ & 0 & $\alpha_2$ & $-\alpha_2$ & $r_2, r_1$ \\
\hline
\end{tabular}
\caption{The N=2 supersymmetries shared by the different intersections for the six stack N=1 SUSY Model-III.}
\end{table}

In this case the full gauge group becomes

$$SU(5) \times U(1)_b \times U(1)_c \times U(1)_f$$ \quad (4.9)

or even further enhanced to

$$SU(6) \times U(1)_b \times U(1)_c$$ \quad (4.10)

if brane $f$ is brought on top of brane $a$. Alternatively, we could choose to further enhance the gauge group by bringing together the $b, c$ branes on top of their orientifold images.
The final gauge group in this case will be

$$SU(5) \times SU(2)_L \times SU(2)_R \times U(1)_f \quad (4.11)$$

or if brane $f$ is located on top of brane $a$ to

$$SU(6) \times SU(2)_L \times SU(2)_R \quad (4.12)$$

If brane $b$ (or alternatively brane $c$) is not brought on top of its respective orientifold image $b^*$, then the gauge group may be $SU(5) \times U(1)_b \times SU(2)_R \times U(1)_f$ ($SU(5) \times SU(2)_L \times U(1)_c \times U(1)_f$). The reverse procedure of the gauge group enhancement that is the actual splitting of i.e. the $SU(5) \rightarrow SU(3) \times U(1) \times U(1)$ corresponds to adjoint breaking.

### 4.1 Yukawa Couplings

The tree level Yukawa couplings for the N=1 SM’s I, II III have several differences as the charged lepton and neutrino species are localized to different intersection points. The Higgs fields in all models are localized between the $b,c$ and $b^*, c^*$, branes and are shown in table (10). We assume the existence of the D-brane configuration of tables (1), (4), (7).

| Intersection | EW Higgs | $Q_c$ | $Y$ |
|--------------|----------|-----|-----|
| $\{bc\}$    | $h_1$    | $-1$ | $1/2$ |
| $\{bc\}$    | $h_2$    | $1$  | $-1/2$ |
| $\{bc^*\}$  | $H_1$    | $-1$ | $1/2$ |
| $\{bc^*\}$  | $H_2$    | $1$  | $-1/2$ |

Table 10: Higgs fields responsible for electroweak symmetry breaking in the N=1 SUSY type I, II and III models.

The tree level Yukawa couplings, that are being allowed by gauge and charge conservation invariance, are common for the quark sector to models I, II and III - and they are given by

$$\mathcal{L}^I \propto h^i_j Q_L U_R^i h_1 + h^j_D Q_L D_R^j H_2 \quad (4.13)$$

$^6$When $\beta_1 = \beta_2$ one may locate the brane $f$ on top of its orientifold image, extending the gauge group to $SU(5) \times SU(2)_L \times SU(2)_R \times SU(2)_f$. 

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Also the tree level Yukawa’s for the charged lepton and neutrino sectors of models I, II, III, are given respectively by

\[ \mathcal{L}^I \propto y_{iL}^j L^i E_R^j H_2 + Y_{N_R}^j N_R^i h_1, \]

\[ \mathcal{L}^{II} \propto y_{iL}^j L^i E_R^i H_2 + y_{\nu_L}^i \nu_L h_2 + y_{N_R}^j L N_R^i h_1 + y_{\nu_R}^i \nu R h_1, \]

\[ \mathcal{L}^{III} \propto \left( \sum_{k=1}^3 L^k E_R^i h_2 \right) + \left( \sum_{k=1}^3 L^k N_R^i h_1 \right) \] (4.14)

where \( i=1,2; \ j =1,2,3. \)

5 Supersymmetry breaking and sparticle masses

As have been noted earlier for special values of the complex structure \( \chi \) moduli all intersections in the models may share the same \( N=1 \) SUSY with the orientifold plane, that is the same \( N=1 \) SUSY is preserved at each intersection. In general, as has been noted in [36] where configurations with non-zero B-field and two D-branes each one preserving \( N=1 \) supersymmetry were used, by turning on a Fayet-Iliopoulos (FI) term one may able to break supersymmetry. These FI terms may be computed by the use of the supersymmetric completion of the Chern-Simon

\[ \sim \sum_{i=1}^{b_3} \sum_{\alpha=1}^{l} \int d^4x \; Z_{ij} B_i \wedge F_{\alpha} \] (5.1)

couplings [38] involved in the Green-Schwarz anomaly cancellation mechanism appearing with the usual auxiliary \(^7\) field \( D_{\alpha}. \)

\[ \sim \sum_{i=1}^{b_3} \sum_{\alpha=1}^{l} \int d^4x \; Z_{ij} \frac{\partial K}{\partial \phi^i} D_{\alpha} \] (5.2)

