

Modulus-dominated SUSY-breaking soft terms in F-theory and their test at LHC

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ABSTRACT: We study the general patterns of SUSY-breaking soft terms arising under the assumption of Kahler moduli dominated SUSY-breaking in string theory models. Insisting that all MSSM gauginos get masses at leading order and that the top Yukawa coupling is of order the gauge coupling constant identifies the class of viable models. These are models in which the SM fields live either in the bulk or at the intersection of local sets of Type IIB $D7$ -branes or their F-theory relatives. General arguments allow us to compute the dependence of the Kahler metrics of MSSM fields on the local Kahler modulus of the brane configuration in the large moduli approximation. We illustrate this study in the case of toroidal/orbifold orientifolds but discuss how the findings generalize to the F-theory case which is more naturally compatible with coupling unification. Only three types of 7-brane configurations are possible, leading each of them to very constrained patterns of soft terms for the MSSM. We study their consistency with radiative electroweak symmetry breaking and other phenomenological constraints. We find that essentially only the configuration corresponding to intersecting 7-branes is compatible with all present experimental constraints and the desired abundance of neutralino dark matter. The obtained MSSM spectrum is very characteristic and could be tested at LHC. We also study the LHC reach for the discovery of this type of SUSY particle spectra.

KEYWORDS: Supersymmetry Breaking, Compactification and String Models, F-Theory, Intersecting branes models.

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1. Introduction

In the very near future LHC will start operating and a window for new physics will be open. Low energy supersymmetry is one of the most solid candidates for that new physics. If indeed low energy SUSY exists, a plethora of SUSY particles (gluinos, squarks, sleptons etc.) will be produced at LHC and their mass will be measured. Discovering SUSY at LHC would certainly be a revolution. On the other hand there is not at the moment a unique theory predicting the values of masses for squarks, gluinos, sleptons. . . There are plenty of proposals for implementing the breaking of SUSY giving rise to different patterns of SUSY

masses. In any event, if low energy SUSY is found, measurements (however imprecise) of, e.g. the value of squark and gluino masses would give very important information about the underlying theory.

In the present article we argue that a measurement of the SUSY spectrum of sparticles would also provide an experimental test for large classes of string compactifications giving rise to standard model physics.

In addressing these issues we will

1. Assume that the effective low-energy $N = 1$ supergravity approximation is valid.
2. Assume that the low-energy theory has the structure of the MSSM.
3. Assume that eventually all moduli are stabilized with a vacuum energy close to zero.
4. Assume that SUSY-breaking predominantly originates in the vacuum expectation values of auxiliary fields of the Kahler moduli T_i of the compactification.
5. Insist that all SM gauginos get a mass through the above mechanism at leading order.
6. Assume standard MSSM gauge coupling unification.
7. Assume there is at least one Yukawa coupling (that of the top-quark) of order g (the gauge coupling constant).

Under these reasonable assumptions one can identify a large class of string compactifications for which, in the leading large volume approximation, one can compute the structure of SUSY-breaking soft terms which in turn could be compared with data (if SUSY is actually found).

Let us briefly explain in more detail what underlies the above assumptions and what class of string compactifications they identify. The first two conditions are just the expectation of being able to reproduce the MSSM coupled to $N = 1$ supergravity in a perturbative regime (both in α' (string tension) and string coupling g_s) from some successful string theory compactification. The third assumption, that eventually all moduli gets dynamically fixed with an almost vanishing cosmological constant, has been the subject of much work in recent years [1, 2]. It has been shown that indeed the presence of Ramond-Ramond (RR) and Neveu-Schwarz (NS) antisymmetric fluxes in Type IIB orientifolds combined with non-perturbative effects may stabilize all string moduli at weak coupling. Specific examples have been worked out (for reviews and references see [2]).

The fourth assumption needs more explanation. It is well known that, within the mentioned approximations the Kahler moduli in string compactifications have a classical no-scale structure [3] in such a way that one can obtain SUSY breaking with a vanishing (tree level) cosmological constant. This is true both in the Heterotic as well as Type I, Type IIB orientifold and other compactifications and corresponds to having non-vanishing vacuum expectation values for the auxiliary fields of the Kahler moduli. In Type IIB Calabi-Yau (CY) orientifolds this corresponds to string compactifications with spontaneously broken SUSY solving the classical equations of motion [4]. It has also been shown [4–6] that in

those Type IIB orientifolds such SUSY-breaking corresponds to the presence of RR and NS antisymmetric tensor fluxes in the compactification. Such compactifications will *generically* have non-vanishing fluxes so it looks like a most natural source of supersymmetry breaking in string theory. The scale of SUSY breaking may be hierarchically small if the flux in the compact region where the SM sector resides is appropriately small.

The 5-th condition, asking for the presence of SUSY-breaking gaugino masses at leading order has a phenomenological motivation, as we will discuss in the main text. It turns out to be quite restrictive. Indeed, as we will argue, it excludes large classes of string compactifications including perturbative heterotic compactifications, Type I string compactifications and Type IIB orientifolds with the SM residing on $D3$ -branes at local singularities. In all these classes of compactifications the gauge kinetic function is independent of the (untwisted) Kahler moduli to leading order so that modulus dominance does not give rise to gaugino masses. This argumentation will leave us just with Type IIB orientifold compactifications with MSSM fields residing on $D7$ -branes (or their intersections) and their non-perturbative F-theory extensions. Alternatively one may consider the mirror description in terms of Type IIA orientifolds with the MSSM fields living at $D6$ -branes (or their intersections) or their non-perturbative extensions from M-theory compactified on manifolds with G_2 holonomy. Being in principle equivalent we will concentrate in the case of the MSSM chiral fields residing in the bulk of $D7$ -branes or at the intersection of $D7$ -branes (or F-theory 7-branes) since at present little is known about the effective action of M-theory compactifications in manifolds with G_2 holonomy (see [7] for work in this direction).

The unification of couplings is one of the most solid arguments in favour of the MSSM [8]. We will thus also impose that gauge couplings of the MSSM unify (6-th assumption). A simple way in which this is achieved is having at some level some $SU(5)$, $SO(10)$ or E_6 symmetry relating the couplings. This unification structure, as recently emphasized in [9, 10], may be obtained most naturally in F-theory.¹ It requires that the 7-branes associated to the SM gauge group all reside in the same stack. Note that this does not imply the existence of an explicit GUT structure since the symmetry may be directly broken to the SM at the string scale. The value of the gauge coupling is related to the (inverse) size of the 4-volume V_7 wrapped by the 7-branes. There is a Kahler modulus whose real part t is proportional to this volume V_7 . This is the local modulus which will be relevant for the generation of the MSSM SUSY-breaking soft terms. In particular we will assume that a non-vanishing VEV for the auxiliary field of this *local* Kahler modulus t exists giving rise to gaugino masses and other SUSY-breaking soft terms.

The final assumption requires the existence of at least one Yukawa coupling (that of the top quark) of order g (the MSSM gauge coupling constant), and has an obvious phenomenological motivation. It turns out to be quite restrictive too. Chiral fields in $D7$ -branes come in three classes depending on their geometric origin: A-fields (from the vector multiplets A in the $D = 8$ worldvolume action of the 7-brane and fermionic partners),

¹It is also naturally obtained in heterotic string compactifications. However in the perturbative heterotic case modulus dominated SUSY breaking does not give rise to SUSY breaking soft terms to leading order, due to the classical no-scale structure.

ϕ -fields (from $D = 8$ scalar ϕ multiplets and fermionic partners) and I-fields (fields from the intersection of different 7-branes). As we will justify, this classification is valid both for simple (i.e. toroidal) Type IIB orientifold models as well as for more general F-theory configurations. As we will see cubic Yukawa couplings of order one may only be present for couplings of type (A-A- ϕ), (I-I-A) or (I-I-I). This substantially reduces the possible 7-brane configurations with a potentially realistic phenomenology.

In order to compute the SUSY-breaking scalar masses m_α and trilinear couplings $A_{\alpha\beta\gamma}$ one needs to know the dependence of the matter metrics of the MSSM chiral fields on the local Kahler modulus t . One can obtain this t -dependence for the different type of fields A, ϕ, I which depends only on the local geometry of the configuration of 7-branes. The results may explicitly be obtained in simple toroidal orientifolds but one can argue for the validity of analogous results in the more general F-theory case. One finds Kahler metrics depending like $K_\alpha = t^{-\xi_\alpha}$ with ξ_α (the 'modular weight' for the chiral field Φ_α) equal to 0, 1/2, 1 for ϕ, I, A respectively. We also estimate subleading (in α') corrections which may arise from gauge magnetic fluxes in the worldvolume of 7-branes (generically needed to get chirality) and argue that the Kahler metrics are flavour blind to leading order in α' . Armed with the knowledge of this t -dependence for the Kahler metrics and the gauge kinetic function one can then explicitly compute the SUSY-breaking soft terms as a function of the auxiliary field of the local Kahler modulus t . One can do that for the three viable distributions of matter fields consistent with Yukawa couplings i.e. (A-A- ϕ), (I-I-A) or (I-I-I). The obtained scalar soft terms are flavour blind to leading order, which is good in order to avoid too large flavour changing neutral currents (FCNC).

It is important to emphasize that our philosophy in this paper will be analogous to that in e.g. [11]. We will not assume a definite mechanism which fixes all moduli and then compute the soft terms but rather the reverse. We will rather assume as above that the auxiliary field of a local Kahler modulus is the dominant source of SUSY-breaking in the MSSM sector. Under these assumptions we will end up with very definite patterns for the low energy SUSY spectrum. If they are confirmed at LHC this would tell us that the set of simple assumptions here taken could be in the right track. This would also suggest that one should concentrate in moduli fixing models leading to modulus dominance in a IIB/F-theory setting (or dual models), gauge coupling unification and a standard large Fermi scale/Planck scale hierarchy. If they are not, a large class of string compactifications with the above simplest properties would be ruled out. In any of both cases, if low energy SUSY is found, we will extract important information about the underlying string theory compactification.

In the second part of this article we study the phenomenological viability of these three options. To do this we run the soft terms from the string/unification scale down to the electroweak (EW) scale according to the renormalization group equations and impose radiative breaking of the EW symmetry [12] in the standard way. We also compute the full spectrum of sparticles and impose experimental constraints, coming most notably from the bounds on the Higgs boson mass, branching ratios of rare decays ($b \rightarrow s\gamma$, $B_s^0 \rightarrow \mu^+\mu^-$), and the muon anomalous magnetic moment $(g - 2)_\mu$. Present limits on the masses of supersymmetric particles are also imposed. We also comment on the existence of danger-

ous charge and colour-breaking minima in the parameter space. Finally, the viability of the corresponding lightest supersymmetric particle (LSP) as a dark matter candidate is examined.

We present results for the phenomenological viability of the three string options and find that the first ('non-intersecting') option with couplings (A-A- ϕ) is very constrained. The most severe bound is obtained by demanding the absence of tachyons in the scalar sector, which imposes $\tan\beta \leq 25$. Also, this scenario fails to provide a viable dark matter candidate. The LSP is always the stau, which is charged and should be unstable to be cosmologically consistent. The simplest way out in this case would be that R-parity would be broken with lepton violating dimension four couplings. This would lead to very specific signatures at LHC.

On the other hand the other two ('intersecting') options with couplings (I-I-A) or (I-I-I) give rise to similar results. However, obtaining the appropriate amount of neutralino dark matter indicates a preference for the (I-I-I) option in which both fermions and Higgs multiplets reside at intersecting 7-branes. Agreement with the WMAP results is obtained for $\tan\beta \simeq 40$. The spectrum of sparticles may be accessible to LHC with a luminosity of $1 - 10 \text{ fb}^{-1}$ by means of the missing energy signature. The structure of the SUSY spectrum is quite specific and could also be tested at LHC.

The structure of this article is as follows. In the next section we study the Kahler moduli dependence of the metrics of MSSM chiral fields both in IIB toroidal orientifold models and more general 'swiss cheese' CY and F-theory compactifications. We classify the possible modular weights of matter fields and also the possible 7-brane configurations consistent with the existence of trilinear couplings. Subleading corrections to the metrics of matter fields coming from magnetic fluxes in the worldvolume of 7-branes are also evaluated. In section 3 the MSSM SUSY breaking soft terms are computed for the different possible 7-brane configurations in terms of the modular weights under the assumption of a single local Kahler modulus domination. The low energy SUSY spectrum and radiative Electroweak (EW) symmetry breaking is obtained in section 4 by numerically solving the renormalization group equations from the string/GUT scale down to the EW scale. We examine the experimental and dark matter constraints on the three different options for MSSM fields at 7-brane configurations and discuss how the scenario with all MSSM chiral fields at intersecting 7-branes can pass all tests. We describe in section 5 the possible implications for LHC physics, in particular the LHC reach for detecting squark and gluinos in this scheme through the missing energy signature. Section 6 is left for our final comments and conclusions.

2. Effective action of Type IIB orientifold models

As we mentioned, our assumption of Kahler modulus dominance for SUSY breaking combined with the requirement of non-vanishing leading order gaugino masses narrows very much the possible string compactifications. Present collider bounds on the mass of gluinos \tilde{g} imply $m_{\tilde{g}} > 195 \text{ GeV}$ at 95% c.l. [13]. Such large value for gaugino masses is difficult to understand if the origin of gaugino masses were loop or other type of subleading effects.

Even taking into account the low-energy running of the gluino mass, such large values seem to require the existence of a large unsuppressed tree level contribution to gaugino masses.

This seems to imply that only models in which the MSSM gauge kinetic function depends on the Kahler moduli at tree level are viable. This would exclude Type IIB orientifold models with the MSSM fields residing at either $D3$ or $D9$ -branes. In both cases the gauge kinetic function is given by the universal dilaton S -field.² It also excludes perturbative heterotic vacua, whose gauge kinetic functions have the same property to leading order. Of course we do not claim by any means that those widely studied kind of vacua are not phenomenologically viable. We just concentrate in the subset of string vacua in which the simplest assumption of Kahler dominance gives consistent soft terms to leading order. We are just left in the Type IIB case with orientifolds with SM fields residing on $D7$ -branes³ and their non-perturbative F-theory generalizations with SM fields at F-theory 7-branes. We will mostly concentrate in these classes of vacua in what follows.

Before turning to that class of vacua we should emphasize again that there are other dual vacua which would be equivalent from the point of view of the effective low-energy action and hence we will not discuss these explicitly. In particular that would be the case of Type IIB orientifolds with the SM fields residing at $D5$ -branes or the S-dual corresponding to non-perturbative heterotic vacua with the SM living on NS 5-branes. In both cases the gauge kinetic function depends on the Kahler moduli at leading order. This is also the case for Type IIA orientifolds with the SM fields residing at intersecting $D6$ -branes or coisotropic $D8$ -branes [14] which are mirror to the mentioned Type IIB orientifolds with $D7$ -branes. The same (mirror) effective action is obtained although in this case the role of the Kahler moduli is played by the complex structure fields.

2.1 Generalities

We will concentrate in this section in the case of $N = 1$ Type IIB orientifold Calabi-Yau compactifications.⁴ In this class of models one projects the string spectrum over ΩP invariant states, Ω being the worldsheet parity operator and P some particular order-2 geometric action on the compact CY coordinates. That operation leaves invariant certain submanifolds of the compact CY variety corresponding to orientifold planes. We further concentrate on the case in which these orientifold planes are $O(3)$ and $O(7)$ planes. In the first case the $O(3)$ planes are volume filling (i.e. contain Minkowski space) and are pointlike on compact space. The $O(7)$ planes will also fill Minkowski space and wrap orientifold invariant 4-cycles in the CY space. Global consistency of the compactification will in general require the presence of volume filling $D3$ - and $D7$ -branes, the latter wrapping 4-cycles Σ_4 in the CY. MSSM fields will appear from open strings exchanged among these D-branes.

²See footnotes 3 and 6 for some qualifying remarks to this point.

³Equivalent T-dual models may also be obtained from IIB orientifolds with appropriately magnetized $D9$ -branes playing the role of $D7$ -branes. Those models would have an equivalent effective action and will not be discussed here separately.

