

Soft SUSY breaking terms for chiral matter in IIB string compactifications

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ABSTRACT: This paper develops the computation of soft supersymmetry breaking terms for chiral D7 matter fields in IIB Calabi-Yau flux compactifications with stabilised moduli. We determine explicit expressions for soft terms for the single-modulus KKLT scenario and the multiple-moduli large volume scenario. In particular we use the chiral matter metrics for Calabi-Yau backgrounds recently computed in hep-th/0609180. These differ from the better understood metrics for non-chiral matter and therefore give a different structure of soft terms. The soft terms take a simple form depending explicitly on the modular weights of the corresponding matter fields. For the large-volume case we find that in the simplest D7 brane configuration, scalar masses, gaugino masses and A-terms are very similar to the dilaton-dominated scenario. Although all soft masses are suppressed by $\ln(M_P/m_{3/2})$ compared to the gravitino mass, the anomaly-mediated contributions *do not* compete, being doubly suppressed and thus subdominant to the gravity-mediated tree-level terms. Soft terms are flavour-universal to leading order in an expansion in inverse Kähler moduli. They also do not introduce extra CP violating phases to the effective action. We argue that soft term flavour universality should be a property of the large-volume compactifications, and more generally IIB flux models, in which flavour is determined by the complex structure moduli while supersymmetry is broken by the Kähler moduli. For the simplest large-volume case we run the soft terms to low energies and present some sample spectra and a basic phenomenological analysis.

KEYWORDS: Supersymmetry Phenomenology, Supersymmetry Breaking, Superstring Vacua, Superstrings and Heterotic Strings.

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1. Introduction

The imminent advent of the LHC is an excellent motivation to develop techniques to relate high energy string compactifications to observable low-energy collider physics. The LHC will be an unprecedented probe of the terascale and of the physics that stabilises the electroweak hierarchy. If supersymmetry is discovered at the LHC and the spectrum of superpartners measured, it will be necessary to connect the features of this spectrum to a more fundamental theory such as string theory. This connection is provided by the soft terms, that parametrise supersymmetry breaking and whose RG running determines the low-scale spectrum.

Supersymmetry breaking and the effective supergravity Lagrangian are among the most studied areas of string phenomenology [1, 2]. There are two principal tasks here, the first to determine how supersymmetry is broken and the second to determine how this breaking is transmitted to the matter fields. The first task primarily requires control of the moduli potential. This determines the vacuum structure of the theory. By studying the minima of the potential, it is possible to work out which fields break supersymmetry and the

magnitude of the relevant F-terms. The second task requires a knowledge of the couplings of the fields breaking supersymmetry to the matter sector. In a supergravity context, this is equivalent to knowing the modular dependence of the gauge kinetic functions and matter metrics. These are necessary to go from the moduli F-terms to soft masses for squarks and sleptons.

In recent years much progress has been made on the first problem. In the original studies of soft terms in string compactifications, it was necessary to classify supersymmetry breaking as S -, T - or U -dominated, depending on the dominant F-term, without attempting to specify the underlying microscopic source of supersymmetry breaking [2]. These studies were restricted to the heterotic string and were limited by the lack of controlled potentials in which all moduli were stabilised. In recent years much progress has been made in the field of moduli stabilisation [3–5], particularly in the context of type II compactifications. The most substantial advance is that it is now possible to compute potentials that can stabilise all the moduli in a controlled fashion. By studying the vacuum structure of these potentials, it becomes possible to determine from first principles how supersymmetry is broken and by which moduli.

Progress has also been made on the second problem. At a practical level, this requires knowledge of the gauge kinetic functions $f_a(\Phi_i)$ and the matter metrics $\tilde{K}_{\alpha\bar{\beta}}(\Phi_i)$. The former, being holomorphic, are relatively easy to compute but it is in general a hard problem to compute the latter. However both are necessary to go from a supersymmetry-breaking vacuum to soft masses and trilinear A-terms. On flat toroidal backgrounds, the Kähler metrics for chiral bifundamental matter can be computed through explicit string scattering computations [6, 7]. In Calabi-Yau backgrounds, it is possible to use dimensional reduction to compute Kähler metrics for D7 adjoint and Wilson line matter, but this does not apply to bifundamental matter. However it was recently described, in a companion paper [8] to the present one, how to compute bifundamental Kähler metrics for Calabi-Yau backgrounds. The purpose of this paper is to apply these results to the computation of soft terms.¹

For realistic models of supersymmetry breaking, it is necessary that the scale of supersymmetry breaking be hierarchically small, in order to generate a TeV-scale gravitino mass and explain the hierarchy between the weak and Planck scales. One class of models in which this occurs naturally are the large-volume models of [11, 12]. These arise in IIB flux compactifications once the leading α' correction is used to lift the no-scale structure. They fix all the moduli and are characterised by the resulting exponentially large volume \mathcal{V} . The large volume arises from the combination of α' Kähler corrections and non-perturbative superpotential corrections. The large volume lowers both the string scale m_s and the gravitino mass $m_{3/2}$,

$$m_s \sim \frac{M_P}{\sqrt{\mathcal{V}}}, \quad m_{3/2} \sim \frac{M_P}{\mathcal{V}} W_0.$$

Here M_P is the Planck mass and W_0 the flux-induced superpotential. For generic values of $W_0 \sim 1$, a volume $\mathcal{V} \sim 10^{15} l_s^6$ corresponds to TeV-scale soft terms. This gives a string

¹The soft terms for chiral matter fields in toroidal models were considered in [9, 10].

scale $m_s \sim 10^{11} \text{ GeV}$. A notable feature of these models is that the large volume occurs dynamically and there is no need to fine-tune W_0 . If we do fine-tune the flux superpotential $W_0 \ll 1$, the value of the string scale can be increased to the GUT scale. Fine-tuned flux superpotentials are characteristic of the KKLT scenario, for which $W_0 \sim 10^{-13}$ gives rise to a TeV gravitino mass and a GUT string scale.

Supersymmetry breaking and soft terms for the large volume models have been discussed in [12–14]. The first of these papers computed the overall scale of the soft terms and the second identified a logarithmic suppression in the gaugino masses compared to the gravitino mass, with the scalar masses generically not suppressed. However these results were for non-chiral D7 matter fields which are less relevant for phenomenology. Chiral matter was not considered in [14] due to the fact that the appropriate Kähler potentials were not known.

An alternative scenario of soft terms has also been put forward in [15–17]. This uses the simplest KKLT scenario with one Kähler modulus.² A suppression of all soft terms was also found in such a way that the anomaly-mediated contribution to the soft terms was of a similar order to the soft terms induced by the moduli F-terms. This gave rise to an interesting scenario that has been named mirage mediation. Having a single modulus simplifies some of the calculations and the structure of soft terms was found for all possible modular weights of the D7 matter fields. However the phenomenological analysis was performed only for the modular weights corresponding to non-chiral matter, again because those for chiral fields were unknown.

In [8] the functional dependence of the chiral matter metric on the moduli fields was recently computed for the above models, and in this paper we will use these results to further develop the computation of soft terms. We will then obtain the expressions for soft terms for chiral fields in both scenarios. These are the phenomenologically relevant expressions since squarks and sleptons come from chiral matter fields.

Both the above models involve gravity-mediated supersymmetry breaking. A common criticism of gravity-mediation is that generic supersymmetry breaking leads to non-universal squark masses and large unobserved flavour-changing neutral currents. This arises because in gravity mediation both flavour and supersymmetry breaking are UV-sensitive, and thus one would expect that the two cannot be decoupled - the physics breaking supersymmetry is also sensitive to flavour. We will consider this point and argue that it does not hold in the large-volume models, or more generally in IIB flux compactifications, as the sector that gives flavour - the complex structure moduli - and the sector that breaks supersymmetry - the Kähler moduli - are decoupled at leading order.

The structure of this paper is as follows. In section 2 we review the basics of gravity-mediated supersymmetry breaking and recall the results of [8] on bifundamental matter metrics. In section 3 we start by reviewing the large volume model and its geometry. In

²Actually the separation of these as two alternative scenarios is not really appropriate. The reason is that starting with the minimum of the scalar potential in the large volume scenario, which has a natural $\mathcal{O}(1)$ flux superpotential W_0 , it has been shown that, reducing the value of W_0 , in the limit of $W_0 \ll 1$ one reproduces the KKLT minimum that requires very small W_0 . It is instead more appropriate to talk about two limits of the same scenario.

section 3.2 we use these matter metrics to perform the computations of soft terms for diagonal matter metrics, in the limit of large cycle volume and dilute fluxes. In section 3.3 we repeat these computations for the case of non-diagonal matter metrics, obtaining identical results. In section 3.4 we give a general discussion of universality in gravity-mediated models and explain why the decoupling of Kähler and complex structure moduli naturally makes the soft terms of IIB flux models flavour-universal. We also compute the anomaly mediated contribution within supergravity and find it to be suppressed compared to the tree-level soft terms. In section 4 we perform a simple phenomenological analysis of the soft terms, carrying out the running and illustrating this with sample spectra. In section 5 we discuss the simplest KKL_T one-modulus scenario, mostly to illustrate how our results modify the previous analysis of this scenario. Finally, in section 6 we conclude.

2. Moduli SUSY breaking

2.1 Effective $N = 1$ supergravity Lagrangian

A four dimensional $N = 1$ supergravity Lagrangian is specified at two derivatives by the Kähler potential K , superpotential W and gauge kinetic function f_a . The computation of soft terms starts by expanding these as a power series in the matter fields,

$$W = \hat{W}(\Phi) + \mu(\Phi)H_1H_2 + \frac{1}{6}Y_{\alpha\beta\gamma}(\Phi)C^\alpha C^\beta C^\gamma + \dots, \quad (2.1)$$

$$K = \hat{K}(\Phi, \bar{\Phi}) + \tilde{K}_{\alpha\bar{\beta}}(\Phi, \bar{\Phi})C^\alpha C^{\bar{\beta}} + [Z(\Phi, \bar{\Phi})H_1H_2 + h.c.] + \dots, \quad (2.2)$$

$$f_a = f_a(\Phi). \quad (2.3)$$

C^α denotes a matter field and we have here, for convenience, separated the Higgs fields $H_{1,2}$ from the rest of the matter fields and specialised to the MSSM by assuming two Higgs doublets. We use Φ to denote an arbitrary modulus field and do not specify the total number of moduli.

In IIB compactifications the Kähler potential and superpotential for the moduli take the standard form [4, 18–20],

$$\hat{K}(\Phi, \bar{\Phi}) = -2 \ln \left(\mathcal{V} + \frac{\hat{\xi}}{2g_s^{3/2}} \right) - \ln \left(i \int \Omega \wedge \bar{\Omega} \right) - \ln(S + \bar{S}). \quad (2.4)$$

$$\hat{W}(\Phi) = \int G_3 \wedge \Omega + \sum_i A_i e^{-a_i T_i}. \quad (2.5)$$

\mathcal{V} is the Einstein-frame volume of the Calabi-Yau. The first term in \hat{K} is the Kähler moduli dependence, including the leading α' correction, while the second and third give the complex structure and dilaton dependence. In the superpotential, the first term is the flux-induced superpotential [18] that depends on dilaton and complex structure moduli. When evaluated at the minimum of the potential with respect to these moduli, this is denoted by W_0 . The second term is the nonperturbative superpotential responsible for fixing the Kähler moduli. Note the separate dependences of \hat{W} and \hat{K} on the complex

and Kähler moduli. This will play an important role when we subsequently discuss flavour universality.

The gauge kinetic functions $f_a(\Phi)$ depend on whether the gauge fields come from D3 or D7 branes and, in the latter case, on the 4-cycle wrapped by the D7 brane. If T_i is the Kähler modulus corresponding to a particular 4-cycle, reduction of the DBI action for an unmagnetised brane wrapped on that cycle gives

$$f_i = \frac{T_i}{2\pi}. \tag{2.6}$$

We are interested in magnetised branes wrapped on 4-cycles. The magnetic fluxes alter (2.6) to

$$f_i = h_i(F) S + \frac{T_i}{2\pi}, \tag{2.7}$$

where h_i depends on the fluxes present on that stack. This can be understood microscopically. The gauge coupling is given by an integral over the cycle Σ wrapped by the brane,

$$\frac{1}{g^2} = \int_{\Sigma} d^4y e^{-\phi} \sqrt{g + (2\pi\alpha')\mathcal{F}}. \tag{2.8}$$

Thus in the presence of flux, the gauge coupling is no longer determined simply by the cycle volume. The factors $h_a(F)$ have been explicitly computed for toroidal orientifolds [6, 9].

The presence of the term $h_a(F)$ can also be understood from the mirror IIA description. In the mirror description, IIB D7 branes become IIA D6 branes (see for example [21]). The different choices of magnetic flux on D7-branes correspond to different cycles wrapped by D6 branes. As different cycles have different sizes, it is manifest that the different branes should have different gauge couplings. The functions $h_a(F)$ are topological and so in principle may be computed even on Calabi-Yau geometries. However a sensible computation would require a geometry already yielding the Standard Model matter content - a hard and unresolved problem. Notice that (2.7) allows for non-unified gauge couplings, which are consistent with an intermediate string scale.

The computation of soft scalar masses and trilinears depends crucially on the Kähler potential for matter fields $\tilde{K}_{\bar{\alpha}\beta}(\Phi, \bar{\Phi})$ and $Z(\Phi, \bar{\Phi})$. However these are also the least understood part of the supergravity effective action, as being nonholomorphic they are not protected in the same way that the superpotential is. For D3 and D7 fields living in the adjoint of the gauge group the matter metrics have been computed in [22], and for chiral bifundamental matter fields the Kähler potential has been fully computed only in particular toroidal orientifold models [9].

This paper is a companion to [8], where it was recently described how for the geometry of the large-volume models the dependence of the bifundamental matter metric on the Kähler moduli may be computed using scaling arguments involving the physical Yukawa couplings,³

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_{\alpha}\tilde{K}_{\beta}\tilde{K}_{\gamma})^{\frac{1}{2}}}. \tag{2.9}$$

³Formula (2.9) specialises to the case of diagonal matter metrics, $\tilde{K}_{\alpha\beta} = \tilde{K}_{\alpha}\delta_{\alpha\beta}$ (no sum on right hand side), but the results hold equally well in the general, non-diagonal case.

As the combination of shift symmetry and holomorphy implies that the Kähler moduli do not appear perturbatively in $Y_{\alpha\beta\gamma}$, it is possible to extract the behaviour of the matter metrics $\tilde{K}_{\alpha\bar{\beta}}(T_i)$ through the scaling of the physical Yukawa couplings $\hat{Y}_{\alpha\beta\gamma}(T_i)$. This scaling may be computed through wavefunction overlap or direct dimensional reduction of higher-dimensional Yang-Mills.

