MSSM GUT String Vacua, Split Supersymmetry and Fluxes

E. G. Floratos\textsuperscript{(1,2)} and C. Kokorelis\textsuperscript{(1,2)}

\textsuperscript{1} Institute for Nuclear \& Particle Physics, N.C.S.R. Demokritos,GR-15310, Athens, Greece
\textsuperscript{2} Nuclear and Particle Physics Sector, Univ. of Athens, GR-15771 Athens, Greece

Abstract

We show that previous proposals to accommodate the MSSM with string theory N=0 non-supersymmetric compactifications coming from intersecting D6-branes may be made fully consistent with the cancellation of RR tadpoles. In this respect we present the first examples of non-supersymmetric string Pati-Salam model vacua with starting observable gauge group $SU(4)_c \times SU(2)_L \times SU(2)_R$ ($SU(2)$ from Sp(2)'s) that accommodate the spectrum of the 3 generation MSSM with a gauged baryon number with all extra exotics (either chiral or non-chiral) becoming massive and all MSSM Yukawas realized. These constructions include models with $\sin^2(\theta_W) = 3/8$ (SU(5) type) and can have 1, 2 or 4 pairs of higgsinos depending on the # of tilted tori. We work within four dimensional compactifications of IIA theory on toroidal orientifolds without (and with) fluxes. The MSSM spectrum (together with right handed neutrinos) is realized in the intersections of the visible sector that may contain D6-branes whose intersections share the same N=1 supersymmetry. The N=1 supersymmetry of the visible sector is broken by an extra supersymmetry messenger breaking sector that preserves a different N'=1 susy, exhibiting the first examples of stringy gauge mediated models. Due to the high scale of the models, these models are also the first realistic examples of carriers of stringy split supersymmetry exhibiting universal slepton/squark masses, massive string scale gauginos, unification of $SU(3)$, $SU(2)$ gauge couplings at $2.04 \times 10^{16}$ GeV, a stable proton and the appearance of a landscape split SM with chiral fermions and only Higgsinos below the scale of susy breaking; the LSP neutralino candidate could also be only Higgsino or Higgsino-Wino mixture. We also add RR, NS and metric fluxes as every intersecting D-brane model without fluxes can be accommodated in the presence of fluxes. The addition of metric fluxes in the toroidal lattice also stabilizes the expected real parts of all in AdS closed string moduli (modulo D-term effects), leaving unfixed only the imaginary parts of Kähler moduli.
1 Introduction

Maybe the most serious problem of string theory nowadays is to derive a model of particle physics which will be as close as possible to the Standard Model at low energies and which not only manage to fix all its free parameters, its moduli, but it will also get rid of all its extra exotics (chiral or non-chiral) – by making them massive through appropriate Yukawa couplings or some other mechanism – which always have been a problem in model building attempts. The obvious next step to such an goal is to calculate specific phenomenological quantities that could make some definite predictions for present and future experiments. In this respect recent model building attempts - without the presence of background fluxes - have been focused in the construction of N=1 supersymmetric [1, 2, 8, 63] and non-supersymmetric models [14, 15, 16, 17, 26, 54, 55] and the use of intersecting branes [See also [9], [10] — for some semirealistic attempts in deriving the MSSM from another direction.]. These models make use of the fact that chiral fermions live in the intersections of branes that intersect at angles or in other cases they make use of the T-dual formulation of models with D9-branes and magnetic fluxes [11, 12, 13]. Some comments are in order. In all recent string N=1 supersymmetric models [1, 2, 8, 21, 22, 57, 58, 59]- where NSNS and RR tadpoles cancel - the localization of MSSM (we mean the usual multiplet context MSSM with one or more Higgsino $H_u$, $H_d$ multiplets and also right handed neutrinos) is accompanied by an unwanted problematic large number of non-chiral or chiral open string exotics that survive massless to low energies. These states could be coming from either the adjoint gauge multiplet (chiral ones) or from sectors (non-chiral ones) formed between in brane intersections where the participating branes are parallel in some compact direction. It is also possible that adjoint matter is formed from open string states that are accommodated in the intersection of a brane with its orbifold images. This is the case of $Z_N$ or $Z_N \times Z_M$ orientifolds of IIA compactifications with intersecting D6-branes. See for example [8], [50].

Following these developments during the era of development of intersecting brane models the initial expectation was that possibly IIB string backgrounds in the presence of NS and RR fields [38, 39] will allow for more flexibility into the spectrum, such that the extra exotics will disappear from the spectrum or alternatively the present formulation of N=1 models with intersecting D6-branes will manage somehow to find a vacuum that have less [2] or even no massless exotics at the end. Needless to say that in model building

\[1\] We also note that N=1 model building based on the by now old heterotic string approach suffers from problems like proton decay and unfixed moduli parameters. See for example [34].
with or without fluxes surviving massless exotics to low energies are always present either as a part of a chiral or a non-chiral set of new exotic particles [for some examples see [40, 22]].

On the other hand non-supersymmetric models (NSM) from intersecting branes (without fluxes) [14, 15, 16, 17, 26] in toroidal orientifold compactifications [3] of type IIA theory [see also [54, 55] for the generation of the same SM configurations using D5-branes in different IIB backgrounds], have some successes as it become possible to derive for the first time vacua which have only the SM at low energies and with no extra chiral exotics - and with proton stability - without the presence of any additional chiral fermions at low energies. Hence the original four stack non-susy SM [14] string vacua - that have no supersymmetry preserved at any intersection - have been extended 2 to five and six stack SM string vacua that have one and two intersections preserving a supersymmetry [16, 17] respectively 3. [For one way to get rid in these vacua of non-chiral fermions in any representation see e.g. [61]; one could also try to get rid of only the adjoint matter but using fractional D-branes in 4D IIB chiral compactifications [63]]. These models follow a bottom-up approach as they embed the SM local configuration to a string compactification and thus they are different to the top-bottom MSSM embedding approach of heterotic string compactifications. We also mention at this point the construction of non-susy Pati-Salam vacua with only the SM at low energy [15], where all the accompanying beyond the SM extra chiral exotics become massive through the use of non-renormalizable mass couplings 4. In this work, we will see an identical effect - which is consistent with a high scale - and makes all exotics massive and leaves only the MSSM N=1 context at low energies [using the same orientifold backgrounds [3] without fluxes].

In parallel with the development of non-susy models with only the SM at low energy, a different direction was initiated in [24], and further explored in [25], which localized the MSSM on a 4D four stack toroidal orientifold IIA vacuum (no fluxes present). In this

---

2By spreading the SM particle context to different intersections.

3These models predict the existence of the chiral spectrum of the SM in addition to only one type of supersymmetric particles, namely the susy partners of $\nu_R$’s the sneutrinos.

4One important constraint that is derived in these models - proton is stable - is that the masses of the extra chiral exotics are greater than the electroweak scale only if the string scale is low and below 1.2 TeV. However the latter is unlikely to happen- the string scale is close to the Planck scale - as on these models the intersecting D6-branes wrap the whole of the toroidal orientifold space and there are no transverse dimensions to the D6-branes that could become large such that the string scale could be as lowered to below 1.2 GeV. Some other possibilities which make these models consistent with a high scale will be explored elsewhere.
case, even though an anomaly free configuration was found that locally corresponded to the MSSM N=1 multiplet spectrum nevertheless RR tadpoles could not be satisfied. In [26] we generalized the considerations of [24] to the maximal five and six stack locally supersymmetric MSSM without fluxes and also considered the introduction of tilted tori on four, five and six stack MSSM local configurations 5. Furthermore in [24] it was argued that since the MSSM local configuration was anomaly free in order to satisfy RR tadpoles either an extra anomaly free sector was needed or NS-NS background fluxes that however should add no net chiral content. Additionally, in [26] we argued 6 that the extra sector should play the role of the supersymmetry breaking sector of gauge mediation [35] and the states of the extra RR cancelling sector should be vector-like. We should mention that the same local configuration describing the MSSM [24], [25] has been also used for the studies of unification of gauge couplings [36] and calculation of soft terms without fluxes [37] in intersecting brane models.

The purpose of the present paper is twofold. Initially we present a new method of cancelling RR tadpoles in N=0 (non-supersymmetric) toroidal orientifold models by adding a vector-like \( N' = 1 \) supersymmetric sector to N=1 local MSSM configurations, leaving as the only net chiral context of the theory the multiplet spectrum of the N=1 Standard Model. The vector-like sector respects a different N=1 supersymmetry than the one respected by the sector that localizes the MSSM. Hence overall SUSY is broken and the models are non-supersymmetric. Furthermore, since we show that all extra beyond the MSSM matter either chiral or non-chiral becomes massive this is the first appearance of a realistic string model that finds the N=1 MSSM particle content inside a model where SUSY is already broken 7[overall the model is non-supersymmetric]. At present there is no N=1 model that can localize the MSSM matter context and be able to get rid of its massless exotics chiral and/or non-chiral ones.

We apply this method to toroidal orientifolds by finding the most general solution to RR tadpoles that localize the MSSM multiplet spectrum inside a Pati-Salam \( SU(4)_c \times Sp(2)_L \times Sp(2)_R \) construction. *What we call the MSSM is the usual MSSM chiral multiplet spectrum with right handed neutrinos - with gauged baryon number and hence proton stability - and either one (1), two (2) or four (4) pairs of Higgsinos \( H_u, H_d \); the latter choices depending*

5By brane recombination the five, six stack local MSSM models of [26] flow to the four stack MSSM local configurations of [24, 25].

6The 5- and 6-stack MSSM local configurations of [26] flow - under brane recombination of the U(1)’s - to the 4-stack models of [24, 25].

7At present there is no N=1 supersymmetric models which manage in some way to get rid of its extra exotics chiral and/or non-chiral ones that accompany the MSSM spectrum.
on the number of tilted tori. Furthermore we will show that all the extra beyond the MSSM matter - all the states of the extra vector-like messenger sector will get masses from Yukawa couplings. Adjoint matter will get masses from the introduction of Scherk-Schwarz breaking. On the other hand - following recent important developments on moduli stabilization either in N=1 supersymmetric AdS vacua with NS/R fluxes [41] and N=1 supersymmetric vacua with NS/R and metric fluxes [44] recently, where all the real parts of the moduli could get fixed and only some combinations of axions are being left unfixed - we examine the issue of moduli stabilization by adding metric fluxes in the present models. We recall that in non-susy compactifications [14, 15, 16, 17] from toroidal orientifolds only complex moduli could get fixed through the use of supersymmetry conditions on intersections; see for example [15]. Moreover the presented Pati-Salam models - when Scherk-Schwarz breaking is included - have all the necessary ingredients to be the first realistic split susy [47] models from string theory [48], [49], [50], [51] as the models exhibit partial unification of the strong SU(3) and weak SU(2) force gauge couplings at the famous value \(2.04 \times 10^{16} \text{ GeV}\) as at the string scale only the MSSM particle content remains massless.

The structure of the present paper is as follows. In section 2 we describe the local MSSM configurations as described in [25], [26] with zero and non-zero NS B-field respectively which did not satisfy at the time RR tadpoles. In section 3 we present a class of RR tadpole solutions that describes the Pati-Salam embedding \(SU(4) \times Sp(2)_L \times Sp(2)_R\) type of models where all the additional exotics become massive and where the extra messenger - beyond the MSSM gauge group - is just two broken U(1)’s. These RR tadpole solutions embed consistently the MSSM spectrum in the string compactification of four dimensional toroidal orientifolds [3]. We also show how all the chiral and non-chiral beyond the MSSM exotics become massive. In section 4, the breaking of these models to its left-right counterpart symmetric models and the breaking to the MSSM is described also describing the way all the chiral and non-chiral beyond the MSSM exotics become massive. We also use complex structure moduli in Fayet-Iliopoulos terms to show that all sparticles/squarks get massive from the breaking of \(N = 1, N' = 1\) supersymmetries. In section 5, the relation of the present constructions to split supersymmetry models and the satisfaction of all criteria for its application to strings set out in [48, 49] is described. We also present a) constraints on the complex structure moduli derived from the requirement that the Pati-Salam models contain an SU(5) type of \(\sin^2(\theta)\) at \(M_s\), as well consequences

\(^8\)Even though in a revised version of [26] a new RR tadpole canceling model may be studied.
for dark matter candidates. In section 6 we describe the addition of R, NS and metric fluxes as described in [43], [44] and the fixing of most of the moduli in the Ads vacuum of the present models. We present our final comments and conclusions in section 7.

2 MSSM and \( N=1 \) Supersymmetry in \( N=0 \)

Toroidal Orientifolds from Intersecting D6-branes

2.1 Preliminaries

In string models from intersecting branes chiral fermions get localized in the intersection between branes. Let us describe in some detail one such construction that involves four dimensional intersecting brane models coming from toroidal orientifolds of type IIA [3] theory. Consider IIA theory with D9-branes with fluxes compactified on a six dimensional torus, which is acted upon the worldsheet parity \( \Omega \). Performing a T-duality on the \( x^4, x^6, x^8 \) compact directions the worldsheet parity action symmetry is mapped to \( \Omega R \), where the action \( R : z \rightarrow \bar{z} \) and the D9-branes with magnetic flux \( F \) get mapped to D6-branes intersecting at angles \( \theta = \tan^{-1}\frac{mRa}{nRa} \), where \((n_a, m_a)\) the number of times the D6-brane \( a \) is wrapping a one cycle on the i-th two torus. The \( \Omega R \) action introduces an \( \Omega R \) image of the D6-brane that is wrapping the cycle \((n, -m)\). The introduction of discrete B-field [23] tilts the torus and maps the wrapping numbers according to the rule \((n, m) \rightarrow (n, m + (n/2))\). For this reason we introduce K-stacks of \( N_a, a = 1, \ldots K \) of D6-branes wrapped on \((n_a^i, m_a^i)\) cycle on the i-th two torus. The net number of chiral fermions localized in the intersection between the branes \( a, b \) and \( a, b^* \) is then given respectively \(^9\)

\[
I_{ab} = [\Pi_a] \cdot [\Pi_b] = \prod_{i=1}^{3} (n_a^i m_b^i - n_b^i m_a^i),
\]

\[
I_{ab^*} = [\Pi_a] \cdot [\Pi_b^*] = -\prod_{i=1}^{3} (n_a^i m_b^i + n_b^i m_a^i),
\]

for open string states starting from brane \( a \) and ending on brane \( b \). Such a state belongs to a bifundamental \((N_a, \bar{N}_b)\) of the gauge group, e.g a left handed quark. Chirality is

\(^9\)The intersection number is the product of the homology classes of the D6\(_a\)-branes \( \Pi_a \) and their orientifold images \( \Pi_a^* \), where we define the homology classes of the three cycles \([a_i], [b_i] \) of the i-th-torus as

\[
[\Pi_a] = \prod_{i=1}^{3} (n_a^i [a_i] + m_a^i [b_i]), \quad [\Pi_a^*] = \prod_{i=1}^{3} (n_a^i [a_i] - m_a^i [b_i])
\]
defined by choosing an intersection sign, negative sign implies right handed particles. The spectrum of toroidal orientifolds could also accommodate fermions in the $S + A$ representations \(^{10}\) of the gauge group. A gauge group $U(N)$ appears from sectors that involve open strings that start and end on the same stack of $N$ coincident branes. If a brane is its own orientifold image then an $Sp(2)$ gauge group can appear.

| Sector | Multiplicity | Representation |
|--------|--------------|----------------|
| $aa$   | 3            | $U(N_a)$ vector multiplet |
| $a(b)$ | $I_{ab}$     | Adj. chiral multiplets |
| $a(b')$| $I_{ab'}$    | (\(\square, \Box\)) fermions |
| $a(a')$| $4m_a^1m_a^2m_a^3(n_1^a n_2^a n_3^a - 1)$ | (\(\square\)) fermions |
|        | $4m_a^1m_a^2m_a^3(n_1^a n_2^a n_3^a + 1)$ | (\(\Box\)) fermions |

Table 1: General spectrum on D6-branes at generic angles (namely, not parallel to any O6-plane in all three tori) in toroidal orientifolds. The models contain additional non-chiral pieces in the $aa'$, $ab$, $ab'$ sectors with zero intersection, if the relevant branes have a parallel direction. The latter could become massive in principle when SS breaking is included.