⁴For reviews on orientifold models see [15].

The closed string spectrum will include complex scalars, the axidilaton S , Kahler moduli T_i and complex structure U_n fields defined as

$$S = \frac{1}{g_s} + iC_0 ; T_j = \frac{1}{g_s(\alpha')^2} \text{Vol}(\Sigma_j) + iC_4^j ; U_n = \int_{\sigma_n} \Omega, \quad (2.1)$$

where g_s is the $D = 10$ IIB dilaton, $\Sigma(\sigma)$ are 4(3)-cycles in the CY and Ω is the holomorphic CY 3-form. Here C_p , $p = 0, 4$ are Ramond-Ramond (RR) scalars coming from the dimensional reduction of antisymmetric p-forms. In terms of these the $N = 1$ supergravity Kahler potential of the moduli is given by

$$K(S, T_i, U_n) = -\log(S + S^*) - 2 \log(\text{Vol}[\text{CY}]) - \log \left[-i \int \Omega \wedge \bar{\Omega} \right]. \quad (2.2)$$

To get intuition it is useful to consider the case of the (diagonal) moduli in a purely toroidal $T^2 \times T^2 \times T^2$ orientifold. There are 3 diagonal T_i Kahler moduli with $\text{Re}T_i = \frac{1}{g_s(\alpha')^2} A_j A_k$, $i \neq j \neq k$, A_i being the area of the i -th 2-torus. The complex structure moduli U_n , $n=1,2,3$ coincide with the three geometric complex structure moduli τ_i of the three 2-tori. One then obtains

$$K = -\log(S + S^*) - \log(\Pi_i(T_i + T_i^*)) - \log(\Pi_n(\tau_n + \tau_n^*)). \quad (2.3)$$

The reader familiar with $N = 1$, $D = 4$ string effective actions will recognize the standard log (no-scale) structure of the Kahler potential. The gauge interactions and matter fields to be identified with the SM reside at the $D7$, $D3$ -branes and/or their intersections. The gauge kinetic function of the different gauge groups may be obtained from the expansion of the Dirac-Born-Infeld action of the corresponding brane. For a $D3$ -brane and a $D7$ brane wrapping a 4-cycle Σ_j one gets the simple result (see e.g. [16] for a derivation making use of dualities)

$$f_{D7j} = T_j \quad ; \quad f_{D3} = S. \quad (2.4)$$

This difference is important: in modulus dominance the $D7$ -brane gauginos will get mass but those in $D3$ -branes will remain massless.⁵ As we emphasized, this makes it phenomenologically problematic to build models with the SM residing at $D3$ -branes, as long as SUSY-breaking originates on the Kahler moduli.⁶

In order to get a semirealistic model we need D-brane configurations giving rise to chiral gauge theories. In the classes of IIB models here considered one can classify the origin of chiral matter fields in terms of five possibilities which are pictorially summarized in figure 1.

⁵In terms of fluxes imaginary selfdual (ISD) fluxes give masses to $D7$ but not to $D3$ gauginos, see e.g. [5, 6].

⁶In the case of $D7$ -branes the gauge kinetic function can get S -dependent corrections if there are gauge fluxes in their worldvolume. For large Kahler moduli the fluxes are diluted and the correction suppressed, see the discussion below. On the other hand e.g. for $D3$ -branes sitting at singularities the gauge kinetic function does depend on the blowing-up Kahler moduli associated to the singularity. However it is difficult to make this compatible with gauge coupling unification unless one stays in the singular limit. Furthermore, to assume that SUSY-breaking is dominated by the blowing-up Kahler moduli would lead to highly non-universal soft masses, leading generically to large FCNC transitions.

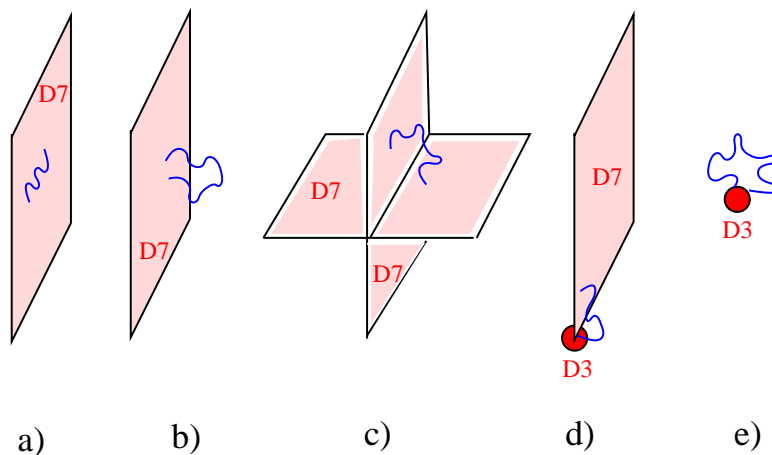


Figure 1: Different origin of matter fields in $D7/D3$ configurations. a) States from reduction of gauge fields A within the $D7$ worldvolume; b) From reduction of $D=8$ fields ϕ parametrizing the $D7$ position; c) From two intersecting $D7$ -branes; d) From open strings between a $D3$ and a $D7$; e) From open strings starting and ending on $D3$ branes.

The worldvolume theory of $D7$ -branes is 8-dimensional supersymmetric Yang-Mills and contains both vector A and matter scalar ϕ multiplets (plus fermionic partners) before compactification to $D = 4$. The first class of matter fields a) come from massless modes of the gauge multiplet fields inside the $D7$ worldvolume (plus fermionic partners). The case b) corresponds to massless modes coming from the $D = 8$ scalars ϕ which parametrize the position of $D7$ -branes in transverse space (plus fermionic partners). The case c) corresponds to massless fields coming from the exchange of open strings between intersecting $D7$ -branes. Open strings between $D3$ and $D7$ fields give rise to matter fields of type d) whereas matter fields living fully on $D3$ branes correspond to type e). Generically in order for those zero modes to be chiral in $D = 4$ additional ingredients are required. In particular there should be some non-vanishing magnetic field in the worldvolume of the $D7$ -branes. In the case of $D3$ -branes one may obtain chirality if they are located at some (e.g. orbifold) singularity in the compact CY space.

In principle SM fields could come from any of these five configurations but, as we will see, there are a number of constraints which reduce considerably the possibilities. For the moment let us state that there are specific semirealistic constructions in which SM fields live in any of these five configurations. For example, in section 9.1 of [6] there are examples in which SM particles reside in the a,b and d type of open strings above. Models with SM particles in (3,3) branes at singularities have been discussed in [17–19]. Concerning the c) configurations, they appear in magnetized $D7$ -brane models which are mirror to IIA intersecting D6-brane models [20–22, 14, 15]

2.2 Classification of 'modular weights'

An important role in the computation of SUSY-breaking soft terms is played by the metric $K_{\alpha\beta}$ of the SM matter fields C_α , which gives their normalization. This metric is in general a

non-holomorphic function of the moduli $K_{\alpha\beta}(S, S^*, T_i, T_i^*, U_n, U_n^*)$. Since we are interested in the case of modulus dominance we are particularly interested in the Kahler moduli dependence of these metrics. Such metrics have been computed in a number of string compactifications both in Heterotic and Type II cases and for large volume (i.e. to leading order in α') have a Kahler moduli dependence of the form

$$K_{\alpha\beta} = \frac{\delta_{\alpha\beta}}{\prod_i (T_i + T_i^*)^{\xi_i}}, \tag{2.5}$$

where ξ_i are often called 'modular weights' following the heterotic terminology [23] and are in general non-negative rational numbers. The precise dependence on the moduli depends on the geometry and the origin of each field. We would like to discuss now the structure of these Kahler metrics in Type IIB orientifold models as well as their non-perturbative F-theory extensions.

2.2.1 Kahler metrics in toroidal IIB orientifolds

To be more explicit concerning the metrics it is useful to consider first the case of a $Z_2 \times Z_2$ projection of a Type IIB orientifold of a factorized 6-torus $T^2 \times T^2 \times T^2$. The orientifold operation is $\Omega(-1)^{FL} R_1 R_2 R_3$ with R_i being a reflection on the i -th complex plane. Tadpole cancellation conditions require the presence of 16 $D3$ branes and 3 different sets of 16 $D7^i$ branes, each wrapping a 4-torus transverse to the i -th complex dimension. As we mentioned, a variation of this orientifold with magnetic fluxes in the worldvolume of some of the $D7$ branes was used to construct explicit semirealistic models [22, 20, 21]. A T-dual of this orientifold model with $D9$ and $D5$ -branes (without magnetic fluxes) was first constructed in [24] and the structure of Kahler metrics for the different matter fields was given in [16] using duality arguments and in [25] through an explicit string computation (see also [26] and [27]).

The five type of matter fields discussed above correspond in the toroidal/orbifold case to states with a brane structure

$$a = (7^i 7^i)_j ; b = (7^i 7^i)_i ; c = (7^i 7^j) ; d = (37^i) ; e = (33)_i \quad , \quad i \neq j, \tag{2.6}$$

where the superindices label the three types of $D7^i$ -branes whereas the subindex in the 1-st, 2-nd and 5-th cases correspond to the complex planes $j = 1, 2, 3$.

As discussed in [16, 25] these matter fields have the following Kahler metrics respectively:

$$K_{(7^i 7^i)_j} = \frac{1}{t_k} ; K_{(7^i 7^i)_i} = \frac{1}{s} ; K_{(7^i 7^j)} = \frac{1}{s^{1/2} t_k^{1/2}} ; K_{(37^i)} = \frac{1}{t_j^{1/2} t_k^{1/2}} ; K_{(33)_i} = \frac{1}{t_i}. \tag{2.7}$$

Here s and t_i are (twice) the real part of the complex dilaton S and Kahler moduli T_i ($i=1,2,3$). Note that we have not included here the dependence on the complex structure fields U_n , which will play no role in modulus dominated SUSY-breaking. Note also that all metrics are flavour diagonal. It is worth remarking the difference between the modular weights of the fields of type a) and b). This may be easily understood from the point of view

of the 8-dimensional $D7$ -brane action. As we said the first class of matter fields a) come from massless modes of the $D = 8$ gauge multiplet fields A inside the $D7$ worldvolume. The case b) corresponds to massless modes coming from the $D = 8$ scalars ϕ which parametrize the position of $D7$ -branes in transverse space. The dimensional reduction of this 8-dimensional action gives rise to kinetic terms for scalars A_i inside the vector multiplet and scalars ϕ

$$(\partial_\mu A_i)(\partial^\mu A_j) g^{ij} + (\partial_\mu \phi)^2, \tag{2.8}$$

where g^{ij} is the inverse metric corresponding to the wrapped 4-cycle. The latter scales like $g^{ii} \simeq 1/R^4 \simeq 1/t$ and gives the announced t -dependence for the metric of those fields. On the other hand the kinetic term of the scalar involves no internal contraction and is hence t -independent. Note that the massless modes from ϕ are in general not particles transforming in the adjoint of the $D7$ gauge group as sometimes stated in the literature. They are in general chiral and may accommodate SM fields (see e.g. the examples in [6]).

It is worth comparing this structure with that of perturbative heterotic orbifold compactifications. In that case there is no dependence on the heterotic axidilaton S (to leading order) and the modular weights depend on the orbifold symmetry and are 0,1 for untwisted matter and fractional numbers multiples of $(N - 1)/N$ for a Z_N twisted sector [23]. This is related to the fact that, being a closed string theory, there are matter fields which belong to Z_N twisted sectors. Nevertheless the modular weights corresponding to the overall T -modulus is always integer. On the other hand in a Type II orientifold matter fields correspond to open strings which are never twisted. The only twist possible for open strings is that coming from strings starting(ending) on Dirichlet boundary conditions and ending(starting) on Neumann boundary conditions. This gives rise to an effective twist of order 2 leading to the presence of possible modular weights 1/2 in the above Kahler metrics.⁷

In terms of the overall Kahler modulus $t = t_i$ we will thus have for the gauge kinetic function and the chiral field metrics in this toroidal case:

$$f_a = T \quad ; \quad K = \frac{1}{s^{1-\xi} t^\xi} \quad ; \quad \xi = 0, 1/2, 1. \tag{2.9}$$

Fields of type a), d) and e) have $\xi = 1$, those of type c) have $\xi = 1/2$ and type b) has $\xi = 0$.

2.2.2 Modular weight possibilities and Yukawa couplings

One could then be tempted to conclude that the SM fields may have any of these modular weights ξ freely. This is not the case if the SM construction is to be realistic. Indeed one important condition that one has to impose is that there should be at least one Yukawa coupling of order the gauge coupling constant g , the one giving mass to the top quark.⁸

⁷Note that the presence of s on the Kahler metrics makes that the dependence on the $D = 10$ dilaton g_s through the definitions (2.1) does not involve fractional powers. This possible fractional dependence is absent in the heterotic case since then, unlike the Type IIB case, the supergravity Kahler moduli do not depend on the $D = 10$ dilaton.

⁸Masses for the rest of quarks and leptons, being much smaller, could have their origin alternatively in non-renormalizable Yukawa couplings and/or non-perturbative (e.g. string instanton) effects. That is extremely unlikely in the case of the top-quark.

That means that *there should exist a trilinear superpotential* coupling a Higgs field to the left and right-handed top quarks. Now, the modular weights of fields with a trilinear coupling are not arbitrary. Consider again the toroidal $Z_2 \times Z_2$ $N = 1$ orbifold of [24]. The possible Yukawa couplings among the five types of matter fields may be easily understood in geometric terms and are classified in [24, 16]. One finds that Yukawa couplings of the following seven types exist:

$$\begin{aligned}
 A) \quad & (7^1 7^2)(7^2 7^3)(7^3 7^1) \quad [1/2, 1/2, 1/2] & B) \quad & (7^i 7^i)_j (7^i 7^k)(7^k 7^i) \quad [1, 1/2, 1/2] \\
 C) \quad & (7^i 7^i)_1 (7^i 7^i)_2 (7^i 7^i)_3 [0, 1, 1] & D) \quad & (7^i 7^i)_i (7^i 3)(3 7^i) \quad [0, 1, 1] \\
 E) \quad & (7^i 7^j)(3 7^i)(7^j 3) \quad [1/2, 1, 1] & F) \quad & (33)_i (3 7^i)(7^i 3) \quad [1, 1, 1] \\
 G) \quad & (33)_1 (33)_2 (33)_3 \quad [1, 1, 1] & &
 \end{aligned} \tag{2.10}$$

Again subindices label the three complex planes and superindices the three types of $D7^i$ -branes in this toroidal setting. In square brackets are shown the values of ξ for each of the three fields involved. One observes that many possibilities for modular weights (like $[0,0,0]$, $[0,0,1]$, etc.) are not possible. In fact this is not just a property of the described toroidal setting but, as we argue below, it also applies to more general CY IIB orientifolds and F-theory extensions.

Now, since the physical top-quark Yukawa coupling is large, it cannot have a non-renormalizable or non-perturbative origin. It should come from one of the seven tree level possibilities above. Let us consider these options in turn.

A) This corresponds to 3 sets of $D7$ -branes, each pair overlapping over one complex compact dimension. This gives rise to modular weights of type $(1/2, 1/2, 1/2)$. Semirealistic Models of this class (with magnetic fluxes in some or all the branes) have been constructed [20–22, 14, 15]. Those models are mirror to IIA orientifolds with intersecting D6-branes.

B) This case corresponds to states with couplings of type $(7^i 7^i)_j (7^i 7^k)(7^k 7^i)$ with modular weights $(1, 1/2, 1/2)$. Here one can identify the first field with the Higgs multiplet. Then quarks and leptons would come from overlapping $D7$ -branes as in the case A).