For an effective gauge theory involving D-branes the presence of a FI term is then interpreted as coupling the complex structure moduli as FI-terms in the effective theory [35, 36]. Assuming a small departure of the complex structure from the supersymmetric situation - the SUSY wall - the FI term

\[ V_{FI} = \frac{1}{2g_4^2} \left( \sum_{\alpha} \phi_{\alpha}^i |\phi^i| + \xi \right)^2 \] (5.3)

captures the leading order effect to the mass of the scalar - we assume the existence of two branes \( D6_\sigma, D6_\tau \)

\[ \alpha' m_{\sigma\tau}^2 = -q_\sigma \xi_\sigma - q_\tau \xi_\tau \] (5.4)

\(^7\)By \( \phi_{\alpha} \) we denote the superpartners of the Hodge duals of the RR fields; \( K \) is the Kähler potential.
leaving at an intersection $\sigma \tau$ - as it is described from string theory [7]. Related issues have also been examined in [15, 16]. In this section, we will make use of FI terms in order to break N=1 SUSY at the various intersections of the N=1 models I, II, III.

**N=1 SUSY Model I & II**

| Brane | $\theta^1_a$ | $\theta^2_a$ | $\theta^3_a$ | approx. SUSY preserved |
|-------|-------------|-------------|-------------|------------------------|
| $a$   | $\frac{\pi}{2}$ | $\alpha_1 + \delta_a$ | $-\alpha_1$ | $r_1, r_4$ |
| $b$   | $\frac{\alpha_1 + \delta_b}{\alpha_1}$ | 0 | $-\frac{\pi}{2}$ | $r_2, r_4$ |
| $c$   | $\frac{\alpha_1 + \delta_c}{\alpha_1}$ | $-\frac{\alpha_1}{2}$ | 0 | $r_3, r_4$ |
| $d$   | 0 | $\gamma_1 + \delta_d$ | $-\gamma_1$ | $r_1, r_4$ |
| $e$   | 0 | $\alpha_2 + \delta_e$ | $-\alpha_2$ | $r_1, r_4$ |

Table 11: Angle structure; coupling the complex structure moduli to open string modes as Fayet-Iliopoulos terms for the five stack SUSY Model II describing the D-brane configuration of table (4).

As the N=1 SUSY models I, II, III are supersymmetric for the same value of complex structure moduli (2.9) we assume that the complex structure moduli departs only slightly from its line of marginal stability as

$$U^2 = \frac{\beta_2}{\beta_1} U^3 + \delta_2,$$  \hspace{1cm} (5.5)

where $\delta_2$ parametrizes the deviation from the supersymmetric situation. In this case N=1 supersymmetry will be broken.

For N=1 SUSY Model I, assuming that the departure from the supersymmetric limit is described by table (13), one finds

$$\delta'_a = \delta'_b = \delta_a, \hspace{0.5cm} \delta'_b = \delta'_c = 0.$$  \hspace{1cm} (5.6)

The sparticle masses for N=1 SUSY model I may be seen in table (12)

In table (11) the deformed angles at the various intersections of SUSY Model II are shown. Therefore one may find that

$$\delta_a = \frac{3\rho^2 \beta_1 \delta_2}{[1 + (3\rho^2 \beta_2 U^3)^2]}, \hspace{0.5cm} \delta_d = \frac{2\beta_1 \delta_2}{[1 + (2\beta_2 U^3)^2]}, \hspace{0.5cm} \delta_c = \frac{\beta_1 \delta_2}{[1 + (\beta_2 U^3)^2]}, \hspace{0.5cm} \delta_b = \delta_c = 0.$$  \hspace{1cm} (5.7)

The Fayet-Iliopoulos terms can be expressed in terms of the angle deviations from the marginal line. Thus for N=1 Model II, we find that

$$\xi_a = -\frac{3\rho^2 \beta_1 \delta_2}{2[1 + (3\rho^2 \beta_2 U^3)^2]}, \hspace{0.5cm} \xi_d = -\frac{\beta_1 \delta_2}{[1 + (2\beta_2 U^3)^2]}, \hspace{0.5cm} \xi_c = -\frac{\beta_1 \delta_2}{2[1 + (\beta_2 U^3)^2]}.$$  \hspace{1cm} (5.8)
Table 12: Sparticle masses from Fayet-Iliopoulos terms for the four stack quiver of model I seen in figure 2.

Table 13: Angle structure; coupling the complex structure moduli to open string modes as Fayet-Iliopoulos terms for the four stack SUSY Model I.

**N=1 SUSY Model III**

For the N=1 SUSY Model III we assume that the departure from the supersymmetric locus is described in terms of the parameters seen in appendix A - table (17). The dependence of the squark masses on the Fayet-Iliopoulos terms can be extracted easily and it can be seen in table (18). The resulting deviations are:

\[
\tilde{\delta}_a = \delta_a, \quad \tilde{\delta}_b = \tilde{\delta}_c = 0, \quad \tilde{\delta}_d = \tilde{\delta}_e = \delta_e \quad (5.9)
\]

Clearly in all the above models the FI terms alone are not enough to give to all the squarks positive (mass)\(^2\), in order to avoid unwanted charge and colour breaking minima. The resolution of this puzzle may be found by taking into account loop effects for the N=1 local SUSY models I, II, III.
In this section, we will show that gauge breaking transitions (GBT) that allow for the 
breaking of the gauge group without changing its chiral content exist in models I, II, III. In particular one is able to switch between models with different number of stacks. 
The end point of the recombination process is the - four stack - N=1 local model of 
table (1), where all U(1)'s beyond hypercharge are massive\(^8\). That means that the 
breaking of the extra U(1)'s in the five, six stack local N=1 models occurs at the string 
scale.