The above rules of determining the gauge group and chiral spectrum are not enough when deriving an extension of the Standard Model from a string compactification. Another consistency condition that may be satisfied by the D6-branes wrapping the compact space are the RR tadpole cancellation conditions for compactifications which is the cancellation of the RR charge in homology

$$
\sum_a N_a[\Pi_a] + \sum_a N_a[\Pi_{a*}] = 32[\Pi_{O6}] \tag{2.3}
$$

for the D-branes and their orientifold images and their O-planes. For four dimensional compactifications of type IIA on toroidal orientifolds with D6-branes intersecting at angles they are given by

$$
\sum_a N_a n_a^1 n_a^2 n_a^3 = 16,
$$

$$
\sum_a N_a n_a^1 m_a^2 m_a^3 = 0.
$$

\(^{10}\)The sector denoted as $aa'$ in table (1) also involves intersections of the $a$-brane with the orientifold O6 plane.
\[ \sum_{a} N_a m_a^1 n_a^2 m_a^3 = 0, \]
\[ \sum_{a} N_a m_a^1 m_a^2 n_a^3 = 0 \] (2.4)

Another consistency condition which affects the chiral spectrum is the existence of some amount of supersymmetries carried out by the D-branes.

The notation for supersymmetries shared by the intersections and the orientifold planes is as follows. A pair of intersecting D6-branes that wraps across the \( T^6 \) factorizable tori and having angles \( \theta_i \) across the \( T^6 \), \( i=1,2,3 \) preserves \( N=1 \) supersymmetry if

\[ \pm \theta_1 \pm \theta_2 \pm \theta_3 = 0 \] (2.5)

for some choice of the signs, where the angles \( \theta_i \) are the relative angles between a pair of branes across the three 2-tori. We distinguish the different susys shared by the branes and the orientifold planes by the sign choices \(^{11}\)

\[ r_1 = \left( \frac{1}{2} \right) (- + + -) \]
\[ r_2 = \left( \frac{1}{2} \right) (+ - + -) \]
\[ r_3 = \left( \frac{1}{2} \right) (+ + - -) \]
\[ r_4 = \left( \frac{1}{2} \right) (- - - -) \] (2.6)

which correspond respectively to the angles choices

\[ - \theta_1 + \theta_2 + \theta_3 = 0, \]
\[ + \theta_1 - \theta_2 + \theta_3 = 0, \]
\[ + \theta_1 + \theta_2 - \theta_3 = 0, \]
\[ + \theta_1 + \theta_2 + \theta_3 = 0. \] (2.7)

The masses of the lightest scalar states appearing in the NS sector in an intersection are

\[ m^2 = \frac{M^2}{2} (- \theta_1 + \theta_2 + \theta_3), \quad m^2 = \frac{M^2}{2} ( \theta_1 - \theta_2 + \theta_3), \]
\[ m^2 = \frac{M^2}{2} ( \theta_1 + \theta_2 - \theta_3), \quad m^2 = \frac{M^2}{2} (1 - \frac{1}{2} ( \theta_1 + \theta_2 + \theta_3)). \] (2.8)

At present there are three ways to embed the Standard Model (or the MSSM) gauge group and spectrum into a unitary or symplectic configuration.

\(^{11}\)We follow the relevant discussion in [19].
They are classified according to the "observable" gauge group that appears at the string scale and they are grouped into three classes of models, namely the first one \((2.9), (2.10)\); the second class \((2.11), (2.12)\) and the third class \((2.13)\) which read

- \(U(3)_c \times U(2)_L \times U(1)^n; \ n = 2, 3, 4\) \hspace{1cm} (2.9)
- \(U(4)_c \times U(2)_L \times U(2)_R \times U(1)^m; \ m \neq 0\), \hspace{1cm} (2.10)

and

- \(U(3)_c \times Sp(2) \times U(1) \times U(1)\), \hspace{1cm} (2.11)
- \(U(4)_c \times Sp(2) \times Sp(2)\) \hspace{1cm} (2.12)

and

- \(U(4)_c \times Sp(2)_L \times U(2)_R, \ or \ U(4)_c \times U(2)_L \times Sp(2)_R,\) \hspace{1cm} (2.13)

The first class of three generation models, namely \((2.9), (2.10)\), uses only bifundamental fields for the chiral field description of the open string spectrum and treats the weak group \(SU(2)_L\) as the one coming from the decomposition \(U(2) \supset SU(2)_L \times U(1)\).

Such constructions have been discussed in the context of non-supersymmetric constructions with the chiral spectrum of the SM at low energy in \([14]\) \((n=2)\) for four-stack models and in \([16, 17]\) \((n = 3, 4)\) respectively for related constructions that represent the maximal SM embedding in the five- and six-stack extensions of \([14]\) respectively. On the other hand the Pati-Salam \([53]\) embedding \((2.10)\) was discussed in the construction of non-supersymmetric GUTS in \([15]\). The SM embedding \([14, 16, 17, 15]\) follows a bottom-up approach in the sense that they embed anomaly free configurations that describe the SM chiral spectrum \([14, 16, 17, 15]\) into overall \(N=0\) non-supersymmetric string models \(^{12}\) with complete cancellation of RR tadpoles. Hence these are genuine string \(^{13}\) models \(^{14}\) embedded in four dimensional toroidal orientifolds of type IIA theory with intersecting

\(^{12}\)N=1 string models will not be mentioned explicitly in the following considerations but we will comment where necessary. In general these models are not very appealing at the moment as they are full of extra beyond the MSSM exotics which survive massless to low energies. At present there is no know way to get rid of these exotics in any N=1 construction that involves D-branes intersecting or not.

\(^{13}\)See \([56]\) for RG gauge group studies of non-string models in a D-brane inspired context.

\(^{14}\)An embedding that localizes the chiral SM spectrum have been also found in compactifications of \(Z_3\) \([18]\) or \(Z_3 \times Z_3\) \([50]\) orientifolds of type IIA using intersecting D6-branes. In this case the SM is made of mixtures from bifundamental and antisymmetric representations.
D6-branes [3]. These models have a stable proton, as baryon number is a gauged symmetry and the corresponding global symmetry survives to low energy as the corresponding gauge boson becomes massive, while in some of these models neutrinos get masses from quark condensates (QCs) [14], [15] (for recent discussions on QCs see [68]).

In the second class of models, namely the constructions given by (2.11), (2.12), where three generation non supersymmetric constructions are realized, the weak $SU(2)_L$ gauge symmetry is delivered within an $Sp(2) \equiv SU(2)$ factor. In this case, also the left-right construction utilizes the $SU(2)_R$ factor as an $Sp(2)_R$.

Let us note that while the non-supersymmetric (non-susy) models of the first class (2.9), (2.10), are useful in deriving the Standard Model spectrum, non-susy models of the second class (2.11), (2.12), are useful to build the MSSM intersection numbers that are associated with the N=1 chiral multiplet spectrum of the MSSM [24, 25, 26] in the context of a non-supersymmetric string construction. Models of those types have been discussed in [24, 25, 26] and at present they do not satisfy RR tadpoles. The purpose of his work is to bridge this literature gap and to show that the embedding of these models in a string construction is possible; so that model building will be elevated from the local [24, 25, 26] to a string level.

The third class of constructions seen in eqn. (2.13), do not exist in the context of non-supersymmetric string constructions from intersecting branes and the embedding of the three generation Pati-Salam non-susy constructions - into a toroidal orientifold compactification of type IIA with intersecting D6-branes [3] - which localizes the MSSM chiral multiplet spectrum may be described in [26].

We mention that N=1 constructions from intersecting D6-branes where part of the gauge group is as in eqn. (2.12) have been considered in [65] in the presence of only RR and NS fluxes. Unfortunately, at present there is no either compactification or known way to get rid from the rich varieties of extra chiral fields that appear in all N=1 4D supersymmetric models from intersecting D6-branes with or without general fluxes. Also other MSSM-like N=1 constructions with the gauge group as in (2.12) have been considered [57] in the context of Gepner-type IIB orientifold compactifications.

Explicitly, the first attempt $^{15}$ to localize the MSSM within string theory D-brane non-supersymmetric models has been initiated in [24, 25] where only the intersection numbers - in the absence of a discrete NS B-flux - that localize the spectrum of the MSSM were provided (in a four stack construction) as they are seen in table (2). These wrappings $^{15}$

\[ 15 \text{In N}=1 \text{ models the MSSM chiral spectrum appears as part of the spectrum but such attempts produce exotics surviving massless to low energies.} \]
Table 2: D6-brane wrapping numbers [24, 25] with orthogonal tori that gives rise to a three generation N=1 MSSM spectrum via the Pati-Salam $SU(4)_c \times SU(2)_L \times SU(2)_R$ gauge group. The general solution to the intersection numbers is parametrized by the a parameter $\rho = 1, 1/3$.

| $N_i$ | $(n_1^i, m_1^i)$ | $(n_2^i, m_2^i)$ | $(n_3^i, m_3^i)$ |
|-------|------------------|------------------|------------------|
| $N_a = 3$ | $(1, 0)$ | $(1/\rho, 3\rho)$ | $(1/\rho, -3\rho)$ |
| $N_b = 1$ | $(0, 1)$ | $(1, 0)$ | $(0, -1)$ |
| $N_c = 1$ | $(0, 1)$ | $(0, -1)$ | $(1, 0)$ |
| $N_a = 1$ | $(1, 0)$ | $(1/\rho, 3\rho)$ | $(1/\rho, -3\rho)$ |

Table 3: General D6-brane wrapping numbers [26] that gives rise to a three generation N=1 MSSM spectrum via the Pati-Salam $SU(4)_c \times SU(2)_L \times SU(2)_R$ gauge group when $\epsilon = \tilde{\epsilon} = +1$, $\epsilon = -\tilde{\epsilon} = +1$. The general solutions to the intersection numbers are parametrized by two phases $\epsilon = \pm 1$, $\tilde{\epsilon} = \pm 1$, the NS background on the second and third tori respectively $\beta_1$, $\beta_2$ and a parameter $\rho = 1, 1/3$.

| $N_i$ | $(n_1^i, m_1^i)$ | $(n_2^i, m_2^i)$ | $(n_3^i, m_3^i)$ |
|-------|------------------|------------------|------------------|
| $N_a = 3$ | $(1, 0)$ | $(1/\rho, 3\rho \beta_1)$ | $(1/\rho, -3\rho \tilde{\beta}_2)$ |
| $N_b = 1$ | $(0, \epsilon\tilde{\epsilon})$ | $(1/\beta_1, 0)$ | $(0, -\tilde{\epsilon})$ |
| $N_c = 1$ | $(0, \epsilon)$ | $(0, -\epsilon)$ | $(\tilde{\epsilon}/\beta_2, 0)$ |
| $N_a = 1$ | $(1, 0)$ | $(1/\rho, 3\rho \epsilon \beta_1)$ | $(1/\rho, -3\rho \tilde{\epsilon} \beta_2)$ |

give rise to an $U(3) \times Sp(2) \times Sp(2) \times U(1)$ gauge group from the intersecting D6-brane stacks a, b, c, d respectively. Also the choice of parallel stacks in table (2) suggests the gauge symmetry enhancement $U(3)_a \times U(1)_d \rightarrow U(4)$ by turning on suitable vevs for the adjoint multiplets of the model. However this choice do not satisfy RR tadpoles in the simplest of the string constructions that is toroidal orientifolds $T^6/\Omega R$ of IIA [3]. The four stack intersection number configurations of table (2) have been generalized in the presence of the discrete NS B-flux in [26] hence suggesting the maximal five and six stack
extensions\textsuperscript{16} of table (2), again in the absence of any solutions to RR tadpole cancellation conditions; they are shown in table (3). The latter solutions depend on several parameters, the NS-B-fields $\beta_1, \beta_2$ of the second and third tori; the phases $\epsilon, \bar{\epsilon}$ that can take the values $\pm 1$ and the parameter $\rho = 1, 1/3$.

Notice that the solutions of table (3) for the special values of the parameters

$$\beta_1 = \beta_2 = 1, \quad \epsilon = \bar{\epsilon} = 1, \quad \rho = \frac{1}{3}$$

(2.14)

with all the tori orthogonal, as they appear in the RR tadpole solution of the Pati-Salam models of table (7), reproduce in the top part the wrappings numbers of the MSSM Pati-Salam embedding of \cite{24, 25} and also for the values of the parameters (2.14) the observable Pati-Salam sector embedding of the MSSM intersection numbers of the IIB MSSM-like fluxed embedded model of \cite{40}.

Explicitly, the starting gauge group for the parameter values (2.14), is an

$$U(4)_c \times Sp(2)_L \times Sp(2)_R$$

(2.15)

where the first stack of D6-branes gives rise to a $U(4)_c$ factor, the second stack to an $Sp(2)_L \equiv SU(2)_L$ gauge group and the third stack to an $Sp(2)_R \equiv SU(2)_R$ gauge group, since the branes b, c are invariant under the $\Omega R$ action. The spectrum associated with the parameters (2.14) is seen in table (4). By adjoint splitting of the $U(4)$ factor we get an $U(3) \times U(1)_d$ - which could be identified with $U(1)_{B-L}$ - thus recovering a left-right extension $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. Subsequently by also considering arbitrary positions and Wilson lines for the c D6-brane (see also \cite{25}) $SU(2)_R \to U(1)_c$, the original Pati-Salam model of table (4) gives rise to the MSSM spectrum of table (5)(as the spectrum exhibits N=1 supersymmetry as we will comment below), where only the

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|}
\hline
Matter & $SU(4) \times Sp(2)_L \times Sp(2)_R$ & $I_{ij}$'s & $Q_a$ & SY SY \\
\hline
$F_L$ & 3(4, 2, 1; 1, 1) & (ab) & 1 & $r_4$ \\
\hline
$\bar{F}_R$ & 3(4, 1, 2; 1, 1) & (ac) & $-1$ & $r_4$ \\
\hline
$\bar{h}$ & $\frac{1}{2} \epsilon_{abc}$ & (bc) & 0 & $r_1, r_4$ \\
\hline
\end{tabular}
\caption{Chiral spectrum of a 3-stack D6-brane Pati-Salam extension of the MSSM with three generations of chiral multiplets. We have chosen $\rho = 1/3, \epsilon = \bar{\epsilon} = 1, \beta_1 = \beta_2 = 1$ in table (2).}
\end{table}

\textsuperscript{16}Represent deformations of the models of [24, 25].
Table 5: Chiral spectrum of the four stack D6-brane N=1 Supersymmetric Standard Model with its $U(1)$ charges. The general form of the spectrum for non-trivial tilding along the second and third torus has been reproduced from [26]. For $\beta_1 = \beta_2 = 1$, it gives the local MSSM-like models of [25].

\[
U(1)^Y = \frac{1}{6} Q_a - \frac{1}{2} Q_c - \frac{1}{2} Q_d
\]

survives the Green-Schwarz mechanism massless to low energies.

Let us now describe some properties of tables (2), (3), (4), (5).

- Tables (2), (3): The intersection numbers are $I_{ab} = 3$, $I_{bc} = 3$, give rise to the usual multiplets of the Pati-Salam $G_{422}$ structure that accommodates three generations quarks and leptons into the following representations

\[
F_L = (4, 2, 1) = q(3, \frac{2}{6}, \frac{1}{2}) + l(1, \bar{2}, -\frac{1}{2}) \equiv (u, d, l)
\]

\[
\bar{F}_R = (\bar{4}, 1, 2) = u^c(3, 1, -\frac{2}{3}) + d^c(3, 1, \frac{1}{3}) + e^c(1, 1, 1) + N^c(1, 1, 0) \equiv (u^c, d^c, l^c)
\]

The quantum numbers on the right hand side of (2.17) correspond to the decomposition of $SU(4)_C \times SU(2)_L \times SU(2)_R$ under $SU(3)_C \times SU(2)_L \times U(1)_Y$; $l = (\nu, e)$ the lepton doublet of the SM and $l^c = (N^c, e^c)$ the right handed leptons.

- Table (3): Note that the brane b-wrappings give rise to either the group $Sp(2)_b$, when the brane b is its own orientifold image as happens in the case $\epsilon = \tilde{\epsilon} = +1$, $\epsilon = -\tilde{\epsilon} = +1$. 

| 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    |
| $H_d$         | $\frac{1}{\beta_1 \beta_2}(1, 2)$ | $(cb^*)$ | 0     | 1     | 0     | $-1/2$ |
| $H_u$         | $\frac{1}{\beta_1 \beta_2}(1, 2)$ | $(cb)$    | 0     | $-1$  | 0     | $1/2$ |
• Tables (3), (4): The spectrum exhibits N=1 supersymmetry as long as the following susy condition holds

\[ \beta_1 \chi_2 = \beta_2 \chi_3. \tag{2.18} \]

Also notice that the branes separately, share some susy with the orientifold plane as seen in table (4).