C),D)

The case C) corresponds to couplings $(7^i 7^i)_1 (7^i 7^i)_2 (7^i 7^i)_3$ among open string states within the same stack of $D7$ -branes. They have modular weights $(1, 1, 0)$ (or permutations). Note that in this case in which the particles originate in the same stack of 7-branes there is a built-in asymmetry: two of the states must have $\xi = 1$ and the other $\xi = 0$. Then it is natural to assign $\xi = 0$ to the Higgs field and $\xi = 1$ to the squarks/sleptons. Semirealistic models of this type corresponding to $D7, D3$ -branes at a Z_3 singularity may be constructed (see e.g. section 9.1 of [6]). The case D), with similar overall modular weights, correspond to models with both $D3$ and $D7$ -branes. There are couplings $(7^i 7^i)_i (7^i 3)(3 7^i)$ whose fields have analogous Kahler metrics in terms of the overall Kahler modulus than the case C). An example of these couplings are the leptonic couplings in the same example of [6]. Quarks and leptons could live

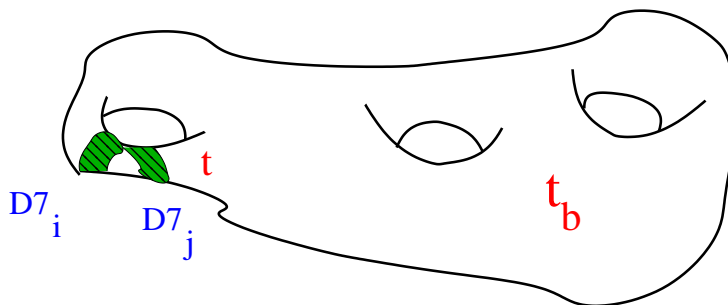


Figure 2: The SM 7-branes wrap local 4-cycles of size t in a CY whose (larger) volume is controlled by t_b .

in the $(7^i 3), (37^i)$ open strings and Higgs fields on $(7^i 7^i)_i$. Having the same modulus dependence, we will not consider this case D) as a separate option in what follows.

E), F), G)

These cases have in common that all or some of the SM gauge groups would have to reside on $D3$ -branes rather than $D7$ -branes. Since we are all the time restricting ourselves to modulus dominated SUSY-breaking, that means that the corresponding gauginos would be massless to leading order. Thus, e.g., in the first case with a $(7^i 7^j)(37^i)(7^j 3)$ coupling it is natural to locate the QCD gauge group on $D3$ -branes and that of the electroweak group on $D7$ -branes. That would imply having massless gluinos at this level. The opposite happens with case F) with $(33)_i(37^i)(7^i 3)$ couplings, which suggest locating the electroweak sector on the $D3$ -branes. Both situations are not viable phenomenologically. Concerning the last case with all SM fields on $D3$ -branes, semirealistic models of this type have been constructed in [17–19]. In this case all gauginos and scalars would be massless to leading order. This yields equivalent results to modulus dominance in Heterotic compactifications. Again this is not viable phenomenologically unless we give up our main assumption of modulus dominance and/or assume that all SUSY-breaking terms come from subleading effects. The latter, although possible, is extremely model dependent and would make the identification of any stringy pattern on the MSSM SUSY-breaking spectrum quite difficult.

We thus conclude that under the assumption of an overall Kahler modulus dominance, there are essentially only three options, those corresponding to the modular weight distributions $(1/2, 1/2, 1/2)$, $(1/2, 1/2, 1)$ or $(1, 1, 0)$ with couplings of type A), B) and C) respectively. From now on we will denote these three options by (I-I-I), (I-I-A) and (A-A- ϕ) respectively.

2.3 Modular weights beyond the toroidal setting

Toroidal orientifolds are very special in some ways and we would like to see to what extent the above findings generalize to more general IIB CY orientifolds. In particular $D7$ -branes

wrap 4-tori whose volume are directly related to the overall volume of the compact manifold and then the Kahler moduli appearing in the MSSM effective action are directly related to the overall volume of the tori. This is unnecessarily restrictive. One would expect more general situations in which the $D7$ -branes wrap local cycles whose volume need not be directly related to the overall volume of the CY. An example of this is provided by the 'swiss cheese' type of compactifications discussed in [28–32] (see also [33]). In this class of more general models one assumes that the SM resides at stacks of $D7$ -branes wrapping 'small cycles' in a CY whose overall volume is controlled by a large modulus t_b (see figure2) [28]:

$$\text{Vol}[\text{CY}] = t_b^{3/2} - h(t_i), \tag{2.11}$$

where h is a homogeneous function of the 'small' Kahler moduli t_i of degree $3/2$. The simplest example of CY with these characteristics is the manifold $\mathbf{P}_{[1,1,1,6,9]}^4$ which has two Kahler moduli with [34]⁹

$$\text{Vol}[\text{CY}] = t_b^{3/2} - t^{3/2}. \tag{2.12}$$

One important difference with the assumptions in [28] is that we will not necessarily assume that t_b is hierarchically larger than the t_i , rather we will just assume that t_b is big enough so that an expansion in powers of t_i/t_b makes sense. We will consider for simplicity the dependence on a single small Kahler modulus t and take both t_b and t sufficiently large so that the supergravity approximation is valid. In this case one can write a large volume expansion for the Kahler metrics of the form [28, 29]

$$K_\alpha = \frac{t^{(1-\xi_\alpha)}}{t_b}, \tag{2.13}$$

with ξ_α the modular weights discussed above. Note that for $t_b = t$ one would recover the metric we found in toroidal cases. Although the explicit computation of the t -dependence in the toroidal models is not directly applicable, one can use arguments based on the scaling of Yukawa couplings to indirectly compute the ξ_α 's. The idea [29] is to recall that the physical Yukawa couplings $\hat{Y}_{\alpha\beta\gamma}$ among three fields Φ_α after normalization is given in $N = 1$ supergravity by

$$\hat{Y}_{\alpha\beta\gamma} = e^{K/2} \frac{Y_{\alpha\beta\gamma}^{(0)}}{(K_\alpha K_\beta K_\gamma)^{1/2}}, \tag{2.14}$$

where $Y^{(0)}$ is the unnormalized Yukawa coupling which in Type IIB orientifolds is known to be holomorphic on the complex structure fields U_n and independent of the Kahler moduli T_i (at least at the perturbative level we are considering).¹⁰ From that we conclude (using eqs. (2.11)–(2.13)) that the physical Yukawa coupling scales with t like

$$\hat{Y}_{\alpha\beta\gamma} \simeq t^{1/2(\xi_\alpha + \xi_\beta + \xi_\gamma - 3)}. \tag{2.15}$$

Note that the dependence on the large modulus t_b has dropped to leading order in t/t_b . This is expected since the wave functions of the matter fields are localized at the small

⁹For discussions for the multi Kahler moduli case see [30, 32, 33].

¹⁰This also justifies a posteriori the $1/t_b$ dependence in eq. (2.13), it is required to get a holomorphic Yukawa $Y_{\alpha\beta\gamma}$ non dependent on the large Kahler modulus t_b .

cycles and the Yukawa couplings are local quantities, expected to be independent of the overall volume. On the other hand one can estimate the t -scaling of the physical Yukawa couplings by noting that in Type IIB string theory those may be found by computing overlap integrals of the form

$$\widehat{Y}_{\alpha\beta\gamma} \simeq \int dx^6 \Psi_\alpha \Psi_\beta \Psi_\gamma, \tag{2.16}$$

where the integration is dominated by the overlap region of the three extra dimension zero mode states Ψ .¹¹ This assumes those internal wave functions to be normalized:

$$\int |\Psi_\alpha|^2 = \int |\Psi_\beta|^2 = \int |\Psi_\gamma|^2 = 1. \tag{2.17}$$

Consider for example the case of three sets of $D7_i$ branes $i=1,2,3$, each overlapping with each other on a complex curve of real dimension two. This would be analogous to the case A) in the toroidal discussion above. Then the matter fields at each intersection overlap over a 2-dimensional space and, given eq. (2.17), their wave function would scale like $1/R \simeq t^{-1/4}$. For a trilinear coupling to exist the three $D7$ -branes overlap in this case at a point so that the physical Yukawa should scale like $\widehat{Y} \simeq (1/t^{1/4})^3 = t^{-3/4}$. This agrees with (2.15) if the three chiral fields at intersecting $D7$ -branes have $\xi_\alpha = 1/2$. Using a similar argument, this time using the Yukawa coupling of type (I-I-A) one finds that bulk fields of type A have $\xi_\alpha = 1$. Finally from the existence of couplings of type $A - A - \phi$ one gets that the zero modes from fields of type ϕ have $\xi_\alpha = 0$. Note that for $t_b \propto t$ the Kahler metrics from (2.13) will then have the same t -scaling that we found for the different type of zero modes A, ϕ , I in the toroidal case.

The above results are rather general and independent of the particular geometry as long as one can approximate the scaling behaviour with a single local Kahler modulus. Note also that one can only extract information on the sum of the three modular weights (not the individual ones) and obviously only in models in which those trilinear couplings actually exist

2.4 F-theory, coupling unification and modular weights

Our discussion in the previous section applied to general Type IIB CY orientifolds in which the SM fields reside at $D7$ -branes wrapping 'small' cycles characterized by an overall local modulus t . Prototype models in this class would include intersecting $D7$ -brane models in which the gauge group factors $SU(3)$, $SU(2)$ and $U(1)$ of the SM live on different stacks of branes. In this case each gauge coupling is independent and the observed joining of the MSSM coupling constant at a scale of order 10^{16} GeV would be a coincidence. It would be interesting to have a Type IIB framework with the SM on 7-branes in which gauge coupling unification could be a natural property from the beginning. This would also be more consistent with our approach which uses a single Kahler modulus t to describe the Kahler metrics and gauge kinetic function.

¹¹See [21] for an explicit computation of Yukawa couplings in a semirealistic D-brane model by the overlap integral of three internal wave functions.

It has recently been pointed out [9, 10] that one can naturally achieve gauge coupling unification within a type IIB scheme by going to F-theory non-perturbative constructions. This is not the place to give an introduction to these constructions but we will discuss here briefly how our discussion above is extended to F-theory compactifications of the class introduced in [9, 10]. For an introduction to F-theory and references see the first paper in [2].

F-theory [35] may be considered as a non-perturbative version of Type IIB orientifolds. Type IIB string theory in $D = 10$ has a non-perturbative $SL(2, \mathbf{Z})$ S-duality symmetry under which the complex field $\tau = \frac{1}{g_s} + iC_0$ transforms. Here $1/g_s$ is the Type IIB dilaton and C_0 is the RR scalar. The idea in F-theory is to identify locally this τ with the complex structure of a 2-torus (also having this $SL(2, \mathbf{Z})$ symmetry) living in extra 11-th and 12-th dimensions. Thus F-theory gives a geometric description of the S-duality symmetry in compactifications of Type IIB theory. One considers compactifications of this 12-dimensional theory on a CY complex 4-fold M down to $D = 4$. The CY 4-fold must be elliptically fibered over a complex 3-dimensional CY B , meaning that locally one can write $M = T^2 \times B$ with the complex structure modulus of the T^2 identified with τ . These are clearly non-perturbative vacua since the $SL(2, \mathbf{Z})$ symmetry includes transformations under which $g_s \rightarrow 1/g_s$. Still the effective $D = 4$, $N = 1$ low-energy action will admit a standard perturbative supergravity description in the limit of large Kahler moduli.

In [10] specific F-theory models with 7-brane content with GUT-like structure were discussed. These are *local* models in which only the local GUT physics associated to a chain of singularities is specified. In these models F-theory is compactified on complex 4-folds with the local structure

$$M = T^2 \times B = K3 \times S, \tag{2.18}$$

in which the $K3$ surface is itself elliptically fibered. The theory contains F-theory 7-branes which wrap the complex 2-fold S . Inside the 3-fold B these 7-branes correspond to complex codimension 1 singularities. Depending on the canonical ADE classification of the singularities the gauge groups of 7-branes have A_n , D_n or exceptional $E_{6,7,8}$ algebras. This corresponds to gauge groups $SU(n+1)$, $SO(2n)$ and E_6, E_7, E_8 . Thus the gauge group in F-theoretical 7-branes goes beyond what one can get in perturbative Type IIB orientifold $D7$ -branes in which only $SU(n+1)$ and $SO(2n)$ gauge groups may be obtained. Furthermore in F-theory the matter content in models with $SO(2n)$ gauge symmetry may include spinorial representations which are not present in perturbative IIB orientifold compactifications. This is important since e.g. it allows for an underlying $SO(10)$ GUT structure which contains SM fields in 16-dimensional spinorial representations.

The $D = 4$ chiral matter fields in F-theory have the same qualitative origin as in perturbative IIB orientifolds. There are chiral fields of type A and ϕ which originate from zero modes upon compactification of the $D = 8$ 7-brane worldvolume gauge field A and scalar ϕ (plus fermionic partners). As in the perturbative case, in order to get chirality one needs to add additional magnetic backgrounds which break the original 7-brane gauge

symmetry as $G \rightarrow H$. Then superpotential couplings of the schematic form

$$Y_{\alpha\beta\gamma} = \int_S h_{nlm} A^{\alpha n} A^{\beta l} \phi^{\gamma m} \tag{2.19}$$

are in general present [9, 10]. Here the coupling constant is evaluated as an overlap integral of the three wave functions over the 4-dimensional manifold S which is wrapped by the corresponding 7-brane. The h_{nlm} are gauge structure constants in the branching $G \rightarrow H$. This is analogous to the couplings of type C) which we discussed for the perturbative case. One interesting point noticed in [10] is that in fact for the simplest class of F-theory compactifications (with S either Hirzebruch or del Pezzo surfaces) there are no ϕ zero modes and hence no Yukawa couplings are possible. This is important since it indicates that in F-theory Yukawa couplings of type (A-A- ϕ) are difficult to obtain and hence this class of models is phenomenologically problematic.

In addition, chiral zero modes may appear at the intersection of F-theory 7-branes over a complex Riemann curve Σ inside S . In the F-theory language this corresponds to the case in which the rank of the singularity type of a 7-brane can increase along this complex curve Σ (e.g. from D_5 to E_6). This is analogous to the case we mentioned for the toroidal setting with two $D7$ branes intersecting over a common $\Sigma = T^2$. The overlap takes place in 6 real dimensions (Minkowski $\times \Sigma$) and hypermultiplet matter fields Λ, Λ^c appear coming from open strings going from one 7-brane with base S to another with base S' . Each 7-brane carries its own gauge group G_S and $G_{S'}$ and the hypermultiplets transform as 'bifundamentals' with respect to the $G_S \times G_{S'}$ group. We talk of bifundamentals here in a generalized sense since e.g. fields may transform like spinorials or antisymmetric tensors with respect G_S or $G_{S'}$.¹² In this case there are trilinear Yukawa couplings which involve an overlap integral over the complex curve Σ

$$Y_{\alpha\beta\gamma} = \int_{\Sigma} c_{ijk} \Lambda^{c;\alpha i} A^{\beta j} \Lambda^{\gamma k}, \tag{2.20}$$

where the $A^{\beta j}$ are zero modes originated from $D = 8$ gauge fields either in one 7-brane or the other. Here c_{ijk} are structure constants associated to the branching $G_{\Sigma} \rightarrow G_S \times G_{S'}$, where here G_{Σ} is the group corresponding to the enhanced singularity. Note that this is analogous to the Yukawa couplings of type (I-A-I) which we described in the toroidal case.

There is also the possibility of two 'matter curves' Σ_1 and Σ_2 in the base S to collide in a point. This would correspond to a triple intersection of 7-branes. At the intersection of Σ_1 and Σ_2 the rank of the singularity type would be increased in two units. An example in which a singularity of type $SO(12)$ is enhanced to a E_8 type singularity is described in [10]. Chirality is obtained by adding a $U(1)$ magnetic background in the underlying 7-brane with gauge group $SO(12)$. The effective low-energy group is $SO(10)$ with 3 generations of spinorial 16-plets ('bifundamentals' under the E_8 subgroup $SO(12) \times U(1)_1 \times U(1)_2$) living at two intersection matter curves Σ_1 and Σ_2 . The Higgs 10-plets arise from an additional

¹²Here G_S and $G_{S'}$ need not be SM gauge factors, one could have e.g. $G_S = SO(10)$, $G_{S'} = U(1)$ with G_S broken down to the SM with the addition of gauge backgrounds on the complex curve Σ .

matter curve Σ_3 . In these multiple 7-brane intersections trilinear Yukawa couplings involving three 'bifundamentals' exist. In this coupling wave functions overlap in a point and one has a structure

$$Y_{\alpha_j\beta_k\gamma_l} = f_{abc}\Lambda^{\alpha_j a}\Lambda^{\beta_k b}\Lambda^{\gamma_l c}. \tag{2.21}$$

This is the F-theory analogue of the (I-I-I) couplings that we described in the toroidal setting.

