Published by Institute of Physics Publishing for SISSA

Received: *July 27, 2005*
Revised: *September 30, 2005*
Accepted: *November 4, 2005*
Published: *November 22, 2005*

SUSY phenomenology of KKLT flux compactifications

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Abstract: We study SUSY phenomenology of the KKLT (Kachru-Kallosh-Linde-Trivedi) type scenarios of string theory compactifications with fluxes. This setup leads to a specific pattern of soft masses and distinct phenomenological properties. In particular, it avoids the cosmological gravitino/moduli problems. Remarkably, the model allows for the correct abundance of SUSY dark matter consistently with all experimental constraints including the bound on the Higgs mass, $b \to s\gamma$, etc. This occurs for both small and large $\tan\beta$, and requires the SUSY spectrum above 1 TeV.

Keywords: Supergravity Models, Superstring Vacua, Supersymmetry Phenomenology, Supersymmetry Breaking.

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1. Introduction

String compactifications with fluxes have recently attracted considerable attention. The presence of fluxes allows to stabilize most moduli and eliminate these unwanted scalars from the low energy action [1]. One of the most attractive setups in which all the moduli are fixed and the cosmological constant is zero or small is a model due to Kachru, Kallosh, Linde and Trivedi (KKLT) [2]. The consequent SUSY spectrum exhibits a number of interesting features [3, 4]. In particular, the soft terms receive comparable contributions from gravity (modulus) mediated [5] and anomaly mediated [6] SUSY breaking.[1] Another robust feature is a hierarchy among the MSSM soft masses, the gravitino and moduli masses,

$$m_{\text{MSSM}} \ll m_{3/2} \ll m_{\text{moduli}} \,. \tag{1.1}$$

Some phenomenological aspects of this class of models have recently been studied in refs. [9, 10]. In particular, it was observed that the heavy gravitino and moduli alleviate cosmological problems associated with late decays of these particles [10]. Also, the pattern of soft masses was found to be quite distinct [9].

In the present work, we undertake a comprehensive study of phenomenological properties of the model. We analyze experimental constraints on the spectrum from collider bounds on sparticle and Higgs masses, BR($b \to s\gamma$), etc. as well those imposed by correct electroweak symmetry breaking and absence of charge and color breaking minima in the scalar potential. Then we study compatibility of these constraints with the requirement of the correct SUSY dark matter abundance. Although the spectrum is very constrained and parametrized in terms of three continuous quantities only ($m_{3/2}$, α and $\tan\beta$), we find that the right amount of dark matter can be produced in considerable regions of parameter space. Unlike in the common mSUGRA model, both low and high values of $\tan\beta$ are allowed.

The outline of the paper is as follows. In section 2 we introduce the KKLT model, in section 3 we analyze the consequent soft SUSY breaking terms and in section 4 we study relevant phenomenological constraints. Our conclusions are presented in section 5. Some technical details concerning the anomaly mediated soft terms are given in the appendix.

2. The KKLT setup

In this section we discuss the KKLT construction and its main features. The KKLT setup is based on Calabi-Yau compactifications of type-IIB string theory with fluxes [11]. The presence of background fluxes in the compactified space, that is non-zero vacuum expectation values of certain field strengths, allows one to fix all complex structure moduli as well as the dilaton [1]. The former parametrize the shape of the internal manifold and in the absence of fluxes have a zero potential to all orders in perturbation theory. Internal fluxes create a potential for moduli thereby mitigating a number of phenomenological problems associated with light or massless moduli. This mechanism, however, does not apply to the overall T-modulus parametrizing the size of the compact manifold. The KKLT proposal is to invoke nonperturbative mechanisms such as gaugino condensation on D7 branes to stabilize the remaining modulus. As a result, the vacuum energy in such a theory is negative which requires further modifications of the setup. To this end, KKLT add a contribution from a non-supersymmetric object (anti-brane) which does not significantly affect moduli stabilization. Thus the setup requires the presence of a number of D7/D3

[1] A similar pattern also appears in the heterotic string [7]. For phenomenology of compactifications with fluxes, see also [8].

branes and an anti D3 brane. The final outcome is that (i) all moduli are fixed, (ii) the cosmological constant is small and positive. This is the major achievement of the model.

Let us now consider the KKLT model in more detail. We start with a 4D supergravity scalar potential. A supergravity model is defined in terms of three functions: the Kähler potential K, the superpotential W, and the gauge kinetic function f. The scalar potential is given by

$$V_{\rm SUGRA} = M_{\rm Pl}^{-2}\, e^K \left(K^{I\bar{J}} D_I W D_{\bar{J}} W^* - 3|W|^2 \right). \tag{2.1}$$

Here $D_I W = \partial_I W + W \partial_I K$ is the Kähler covariant derivative of the superpotential and $K^{I\bar{J}} = (\partial_I \partial_{\bar{J}} K)^{-1}$. The gravitino mass is given by

$$m_{3/2} = M_{\rm Pl}^{-2}\, e^{K/2} W \tag{2.2}$$

and the SUSY breaking F-terms are

$$F^I = -M_{\rm Pl}^{-2}\, e^{K/2} K^{I\bar{J}} D_{\bar{J}} W^* \,. \tag{2.3}$$

Given K and W as functions of the fields in the system, one minimizes the scalar potential $V_{\rm SUGRA}$ and finds whether supersymmetry is broken ($F_I \neq 0$) in the vacuum or not. In supergravity, vanishing of the cosmological constant imposes the relation $m_{3/2}^2 \sim K^{I\bar{J}} F_I F_{\bar{J}}^*$, therefore $m_{3/2}$ serves as a measure of SUSY breaking. The gravitino acquires its mass through the super-Higgs effect, that is, it absorbs the spin 1/2 Goldstino associated with spontaneous SUSY breaking. The MSSM soft masses are controlled by the F-terms such that typically one expects the soft masses to be of the order of the gravitino mass. The moduli masses are found from derivatives of $V_{\rm SUGRA}$ at the minimum and are also within one-two orders of magnitude from $m_{3/2}$ (cf. [12]).