• Tables (3), (4), (5): The Higgs sector arises from open strings stretching between the branes b, c. The branes b, c are parallel in the first tori giving rise to a non-chiral sector with N=2 supersymmetry. Note that the intersection number \( I_{bc} \) vanishes thus the net chirality in this sector is zero. Hence we recover from this sector \( 1/(\beta_1 \beta_2) \) chiral multiplets in the (2, 2) representation of \( SU(2)_L \times SU(2)_R \). When breaking the \( SU(2)_R \rightarrow U(1)_c \) by adjoint breaking or else, one (2, 2) multiplet will split into two N=1 multiplets (2, 1), and (2, -1) under \( SU(2)_b \times U(1)_c \) which are identified as the MSSM Higgs multiplets \( H_u, H_d \). Hence after the breaking of the left-right symmetry, we will get \( 1/(\beta_1 \beta_2) \) \( H_u \)'s and an equal number of \( H_d \) N=1 multiplets, which gives us three (3) versions of the MSSM with one (1), two (2) and four (4) pairs of N=1 Higgs multiplets \( H_u, H_d \).

• Global symmetries of the SM could be identified with some of the U(1)’s appearing in table (5). Hence the baryon number may be identified as \( 3B = Q_a \), lepton number \( L = Q_d \). As a result of the couplings of the U(1)’s to RR fields since baryon number is a gauged symmetry and the corresponding gauge boson is getting massive; proton is stable perturbatively. Issues on proton stability in intersecting branes can be found in [28, 29, 30, 31, 32].

Summarizing the number of Higgs multiplets which depends on the number of tilted tori reads

\[
\frac{1}{\beta_1 \beta_2} = \begin{cases} 
1, & \beta_1 = \beta_2 = 1 \\
2, & (\beta_1, \beta_2) = (1, 1/2) \text{ or } (\beta_1, \beta_2) = (1/2, 1) \\
4, & \beta_1 = \beta_2 = 1/2.
\end{cases} \tag{2.19}
\]

\(^{17}\)where \( \chi_i = R_i^b/R_i^1 \) the complex structure moduli in the i-th torus.
3 Embedding the MSSM in a Pati-Salam String Compactification

The local construction [25], [26] of tables (2), (3), (4), (5) respectively, represent the observable Pati-Salam (PS) extension of MSSM that in this section may be embedded globally in a 4D toroidal IIA orientifold string compactification since RR tadpoles will be shown to be cancelled appropriately. At this point we introduce a new way to cancel RR tadpoles.

- The observable sector $\mathcal{O}$ made of $a, b, c, d$ D6-branes seen in table (4) is made from intersections that localize the MSSM chiral multiplets and where all intersections respects the same N=1 supersymmetry. We now cancel the RR tadpoles (2.4) by the addition of an extra sector that respects a $N' = 1$ supersymmetry and is made from the branes $h_1, h_2, h_3, h_4, h_5$ and which it is different than the N=1 respected by the ”observable” Pati-Salam MSSM sector.

- The wrapping numbers of the extra sector can be seen in table (6). Observe that a number of extra branes $h_1, h_2, ..., h_5$ have to be added to cancel RR tadpoles. From these branes, only the $h_3, h_4$ one’s have a non-zero intersection number with the observable sector branes on all three tori and thus have extra net chiral exotic fermions localized in their intersection. Also we notice that since all intersections preserve either a $N = 1$ or a $N' = 1$ supersymmetry, each fermion will also accompanied by its massless boson part of the chiral multiplet. Non-chiral matter from sectors where the branes are parallel in some tori may be made massive as we discuss shortly. These RR tadpole solutions generate a PS embedding of the MSSM which may be studied in detail in the next section.

3.1 Three generation Pati-Salam models with a Messenger Sector

We will now derive unique PS string vacua where all the extra exotics become massive. In the wrapping numbers of table (3) we set $\epsilon = \tilde{\epsilon}$. The solution to the RR tadpoles may be seen in table (13) where the $\rho$ parameter takes the value $1, 1/3$.

Substituting on table (13) the value $\rho = 1/3$, we obtain in table (7) the much wanted solutions to the RR tadpoles for the local construction of [25]. The chiral spectrum can be seen in table (8). The observable gauge group under which chiral fermions exist is

$$U(4) \times U(2)_b \times U(2)_c \times U(1)_{h_3} \times U(1)_{h_4}$$ (3.1)
Table 6: Solution to the RR tadpoles for toroidal orientifold models. The N=1 MSSM chiral spectrum arises as part of Pati-Salam models in the top part of the table from intersections between a, b, c, d branes. Messenger multiplet states respecting a $N' = 1$ supersymmetry arise from the intersections of the a, d branes with the $h_3, h_4$ branes. We have set $\rho = 1/3, \epsilon = \tilde{\epsilon}$ in table (3).

\[
SU(4) \times U(1)_a \times SU(2)_b \times SU(2)_c \times U(1)_{h_3} \times U(1)_{h_4} ,
\]

where the SU(2)'s come from Sp(2)'s; the chiral spectrum can be seen in table (8). and the gauge group (G.G) transformations and charges are under the (3.1) G.G. In the top part of table (8) we find the usual Pati-Salam part that embeds the MSSM chiral spectrum while the extra sector that cancels RR tadpoles is seen in the bottom part.

The branes, namely $h_1, h_2, h_5$ have no net chiral particle context with the rest of the branes and thus they constitute a "hidden" sector. Also the rest of gauge group factors associated with the "hidden" branes $h_1, h_2, h_5$ do not give rise to chiral particles as they
have no intersections with the Standard model particles and the rest of the branes. As we will see later the non-chiral matter arising from these branes is made massive by the introduction of Scherk-Schwarz deformations. Regarding the notation of the gauge groups related to the $h_1, h_2, h_5$ stacks we have considered that each one of these stacks is made from single stacks. To this end, non-chiral states that are coming from $h_1 h_1, h_5 h_5, h_1 h_5$ intersections should be properly considered as corresponding to the intersections $h_1^j h_1^j, h_5^j h_5^j, h_1^j h_5^j$, where $j = 1, ..., 36 \beta_1 \beta_2, v = 1, ..., 16 \beta_1 \beta_2$.

- Neutrino Masses

The bi-doublet ”h” Higgs multiplets of table (8) may be used in the process of electroweak symmetry breaking as the relevant Yukawa read

$$F_L \bar{F}_R h = \nu (u \bar{u} + \nu N) + \tilde{\nu} (d \bar{d} + e \bar{e}) , \ h = \text{diag}(\nu, \tilde{\nu})$$

(3.3)

giving masses to the up-quarks and neutrinos.

- Gauge mediation

The extra sector that is needed to cancel RR tadpoles introduces a vector-like sector that plays the role of the messenger supersymmetry breaking sector of the gauge mediated models [35] and transforms under the observable $O$ and the extra $h = \{h_1, ..., h_5\}$ sectors. The messenger extra sector needed for the RR tadpole cancellation is anomaly free and possess a different N=1 supersymmetry that the one respected by the top MSSM embedding. Lets us mention that the spectrum at the top of table (8) for $\beta_1 = \beta_2 = 1$, recovers the observable local MSSM spectrum of [24, 25] and also the “observable” sector of the MSSM spectrum of [40]. However, in our case the messenger exotic multiplets may become massive.

The questions that at this point remain are :

a) in which way we can get rid off the chiral fermions that appear in the messenger sector and which make vector-like pairs of N=1 multiplets with respect to the ”observable” SM $SU(3) \times SU(2)_w \ U(1)_Y$

b) how we can get rid off the adjoint fermions and gauginos that appear from brane-brane sectors

c) the non-chiral non-adjoint fermions that appear in intersections where the different participating intersecting branes are parallel in at least one complex plane.
Table 7: Solution to the RR tadpoles for the toroidal orientifold models of [25]. The N=1 MSSM chiral spectrum arises as part of Pati-Salam models in the top part of the table from intersections between a, b, c, d branes. Messenger multiplet states respecting a $N'=1$ supersymmetry arise from the intersections of the a, d branes with the $h_3, h_4$ branes.

### 3.2 Mass couplings for chiral fermions - in the Messenger Sector

Before discussing in this section the mass couplings of the extra chiral fermions that make vector-like exotics of table (8), let us examine the sector formed from open strings stretching between the branes $h_3, h_4$. It is easy to see that in this sector the intersection number $I_{h_3 h_4}$ vanishes and also that the N=2 supersymmetries preserved by this sector are the $r_2, r_3$. This implies - since also these is no gauged hypercharge for this multiplets - that this sector acts as the usual Higgs sector between the b, c branes and the net number of chiral fermions in this sector is zero. Hence there are $1/\beta_1\beta_2$ equal numbers (the intersection number across the remaining non-parallel tori ) of the N=1 massless
Table 8: Chiral spectrum of the Pati-Salam D6-brane N=1 Supersymmetric Standard Model. The bottom part, that includes the massive $N' = 1$ messenger fields that communicate supersymmetry breaking to the visible $N=1$ sector, completes the missing RR canceling sector of [24, 25]. Note that the notation L,R at the bottom of the table is related to chirality and $\rho = 1/3$, $\epsilon = \bar{\epsilon} = 1$.

Where the notation regards the gauge group representations $SU(4) \times SU(2)_{L} \times SU(2)_{R} \times U(1)_{h_{3}} \times U(1)_{h_{4}}$ and charges follow the table (8).

\[ h_{34}^{(1)} = (1, 1, 1; 1, 1)_{(0;1,-1)} \, , \, h_{34}^{(2)} = (1, 1, 1; 1, 1)_{(0;-1,1)} \, , \]  

where the branes are parallel across the first tori\(^{18}\). These fermions are singlets under the SM gauge group and thus can receive vevs. Observe the rather striking fact that these Higgs multiplets that help us to get rid of the extra vector-like exotics are equal in number to the usual MSSM multiplets H that appear between the b, c branes and their number depends also on the shape of the tori of the compactification.

\(^{18}\)Where the notation regards the gauge group representations $SU(4) \times SU(2)_{L} \times SU(2)_{R} \times U(1)_{h_{3}} \times U(1)_{h_{4}}$ and charges follow the table (8). 

| Matter | Repr. | $I_{ij}$'s | $Q_{a}$ | $Q_{h^{3}}$ | $Q_{h^{4}}$ |SYSY |
|---|---|---|---|---|---|---|
| $F_{L}$ | $3(4, 2, 1; 1, 1)$ | $(ab)$ | 1 | 0 | 0 | $r_{4}$ |
| $\tilde{F}_{R}$ | $3(\bar{4}, 1, 2; 1, 1)$ | $(ac)$ | $-1$ | 0 | 0 | $r_{4}$ |
| $\tilde{h}$ | $L_{1,2}^{-1}(1, 2, 2; 1, 1)$ | $(bc)$ | 0 | 0 | 0 | $r_{4}, r_{1}$ |
| $q_{R}^{1}$ | $3(\bar{4}, 1, 1; 1, 1)$ | $(ah_{3})$ | $-1$ | 1 | 0 | $r_{1}$ |
| $q_{R}^{2}$ | $3(\bar{4}, 1, 1; 1, 1)$ | $(ah_{3}*)$ | $-1$ | $-1$ | 0 | $r_{1}$ |
| $q_{L}^{1}$ | $3(4, 1, 1; 1, 1)$ | $(ah_{4})$ | 1 | 0 | $-1$ | $r_{1}$ |
| $q_{L}^{2}$ | $3(4, 1, 1; 1, 1)$ | $(ah_{4}*)$ | 1 | 0 | 1 | $r_{1}$ |
masses from the following superpotential couplings

\[(q^1_R)^{-1;1}_3 (q^3_L)^{1;0,-1}_H \cdot (q^2_R)^{-1;0}_3 (q^4_L)^{1;0,1}_H \]

The \(h^{(1)}_{34}, h^{(2)}_{34}\) represent flat directions in the effective potential. By assuming that their scalar components receive a vev all messenger multiplets become massive and disappear from the low energy spectrum. At this point - the models are at the string scale - the observable gauge group is that of (3.2), and the massless chiral multiplet fields are those of the top part of table (8) and also present there are non-chiral multiplets and the adjoint multiplets. Next we discuss how the non-chiral multiplets become massive.

### 3.3 Masses for non-chiral fermions

**All non-chiral fermions in the present Pati-Salam models get massive from the combined use of flat directions and Schwerk-Schwarz (SS) deformations.**

Non-chiral fermions (NCM) arise in sections that branes are parallel in at least one tori direction in some complex plane. The introduction of SS deformations in directions parallel to the directions that the intersecting D6-branes wrap give masses to non-chiral fermions that have odd n- “electric” wrapping numbers [61] in any representation. In this work, we introduce SS deformations in all three complex tori. We will examine the case that all the tori are not tilted, \(\beta_1 = \beta_2 = 1\).

NCM’s in sectors \(aa, aa^*, bb, bb^*, cc, cc^*, h^i h^i, h^i h^{i*}\), \(i=1,2,3,4,5\) get a mass from SS deformations from n-wrappings that are odd in at least one tori, namely the first, first, second, second, third, third, first, first, first, first, second, second, third, third, third, third, third, third ones respectively.

- **Masses for NCM’s in messenger sector**

There are also NCM from the sectors \(ah^1, ah^{1*}, ah^2, ah^{2*}, ah^5, ah^{5*}, bh^3, bh^{3*}, bh^4, bh^{4*}, ch^3, ch^{3*}, ch^4, ch^{4*}\), where the participating intersecting branes (PIB’s) are parallel in the first tori; \(bh^1, bh^{1*}, bh^3, bh^{3*}, bh^5, bh^{5*}, ch^2, ch^{2*}, ch^4, ch^{4*}\), where the PIB’s are parallel in the second tori; \(bh^2, bh^{2*}, bh^3, bh^{3*}, ch^1, ch^{1*}, ch^4, ch^{4*}, ch^5, ch^{5*}\), where the PIB’s are parallel in the third tori. All these NCM’s get masses – from SS deformations as they have odd n’s – but the ones from the intersections \(bh^2, bh^4, ch^2, ch^3\) (and their orientifold images) that appear in table (9). The latter ones’s could get masses from tree level flat directions.
Table 9: Non-chiral multiplet states from Pati-Salam models that also persist in their – by adjoint breaking – L-R symmetric models of section (4). The bottom part of the table includes the orientifold images of the top part multiplet states.