We thus see that the three classes of Yukawa couplings that we discussed for perturbative Type IIB orientifolds are the ones which are present in the local F-theory non-perturbative models of [9, 10]. Note that the Yukawa couplings are again obtained as overlap integrals with support on spaces of real dimension 4, 2 and 0 respectively. Thus our arguments to extract the leading large Kahler modulus behaviour of the Kahler metrics of matter fields go through in these local F-theory models. We thus again have that, restricting ourselves to the dependence on an overall local Kahler modulus t the matter fields of type A, I and ϕ have respectively modular weight 1, 1/2 and 0.

In many aspects the results are quite analogous to the perturbative orientifold case in which gauge groups reside on standard $D7$ s. The main difference is that 'bifundamentals' in F-theory may contain full SM generations (like in the spinorial of $SO(10)$) and the full gauge group of the SM may reside in a single F-theory 7-brane stack. Thus in these more general schemes gauge coupling unification is a natural consequence.

2.5 Subleading effects from magnetic gauge fluxes

In previous sections we have mentioned that getting chirality of the massless $D = 4$ spectrum in general requires the presence of magnetic fluxes¹³ on the worldvolume of $D7$ or F-theory 7-branes. So a natural question is how these fluxes may modify the Kahler metrics of matter fields and the gauge kinetic functions associated to 7-branes. In the large t limit the magnetic fluxes are diluted and one expects this effect to be subleading. To flesh out this expectation we study in this subsection how the presence of magnetic fluxes affects the effective action in the well studied case of toroidal/orbifold perturbative IIB orientifolds. The presence of fluxes affects in the same manner perturbative and F-theory compactifications so one expects similar corrections to appear in more general CY orientifolds or F-theory compactifications.

In configurations with sets of overlapping $D7$ -branes in generic 4-cycles, chirality is obtained by the addition of magnetic fluxes. In the case of toroidal/orbifold IIB orientifold models the corrections to the Kahler metrics coming from magnetic fluxes have been computed in [25]. The results for states coming from strings within $D7$ s wrapping the same cycle is

$$K_{(\tau^i\tau^j)_j} = \frac{1}{t^k} \left| \frac{1 + iF^k}{1 + iF^j} \right| ; \quad K_{(\tau^i\tau^i)_i} = \frac{1}{s} (1 + |F^j F^k|), \tag{2.22}$$

¹³These open string gauge fluxes should not be confused with antisymmetric RR and NS fluxes which may break SUSY and come from the closed string sector.

where $i \neq j \neq k$ label the 3 2-tori and F^i is the magnetic flux going through the i -th 2-torus which may be written as

$$F^i = n^i \left(\frac{st_i}{t_j t_k} \right)^{1/2}, \quad (2.23)$$

with n^i quantized integer fluxes. For states coming from open strings in between (magnetized) branes $D7^a, D7^b$ wrapping different 4-tori one has

$$K_{7^a 7^b} = \frac{1}{(st_1 t_2 t_3)^{1/4}} \prod_{j=1}^3 u_j^{-\theta_{ab}^j} \sqrt{\frac{\Gamma(\theta_{ab}^j)}{\Gamma(1 - \theta_{ab}^j)}}, \quad (2.24)$$

where u_j are the real parts of the complex structure moduli, Γ is the Euler Gamma function and

$$\theta_{ab}^j = \arctan(F_b^j) - \arctan(F_a^j). \quad (2.25)$$

The gauge kinetic functions are also modified in the presence of magnetic fluxes as

$$Ref_i^a = T_a (1 + |F_a^j F_a^k|). \quad (2.26)$$

These results apply for the case of a toroidal and/or orbifold setting. However we can try to model out what could be the effect of fluxes in a more general setting following the above structure. To model out the possible effect of fluxes we consider again the limit with $t_i = t$ and diluted fluxes $|F_i| = F$, i.e. large t and ignore the dependence on the complex structure u_i fields. Then from the above formulae one obtains the dilute flux results

$$K_{(7^i 7^i)_j}^i = \frac{1}{t}; \quad (2.27)$$

$$K_{(7^i 7^i)_i} = \frac{1}{s} \left(1 + a_i \frac{s}{t} \right), \quad (2.28)$$

$$K_{7^a 7^b} = \frac{1}{(s^{1/2} t^{1/2})} \left(1 + c_{ab} \frac{s^{1/2}}{t^{1/2}} \right), \quad (2.29)$$

$$Ref_i = t + a_i s, \quad (2.30)$$

where a_i, c_{ab} are constants (including the flux quanta) of order one. The limit of eq. (2.24) has been obtained by expanding the Gamma functions for dilute flux and ignoring the dependence through the complex structure fields u_i .¹⁴ The above three formulae may be summarised by:

$$K_{matter} = \frac{1}{s^{(1-\xi)} t^\xi} \times (1 + c_\xi (s/t)^{1-\xi}) = \frac{1}{s^{(1-\xi)} t^\xi} + \frac{c_\xi}{t}, \quad (2.31)$$

with c_ξ some flux-dependent constant coefficient whose value will depend on the modular weight ξ and the magnetic quanta. This means that the addition of fluxes has the effect in all cases to add a term which goes like $1/t$. Heuristically, we know that with the addition

¹⁴As an example, for the specific semirealistic model analyzed in detail in [36] one has $F_1^2 = F_1^3 = 3^2(s/t)$, $F_2^i = F_3^j = 0$ with $a_1 = 3^2$, $c_{12} = c_{13} = 3 \ln 4$, $a_2 = a_3 = c_{23} = 0$.

of large fluxes the $D7$ -branes turn into localized branes, very much like $D3$ -branes and we know that matter fields for $D3$ -branes have kinetic terms which go like $1/t$. This would explain the modulus dependence of the last term in eq. (2.31).

Note that in the diluted flux limit ($s/t \rightarrow 0$) one recovers the metrics we discussed above. Note also that the metric for the matter fields of type $(7^i 7^j)_{j \neq i}$ does not get corrections to leading order. As we will see, this will imply that fields with modular weight $\xi = 1$ will remain massless even after SUSY-breaking. This will be phenomenologically relevant, as we will discuss in the numerical analysis.

For a non-toroidal compactification one expects analogous flux corrections to exist with $\xi = 0, 1/2, 1$ corresponding to chiral fields of types ϕ , I and A respectively. One can easily generalize this to the case in which t is the local modulus coupling to the MSSM in a CY whose volume is controlled by a larger modulus t_b , as we discussed in section (2.3). From the above expressions one then finds a Kahler moduli dependence of the matter metrics of the form

$$K_{matter} = \frac{t^{1-\xi}}{t_b} (1 + c_\xi t^{\xi-1}) = \frac{t^{(1-\xi)}}{t_b} + \frac{c_\xi}{t_b}. \quad (2.32)$$

as a generalization of eq. (2.13). We have ignored here the explicit dependence on s and the complex structure, not relevant for the computation of soft terms. Note that the flux correction actually will depend on the large modulus t_b rather than on the local modulus t .

An additional comment is in order concerning flavour dependence of the Kahler metrics. In the three schemes with viable Yukawa couplings described above all quarks and leptons have the same t -dependent Kahler metrics. Magnetic fluxes may introduce some flavour dependence but that will be suppressed in the dilute limit here considered. On the other hand this is still compatible with the existence of a non-trivial flavour dependence in the Yukawa coupling sector. Indeed in Type IIB orientifolds the holomorphic Yukawa (perturbative) superpotentials are independent from the Kahler moduli and depend on the complex structure fields. Thus flavour dependence in these compactifications resides in the complex structure sector. This fact has also been recently emphasized in [29, 37]. Since we are assuming that SUSY-breaking is taking place only in the Kahler moduli sector, no flavour dependence will appear to leading order in the Kahler metrics and, consequently, in the obtained SUSY-breaking soft terms.

As a final observation, note that with SUSY-breaking controlled by the auxiliary field F_t of a particular local modulus t , the complex phases which might appear in soft terms can always be rotated away [38]. So there will be no SUSY-CP problem and there will be no SUSY contribution to the electric dipole moment of the neutron coming from soft terms.

3. Supersymmetry breaking and soft terms

With the knowledge of the matter metrics and gauge kinetic function we can now address the computation of SUSY-breaking soft terms. We have argued that similar dependence for Kahler metrics on the local Kahler modulus t is obtained both for perturbative IIB CY

orientifolds and F-theory compactifications so that our results will apply to both in the large moduli limit.

3.1 Soft terms

For definiteness we will assume here that the SM resides at stacks of 7-branes wrapping locally small cycles (of size t) in a CY whose overall volume is controlled by a large modulus t_b . We can model out this structure with a Kahler potential of the form [28, 29]

$$G = -2\log(t_b^{3/2} - t^{3/2}) + \log|W|^2, \tag{3.1}$$

with $t = T + T^*$ being the relevant local modulus. The gauge kinetic function and Kahler metrics of a MSSM matter field C_α will be given to leading order by¹⁵

$$f_a = T \quad ; \quad K_\alpha = \frac{t^{(1-\xi_\alpha)}}{t_b}. \tag{3.2}$$

It is important to emphasize that what is relevant here is the no-scale structure of the moduli Kahler potential and that analogous results would be obtained in a more general CY or F-theory compactification with more Kahler moduli as long as we assume that our local modulus t is smaller than the overall modulus t_b so that an expansion on t/t_b is consistent. As emphasized in [28, 29] the soft terms obtained for the local MSSM 7-branes will only depend on the local modulus t and the value of its auxiliary field F_t and will not directly depend on the large moduli. Note that, being a Kahler modulus, assuming $F_t \neq 0$ corresponds to the assumption of the presence of non-vanishing SUSY-violating antisymmetric ISD fluxes [4] in the region in compact space where the SM 7-branes reside.

The MSSM superpotential W has the general form

$$W = W_0 + \frac{1}{2}\mu H_u H_d + \frac{1}{6}Y_{\alpha\beta\gamma}^{(0)} C^\alpha C^\beta C^\gamma + \dots \tag{3.3}$$

Here W_0 is a moduli dependent superpotential (including typically a flux-induced constant term which controls the scale of SUSY breaking). The C^α are the SM chiral superfields and $H_{u,d}$ the minimal Higgs multiplets. In our computations below we are going to assume that there is an explicit μ -term in the Lagrangian which can be taken (at least in some approximation) as independent of the Kahler moduli t_b, t . A simple origin for such a term is again RR and NS fluxes. Indeed, it is known that certain SUSY preserving fluxes may give rise to explicit supersymmetric mass terms to chiral fields, μ -terms (see e.g. [39, 5, 6, 38]). If in addition there are SUSY-breaking fluxes a B -term and the rest of soft terms appear. These two kind of fluxes are in principle independent so that one can simply consider an explicit constant (moduli independent) μ -term in the effective action as a free parameter. On the other hand, if both the origin of the μ -term and SUSY-breaking soft terms is fluxes, it is reasonable to expect both to be of the same order of magnitude, solving in this way

¹⁵We will ignore here the dependence of the Kahler potential on the axidilaton and complex structure fields since they play no role in the computation of the soft terms under consideration, at least to leading order in α' .

the μ problem.¹⁶ However, in the general formulae below we will allow for the presence of a Giudice-Masiero term of the form $Z(T, T)H_u H_d + h.c.$ in the Kahler potential [40]. A discussion on alternative origins for a μ -term in string models is also given in the appendix. As we describe there, these alternatives do not look particularly promising concerning correct radiative EW symmetry breaking and hence in the phenomenological analysis below we will set $Z = 0$.

Now the form of the effective soft Lagrangian obtained is given by (see e.g. [41])

$$\begin{aligned} \mathcal{L}_{\text{soft}} = & \frac{1}{2}(M_a \hat{\lambda}^a \hat{\lambda}^a + h.c.) - m_\alpha^2 \hat{C}^{*\bar{\alpha}} \hat{C}^\alpha \\ & - \left(\frac{1}{6} A_{\alpha\beta\gamma} \hat{Y}_{\alpha\beta\gamma} \hat{C}^\alpha \hat{C}^\beta \hat{C}^\gamma + B \hat{\mu} \hat{H}_u \hat{H}_d + h.c. \right), \end{aligned} \quad (3.4)$$

with

$$M_a = \frac{1}{2} (Ref_a)^{-1} F^m \partial_m f_a, \quad (3.5)$$

$$m_\alpha^2 = \left(m_{3/2}^2 + V_0 \right) - \overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \log K_\alpha, \quad (3.6)$$

$$A_{\alpha\beta\gamma} = F^m \left[K_m + \partial_m \log Y_{\alpha\beta\gamma}^{(0)} - \partial_m \log(K_\alpha K_\beta K_\gamma) \right], \quad (3.7)$$

$$\begin{aligned} B = & \hat{\mu}^{-1} (K_{H_u} K_{H_d})^{-1/2} \left\{ \frac{W^*}{|W|} e^{K/2} \mu \left(F^m \left[\hat{K}_m + \partial_m \log \mu \right. \right. \right. \\ & \left. \left. \left. - \partial_m \log(K_{H_u} K_{H_d}) \right] - m_{3/2} \right) \right. \\ & + \left(2m_{3/2}^2 + V_0 \right) Z - m_{3/2} \overline{F}^{\overline{m}} \partial_{\overline{m}} Z \\ & + m_{3/2} F^m \left[\partial_m Z - Z \partial_m \log(K_{H_u} K_{H_d}) \right] \\ & \left. - \overline{F}^{\overline{m}} F^n \left[\partial_{\overline{m}} \partial_n Z - (\partial_{\overline{m}} Z) \partial_n \log(K_{H_u} K_{H_d}) \right] \right\}, \end{aligned} \quad (3.8)$$

where V_0 is the vacuum energy which is assumed to be negligibly small from now on. Here \hat{C}^α and $\hat{\lambda}^a$ are the scalar and gaugino canonically *normalized* fields respectively

$$\hat{C}^\alpha = K_\alpha^{1/2} C^\alpha, \quad (3.9)$$

$$\hat{\lambda}^a = (Ref_a)^{1/2} \lambda^a, \quad (3.10)$$

and the rescaled Yukawa couplings and μ parameter

$$\hat{Y}_{\alpha\beta\gamma} = Y_{\alpha\beta\gamma}^{(0)} \frac{W^*}{|W|} e^{K/2} (K_\alpha K_\beta K_\gamma)^{-1/2}, \quad (3.11)$$

$$\hat{\mu} = \left(\frac{W^*}{|W|} e^{K/2} \mu + m_{3/2} Z - \overline{F}^{\overline{m}} \partial_{\overline{m}} Z \right) (K_{H_u} K_{H_d})^{-1/2}, \quad (3.12)$$

have been factored out in the A and B terms as usual.

¹⁶An alternative to get a Higgs bilinear could be the presence of an additional SM singlet chiral field X coupling to the Higgs fields like $X H_u H_d$. It is trivial to generalize all the above formulae to this NMSSM case.