Let us review the simplest argument of [8]. The notation used here will be more fully explained in section 3. We suppose the Standard Model is supported on a small cycle τ_s within a large bulk τ_b . This geometry is that encountered in models of branes at (resolved) singularities or the large volume models which will be reviewed in section 3. In this case locality implies the physical Yukawa couplings are determined only by the local geometry and are independent of the overall volume. As $\tilde{K} = -2 \ln \mathcal{V}$, for this to apply in (2.9) we must have

$$\tilde{K}_\alpha \sim \frac{k_{\alpha\bar{\beta}}(\phi, \tau_s)}{\mathcal{V}^{2/3}}, \tag{2.10}$$

where ϕ are complex structure moduli. Expanding $k_{\alpha\bar{\beta}}(\phi, \tau_s)$ in a power series in τ_s , we can write the resulting metric as

$$\tilde{K}_{\alpha\bar{\beta}} = \frac{\tau_s^\lambda}{\mathcal{V}^{2/3}} k_{\alpha\bar{\beta}}(\phi), \tag{2.11}$$

This expression holds in the limit of dilute fluxes and large cycle volume τ_s and will receive corrections subleading in τ_s . For the minimal model in which all branes wrap the same cycle, it was shown in [8] that $\lambda = 1/3$. For other cases λ may take values between 0 and 1. However in computations we will often start by using the more general form

$$\tilde{K}_{\alpha\bar{\beta}} = \frac{k_{\alpha\bar{\beta}}(\tau_s, \phi)}{\tau_b^p}. \tag{2.12}$$

The purpose of this is to illustrate the special cancellations that occur uniquely for the form (2.11).

For KKLТ models with a single modulus, similar arguments [8] allow us to likewise write

$$\tilde{K}_{\alpha\bar{\beta}} = \frac{k_{\alpha\bar{\beta}}(\phi)}{\tau^{2/3}}, \tag{2.13}$$

where $\tau = \text{Re}(T)$ is the size of the single 4-cycle.

We refer to [8] for the full derivations of (2.11) and (2.13), as our focus is on using these formulae to compute soft terms.

2.2 Soft breaking terms

The full $\mathcal{N} = 1$ scalar potential is

$$V = e^K \left(K^{i\bar{j}} D_i W D_{\bar{j}} \bar{W} - 3|W|^2 \right), \tag{2.14}$$

where $D_i W = \partial_i W + (\partial_i K) W$. The soft supersymmetry breaking terms are found by expanding (2.14) in powers of the matter fields using the expansions (2.1). We will give

the formulae for these below for diagonal matter metrics and in section 3.3 for nondiagonal matter metrics, but first we make some general remarks.

Gravity-mediated supersymmetry breaking is quantified through the moduli F-terms, given by

$$F^m = e^{\hat{K}/2} \hat{K}^{m\bar{n}} D_{\bar{n}} \bar{\hat{W}}. \quad (2.15)$$

We can assume without loss of generality that the superpotential \hat{W} is real. To see this, note from the Lagrangian in appendix G of [23] that a phase rotation of the superpotential can be achieved through phase redefinitions of the gauginos, chiral fermions and gravitino. In this case the gravitino mass is real while both the gaugino masses and A-terms are generically complex.

Given the F-terms, the easiest soft parameters to compute are the gaugino masses. In terms of the unnormalised field λ^a , the canonically normalised gaugino field $\hat{\lambda}^a$ is

$$\hat{\lambda}^a = (\text{Re} f_a)^{\frac{1}{2}} \lambda^a. \quad (2.16)$$

The canonically normalised gaugino masses are then given by

$$M_a = \frac{1}{2} \frac{F^m \partial_m f_a}{\text{Re} f_a}. \quad (2.17)$$

The gaugino masses that follow from (2.7) are given by

$$M_i = \frac{F^s}{2} \frac{1}{T_s + 2\pi h_a(F) S}. \quad (2.18)$$

In the limit of large cycle volume, the flux becomes diluted and the gauge coupling is determined solely by the cycle size. We shall mostly work in this dilute flux approximation, in which the gaugino masses become

$$M_i = \frac{F^s}{2\tau_s}. \quad (2.19)$$

The fractional non-universality of gaugino masses is set by the flux contribution to the gauge couplings. The quasi-universal relation (2.19) holds for the minimal geometry: if there are several small cycles involved the expressions for gaugino masses may involve several moduli and be more complicated.

For a diagonal matter metric the soft scalar Lagrangian can be written as

$$\mathcal{L}_{\text{soft}} = \tilde{K}_{\alpha} \partial_{\mu} C^{\alpha} \partial^{\mu} \bar{C}^{\bar{\alpha}} - m_{\alpha}^2 C^{\alpha} \bar{C}^{\bar{\alpha}} - \left(\frac{1}{6} A_{\alpha\beta\gamma} \hat{Y}_{\alpha\beta\gamma} C^{\alpha} C^{\beta} C^{\gamma} + B \hat{\mu} \hat{H}_1 \hat{H}_2 + h.c. \right), \quad (2.20)$$

with the scalar masses, A-terms and B-term being given by [2]

$$m_\alpha^2 = (m_{3/2}^2 + V_0) - F^{\bar{m}} F^n \partial_{\bar{m}} \partial_n \log \tilde{K}_\alpha. \quad (2.21)$$

$$A_{\alpha\beta\gamma} = F^m \left[\hat{K}_m + \partial_m \log Y_{\alpha\beta\gamma} - \partial_m \log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma) \right]. \quad (2.22)$$

$$\begin{aligned} B\hat{\mu} = & (\tilde{K}_{H_1} \tilde{K}_{H_2})^{-\frac{1}{2}} \left\{ e^{\hat{K}/2} \mu \left(F^m \left[\hat{K}_m + \partial_m \log \mu - \partial_m \log(\tilde{K}_{H_1} \tilde{K}_{H_2}) \right] - m_{3/2} \right) \right. \\ & + \left(2m_{3/2}^2 + V_0 \right) Z - m_{3/2} \bar{F}^{\bar{m}} \partial_{\bar{m}} Z + m_{3/2} F^m \left[\partial_m Z - Z \partial_m \log(\tilde{K}_{H_1} \tilde{K}_{H_2}) \right] - \\ & \left. \bar{F}^{\bar{m}} F^n \left[\partial_m \partial_n Z - (\partial_{\bar{m}} Z) \partial_n \log(\tilde{K}_{H_1} \tilde{K}_{H_2}) \right] \right\}. \quad (2.23) \end{aligned}$$

There exist similar formulae for non-diagonal matter metrics that we will present in section 3.3. The computation of soft scalar masses requires knowledge of the matter metric, which we take from (2.11) and (2.12). As above these expressions should be understood as holding in the limit of large cycle volumes and dilute fluxes. We will perform our computation of MSSM soft parameters for both diagonal and non-diagonal matter metrics, before subsequently explaining why the latter reduces to the first and how the decoupling of Kähler and complex structure moduli gives rise to soft term universality.

In order to compute the soft breaking terms - gaugino and scalar masses, A-terms and the B-term, we need a class of models for which we can compute the corresponding F-terms. We can then combine these and the expression for $\tilde{K}_{\alpha\bar{\beta}}$ into the above formulae to compute these quantities. We will first discuss the structure of soft terms in the large volume scenario which we need to introduce in detail. We later discuss the implications of our matter metric for the one-modulus KKLT scenario.

3. The large-volume model

3.1 Geometry

We start this section with a brief description of the geometry of the large-volume models. These models exist within the framework of IIB flux compactifications with D3 and D7 branes.

The dilaton and complex structure moduli are stabilised by fluxes. The Kähler moduli are stabilised by a combination of α' corrections and nonperturbative superpotentials. As shown in [11, 12], these very generally interact to produce one exponentially large cycle controlling the overall volume together with $h^{1,1} - 1$ small cycles. We denote the large and small moduli by $T_b = \tau_b + ic_b$ and $T_i = \tau_i + ic_i$ respectively, with $i = 1 \dots h^{1,1} - 1$. Consistent with this, we assume the volume can be written as

$$\mathcal{V} = \tau_b^{3/2} - h(\tau_i), \quad (3.1)$$

where h is a homogeneous function of the τ_i of degree 3/2.

The large volume lowers both the string scale and gravitino mass,

$$m_s \sim \frac{M_P}{\sqrt{\mathcal{V}}} \quad \text{and} \quad m_{3/2} \sim m_{\text{soft}} \sim \frac{M_P}{\mathcal{V}}.$$

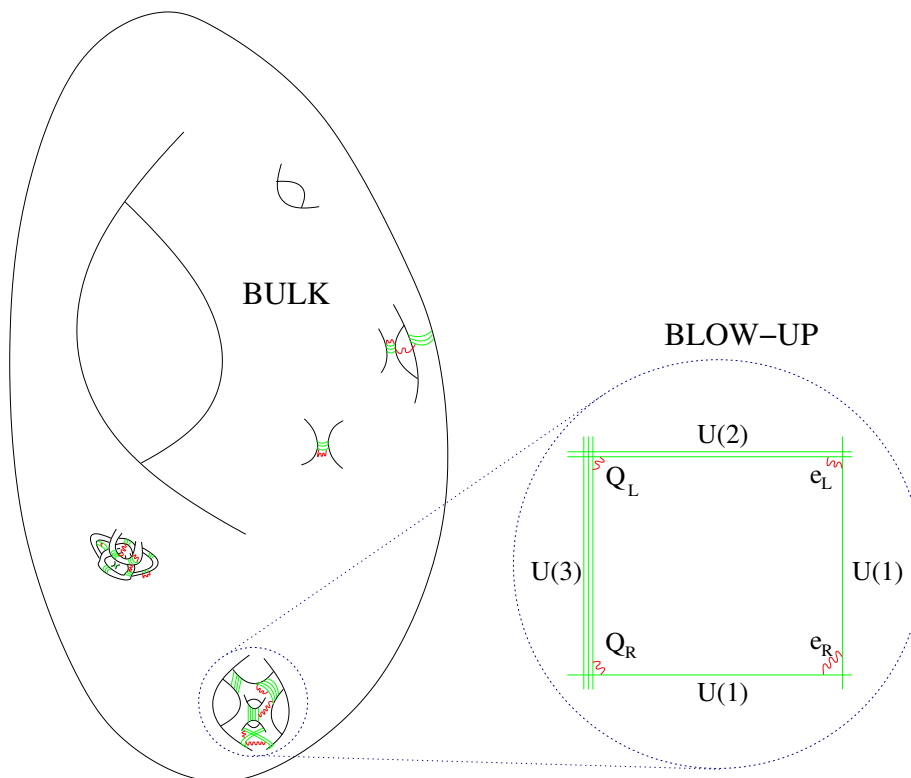


Figure 1: The physical picture: Standard Model matter is supported on a small blow-up cycle located within the bulk of a very large Calabi-Yau. The volume of the Calabi-Yau sets the gravitino mass and is responsible for the weak/Planck hierarchy.

The stabilised volume is exponentially sensitive to the stabilised string coupling, $\mathcal{V} \sim e^{\frac{c}{g_s}}$, and may thus take arbitrary values. A volume $\mathcal{V} \sim 10^{15} l_s^6 \equiv 10^{15} (2\pi\sqrt{\alpha'})^6$ is required to explain the weak/Planck hierarchy and give a TeV-scale gravitino mass. As $m_s \gg m_{3/2}$, the low-energy phenomenology is that of the MSSM and thus the computation of soft terms is central to the study of the low energy phenomenology.

The simplest model is that of $\mathbb{P}^4_{[1,1,1,6,9]}$, which we use as our working example. For this the volume can be written as [24]

$$\mathcal{V} = \frac{1}{9\sqrt{2}} \left(\tau_b^{3/2} - \tau_s^{3/2} \right). \tag{3.2}$$

τ_b and τ_s denote the big and small cycles. Moduli stabilisation for this geometry has been studied very explicitly in [11, 12, 14], and we shall not repeat this discussion here. Our prime interest here is the computation of soft terms and the analysis of universality. The geometry of the large volume models is illustrated in figure 1. The stabilised volumes of the small cycles are $\tau_s \sim \ln \mathcal{V}$. For $\mathcal{V} \sim 10^{15}$,⁴ D7-branes wrapped on such cycles have gauge couplings qualitatively similar to those of the Standard Model. If the branes are magnetised, Standard Model chiral matter can arise from strings stretching between

⁴If no dimensions are given the volume can be assumed to be measured in string units.

stacks of D7 branes. We assume the Standard Model arises from a stack of magnetised branes wrapping one (or more) of the small cycles. We do not attempt to specify the local geometry - for work in this direction, see [25].

In what we call the ‘minimal model’, there exists only one small blow-up 4-cycle on which a stack of magnetised D7 branes are wrapped. For the $\mathbb{P}^4_{[1,1,1,6,9]}$ model, this corresponds to wrapping D7 branes on the small cycle. The existence of only one small cycle need not be incompatible with the several different gauge factors and couplings of the Standard Model. The spectrum of chiral fermions depends on the magnetic flux F present on the brane worldvolume. This is quantised on 2-cycles Σ_i ,

$$\int_{\Sigma_i} F \in \mathbb{Z}. \tag{3.3}$$

If several such 2-cycles exist within the 4-cycle, different brane stacks can be realised through different choices of 2-form flux on these 2-cycles. This is consistent with there being only one small Calabi-Yau 4-cycle, as these 2-cycles may be homologically trivial within the Calabi-Yau and only non-trivial when restricted to the 4-cycle.

There is also no obstruction to the presence of chiral matter - if the cycle is a blow-up cycle, it is by construction localised in the Calabi-Yau. The branes thus cannot move off the cycle and there is no adjoint matter to give masses to strings stretching between branes. This permits a chiral spectrum, as is found in explicit models of branes at singularities. The geometry of this minimal model is shown in figure 2.

3.2 Soft term computations

We now want to use the above expression (2.11) for the matter metric in order to compute the soft terms. However, for much of the calculations we will instead use the more general expression (2.12). This will allow greater clarity of expression and will make apparent certain cancellations that only hold for the particular powers present in (2.11). We will first compute the soft terms assuming a diagonal metric and in section 3.3 generalise it for the non-diagonal case.