The following allowed superpotential couplings

\[
W \sim k_1 \Omega_1 \Omega_4 h_{24}^{(2)} + k_2 \Omega_2 \Omega_3 h^{(1)}_{23} + k_3 \Omega_5 \Omega_6 h_{24}^{(2)} + k_4 \Omega_6 \Omega_7 h_{23}^{(1)} + k_5 \Omega_9 \Omega_{12} h^{(1)}_{24} + k_6 \Omega_{10} \Omega_{11} h_{23}^{(2)} + k_7 \Omega_{13} \Omega_{16} h_{23}^{(1)} + k_8 \Omega_{14} \Omega_{15} h_{23}^{(2)} \quad (3.6)
\]

generate Dirac masses for the fermion pairs \(\Omega_1 \Omega_4, \Omega_2 \Omega_3, \Omega_5 \Omega_8, \Omega_6 \Omega_7, \Omega_9 \Omega_{12}, \Omega_{10} \Omega_{11}, \Omega_{13} \Omega_{16}, \Omega_{14} \Omega_{15}\). We have used the notation \(h_{ij}^{(l)}\) where by \(i, j\) we denote the N=2 multiplet appearing in the intersection between \(h^i\) and \(h^j\). By the superscript \(l\) when \(l = 1, l = 2\) we denote the states associated with the positive, negative intersection

| Intersection | States |
|--------------|--------|
| \(bh^2\)   | \(\Omega_1 : (1, 2, 1; 1, 1, 1, 1)(0; -1, 0, 0)\) |
|             | \(\Omega_2 : (1, 2; 1, 1, 1, 1)(0; 1, 0, 0)\) |
| \(bh^4\)   | \(\Omega_3 : (1, 2, 1; 1, 1, 1, 1)(0; 0, 0, -1)\) |
|             | \(\Omega_4 : (1, 2; 1, 1, 1, 1)(0; 0, 0, 1)\) |
| \(ch^2\)   | \(\Omega_5 : (1, 1, 2; 1, 1, 1, 1)(0; -1, 0, 0)\) |
|             | \(\Omega_6 : (1, 1, 2; 1, 1, 1, 1)(0; 1, 0, 0)\) |
| \(ch^3\)   | \(\Omega_7 : (1, 1, 2; 1, 1, 1, 1)(0; 0, -1, 0)\) |
|             | \(\Omega_8 : (1, 1, 2; 1, 1, 1, 1)(0; 0, 1, 0)\) |
| \(bh^{2*}\) | \(\Omega_9 : (1, 2; 1, 1, 1, 1)(0; -1, 0, 0)\) |
|             | \(\Omega_{10} : (1, 2; 1, 1, 1, 1)(0; 1, 0, 0)\) |
| \(bh^{4*}\) | \(\Omega_{11} : (1, 2; 1, 1, 1, 1)(0; 0, 0, -1)\) |
|             | \(\Omega_{12} : (1, 2; 1, 1, 1, 1)(0; 0, 0, 1)\) |
| \(ch^{2*}\) | \(\Omega_{13} : (1, 1, 2; 1, 1, 1, 1)(0; -1, 0, 0)\) |
|             | \(\Omega_{14} : (1, 1, 2; 1, 1, 1, 1)(0; 1, 0, 0)\) |
| \(ch^{3*}\) | \(\Omega_{15} : (1, 1, 2; 1, 1, 1, 1)(0; 0, 1, 0)\) |
|             | \(\Omega_{16} : (1, 1, 2; 1, 1, 1, 1)(0; 0, -1, 0)\) |
numbers $I_{ij}$ respectively. Also $k_i$ are numerical coefficients that can be calculated by the use of string amplitudes and give rise to Yukawa couplings for fermions in the effective theory. Let us consider for example the superpotential term $k_1 \Omega_1 \Omega_4 h_{24}^{(2)}$ involving the multiplets $h_{24}^{(i)}$ that respects a N=2 susy. The $\Omega_1$ represents a N=1 multiplet preserving the supersymmetry $r_2$; $\Omega_4$ represents a N=2 multiplet preserving the supersymmetries $r_2, r_3$; $h_{24}^{(2)}$ is preserving the susys $r_2, r_4$; they all share the common N=1 susy $r_2$. Also all gauginos and also adjoint multiplets get massive as the SS deformations act in all tori and there are odd $n_i$ in all branes in at least one tori direction. Hence, only the observable MSSM multiplets seen at the top of table (8) remain massless at the string scale.

- **Breaking to the MSSM**

The initial U(4) gauge symmetry breaks to $SU(4) \times U(1)_a$. The $U(1)_a$ gauge symmetry is getting massive from its couplings to RR fields and subsequently the $SU(4) \times SU(2)_L \times SU(2)_R$ symmetry breaks with the help of the Higgs multiplets states - localized in the intersection $ac^*$

$$(4,1,2), \ (4,1,2)$$

(3.7)

to the $SU(3) \times SU(2)_L \times U(1)_Y$. We notice that $SU(4)$ breaks $^{19}$ to $SU(3) \times U(1)_{B-L}$ and $SU(2)_R$ to $I_{3R}$ by (3.7). We also note that the non-chiral multiplet states of table (9) also appear in the left-right models of the next section - with the addition of non-chiral states states that are formed between the intersections of the d- U(1) brane and the other branes - and they get masses from the same mechanism described in this section.

### 4 Flow to three generation Left-Right Symmetric Models and the MSSM

We have seen that three generation Pati-Salam (PS) models of the previous section can have all its chiral and non-chiral states (beyond the MSSM) made massive leaving only the usual Pati-Salam spectrum at $M_s$ and the SM at low energies. In this section, the breaking of these PS models to a three generation left-right $SU(3) \times SU(2)_L \times SU(2)_R$ (L-R) symmetric model with all its exotics made massive and the subsequent breaking of

$^{19}$See also [15] for the breaking of Pati-Salam models that have sextets in their spectrum; also 2nd ref. of [2] for the breaking of the PS MSSM-like models in the context of $Z_2 \times Z_2$ orientifolds with intersecting D6-branes, where however massless exotics also survive to low energies.
the L-R to the SM is described.

- Pati-Salam adjoint breaking to left-right symmetric models

The L-R model can be derived from the Pati-Salam models of the previous section by considering the adjoint breaking of the $U(4)_c \to U(3)_c \times U(1)_d$. The spectrum of these models can be seen in table (10) for $\epsilon = \tilde{\epsilon} = 1$ and $\rho = 1/3$ while the $U(1)$'s appearing in tables (10), namely $Q_a$, $Q_d$ are the $U(1)$'s that are inside $U(3)$, $U(4)$ respectively.

After examining the Green-Schwarz couplings

$$n' \eta K m^I \int B_2^I \wedge F_a, \quad n_a \eta_n a \int C^o \wedge F_a \wedge F_a, \quad n' m^J m^K \int C^I \wedge F \wedge F$$

we find that the mixed $U(1)$ gauge anomalies cancel since the following couplings of the RR fields to $U(1)$'s exist: $B_2^3 \wedge (3\epsilon \beta)(3F_a + F_d)$; $(3\epsilon \beta_1)(3F_a + F_d) \wedge B_2^2$. We find that the surviving massless the Green-Schwarz mechanism $U(1)$'s are the

$$U(1)^{(3)} = F^{h_3} - F^{h_4}, \quad U(1)^{(4)} = F^{h_3} + F^{h_4}$$

and the hypercharge

$$Q_Y \equiv U(1)^{(1)} = (1/6) F_a - (1/2) F_d$$

In table (10) we can see the hypercharge assignment of the left-right symmetric models. [In appendix A there is a different hypercharge assignment for the left-right symmetric models which breaks to the SM at low energies and also possess massive chiral and non-chiral exotics.]. The $U(1)^{(3)}$, $U(1)^{(4)}$ may be broken by vevs of the gauge singlets $\langle \tilde{\phi}_R^1 \rangle$, $\langle \tilde{\phi}_L^3 \rangle$ respectively.

- Global symmetries

The global symmetries that exist in the Standard Model do exist in the left-right symmetric models and get identified in terms of the $U(1)$ symmetries set out by the D6-brane configurations of table (10). Hence baryon number ($B$) is identified as $Q_a = 3B$ and the lepton number ($L$) could practically identified as $Q_d = L$ as the $R$-antimatter multiplet states accommodate both the right leptons.

Note that the spectrum of table (10) is also valid for non-zero NS B-field, that makes the tori tilted along the second and the third tori. We also note that the higgsino N=1 multiplet states $\tilde{H}_a$, $\tilde{H}_d$ are part of the non-chiral spectrum in the $bc$, $bc^*$ sectors respectively and they belong to of a $(2,2)$ representation of $SU(2)_b \times SU(2)_c$. The superpotential

\[20\text{As it is already included in the solutions to the RR tadpoles.} \]
part describing the generation of masses in the “observable” part reads

\[ W^{obs} \sim l_1 Q_L H(i\tau_2)Q_R + l_2 LH(i\tau_2)R \]  

(4.4)

where \( l_1, l_2 \) may be determined by string amplitudes \([5]\); \( i\tau_2 \) is used for conventional reasons in relation to the SM physics (see for example \([70]\)). The representation content of the SM matter is accommodated as

\[ Q = \begin{pmatrix} u \\ d \end{pmatrix}, \quad R = \begin{pmatrix} d^c \\ u^c \end{pmatrix} \]  

(4.5)

where the Higgs fields are accommodated as

\[ H = \begin{pmatrix} \phi^0_{11} & \phi^+_{12} \\ \phi^0_{21} & \phi^0_{22} \end{pmatrix} \]  

(4.6)

and may obtain vevs as \( \langle \phi^0_{11} \rangle = u_1, \langle \phi^0_{22} \rangle = u_2 \).
• Masses for non-chiral matter (NCM)

NCM for the present L-R models appears in bifundamentals arises in $ad$, $ah_1$, $ah_2$, $ah_5$, $dh_1$, $dh_2$, $dh_5$, $h_1h_2$, $h_1h_5$, $h_2h_5$ sectors and it is getting massive by the introduction of SS deformations as described in the previous section. The left-right symmetric models have - compared to the Pati-Salam models of tables (8), (7), (6) additional non-chiral sectors coming from the $ad$, $dh_1$, $dh_2$, $dh_5$ sectors. The latter extra matter is also getting massive by the introduced SS deformations in all tori as there are odd n-wrappings in at least one complex tori in each intersection.

Other NCM that is also present are the one accommodated in the states of table (9) which are getting massive from the same superpotential terms as the ones of eqn. (3.6)

• Masses for the exotic vector-like multiplets in SUSY breaking messenger sector

The present L-R models have also (as the PS models of the previous section) a stable proton as baryon number is a gauged symmetry. The most obvious phenomenological problem that any string model from D-branes is facing is the presence of extra massless exotics that survive massless to low energies.

In the present L-R constructions there are superpotential mass terms for all the chiral beyond the MSSM fermions – namely the extra messenger exotics of the bottom of table (10) $\tilde{q}_R^1, \tilde{q}_R^2, \tilde{q}_L^3, \tilde{q}_L^4, \tilde{\phi}_R^1, \tilde{\phi}_R^2, \tilde{\phi}_L^3, \tilde{\phi}_R^4$. They appear as follows

$$W_{\text{messenger}} = \lambda_1 \tilde{q}_R^1 \tilde{q}_L^3 h_{34}^{(2)} + \lambda_2 \tilde{q}_R^2 \tilde{q}_L^4 h_{34}^{(1)} + \lambda_3 \tilde{\phi}_R^1 \tilde{\phi}_L^3 h_{34}^{(2)} + \lambda_4 \tilde{\phi}_R^2 \tilde{\phi}_R^4 h_{34}^{(1)} \quad (4.7)$$

The multiplets $h_{34}^{(1)}, h_{34}^{(2)}$ have a gauge singlet scalar direction. Assuming that it gets a vev the couplings (4.7) helps the fermion pairs $\tilde{q}_R^1 \tilde{q}_L^3, \tilde{q}_R^2 \tilde{q}_L^4, \tilde{\phi}_R^1 \tilde{\phi}_L^3, \tilde{\phi}_R^2 \tilde{\phi}_R^4$ to form massive Dirac mass eigenstates with masses respectively $^{21}$

$$m_1 = \lambda_1 \langle 0 | h_{34}^{(2)} | 0 \rangle, \quad m_2 = \lambda_2 \langle 0 | h_{34}^{(1)} | 0 \rangle, \quad m_3 = \lambda_3 \langle 0 | h_{34}^{(2)} | 0 \rangle, \quad m_4 = \lambda_4 \langle 0 | h_{34}^{(1)} | 0 \rangle. \quad (4.8)$$

We note that the sector made of the H-multiplet preserves a N=2 SUSY. In particular it preserves the N=1 SUSY preserved by the observable MSSM and also the $N' = 1$ respected by the messenger sector.

• Breaking to the MSSM

$^{21}$We assume that our vacuum is stable locally
In the L-R symmetric models of table (10), the gauge group $SU(2)_R$ arises from the c-brane which is placed at a point of the moduli space where it is its own orientifold image. The breaking to the $U(2)_R \rightarrow U(1)_c$ and thus the L-R symmetric models to the MSSM can be achieved by considering general positions and Wilson lines for brane c that correspond to geometrical separation along the first torus (where the branes are parallel) and Wilson line "phases" along the one cycle in the first torus respectively; in turn they are associated to the real and imaginary part of the $\mu$-parameter. It is easily confirmed that in this case, the $(2, 2)$ N=2 chiral multiplet $Q_R$ gives rise to two N=1 multiplets $(2, 1)$, $(2, -1)$ charged under $SU(2) \times U(1)_Y$ that gets identified with $U_R$, $D_R$ while the $R$ multiplet in table (10) gives rise to $E_R$, $N_R$ as they appear in table (5).

The extra messenger chiral sector remains the one given in the bottom of table (10); the corresponding states are getting massive once the $h^{(l)}_{34}$, $l=1, 2$ multiplets receive a vev. Thus at this stage – being at the string scale – only the MSSM multiplets, seen at the top part of table (5), survive massless. The only question remaining at this point is to show in which way the MSSM sparticle masses after susy breaking. This may be achieved in the next section by the use of D-term breaking.

### 4.1 D-term Supersymmetry breaking and Gauge Mediation

In the usual N=1 SUSY models of gauge mediation [35], the spectrum consists of the MSSM together with an extra hidden sector $G_H$ that contains vector-like messenger particles $\phi_H$, $\bar{\phi}_H$ charged under the observable MSSM and the $G_H$ gauge group. SUSY is broken in the hidden sector when a spurion superfield $X_H$ gets a vev through the superpotential

$$W^{(N=1)} = W_{obs} + \phi_H \bar{\phi}_H X_H$$

and where the $W_{obs}$ and the susy breaking part of the superpotential W respect the same N=1 susy. In field theory terms, the squarks and sleptons get masses from two loop effects while gauginos get masses from one loop effects where messenger states circulate in the loops. On the other hand in the context of string theory models from intersecting branes it was argued in a field theoretical basis that in non-susy models - localized within 4D toroidal orientifolds of IIA with intersecting D6-branes - that exhibit the quasi-susy spectrum (By definition models where MSSM particles are subject to different N=1 supersymmetries) similar effects are expected [19].

---

22See also related comments [25].
The supersymmetry observed in the gauge mediated construction of table (10) can be broken by slightly varying the complex structure of the models. In the low energy effective theory this procedure can be seen as turning on Fayet-Iliopoulos terms for the U(1) fields [19]. Such a procedure has been applied to quasi-susy models in [27]. In the present modes, which can avoid the non-chiral fermions once Scherk-Schwarz breaking is included it appears that SS deformations do not affect the angles between the branes in the RR tadpoles. In general for two D6-branes, that exhibit N=1 SUSY limit at their intersection, the scalar potential contribution appears as

$$V_{FI} = \frac{1}{2\tilde{g}_a^2}(\sum_i q_a^i|\phi|^2 + \xi_a)^2 .$$

(4.10)

Lets us be more specific and consider a slight departure in the complex structure of eqn.(2.18) that defines the supersymmetric limit

$$\beta_1 \chi_2' = \beta_1 \chi_2 + k_a ,$$

(4.11)

where $k_a << 1$ and where $k_a$ could be also written as $k_a = \beta_1 p_a$, where $p_a$ represents the actual variation in the complex structure. When the torus is untilted then $k_a \equiv p_a$. Lets us calculate as an example the angle $^{23}$ between the D6-branes $a, b$ along the second torus (seen in table (12). It is given by

$$\theta_{ab}^2 = tan^{-1}((3\epsilon\rho^2) \beta_1 \chi_2) \approx tan^{-1}((3\epsilon\rho^2) \beta_1 \chi_2) + \frac{(3\epsilon\rho^2)k_a}{1 + (3\epsilon\rho^2\beta_1 \chi_2)^2}$$

(4.12)

$$= a_1 + \delta_a ,$$

(4.13)

where $a_1 = tan^{-1}(3\epsilon\rho^2\beta_1 \chi_2)$, $\delta_a = ((3\epsilon\rho^2)k_a)/(1 + (3\epsilon\rho^2\beta_1 \chi_2)^2)$.

- Universal squark and slepton masses

We observe that all string (mass)$^2$ of the squark and slepton masses are positive and hence colour and charge breaking minima may be avoided. The mass$^2$ value of the Higgses does not get corrected by the variation in the complex structure (or the variation of the magnetic fields in the T-dual language with D9-branes and fluxes) but retains only its dependence on the distance between the parallel D6-branes in the second tori. We notice that between the SU(2) $b, c$ branes there is a N=2 preserving sector which possesses two scalar Higgs states and also two fermionic Higgsino partners, all with masses$^2 \frac{M_s^2}{4\pi^2}l_2$ at tree level. The masses of the Higgsinos depend on the distance $l_2$ between the branes

---

$^{23}$We use the D6-brane wrappings as they are seen in table (13) of appendix B.
Table 11: Angle structure between the D6-branes and the orientifold axis for the left-right symmetric model. We exhibit only branes that give rise to chiral matter structure.

| Brane | $(\theta^1_a)$ | $(\theta^2_a)$ | $(\theta^3_a)$ |
|-------|----------------|----------------|----------------|
| $a$   | 0              | $a_1 + \delta_a$ | $-a_1$         |
| $b$   | $\frac{\pi}{2}$ | 0              | $-\frac{\pi}{2}$ |
| $c$   | $\frac{\pi}{2}$ | $-\frac{\pi}{2}$ | 0              |
| $d$   | 0              | $a_1 + \delta_d$ | $-a_1$         |
| $h_3$ | $\frac{\pi}{2}$ | 0              | $\frac{\pi}{2}$ |
| $h_4$ | $\frac{\pi}{2}$ | $\frac{\pi}{2}$ | 0              |

Table 11: Angle structure between the D6-branes and the orientifold axis for the left-right symmetric model. We exhibit only branes that give rise to chiral matter structure.

b, c which is a open string modulus and thus can be be made small. Furthermore, the sparticle masses get a universal value and as a result we expect that flavour changing neutral currents (FCNC) to be strongly suppressed \(^{24}\). The squarks and sleptons masses seen in table (12) are the exact string expressions as they are calculated by the use of the appropriate mass operators (2.8). These masses could be calculated in the limit of small deviation in the complex structure of the second torus. In this case, one can also calculate the field theory limit of the supergravity (sugra) approximation \(k_a << 1\) (e.g. if for the radii \(R_1 >> R_2\)) that is coming from Fayet-Iliopoulos (FI) terms and find that the string and sugra approximations of the scalar masses at the intersections coincide \([19]\). The string expression is more generic. Thus the scalar masses at an intersection between two branes coming in the limit of small variation in the complex structure could be understood as coming from FI terms in the effective theory.