In Type IIB orientifolds the holomorphic perturbative superpotential is independent of the Kahler moduli so that the derivatives of $\widehat{Y}^{(0)}$ in the expression for A vanish. Using these formulae one then obtains general soft terms (for $t_b \gg t$) as follows:

$$M = \frac{F_t}{t}, \quad (3.13)$$

$$m_\alpha^2 = (1 - \xi_\alpha)|M|^2, \quad \alpha = Q, U, D, L, E, H_u, H_d, \quad (3.14)$$

$$A_U = -M(3 - \xi_{H_u} - \xi_Q - \xi_U),$$

$$A_D = -M(3 - \xi_{H_d} - \xi_Q - \xi_D),$$

$$A_L = -M(3 - \xi_{H_d} - \xi_L - \xi_E),$$

$$B = -M(2 - \xi_{H_u} - \xi_{H_d}).$$

This is analogous to results found in [38] for the case of a single overall modulus T . Note that the dependence on the SUSY-breaking from the large modulus t_b disappears to leading order and the size of soft terms is rather controlled by the local modulus t . In particular, gaugino masses corresponding to the SM gauge groups depend only on the modulus of the local 4-cycles they wrap, rather than the large volume modulus t_b . These gaugino masses set the scale of the SUSY-breaking soft terms. As in [28, 29] here the gravitino mass will be given approximately by the auxiliary field of large moduli, $m_{3/2} \simeq -F_{t_b}/t_b$, with corrections suppressed by the large volume t_b . One can check that all the bosonic soft terms above may be understood as coming from the positive definite scalar potential [38]

$$V_{SB} = \sum_\alpha (1 - \xi_\alpha)|\partial_\alpha W - M^* C_\alpha^*|^2 + \sum_\alpha \xi_\alpha |\partial_\alpha W|^2. \quad (3.15)$$

The positive definite structure may be seen as a consequence of the 'no-scale' structure of modulus domination and is expected to apply also in more general F-theory compactifications.

Within the philosophy of gauge coupling unification we will assume unified modular weights:

$$\xi_f = \xi_Q = \xi_U = \xi_D = \xi_L = \xi_E. \quad (3.16)$$

This is also reasonable within e.g. an F-theory approach with an underlying GUT-like symmetry like $SO(10)$ in which one expects all fermions to have the same modular weight. In us much as the effect of magnetic fluxes giving rise to chirality for SM fermions is negligible (see next section) one also expects flavour independence for them. So we will assume a universal modular weight ξ_f for all quarks and leptons. Concerning the Higgs multiplets we have seen that they can have $\xi_H = 0, 1, 1/2$ and hence there is no reason why they should have the same modular weight as chiral fermions. We will however assume that both Higgses have the same modular weight $\xi_H = \xi_{H_u} = \xi_{H_d}$, as would also be expected in models with an underlying left-right symmetry.

Under these conditions the summary of the soft terms obtained for the three possibilities for brane distributions with consistent Yukawa couplings, (A-A- ϕ), (I-I-A) and (I-I-I) are shown in table 1.

(ξ_L, ξ_R, ξ_H)	Coupling	M	m_L^2	m_R^2	m_H^2	A	B
(1, 1, 0)	(A-A- ϕ)	M	0	0	$ M ^2$	$-M$	$-2M$
(1/2, 1/2, 1)	(I-I-A)	M	$\frac{ M ^2}{2}$	$\frac{ M ^2}{2}$	0	$-M$	0
(1/2, 1/2, 1/2)	(I-I-I)	M	$\frac{ M ^2}{2}$	$\frac{ M ^2}{2}$	$\frac{ M ^2}{2}$	$-3/2M$	$-M$

Table 1: Modulus dominated soft terms for choices of modular weights ξ_α which are consistent with the existence of trilinear Yukawa couplings in 7-brane systems.

Note that in the scenarios with couplings (A-A- ϕ) and (I-I-A) it is natural to assume that the Higgs field is identified with fields of type ϕ and A respectively and these are the cases displayed in the table. Concerning the B parameter it is obtained assuming an explicit μ -term.

We will study the phenomenological viability of these three options in the next section. Since the cases with couplings (I-I-A) and (I-I-I) only differ in the origin of the Higgs fields, and hence in the values of their modular weights, they will give rise to a similar phenomenology. For this reason, we will analyse the general case in which the Higgs modular weight is a free parameter, ξ_H , and regard examples (I-I-A) and (I-I-I) as the limiting cases with $\xi_H = 1$ and $\xi_H = 1/2$, respectively. Notice that this can be understood as if the physical MSSM Higgs was a linear combination of two fields with the correct quantum numbers, one of them living in the intersection of two $D7$ -branes and the other one in the bulk of one of them.¹⁷

On the other hand, the case with couplings (A-A- ϕ) is unrelated to the previous ones and will be studied separately.

3.2 Corrections to soft terms from magnetic fluxes

The Kahler metrics and gauge kinetic functions used in the computations above correspond to the leading behaviour in α' . It is interesting to estimate what could be the effect of subleading terms coming from possible magnetic fluxes in the 7-branes, as discussed in subsection 2.5. We know that the presence of such fluxes is required to get chirality. To do that we can use the results for the Kahler metrics given in the diluted flux approximation in eq. (2.32). One finds for the soft terms the results:

$$M = \frac{F_t}{t + as}, \tag{3.17}$$

$$m_\alpha^2 = \frac{|F_t|^2}{t^2} (1 - \xi_\alpha) \frac{(1 + \frac{c_\alpha \xi_\alpha}{t^{(1-\xi_\alpha)}})}{(1 + \frac{c_\alpha}{t^{(1-\xi_\alpha)}})^2}, \tag{3.18}$$

$$A_{\alpha\beta\gamma} = -\frac{F_t}{t} \sum_{i=\alpha,\beta,\gamma} (1 - \xi_i) \left(1 - \frac{c_i}{t^{(1-\xi_i)}} \right), \tag{3.19}$$

$$B = -\frac{F_t}{t} \sum_{i=H_u, H_d} (1 - \xi_i) \left(1 - \frac{c_i}{t^{1-\xi_i}} \right). \tag{3.20}$$

¹⁷In chapter 6 in ref. [10] an F-theory SO(10) model is constructed in which indeed there are massless Higgs multiplets both in the bulk and at intersections.

Note that for matter fields coming from a $D = 8$ vector multiplet (modular weight $\xi = 1$) the scalar terms are still zero and get no flux correction. In the computation of the B -term an explicit μ term independent on t, t_b has been assumed. We will make use of these formulae to try to estimate the effect of fluxes on the obtained low energy physics below.

3.3 Soft terms and moduli fixing

In the above computations we have tacitly assumed that all moduli have been fixed at a (slightly) de Sitter vacuum of the scalar potential. The issue of moduli fixing in Type IIB orientifold compactifications has been addressed by many authors in recent years starting with the work in [1] (see [2] for reviews and references). Concerning the computation of soft terms in a model with moduli stabilized, two situations have received special attention, 'Mirage mediation' and 'Large Volume Compactification' (LVC) models. In the case [1] of a KKLT model with a single overall modulus, it has been argued [42] that the size of the modulus dominance contribution to SUSY-breaking m_{soft} is suppressed relative to the gravitino mass as $m_{\text{soft}} \simeq m_{3/2}/\log(M_p/m_{3/2})$. Then one cannot neglect the anomaly mediation contribution to SUSY-breaking which is competitive with modulus dominance. The hierarchy $M_p/m_{3/2}$ comes from an assumed small value for the (flux-induced) superpotential W_0 . This situation goes under the name of 'mirage mediation'. In this situation one would have to add to the soft terms here computed the well known universal contribution from anomaly mediation.

In the LVC scenario the prototype of CY considered in these large volume models [28, 31] has a Kahler potential like that in eq. (2.12). In this scheme the hierarchy $M_p/m_{3/2}$ is large due to the presence of the large volume cycle controlled by t_b and the string scale is at an intermediate value, $M_s \propto 10^{11}$ GeV. Such a large value for t_b is obtained as a result of a minimization of a scalar potential in which a one loop term in the α' expansion is also included. The flux induced superpotential W_0 is of order one and the large value for t_b is due to a combination of non-perturbative instanton corrections and one-loop terms in the α' expansion. In this scheme SUSY-breaking is dominated by a modulus dominance pattern as discussed in the previous sections.

As discussed in [31], as one decreases $|W_0|$ in the LVC scheme one ends up in a situation rather analogous to mirage mediation. There are however other small $|W_0|$ minima which are compatible with a large value of the string scale (of order of the canonical GUT scale, 10^{16} GeV) and modulus dominated SUSY-breaking. Such a large string scale scenario would be interesting if one wants to understand the joining of coupling constants in a simple way. In the numerical analysis in the next section we will be tacitly assuming a situation similar to that, with modulus dominance and a string scale of order the GUT scale. In particular, in this case the anomaly mediation contribution is suppressed by the same no-scale cancellation which makes the dependence on the large modulus t_b to disappear from the final expression for soft terms [28, 31, 29],

Before presenting our study of the viability of the three different options for modulus domination that were introduced in the previous sections let us make a few comments concerning some recent computations of soft terms in string models. In [38] the cases with $(\xi_H, \xi_L, \xi_R) = (0, 0, 0)$ and $(1/2, 1/2, 1/2)$ were discussed. The phenomenological viability

of the first was analyzed in detail in [43]. However, there is no known perturbative compactification with such modular weight distribution which is compatible with the existence of trilinear couplings. Concerning the $(1/2, 1/2, 1/2)$ choice, no phenomenological analysis has been yet reported.¹⁸ As we said, in [42] it was found that in one-modulus KKLT-like vacua one is lead to a mixed situation with soft terms coming from competing contributions from modulus dominance and anomaly mediation. Phenomenological analysis have been made (see e.g. [44, 45]). It would be interesting to reconsider those analysis by including the different allowed options for modular weights. Finally, in [28] soft terms have been computed in a LVC 2-modulus scenario with an intermediate string scale. In this case one has modulus domination and our results would apply. The anomaly mediation contribution is negligible. The authors however concentrate on universal values of the modular weights $\xi = 2/3$ and again no example of compactification with such modular weights is known. Another difference with our analysis is that in that case the string scale is at an intermediate scale and the soft terms are then run from an intermediate scale down to the weak scale. In our analysis we will assume a large string scale of order the GUT scale. For a recent phenomenological study of LHC signatures for soft terms coming from different options for moduli fixing see also [47].

4. EW symmetry breaking and SUSY spectrum

We are now ready to extract the low energy implications of the MSSM soft terms listed in table 1. We will use those values as boundary conditions at the string scale which we will identify with the standard GUT scale at which the MSSM gauge couplings unify. We will solve numerically the Renormalization Group Equations (RGEs) and calculate the low energy SUSY spectrum. We will also impose standard radiative electroweak symmetry breaking.

4.1 Radiative EW symmetry breaking and experimental constraints

The minimization of the loop-corrected Higgs potential leaves the following two conditions, which are imposed at the SUSY scale,

$$\mu^2 = \frac{-m_{H_u}^2 \tan^2 \beta + m_{H_d}^2}{\tan^2 \beta - 1} - \frac{1}{2} M_Z^2, \tag{4.1}$$

$$\mu B = \frac{1}{2} \sin 2\beta (m_{H_d}^2 + m_{H_u}^2 + 2\mu^2), \tag{4.2}$$

where $\tan \beta \equiv \langle H_u \rangle / \langle H_d \rangle$ is the ratio of the Higgs vacuum expectation values and $m_{H_{u,d}}^2$ correspond to the Higgs mass parameters shifted through tadpole terms. It should be noted that our choice for the sign of the μ parameter, consistent with our convention for the superpotential (3.3), is opposite to the usual one.

A usual procedure consists then in fixing the value of $\tan \beta$ and use the experimental result for M_Z to determine the modulus of the μ parameter via eq. (4.1). The B parameter

¹⁸An intersecting brane model with a 'dilute fluxed limit' corresponding to those modular weights was analyzed in [36] but no phenomenological analysis was made. See however [46].

is then obtained by solving eq. (4.2). In this approach, once the modular weights which describe the specific model are given, the only free parameters left are the common gaugino mass, M , the value of $\tan\beta$ and the sign of the μ parameter (not fixed by condition (4.1)).

On the other hand, given that the value of the bilinear soft term, $B(M_{\text{GUT}})$, is also a prediction in these constructions (dependent on the source of the μ term) a more complete approach consists in imposing it as a boundary condition for the RGEs at the string scale. Conditions (4.1) and (4.2) can then be used to determine both $\tan\beta$ and μ as a function of the only free parameter, M . Note that this is an extremely constrained situation and it is not at all obvious a priori that solutions passing all experimental constraints exist.

It is not possible to derive an analytical solution for $\tan\beta$ from eqs. (4.1) and (4.2), since tadpole corrections to the Higgs mass terms depend on $\tan\beta$ in a non-linear way. Furthermore, $\tan\beta$ is needed in order to adjust the Yukawa couplings at the GUT scale so that they agree with data. Thus, to find a solution an iterative procedure has to be used where the RGEs are numerically solved for a first guess of $\tan\beta$ using the soft parameters as boundary conditions at the GUT scale. The resulting B at the SUSY scale is then compared with the solution of eq. (4.2). In subsequent iterations, the value of $\tan\beta$ is varied, looking for convergence of the resulting $B(M_{\text{SUSY}})$. It is not always possible to find a solution with consistent REWSB, and this excludes large areas of the parameter space.¹⁹ In our analysis we have implemented such iterative process through a modification of the SPheno2.2.3 code [49]. The constraints imposed by the full REWSB conditions have been analysed for models with universal soft parameters [50, 51], as well as string-inspired scenarios [43].

Once the supersymmetric spectrum is calculated, compatibility with various experimental bounds has to be imposed. We have taken into account the constraints obtained by LEP on the masses of supersymmetric particles, as well as on the lightest Higgs boson [13]. Also, the most recent experimental limits on the contributions to low-energy observables have been included in our analysis. More specifically, we impose the experimental bound on the branching ratio of the rare $b \rightarrow s\gamma$ decay, $2.85 \times 10^{-4} \leq \text{BR}(b \rightarrow s\gamma) \leq 4.25 \times 10^{-4}$, obtained from the experimental world average reported by the Heavy Flavour Averaging Group [52], and the theoretical calculation in the Standard Model [53], with errors combined in quadrature. We also take into account the upper constraint on the $(B_s^0 \rightarrow \mu^+\mu^-)$ branching ratio obtained by CDF, $\text{BR}(B_s^0 \rightarrow \mu^+\mu^-) < 5.8 \times 10^{-8}$ at 95% c.l. [54] (which improves the previous one from D0 [55]).

Regarding the muon anomalous magnetic moment, a constraint on the supersymmetric contribution to this observable, a_μ^{SUSY} , can be extracted by comparing the experimental result [56], with the most recent theoretical evaluations of the Standard Model contributions [57–59]. When e^+e^- data are used the experimental excess in $a_\mu \equiv (g_\mu - 2)/2$ would constrain a possible supersymmetric contribution to be $a_\mu^{\text{SUSY}} = (27.6 \pm 8) \times 10^{-10}$, where theoretical and experimental errors have been combined in quadrature. However, when tau data are used, a smaller discrepancy with the experimental measurement is found. Due to

¹⁹Moreover, the same input values for the soft parameters can lead to multiple solutions in terms of $\tan\beta$ [48]. We have explicitly checked this is not the situation in any of our examples for $\mu < 0$, but usually occurs for $\mu > 0$.

this reason, in our analysis we will not impose this constraint, but only indicate the regions compatible with it at the 2σ level, this is, $11.6 \times 10^{-10} \leq a_\mu^{\text{SUSY}} \leq 43.6 \times 10^{-10}$.

Assuming R-parity conservation, and hence the stability of the LSP, we also investigate the possibility of obtaining viable neutralino dark matter. This is, in the regions of the parameter space where the neutralino is the LSP we compute its relic density by means of the program `micrOMEGAs` [60], and check compatibility with the data obtained by the WMAP satellite [61], which constrain the amount of cold dark matter to be $0.1037 \leq \Omega h^2 \leq 0.1161$.

The value of the mass of the top quark is particularly relevant. In our computation we have used the central value corresponding to the recent measurement by CDF [62], $m_t = 172 \pm 1.4 \text{ GeV}$. We will briefly comment on the effect that deviations from this quantity may have on REWSB.