Several results follow simply from the large-volume geometry illustrated in figure 1. We first note that in the large-volume scenario, the F-term for the large modulus can be computed directly.

$$F^b = e^{\hat{K}/2} \hat{K}^{b\bar{n}} \left(\partial_{\bar{n}} \bar{W} + (\partial_{\bar{n}} \hat{K}) \bar{W} \right). \tag{3.4}$$

It is a property of the Kähler potential $\hat{K} = -2 \ln \mathcal{V}$ that

$$\hat{K}^{m\bar{n}} \partial_{\bar{n}} \hat{K} = -2\tau_m. \tag{3.5}$$

In deriving this it is useful to recall that, as $\hat{K} = \hat{K}(T_i + \bar{T}_i)$, any function $\Omega(K)$ satisfies

$$\partial_j \Omega(\hat{K}) \equiv \frac{\partial}{\partial T_j} \Omega(\hat{K}) = \partial_{\bar{j}} \Omega(\hat{K}) \equiv \frac{\partial}{\partial \bar{T}_j} \Omega(\hat{K}) = \frac{1}{2} \frac{\partial}{\partial \tau_j} \Omega(\hat{K}). \tag{3.6}$$

The stabilisation of the small moduli is such that $\partial_{\bar{i}} \bar{W} \sim e^{-\bar{T}_i} \sim \mathcal{V}^{-1}$, while $\hat{K}^{b\bar{i}} \sim \mathcal{V}^{2/3}$. Thus $\hat{K}^{b\bar{i}} \partial_{\bar{i}} \bar{W} \sim \hat{K}^{b\bar{i}} e^{-T_i} \sim \mathcal{V}^{-1/3}$, and so

$$F^b = -2\tau_b m_{3/2} (1 + \mathcal{O}(\mathcal{V}^{-1})). \tag{3.7}$$

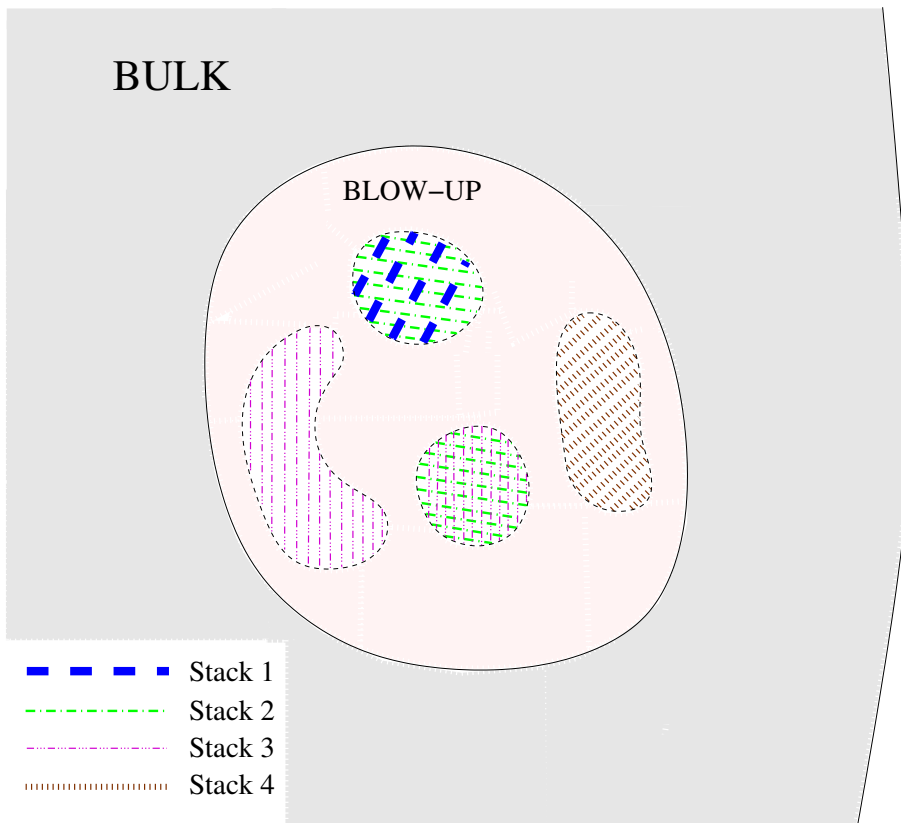


Figure 2: In the minimal geometry, there is only one small 4-cycle. The different brane stacks of the Standard Model are distinguished by having different magnetic fluxes on the internal 2-cycles of the 4-cycle. In the minimal model above, these 2-cycles are not inherited from the Calabi-Yau and only exist as cycles in the geometry of the 4-cycle. Four distinct brane stacks are required to realise the Standard Model, and we schematically show how these stacks are distinguished by different choices of magnetic flux.

This implies that any soft breaking term solely depending on F^b is naturally of order the gravitino mass $m_{3/2}$. However, it was established in [14] that the F-terms of the smaller Kähler moduli are hierarchically smaller:

$$F^i \sim 2\tau^i \frac{m_{3/2}}{\log m_{3/2}}, \tag{3.8}$$

in units of $M_P = 1$. Therefore if (as will occur) the dependence of the soft terms on F^b is cancelled, the soft terms will naturally be smaller than the gravitino mass by a factor of $\log(M_P/m_{3/2}) \sim 30 - 40$.

Using (3.7) and (2.12) the expression (2.21) for the soft scalar masses is easily seen to become

$$m_\alpha^2 = (1 - p_\alpha)m_{3/2}^2 - \bar{F}^{\bar{i}} F^j \partial_{\bar{i}} \partial_j \log k_\alpha(\tau_i, \phi), \tag{3.9}$$

where as a reminder the index i runs over small moduli only.

Now consider the A-terms (2.22). One simplification that can be made is that, at least in perturbation theory,

$$F^m \partial_m \log Y_{\alpha\beta\gamma} = 0. \tag{3.10}$$

The reason for this is that the superpotential does not depend on the Kähler moduli T_i . These have perturbative shift symmetries $T \rightarrow T + i\epsilon$ that preclude any polynomial power of T appearing in the superpotential Yukawas $Y_{\alpha\beta\gamma}$, which thus depend only on complex structure moduli.⁵ However, it is only the Kähler moduli that have non-vanishing F-terms and thus the sum in (3.10) vanishes. (2.22) therefore becomes

$$A_{\alpha\beta\gamma} = F^m \left[\hat{K}_m - \partial_m \log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma) \right]. \tag{3.11}$$

Now,

$$F^m \hat{K}_m = e^{\hat{K}/2} \hat{K}^{m\bar{n}} D_{\bar{n}} \bar{W} \hat{K}_m \tag{3.12}$$

$$= e^{\hat{K}/2} \left[\sum_n -2\tau_n (\partial_{\bar{n}} \bar{W} + (\partial_{\bar{n}} \hat{K}) \bar{W}) \right], \tag{3.13}$$

where we have used $\hat{K}^{m\bar{n}} \hat{K}_m = -2\tau_n$. As $\hat{K} = -2 \log \mathcal{V}$, $\partial_{\bar{n}} \hat{K} = -2 \frac{\partial_{\bar{n}} \mathcal{V}}{\mathcal{V}}$. \mathcal{V} is homogeneous of degree 3/2 in the τ_n , and thus

$$\sum \tau_n \partial_{\bar{n}} \mathcal{V} = \frac{3\mathcal{V}}{4}.$$

Recalling that $\partial_{\bar{n}} \bar{W} \sim \mathcal{V}^{-1}$ for the large-volume models, we get

$$F^m \hat{K}_m = 3m_{3/2} (1 + \mathcal{O}(\mathcal{V}^{-1})). \tag{3.14}$$

This gives

$$A_{\alpha\beta\gamma} = 3m_{3/2} - F^m \partial_m \log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma) \tag{3.15}$$

$$= 3m_{3/2} - F^b \partial_b \log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma) - F^i \partial_i \log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma). \tag{3.16}$$

Now,

$$\log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma) = -(p_\alpha + p_\beta + p_\gamma) \log \tau_b + \log(k_\alpha k_\beta k_\gamma(\tau_i, \phi)),$$

and so

$$\partial_b \log(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma) = -\frac{p_\alpha + p_\beta + p_\gamma}{2\tau_b}, \tag{3.17}$$

(where we have used $\frac{\partial}{\partial T_b} = \frac{1}{2} \frac{\partial}{\partial \tau_b}$). Thus we obtain

$$A_{\alpha\beta\gamma} = \left(3 - (p_\alpha + p_\beta + p_\gamma) \right) m_{3/2} - F^i \partial_i \log(k_\alpha k_\beta k_\gamma). \tag{3.18}$$

We finally want to derive an expression for the B-term (2.23). In doing so we will use the result of [8] that $\mu = 0$ - this considerably simplifies expression (2.23) as many terms

⁵We neglect any nonperturbative e^{-T} dependence; this is volume-suppressed and tiny.

vanish. We expand in full:

$$\begin{aligned}
 B\hat{\mu} &= \frac{1}{(\tilde{K}_{H_1}\tilde{K}_{H_2})^{\frac{1}{2}}} \left[(2m_{3/2}^2 + V_0)Z \right. \\
 &\quad \left. + m_{3/2} (F^m \partial_m Z - F^{\bar{m}} \partial_{\bar{m}} Z) - m_{3/2} Z F^m \partial_m \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) - \right. \\
 &\quad \left. \bar{F}^{\bar{m}} F^n \left[\partial_{\bar{m}} \partial_n Z - (\partial_{\bar{m}} Z) \partial_n \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \right] \quad (3.19)
 \end{aligned}$$

$$\begin{aligned}
 &= \frac{1}{(\tilde{K}_{H_1}\tilde{K}_{H_2})^{\frac{1}{2}}} \left[(2m_{3/2}^2 + V_0)Z + m_{3/2} (F^m \partial_m Z - F^{\bar{m}} \partial_{\bar{m}} Z) \right. \\
 &\quad - m_{3/2} Z \left[F^b \partial_b \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) + F^i \partial_i \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \\
 &\quad - \bar{F}^{\bar{b}} F^b \left[\partial_{\bar{b}} \partial_b Z - (\partial_{\bar{b}} Z) \partial_b \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \\
 &\quad - \bar{F}^{\bar{b}} F^i \left[\partial_{\bar{b}} \partial_i Z - (\partial_{\bar{b}} Z) \partial_i \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \\
 &\quad - \bar{F}^{\bar{j}} F^b \left[\partial_{\bar{j}} \partial_b Z - (\partial_{\bar{j}} Z) \partial_b \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \\
 &\quad \left. - \bar{F}^{\bar{j}} F^i \left[\partial_{\bar{j}} \partial_i Z - (\partial_{\bar{j}} Z) \partial_i \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \right], \quad (3.20)
 \end{aligned}$$

which gives

$$\begin{aligned}
 B\hat{\mu} &= \frac{\tau_b^{\frac{p_1+p_2}{2}}}{(k_{H_1}k_{H_2})^{\frac{1}{2}}} \left[(2m_{3/2}^2 + V_0)Z + \tau_b^{-p_z} m_{3/2} (F^i \partial_i z - F^{\bar{i}} \partial_{\bar{i}} z) \right. \\
 &\quad - m_{3/2} Z \left[-2\tau_b m_{3/2} \partial_b \left(-(p_1 + p_2) \log \tau_b \right) + F^i \partial_i \log(k_{H_1}k_{H_2}) \right] \\
 &\quad - (2\tau_b)^2 m_{3/2}^2 \left[\partial_{\bar{b}} \partial_b (\tau_b^{-p_z} z) - \partial_{\bar{b}} (\tau_b^{-p_z} z) \partial_b \left(-(p_1 + p_2) \log \tau_b \right) \right] \\
 &\quad + (2\tau_b) m_{3/2} F^i \left[\partial_{\bar{b}} \partial_i (\tau_b^{-p_z} z) - \partial_{\bar{b}} (\tau_b^{-p_z} z) \partial_i \log(k_{H_1}k_{H_2}) \right] \\
 &\quad + (2\tau_b) m_{3/2} \bar{F}^{\bar{j}} \left[\partial_{\bar{j}} \partial_b (\tau_b^{-p_z} z) - \partial_{\bar{j}} (\tau_b^{-p_z} z) \partial_b \log \left(-(p_1 + p_2) \log \tau_b \right) \right] \\
 &\quad \left. - \bar{F}^{\bar{j}} F^i \left[\partial_{\bar{j}} \partial_i Z - (\partial_{\bar{j}} Z) \partial_i \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \right]. \quad (3.21)
 \end{aligned}$$

This becomes

$$\begin{aligned}
 B\hat{\mu} &= \frac{\tau_b^{\frac{p_1+p_2}{2} - p_z}}{(k_{H_1}k_{H_2})^{\frac{1}{2}}} \left[\left(2 - (p_1 + p_2) + p_z(p_1 + p_2) - p_z(p_z + 1) \right) m_{3/2}^2 z + V_0 z \right. \\
 &\quad + m_{3/2} \left(F^i \partial_i z - F^{\bar{i}} \partial_{\bar{i}} z \right) - m_{3/2} z F^i \partial_i \log(k_{H_1}k_{H_2}) \\
 &\quad \left. - m_{3/2} p_z (F^i \partial_i z + \bar{F}^{\bar{j}} \partial_{\bar{j}} z) + p_z z F^i \partial_i \log(k_{H_1}k_{H_2}) + (p_1 + p_2) \bar{F}^{\bar{j}} \partial_{\bar{j}} z \right] \\
 &\quad + \frac{\tau_b^{\frac{p_1+p_2}{2}}}{(k_{H_1}k_{H_2})^{\frac{1}{2}}} \left[-\bar{F}^{\bar{j}} F^i \left[\partial_{\bar{j}} \partial_i Z - (\partial_{\bar{j}} Z) \partial_i \log(\tilde{K}_{H_1}\tilde{K}_{H_2}) \right] \right], \quad (3.22)
 \end{aligned}$$

giving

$$\begin{aligned}
 B\hat{\mu} = & \frac{\tau_b^{\frac{p_1+p_2}{2}-p_z}}{(k_{H_1}k_{H_2})^{\frac{1}{2}}} \left[\left(2 - (p_1+p_2)(1-p_z) - p_z(p_z+1) \right) m_{3/2}^2 z + V_0 z \right. \\
 & \left. + (p_z-1)zF^i\partial_i \log(k_{H_1}k_{H_2})m_{3/2} + m_{3/2} \left((1-p_z)F^i\partial_i z + (p_1+p_2-p_z-1)F^{\bar{j}}\partial_{\bar{j}} z \right) \right] \\
 & - \frac{\tau_b^{\frac{p_1+p_2}{2}-p_z}}{(k_{H_1}k_{H_2})^{\frac{1}{2}}} \left[\bar{F}^{\bar{j}}F^i \left[\partial_{\bar{j}}\partial_i z - \partial_{\bar{j}}z\partial_i \log(k_{H_1}k_{H_2}) \right] \right]. \tag{3.23}
 \end{aligned}$$

We now note that when $p_\alpha = 1$, there are cancellations in each of equations (3.9), (3.18) and (3.23). This is an independent illustration of the the critical nature of the value $p_\alpha = 1$ that appears in equation (2.11).