GBT transitions first appeared in brane configurations [14] with D5-branes wrapping 
two-cycles in a \(T^4 \times C/Z_N\) orientifold of type IIB [13]. As has been observed in 
[14] GBT's also hold for the toroidal examples of [5, 6]. In these models [14, 5, 6] one 
is able to show that GBT transitions are able to switch on - field directions - between 
different D-brane stack configurations that have the distinct characteristic that they 
all result in classes of models, which have only the (non-SUSY) SM at low energy. 
These field directions cannot be described using only renormalizable couplings, thus 
necessarily their description involves higher order couplings. For some recent work on 
the field theoretical description of BR effects see [39]. 

We note that GBT transitions appear as a direct consequence of the specific way we

\(^8\)due to the Green-Schwarz mechanism

| Sparticle | \((\theta^1, \theta^2, \theta^3)\) | Sector | \((\text{mass})^2\) |
|----------|-------------------------------|--------|-----------------|
| \(Q_L\)  | \((\frac{\pi}{2} + \delta_b, -\alpha_1 - \delta_a, \alpha_1 - \frac{\pi}{2})\) | \((ab)\) | \(\frac{1}{2}(\delta_a - \delta_b)\) |
| \(U_R\)  | \((\frac{\pi}{2} + \delta_c, -\alpha_1 - \delta_a - \frac{\pi}{2}, \alpha_1)\) | \((ac)\) | \(\frac{1}{2}(\delta_c - \delta_a)\) |
| \(D_R\)  | \((-\frac{\pi}{2} - \delta_c, -\alpha_1 - \delta_a + \frac{\pi}{2}, \alpha_1)\) | \((ac*)\) | \(-\frac{1}{2}(\delta_c + \delta_a)\) |
| \(L\)    | \((\frac{\pi}{2} + \delta_b, -\gamma_1 - \gamma_d, \gamma_1 - \frac{\pi}{2})\) | \((db)\) | \(\frac{1}{2}(\delta_d - \delta_b)\) |
| \(l_L\)  | \((-\frac{\pi}{2} - \delta_b, \alpha_2 + \delta_c, \frac{\pi}{2} - \alpha_2)\) | \((be)\) | \(\frac{1}{2}(\delta_c - \delta_b)\) |
| \(N_R\)  | \((\frac{\pi}{2} + \delta_c, -\gamma_1 - \gamma_d - \frac{\pi}{2}, \gamma_1)\) | \((dc)\) | \(\frac{1}{2}(\delta_c - \delta_d)\) |
| \(E_R\)  | \((-\frac{\pi}{2} - \delta_c, -\gamma_1 - \gamma_d + \frac{\pi}{2}, \gamma_1)\) | \((dc*)\) | \(-\frac{1}{2}(\delta_c + \delta_d)\) |
| \(v_R\)  | \((-\frac{\pi}{2} - \delta_c, \frac{\pi}{2} + \alpha_2 + \delta_c, -\alpha_2)\) | \((ce)\) | \(\frac{1}{2}(\delta_c - \delta_e)\) |
| \(e_R\)  | \((-\frac{\pi}{2} - \delta_c, \frac{\pi}{2} - \alpha_2 - \delta_c, \alpha_2)\) | \((ce*)\) | \(-\frac{1}{2}(\delta_c + \delta_e)\) |

Table 14: Sparticle masses from Fayet-Iliopoulos terms for the five stack quiver
construct these D6-brane (or D5) configurations. In this respect, these transitions, that allow movement between models with different number of stacks, differ from the small instanton/chirality changing transitions (in the T-dual picture) that were considered in [7]. In the latter models, brane recombination (BR) examples were causing a change in the chiral content among the different BR phases. We note that BR as an alternative for electroweak Higgs breaking mechanism have been considered in [15].

### 6.1 Six stack BRSM's

Let us recombine the branes $e, f$ in the six stack $N=1$ SUSY Model III. That is we assume that the branes $e, f$ recombine into a new brane $\tilde{e}$,

$$e + f = \tilde{e} \quad (6.1)$$

The intersection numbers appearing after brane recombination, may be found by assuming linearity under homology change, may be seen in table (16) and correspond exactly to the U(1) particle assignments of table (15).