Finally, the presence in SUSY theories of scalar fields which carry electric and colour charges can lead to the occurrence of minima of the Higgs potential where charge and/or colour symmetries are broken when these scalars take non-vanishing VEVs. If these minima are deeper than the physical (Fermi) vacuum, the latter would be unstable. The different directions in the field space that can lead to this situation were analysed and classified in [63]. It was found there that the most dangerous direction corresponds to the one labelled as UFB-3, where the stau and sneutrino take non-vanishing VEVs, since the tree-level scalar potential could even become unbounded from below. These UFB constraints were found to impose stringent constraints on the parameter space of general supergravity theories [64], as well as in different superstring and M-theory scenarios [65, 66]. In our study we will comment on the constraints which are derived when the absence of such charge and/or colour-breaking minima is imposed.²⁰

In order to understand the effect of all these experimental and astrophysical constraints we have performed a scan over the gaugino mass parameter, M , and $\tan\beta$ for the three different consistent models that were specified in table 1. We will discuss first the results obtained without imposing the predicted boundary condition for B and analyse the effect of this constraint in the following subsection.

4.2 The intersecting 7-brane (I-I-I)-(I-I-A) configurations

As commented at the end of section 3.1, we will analyse these two cases together within the framework of a generic scenario in which the modular weight for the Higgs is a free parameter. In this approach, ξ_H can take any value between $\xi_H = 1/2$, which corresponds to case (I-I-I), and $\xi_H = 1$, as in case (I-I-A). The soft parameters for such a model can be

²⁰Strictly speaking, the existence of a global charge and/or colour-breaking vacuum cannot be excluded if the lifetime of the metastable physical minimum is longer than the age of the Universe, as is usually the case.

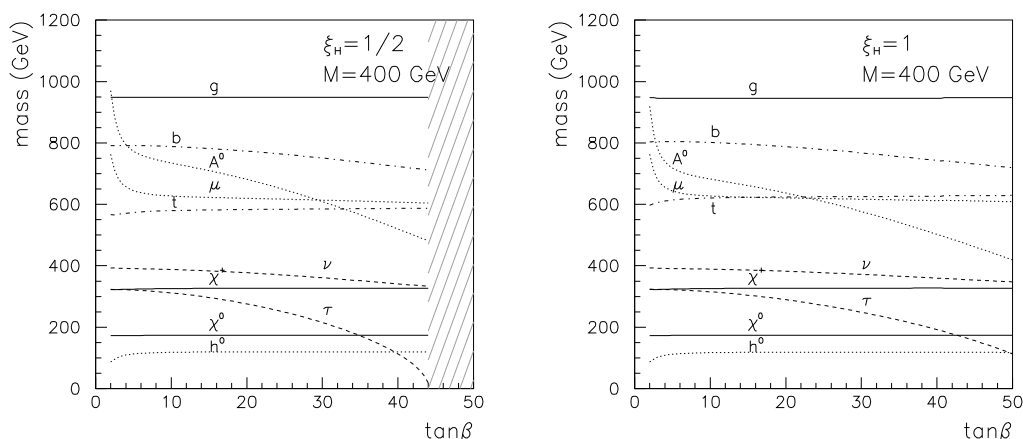


Figure 3: Low-energy supersymmetric spectrum as a function of $\tan\beta$ for $\xi_H = 1/2$, (left) and $\xi_H = 1$ (right) with $M = 400$ GeV and $\mu < 0$. From bottom to top, the solid lines represent the masses of the lightest neutralino, the lightest chargino, and the gluino. Dashed lines display the masses of the lightest stau and lightest sneutrino. Dot-dashed lines correspond to the stop and sbottom masses. Finally, the lightest Higgs mass, the pseudoscalar Higgs mass and the absolute value of the μ parameter are displayed by means of dotted lines. The ruled area for large $\tan\beta$ is excluded by the occurrence of tachyons in the slepton sector.

extracted from (3.15) and read

$$\begin{aligned}
 m_{L,E,Q,U,D}^2 &= |M|^2/2, \\
 m_{H_u,H_d}^2 &= (1 - \xi_H)|M|^2, \\
 A_{U,D,L} &= -M(2 - \xi_H), \\
 B &= -2M(1 - \xi_H).
 \end{aligned}
 \tag{4.3}$$

A sample of the resulting supersymmetric spectrum is represented in figure 3 as a function of the value of $\tan\beta$ for $M = 400$ GeV and $\mu < 0$ for the two limiting cases $\xi_H = 1/2$ and $\xi_H = 1$. As evidenced by the plot, despite the fact that the slepton mass-squared terms are always positive at the GUT scale, the running down to the EW scale can lead to the occurrence of tachyonic eigenstates for large $\tan\beta$. This is typically the case of the lightest stau. The negative contribution to the RGEs governing the stau mass parameters increases with $\tan\beta$, since it is proportional to the lepton Yukawa, which varies as $1/\cos\beta$. As a consequence, the stau mass decreases towards large values of $\tan\beta$, first becoming the LSP (which in the $\xi_H = 1/2$ example occurs for $\tan\beta \gtrsim 35$), then falling below its experimental lower bound and eventually turning tachyonic. This sets an upper constraint on the possible value of $\tan\beta$ (which obviously increases for larger values of M). The effect of this constraint is more important when ξ_H is small, since the trilinear parameter A_L is larger (more negative) and enhances the negative contribution in the RGEs of the stau mass parameters. Thus, the bound $\tan\beta \lesssim 45$ for $\xi = 1/2$ is relaxed to $\tan\beta \lesssim 55$ with $\xi = 1$.

Another interesting feature is that the resulting value for $|\mu|$, calculated from eq. (4.1), turns out to be relatively large, of order $1.5 M$. Similarly, the pseudoscalar Higgs (as well as the heavy neutral and charged Higgses) is also heavy, decreasing for large values of $\tan \beta$. When the value of ξ_H increases the predicted pseudoscalar mass is slightly smaller.

The rest of the properties of the spectrum are less sensitive to variations in the Higgs modular weight. For small values of $\tan \beta$ the lightest neutralino is typically the LSP in this example. Since the value of the μ parameter is always large, the lightest neutralino is mostly bino-like. The universality of gaugino masses at the GUT scale and the large values of $|\mu|$ also lead to the well known low-energy relation among the masses of the lightest neutralino, the lightest chargino and the gluino, $m_{\tilde{\chi}^0} : m_{\tilde{\chi}^\pm} : m_{\tilde{g}} \approx 1 : 2 : 5.5$. The squark sector is rather heavy, another consequence of the universality of the soft masses.

Notice finally that the lightest Higgs mass decreases towards small $\tan \beta$ (through the decrease of its tree-level value). Hence, the LEP lower constraint on the Higgs mass leads to a lower bound on the phenomenologically viable values of $\tan \beta$. This bound also depends on the value of M , as we will see, since M sets the overall scale for the soft parameters and therefore determines the size of loop corrections to m_h .

For a better understanding of the effect of the various experimental constraints a full scan on the two free parameters M , and $\tan \beta$ (for $\mu < 0$) is presented in figure 4 for four examples of Higgs modular weight,²¹ $\xi_H = 0.5, 0.6, 0.8$ and 1 . Some of the features of the supersymmetric spectrum we have just described are also clearly displayed. For example, the ruled area for large $\tan \beta$ and small M corresponds to the area excluded due to tachyons in the stau eigenstates. This area becomes smaller as the modular weight for the Higgs increases as a consequence of the increase in the stau mass. At the same time, the region with stau LSP (which is represented by light grey) is also shifted towards larger $\tan \beta$, thereby enlarging the area in which the neutralino is the LSP.

The region below the thin dashed line is excluded by the LEP constraint on the Higgs mass, and corresponds to $\tan \beta \lesssim 5$ and $M \lesssim 300$ GeV. This constitutes the stronger lower limit on M for small values of $\tan \beta$. For $\tan \beta \gtrsim 15$, however, the experimental bound on $\text{BR}(b \rightarrow s\gamma)$ sets a more constraining lower limit on M , which increases with $\tan \beta$ and reaches $M \gtrsim 500$ GeV.

The theoretical predictions for $\text{B}(B_s^0 \rightarrow \mu^+ \mu^-)$ only exceed the experimental upper bound for small values of M and very large $\tan \beta$. The reason is that, as already explained, the pseudoscalar Higgs is always very heavy in these scenarios. Only in the cases with $\xi_H = 0.8$ and $\xi_H = 1$ there are regions excluded for this reason. In any case, these areas are already disfavoured by a number of other experimental bounds.

Regarding the supersymmetric contribution to the muon anomalous magnetic moment, the region of the parameter space which is favoured by the experimental constraint corresponds to the area between the thin dot-dashed lines. The lower(upper) bound on a_μ^{SUSY} sets an upper(lower) limit on M . Since a_μ^{SUSY} increases for large $\tan \beta$ (through an enhancement of the contribution mediated by charginos and sneutrinos), both the upper and

²¹In these figures the constraint for the B-parameter has not yet been imposed. We will see below that the dark matter constrained combined with the prediction for B singles out the case with $\xi_H \simeq 0.5 - 0.6$.

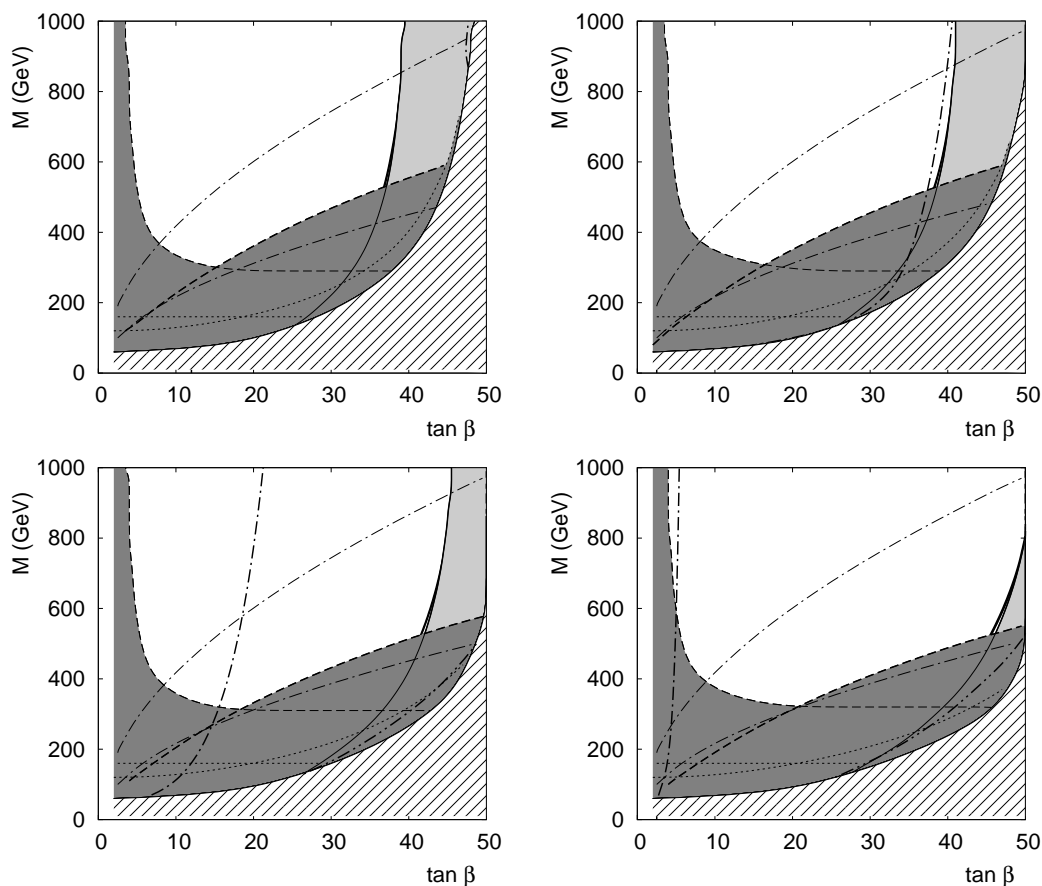


Figure 4: Effect of the various experimental constraints on the $(M, \tan \beta)$ plane for cases with $\xi_H = 0.5, 0.6, 0.8,$ and 1 , from left to right and top to bottom. Dark grey regions correspond to those excluded by any experimental bound. Namely, the area below (and to the left of) the thin dashed line is ruled out by the lower constraint on the lightest Higgs mass. The region below the thin dotted lines is excluded by the lower bounds on the stau and chargino masses. The area below the thick dashed line is excluded by $b \rightarrow s\gamma$. The region below the double dot-dashed line is excluded by $B_s^0 \rightarrow \mu^+\mu^-$. The thin dot-dashed lines correspond, from top to bottom, to the lower and upper constraint on a_μ^{SUSY} . The area contained within solid lines corresponds to the region in which the stau is the LSP, and is depicted in light grey when experimental constraints are fulfilled. In the remaining white area the neutralino is the LSP. The thin black area, in the vicinity of the region with stau LSP, corresponds to the region where the neutralino relic density is in agreement with the WMAP bound.

lower limits on M also increase. As we already indicated, in our analysis this constraint is not imposed.

As a result, there are vast regions of the parameter space compatible with all the experimental constraints and in which the lightest neutralino is the LSP. In order to determine its viability as a dark matter candidate, its relic density has to be computed and compared with existing bounds on the abundance of cold dark matter.

As mentioned above, the neutralino is mostly bino in these constructions, and for this

reason its relic density easily exceeds the recent WMAP constraint. The correct neutralino abundance is only found in those regions of the parameter space where the neutralino mass is very close to the stau mass, since then a coannihilation effect [67] takes place in the early Universe which reduces very effectively the neutralino abundance.²² After imposing the constraint on the neutralino relic density, the only regions of the parameter space which are left correspond to very narrow bands in the vicinity of the area with stau LSP. Interestingly, these favour a very narrow range of values for $\tan\beta$, which is always large. Also, as ξ_H increases, the allowed region is shifted towards larger $\tan\beta$. Thus, while $35 \lesssim \tan\beta \lesssim 40$ for $\xi_H = 1/2$ (case (I-I-I)), $45 \lesssim \tan\beta \lesssim 55$ is needed for case (I-I-A) with $\xi_H = 1$.

So far we have not commented on the effect of UFB constraints. Most of the parameter space turns out to be disfavoured on these grounds. The reason for this is the low value of the slepton masses, and more specifically, of the stau mass. Indeed, the lighter the stau, the more negative the scalar potential along the UFB-3 direction becomes, thus easily leading to a minimum deeper than the realistic (physical) vacuum. In particular, the whole $(M, \tan\beta)$ plane with $M < 1000$ GeV is disfavoured for these reason in all the cases represented in figure 4. This is consistent with what was found for other superstring and M-theory scenarios [65, 66].

In the analysis above we have only considered the negative sign for the μ parameter. If $\mu > 0$ was taken the experimental constraint on the branching ratio of $b \rightarrow s\gamma$ would become significantly more stringent, excluding regions of the parameter space with $M < 1000$ GeV for $\tan\beta \gtrsim 30$. This rules out all the regions where viable neutralino dark matter is found. Moreover, $\mu > 0$ leads to a negative SUSY contribution to the muon anomalous magnetic moment, hence being further disfavoured.

4.3 B parameter constraint on the (I-I-I)-(I-I-A) configuration

So far we have not imposed the boundary condition on the value of the B parameter at the GUT scale, which was also predicted in terms of the modular weights and related to the rest of the soft parameters by eq. (4.3). As we already explained, in this approach the REWSB condition (4.2) must be used in order to determine the value of $\tan\beta$ by means of an iterative procedure. This leaves only one free parameter, M , to describe all the soft terms and, if a solution for the REWSB equations is found, a value of $\tan\beta$ is predicted for each M .

In figure 5 (left panel) we display the value of B/M at the GUT scale as a function of $\tan\beta$ for several values of M and for the two possible choices of the sign of the μ parameter. The solutions of the REWSB equations correspond to the values of $\tan\beta$ where the different lines intersect the dotted line, which represents the boundary condition $B/M = -1$ in the $\xi_H = 1/2$ case. As we can see, solutions are found for $\mu < 0$ only when $\tan\beta$ is very large. In the $\mu > 0$ case (disfavoured by $b \rightarrow s\gamma$ limits) solutions are found for $\tan\beta \sim 4$ and $30 \lesssim \tan\beta \lesssim 40$.