If we now use the correct value $p_\alpha = 1$, and also assume a vanishing cosmological constant $V_0 = 0$, we obtain

$$\hat{Y}_{\alpha\beta\gamma} = \frac{Y_{\alpha\beta\gamma}}{(k_\alpha k_\beta k_\gamma)^{\frac{1}{2}}}, \tag{3.24}$$

$$\hat{\mu} = -\frac{F^{\bar{i}}\partial_i z}{(k_{H_1}k_{H_2})^{\frac{1}{2}}}, \tag{3.25}$$

$$M_i = \frac{F^i}{2\tau_i}, \tag{3.26}$$

$$m_\alpha^2 = -\sum_{i,\bar{j}} F^i \bar{F}^{\bar{j}} \partial_i \partial_{\bar{j}} \log k_\alpha(\tau_i), \tag{3.27}$$

$$A_{\alpha\beta\gamma} = -\sum_i F^i \partial_i \log(k_\alpha k_\beta k_\gamma), \tag{3.28}$$

$$B\hat{\mu} = -\frac{1}{(k_{H_1}k_{H_2})^{\frac{1}{2}}} \left[\bar{F}^{\bar{j}}F^i \left[\partial_{\bar{j}}\partial_i z - \partial_{\bar{j}}z\partial_i \log(\tilde{k}_{H_1}\tilde{k}_{H_2}) \right] \right]. \tag{3.29}$$

We further simplify these expressions by expanding $k_\alpha(\tau_s, \phi)$ and $z(\tau_s, \phi)$ as a power series in τ_s ,

$$k_\alpha(\tau_s, \phi) = \tau_s^\lambda k_\alpha^0(\phi) + \mathcal{O}(\tau_s^{\lambda-1})k_\alpha^1(\phi) + \dots, \tag{3.30}$$

$$z(\tau_s, \phi) = \tau_s^\lambda z^0(\phi) + \mathcal{O}(\tau_s^{\lambda-1})z^1(\phi) + \dots \tag{3.31}$$

We assume k_α and z have the same scaling power λ . The soft parameters then become

$$\hat{Y}_{\alpha\beta\gamma} = \frac{Y_{\alpha\beta\gamma}}{\tau_s^{\frac{3\lambda}{2}} (k_\alpha^0 k_\beta^0 k_\gamma^0(\phi))^{\frac{1}{2}}}, \quad (3.32)$$

$$\hat{\mu} = - \left(\frac{F^{\bar{s}}}{2\tau_s} \right) \lambda \frac{z^0(\phi)}{(k_{H_1}^0 k_{H_2}^0)^{\frac{1}{2}}}, \quad (3.33)$$

$$M_i = \frac{F^s}{2\tau_s}, \quad (3.34)$$

$$m_\alpha^2 = \lambda \left(\frac{F^s}{2\tau_s} \right) \left(\frac{\bar{F}^{\bar{s}}}{2\tau_s} \right), \quad (3.35)$$

$$A_{\alpha\beta\gamma} = -3\lambda \left(\frac{F^s}{2\tau_s} \right) \quad (3.36)$$

$$B\hat{\mu} = \frac{z^0(\phi)}{(k_{H_1}^0 k_{H_2}^0(\phi))^{\frac{1}{2}}} \lambda(\lambda+1) \left(\frac{F^s}{2\tau_s} \right) \left(\frac{\bar{F}^{\bar{s}}}{2\tau_s} \right) \equiv - \left(\frac{F^s}{2\tau_s} \right) (\lambda+1)\hat{\mu}. \quad (3.37)$$

The expressions (3.32) to (3.37) are appealingly simple. Note that the supersymmetric parameters, namely $Y_{\alpha\beta\gamma}$ and μ , both require knowledge of the flavour sector through the complex structure moduli. However the soft parameters are entirely set by λ and F^s . As discussed in [8], $0 < \lambda < 1$, and for the geometry of the minimal model $\lambda = 1/3$. In this case, the pure soft parameters in the dilute flux approximation become

$$M_i = \frac{F^s}{2\tau_s}, \quad (3.38)$$

$$m_\alpha = \frac{1}{\sqrt{3}} M_i, \quad (3.39)$$

$$A_{\alpha\beta\gamma} = -M_i, \quad (3.40)$$

$$B = -\frac{4}{3} M_i. \quad (3.41)$$

It is amusing to note that the scalars, gauginos and A-terms in expressions (3.38) to (3.41) are identical to that of the dilaton-dominated scenario that was much studied in heterotic models. Notice that all soft terms are proportional to $\frac{F^s}{2\tau_s} \sim m_{3/2}/\log(M_P/m_{3/2})$ and are therefore reduced with respect to the gravitino mass.

3.3 General non-diagonal matter metrics

We now want to extend the above formulae to the case of arbitrary non-diagonal matter metrics. The expression for the normalised gaugino masses is unaltered

$$M_i = \frac{F^i}{2\tau_i}. \quad (3.42)$$

The soft scalar Lagrangian is

$$\mathcal{L}_{\text{soft}} = \tilde{K}_{\alpha\bar{\beta}} \partial_\mu C^\alpha \partial^\mu \bar{C}^{\bar{\beta}} - \tilde{m}_{\alpha\bar{\beta}}^2 C^\alpha \bar{C}^{\bar{\beta}} - \left(\frac{1}{6} A'_{\alpha\beta\gamma} C^\alpha C^\beta C^\gamma + B\hat{\mu} H_1 H_2 + c.c. \right), \quad (3.43)$$

where [2]

$$\tilde{m}_{\alpha\bar{\beta}}^2 = (m_{3/2}^2 + V_0)\tilde{K}_{\alpha\bar{\beta}} - \bar{F}^{\bar{m}}F^n \left(\partial_{\bar{m}}\partial_n\tilde{K}_{\alpha\bar{\beta}} - (\partial_{\bar{m}}\tilde{K}_{\alpha\bar{\gamma}})\tilde{K}^{\bar{\gamma}\delta}(\partial_n\tilde{K}_{\delta\bar{\beta}}) \right), \quad (3.44)$$

$$A'_{\alpha\beta\gamma} = e^{\hat{K}/2}F^m \left[\hat{K}_m Y_{\alpha\beta\gamma} + \partial_m Y_{\alpha\beta\gamma} - \left((\partial_m\tilde{K}_{\alpha\bar{\rho}})\tilde{K}^{\bar{\rho}\delta}Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right]. \quad (3.45)$$

As with the diagonal case, we can extract the overall volume dependence to write

$$\tilde{K}_{\alpha\bar{\beta}} = \tau_b^{-p}k_{\alpha\bar{\beta}}(\tau_i, \phi_i). \quad (3.46)$$

Again $p = 1$, but as above we for now leave p unspecified. We first compute the Lagrangian parameters $\tilde{m}_{\alpha\bar{\beta}}^2$ and $A'_{\alpha\beta\gamma}$, before subsequently considering the physical masses and couplings.

The expression for soft masses becomes

$$\begin{aligned} \tilde{m}_{\alpha\bar{\beta}}^2 &= (m_{3/2}^2 + V_0)\tilde{K}_{\alpha\bar{\beta}} \\ &\quad - \bar{F}^{\bar{b}}F^b \left[\partial_{\bar{b}}\partial_b \left(\tau_b^{-p}k_{\alpha\bar{\beta}} \right) - \partial_{\bar{b}} \left(\tau_b^{-p}k_{\alpha\bar{\gamma}} \right) \left(\tau_b^p k^{\bar{\gamma}\delta} \right) \partial_b \left(\tau_b^{-p}k_{\delta\bar{\beta}} \right) \right] \\ &\quad - \bar{F}^{\bar{b}}F^i \left[\partial_{\bar{b}}\partial_i \left(\tau_b^{-p}k_{\alpha\bar{\beta}} \right) - \partial_{\bar{b}} \left(\tau_b^{-p}k_{\alpha\bar{\gamma}} \right) \left(\tau_b^p k^{\bar{\gamma}\delta} \right) \partial_i \left(\tau_b^{-p}k_{\delta\bar{\beta}} \right) \right] \\ &\quad - \bar{F}^{\bar{i}}F^b \left[\partial_{\bar{i}}\partial_b \left(\tau_b^{-p}k_{\alpha\bar{\beta}} \right) - \partial_{\bar{i}} \left(\tau_b^{-p}k_{\alpha\bar{\gamma}} \right) \left(\tau_b^p k^{\bar{\gamma}\delta} \right) \partial_b \left(\tau_b^{-p}k_{\delta\bar{\beta}} \right) \right] \\ &\quad - \bar{F}^{\bar{i}}F^j \left[\partial_{\bar{i}}\partial_j \left(\tau_b^{-p}k_{\alpha\bar{\beta}} \right) - \partial_{\bar{i}} \left(\tau_b^{-p}k_{\alpha\bar{\gamma}} \right) \left(\tau_b^p k^{\bar{\gamma}\delta} \right) \partial_j \left(\tau_b^{-p}k_{\delta\bar{\beta}} \right) \right]. \end{aligned} \quad (3.47)$$

Putting $V_0 = 0$, this becomes

$$\begin{aligned} \tilde{m}_{\alpha\bar{\beta}}^2 &= m_{3/2}^2\tau_b^{-p}k_{\alpha\bar{\beta}}(\tau_i) \\ &\quad - (2\tau_b)^2 m_{3/2}^2 \left[\frac{p(p+1)}{2\tau_b} \left(\tau_b^{-p}k_{\alpha\bar{\beta}} \right) - \left(\frac{(-p)}{2\tau_b}k_{\alpha\bar{\gamma}}\tau_b^{-p}(\tau_b^p k^{\bar{\gamma}\delta})\frac{(-p)}{2\tau_b}\tau_b^{-p}k_{\delta\bar{\beta}} \right) \right] \\ &\quad - (2\tau_b)m_{3/2}F^i \left[\frac{(-p)}{2\tau_b}\tau_b^{-p}\partial_i k_{\alpha\bar{\beta}} - \left(\frac{(-p)}{2\tau_b}k_{\alpha\bar{\gamma}}\tau_b^{-p}(\tau_b^p k^{\bar{\gamma}\delta})\tau_b^{-p}\partial_i k_{\delta\bar{\beta}} \right) \right] \\ &\quad - (2\tau_b)m_{3/2}F^{\bar{j}} \left[\frac{(-p)}{2\tau_b}\tau_b^{-p}\partial_{\bar{j}} k_{\alpha\bar{\beta}} - \left(\frac{(-p)}{2\tau_b}k_{\alpha\bar{\gamma}}\tau_b^{-p}(\tau_b^p k^{\bar{\gamma}\delta})\tau_b^{-p}\partial_{\bar{j}} k_{\delta\bar{\beta}} \right) \right] \\ &\quad - F^i \bar{F}^{\bar{j}} \left[\tau_b^{-p}\partial_i\partial_{\bar{j}} k_{\alpha\bar{\beta}} - \tau_b^{-p}(\partial_i k_{\alpha\bar{\gamma}})(k^{\bar{\gamma}\delta}\tau_b^p)(\tau_b^{-p}\partial_{\bar{j}} k_{\delta\bar{\beta}}) \right]. \end{aligned} \quad (3.48)$$

Contracting indices we obtain

$$\begin{aligned} \tilde{m}_{\alpha\bar{\beta}}^2 &= m_{3/2}^2(\tau_b^{-p}k_{\alpha\bar{\beta}}) - m_{3/2}^2(p(p+1) - p^2)(k_{\alpha\bar{\beta}}\tau_b^{-p}) \\ &\quad - m_{3/2}F^i\tau_b^{-p}(-p\partial_i k_{\alpha\bar{\beta}} + p\partial_i k_{\alpha\bar{\beta}}) - m_{3/2}\bar{F}^{\bar{j}}\tau_b^{-p}(-p\partial_{\bar{j}} k_{\alpha\bar{\beta}} + p\partial_{\bar{j}} k_{\alpha\bar{\beta}}) \\ &\quad - F^i\bar{F}^{\bar{j}}\tau_b^{-p}(\partial_i\partial_{\bar{j}} k_{\alpha\bar{\beta}} - (\partial_i k_{\alpha\bar{\gamma}})k^{\bar{\gamma}\delta}(\partial_{\bar{j}} k_{\delta\bar{\beta}})). \end{aligned} \quad (3.49)$$

This simplifies to

$$\tilde{m}_{\alpha\bar{\beta}}^2 = \tau_b^{-p}m_{3/2}^2 k_{\alpha\bar{\beta}}(1-p) - F^i\bar{F}^{\bar{j}}\tau_b^{-p}(\partial_i\partial_{\bar{j}} k_{\alpha\bar{\beta}} - (\partial_i k_{\alpha\bar{\gamma}})k^{\bar{\gamma}\delta}(\partial_{\bar{j}} k_{\delta\bar{\beta}})). \quad (3.50)$$

For the critical value of $p = 1$, (3.50) simplifies to

$$m'_{\alpha\bar{\beta}} = -\frac{F^i \bar{F}^{\bar{j}}}{\tau_b} \left(\partial_i \partial_{\bar{j}} k_{\alpha\bar{\beta}} - (\partial_i k_{\alpha\bar{\gamma}}) k^{\bar{\gamma}\delta} (\partial_{\bar{j}} k_{\delta\bar{\beta}}) \right). \quad (3.51)$$

Now consider the A-terms (3.45). As before, the Kähler moduli do not appear in the superpotential and thus

$$F^m \partial_m Y_{\alpha\beta\gamma} = 0. \quad (3.52)$$

This simplifies (3.45) to

$$A'_{\alpha\beta\gamma} = e^{\hat{K}/2} F^m \left[\hat{K}_m Y_{\alpha\beta\gamma} - \left((\partial_m \tilde{K}_{\alpha\bar{\rho}}) \tilde{K}^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right]. \quad (3.53)$$

The calculation of $F^m \hat{K}_m$ depends only on the moduli fields and thus the previous result (3.14) is unaltered,

$$F^m \hat{K}_m = 3m_{3/2} (1 + \mathcal{O}(\mathcal{V}^{-1})), \quad (3.54)$$

giving

$$A'_{\alpha\beta\gamma} = e^{\hat{K}/2} \left[3m_{3/2} Y_{\alpha\beta\gamma} - F^m \left((\partial_m \tilde{K}_{\alpha\bar{\rho}}) \tilde{K}^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right]. \quad (3.55)$$

We split the sum over moduli in (3.55),

$$\begin{aligned} A'_{\alpha\beta\gamma} &= e^{\hat{K}/2} \left[3m_{3/2} Y_{\alpha\beta\gamma} \right. \\ &\quad \left. - F^b \left((\partial_b \tilde{K}_{\alpha\bar{\rho}}) \tilde{K}^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right] \\ &\quad - F^i \left((\partial_i \tilde{K}_{\alpha\bar{\rho}}) \tilde{K}^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \end{aligned} \quad (3.56)$$

$$\begin{aligned} &= e^{\hat{K}/2} \left[3m_{3/2} Y_{\alpha\beta\gamma} \right. \\ &\quad \left. - (-2\tau_b m_{3/2}) \left(\frac{(-p)}{2\tau_b} \tau_b^{-p} k_{\alpha\bar{\rho}} (\tau_b^p k^{\bar{\rho}\delta}) Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right. \\ &\quad \left. - F^i \left((\partial_i \tilde{K}_{\alpha\bar{\rho}}) \tilde{K}^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right]. \end{aligned} \quad (3.57)$$

Simplifying (3.57) we obtain

$$A_{\alpha\beta\gamma} = e^{\hat{K}/2} \left[3(1-p)m_{3/2} Y_{\alpha\beta\gamma} - F^i \left((\partial_i k_{\alpha\bar{\rho}}) k^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right]. \quad (3.58)$$

Imposing the critical value $p = 1$, we get

$$A'_{\alpha\beta\gamma} = -e^{\hat{K}/2} F^i \left((\partial_i k_{\alpha\bar{\rho}}) k^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right). \quad (3.59)$$

There is no family replication in the Higgs sector, and so for the μ and B terms the issues of non-diagonal metrics do not apply. As with the diagonal case, the unnormalised physical μ term is

$$\mu' = e^{\hat{K}/2} \mu + m_{3/2} (1 - p_z) z - (F^{\bar{i}} \partial_{\bar{i}} z) \tau_b^{-p_z}. \quad (3.60)$$

Imposing $p = 1$ and the vanishing of the superpotential μ term, we obtain

$$\mu' = -\frac{F^{\bar{i}} \partial_{\bar{i}} z}{\tau_b}. \quad (3.61)$$

The normalised μ term remains

$$\hat{\mu} = -\frac{F^{\bar{i}}\partial_i z}{(k_{H_1}k_{H_2}(\tau_i, \phi))^{\frac{1}{2}}}. \quad (3.62)$$

The B-term calculation is likewise unaltered and the general un-normalised B term can be read off from (3.23). Putting $p = 1$, we again obtain

$$B\hat{\mu} = -F^{\bar{j}}F^i [\partial_{\bar{j}}\partial_i z - (\partial_{\bar{j}}z)\partial_i \log(k_{H_1}k_{H_2})]. \quad (3.63)$$

where we have assumed $p_z = p_1 = p_2 = 1$. Finally, the unnormalised Yukawa couplings are

$$Y'_{\alpha\beta\gamma} = e^{\hat{K}/2}Y_{\alpha\beta\gamma}. \quad (3.64)$$

Gathering together the results on soft terms, we have

$$\tilde{K}_{\alpha\bar{\beta}} = \frac{k_{\alpha\bar{\beta}}(\tau_i, \phi)}{\tau_b}, \quad (3.65)$$

$$Y'_{\alpha\beta\gamma} = e^{\hat{K}/2}Y_{\alpha\beta\gamma}, \quad (3.66)$$

$$\mu' = -(F^{\bar{i}}\partial_i z)\tau_b^{-p_z}, \quad (3.67)$$

$$M_i = \frac{F^i}{2\tau_i}, \quad (3.68)$$

$$m'_{\alpha\bar{\beta}} = -\frac{F^i\bar{F}^{\bar{j}}}{\tau_b} \left(\partial_i\partial_{\bar{j}}k_{\alpha\bar{\beta}} - (\partial_i k_{\alpha\bar{\gamma}})k^{\bar{\gamma}\delta}(\partial_{\bar{j}}k_{\delta\bar{\beta}}) \right), \quad (3.69)$$

$$A'_{\alpha\beta\gamma} = -e^{\hat{K}/2}F^i \left((\partial_i k_{\alpha\bar{\rho}})k^{\bar{\rho}\delta}Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right), \quad (3.70)$$

$$B\hat{\mu} = -F^{\bar{j}}F^i [\partial_{\bar{j}}\partial_i z - (\partial_{\bar{j}}z)\partial_i \log(k_{H_1}k_{H_2})]. \quad (3.71)$$

Notice again that having imposed $p = 1$ the F-term of the large modulus F^b does not appear in these expressions and that all the soft-terms are proportional to $\frac{F^i}{2\tau_i} \sim m_{3/2}/\log(m_{3/2})$ and are therefore reduced with respect to the gravitino mass.