The reader can be easily convinced that these are exactly the U(1) particle intersection numbers of table (4), but where the branes $d, e$ have been interchanged. Thus no new examples are being generated by this BR. Instead, let us try the BR direction;

---

| Matter Fields | Representation | Intersection | $Q_a$ | $Q_c$ | $Q_d$ | $Q_e$ |
|--------------|----------------|--------------|-------|-------|-------|-------|
| $Q_L$        | $3(3, 2)$      | $(ab), (ab^*)$ | 1     | 0     | 0     | 0     | $1/6$ |
| $U_R$        | $3(3, 1)$      | $(ac)$       | $-1$  | 1     | 0     | 0     | $-2/3$ |
| $D_R$        | $3(3, 1)$      | $(ac^*)$     | $-1$  | $-1$  | 0     | 0     | $1/3$ |
| $L$          | $2(1, 2)$      | $(eb), (eb^*)$ | 0     | 0     | 0     | 1     | $-1/2$ |
| $l_L$        | $(1, 2)$       | $(be), (be^*)$ | 0     | 0     | 1     | 0     | $-1/2$ |
| $N_R$        | $2(1, 1)$      | $(dc)$       | 1     | 0     | 0     | $-1$  | 0     |
| $E_R$        | $2(1, 1)$      | $(dc^*)$     | $-1$  | 0     | 0     | $-1$  | 1     |
| $\nu_R$      | $(1, 1)$       | $(ce)$       | 0     | 1     | $-1$  | 0     | 0     |
| $e_R$        | $(1, 1)$       | $(ce^*)$     | 0     | $-1$  | $-1$  | 0     | 1     |
| $H_{u,d}$    | $(1, 2)$       | $(cb^*), (cb)$ | $\pm 1$ | 0     | 0     | 0     | $\mp 1/2$ |

Table 15: Fermion spectrum of the five stack D6-brane Model-I with its U(1) charges appearing after the brane recombination (6.1).
recombining the branes d, f; with

\[ d + f = \tilde{d} \]  

(6.2)
to see if there are any new examples of vacua left unexplored. In this case, we get exactly

\[
\begin{array}{cccccc}
I_{ab} & = & 3 & I_{ahs} & = & 3 \\
I_{ac} & = & -3 & I_{acs} & = & -3 \\
I_{ec} & = & -2 & I_{ecs} & = & -2 \\
I_{bc} & = & -1 & I_{bcs} & = & -1 \\
I_{db} & = & 1 & I_{bsd} & = & 1 \\
I_{dc} & = & 1 & I_{dcs} & = & 1 \\
I_{dcb} & = & 1 & I_{dcs} & = & 1 \\
I_{dcb} & = & 1 & I_{dcs} & = & 1
\end{array}
\]

Table 16: Intersection numbers appearing in the five stack model coming after the brane recombination (6.1).

the intersection numbers of the five stack N=1 SUSY model of table (4). Moreover, we further assume the BR combination

\[ d + e + f = d' \]  

(6.3)
Our new brane content consists of the \(a, b, c, d'\). In this case, the chiral content appearing after BR is that if the four stack SUSY model of table (1).
6.2 Five stack SM's

Lets us now consider the five stack N=1 SUSY Model II and recombine the branes d, e. That is we assume that

\[ d + e = \tilde{d} \]  

(6.4)

Calculating the intersection numbers of the D-brane models surviving the recombination process we find that the models flow to their four stack counterparts. In terms of the maximal gauge symmetry in the models in this case, the following gauge symmetry pattern is realizable \(^9\)

\[ SU(5) \times SU(2)_L \times SU(2)_R \xrightarrow{d+e \to \tilde{d}} SU(4) \times SU(2)_L \times SU(2)_R \]  

(6.5)

The homology flows allowing GBT between Models I, II, III may be depicted nicely in terms of quiver diagrams [37]. Each quiver will be depicted in terms of the nodes and their associated links. Each node represents the brane content of the brane and its mirror image; the link between two nodes representing the non-trivial intersection between the nodes. For an intersection, respecting N=1 SUSY, between two branes a', b', the link will be associated with the presence of \( I_{ab}(N_a', \bar{N}_b') + I_{ab^*}(N_{a'}, N_{b'}) \) N=1 chiral multiplets. The relevant GBT transitions may be seen in figure 2.

7 Gauge coupling constants

In this section we will calculate the gauge couplings for the models I, II, III. The tree level gauge couplings at the string scale are defined \(^{10}\) by

\[ \frac{4\pi}{g_\alpha^2} = \frac{M_{Pl}}{2\sqrt{2}\kappa_\alpha M_s \sqrt{V_6}}, \]  

(7.1)

where \( V_\alpha \) is the volume of the three-cycle that the \( D6_a \) branes wrap and \( V_6 \) is the volume of the Calabi-Yau manifold. A universal property of the models examined in this work, is that the volumes of the three cycles of the \( b, c \) branes agree as a consequence of the N=1 SUSY condition (2.9) since

\[ \frac{V_b}{V_c} = \frac{\beta_2}{\beta_1} \cdot \frac{U^3}{U^2} \stackrel{N=1 \text{ SUSY}}{=} 1. \]  

(7.2)

- Four stack models

\(^9\)where we have assumed that the brane e has been brought on top of brane a.