5 Split Supersymmetry in String theory and MSSM

The Split supersymmetry scenario (SSS) \([46]\) was proposed as an alternative possibility for generating a signal for LHC. As in a global susy or supergravity context there is no underlying principle for constraining the mass of the particles, in the original SSS

\(^{24}\) Universal soft terms masses for the sparticles also appear (as a result of the hypothesis of parametrizing the unknown susy breaking effects in the scalar potential expansion) on another occasion in string theory, in N=1 heterotic orbifold compactifications of the heterotic string where the dilaton dominant source of susy breaking \([69]\); and also in models with fluxes - 1st ref. of \([62]\).
| Sector | $(\theta_a^1, \theta_a^2, \theta_a^3)$ | sparticle | (mass)$^2$ |
|--------|----------------------------------|-----------|-----------|
| $(ab)$ | $(-\frac{\pi}{2}, \alpha_1 + \delta_a, -\alpha_1 + \frac{\pi}{2})$ | $1 \times \tilde{Q}_L$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $ac$   | $(-\frac{\pi}{2}, \alpha_1 + \delta_a + \frac{\pi}{2}, -\alpha_1)$ | $3 \times \tilde{Q}_R$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $db$   | $(-\frac{\pi}{2}, \alpha_1 + \delta_a, -\alpha_1 + \frac{\pi}{2})$ | $3 \times \tilde{L}$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $dc$   | $(-\frac{\pi}{2}, \alpha_1 + \delta_d + \frac{\pi}{2}, -\alpha_1)$ | $3 \times \tilde{R}$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $bc$   | $(0, \frac{\pi}{2}, -\frac{\pi}{2})$ | $\frac{1}{\beta_1 \beta_2} \times \tilde{H}$ | $\frac{M_a^2 l_2}{4\pi^2}$ |
| $ah_3$ | $(-\frac{\pi}{2}, \alpha_1 + \delta_a, -\alpha_1 - \frac{\pi}{2})$ | $3 \times \tilde{q}^1_R$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $dh_3$ | $(-\frac{\pi}{2}, \alpha_1 + \delta_d, -\alpha_1 - \frac{\pi}{2})$ | $3 \times \tilde{\phi}^1_R$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $ah^*_3$ | $(\frac{\pi}{2}, \alpha_1 + \delta_a, -\alpha_1 + \frac{\pi}{2})$ | $3 \times \tilde{q}^2_R$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $dh^*_3$ | $(\frac{\pi}{2}, \alpha_1 + \delta_d, -\alpha_1 + \frac{\pi}{2})$ | $3 \times \tilde{\phi}^2_R$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $ah_4$ | $(\frac{\pi}{2}, \alpha_1 + \delta_0 - \frac{\pi}{2}, -\alpha_1)$ | $3 \times \tilde{q}^3_L$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $dh_4$ | $(\frac{\pi}{2}, \alpha_1 + \delta_d - \frac{\pi}{2}, -\alpha_1)$ | $3 \times \tilde{\phi}^3_L$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $ah^*_4$ | $(\frac{\pi}{2}, \alpha_1 + \delta_a + \frac{\pi}{2}, -\alpha_1)$ | $3 \times \tilde{q}^4_L$ | $\frac{M_a^2}{2}(\delta_a)$ |
| $dh^*_4$ | $(\frac{\pi}{2}, \alpha_1 + \delta_d + \frac{\pi}{2}, -\alpha_1)$ | $3 \times \tilde{\phi}^4_L$ | $\frac{M_a^2}{2}(\delta_a)$ |

Table 12: Universal Squark and Slepton masses from FI-terms in the Left-Right Symmetric model.
appearance [46], it was assumed that sparticles of the MSSM are massive at the UV high scale \( m_s^H \) where supersymmetry is getting broken, while gauginos and higgsinos are the TeV scale in order to retain the unification of the gauge couplings at the \( 10^{16} \) GeV. Furthermore it was assumed [46] that one of the Higgs doublets of the MSSM is finely tuned to be light and below \( m_s^H \). Exactly at the same time it was argued that in the context of string theory split susy models should have slightly different characteristics respectively [48, 49, 50]. These properties that can still keep the nice features that are needed for realistic string model building are: proton stability, partial unification of two of the gauge couplings at the high scale [48, 49, 50] and either light gauginos [48]; and higgsinos that be either light [48] or anywhere in the range \( M_s \) to \( M_s [49, 50] \); and in all cases massive squarks and sleptons [48, 49, 50]. In stringy MSSM models of this work proton is stable as baryon number is a gauged symmetry. Next we discuss the rest of the Stringy Split SUSY criteria set out in [48, 49, 50] and which concern gauge coupling unification of SU(3) and SU(2) gauge couplings accompanied by appropriate Weinberg angle \( \sin^2(\theta_W) \) at the unification scale; it turns out that gaugino and higgsino masses are different from the standard split susy scenario of the local theory [46].

### 5.1 Gauge Coupling Unification

Necessary ingredient of any split theory is that the successful prediction of unification of gauge interactions [46] is retained. In [47], based on examples on D-brane inspired models, it was argued that models of split susy in a string theory content should also accommodate the successful GUT prediction \( \sin^2(\theta_W) = 3/8 \) with equal SU(3) and SU(2) gauge couplings at the unification scale. Explicit realizations of non-supersymmetric string models with intersecting D6-branes where \( \sin^2(\theta_W) = 3/8 \) and equality of SU(3) and SU(2) gauge couplings had appear in [49], [50]. The latter models are based on \( Z_3 \times Z_3 \) orientifold compactifications of type IIA theory. These models are non-susy and have no supersymmetry preserved at the different intersections in sharp contrast with the models presented in this work where the appearance of N=1 susy on intersections is explicit. In the present models before the breaking of the left-right symmetric models to the MSSM at \( M_s \), the gauge group at the observable sector is an \( SU(3)_c \times Sp(2)_W \times U(1)_a \times U(1)_d \times SU(2)_R \). When all the tori along the three complex dimensions are orthogonal the unification of the gauge couplings of the “observable” SM spectrum is described by the wrappings numbers of the top part of table (7) when \( \beta_1 = \beta_2 = 1 \), and has been studied in [36]. In this case it was found that \( \sin^2(\theta_W) = \frac{3}{5} \) at the string scale as in SUSY SU(5).
In this section, we study the gauge coupling unification of the MSSM’s coming from the breaking of $SU(4)_C \times SU(2)_W \times SU(2)_R$; the tori could be also tilted. The general background is a four dimensional IIA orientifold compactification on a six dimensional tori in the presence of D6$_a$-brane wrapping at angles on a general factorizable 3-cycle described by the wrappings $(n,m)_a$. Using the general RR solutions of table (6) we find that the volumes obey [24]

$$V_a = V_d, \quad V_b = V_c$$

(5.1)

The relation of gauge couplings in terms of wrappings and the complex moduli $\chi_i$ along the three tori, are given by the generic relation [36]

$$\frac{1}{a_a} = \frac{1}{k_a} \cdot \frac{M_{pl}}{2\sqrt{2}} \cdot \frac{V_a}{M_s \sqrt{V_6}},$$

(5.2)

where

$$a_a^{-1} = \frac{1}{\sqrt{2k_a}} \frac{M_{pl}}{M_s} \times \left( (n^1 n^2 n^3)_a \sqrt{\frac{1}{U_1 U_2 U_3}} - (n^1 m^2 m^3)_a \sqrt{\frac{U_2 U_3}{U_1}} - (m^1 n^2 m^3)_a \sqrt{\frac{U_1 U_3}{U_2}} - (m^1 m^2 n^3)_a \sqrt{\frac{U_1 U_3}{U_2}} \right)$$

(5.3)

and the value of $k_a$ is set out by the gauge group on the brane $a$ and is $k_a = 1$ when the gauge group is a U(N) one and $k_a = 2$ when the gauge group is an Sp(N). We derive in the general case of tilted tori, the gauge couplings at the string scale of the MSSM. We use the more general solution to the RR tadpoles seen in table (6) of appendix B for the Pati-Salam models that accommodates the parameter $\rho = 1, 1/3$. We set

$$\frac{1}{\sqrt{2k_a}} \frac{M_{pl}}{M_s} = K_0, \quad \beta_2 \chi_2 = \beta_3 \chi_3$$

(5.4)

The gauge couplings are

$$\frac{1}{a_s} = K_0 \sqrt{\beta_1} \frac{1}{\sqrt{\beta_2 \chi_1}} \left( \frac{1}{\rho^2 \chi_3} + 9\rho^2 \beta_2 \chi_3 \right)$$

$$\frac{1}{a_W} = \frac{K_0}{2} \frac{\epsilon}{\beta_1 \beta_2 \sqrt{\chi_1}}$$

$$\frac{1}{a_c} = K_0 \frac{\epsilon}{\sqrt{\beta_1 \beta_2 \chi_1}}$$

(5.5)

where $a_c = \frac{1}{2} a_W$ and the gauge group for general positions and Wilson lines for c-brane is an $SU(3)_C \times Sp(2)_W \times U(1)_Y$ ($Sp(2) \equiv SU(2)$).
The gauge couplings of the $SU(3)_c \times SU(2)_w \times U(1)_Y$ gauge group are

$$a_s = a_d, \quad a_c = \frac{1}{2} a_b = \frac{1}{2} a_w$$  \hspace{1cm} (5.6)

Also since

$$\frac{1}{a_Y} = \frac{1}{6} a_s + \frac{1}{2} a_c + \frac{1}{2} a_d$$  \hspace{1cm} (5.7)

we find that at the string scale the gauge couplings obey

$$\frac{1}{a_Y} \equiv \frac{2}{3} a_s + \frac{1}{a_w}.$$  \hspace{1cm} (5.8)

From (5.8) we derive the value of the Weinberg angle at the string scale

$$\sin^2(\theta_W)_{M_s} = \frac{3 a_s}{6 a_s + 2 a_w}.$$  \hspace{1cm} (5.9)

The free parameters in this equation are the values of the complex structure moduli $\chi_i, \ i=1, 2, 3$. Complex structure moduli, as was first noticed in [15] [see eqn.(4.37) in hep-th/0203187] in the context of 4D toroidal orientifold intersecting brane worlds [3], the supersymmetry conditions in chiral fermion intersections could be all fixed to a certain value without the presence of fluxes. These models were based on Pati-Salam $SU(4) \times SU(2)_L \times SU(2)_R$ type of constructions and in the same 4D string backgrounds of the present work.

- **Fixing complex structure moduli via $\sin^2(\theta_w)$ at the String scale** $M_s$

In the present constructions complex structure moduli could be further constrained - but not completely fixed - using the values of the gauge couplings at the string unification scale. From (5.9) we observe that successful unification relations such that of SU(5) GUT could be also derived within the framework of Pati-Salam GUTs from intersecting branes. Hence by demanding that

$$a_s = a_w,$$  \hspace{1cm} (5.10)

using appendix C, we can reproduce the GUT predictions for $\sin^2(\theta_w)$ at the unification/string scale as

$$a_y = a_w = \frac{5}{3} a_Y, \quad \sin^2(\theta_w)(M_s) = \frac{3}{8}.$$  \hspace{1cm} (5.11)

So the MSSM models behave as a hidden SU(5) at the unification GUT/String scale of the SU(3) and SU(2) couplings.

Eqn. (5.10) could be used to constrain the complex structure moduli $\chi_1$ - along the first
torus - in terms of the values of the rest of moduli along the 2nd and 3rd tori that are related by N=1 SUSY as in (2.18). From (5.10) we get a second order equation in the form \( \alpha \chi^2 + b \chi + \gamma = 0 \), on the \( \chi_3 \) modulus, where

\[
\hat{a} = 9 \rho^4 \beta_2, \quad \hat{b} = \chi_1 \rho^2 (2 \beta_1 \sqrt{\beta_2})^{1/2}, \quad \hat{\gamma} = 1
\] (5.12)

Demanding positivity \( \hat{b}^2 - 4\hat{\alpha}\hat{\gamma} \) of the square root (to demand only real values for the modulus \( \chi_1 \)) so at least one root is positive \( ^{25} \), we get the constraint

\[
\chi_1 = \frac{R_2}{R_1} > \frac{\beta_1 \beta_2}{3}
\] (5.13)

The present models reveal the presence of the successful SU(5) GUT prediction \( \sin^2(\theta_W) = 3/8 \), when all tori were orthogonal (this was also noticed in \( ^{26} \)) and also when tilted tori are involved. It appears that the presence of a ”hidden” SU(5) inside Pati-Salam models that is behind (5.11) is independent of the existence or not of tilted tori.

- **Gauge coupling unification at \( M_s \)**

At this part, we determine the value of the string scale in which the gauge couplings of the SM unify \( ^{27} \). The tree level relations of the gauge couplings of the SM at the string scale read

\[
\frac{1}{a_s(\mu)} = \frac{1}{a_s} + \frac{b_3}{2\pi} \ln(\frac{\mu}{M_s})
\]

\[
\frac{\sin^2(\theta_w)(\mu)}{a(\mu)} = \frac{1}{a_w} + \frac{b_2}{2\pi} \ln(\frac{\mu}{M_s})
\]

\[
\frac{\cos^2(\theta_w)(\mu)}{a(\mu)} = \frac{1}{a_Y} + \frac{b_1}{2\pi} \ln(\frac{\mu}{M_s})
\] (5.14)

By using the relation (2.18) we obtain that at the string scale

\[
\frac{2}{3} \frac{1}{a_s(\mu)} + \frac{2\sin^2(\theta_w)(\mu) - 1}{2\pi} = \frac{\hat{B}}{2\pi} \ln \left( \frac{\mu}{2\pi} \right)
\] (5.15)

where

\[
\hat{B} = \frac{2}{3} b_3 + b_2 - b_1
\] (5.16)

\( ^{25} \)namely the \( -\hat{b}/2\hat{a} + 1/(2\hat{a})\sqrt{\hat{b}^2 - 4\hat{\alpha}\hat{\gamma}} \)

\( ^{26} \)See also first ref. of \([8]\) for an attempt to construct N=1 Pati-Salam MSSM-like models where however \( \sin^2(\theta_w) \) at the GUT/string scale does not possess a hidden SU(5), e.g. \( \sin^2(\theta_w) \neq 3/8 \).

\( ^{27} \)We briefly reproduce a relevant discussion from \([36]\).
After the breaking of the left-right symmetric model at the MSSM, the massless content of our models at the string scale is the one given by the MSSM in addition to the extra matter as it seen in the bottom of table (10). As the extra exotic non-chiral matter is getting massive by the vevs of the flat directions $h_k^{(l)}$, $l=1,2$ which are of the order of the string scale, at $M_s$ only the MSSM remain massless for which $\hat{B} = 12$. Thus the unification scale (US) - at which the SU(3) and SU(2) couplings unify for the present MSSM models is just

$$M_s = M_{GUT} = 2.04 \times 10^{16} \text{ GeV} \quad (5.17)$$

This scale is independent of the number of Higgsino multiplets in the MSSM as their contribution within the combination $\hat{B}$ cancels out as

$$b_3 = -2n_G + 9 \quad , \quad b_2 = -2n_G - 6 - n_u \frac{1}{2} - n_d \frac{1}{2} \quad , \quad b_1 = -n_G \frac{10}{3} - n_u \frac{1}{2} - n_d \frac{1}{2} \quad (5.18)$$

and $n_G = 3$ in the present models.