Following the same calculational steps as for the diagonal case we now perform a power series expansion of $k_{\alpha\bar{\beta}}$ in terms of the small modulus τ_s . We write

$$k_{\alpha\bar{\beta}}(\tau_s, \phi_j) = \tau_s^\lambda \tilde{k}_{\alpha\bar{\beta}}(\phi) + \mathcal{O}(\tau_s^{\lambda-1}) \tilde{k}_{\alpha\bar{\beta}}^1(\phi) + \dots, \\ \tilde{K}_{\alpha\bar{\beta}} = \frac{\tau_s^\lambda \tilde{k}_{\alpha\bar{\beta}}(\phi)}{\tau_b}. \quad (3.72)$$

As discussed further in [8] the subleading terms in (3.72) should be interpreted as loop corrections. For the geometry of the minimal model, $\lambda = 1/3$.

As discussed in [8] the power λ is flavour-universal. This is because flavour is defined through the superpotential and in particular through the Yukawa couplings. It is the Yukawa couplings, and only the Yukawa couplings, that distinguish the top quark from the up quark. These only depend on the complex structure moduli, and thus it is the values of the complex structure moduli - not the Kähler moduli - that distinguish the different flavours. By neglecting the complex structure moduli and focusing only on the Kähler

moduli dependence through the power λ , we are insensitive to flavour. λ will thus be universal for fields with different flavours but the same gauge charges. By contrast, the function $k_{\alpha\bar{\beta}}$ is flavour-dependent, as it is sensitive to the complex structure moduli which source flavour.

However, the form (3.72) of the matter metric in fact gives soft term universality! To see this, note that the soft terms $\tilde{m}_{\alpha\bar{\beta}}^2$ and $A'_{\alpha\beta\gamma}$ become

$$A'_{\alpha\beta\gamma} = -3\lambda \left(\frac{F^s}{2\tau_s} \right) e^{\hat{K}/2} Y_{\alpha\beta\gamma}, \tag{3.73}$$

$$\begin{aligned} m_{\alpha\bar{\beta}}'^2 &= -\frac{1}{\tau_b} \frac{F^s F^{\bar{s}}}{(2\tau_s)^2} (\lambda(\lambda - 1) - \lambda^2) k_{\alpha\bar{\beta}}, \\ &= \lambda \left(\frac{F^s F^{\bar{s}}}{(2\tau_s)^2} \right) \tilde{K}_{\alpha\bar{\beta}}. \end{aligned} \tag{3.74}$$

As the unnormalised soft scalar masses are proportional to the matter metric, the physical masses are universal and flavour-independent. Furthermore the A-terms are also proportional to the Yukawa couplings, and thus the scalar superpartners are diagonalised by the same transformation that diagonalises the Standard Model fermions. Thus the physical soft terms in the non-diagonal case are in fact identical to those found for the diagonal case in equations (3.32) to (3.37). For the minimal model with $\lambda = 1/3$, this again gives

$$M_i = \frac{F^s}{2\tau_s}, \tag{3.75}$$

$$m_\alpha = \frac{1}{\sqrt{3}} M_i, \tag{3.76}$$

$$A_{\alpha\beta\gamma} = -M_i, \tag{3.77}$$

$$B = -\frac{4}{3} M_i. \tag{3.78}$$

In the next section we discuss universality and explain on a more conceptual level why the above computations led to universal soft terms.

3.4 Universality and gravity mediation

Gravity mediation is commonly held to have a generic problem with non-universality, being expected to lead to non-diagonal squark masses and large unobserved flavour-changing neutral currents. The intuition for this is that flavour physics is Planck-scale physics, sensitive to the details of the compactification. In gravity mediation, supersymmetry breaking is also Planck scale physics, and thus we expect the physics that breaks supersymmetry to also be sensitive to flavour, thus generating non-universality. This has encouraged the development of other approaches to supersymmetry breaking, such as gauge [26] or anomaly mediation [28, 27]. These have in common the property that the supersymmetry breaking sector is decoupled or sequestered from the flavour sector.

A general pre-requisite for universality is that supersymmetry breaking does not see flavour. With this in mind, let us reconsider universality in gravity-mediated models. Gravity mediation is characterised by a hidden sector that only couples to visible matter through

non-renormalisable operators. For supersymmetry breaking to be flavour-universal, it is necessary to be able to divide the hidden sector (i.e. the moduli) into (at least) two decoupled sectors

$$\Phi = \Psi_{\text{susy-breaking}} \oplus \chi_{\text{flavour}}. \tag{3.79}$$

One sector (Ψ) is responsible for breaking supersymmetry, while the other (χ) is responsible for breaking universality and generating the flavour structure through the Yukawa couplings. As the Ψ sector breaks supersymmetry and generates soft terms, we have

$$\begin{aligned} D_{\Psi}W &\neq 0, \\ F^{\Psi} &\neq 0. \end{aligned} \tag{3.80}$$

As the χ sector breaks universality and generates flavour, in order for soft terms to be flavour-universal the χ sector must not break supersymmetry. We require

$$\begin{aligned} D_{\chi}W &= 0, \\ F^{\chi} &= 0. \end{aligned} \tag{3.81}$$

As F^{χ} contains a term $e^{K/2}K^{\chi\bar{\Psi}}D_{\bar{\Psi}}W$ and $D_{\bar{\Psi}}W \neq 0$, for these equations to be satisfied we require $K^{\chi\bar{\Psi}} = 0$. That is, the two sectors Ψ and χ must be entirely decoupled and the hidden sector metric must be block-diagonal,

$$K_{\Phi\bar{\Phi}} = \begin{pmatrix} K_{\Psi\bar{\Psi}} & 0 \\ 0 & K_{\chi\bar{\chi}} \end{pmatrix}. \tag{3.82}$$

We suppose such a decoupling can be achieved and consider, within supergravity, the dependence of the superpotential and Kähler potential on Ψ and χ . As Ψ does not see flavour, it must be absent from the superpotential Yukawa couplings,

$$Y_{\alpha\beta\gamma}(\chi, \Psi) = Y_{\alpha\beta\gamma}(\chi). \tag{3.83}$$

Furthermore, as Ψ does not see flavour the matter Kähler metric can only depend on Ψ as an overall, flavour-universal, prefactor,

$$\tilde{K}_{\alpha\bar{\beta}}(\chi, \Psi) = f(\Psi)\tilde{k}_{\alpha\bar{\beta}}(\chi). \tag{3.84}$$

One may confirm through the formulae (3.44) and (3.45) that the above metric and superpotential with the supersymmetry breaking of (3.80) and (3.81) do indeed give physical soft terms that are flavour universal.

In the context of effective field theory, the above decoupling is obviously unnatural. There is no reason for certain moduli to be excluded from the superpotential, and in any case a generic hidden sector cannot be written as a direct sum of two distinct sectors. However, this is not the case in string theory. In string compactifications the moduli can, at leading order, be separated into distinct sectors,

$$\Phi = (\text{dilaton}) \oplus (\text{complex structure moduli}) \oplus (\text{Kähler moduli}). \tag{3.85}$$

The leading order Kähler potential is⁶

$$\hat{K} = -\ln(S + \bar{S}) - \ln\left(i \int \Omega \wedge \bar{\Omega}\right) - 2\ln(\mathcal{V}), \quad (3.86)$$

and generates a block-diagonal metric,

$$\hat{K}_{\Phi\bar{\Phi}} = \begin{pmatrix} \hat{K}_{S\bar{S}} & 0 & 0 \\ 0 & \hat{K}_{U_i\bar{U}_j} & 0 \\ 0 & 0 & \hat{K}_{T_i\bar{T}_j} \end{pmatrix}. \quad (3.87)$$

Furthermore, it is well-known that brane worldvolume theories do exhibit a decoupling of Kähler and complex structure moduli. For B-type branes, the superpotential describing the theory on the brane worldvolume depends only on complex structure moduli, whereas for A-type branes the superpotential depends only on Kähler moduli [29]. The genericity arguments for flavour non-universality in gravity-mediated supersymmetry breaking therefore do not necessarily apply to string compactifications, where the moduli do separate into distinct sectors. Flavour universality can be achieved if one of these sectors - the Kähler moduli - break supersymmetry while a second sector - the complex structure moduli - source flavour. The complex structure and Kähler sectors are separate but related: the symmetry that relates the sectors but also keeps them distinct is mirror symmetry.

Our contention is that is exactly what occurs for IIB flux models. As argued above, the constraints of holomorphy and shift symmetry imply that the Kähler moduli cannot appear in the superpotential and do not see flavour. The complex structure moduli do appear in $Y_{\alpha\beta\gamma}$ and source flavour. However, it is the Kähler moduli that break supersymmetry - $D_T W \neq 0$ while $D_U W = 0$. As the metric of (3.87) is block-diagonal, this is equivalent to $F^T \neq 0$ and $F^U = 0$. The moduli controlling supersymmetry breaking are decoupled from the moduli controlling flavour, and thus the soft terms generated are naturally flavour-universal. IIB flux models, and in particular the large-volume models used above, naturally map into the general analysis above,

$$\begin{aligned} \Psi &\iff \text{Kähler moduli,} \\ \chi &\iff \text{Complex structure moduli.} \end{aligned}$$

We can now understand the calculations of sections 3.2 and 3.3 as simply the calculational illustration of this decoupling.

We also note that the above arguments also apply to the dilaton-dominated scenario in heterotic string theory, which is well known to give universality [2]. In the heterotic string the dilaton has a shift symmetry,

$$\text{Im}(S) \rightarrow \text{Im}(S) + \epsilon_i, \quad (3.88)$$

which corresponds to a shift in the universal axion and is unbroken in perturbation theory. This implies the dilaton cannot appear in the tree-level superpotential and thus cannot

⁶for IIB compactifications - very similar expressions hold for other cases.

enter the Yukawa couplings. If the dilaton dominates the supersymmetry breaking, the supersymmetry breaking sector is again decoupled from the flavour sector and soft term universality is guaranteed. In this case the correspondence is

$$\begin{aligned}\Psi &\iff \text{Dilaton,} \\ \chi &\iff \text{Complex structure moduli.}\end{aligned}$$

In the IIB context the product matter metric

$$\tilde{K}_{\alpha\bar{\beta}} = f(T + \bar{T})\tilde{k}_{\alpha\bar{\beta}}(\phi) \tag{3.89}$$

gives soft terms with exact flavour universality. We noted that we expect this product form to break down at subleading order in τ_s . As τ_s controls the gauge coupling, this expansion corresponds to loop corrections. We would naturally expect these to be flavour sensitive - loop corrections are sensitive to particle masses, and thus different flavours should feel these corrections in different ways. The subleading terms that destroy the product form will also destroy exact universality, which will only hold at leading order. Another way of stating this is to say that in the presence only of $\mathcal{N} = 1$ supersymmetry, the decoupling of the T and U moduli will hold only at leading order and will not be preserved to all orders in perturbation theory.

However these are side issues which do not affect our main point, that soft term flavour universality is naturally expected in IIB flux models: the complex structure moduli source flavour, the Kähler moduli source supersymmetry breaking, and at leading order the two do not talk to each other.

3.5 CP violation

Let us also discuss CP violation in the context of the above soft terms. The MSSM adds a set of potentially dangerous additional sources of CP violation to the Standard Model. These can be parametrised as follows [30, 31]:

$$\phi_A = \left\{ \arg \left(\frac{A_{\alpha\beta\gamma}}{Y_{\alpha\beta\gamma}} \right) \right\}, \phi_B = \{ \arg B \}, \phi_C = \{ \arg(M_a) \}. \tag{3.90}$$

The physically relevant combinations $\phi = \{ \phi_A - \phi_C, \phi_B - \phi_C \}$ are strongly constrained by the measurements of the electric dipole moment, $\phi \leq 10^{-2} - 10^{-3}$.

Let us now consider the soft masses obtained in (3.73) and (3.75)–(3.78). The phase of $A_{\alpha\beta\gamma}/Y_{\alpha\beta\gamma}$ is aligned with the phase of $F^s/(2\tau_s)$, i.e. the gaugino mass. Therefore $\phi_A - \phi_C$ vanishes. The same argument can be applied to deduce that $\phi_B - \phi_C$ also vanishes. Therefore there are no extra CP violating contributions, and the experimental constraint on the phases ϕ is satisfied in the scenario under consideration.

3.6 Anomaly-mediated contributions

In this section we compute the magnitude of anomaly-mediated contributions to gaugino masses following [32] (also see [33]). In that paper a formula is given for anomaly-mediated

contributions to gaugino masses in supergravity,

$$m_{1/2} = -\frac{g^2}{16\pi^2} \left[(3T_G - T_R)m_{3/2} - (T_G - T_R)K_i F^i - \frac{2T_R}{d_r} (\ln \det K|''_R)_i F^i \right]. \quad (3.91)$$

Here $K|''_R$ is the Kähler metric restricted to the matter fields in representation R (i.e. we only take derivatives with respect to the matter fields). d_R is the dimension of the representation and T_R is the Dynkin index of the representation. There is an implicit sum over matter fields and representations in (3.91). T_G is the Dynkin index of the adjoint representation, and $3T_G - T_R$ is the coefficient of the β -function. In minimal anomaly-mediation, the latter two terms of equation (3.91) are absent. The F^i are the normal supergravity F-terms for the moduli fields.