\(^{10}\)where \( \kappa_\alpha = 1 \) for \( U(N_a) \) and \( \kappa_\alpha = 2 \) for \( Sp(2N_a)/SO(2N_a) \) and the gauge fields have been \( tr(T_a T_b) = 1/2 \).
The four stack model version of model I when there is no antisymmetric NS B-field present, that is \( \beta_1 = \beta_2 = 1 \), have been proposed in [16]. Some consequences for
gauge coupling unification for these models have been examined in [25]. Let us assume
that the gauge group in the four stack model I is that of
\[ U(3)_c \times SP(2)_b \times U(1)_c \times U(1)_d. \]
Then as the volumes of the \( a, d \) branes are equal, and the gauge

group of the weak SU(2) comes from an \( SP(2) = SU(2)_b \) factor the following relations
among the gauge coupling constants hold
\[
\alpha_d = \alpha_a = \alpha_s, \quad \alpha_c = \frac{1}{2} \alpha_b = \frac{1}{2} \alpha_w \quad (7.3)
\]
and as a consequence the relation
\[
\alpha_Y^{-1} = \frac{2}{3} \alpha_s^{-1} + \frac{1}{\alpha_w} \quad (7.4)
\]
holds at the string scale [25], where \(^{11} \)
\[
\alpha_Y^{-1} = \frac{1}{6} \alpha_a^{-1} + \frac{1}{2} \alpha_c^{-1} + \frac{1}{2} \alpha_d^{-1} \quad (7.5)
\]
\[
\alpha_w^{-1} = \frac{1}{2 \sqrt{2}} \frac{M_{pl}}{M_s} \cdot \frac{\epsilon \sqrt{U_1}}{2 \sqrt{\beta_1 \beta_2}}, \quad (7.6)
\]
\[
\alpha_s^{-1} = \frac{1}{2 \sqrt{2}} \cdot \frac{M_{pl}}{M_s} \cdot \left( \frac{1}{\rho^2} \sqrt{\frac{\beta_1}{\beta_2} U_1 U_3} + 9 \rho^2 (\epsilon \tilde{\epsilon}) \sqrt{\beta_2 \frac{1}{\beta_1} U_3} \right). \quad (7.7)
\]
Note that we have assumed that the brane \( b \) have been placed on top of its orientifold
image that is their homology classes agree, \( \pi_b = \pi_b^* \). The relation (7.4) - which is
compatible with the SU(5) relation \( 1/\alpha_s = 1/\alpha_w = (3/5)1/\alpha_Y \) [ has been derived before
in [25] for the models of [16] ] - is not new in the context of model building in string
compactifications as it has been also derived in the context of type IIB 4D orientifold
compactifications [41]. We can further assume that the initial gauge symmetry in the
models is \( U(4) \times SU(2)_b^L \times SU(2)_b^R \), that is the brane \( d \) has been placed on top of brane
\( a \), and the homology classes of \( b, c \) branes agree \( (V_b \rightarrow V_c) \). It follows
\[
\alpha_d = \alpha_a = \alpha_s, \quad \alpha_c = \alpha_b = \alpha_w \quad (7.8)
\]
and as a consequence the relation
\[
\alpha_Y^{-1} = \frac{2}{3} \alpha_s^{-1} + \frac{1}{2} \alpha_w^{-1} \quad (7.9)
\]
holds at the string scale.

\(^{11}\)we have introduced a non-trivial NS B-field in the results of [25].
Five stack models

Assuming that the initial gauge symmetry at the string scale is given by (3.2) the following relations among the gauge couplings hold

\[ \alpha_{\text{Y}}^{-1} = \frac{1}{6} \alpha_{a}^{-1} + \frac{1}{2} \alpha_{c}^{-1} + \frac{1}{2} \alpha_{d}^{-1} + \frac{1}{2} \alpha_{e}^{-1} \]  
(7.10)

\[ \alpha_{d}^{-1} = \frac{1}{2 \sqrt{2}} \cdot \frac{M_{\text{pl}}}{M_{s}} \cdot \left( \sqrt{\frac{2}{\beta_1} \frac{1}{U_1 U_2 U_3}} + 4(\epsilon \bar{\epsilon}) \beta_1 \beta_2 \sqrt{\frac{2}{\beta_1}} \frac{1}{U_1^3} \right) \]  
(7.11)

\[ \alpha_{e}^{-1} = \frac{1}{2 \sqrt{2}} \cdot \frac{M_{\text{pl}}}{M_{s}} \cdot \left( \sqrt{\frac{2}{\beta_1} \frac{1}{U_1 U_2 U_3}} + (\epsilon \bar{\epsilon}) \beta_1 \beta_2 \sqrt{\frac{2}{\beta_1}} \frac{1}{U_1^3} \right) \]  
(7.12)

\[ \alpha_{c} = \frac{1}{2} \alpha_{b} = \frac{1}{2} \alpha_{w}, \]  
(7.13)

where \( \tilde{\alpha}_{w}, \tilde{\alpha}_{a} \) are given in (7.6), (7.7) respectively.