### 5.2 Gaugino, Higgsino and Dark Matter Split SUSY candidates

• Gaugino masses

In [48] it was argued - following arguments on Scherk-Schwarz compactifications in the absence of branes intersecting at angles [45] 28 - that gauginos in string theory should receive TeV scale masses. On the contrary in [48, 50], it was suggested that gauginos should only receive masses of the order of the string scale even though at tree level in intersecting brane models appear to be massless. The latter suggestion originated from field theory arguments for intersecting brane models in toroidal orientifolds 29, where the $D6_a$ gauginos receive non-zero masses from one loop corrections from massive N=1 hypermultiplets running in the loops, while the extension to the $Z_N, Z_N \times Z_M$ orbifolds case comes by considering the contribution from the different intersections when taking into account the orbits in the orbifold brane structure. The stringy one loop correction for the toroidal orientifold case was verified recently by a string calculation [73].

• Higgsino masses

28 where repeating the argument [45] $m_{1/2} \sim m_3^{1/2}/M_p^2$, $m_{1/2} \sim \text{TeV}$ when $m_3/2 \sim 10^{13-14}$ GeV. Notice that it is also allowed $m_{1/2} \sim 10^{10}, 10^{19}$ GeV when $m_3/2 \sim 10^{16}, 10^{19}$ GeV respectively.

29 As they appear in the appendix of [14].
It was suggested [48] that in the context of string theory, a necessary condition for having a light Higgs multiplet in the spectrum would be to have unbroken supersymmetry in the Higgs sector even after supersymmetry breaking, that is tree level massless Higgs multiplets in the spectrum. Moreover it was also suggested [48] that it was likely - in the sense when such explicit models will be constructed - given the similarity of the Higgs doublets to sleptons that even sleptons could be found massless at tree level. Different suggestions for the exact mass of the Higgsinos in string models had appear in [49, 50]. In [49, 50] it was suggested that in non-susy models with no susy at intersections on toroidal orientifolds Higgsinos could be in principle anywhere between the weak and the string scale. In the latter orbifolded orientifolds the chiral spectrum consists of only the MSSM fermions - in addition to three pairs of massive non-chiral colour exotics that become massive from existing Yukawa couplings - and where the Higgsinos make Dirac pairs and couple to gauge singlets at tree level. See for example section (8.1) in [49] and eqn. (6.7) in [50]. Thus the models of [49, 50] - where supersymmetry is already broken by construction at the string scale $M_s$ in the open string sector where the SM is localized, accommodate all the split susy criteria namely:

a) equality of SU(3) and SU(2) gauge couplings and correct $\sin^2\theta_W = 3/8$ at the unification/string scale,

b) accommodating massive gauginos and higgsinos that could be anywhere from $M_z$ to $M_s$ ($M_z \leq \mu \leq M_s$).

Thus interesting conclusions on the magnitude of Higgsino masses in string models with the split susy characteristics were obtained. However since in these models susy is broken by construction in the open sector high at $M_s$ and there is no explicit N=1 SUSY preserved in their Higgs sector, the solution of the gauge hierarchy problem in the Higgs sector of these models is not guaranteed.

On the contrary in the MSSM models of this work - where N=2 supersymmetry is preserved in the N=2 sector - after the breaking of supersymmetry, Higgsinos do not receive any correction (modulo the distance $l_2$ between the branes b, c that can be made very small and is an open string modulus) at tree level due to the D-term FI breaking as we have seen in the previous section. They do however receive a one loop correction of the order $\sim \delta_a^2 M_s$ [73].

\textsuperscript{30}neglecting the non-chiral matter
\textsuperscript{31}In addition the models of [49, 50] have also present the non-chiral matter that appears in models with intersecting D6-branes.
• Dark matter candidates

One of the most difficult cosmological problems is the composition of the present day dark matter (DM) density $\Omega_{DM}$ in the universe. Weakly interacting massive particles (WIMPS) are good candidates for generating $\Omega_{DM}$. In the context of supersymmetric or SUGRA theories where the breaking of supersymmetry used to be at the TeV to avoid the hierarchy problem, neutralinos $X^{\text{neu}}$ \(^{32}\) are popular candidates for generating the cold dark matter of the universe (see [74]). They are expected to be non-relativistic in the present epoch and are eigenstates of a mixture of binos $\tilde{B}$, winos $\tilde{W}_3$ (the superpartners of $B, W_3$) and neutral components of higgsinos where

$$X^{\text{neu}} = N_{11} \tilde{B} + N_{13} \tilde{W}_3 + N_{13} H^0_1 + N_{14} H^0_2$$  \hspace{1cm} (5.19)$$

E.g. the higgsino fraction is defined as $N_H = |N_{13}|^2 + |N_{14}|^2$ and is used to distinguish the amount of higgsinos mixing in the $X^{\text{neu}}$ composition, e.g. if one could find that e.g. $N_H > 0.9$ then such a mixing could be considered as mostly higgsino.

In the present models, the introduction of SS breaking which makes massive all the non-chiral matter fields of the models, make also massive and of the order of the string scale the gauginos which belong to branes with odd "electric" n-wrappings. that are charged under the SM gauge group. Which gauginos could become massive also depends on the number of tilted tori (or the NS B-field) that enter the RR tadpoles. Gauginos that are associated with a D6-brane for which SS deformations act in some complex tori receive \(^{33}\) a string scale mass of order

$$m \sim \frac{1}{2\sqrt{n_1^2 R_1^2 + m^2 R_2^2}}$$  \hspace{1cm} (5.20)$$

where $(n,m)$ the wrapping numbers of the tori for which the SS deformation is acting.

• **Neutralinos**: Only Higgsino $\beta_1 = \beta_2 = 1$

In the PS models of section 3 and the left right symmetric models of section 4, all gauginos may receive masses from SS breaking due to their odd-wrappings. Thus the neutralino could be only made of Higgsinos. Also charginos are made only from only charged Higgsinos.

• **Neutralinos: only Higgsino-Wino**: $\beta_1 = 1/2; \beta_2 = 1$

\(^{32}\)we denote them by capital $X^{\text{neu}}$ to avoid confusion with $\chi$ which is used to describe the complex structure.

\(^{33}\)Non-chiral matter also receives a similar mass.
In this case, the second torus is tilted and the b-gauginos are massless as they receive no mass from SS deformations while we have present two pairs of tree level massless Higgsinos \((H_u, H_d)\). Thus at low energy neutralinos are made only from neutral components of higgsinos and "electroweak" Winos. Mixing of the higgsinos with the Winos perhaps is necessary in order to break the degeneracy of the two Dirac type lightest Higgsino newtralinos and to scatter them inelastically off the nucleus \([75]\). Also charginos which in general is a mixture of the charged components of higgsinos and charged gauginos could be also made from charged higgsinos and Winos. Binos may be massive as their gauginos are associated with the ones coming from the adjoint breaking of the the U(4) Pati-Salam branes that create the left-right symmetric models.

The introduction of the tilt in the second torus reintroduces the massless non-chiral fermions, as the SS deformation no more affects their masses. Thus massless adjoint fermions from the weak brane are reintroduced and at this level we have no suggestion how to give them a mass. However, the non-chiral non-adjoint matter from the intersections \(bh_1, bh_3, bh_5, ch_1, ch_4, bh_5\); where we denote the associated fermions as \(\chi^{(bh_1)}, \chi^{(bh_3)}, \chi^{(bh_5)}, \chi^{(ch_1)}, \chi^{(ch_4)}, \chi^{(bh_5)}\) are getting massive from the superpotential couplings \(^{34}\)

\[
\chi^{(bh_1)}\chi^{(bh_3)}\chi^{(h_1h_3)}, \chi^{(bh_5)}\chi^{(ch_1)}\chi^{(bc)}\chi^{(h_4h_5)}, \chi^{(ch_4)}\chi^{(ch_5)}\chi^{(h_4h_5)};
\]

where \(\chi^{(h_1h_3)}, \chi^{(h_4h_5)}, \chi^{(h_4h_5)}\), the N=2 matter that possess gauge singlet directions under the SM hypercharge. In the last term in (5.21), we have denote by underscript the corresponding charges under the gauge groups of the \((c; h_4, h_5)\) D6-branes respectively. By \(\chi^{(bc)}\) we denote the difundamental Higgs multiplets appearing as a basic ingredient in the spectrum of any PS or left-right symmetric model. We also mention of the possibility for the gravitino to be the LSP as it happens in local theories of gauge mediation with low scale susy breaking. See also \([47]\) for some relevant work in a split susy non-string context. Further consequences of the present models for dark matter will be considered elsewhere.

6 Addition of RR, NS and Metric fluxes

In this section, we will examine the classes of Pati-Salam models of this work, by studing their low energy effective action in the presence of background fluxes. In particular we study the question of moduli stabilization in the presence of RR/NS and also metric fluxes.

\(^{34}\)Superscripts denote the corresponding intersection.
as they have been recently incorporated in [43] and further studied in [44]. In type IIA Calabi-Yau orientifolds such as the $T^6$ torus the introduction of fluxes makes possible to fix most of the moduli of the theory. Given the unique string Pati-Salam models of the previous sections that leave only the MSSM at the string scale the question is if these models can be also reproduced in the context of fluxes where also moduli could be fixed.

6.1 Preliminaries - RR tadpoles

The general superpotential involving RR/NS and metric fluxes is the sum of the terms generating the Kähler and the complex structure moduli, $W_K, W_c$ respectively. We define

$$W_K = \int_Y e^{J_c} \wedge \tilde{F}_{RR}$$

where $\tilde{F}_{RR}$ represents a sum of even RR fluxes. The complex structure moduli $W_c$ is a sum of two terms, the first one that corresponds to the one generating a flux due to the NS $\mathcal{H}_3$ and the second one that is generated by the metric fluxes inside $dJ_c$; $W_c$ was computed in [43] as

$$W_c = \int_Y \Omega_c \wedge (\mathcal{H}_3 + dJ_c) = - \sum_{I,J=0}^{3} A_{IJ} \tilde{T}_I \tilde{U}_J$$

In the case of our models that are accommodated in the presence of RR/NS and metric fluxes – within four dimensional toroidal compactifications of type IIA theory – (and intersecting D6-branes) the full superpotential is computed to be [44]

$$W = W_c + W_Q = e_0 + ih_0 S + \sum_{i=1}^{3} [(ie_i - a_i S - b_{ii} U_i - \sum_{j \neq i} b_{ij} U_j)T_i - ih_i U_i]$$

$$- q_1 T_2 T_3 - q_2 T_1 T_3 - q_3 T_1 T_2 + imT_1 T_2 T_3.$$  

This is a result also obtained in [43]. Recall that the metric fluxes are constrained by the Jacobi identities which imply the twelve constraints

$$b_{ij} a_j + b_{jj} a_i = 0 ; \quad i \neq j$$

$$\tilde{T}_I = (i, T_1, T_2, T_3) ; \quad A_{IJ} = \begin{pmatrix} -h_0 & h_1 & h_2 & h_3 \\ a_1 & b_{11} & b_{12} & b_{13} \\ a_2 & b_{21} & b_{22} & b_{23} \\ a_3 & b_{31} & b_{32} & b_{33} \end{pmatrix}$$

$$\tilde{U}_I = (S, U_1, U_2, U_3).$$

35Detail definitions can be found e.g. in [44].
\[ b_{ik}b_{kj} + b_{kk}b_{ij} = 0 \quad ; \quad i \neq j \neq k. \] (6.5)

The RR tadpole conditions become

\[
\begin{align*}
\sum_a N_a n_a^1 n_a^2 n_a^3 + \frac{1}{2}(h_0 m + \sum_{i=1}^{3} a_i q_i) &= 16, \\
\sum_a N_a n_a^1 m_a^2 m_a^3 + \frac{1}{2}(h_1 m - \sum_{i=1}^{3} q_i b_{i1}) &= 0, \\
\sum_a N_a m_a^1 n_a^2 m_a^3 + \frac{1}{2}(h_2 m - \sum_{i=1}^{3} q_i b_{i2}) &= 0, \\
\sum_a N_a m_a^1 m_a^2 n_a^3 + \frac{1}{2}(h_3 m - \sum_{i=1}^{3} q_i b_{i3}) &= 0. \quad (6.6)
\end{align*}
\]

Hence for instance, a viable solution is \( b_{ji} = b_i, \ b_{ii} = -b_i, \ a_i = a. \) Further choosing RR fluxes

\[ q_i = -c_2, \quad e_i = c_1 \] (6.7)

we allow a configuration with \( T_1 = T_2 = T_3 = T. \) Then the superpotential (6.4) reduces to

\[ W = e_0 + 3ic_1 T + 3c_2 T^2 + imT^3 + ih_0 S - 3aST - \sum_{k=0}^{3}(ih_k + b_k T) U_k. \] (6.8)

If the fluxes \( h_k \) and \( b_k \) are independent of \( k, \) we can also set \( U_1 = U_2 = U_3 = U. \) Given the fluxes leading to (6.8), the tadpole conditions (6.6) become [43, 44]

\[
\begin{align*}
\sum_a N_a n_a^1 n_a^2 n_a^3 + \frac{1}{2}(h_0 m - 3ac_2) &= 16, \\
\sum_a N_a n_a^1 m_a^2 m_a^3 + \frac{1}{2}(h_1 m + b_1 c_2) &= 0, \\
\sum_a N_a m_a^1 n_a^2 m_a^3 + \frac{1}{2}(h_2 m + b_2 c_2) &= 0, \\
\sum_a N_a m_a^1 m_a^2 n_a^3 + \frac{1}{2}(h_3 m + b_3 c_2) &= 0. \quad (6.9)
\end{align*}
\]

We will examine the case of AdS vacua with negative cosmological constant in the non-supersymmetric models considered in the previous sections. We recall that non-susy intersecting models without fluxes [14, 15, 16, 17] are unstable due to uncanceled NSNS tadpoles. However some of the tadpoles vanish as all the complex structure moduli could get fixed by the supersymmetry conditions present in the Pati-Salam models of [15].

The important observation to make is that the Pati-Salam models considered in the previous section can be also be considered as possible solutions to models that accommodate R/NS and metric fluxes. A similar observation has already been made in the
appendix of [44], where it was shown that the N=0 models of [14](that had no explicit susy on intersections) could stabilize some moduli in the presence of metric fluxes.

We consider the case of isotropic Kähler moduli by choosing $T_k = T$ and look for minima of the effective supergravity potential with

$$D_{U_k} W = 0 \ ,\ D_{S} W = 0$$

(6.10)

where $W$ given in (6.8) and as usual $D_X = \partial_X W + W \partial_K W$. The above conditions imply the following constraint [44]

$$3as = b_k u_k .$$

(6.11)

where $a$ and $b_k$ must be both non-zero and of the same sign. Moreover, the following consistency conditions also hold

$$3h_k a + h_0 b_k = 0 \ ; \ k = 1, 2, 3 .$$

(6.12)

Before considering the issue of moduli stabilization let us mention one further constraint that should be imposed on the models. In general U(1) fields coupled to RR fields give rise to massive U(1)'s that are generated by the couplings

$$F^a \ &\ N_a \ \sum_{\ i = 0}^{3} \ c_i^a C^{(2)}_I$$

(6.13)

with

$$c_0^a = m_1^a m_2^a m_3^a ; \ c_1^a = m_1^a n_2^a m_3^a ; \ c_2^a = m_1^a m_2^a n_3^a ; \ c_3^a = n_1^a n_2^a m_3^a .$$

(6.14)

In addition, metric backgrounds give rise - to avoid inconsistencies resulting from loss of gauge invariance - to the constraint [44]

$$\int_{\Pi_a} (\overline{H}_3 + \omega J_c) = 0 ,$$

(6.15)

evaluated at the vacuum under examination, where $\Pi_a$ denotes the 3-cycle wrapped by the D6-brane, $\omega$ are the metric fluxes and $J_c$ is the complexified Kähler 2-form of the torus. Assuming a weaker form of this constraint that neglects metric fluxes, one gets the weaker condition [44]

$$\sum_{t = 0}^{3} c_t^a h_I = 0 .$$

(6.16)

The above condition guarantee that the U(1)'s that are getting massive from their couplings to RR fields are orthogonal to those one's becoming massive from NS fluxes.