²²Another possible mechanism that could help in reducing the neutralino relic density is their resonant annihilation through an s-channel mediated by a pseudoscalar Higgs. However, that is only possible when $2m_{\tilde{\chi}^0} \approx m_{A^0}$. As we saw in figure 3 the pseudoscalar is too heavy in these constructions for such a resonance to take place.

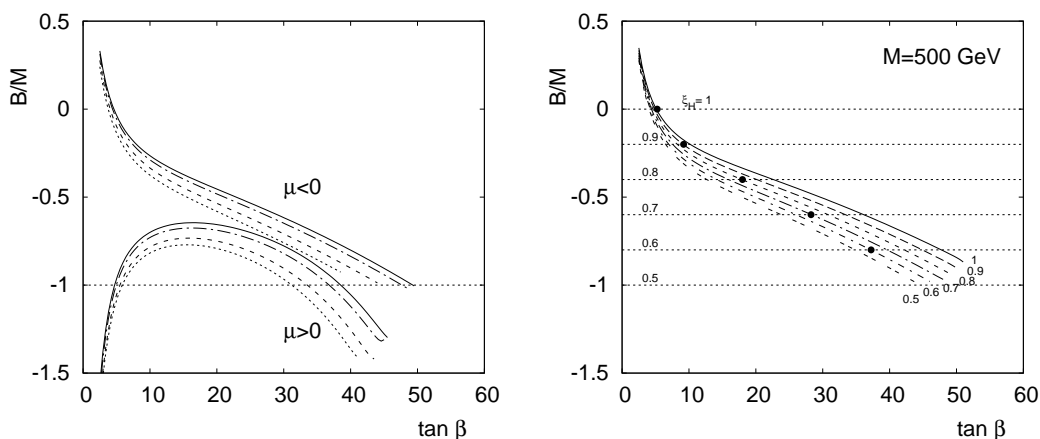


Figure 5: Left) Resulting $B(M_{\text{GUT}})/M$ as a function of $\tan\beta$ for the case with $\xi_H = 0.5$. The dotted, dashed, dot-dashed, and solid lines correspond to $M = 300, 500, 1000,$ and 1500 GeV, respectively, for both signs of the μ parameter. The boundary condition $B = -M$ is indicated with a horizontal dotted line. Right) The same, but for the cases with $\xi_H = 0.5, 0.6, 0.7, 0.8, 0.9,$ and 1 , from bottom to top, with $\mu < 0$ and $M = 1000$ GeV. The corresponding boundary conditions for B are represented with horizontal dotted lines, and the solid circles indicate the values of $\tan\beta$ consistent with these.

When the modular weight for the Higgses increases, the boundary condition for B is seriously affected. For instance, when $\xi_H = 1$ as in case (I-I-A), one obtains $B = 0$. As a consequence, the ranges of $\tan\beta$ which are solutions of the REWSB conditions change significantly. This is illustrated on the right hand-side of figure 5, where B/M at the GUT scale is represented as a function of $\tan\beta$ (with $\mu < 0$ and $M = 500$ GeV) for the cases with $\xi_H = 0.5, 0.6, 0.7, 0.8, 0.9,$ and 1 , from bottom to top, respectively. The boundary conditions corresponding to these values of the Higgs modular weights are represented by means of dotted lines, also from bottom ($\xi_H = 0.5$) to top ($\xi_H = 1$). The solutions for $\tan\beta$ for each choice of modular weight correspond to the intersection of the B/M line with the corresponding boundary condition, and are indicated with filled circles. As we see in the figure, already slightly above $\xi_H = 1/2$ (e.g. for $\xi_H = 0.52$) correct EW symmetry breaking is obtained with the predicted B -term. This happens around $\tan\beta \simeq 40$. As ξ_H increases from $1/2$ to 1 correct EW symmetry breaking is obtained at the predicted B with lower and lower values of $\tan\beta$. In the (I-I-A) limit with $\xi_H = 1$ solutions for REWSB are obtained for $\tan\beta \simeq 4$.

The resulting values for $\tan\beta$ as a function of the common scale M have been also superimposed on figure 4 by means of a thick dot-dashed line. Consistently with what we just explained, when $\xi_H = 1/2$ solutions are only found for $\tan\beta \approx 45$ and $M \gtrsim 900$ GeV, in the region with stau LSP. However, even with a slight increase in ξ_H , due to the rapid change in the boundary condition for B , solutions of the REWSB equations are found with smaller values of $\tan\beta$. If we want to obtain successful REWSB in a region consistent with appropriate neutralino dark matter, one is lead only to the region with $\xi_H \simeq 0.6$, quite

close to the boundary conditions (I-I-I) with $\xi_f = \xi_H = 1/2$. This may be seen in figure (4) (upper right) in which the thick dot-dashed line is very close to the line marking the separation between neutralino and stau LSP regions, i.e., the coannihilation region. On the other hand, in the (I-I-A) case with $\xi_H \simeq 1$ one obtains appropriate REWSB for $\tan\beta \simeq 4$, far away from the coannihilation region and hence too much dark matter is predicted. Thus insisting in getting neutralino dark matter consistent with WMAP measurements and consistent REWSB selects the region close to the (I-I-I) boundary conditions in which all MSSM matter fields live at intersecting 7-branes.

It is worth mentioning here that variations in the value of the top mass slightly alter the running of the B parameter and, as a consequence, lead to a small shift in the solutions for $\tan\beta$. We have checked that in the previous examples this shift is $\Delta \tan\beta \approx \pm 1$ when the top mass varies from $m_t = 169.2 \text{ GeV}$ to $m_t = 174.8 \text{ GeV}$ (which corresponds to a 2σ deviation from the experimental central value). Although smaller top mass also implies a more stringent constraint from the experimental bound on the lightest Higgs mass, the rest of the constraints are not significantly affected and the coannihilation region remains basically unaltered. One can therefore understand $\Delta \tan\beta$ as a small uncertainty on the trajectories for $\tan\beta$ in figure 4 to be taken into account when demanding compatibility with the regions with viable neutralino dark matter.

4.4 The bulk 7-brane (A-A- ϕ) configuration

Let us consider now the other alternative left, in which the MSSM resides at the bulk of the 7-branes. As seen in table 1, in this case the sfermion soft masses vanish at the GUT scale. This has important implications on the resulting low-energy spectrum. Although squark masses (which receive large positive contributions in the corresponding RGEs from the gluino mass parameter) easily become large enough, slepton masses remain rather light. This is particularly problematic for the lightest stau, due to the negative contribution proportional to the Yukawa in the RGEs. For this reason the stau mass-squared becomes negative even for moderate values of $\tan\beta$. In this example, this sets an upper bound of $\tan\beta \lesssim 25$.

In the remaining allowed areas of the parameter space the stau becomes the LSP. A stable charged LSP would bind to nuclei, forming exotic isotopes on which strong experimental bounds exist. Phenomenological consistency would then require such staus to decay, a possibility which arises if R-parity was broken. There is therefore no viable supersymmetric dark matter candidate in this scenario.

The resulting SUSY spectrum, together with the effect of the rest of the experimental constraints on the parameter space are shown in figure 6, clearly displaying all the above mentioned features. Interestingly, and contrary to what we observed for the intersecting 7-brane configurations, there is a region of the parameter space which satisfies the UFB constraints, corresponding to the area above the thick solid line with $M \gtrsim 500 \text{ GeV}$ and $\tan\beta \lesssim 20$. This is possible because of the increase in the Higgs mass parameters at the GUT scale (remember that now $m_{H_u, H_d}^2 = M^2$), which entails a less negative contribution to the scalar potential along the UFB directions.

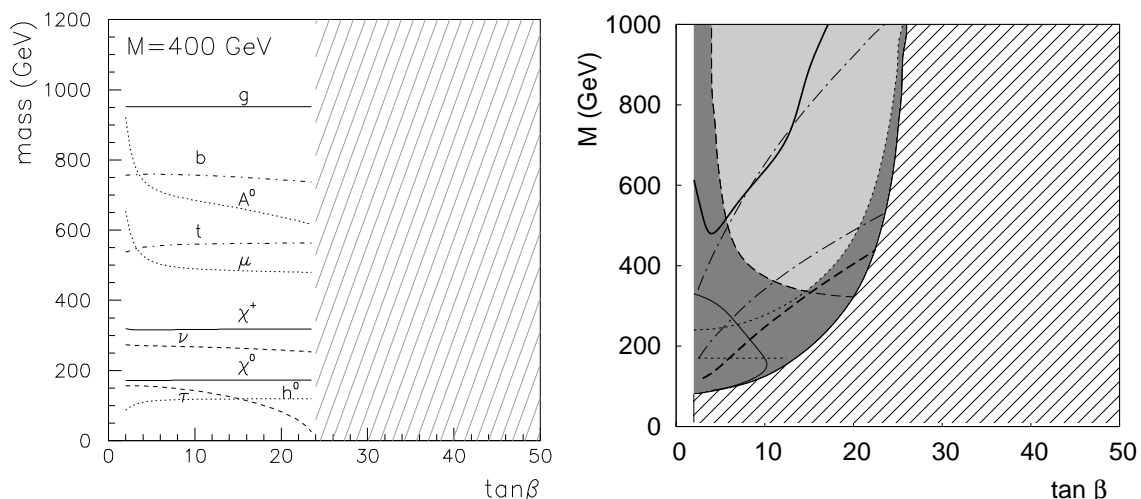


Figure 6: Left) Low-energy supersymmetric spectrum as a function of $\tan\beta$ for case (A-A- ϕ). Right) Effect of the various experimental constraints on the $(M, \tan\beta)$ plane for case (A-A- ϕ). Colour and line conventions are the same as in figure 4. In addition, the UFB constraints are fulfilled in the region above the thick solid line.

If we further impose the prediction for the B-parameter the situation is worse. Indeed in this scenario the boundary condition for B at the string scale is $B = -2M$. No solutions are found for $\tan\beta$ neither for $\mu < 0$, nor for $\mu > 0$ satisfying this boundary condition.

4.5 Effect of magnetic fluxes

We have seen how the intersecting 7-brane configuration (I-I-I)-(I-I-A) with $\xi_H \simeq 0.6$ is consistent with all constraints including appropriate amount of neutralino dark matter. This 'effective modular weight' $\xi_H \simeq 0.6$ could be understood if the Higgs field is a linear combination of fields with modular weights $\xi_H = 1/2$ (predominant) and $\xi_H = 1$.

Alternatively one could think that there could be some higher order correction which could explain the small deviation from the fully intersecting configuration (I-I-I) with $\xi_H = 1/2$. As we discussed, one possible source for such small corrections could be the magnetic fluxes which are anyway required for the spectrum to be chiral. Using the results in section (3.2) we can estimate what could be the structure of such corrections. In agreement with the unification hypothesis, we will assume that all sfermions have the same flux correction in their Kahler metrics proportional to some parameter c_f . We will then parametrize the corrections in terms of parameters defined for the different cases as:

$$\rho = \frac{(c_H - as)}{t} ; \sigma = \frac{as}{t} ; \rho_f = \frac{c_f}{t^{1/2}} ; \rho_H = \frac{c_H}{t^{1/2}}, \quad (4.4)$$

where a is defined in eq. (2.30) (in our case $a_i = a$ for gauge coupling unification). The results for the soft terms are shown in table 2.

Coupling	m_f^2	m_H^2	A	B
(A-A- ϕ)	0	$ M ^2(1 - 2\rho)$	$-M(1 - \rho)$	$-2M(1 - \rho)$
(I-I-A)	$\frac{ M ^2}{2}(1 - \frac{3}{2}\rho_f)$	0	$-M(1 - \rho_f)$	0
(I-I-I)	$\frac{ M ^2}{2}(1 - \frac{3}{2}\rho_f)$	$\frac{ M ^2}{2}(1 - \frac{3}{2}\rho_H)$	$-\frac{1}{2}M(3 - \rho_H - 2\rho_f)$	$-M(1 - \rho_H)$

Table 2: Corrections from magnetic fluxes to the different soft terms. The parameters ρ, ρ_H, ρ_f are defined in the main text.

To get the results we have assumed that $\rho_H, \rho_f \gg \sigma$, given their different large t behaviour. Looking at the table one observes as a general conclusion that the size of scalar soft terms decreases with respect to gaugino masses as fluxes are turned on. This is consistent with the known fact that as fluxes increase 7-branes localize and get boundary conditions more and more similar to 3-branes, whose matter fields are known do not get bosonic soft terms.

Note that, as we mentioned, the scalar masses of fields of A-type (modular weight $\xi = 1$) remain massless even after the addition of magnetic fluxes. So the problem of the scheme (A-A- ϕ) of having too light (or even tachyonic) right-handed sleptons cannot be solved with the addition of fluxes.

Interestingly, the inclusion of magnetic fluxes also alters the boundary condition for the bilinear parameter. In particular, for case (I-I-I), B at the GUT scale becomes less negative. This is welcome, as we already saw in the previous subsection, in order to obtain successful radiative electroweak symmetry breaking.

We have explored the effect of a small flux correction to the (I-I-I) case ($\xi_H = 1/2$) with $\rho_H = 0.1, 0.2$ and $\rho_f = 0$. The results are shown in figure 7. The low-energy supersymmetric spectrum is not very affected by the change on the Higgs mass parameters. The only visible effect is a very slight increase in the stau mass for large ρ_H as a consequence of the decrease in $|A_L|$. Hence, the allowed region with neutralino LSP and correct dark matter abundance barely changes. On the contrary, the line along which the boundary condition for B is satisfied changes significantly, being shifted towards smaller $\tan \beta$ as ρ_H increases. Compatibility with viable neutralino dark matter is found around $\rho_H \approx 0.2$ with $\xi_H = 1/2$. Thus indeed, magnetic fluxes could provide for the small correction required to get full agreement with the appropriate dark matter for the (I-I-I) intersecting 7-branes scheme.

Note that the corrections to sfermion masses, parametrized by ρ_f , imply a decrease of their mass at the GUT scale. This leads to lighter staus at the EW scale, and therefore the region where the neutralino is the LSP (and obviously the region with correct relic density) is shifted towards smaller values of $\tan \beta$. This would make it more complicated to obtain compatibility of successful REWSB and neutralino dark matter (larger ρ_H would be needed).

5. Implications for LHC

In the previous subsections we have seen how case (I-I-I) is singled out as the one in

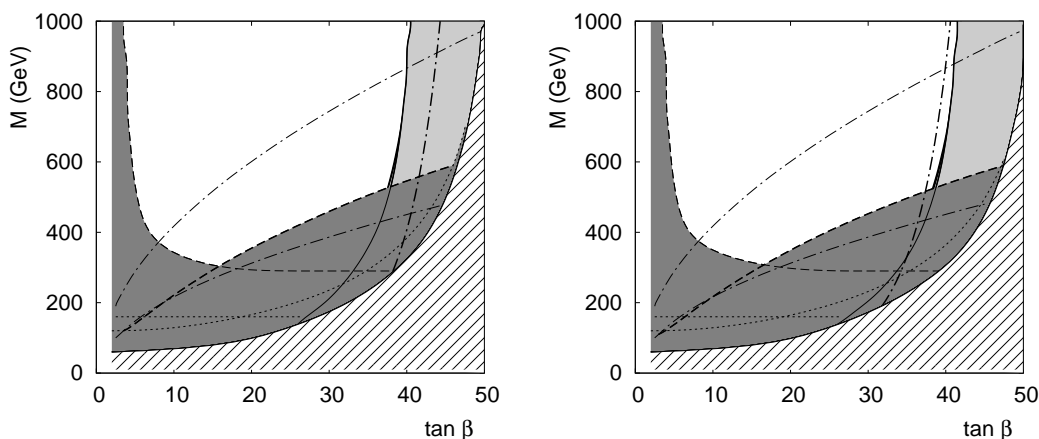


Figure 7: The same as in figure 4 but for case (I-I-I) where corrections coming from fluxes are included with $\rho_H = 0.1$ (left) and $\rho_H = 0.2$ (right). In both of them $\rho_f = 0$.

which REWSB can be reconciled with experimental and astrophysical constraints more easily, either via a small mixing of the Higgs or by the inclusion of small magnetic flux corrections. Indeed all constraints are fulfilled by soft term boundary conditions

$$\begin{aligned}
 m_f^2 &= 1/2 |M|^2, \\
 m_H^2 &\approx (1/2 - 0.1) |M|^2, \\
 A_{U,D,L} &\approx (-3/2 + 0.1) M, \\
 B &\approx -(1 - 0.2) M.
 \end{aligned}
 \tag{5.1}$$

remarkably close to the values obtained in a configuration (I-I-I) in which all quark, lepton and Higgs fields lie at intersecting 7-branes. It is then natural to wonder whether this whole class of string scenarios can be tested at the LHC and what might a typical signal look like. In order to estimate the potential LHC reach on the $(M, \tan \beta)$ planes previously shown we have followed the same approach that was used in [68] to explore the discovery potential of various SUSY scenarios.