The dependence of the matter metric on moduli fields enters (3.91) through the third term and in particular through $\ln \det K|''_R$. As discussed earlier, we write

$$K_{C\bar{C}} = \frac{h(\tau_s, \phi) C \bar{C}}{\tau_b^p}. \quad (3.92)$$

We assume there are d_R matter fields. The metric $K|''_R$ is simply the matter metric and is a $d_R \times d_R$ matrix, with

$$(K|''_R)_{i\bar{j}} = \frac{h_{i\bar{j}}(\tau_s, \phi)}{\tau_b^p}, \quad (3.93)$$

where

$$h_{i\bar{j}}(\tau_s, \phi) = \tau_s^\lambda X_{i\bar{j}}(\phi) + \tau_s^{\lambda-1} X'_{i\bar{j}}(\phi) + \dots \quad (3.94)$$

Then

$$(\det K|''_R) \sim \frac{\tau_s^{\lambda d_R}}{\tau_b^{p d_R}} X(\phi), \quad (3.95)$$

where $X(\phi)$ is a complicated function of complex structure moduli. These expressions hold so long as the matter fields have vanishing vevs. We then have

$$\ln \det K|''_R = \lambda d_R \ln \tau_s - p d_R \ln \tau_b + \dots, \quad (3.96)$$

and

$$\frac{1}{d_R} (\ln \det K|''_R)_i F^i = \frac{\lambda F^s}{2\tau_s} - \frac{p F^b}{2\tau_b}. \quad (3.97)$$

We also recall that

$$K_i F^i = -\frac{3F^b}{2\tau_b} (1 + \mathcal{O}(\mathcal{V}^{-1})) = 3m_{3/2} (1 + \mathcal{O}(\mathcal{V}^{-1})). \quad (3.98)$$

We can then evaluate the anomaly-mediated contribution to gaugino masses to obtain

$$m_{1/2} = -\frac{g^2}{16\pi^2} \left[2T_R(1-p)m_{3/2} - 2\lambda T_R \left(\frac{F^s}{2\tau_s} \right) \right]. \quad (3.99)$$

This expression has a similar form to those previously encountered in our analysis of soft terms: there is a dominant part, which cancels for the critical value $p = 1$ and a suppressed

part depending on the F-term of the small modulus (we recall this is naturally suppressed by a factor 30 compared to the gravitino mass). For the actual values $p = 1$ and $\lambda = 1/3$, we obtain

$$m_{1/2,AMSB} = \frac{g^2}{16\pi^2} \frac{2}{3} T_R \left(\frac{F^s}{2\tau_s} \right). \quad (3.100)$$

Because of the no-scale $p = 1$ cancellation, this is suppressed compared to the tree-level term by the loop factor $\frac{g^2}{16\pi^2}$ and is thus negligible.

Thus, while we may have expected a competition of modulus and anomaly mediation due to the suppressed tree-level gaugino masses, this does not in fact occur. The anomaly-mediated contributions are doubly suppressed: once by the loop factor, and once by the same suppression as occurred at tree-level. The root of this double suppression lies in the no-scale structure, which gives rise to the $p = 1$ cancellation.

Similar behaviour will occur for anomaly-mediated contributions to scalar masses. For a no-scale model, these cancel as for anomaly-mediated contributions to gaugino masses [27, 33]. In our models, the no-scale structure is broken by the F-terms for the small moduli, and these will give rise to anomaly-mediated contributions as in (3.100). However, as in (3.100) such contributions will be doubly suppressed, once by the loop factor and once by the small modulus F-term.⁷

The upshot of these discussions is that, owing to the pseudo-no scale structure, the anomaly-mediated contributions are suppressed by a loop factor compared to the tree level soft terms and not compared to the gravitino mass. This implies that they can be consistently neglected in our analysis of soft terms.

3.7 Comparison with previous results

In [14] a preliminary calculation of scalar masses was performed without knowledge of the matter metric. In particular, the power of the large modulus τ_b in (2.12) was left undetermined. Denoting this power by p , it was there found that

$$M_i = \frac{F^s}{2\tau_s} \sim \frac{m_{3/2}}{\ln m_{3/2}}, \quad (3.101)$$

$$m_i^2 = (1 - p) m_{3/2}^2 + \frac{F^s \bar{F}^s}{(2\tau_s)^2}. \quad (3.102)$$

Since the power p was not known for chiral matter, reference [14] considered the generic value of $p \neq 1$. In this case scalar masses should be comparable to the gravitino mass and thus significantly larger than the (logarithmically suppressed) gaugino masses. This was used for considering universality of scalar masses since scalar masses could be written as

$$m_i^2 = [(1 - p) + \epsilon_i^2] m_{3/2}^2, \quad (3.103)$$

with $\epsilon_i \sim 1/\log(M_P/m_{3/2}) \sim 1/30$ being responsible for the breaking of universality. This was understood by arguing that if the Standard Model lives on a small cycle, and the

⁷This behaviour should also be present in other configurations such as for D3 branes in KKLT models that, to leading order, have a no-scale structure and cancellations of soft terms will happen for both the standard gravity mediation and anomaly mediation.

stronger source of SUSY breaking is the F-term associated to the size of the large cycle, then locality would imply that the source of non-universality could be associated to the suppressed value of the F-term corresponding to the small cycle, that is suppressed.

However the results of [8] show that for bifundamental matter actually $p = 1$. While this does not obviate any of the calculations of [14], it does change the interpretation, as for $p = 1$ the scalar masses are comparable to the gaugino masses rather than logarithmically larger. For matter with $p \neq 1$, the scalar masses would be logarithmically larger than the gaugino masses. An example of such matter would be D7 adjoint matter. It would be interesting to find other examples of matter giving $p \neq 1$.

This modifies the discussion of universality. We however have been able to prove in sections 3.3 and 3.4 that there is still leading-order flavour universality even for the second term in equation (3.103). This is because the source of non-universality is the flavour mixing in the Yukawa couplings that have no perturbative dependence on any of the Kähler moduli, large or small.

4. RG running and sample spectra

We now turn to a phenomenological analysis of the large volume models with soft terms as described in § 3.2. Since the soft terms are computed at the string scale, to obtain the low energy spectra they need to be evolved to the weak scale using the RG equations of the MSSM. For this purpose we use the program `SoftSUSY2.0` [34]. By doing so we are making the strong theoretical assumption that the only charged matter content below the string scale is the MSSM.⁸

The resulting low energy spectra are then subjected to various experimental constraints. These are the size of the $b \rightarrow s\gamma$ branching ratio, the anomalous magnetic moment of the muon $(g - 2)_\mu$ and the dark matter relic density Ωh^2 .

The average measurement of the $b \rightarrow s\gamma$ branching ratio is obtained from [36]: $BR(b \rightarrow s\gamma) = (3.55 \pm 0.26) \times 10^{-4}$. The theoretical uncertainty in this result is [37] 0.30×10^{-4} ; adding the two errors in quadrature, we obtain the 1σ bound, $BR(b \rightarrow s\gamma) = (3.55 \pm 0.40) \times 10^{-4}$.

We also impose limits on the new physics contribution to $a_\mu = (g_\mu - 2)/2$. The average experimental value of a_μ is 11659208×10^{-10} [38]. Following [39], we use the Standard Model computation based on data from $e^+e^- \rightarrow$ hadrons, which predicts 11659186×10^{-10} , rather than that based on tau leptons. Combining the two results (with respective uncertainties) yields the 1σ bound on the non-SM contribution to a_μ , $\delta a_\mu = (22.2 \pm 10.2) \times 10^{10}$. It is important to note that δa_μ usually has the same sign as μ , so $\mu > 0$ is preferred by experiment.

We also impose bounds on the Higgs mass obtained by the LEP2 collaborations [40]. The lower bound is 114.4 GeV at the 95% CL. The theoretical computation of the Higgs

⁸The dilaton-dominated scenario with an intermediate string scale was also investigated in [35]. In addition to the improved software and constraints since [35], this work also differs from [35] in that the soft terms are derived after a process of moduli stabilisation rather than assumed *ad hoc*.

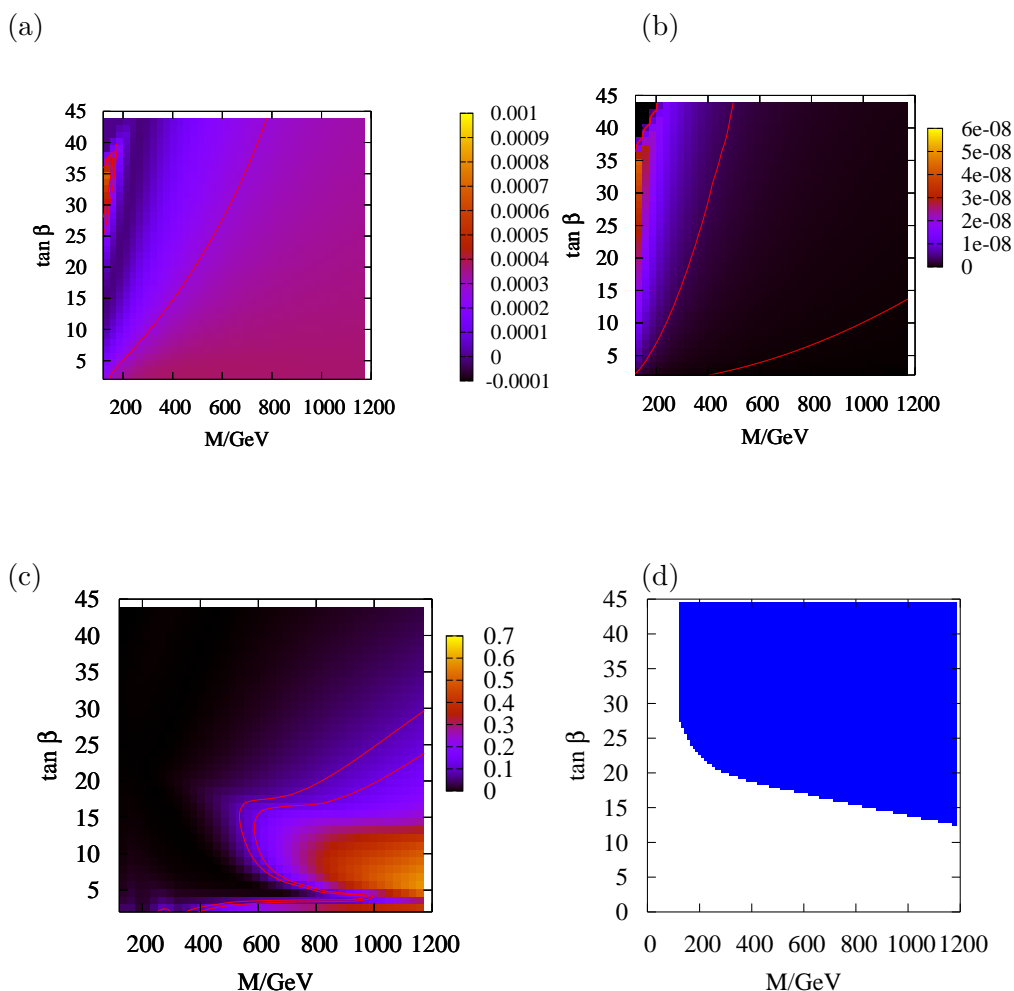


Figure 3: Contour plots of (a) $BR(b \rightarrow s\gamma)$ (b) δa_μ and (c) Ωh^2 on $\tan \beta$ and M , for $\mu > 0$ and $m_s \sim 10^{11}$ GeV. 2σ bounds are used. The blue region in (d) denotes that the LSP is stau, the white that it is neutralino.

mass in supersymmetric scenarios is subject to an estimated error of 3GeV, and so we impose $m_h > 111\text{GeV}$ on the result obtained from `SoftSUSY`.

The WMAP [41] constraint on the relic density of dark matter particles (at the 2σ level) is

$$0.085 < \Omega h^2 < 0.125. \tag{4.1}$$

We compute the neutralino contribution to dark matter using `micrOmegas1.3` [42]. The lower bound in (4.1) is only applicable if we require that the neutralino is the only contribution to dark matter. Since there may be other dark matter constituents, such as axions or for the large volume models the volume modulus, the upper bound on Ωh^2 is the most relevant.

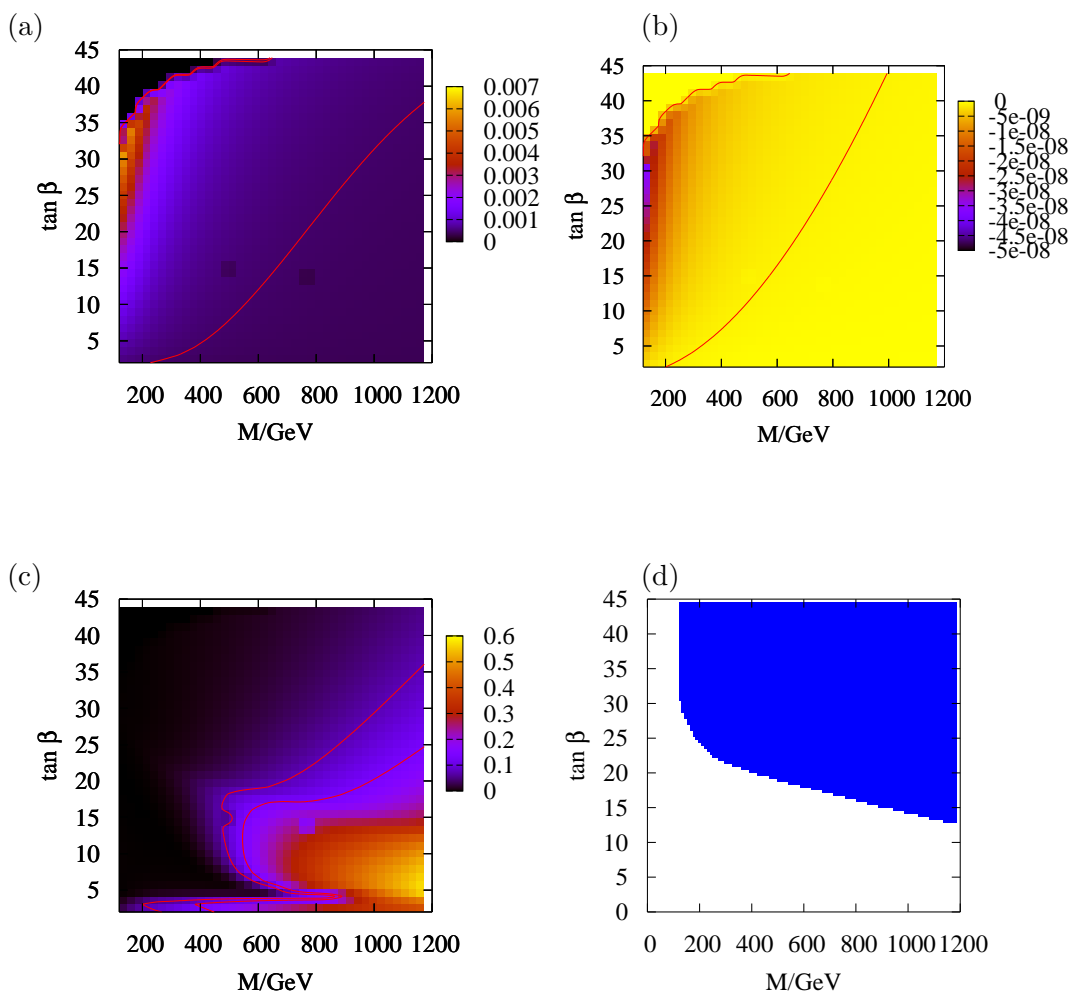


Figure 4: Contour plots of (a) $BR(b \rightarrow s\gamma)$ (b) δa_μ and (c) Ωh^2 on $\tan\beta$ and M for $\mu < 0$ and $m_s \sim 10^{11}$ GeV. Note that 3σ bounds are used. The blue region in (d) denotes that the LSP is stau, the white that it is neutralino.