Six stack models

Let us assume that the initial gauge symmetry in the models is as in (4.2). Then the gauge couplings are as follows:

\[ \tilde{\alpha}_{\text{Y}}^{-1} = \frac{1}{6} \tilde{\alpha}_{a}^{-1} + \frac{1}{2} \tilde{\alpha}_{c}^{-1} + \frac{1}{2} \tilde{\alpha}_{d}^{-1} + \frac{1}{2} \tilde{\alpha}_{e}^{-1} + \frac{1}{2} \tilde{\alpha}_{f}^{-1}, \]  
(7.14)

\[ \tilde{\alpha}_{c} = \frac{1}{2} \alpha_{b} = \frac{1}{2} \alpha_{w}, \quad \tilde{\alpha}_{d}^{-1} = \tilde{\alpha}_{e}^{-1} = \tilde{\alpha}_{f}^{-1}, \]  
(7.15)

\[ \alpha_{d}^{-1} = \frac{1}{2 \sqrt{2}} \cdot \frac{M_{\text{pl}}}{M_{s}} \left( \sqrt{\frac{2}{U_1 U_2 U_3}} + \epsilon \bar{\epsilon} \beta_1 \beta_2 \sqrt{\frac{2}{U_1^3}} U^3 \right) \]  
(7.16)

and \( \alpha_{w}, \alpha_{a} \) are given in (7.6), (7.7) respectively.

As the D6-brane configurations of tables (1), (4), (7) localize the spectrum of N=1 MSSM in cases where \( \beta_1 = \beta_2 = 1 \), the five and the six stack models they may be expected to exhibit \(^{12}\) unification of the gauge coupling constants at a scale of the order of \( 10^{16} \) GeV and for appropriate values of the complex moduli. This is possible as the models possess enough freedom - through their dependence on the string scale and the complex structure moduli.

\(^{12}\)As it has also been exhibited for model I, when \( \beta_1 = \beta_2 = 1 \), in [25].
8 Gauge mediated N=1 SUSY breaking

The models we have constructed in the present work do not satisfy the RR tadpoles (2.5). However as the gauge and mixed U(1) anomalies for the D-brane configurations of tables (1), (4), (7) cancel, RR tadpoles may be cancelled by adding a non-factorizable D6 brane (NF). As the presence of the NF brane is not expected to create fields on intersections that will survive massless to low energies, no new open string 'visible' sectors may be created but massive fields. Also as $StrM^2$ in a supersymmetric theory is expected to vanish one may expect that SUSY breaking is generated beyond tree level. In this case the couplings responsible for breaking N=1 SYSYM may be non-renormalizable. As a consequence of this remark, this phase of the theory is similar to the one we have found in the Pati-Salam models - that possess N=1 supersymmetric subsectors - of [8], where the presence of extra branes was breaking the extra U(1)'s surviving the Green-Schwarz mechanism. In the latter case the extra branes create gauge singlets that make massive through non-renormalizable couplings all exotic fermions that transform under both the extra -plays the role of an 'NF' - brane and the visible SM branes; the latter fields playing the role of messenger fields.

The NF D6-brane in the present constructions is expected to break the N=1 SUSY respected by the D6-branes on intersections and thus plays the role of a supersymmetry breaking (SB) sector. This SB sector may be expected to break N=1 supersymmetry by gauge mediated supersymmetry breaking (GMSB)[see for example [40].] The NF brane, as the 4D models we examine are toroidal orientifolds of type IIA, may be expected, in the general case, to preserve $N = 0$ SUSY. In this case the one-loop contribution to gauginos may be depicted as in figure (3). As usual in gauge mediation, the mass of the gauginos may be of order

$$m_{\lambda_i} \approx c_i \frac{\alpha_i}{4\pi} M_s ,$$

(8.1)

where we have assumed that the supersymmetry breaking scale is at the string scale. As one-loop contributions to squarks and sleptons may vanish Higgs masses will get corrections from two loop effects of the order $\sim (\alpha/4\pi) M_s$, which limits the scale of electroweak symmetry breaking at small values as long as $M_s \leq 30$ TeV. Precise bounds on $M_s$ will be derived once the NF brane is available. The models constructed in this work that localize the spectrum of the N=1 SM use intersecting D6-branes. One can imagine a scenario of this sort in the presence of a mixture of orthogonal D3, D7 branes.

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13 as it has been noticed in the first reference of [15].
14 see relevant comments in the first reference of [15].
branes, as in the type I compactifications described in [41], where the gauge group of the N=1 SM may originate from the D3_aD3_a sector and the role of the NF brane in these models is played by the D7_aD7_a sector. In this case, supersymmetry breaking will be transmitted to the visible N=1 SM sector by the gauge interactions of the massive fields from the D3_aD7_b sector.