Namely, we assume the canonical signature for SUSY searches, consisting in the observation of missing transverse energy and multijet events. These effects can be produced in the decays of squarks and/or gluinos into a final state with a pair of neutralinos plus the corresponding hadronic products. The expected sensitivity of the CMS detector to this kind of signal was explored in [69], where the projected 5σ reach contours were extracted for 1 fb^{-1} and 10 fb^{-1} in the (M, m) plane of the Constrained Minimal Supersymmetric Standard Model (CMSSM), where soft parameters are considered to be universal at the GUT scale.

The sensitivity to this kind of signal is mostly dependent on the gluino and squark masses, since these determine their production cross section. For this reason, one can express the former reach contours as functions of $m_{\tilde{g}}$ and $m_{\tilde{q}}$ and thus derive the corresponding contours in the $(M_{\tilde{g}}, m_{\tilde{q}})$ plane. In doing so, and following the procedure of [68],

we have used a numerical fit to approximate the contours of figure 13.5 of [69], then using the relations of the gluino and squark masses with the M and m parameters of the CMSSM.

This is very useful, because it allows us to estimate the potential LHC reach on scenarios beyond the CMSSM, such as the ones obtained in this work and express them in terms of reach contours on the $(M, \tan \beta)$ plane. In fact, the soft terms in eq. (5.1) do not display a large departure from the universal structure of the CMSSM and would correspond to the region with a small common scalar mass. Interestingly, these are conditions under which the resulting squarks are lighter than the gluino, which leads to large rates in the production processes of $\tilde{q}\tilde{q}$, $\tilde{q}\tilde{g}$ and $\tilde{g}\tilde{g}$ (see, e.g., [70]). Consequently, this region of the parameter space is more easily explored through searches for missing transverse energy.

The resulting reach contours show in our case no dependence with $\tan \beta$ (this is consistent with the CMSSM, for which the reach contours for this particular signal were also found to be mostly insensitive to $\tan \beta$, see e.g., [70]). Hence, the reach contours on the $(M, \tan \beta)$ plane are simply horizontal lines for a given value of the common gaugino mass, M . More specifically, the region of the parameter that could be explored by CMS with a luminosity of 1 fb^{-1} corresponds to $M \lesssim 650 \text{ GeV}$ in all the plots of figure 4. Notice that this is enough to start probing some of the regions where consistent REWSB can be obtained while fulfilling all experimental constraints and having viable neutralino dark matter. In these areas the gluino and squark masses are $m_{\tilde{g}} \lesssim 1.5 \text{ TeV}$ and $m_{\tilde{q}} \lesssim 1.3 \text{ TeV}$, corresponding also to neutralinos with a mass of $m_{\tilde{\chi}^0} \lesssim 300 \text{ GeV}$. With a luminosity of 10 fb^{-1} , LHC would be able to explore the whole region of the parameter space with $M \lesssim 900 \text{ GeV}$. This corresponds to gluinos with a mass $m_{\tilde{g}} \lesssim 2 \text{ TeV}$ and squarks with $m_{\tilde{q}} \lesssim 1.8 \text{ TeV}$. The neutralino mass is $m_{\tilde{\chi}^0} \lesssim 400 \text{ GeV}$ in this region.

Due to the constraint on the neutralino relic density, the supersymmetric spectrum is characteristic of the CMSSM in the coannihilation region, featuring a very small mass difference between the lightest neutralino and the lightest stau. A typical signal in this region would consist in looking for missing energy in the decay chain of the second-lightest neutralino into the lightest one, $\tilde{\chi}_2^0 \rightarrow \tau\tilde{\tau}_1 \rightarrow \tau\tau\tilde{\chi}_1^0$. In this case, end-point measurements at the LHC could make it possible to determine the neutralino-stau mass difference [71].

6. Final comments and conclusions

With the advent of LHC the possibility exists of SUSY particles being produced and their masses being measured. If indeed SUSY particles are found and (at least some of) their masses are measured we will have to address the issue of the origin of SUSY-breaking. We have argued in the present article that such measurements could provide important information which could rule out (or favour) large classes of MSSM-like possible string compactifications.

Modulus dominated SUSY breaking is singled out in string theory as a way to obtain SUSY-breaking vacua at the classical level with an approximately vanishing cosmological constant. In Type IIB orientifold compactifications it corresponds to the presence of certain class of SUSY breaking antisymmetric fluxes, a degree of freedom which is generically present in such compactifications. We have addressed in this paper what could be the

possible patterns of SUSY breaking masses under the assumption of modulus dominance in string theory. We have argued that this assumption together with other phenomenologically motivated ones (MSSM spectrum, existence of one large (top) Yukawa coupling and tree level gaugino masses) identify generic classes of possible string compactifications. These are Type IIB orientifold compactifications in which the MSSM particles reside either in the bulk or at intersections of 7-branes. If one further insists in a unification of coupling constants that points in the direction of F-theory 7-branes in which underlying unification structures like $SO(10)$ are possible.

In order to compute the effective action and SUSY breaking soft terms the dependence of the metrics of MSSM matter fields on the Kahler moduli are required. Those matter metrics may be described for the case of a single local Kahler modulus t in terms of 'modular weights' ξ_α . We have argued that both for simple toroidal settings as well as more complicated ones like 'swiss cheese' CY or local F-theory configurations there are three classes of matter fields I,A, ϕ with modular weights $\xi = 1/2, 1, 0$, respectively.

The possible modular weights of MSSM fields are further constrained by the condition that at least one large Yukawa coupling (that of the top) should exist. It turns out that non-vanishing trilinear Yukawa couplings in this setting only exist for combinations of type (I-I-I), (I-I-A) and (A-A- ϕ). In the first two cases MSSM fermions reside at intersections (I) of the 7-branes whereas the Higgs multiplet resides either at an intersection (I) or in the bulk (A) of a 7-brane. In the last case all MSSM particles live in the bulk of a stack of 7-branes but come from zero modes of 8-dimensional vector multiplets (A) or scalar multiplets (ϕ).

With the knowledge of the modular weights and assuming that SUSY-breaking is dominated by the auxiliary field of the local Kahler modulus t one can explicitly compute the corresponding SUSY-breaking soft terms for the three configurations, which are given in table 1. We have also estimated subleading corrections coming from magnetic fluxes in the bulk of 7-branes which are in general required to get chirality.

Taking these soft terms as boundary conditions at the string/GUT scale, we have computed the low energy SUSY spectra by numerically solving the renormalization group equations for soft terms, imposing standard radiative electroweak symmetry breaking. We have performed this computation for the three 7-brane configurations allowing for couplings (I-I-I), (I-I-A) and (A-A- ϕ). The scheme is very predictive, there are only two parameters, the gaugino mass M and $\tan\beta$ which is taken as a free parameter. Once radiative EW symmetry breaking is imposed, essentially only one free parameter is left, M , which sets the overall scale of soft terms.

We impose consistency of the resulting supersymmetric spectrum with the most recent experimental and astrophysical constraints. Namely, we include bounds on the Higgs and sparticle masses, as well as on low-energy observables such as the branching ratios of rare decays ($b \rightarrow s\gamma$, $B_s^0 \rightarrow \mu^+\mu^-$) and the supersymmetric contribution to the muon anomalous magnetic moment, $(g-2)_\mu$. Moreover, we have also studied the viability of neutralino dark matter by requiring its relic density to fulfil the WMAP bound on the abundance of cold dark matter. All these constraints essentially single out the configuration (I-I-I) in which all MSSM chiral fields reside at the intersection of 7-branes. Specifically, the modular weights

required to pass all experimental and dark matter constraints are given in eq. (5.1) and are very close to those corresponding to universal modular weight $\xi = 1/2$ for all MSSM chiral fields. We have argued that the small deviations of eq. (5.1) from those universal boundary conditions may be understood as coming from small effects, such as the presence of magnetic fluxes.

It must be emphasized that the fact that this (I-I) option passes all phenomenological tests is quite non-trivial, since, as we said, after imposing radiative electroweak symmetry breaking there is essentially only one free parameter, M . Furthermore the scheme is quite predictive and could be tested at LHC if SUSY particles are found. The overall structure of the SUSY spectrum is shown in figure 3 (left figure for $\tan\beta \simeq 35$). One of the prominent properties is that a charged stau will be only slightly heavier than the lightest neutralino. This is to be expected because appropriate dark matter here appears because of a coannihilation effect. We have shown that, with a luminosity of 1 fb^{-1} , LHC should be able to test the present scheme for $M < 600 \text{ GeV}$, whereas $M < 900 \text{ GeV}$ could be reached for 10 fb^{-1} , using the missing energy signature. A more detailed study of possible signatures at LHC of the present scheme would be quite interesting. We are looking forward to the forthcoming LHC results. If sparticles are found we might be checking not only SUSY but a large class of string theory compactifications.

Acknowledgments

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A. Alternative sources for a μ -term

In the main text it is assumed that there is an explicit μ -term, possibly induced by fluxes. This seems like the simplest possibility. However, for generality we summarize in this appendix other alternatives or modifications concerning the μ -term. As we will see these other alternatives do not seem to be particularly more promising concerning the generation of μ and B parameters consistent with radiative electroweak symmetry breaking.

An explicit μ -term from fluxes may be forbidden if the Higgs bilinear is charged under some (anomalous and massive) $U(1)$. In this case a possible source for the generation of μ - and B -terms could be some sort of string theory non-perturbative effect. As remarked in [72] string instantons may give rise to Higgs bilinear superpotentials of the form

$$M_s e^{-S_{\text{ins}}} H_u H_d = M_s e^{-\sum_i d_i T_i} H_u H_d, \quad (\text{A.1})$$

where the instanton action S_{ins} is given by a particular linear combination of Kahler moduli characteristic of the contributing instanton. As we said, in string compactifications the Higgs bilinear may be charged under some (anomalous and massive) extra U(1) and the full term is rendered gauge invariant because the particular combination of Kahler moduli in the exponent shifts under gauge transformations appropriately. Such a generation of a μ -term has the shortcoming that it does not solve the μ -problem: there is no reason why μ and the soft terms should end up being of the same order of magnitude. That problem could be solved if, instead of a superpotential μ -term, a Giudice-Masiero Z function is generated. In some heterotic orbifold compactifications [73] (and their Type I duals) bilinear terms with the structure of a Z term do actually appear with the structure

$$\frac{1}{(T_i + T_i^*)} \phi_1 \phi_2 + h.c., \tag{A.2}$$

where $\phi_{1,2}$ are chiral fields forming a vectorlike pair. That happens in Abelian heterotic compactifications with the $\phi_{1,2}$ fields corresponding to untwisted matter of a complex plane only affected by an order two twist. To our knowledge no concrete semirealistic model in which one can identify ϕ_1, ϕ_2 fields with H_u, H_d has been constructed, but still it shows that in principle such a GM term could be present at the tree level. On the other hand in some string compactifications the Higgs bilinear ($H_u H_d$) is charged under some extra (anomalous and massive) U(1) gauge symmetry so that such a tree level Z -term is forbidden perturbatively by gauge invariance. In such a case it is conceivable that string instanton effects could generate a non-perturbative Z -term with a structure

$$e^{-S_{\text{ins}}} (K_{H_u} K_{H_d})^{1/2} H_u H_d + h.c. \tag{A.3}$$

Again the gauge transformation under some extra U(1) of the Higgs bilinear would be compensated by the transformation of the instanton action S_{ins} . If such a term is present and the exponential factor does not provide too much suppression, then both a μ -term and a B -term would be generated of the same order of magnitude as the rest of the soft terms.

Let us discuss for completeness what are the results for the B -parameter obtained for these new cases. One could first consider that μ could have a non-perturbative (e.g. instanton) origin with non-perturbative dependence on the local Kahler parameter t , i.e $\mu = a e^{-bt}$ with T the local modulus and a, b constants (more generally dependent on the complex structure fields, which we assume to be fixed at their VEV). Under those circumstances the expression for B changes slightly:

$$B = -M(2 - \xi_{H_u} - \xi_{H_d}) - Mbt. \tag{A.4}$$

Note that in order to have a physical $\hat{\mu}$ -term hierarchically smaller than a string scale of order M_{GUT} one should have $bt \propto \log(M_s/\tilde{\mu}) \propto 30$. This would give rise to a B -term much larger than the rest of the soft terms and would not be viable phenomenologically. Thus this possibility does not look very promising.

As we said, another option is the presence of a Giudice-Masiero term present already at the perturbative level. We could parametrize it, in analogy with the diagonal Higgs

Kahler metric:

$$Z = \frac{t^{(1-\xi_Z)}}{t_b} ; \quad \xi_Z = \frac{(\xi_{H_u} + \xi_{H_d})}{2}. \quad (\text{A.5})$$

Such a GM term would give rise to result for the μ and B-terms:

$$\hat{\mu} = -M(1 - \xi_Z) ; \quad B_Z = -M(2 - \xi_Z). \quad (\text{A.6})$$

Note that for $\xi_Z = 1$ (corresponding to a (I-I-A) scheme with $\xi_H = 1$) no μ -term is generated so in this case the GM mechanism does not provide for the required term. On the other hand, for $\xi_Z = 1/2$, $\mu = -M/2$ is generated. In our numerical solution of the RGE we have found that the required $\hat{\mu}$ is always substantially larger than M , so that we would not get appropriate radiative symmetry breaking for $M/2$ either. All in all, a GM mechanism does not seem sufficiently flexible to be compatible with radiative symmetry breaking in the present scheme.

In the presence of magnetic fluxes one expects some corrections to these results. In particular if in analogy to the diagonal Higgs metric one assumes a Giudice-Masiero term given by $Z = (t^{1-\xi_Z} + c_Z)/t_b$ with $\xi_Z = \xi_{H_u} = \xi_{H_d} = \xi_H$ and $c_Z = c_{H_u} = c_{H_d} = c_H$ one obtains

$$\mu = -\frac{F_t}{t}(1 - \xi_H)\left(1 - \frac{c_H}{t^{1-\xi_H}}\right) ; \quad B_Z = -\frac{F_t}{t}\left(2 - \xi_H\left(1 - \frac{c_H}{t^{1-\xi_H}}\right)\right). \quad (\text{A.7})$$

Still in this case the generated $\hat{\mu}$ -term can only be smaller so that again it will not be consistent with appropriate radiative EW symmetry breaking.

One could finally consider the possible presence of a GM term as in eq. (A.5) appearing at the non-perturbative level from string instanton effects:

$$Z = (\alpha e^{-\beta t}) \frac{t^{(1-\xi_Z)}}{t_b}, \quad (\text{A.8})$$

with α, β t -independent quantities. With this ansatz one finds for the B and $\hat{\mu}$ parameters

$$\hat{\mu} = -M\alpha e^{-\beta t}(1 - \xi_Z + \beta t) ; \quad B_Z = -M \frac{(2 - \xi_Z)(\xi_Z - 1) + \beta^2 t^2}{(\xi_Z - 1) - \beta t}. \quad (\text{A.9})$$

Unlike the case of an explicit perturbative T -dependent μ -term, here βt needs not be large, and (assuming the exponential factor does not provide for a large suppression) the μ term would be naturally of order the SUSY breaking soft terms. Still, within the spirit of the present paper in which the Kahler moduli are taken to be large, one expects in general non-perturbative correction to be rather small.

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