The values of the Standard Model input parameters used in our computations are the current central values: $m_t = 171.4\text{GeV}$ [43], $m_b(m_b)^{\overline{MS}} = 4.25\text{GeV}$, $\alpha_s(M_Z) = 0.1187$, $\alpha^{-1}(M_Z)^{\overline{MS}} = 127.918$, $M_Z = 91.1187\text{GeV}$ [39].

We first neglect the constraint on the B soft term and scan over various values for the gaugino mass M and $\tan\beta$. The `micrOmegas1.3` package, interfaced with `SoftSUSY2.0` via the SUSY Les Houches Accord [44], is used to compute the $b \rightarrow s\gamma$ branching ratio, δa_μ and Ωh^2 . The constraints are presented on separate plots in figures 3 and 4.

Figures 3, 4 also show regions in the $M - \tan\beta$ plane with stau LSP. The existence of a stable charged LSP is experimentally ruled out [45] and therefore if such a scenario is to be considered, it must include R parity violation to allow the LSPs to decay before

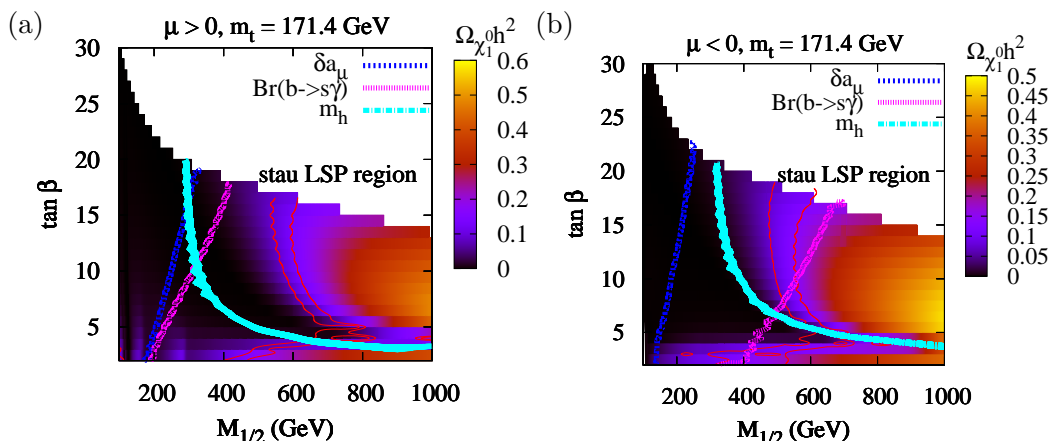


Figure 5: Experimental constraints on the parameter space: If a neutralino LSP is to account for the CDM part of the Universe, then the allowed regions $0.085 < \Omega_{\tilde{\chi}_1^0} h^2 < 0.125$ for (left) $\mu > 0$ and $0.075 < \Omega_{\tilde{\chi}_1^0} h^2 < 0.135$ for (right) $\mu < 0$ are bounded by the red contours. The bounds are at 2σ and 3σ respectively. The white region represents non neutralino (stau) LSP points. The blue and pink curves respectively are the δa_μ and $BR(b \rightarrow s\gamma)$ bounds. The cyan curve is the 111 GeV bound on the Higgs mass.

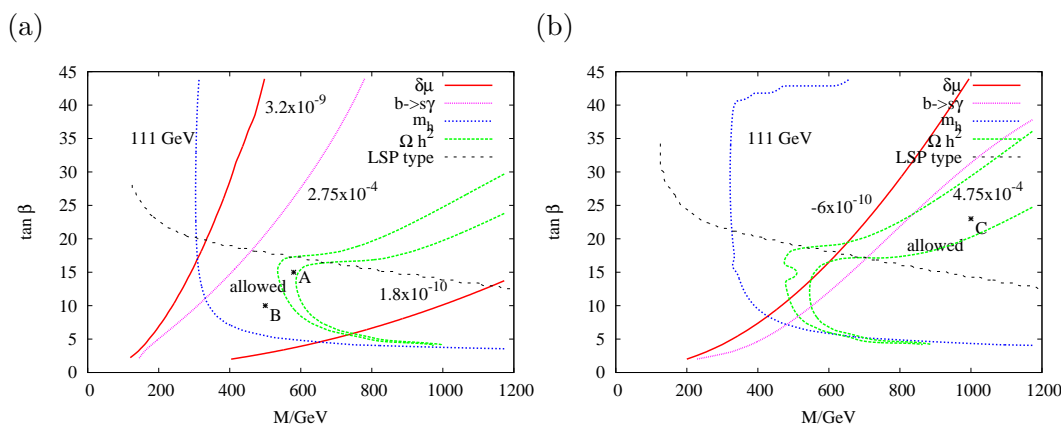


Figure 6: Contour plots of δa_μ , $BR(b \rightarrow s\gamma)$ and Ωh^2 on $\tan \beta$ and M for (a) $\mu > 0$ and (b) $\mu < 0$ and $m_s \sim 10^{11}$ GeV. 2σ bounds are used for $\mu > 0$ while 3σ ones are used for $\mu < 0$.

nucleosynthesis.

We also scanned the $M - \tan \beta$ parameter space for potentially dangerous charge and colour breaking (CCB) minima (see e.g. [1, 48]). This was done by investigating the unbounded from below UFB3 direction in field space. As described in [49], UFB3 is the most restrictive of the UFB and CCB constraints. Other less restrictive CCB minima were

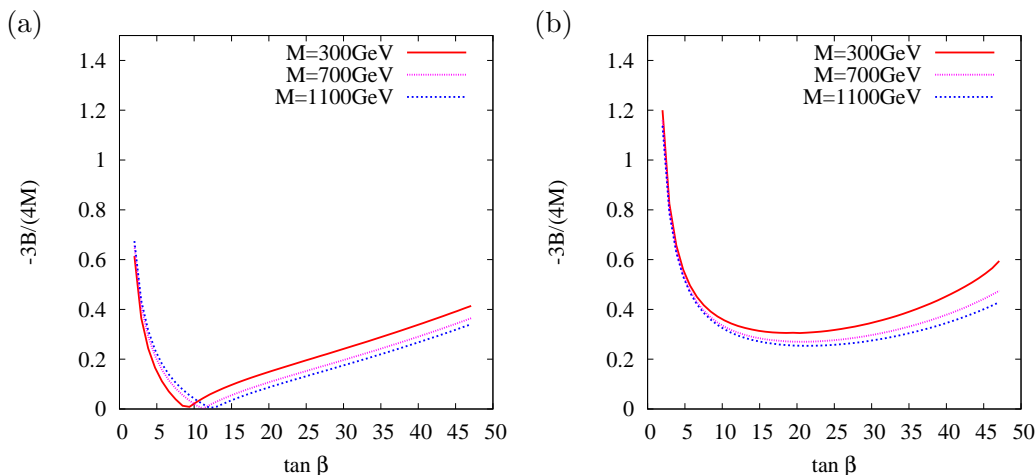


Figure 7: The ratio $B/(-4M/3)$ over a range of values of $\tan \beta$ for (a) $\mu > 0$ and (b) $\mu < 0$ and $m_s \sim 10^{11}$ GeV.

also investigated. In the region of interest we found no such minima, for both the $\mu > 0$ and $\mu < 0$ cases.

The combined results involving all the experimental constraints and displaying stau LSP regions are shown in figures 5, 6.

We can see that in the $\mu > 0$ regime there is a very restricted throat-like region in the $M - \tan \beta$ plane which satisfies all the experimental constraints (including dark matter) and has a neutralino LSP. If the lower bound on Ωh^2 is not imposed the region is much larger. In the $\mu < 0$ regime the experimentally allowed region is entirely stau LSP, so one is forced to consider R-parity violating models (for a recent discussion see [46]).

Table 1 shows a few sample points, one with $\mu > 0$ from the allowed region in figures 5 and 6 and one with $\mu < 0$ with stau LSP, which pass all the experimental constraints we imposed. The corresponding spectra computed using `SoftSUSY` are also shown. The spectra B and C are also pictorially represented in figures 8 and 9.

We next attempt to also impose the high scale constraint $B = -4M/3$. This turns out to be impossible to satisfy for $\mu > 0$, as figure 7 indicates. For $\mu < 0$ it can only be satisfied in the very low $\tan \beta$ region. However, for those values of $\tan \beta$ the Higgs mass is below the LEP bound.

Therefore we are forced to assume that there is some mechanism at work responsible for generating the correct B -term, e.g. a low energy coupling of the form $\alpha N H_1 H_2$ with a gauge invariant scalar N which obtains an appropriate vev.

One necessary caveat on all the above results is that they are all obtained in the dilute flux approximation, where we have neglected the contribution of magnetic fluxes to the gauge couplings. A further caveat is that there may be additional vector-like matter at the TeV scale that will alter the running of the soft terms (for explicit string models where

this effect has been studied see [47]). The results are therefore necessarily approximate and there is a certain amount of ‘theoretical error’ in the soft terms used.⁹ In particular, we have used universal gaugino masses even though the Standard Model gauge couplings are not unified at the intermediate scale. In the large volume model, we interpret this lack of unification as due to the magnetic fluxes. These fluxes will also modify the matter metrics, and in general contribute corrections to all the soft terms considered here. The magnitude of these corrections can be roughly estimated by the fractional non-universality of the SU(2) and SU(3) gauge couplings at the intermediate scale. The above results should all be understood as leading order results in the dilute flux approximation. It would be interesting to perform a careful study of the extent to which the inclusion of magnetic fluxes would alter the results in this section.

5. The simplest KKLT scenario

Let us now briefly recall the one-modulus KKLT scenario. In this case there is only one T field and the volume is given by $\mathcal{V} = (T + T^*)^{3/2}$. Here the flux superpotential is required to be very small, $W_0 \sim 10^{-13}$, in order to address the hierarchy problem. The minimum of the scalar potential is supersymmetric AdS and supersymmetry is only broken by the lifting term, which unfortunately is the least understood ingredient of the KKLT scenario. This has an important effect on the structure of soft breaking terms, leading to the overall suppression by the factor $\ln(M_P/m_{3/2})$. Soft terms have been computed for general values of the modular weight in the matter Kähler potential. However the phenomenological analysis has mainly been performed for vanishing modular weight which is standard for non-chiral D7 adjoint matter.

The expression for soft masses found in [15] takes the form:

$$\begin{aligned}
 m_\lambda^2 = & (1 - n_\lambda) \left| \frac{F^T}{(T + T^*)} \right|^2 - \frac{1}{32\pi^2} \frac{d\gamma_\lambda}{d \log \mu} \left| \frac{F^\Phi}{\Phi_0} \right|^2 + \\
 & \left(\frac{1}{8} \sum_{\mu\nu} (3 - n_\lambda - n_\mu - n_\nu) |Y_{\lambda\mu\nu}|^2 - \frac{1}{2} \sum_a l_a T_a (C^\lambda) g_a^2 \right) \cdot \\
 & \left(\frac{F^T}{(T + T^*)} \left(\frac{F^{*\Phi}}{4\pi^2 \Phi_0^*} \right) + \frac{F^{*T}}{(T + T^*)} \left(\frac{F^\Phi}{4\pi^2 \Phi_0} \right) \right). \tag{5.1}
 \end{aligned}$$

Here Φ is the conformal compensator field, b_a are the one-loop beta function coefficients, g_a the gauge couplings, γ_λ are the anomalous dimensions of the matter fields C^λ , and T_a are group theory factors. Finally, n_λ is the power of $T + T^*$ in the expansion of the Kähler potential in matter fields,

$$K = -3 \log(T + T^*) + \frac{C^\lambda C^{*\lambda}}{(T + T^*)^{n_\lambda}}. \tag{5.2}$$

⁹In order to estimate the changes that our assumptions may have, we also computed the spectrum for a model with the spectrum of the MSSM plus extra doublets such that gauge unification is achieved at the intermediate scale. The spectrum takes a similar structure as the one shown here, with the scalar masses increasing by 10 – 20 GeV, and the lightest neutralino becoming roughly 40 GeV lighter; this has the interesting consequence of expanding the neutralino LSP region in figure 5.

	A	B	C
m_s	10^{11}	10^{11}	10^{11}
$\tan \beta$	15	10	23
M	580	500	1000
$\text{sgn}\mu$	+	+	-
$\tilde{e}_L, \tilde{\mu}_L$	464	401	792
$\tilde{e}_R, \tilde{\mu}_R$	386	333	661
$\tilde{\tau}_L$	463	402	779
$\tilde{\tau}_R$	369	326	618
\tilde{u}_1, \tilde{c}_1	924	806	1527
\tilde{u}_2, \tilde{c}_2	951	829	1580
\tilde{t}_1	679	582	1166
\tilde{t}_2	958	815	1448
\tilde{d}_1, \tilde{s}_1	915	798	1512
\tilde{d}_2, \tilde{s}_2	958	835	1585
\tilde{b}_1	859	752	1405
\tilde{b}_2	903	792	1455
χ_1^0	364	311	643
χ_2^0	469	400	822
χ_3^0	541	479	862
χ_4^0	587	524	927
χ_1^\pm	467	397	821
χ_2^\pm	584	521	924
A_0, H_0	679	610	1042
H^\pm	684	614	1046
\tilde{g}	1048	913	1745
$\tilde{\nu}_{1,2}$	456	392	789
$\tilde{\nu}_3$	451	390	771
h	116	114	118
$B(b \rightarrow s\gamma)/10^{-4}$	3.3	3.4	4.42
$\delta a_\mu/10^{-10}$	7.9	7.0	-4.3
Ωh^2	0.12	0.01	—

Table 1: Sparticle spectra for two sample points with $\mu > 0$ and $\mu < 0$. All masses are in GeV.

The F-terms can be written as

$$F^\Phi \sim 16\pi^2 M_s, \quad F^T/(T + T^*) = \alpha M_s.$$

The parameter α depends on the shape of the potential used to lift the supersymmetric

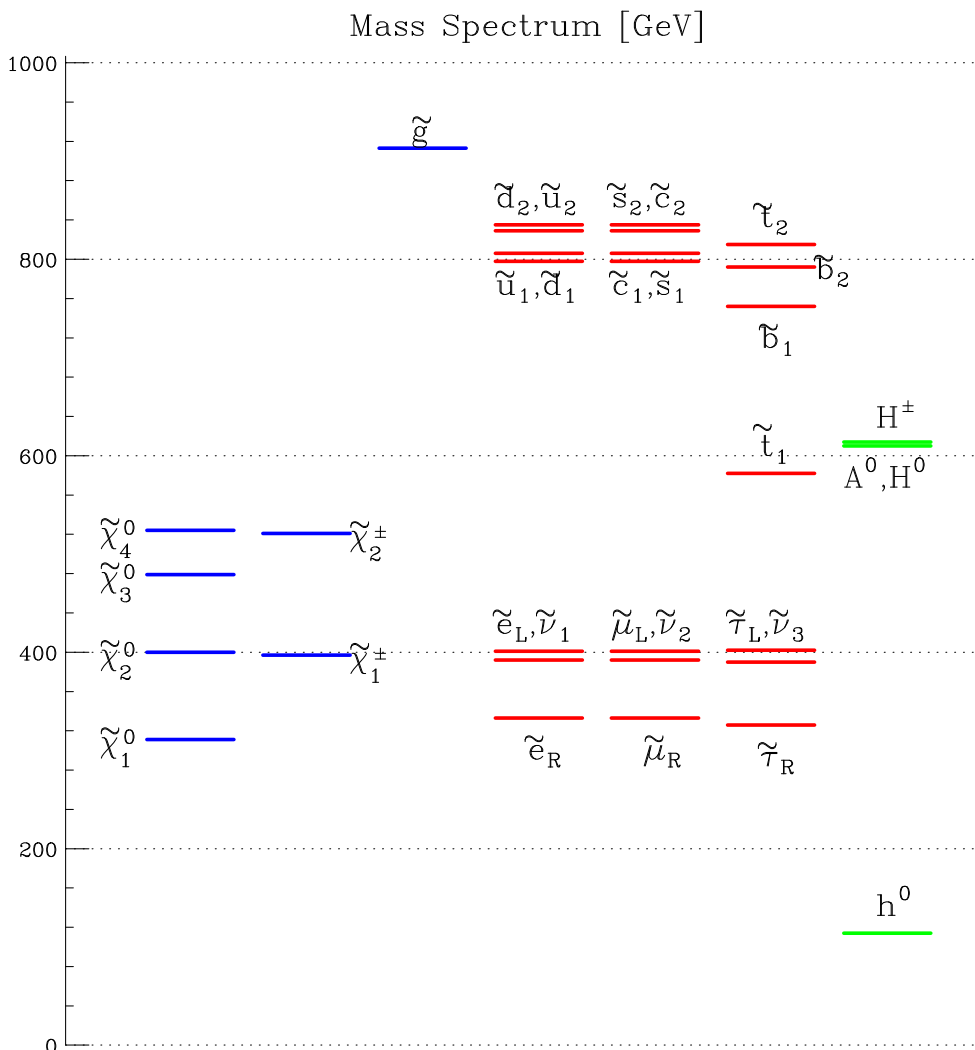


Figure 8: Pictorial representation of sample mass spectra in allowed $\mu > 0$ region. The numerical values are shown in column B of the sample spectra table.