However, the presence of this NF brane is rather out of reach at present. As a consequence of this difficulty we have chosen to break N=1 SUSY from the start by using the presence of Fayet-Iliopoulos terms. We have found that at tree level it is not possible to give positive (mass)^2 to all sparticles. By calculating the two-loop corrections to these masses, in order to guarantee the presence of positive (mass)^2 for all sparticle masses, we may be able to avoid the existence of charge and colour breaking minima. We note that it is also possible that the existence of FI-terms triggers the generation of vev’s for the some scalars which restabilizes the vacuum and restores the local N=1 susy on the visible branes. These issues may be clarified when also two loop corrections may be calculated. We hope to return to these issues in the future.

8.1 Conclusions

In this work, we have constructed four dimensional string theory models, which in some cases, localize the MSSM spectrum when the NS B-field is not present across the six dimensional internal toroidal space. We have also constructed extensions of the MSSM with the same fermion and boson spectrum but where we have present extra - beyond the standard MSSM H_u, H_d - Higgs multiplets (HM’s). The number of HM’s depend

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15 we hope to refer to these issues in the future
16 The same observation was made for the ‘triangular’ quivers of [15] that possess the same N=1 properties as the present constructions
17 We thank Angel Uranga for reminding us of this possibility.
on the existence of the NS B-field $\#_{\text{Higgs}} = 1/(\beta_1 \beta_2)$ \textsuperscript{18}.

The most important property shared by these constructions is that all visible intersections, of the D-brane configurations of tables (1), (4), (7) share the same N=1 SUSY. Once the supersymmetry breaking sector is found - the explicit form of the NF brane - we will be able to show explicitly the breaking of N=1 supersymmetry by gauge mediation.

As we have described in section six we are able to move between models with different number of stacks. This is achieved using brane recombination (BR). BR describes the homology flow deformations of models II, III around the four stack model I. The presence of this mechanism is compatible with the deformation of the flat directions in the line of marginal stability of the D-term potential. To follow the present BR directions we have to deform along the lines of marginal stability. We note that even though there is a homology flow between the different models, they are not equivalent at the quantum level. The Yukawa interactions of models I, II, II are different. In addition, unless the full effective potential is calculated we cannot be sure that a specific BR direction is preferred against another one.

Baryon and lepton number are gauged symmetries in the present models as the corresponding gauge bosons receive a mass through the generalized Green-Schwarz mechanism. Thus proton is stable and only Dirac mass terms are allowed for neutrinos. Using Model II (five stack) against Model I - or Model III against Model II and/or Model I - for the explanation of charged lepton and neutrino mass mixing is certainly more advantageous as localizing each charged lepton and neutrino to different intersections allows for more freedom to generate the quark and lepton hierarchies. For example neutrino hierarchies may most easily accommodated in the five stack model II, as two generations of neutrinos get localized at the intersection cd; the third generation localized at the intersection ce. In this case two generations of neutrinos may be shown to be heavier than the third one, as the SuperKamiokande results suggest \textsuperscript{19}. It will be interesting to discuss extensively the flavour problem, the origin of families and fermion - quarks and lepton - masses as the latter appears to be consistent with the implementation of 'texture zeros' \cite{42} that points towards beyond the SM symmetries.

Finally, we examined the modification of the gauge coupling constants in the presence of a NS B-field. We obtained several interesting relations among the gauge cou-

\textsuperscript{18}where $\{\beta_1, \beta_2\} = \{1, 1/2\}$ the value of the B-field in the second and third $T^2$. The fractional values correspond to the presence of a non-trivial B-field.

\textsuperscript{19}By construction in the present model II, as in the five stack D6-models of \cite{5}, two neutrino and charged lepton species get localized at different intersection points.
pling constants at the string scale that may be used to examine the consequences for the
gauge coupling running unification (GCU) of the MSSM models I, II and III (and
their extensions in the presence of a non-zero NS B-field; its net effect the addition
of extra Higgs MSSM multiplets with identical quantum numbers). The GCU in this
case should appear at a relative low scale and below 30 TeV as a consequence of gauge
mediation and the existence of the SUSY breaking sector. This is to be contrasted
with the ‘observed’ unification of gauge couplings at the gauge theory version of the
MSSM at a GUT scale of $\approx 10^{16}$ GeV and the GCU in models of 4D weak coupling
N=1 heterotic orbifold compactifications [43] at $\approx 10^{17}$ GeV. The gauge coupling de-
pendence of the present models on the complex moduli values allows enough freedom
to possibly accommodate such unification choices. These issues will be examined in
detail elsewhere.

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where part of this work was done.