$^{37}$In general from NS and metric, see (6.15).
6.2 Moduli stabilization from fluxes

Conditions (6.11), (6.12), (6.16) may be used to examine the models of the previous sections in the presence of general flux backgrounds. Imposing the condition (6.16) - on the general solution to the RR tadpoles (2.4) seen in table (13) of appendix B - we derive for the left-right symmetric (the result is also valid for the Pati-Salam models and the MSSM coming from Wilson line breaking of the L-R)

$$h_2 = \tilde{\epsilon} \frac{\beta_2}{\beta_1} h_3$$

(6.17)

Observe the similarity of this condition to the susy condition (2.18). In the present models that possess N=1 supersymmetry in the observable MSSM and a different $N' = 1$ supersymmetry in the supersymmetry breaking messenger sector, the hypercharge does survive massless the GS mechanism. It is given e.g. in the simplest left-right models as we have seen by eqn. (4.3) or by (2.16) when the breaking to the MSSM is achieved. Notice that in the non-supersymmetric models with intersecting D6-branes like [14] in the case that the RR tadpole parameters allow the hypercharge to survive massless to low energies [and thus having the SM chiral spectra at low energies] only the dilaton S may be fixed by the fluxes as $h_1 = h_2 = h_3 = 0$ and $h_0 \neq 0$. Similar results are expected for the non-susy models of [16, 17] which are the maximal SM generalizations of [14] to five and six stacks of intersecting D6-branes. The difference with [14], is that in the latter models [16, 17] we have in addition to the two extra U(1)’s which survive massless to low energies, one or two extra U(1)’s which are expected to get broken by the vev of the sneutrino.

On the contrary in the present models - where N=1 susy makes its presence manifest - all real parts of moduli, S, T, U, namely s, t, $u_k$ are expected to be fixed and also some linear axion combinations will get fixed.

From conditions (6.6), (6.17), (2.18) we observe that

$$\frac{\text{Re}U_3}{\text{Re}U_2} = \frac{h_2}{h_3} = \frac{\tilde{\epsilon} \beta_2}{\epsilon \beta_1}$$

(6.18)

That means that the supersymmetry conditions in the present models could fix the NS flux coefficients in the RR tadpoles.

$^{38}$For the 5-stack SM [16] there are one of two U(1)’s; for the 6-stack SM [17] there are only two U(1)’s. $^{39}$The intersections that accommodate the right handed neutrino in [16, 17] respect N=1 supersymmetry in sharp contrast with the [14] models where the intersections and hence the chiral fermions preserve no susy at all.
Hence if $\beta_1 = \beta_2 = 1$, $h_2 = h_3$. From (6.18) we also conclude \(^{40}\), as we examined the case that there is no tilt in any of our tori, that

$$ReU_2 = ReU_3$$

(6.19)

ans since our torus is orthogonal,

$$ImU_1 = ImU_2 = ImU_3 = 0$$

(6.20)

We are choosing a vacuum with $e_0 = c_1 = 0$, $h_0 = -3a$, $c_2 = -m$,

$$\epsilon = \bar{\epsilon} = \beta_1 = \beta_2 = 1$$

(6.21)

and which satisfies the consistency conditions (6.12)

$$b_k = (6, 6, 6) \quad , \quad h_k = -\frac{h_0}{3a} b_k \quad , \quad k = 1, 2, 3$$

(6.22)

and where also $h_0 < 0$, $a < 0$, $m < 0$. The flux contributions to the RR tadpoles cancel as we have chosen

$$h_0m = 3ac_2, \quad h_im + b_ic_2 = 0$$

(6.23)

e.g. the contribution of metric fluxes cancels against the contribution from R/NS fluxes.

As this is a general result \(^{44}\) it appears that: One can choose any four dimensional non supersymmetric vacuum on toroidal orientifolds which satisfies RR tadpoles in the context of compactifications of type IIA theory with intersecting D6-branes and make it to satisfy the RR tadpoles with fluxes by making the choice (6.23).

The previous observation extends easily to the N=1 supersymmetric orientifold vacua based on $Z_N$, $Z_N \times Z_M$ orbifolds as one has to add the contribution of the corresponding crosscap contributions to the RR tadpoles; one has to add a contribution of $-16$ to the right hand side of the last three lines of (6.9).

Let us here discuss some points regarding the fate of U(1)'s present. In general the U(1) combination \(^{44}\)

$$Im \ U_2 - \frac{\bar{\epsilon}_\beta_2}{\epsilon_\beta_1} \ Im \ U_3$$

(6.24)

gets massive from its coupling to RR BF fields (4.1). Its orthogonal combination

$$\hat{X} \equiv Im \ U_2 + \frac{\epsilon_\beta_1}{\bar{\epsilon}_\beta_2} \ Im \ U_3 \ = \ \frac{1}{h_2} (h_2 \ Im \ U_2 + h_3 \ Im \ U_3)$$

(6.25)

\(^{40}\) For a two dimensional torus where the generating lattice is spanned by the basis vectors $\vec{e}_i = G_{ij} e^j$; the complex structure moduli is given by $U \overset{def}{=} G_{11} (\sqrt{G} - iG_{12})$. 

\(^{44}\) For the supersymmetric vacuum, one must choose $\epsilon_0 = c_1 = 0$, $h_0 = -3a$, $c_2 = -m$, $\epsilon = \bar{\epsilon} = \beta_1 = \beta_2$. This choice cancels all the flux contributions to the RR tadpoles with fluxes.
is expected to receive a mass from fluxes upon minimization as it should appear in the superpotential $W$. Also massive becomes the combination of complex structure axions defined as \[ 3a \text{Im} S + \sum_k b_k \text{Im} U_k = 3c_1 + \frac{3c_2}{a}(3h_o - 7av) - \frac{3m}{v}(3h_o - 8av), \quad (6.26) \]

where $v = \text{Im} T$ takes different values \(^{41}\) according to whether $m = 0$, or $m \neq 0$. For the vacua associated to (6.20), this combination becomes $3a \text{Im} S + b_1 \text{Im} U_1$. Thus given that $\text{Im} U_i$ are fixed from (6.20), the $\text{Im} S$ is fixed in (6.26). In general - neglecting NS tadpole contributions - the real part of the $S, T^i, U^i$ moduli denoted as $s, u_k, t$ respectively get also fixed \(^{42}\) in the present models at

\[
s = \frac{m}{a 10^{1/3}} \frac{-h_0}{3m} t, \quad u_k = \frac{3s a}{b_k}, \quad t = \frac{\sqrt{1510^{2/3}}}{20} \frac{-h_0}{3m} \quad (6.27)
\]

leaving unfixed only the imaginary parts of the Kähler moduli. The latter could be fixed when the NS potential is included. For some other proposals to fix the Kähler moduli see \([13, ?]\). In order to trust the supergravity approximation that we are using it is necessary that $-h_0/3a$ large since the 4-D dilaton $e^{\phi} = (su_1 u_2 u_3)^{1/4}$.

We also note that that as the present models have uncanceled NSNS tadpoles, the exact stabilization problem is solved by considering the minimization of the full potential that includes the NSNS potential in addition to the potential generated by the superpotential fluxes. For simplicity reasons we will not consider this option here leaving the full treatment in future work.

### 7 Conclusions

We have presented the first three generation D-brane string models that localize the MSSM spectrum at the string scale, simultaneously providing us with a particle physics model that accommodates all the long seeking properties of a realistic MMSM incorporating string vacuum, that possess:

- Unification at $2 \cdot 10^{16}$ GeV (of SU(3) and SU(2) gauge couplings).
- All extra chiral and non-chiral beyond the MSSM matter is getting massive at the string scale.

\(^{41}\)In the following we assume $m \neq 0$.

\(^{42}\)Borrowing the results from \([44]\).
• Proton stability, as baryon number is a gauged symmetry.

• The MSSM spectrum is massless at the string scale while the theory breaks to the SM $SU(3)_c \times SU(2) \times U(1)_Y$ at low energies.

The great advantage of using intersecting D-branes against other approaches involving compactifications of the perturbative string theories is that in the present formalism we can guarantee proton stability, moduli fixing (through supersymmetry conditions at intersections) and exotics becoming massive occurring simultaneously.

A comment is in order. Most of the non-chiral states become massive by the introduction of Scherk-Schwarz breaking. For the rest of them, including also the chiral exotics of the vector-like messenger sector, trilinear superpotential couplings (TSC) are employed, where it is assumed that the coupled gauge singlet directions do receive a vev. The current models provide us with a laboratory that can be used to investigate in the future the precise determination of these vevs. We note that the presence of these TSC – in the presence of a stable proton – for the extra beyond the MSSM matter multiplets is absent in all the models from intersecting branes or other recent model building constructions from string theory [1, 2, 8, 9, 21, 22, 66, 57, 59, 40, 80].

In toroidal orientifold vacua with intersecting D6-branes, supersymmetry is already broken by construction as the absence of a variety of orientifold planes does not allow N=1 models but only N=0 models to be constructed. Thus the vacua presented in this work, even though they localize the MSSM multiplet spectrum among the observable sector D6-branes $O_i$, are non-supersymmetric (N=0); the N=1 supersymmetry of the extra messenger sector is different from the one respected by the MSSM multiplets.

The most important future of the present string models is that they manage to make all the extra chiral/non-chiral beyond the MSSM exotics massive, due to the combined use of the introduction of Scherk-Schwarz deformations and flat directions.

These models have the novel property that only the supersymmetric intersections share a susy with the orientifold plane and not the participating intersecting branes, meaning that some of the branes involved can be non-supersymmetric.

The Pati-Salam/left-right symmetric models of this work are the first realistic examples of gauge mediation in string theory. They provide us with a stringy platform for realistic calculations in the context of string theory. Attempts to construct gauge mediated semirealistic models in string theory could be found in [72, 73], while local models of gauge mediation inspired from string theory could be found in [77]. In our models we don’t have to show dynamical supersymmetry breaking in the standard sense as supersymmetry
is broken by construction. The great advantage of gauge mediation against other schemes of supersymmetry breaking — like gravity mediation or anomaly mediation — are the universal squark and slepton (soft term) masses. In the present constructions masses for sparticles which are due to the D-term breaking and originate from FI terms in the limit of small deviations of complex structure result in degenerate squarks and slepton masses and thus may be responsible for automatic flavor changing neutral current suppression.

We also explained the "embedding" of the present models in toroidal orientifold backgrounds of - type IIA 4D compactifications - in the presence of NS/R and metric fluxes in the spirit of [42], [44] (and [43]). The stringy toroidal orientifolds of sections 2,3,4,5 that in the absence of fluxes have NS tadpoles remaining, now become stable, as the real and imaginary parts of dilaton and complex structure moduli and also the real parts of Kähler moduli get fixed. The only remaining unfixed moduli, the imaginary parts of Kähler ones, could be fixed by the complete minimization of the NS scalar potential

\[ V = \frac{T_6}{\lambda} \sum_a (||l_a|| - ||l_o||), \]  

where \( T_6 \) the tension of the D6-branes; \( \lambda \) the string coupling; \( ||l_a||, ||l_o|| \) the \( T^6 \) volume that the D-branes and the O6-planes lie respectively. The uplifting of Ads vacua to de Sitter ones is an interesting possibility once the NS potential is introduced, without the need of any non-perturbative effects [78], and it will be examined elsewhere.

An important comment is in order. When the non-supersymmetric models of [14] were examined in the presence of RR/NS and metric fluxes [44] - in these models there is no supersymmetry present at any intersection - when the hypercharge survives massless to low energies [16, 17], only the dilaton could be fixed by fluxes. On the contrary in the present non-supersymmetric models where \( N=1, N=2 \) (Higgses/higgsinos) supersymmetry makes explicit its appearance at intersections, the hypercharge survives massless and both the dilaton and the complex structure moduli could be fixed. Also in the present constructions more moduli are fixed than in the usual \( N=1 \) Ads vacua of [44] - where only the real parts of the dilaton and complex and Kähler moduli get fixed - as the toroidal lattice fixes also the imaginary parts of the complex structure moduli.

It will be also interesting to examine moduli stabilization in the present models in the spirit of [76]. We note that when passing from the stringy MSSM models of sections 3, 4, 5 to the fluxed models of section 6, the only feature of the string models that is not surviving in the presence of fluxes is the non-chiral matter that is getting massive by the introduction of SS breaking.
We also note that the universal value obtained for the sparticle masses coming from D-term breaking depend on the small variation of the complex structure and thus on the tori radii. For particular values of these radii, the scale of sparticle masses could be lowered well below the string GUT/unification scale of $2.04 \times 10^{16}$ GeV where the $SU(3), SU(2)$ gauge couplings meet.

Our work also confirms that proposals based on the idea of statistical analysis [84] of the existence of a string landscape [83, 84] which implicitly accepts the existence of a stringy Standard model vacuum with no exotics are definitely real. Further studies of the present models including studies of susy soft breaking terms and low energy squark/slepton phenomenology, quark and lepton mass hierarchies [67, 71] may be examined elsewhere.

Acknowledgments

We would like to thank I. Antoniadis, C. Angelantonj, for discussions; C.K would like to thank P. Camara, G. Giudice, S. Dimopoulos for a discussion. C.K would like also to thank the Organizers of the String Phenomenology 2005, SUSY 2005, HEP2006 conferences where the results of this work were announced and in addition the Cern theory division for their warm hospitality where parts of this work were completed. E.G.F and C.K are supported by the research program Pithagoras II, grant no: 70-03-7992 of the Greek Ministry of National Education that is also funded by the European Union.

8 Appendix A -Alternative Hypercharge embedding for the L-R symmetric models

There is alternative hypercharge embedding for the left-right symmetric models discussed in section (3.2) that gives the same spectrum and also the same hypercharge for the observable MSSM related part of the spectrum in table (10). It is defined as

$$U(1)_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_d - \frac{1}{2}Q_{h_3} - \frac{1}{2}Q_{h_4}$$  \hspace{1cm} (8.1)

Notice that the hypercharge assignment of the messenger multiplets are different than the one appeared assigned to the choice of eqn. (4.3) on table (10). The RR couplings reveal that the U(1) defined as $3Q_a + Q_d$, gets massive while the extra U(1)'s $U(1)^{(1*)} = Q_{h_3} - Q_{h_4}$, $U(1)^{(2*)} = (6Q_a - 18Q_d) + 10(Q_{h_3} + Q_{h_4})$ get broken if $\tilde{\phi}_R, \tilde{\phi}_L$ get vevs respectively as they are flat directions of the potential. Also non-chiral matter gets massive by SS deformations
as in section (3.2). However the couplings that are available for the messenger multiplets within the new assignment (8.1) are exactly the same that appear in eqns (4.7). As a result we obtain for the different hypercharge embedding (8.1) the same physics; same unification scale, all exotics massive, etc.

9 Appendix B

We this appendix we provide the solution to the RR tadpoles for the Pati-Salam models of section 3, in the case that

$$\epsilon = \tilde{\epsilon}, \ \rho = 1, 1/3.$$ (9.1)

The case $\rho = 1/3$ was treated in table 6 and is used to facilitate the discussion in the main text of this work. These solutions appear in table (13).

References

[1] M. Cvetic, G. Shiu, A. M. Uranga, Nucl. Phys. B615 (2001) 3, hep-th/0107166; M. Cvetic, G. Shiu, A. M. Uranga, Phys. Rev. Lett. 87 (2001) 20, hep-th/0107143

[2] M. Cvetic, T. Li, T. Liu, Nucl. Phys. B698 (2004) 163, hep-th/0403061; M. Cvetic, P. Langacker, T. Li, T. Liu, Nucl. Phys. B709 (2005) 241, hep-th/0407178

[3] R. Blumenhagen, B. Kors, D. Lust, *Type I Strings with F- and B-Flux*, JHEP 0102 (2001) 030 hep-th/0012156.

[4] R. Blumenhagen, M. Cvetic, P. Langacker, G. Shiu, Ann.Rev.Nucl.Part.Sci. 55 (2005) 71, hep-th/0502005.