AdS vacuum to a de Sitter one, but is typically in the range 4–6 [15]. The soft terms (5.1) can be rewritten as

$$m_\lambda^2 = M_s^2((1 - n_\lambda)\alpha^2 - \dot{\gamma}_\lambda + 2\alpha(T + T^*)\partial_T \gamma_\lambda), \quad (5.3)$$

with $\dot{\gamma}_\lambda = 8\pi^2 \partial \gamma_\lambda / \partial \log \mu$.

In the paper [16] the value $n_\lambda = 0$ corresponding to a field in the adjoint representation was used. However, in [8] it was shown that the correct modular weight for a chiral field is actually $n_\lambda = 2/3$. Using the explicit expressions for $\dot{\gamma}$ and $(T + T^*)\partial_T \gamma_i$ (see [16]), one

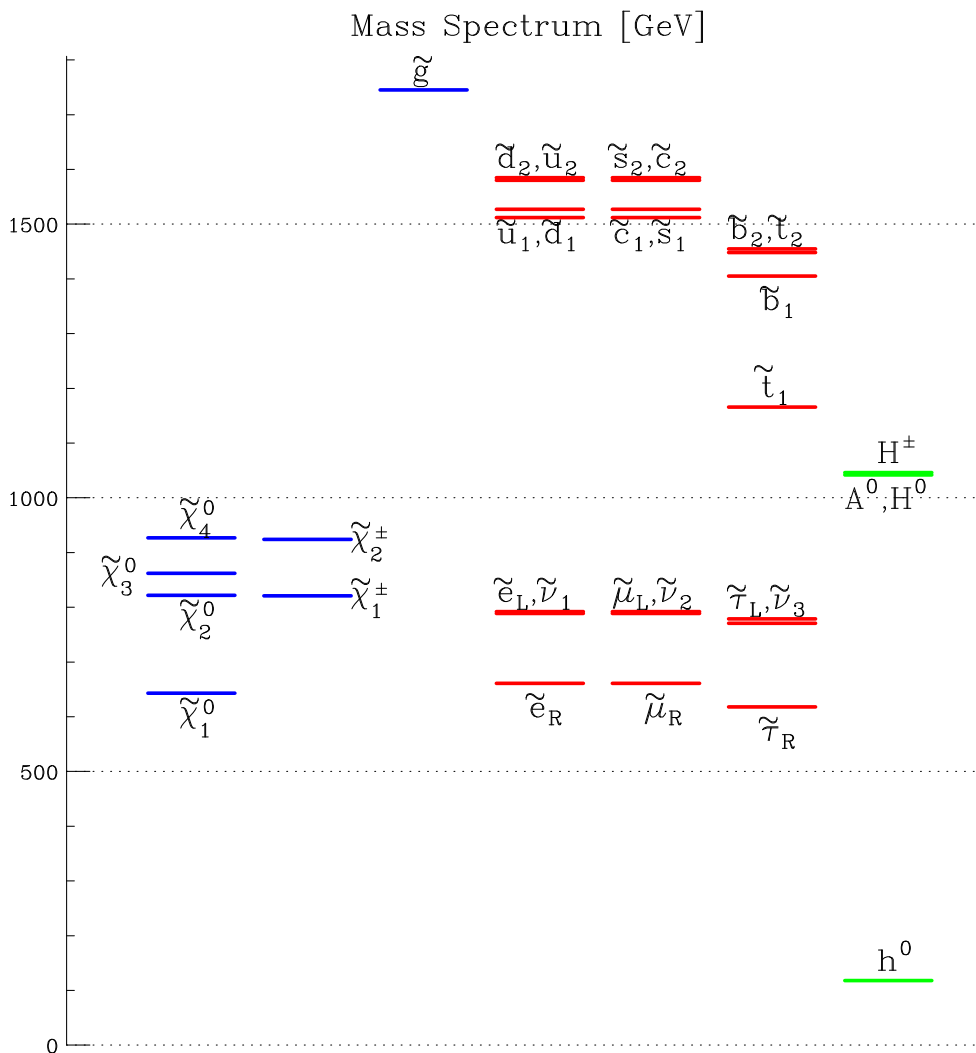


Figure 9: Pictorial representation for sample mass spectra in $\mu < 0$ and stau LSP region. The numerical values are shown in column C of the sample spectra table.

then finds boundary conditions at the GUT scale for the first two generations of particles:

$$\begin{aligned}
 m_L^2 &\sim (-1 - 2\alpha + \alpha^2/3)M_s^2, \\
 m_E^2 &\sim (-2 - \alpha + \alpha^2/3)M_s^2, \\
 m_Q^2 &\sim (2 - 4\alpha + \alpha^2/3)M_s^2, \\
 m_U^2 &\sim (1 - 3\alpha + \alpha^2/3)M_s^2, \\
 m_D^2 &\sim (2 - 3\alpha + \alpha^2/3)M_s^2.
 \end{aligned}
 \tag{5.4}$$

Requiring all the masses to be positive (in order to prevent tachyons and/or charge and colour breaking minima from appearing) gives a lower bound on α , $\alpha \geq 11.5$. Such values are difficult to obtain in the KKLT scenario.

A full reanalysis of the phenomenology of this scenario is beyond the scope of this article. Note that combining gravity and anomaly mediation also leads to some ambiguity of the soft terms. Even though gaugino masses have been fully determined in all generality [32], scalar masses have a Planck scale physics dependence that is not usually under control. Fortunately for the large volume case the cancellation was due to the standard cancellation of soft terms for no-scale supergravity that has been established in [27, 33] and this ambiguity is under control.¹⁰ Furthermore, since in the KKLT scenario the value of T is not very large, α' corrections should play an important role. It would be interesting to consider these corrections using the results of [50].

6. Conclusions

In this paper we have been able to go a long way towards developing a concrete model of low-energy supersymmetry breaking for chiral matter fields from D7 branes. Notice that D3 branes were considered in [12, 13] but there was a need to fine tune in order to obtain low-energy supersymmetry breaking. Previous discussions of soft terms for D7 branes used the matter metrics for non-chiral adjoint fields since the Kähler metric for chiral fields was unknown. Using the recent results of [8], we were able to find the soft supersymmetry breaking terms for chiral fields also, which are the ones most relevant for phenomenology.

We have found several interesting results.

1. There is a small hierarchy between the gravitino mass and all the soft breaking terms. This is given by the suppression factor $\epsilon^{-1} = \log(M_P/m_{3/2})$. This allows the gravitino to be more than one order of magnitude heavier than the scalars and gauginos. Without having to fine-tune W_0 , the large-volume scenario solves the hierarchy problem with an intermediate string scale $m_s \sim 10^{11}\text{GeV}$. The scenario also naturally generates an axion decay constant within the allowed window [51]. We are then led to a hierarchy of scales $m_s \sim 10^{11}\text{GeV} \gg m_{3/2} \sim 10\text{TeV} > m_{\text{soft}} \sim 500\text{GeV}$.
2. For general non-diagonal matter metrics we were able to extract a remarkably simple structure of soft terms, depending only on an overall scale $M_i \sim \epsilon m_{3/2}$ and a parameter λ which is the modular weight of the corresponding matter fields with respect to the small cycles,

$$\begin{aligned}
 M_i &= \frac{F^s}{2\tau_s}, \\
 m_\alpha &= \sqrt{\lambda} M_i, \\
 A_{\alpha\beta\gamma} &= -3\lambda M_i, \\
 B &= -(\lambda + 1) M_i.
 \end{aligned}
 \tag{6.1}$$

This produces a well motivated scenario of soft SUSY breaking at low-energies which can be subjected to a detailed phenomenological analysis.

¹⁰The ultraviolet effects can be controlled for instance by regularising with Pauli-Villars fields that appear in the no-scale combination (we thank E. Poppitz, M.K. Gaillard and B. Nelson for helpful discussions on this point).

3. The structure of the soft terms provides an elegant explanation of universality. The origin of universality can be traced to the fact that the sector that controls flavour (complex structure moduli) is decoupled, at leading order, from the sector that breaks supersymmetry (Kähler moduli). It is promising to have a solution to probably the most serious problem of gravity mediated supersymmetry breaking. Note that universality is only approximate in our scheme, as it is valid only to first order in an expansion of the Kähler moduli. This is interesting because it predicts the breaking of universality at next order in a gauge coupling expansion.
4. For the simplest case, corresponding to the Standard Model gauge and matter sectors coming from magnetised D7 branes wrapping the same cycle, $\lambda = 1/3$. The structure of breaking terms is essentially the same as the dilaton domination scenario and only differs in the value of the soft terms compared to the gravitino mass and the expression for the B term. This is quite curious since the origin of these two mechanisms is completely different. Dilaton domination was considered very special in the heterotic case due to its simplicity, universality and potential connection with finite theories [52, 53]. Furthermore, its simplicity allowed it to be subject to phenomenological constraints, such as the absence of charge and colour breaking, that ruled this scenario out for a string scale of M_{GUT} . Fortunately these constraints were satisfied if the string scale is intermediate [35] as indeed happens with our case. Note that the difference with the previous studies is that in the past the value of the string scale was put in by hand but now it is derived after a well defined process of moduli stabilisation. Furthermore in the past there were no explicit models in which dilaton domination was derived (see also the recent discussion in [54]).
5. Even though the soft breaking terms are suppressed with respect to the gravitino mass, the anomaly mediated contribution is further suppressed. The main reason for suppression in both cases is the underlying no-scale structure of our scenario. In pure no-scale models anomaly mediated soft terms vanish. Since no-scale is only approximate in the large volume models, the AMSB soft terms are non-vanishing but they are suppressed compared to the minimal AMSB soft breaking terms. Recall that for the no-scale model anomaly mediation also gives vanishing soft terms and therefore its contribution is doubly suppressed (by the no-scale approximate cancellation and the standard loop suppression). Therefore, contrary to the one-modulus KKL'T case, anomaly mediation does not play a role in our scenario.
6. The scenario provides interesting low-energy implications. Regarding CP violation, since the A terms are proportional to the gaugino masses, dangerous CP violating phases are not present (see for instance [30]). Second, there are regions in parameter space satisfying the standard constraints of charge and colour breaking minima, neutral LSP, $BR(b \rightarrow s\gamma)$, $(g - 2)_\mu$ and dark matter.
7. The simplest one-modulus KKL'T scenario also gets modified due to the knowledge of the chiral field matter metric. Even though this may not change the broad properties

of the model, the detailed phenomenology is changed due to the particular value of the modular weight, being $2/3$ and not 0 as has been assumed in the phenomenological discussions in [15–17].

There are several open questions that deserve further study. It would be interesting to further explore all the phenomenological implications of this scenario, analysing collider signatures in a similar fashion to the pioneering work of [55]. We have also neglected the contribution of fluxes to the gauge couplings and therefore also to the gaugino masses. This corresponds to magnetic fluxes being diluted in the corresponding 4-cycle. Our analysis in this paper has all been done in the dilute flux approximation. In addition to gaugino masses, the fluxes will also enter the Kähler metrics. For this reason it would be very interesting to extend this work by going beyond the limit of large cycle volumes and dilute fluxes. In general it should be remarked that the running to low-energies can only be seen as an indication of the implications of our soft terms, since some of its details depend on the precise structure of the realistic string model, including spectrum and couplings which are model dependent.

We have considered explicitly a minimal model in which the Kähler metric depends only on the volume and the corresponding ‘small’ modulus. A more general case would include dependence on other small moduli.

One issue with the above large-volume scenario is cosmological. As has been emphasised in [12, 14], our scenario is such that all but one of the Kähler moduli get masses of order $m_{3/2}/\epsilon$. These moduli decay very rapidly and cause no cosmological problem. However the modulus corresponding to the size of the large 4-cycle is lighter than the gravitino and extremely long lived. This may cause cosmological problems which need to be addressed [56, 57].

Finally, this scenario has a dichotomy about preferred values of the string scale. Either we have the GUT scale motivated by gauge unification and inflation, or we have the intermediate scale, motivated by the solution of the hierarchy problem in gravity-mediated scenarios, as well as being the natural scale for axions to solve the strong CP problem. We have taken for our analysis the second option, motivated by the fact that to have the GUT scale we need very small values of W_0 . We consider this to be a very strong fine tuning, which diminishes the significance of low-energy supersymmetry in addressing the hierarchy problem. Natural values of $W_0 \sim 1$ select the intermediate scale. If the use of very small values of W_0 can be justified, an extension of our analysis to include GUT scale string theory may be considered. This would be a replacement for the GUT scale fluxed MSSM models considered in [58, 59, 13], where the adjoint field metrics were used. However, it is well known that the entire $M - \tan\beta$ parameter space in the GUT scale dilaton domination scenario suffers from the presence of CCB minima [60], rendering the scenario rather unattractive.

It would also be interesting to search for scenarios in which the good properties of both the GUT and intermediate scale could be at work.

Acknowledgments

We acknowledge useful conversations with S.P. de Alwis, E. Bassouls, M. Berg, C.P. Burgess, M.K. Gaillard, D. Grellscheid, L.E. Ibáñez, S. Kachru, G. Kane, S. Kom, P. Kumar, B. Nelson, H.P. Nilles, E. Poppitz. We especially thank B. Allanach and D. Cremades for many interesting discussions. SA is supported by a Gates Fellowship. JC is supported by EPSRC and Trinity College, Cambridge. FQ is partially supported by PPARC and a Royal Society Wolfson award. KS is supported by Trinity College, Cambridge. SA thanks the organisers of YETI'06-BSM, IPPP Durham. FQ thanks KITP, Santa Barbara and the organisers of the ‘String Phenomenology’ program, where part of this work was developed and which was funded in part by the National Science Foundation under Grant No. PHY99-07949.

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