9 Appendix A

In this appendix further details for Model III may be found. These two tables are

| Brane | $\theta_1^a$ | $\theta_2^a$ | $\theta_3^a$ | approx. SUSY preserved |
|-------|------------|------------|------------|------------------------|
| a     | 0          | $\alpha_1 + \delta_a$ | $-\alpha_1$ | $r_1, r_4$            |
| b     | $\frac{\pi}{2} + \delta_b$ | 0          | $-\frac{\pi}{2}$ | $r_2, r_4$            |
| c     | $\frac{\pi}{2} + \delta_c$ | $-\frac{\pi}{2}$ | 0          | $r_3, r_4$            |
| d     | 0          | $\alpha_2 + \delta_d$ | $-\alpha_2$ | $r_1, r_4$            |
| e     | 0          | $\alpha_2 + \delta_e$ | $-\alpha_2$ | $r_1, r_4$            |
| f     | 0          | $\alpha_2 + \delta_f$ | $-\alpha_2$ | $r_1, r_4$            |

Table 17: Angle structure, coupling the complex structure moduli to open string modes
as Fayet-Iliopoulos terms for the six stack SUSY Model III.

related to the angle structure and the sparticle masses coming after the introduction
of the Fayet-Iliopoulos terms. We note that in table (18) we have assumed that $|\frac{\pi}{2} - \alpha_1| >$
| Sparticle | $(\theta^1, \theta^2, \theta^3)$       | Sector | $(mass)^2$  |
|----------|----------------------------------------|--------|-------------|
| $Q_L$    | $(-\frac{\pi}{2} - \delta_b, \alpha_1 + \delta_a, -\alpha_1 + \frac{\pi}{2})$ | $(ab)$ | $\frac{1}{2}(\delta_a - \delta_b)$ |
| $U_R$    | $(-\frac{\pi}{2} - \delta_c, \alpha_1 + \delta_a + \frac{\pi}{2}, -\alpha_1)$ | $(ac)$ | $\frac{1}{2}(\delta_a - \delta_c)$ |
| $D_R$    | $(\frac{\pi}{2} + \delta_c, \alpha_1 + \delta_a - \frac{\pi}{2}, -\alpha_1)$ | $(ac*)$| $-\frac{1}{2}(\delta_a + \delta_c)$ |
| $L^1$    | $(-\frac{\pi}{2} - \delta_b, \alpha_2 + \delta_d, -\alpha_2 + \frac{\pi}{2})$ | $(db)$ | $\frac{1}{2}(\delta_d - \delta_b)$ |
| $L^2$    | $(\frac{\pi}{2} + \delta_b, -\alpha_2 - \delta_c, -\frac{\pi}{2} + \alpha_2)$ | $(be)$ | $\frac{1}{2}(\delta_c - \delta_b)$ |
| $L^3$    | $(\frac{\pi}{2} + \delta_b, -\alpha_2 - \delta_f, -\frac{\pi}{2} + \alpha_2)$ | $(bf)$ | $\frac{1}{2}(\delta_f - \delta_b)$ |
| $N^1_{1R}$ | $(\frac{\pi}{2} + \delta_c, -\alpha_2 - \delta_d - \frac{\pi}{2}, \alpha_2)$ | $(cd)$ | $\frac{1}{2}(\delta_c - \delta_d)$ |
| $E^1_{1R}$ | $(\frac{\pi}{2} + \delta_c, \alpha_2 + \delta_d - \frac{\pi}{2}, -\alpha_2)$ | $(cd*)$| $\frac{1}{2}(\delta_c + \delta_d)$ |
| $N^2_{2R}$ | $(\frac{\pi}{2} + \delta_c, -\alpha_2 - \delta_e - \frac{\pi}{2}, \alpha_2)$ | $(ce)$ | $\frac{1}{2}(\delta_c - \delta_e)$ |
| $E^2_{2R}$ | $(\frac{\pi}{2} + \delta_c, \alpha_2 + \delta_e - \frac{\pi}{2}, -\alpha_2)$ | $(ce*)$| $\frac{1}{2}(\delta_c + \delta_e)$ |
| $N^3_{3R}$ | $(\frac{\pi}{2} + \delta_c, -\alpha_2 - \delta_f - \frac{\pi}{2}, \alpha_2)$ | $(cf)$ | $\frac{1}{2}(\delta_c - \delta_f)$ |
| $E^3_{3R}$ | $(\frac{\pi}{2} + \delta_c, \alpha_2 + \delta_f - \frac{\pi}{2}, -\alpha_2)$ | $(cf*)$| $\frac{1}{2}(\delta_c + \delta_f)$ |

Table 18: Sparticle masses from Fayet-Iliopoulos terms for the six stack quiver

0, $|\frac{\pi}{2} - \alpha_1| > |\delta_a - \delta_b|$, are satisfied. Also we assume that the relations $|\frac{\pi}{2} - \alpha_2| > 0$, $|\frac{\pi}{2} - \alpha_2| > |\delta_d - \delta_b|$, $|\frac{\pi}{2} - \alpha_2| > |\delta_e - \delta_b|$, $|\frac{\pi}{2} - \alpha_2| > |\delta_f - \delta_b|$ hold.

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