[5] S.Abel, A.W.Owen, Nucl.Phys. B663 (2003)197 hep-th/0303124; M.Cvetic and I.Papadimitriou, Phys.Rev.D68:046001,2003, Erratum-ibid.D70 (2004) 029903, hep-th/0303083

[6] D.Lust, Class.Quant.Grav.21:S1399 (2004), hep-th/0401156

[7] C. Kokorelis, *Standard Model Building from Intersecting D-Branes*, Appeared in the series *New developments in String Theory Research*, Nova Publishers, NY , 2005, hep-th/0410134
Table 13: Solution to the RR tadpoles for toroidal orientifold models. The N=1 MSSM chiral spectrum arises as part of Pati-Salam models $SU(4) \times Sp(2)_b \times Sp(2)_c$ in the top part of the table from intersections between a, b, c, d branes. Messenger multiplet states respecting a $N' = 1$ supersymmetry arise from the intersections of the a, d branes with the $h_3, h_4$ branes. We have set $\epsilon = \tilde{\epsilon}$ in table (3); the parameter $\rho$ can take the values 1, 1/3.

| $N_n$ | $(n^1_i, m^1_i)$ | $(n^2_i, m^2_i)$ | $(n^3_i, m^3_i)$ | G.G. - $\epsilon = +1$ |
|-------|-----------------|-----------------|-----------------|------------------|
| $Na = 4$ | (1, 0) | (1/3, 3$\rho\epsilon\beta_1$) | (1/3, -3$\rho\epsilon\beta_2$) | $U(4)$ |
| $Nh_1 = 1$ | (0, 1) | (1/3, 0) | (0, $-\epsilon$) | $Sp(2)$ |
| $Nh_2 = 1$ | (0, $\epsilon$) | (0, $-\epsilon$) | ($\epsilon/\beta_2$, 0) | $Sp(2)$ |
| $Nh_3 = 4$ | (1, 0) | (-1/3, 0) | (1/3, 0) | $U(1)^{Nh_3}$ |
| $Nh_4 = 36$$\rho^2\beta_1\beta_2$ | (1, 0) | (0, $\epsilon$) | (0, $\epsilon$) | $U(1)^{Nh_4}$ |
| $Nh_5 = 1$ | (0, 1) | (1/3, 0) | (0, $\epsilon$) | $U(1)$ |
| $Nh_6 = 1$ | (0, $\epsilon$) | (0, $\epsilon$) | ($\epsilon/\beta_2$, 0) | $U(1)$ |
| $Nh_7 = 16\beta_1\beta_2$ | (1, 0) | (1/3, 0) | (1/3, 0) | $Sp(2)^{Nh_7}$ |

[8] R. Blumenhagen, L. Görlich, T. Ott, JHEP 0301 (2003) 021, hep-th/0211059; R. Blumenhagen, J. P. Conlon, K. Suruliz, JHEP 0407:022, 2004, hep-th/0404254

[9] V. Braun, Y.-H. He, B. A. Ovrut, T. Pantev, hep-th/0512177

[10] T.L. Gomez, S. Lukic, I. Sols, Constraining the Kahler moduli in the Heterotic Standard Model, hep-th/0512205

[11] C. Angelantonj and A. Sagnotti, arXiv:hep-th/0010279

[12] M. Larosa and G. Pradisi, Nucl. Phys. B667:261, 2003, hep-th/0305224
[13] M. Bianchi and E. Trevigne, JHEP 0508, 034 (2005), hep-th/0502147

[14] L. E. Ibáñez, F. Marchesano and R. Rabádán, Getting just the standard model at intersecting branes, JHEP, 0111 (2001) 002, hep-th/0105155

[15] C. Kokorelis, JHEP 08 (2002) 018, hep-th/0203187; GUT model Hierarchies from Intersecting Branes, C. Kokorelis, hep-th/0210004; C. Kokorelis, JHEP 0211 (2002) 027, hep-th/0209202; C. Kokorelis, Deformed Intersecting D6-Branes II, hep-th/0210200

[16] C. Kokorelis, New Standard Model Vacua from Intersecting Branes, JHEP 09 (2002) 029, hep-th/0205147

[17] C. Kokorelis, Exact Standard Model Compactifications from Intersecting Branes, JHEP 08 (2002) 036, hep-th/0206108

[18] R. Blumenhagen, B. Kors, D. Lust, T. Ott, Nucl. Phys. B616 (2001) 3, hep-th/0107138

[19] D. Cremades, L. Ibanez, F. Marchesano, JHEP 0207 (2002) 009, hep-th/0201205

[20] D. Cremades, L. Ibanez, F. Marchesano, Computing Yukawa couplings from magnetized extra dimensions, JHEP 0405 (2004) 079, hep-th/0404229.

[21] G. Honecker and T. Ott, Phys. Rev. D70:126010, 2004, Erratum-ibid. D71:069902, 2005, hep-th/0404055

[22] T. P. T. Dijkstra, L. R. Huiszoon and A. N. Schellekens, hep-th/0403196

[23] M. Bianchi, G. Pradisi and A. Sagnotti, Nucl. Phys. B376 (1992) 365; C. Angelantonj, Nucl. Phys. B566 (2000) 126; Z. Kakushadze, Int. J. Mod. Phys. A15 (2000) 3113

[24] D. Cremades, L. E. Ibanez, F. Marchesano, Towards a Theory of Quark Masses, Mixings And CP Violation, hep-ph/0212064

[25] Cremades, Ibanez, Marchesano, JHEP 0307 (2003) 038, hep-th/0302105

[26] C. Kokorelis, N=1 Locally Supersymmetric Standard Models from Intersecting Branes, hep-th/0309070, Revised version to appear

[27] D. Cremades, L. E. Ibanez, F. Marchesano, JHEP 0207 (2002) 022, hep-th/0203160
[28] I. Klebanov and E. Witten, Nucl. Phys. B664 (2003) 3, hep-th/0304079

[29] M. Axenides, E. Floratos and C. Kokorelis, JHEP0310 (2003) 006, hep-th/0307255

[30] P. Burikham, JHEP 0502 (2005) 030, hep-ph/0502102

[31] P. Nath, P. Fileviez Perez, hep-ph/0601023; R. Tatar, T. Watari, hep-th/0602238

[32] M. Cvetic, R. Richter, hep-th/0606001

[33] D. Berenstein, hep-th/0603103

[34] W. Buchmuller, K. Hamaguchi, O. Lebedev, M. Ratz, hep-th/0606187

[35] S. Dimopoulos, S. Raby, Nucl. Phys. B192 (1982) 353; M. Dine, W. Fischler, M. Srednicki, Nucl. Phys. B189 (1981) 575; M. Dine, W. Fischler, Phys. Lett. B110 (1982) 227; Nucl. Phys. B204 (1982) 346; C. Nappi, B. Ovrut, Phys. Lett. B113 (1982) 175; L. Alvarez-Gaume, M. Claudson and B. Wise, Nucl. Phys. B207 (1992) 96; S. Dimopoulos, S. Raby, Nucl. Phys. B219 (1993) 479; M. Dine, A. E. Nelson, Phys. Rev. D48 (1993) 1277; D51 (1995) 1362; M. Dine, A. E. Nelson, Y. Shirman, Phys. Rev. D51 (1995) 1362; M. Dine and A. E. Nelson, Y. Nir, Y. Shirman, Phys. Rev. D53 (1996) 2658; G. F. Guidice and R. Ratazzi, Phys. Rept. 322 (1999) 419; S. L. Dubovsky, D. S. Gorbunov, S. V. Troitsky, Phys. Usp. 42 (1999) 623; S. Dimopoulos, S. Thomas, J. D. Wells, Nucl. Phys. B488 (1997) 39

[36] Blumenhagen, Lust, Stieberger, ”Gauge Unification in supersymmetric intersecting brane worlds”, JHEP 0307 (2003) 036, hep-th/0305146

[37] G. L. Kane, P. Kumar, J. D. Lykken, T. T. Wang, Phys.Rev.D71 (2005) 115017, hep-ph/0411125

[38] R. Blumenhagen, D. Lust, T. Taylor, Nucl.Phys. B663 (2003) 319, hep-th/0303016

[39] J. F. G. Cascales, A. M. Uranga, JHEP 0305 (2003) 011, hep-th/0303024; hep-th/0311250

[40] F. Marchesano, G. Shiu, ”Building MSSM flux vacua”, JHEP 0411 (2004) 041, hep-th/0409132

[41] O. DeWolfe, A. Giryavets, S. Kachru and W. Taylor, Type IIA moduli stabilization, hep-th/0505160.
[42] J. P. Derendinger, C. Kounnas, P. M. Petropoulos and F. Zwirner, Nucl. Phys. B715 (2005) 211, hep-th/0411276; hep-th/0503229.

[43] G. Villadoro, F. Zwirner, JHEP 0506 (2005) 047, hep-th/0503169

[44] P.G. Cámara, A. Font, L.E. Ibáñez, JHEP 0509 (2005) 013, hep-th/0506066

[45] I. Antoniadis and T. Taylor, Nucl. Phys. B695 (2004)103, hep-th/0403293

[46] N.Arkani-Hamed and S. Dimopoulos, JHEP 0506:073,2005, hep-th/0405159

[47] N. Arkani-Hamed, S. Dimopoulos, G.F. Giudice, A. Romanino, ”Aspects of Split Supersymmetry”, Nucl.Phys. B709 (2005) 3, hep-ph/0409232

[48] I. Antoniadis and S. Dimopoulos, Split supersymmetry in string theory, Nucl.Phys.B715 (2005) 120, hep-th/0411032

[49] C.Kokorelis, Standard Models and Split Supersymmetry from Intersecting Brane Orbifolds, hep-th/0406258

[50] C.Kokorelis, Standard Model Compactifications of IIA Z3 x Z3 Orientifolds from Intersecting D6-branes, Nucl.Phys. B732 (2006) 341, hep-th/0412035

[51] B. Kors and P. Nath, hep-th/0411201

[52] N.Haba, N. Okada, hep-ph/0602013; S. Kumar Gupta, B. Mukhopadhyaya, S. Kumar Rai, hep-ph/0510306; E. Dudas, S. K. Vempati, hep-th/0506172; A.Ibarra, hep-ph/0503160; N.G. Deshpande, J. Jiang, hep-ph/0503116; B. Dutta, Y. Mimura, hep-ph/0503052; N. Haba, N. Okada, hep-ph/0502213]; G. Senjanovic, hep-ph/0501244; Z. Lalak, R. Matyszkiewicz, hep-ph/0506223; K.S. Babu, Ts. Enkhbat, B. Mukhopadhyaya, hep-ph/0501079; S. P. Martin, K. Tobe, J. D. Wells,hep-ph/0412424; P. Fileviez Perez, hep-ph/0412347; P. C. Schuster, hep-ph/0412263; A. Datta and X. Zhang, hep-ph/0412255; A. Masiero, S. Profumo, P. Ullio, hep-ph/0412058

[53] J.Pati and A.Salam, Phys.Rev.D10, 275 ,1974; see also J. Pati, hep-ph/0407220 and J.Pati, Probing Grand Unification Through Neutrino Oscillations, Leptogenesis, and Proton Decay, hep-ph/0305221

[54] C. Kokorelis, Exact Standard model Structures from Intersecting D5-branes, Nucl. Phys. B677:115, 2004, hep-th/0207234
[55] D. Cremades, L. E. Ibáñez and F. Marchesano, *Standard model at intersecting D5-Branes: lowering the string scale*, Nucl. Phys. B643 (2002) 93, hep-th/0205074

[56] A. Prikas, N.D. Tracas, hep-ph/0303258; D.V. Gioutsos, G.K. Leontaris, J. Rizos, Eur.Phys.J. C45:241(2006), hep-ph/0508120; D.V. Gioutsos, hep-ph/0605278

[57] G. Aldazabal, E. Andres, J. E. Juknevich, *On Susy Standard-like models from orbifolds of D=6 Gepner orientifolds*, hep-th/0603217

[58] R. Blumenhagen, G. Honecker, T. Weigand, hep-th/0510050

[59] P. Anastasopoulos, T.P.T. Dijkstra, E. Kiritsis, A.N. Schellekens, hep-th/0605226

[60] J. Kumar and J.D. Wells, JHEP 0509, 067 (2005), hep-th/0506252; hep-th/0604203; J.Kumar, hep-th/0601053

[61] C. Angelantonj, M. Cardella and N. Irges, *Scherk-Schwarz breaking and intersecting branes*, Nucl. Phys. B725 (2005) 115 hep-th/0503179.

[62] P. G. Cámara, L. E. Ibáñez, and A. M. Uranga, *Flux-induced SUSY-breaking soft terms*, Nucl. Phys. B689, 195 (2004), hep-th/0311241; M. Graña, T. W. Grimm, H. Jockers, and J. Louis, Nucl. Phys. B690 (2004) 21, hep-th/0312232; L. Görlich, S. Kachru, P. K. Tripathy and S. P. Trivedi, hep-th/0407130; D. Lüst, P. Mayr, S. Reffert and S. Stieberger, hep-th/0501139.

[63] E. Dudas, C. Timirgaziu, *Internal magnetic fields and supersymmetry in orientifolds*, Nucl. Phys. B716 (2005) 65, hep-th/0502085

[64] H. Verlinde and M. Wijnholt, *Building the SM on a D3-brane*, hep-th/0508089

[65] M. Cvetic, T. Li, T. Liu *Standard-like Models as Type IIB Flux Vacua*, Phys.Rev. D71 (2005) 106008, hep-th/0501041

[66] C-M. Chen, T. Li, D.V. Nanopoulos, Nucl. Phys. B732 (2006) 224, hep-th/0509059; C-M. Chen, V. E. Mayes, D.V. Nanopoulos, Phys.Lett. B633 (2006) 618, hep-th/0511135

[67] B. Dutta and Y. Mimura, *Properties of fermion mixing in intersecting D-brane models*, Phys.Lett. B633 (2006) 761, hep-ph/0512171; hep-ph/0604126

[68] F. Wang, W. Wang, J. Yang, hep-ph/0512133
[69] L. Ibanez, C. Munoz and S. Rigolin, Nucl.Phys. B422 (1994) 125, hep-ph/9308271

[70] K. S. Babu, B. Dutta, R. N. Mohapatra, "Up-Down Unification, Neutrino Masses and Rare Lepton Decays", Phys.Lett. B458 (1999) 93, hep-ph/9904366

[71] T. Higaki, N. Kitazawa, T. Kobayashi, K. Takahashi, Phys.Rev. D72 (2005) 086003, hep-th/0504019

[72] D. Diaconescu, B. Florea, S. Kachru, P. Svrcek, JHEP 0602 (2006) 020, hep-th/0512170.

[73] I.Antoniadis, K.Benakli, A.Delgado, M.Quiros, M.Tuckmantel, hep-th/0601003

[74] C. Munoz, Int. J. Mod. Phys. A19 (2004)3093 , hep-ph/0309346; G. Bertone, D. Hooper, J. Silk, Phys. Rept. 405(2004) 279, hep-ph/0404175

[75] D.R.Smith and N.Weiner, Phys. Rev. D64 (2001)043502, hep-ph/0101138

[76] G. Aldazabal, P.G. Camara, A. Font, L.E. Ibanez, "More Dual Fluxes and Moduli Fixing", hep-th/0602089

[77] D. Berenstein, C.P. Herzog, P. Ouyang, S. Pinansky, JHEP 0509:084,2005, hep-th/0505029; K. Intriligator, N. Seiberg, JHEP 0602:031,2006, hep-th/0512347; I. Garcia-Etxebarria, F. Saad, A. M. Uranga, hep-th/0605166

[78] S.Kachru, R.Kallosh, A.Linde, S.P.Trivedi, De Sitter vacua in string theory, hep-th/0301240

[79] I.Antoniadis, A.Kumar, T.Mailard, Moduli stabilization with open and closed string fluxes, hep-th/0505260;
F.Denef, M.Douglas, B.Florea, A.Grassi, S.Kachru, Fixing all moduli in a simple f-theory compactification, hep-th/0503124;
P.Bergrlund, P.Mayr, Non-perturbative superpotentials in F-theory and string duality, hep-th/0504058;
J.Conlon, F.Quevedo, K.Suruliz, Large-volume flux compactifications: Moduli spectrum and D3/D7 soft supersymmetry breaking, hep-th/0505076

[80] R. Blumenhagen, S. Moster, T. Weigand, hep-th/0603015

[81] N. Arkani-Hamed, S. Dimopoulos, S. Kachru Predictive Landscapes and New Physics at a TeV, hep-th/0501082
[82] M. R. Douglas, W. Taylor, *The landscape of intersecting brane models*, hep-th/0606109.

[83] L. Susskind, *The anthropic landscape of string theory*, hep-th/0302219

[84] M. R. Douglas, *The statistics of string M-theory vacua*, JHEP 0305, 046 (2003), hep-th/0303194; S. Ashok, M. R. Douglas, *Counting flux vacua*, JHEP 0401, 060 (2004), hep-th/0307049; F. Denef, M. R. Douglas, *Distributions of string vacua*, JHEP 0405, 072 (2004), hep-th/0404116