In the KKLT setup, the total scalar potential is given by the sum

$$V = V_{\rm SUGRA} + V_{\rm lift}\,, \tag{2.4}$$

where $V_{\rm lift}$ is an explicitly SUSY breaking contribution which serves to lift the minimum of the potential to a Minkowski or de Sitter vacuum. With a general $V_{\rm lift}$, the gravitino mass (2.2) is not related to the F-terms. and is an explicit mass term. Similarly, the moduli masses found by differentiating V are not related to $m_{3/2}$ or the F-terms. This has its advantages since the gravitino and the moduli can be made heavy so as to avoid cosmological problems associated with late decays of these particles. At the same time, the F-terms can be kept small enough to produce a TeV MSSM spectrum required by naturalness in the Higgs sector.

Let us now consider the specifics of the KKLT scenario. $V_{\rm SUGRA}$ is a function of the T-modulus (as well as the MSSM fields which we suppress) with the Kähler potential and the superpotential given by

$$K = -3\ln(T + \bar{T})\,, \qquad W = w_0 - Ce^{-aT}\,. \tag{2.5}$$

Here T is related to the compactification radius R, ${\rm Re}\, T \sim R^4$, w_0 is a constant induced by the fluxes, C is a model-dependent coefficient and a is related to the beta function of gaugino condensation on the D7 branes, $a = 8\pi^2/N_c$ for $\mathrm{SU}(N_c)$. The lifting potential due to the presence of the anti D3 brane is

$$V_{\rm lift} = \frac{D}{(T+\bar{T})^n}\,, \tag{2.6}$$

with n being an integer ($n = 2$ in the original KKLT version) and D is a tuning constant allowing to obtain a Minkowski/de Sitter vacuum.

At $D = 0$, the scalar potential is minimized at $V = -3m_{3/2}^2 M_{\rm Pl}^2$. The addition of the lifting term leaves the value of T at the minimum essentially intact. This is because the supergravity potential is exponentially steep unlike the lifting term. Thus, the effect of the lifting term is simply to change the vacuum energy to a small positive or zero value. This is achieved with $D \sim m_{3/2}^2 M_{\rm Pl}^2 \sim 10^{-26} M_{\rm Pl}^4$. Such a small value may appear unnatural. However, one should remember that the background geometry in the KKLT model is warped, $ds^2 = e^{2A(y)}\eta_{\mu\nu}dx^\mu dx^\nu + \cdots$ with y parametrizing the compact dimensions and $A(y)$ being a flux-dependent warp factor. At the location of the SM fields the warping can be negligible, $e^{2A} \sim 1$, whereas at the location of the anti D3 brane the space can be significantly warped, $e^{2A} \ll 1$. In this case, the natural mass scale on the anti D3 brane is much smaller than the Planck scale and can be chosen to be of the order of the intermediate scale, $m \sim e^A M_{\rm Pl} \sim \sqrt{m_{3/2} M_{\rm Pl}}$. Thus the desired value of D can be obtained by placing the anti-brane at the appropriate point in the compact space.

Minimizing the scalar potential one finds,

$$\begin{aligned}
m_{3/2} &\simeq M_{\rm Pl}^{-2}\frac{w_0}{(2\,\mathrm{Re}\,T)^{3/2}}\,,\\
a\,\mathrm{Re}\,T &\sim -\ln(m_{3/2}/M_{\rm Pl})\,,\\
\frac{F_T}{\mathrm{Re}\,T} &\sim \frac{m_{3/2}}{a\,\mathrm{Re}\,T}\,,\\
m_T &\sim a\,\mathrm{Re}\,T\, m_{3/2}\,.
\end{aligned} \tag{2.7}$$

To get a TeV MSSM spectrum, w_0 should be adjusted to be very small, 10^{-13}, which can be achieved by finetuning fluxes in the underlying string theory. Then, $a\,\mathrm{Re}\,T \sim 25$. This is a moderately large parameter leading to a hierarchy among the gravitino, the modulus and the MSSM soft masses. Indeed, assuming that the MSSM fields live on D7 branes, the soft masses are controlled by $F^T\partial_T \ln K_{\rm MSSM} \sim F^T/\,\mathrm{Re}\,T$ [13] and thus are suppressed by $a\,\mathrm{Re}\,T$ compared to the gravitino mass. On the other hand, the modulus mass is enhanced by the same factor compared to $m_{3/2}$. As mentioned earlier, this moderate hierarchy is highly desired from the cosmological perspective: the modulus and the gravitino produced in the early Universe would decay before the nucleosynthesis and thus would not affect the abundances of light elements.

We note here that the most important effect of the presence of an anti-brane is lifting the vacuum energy. As argued in ref. [4], other effects due to the existence of an explicit gravitino mass term on the anti D3 brane or other explicit SUSY breaking terms are expected to be suppressed by warping.

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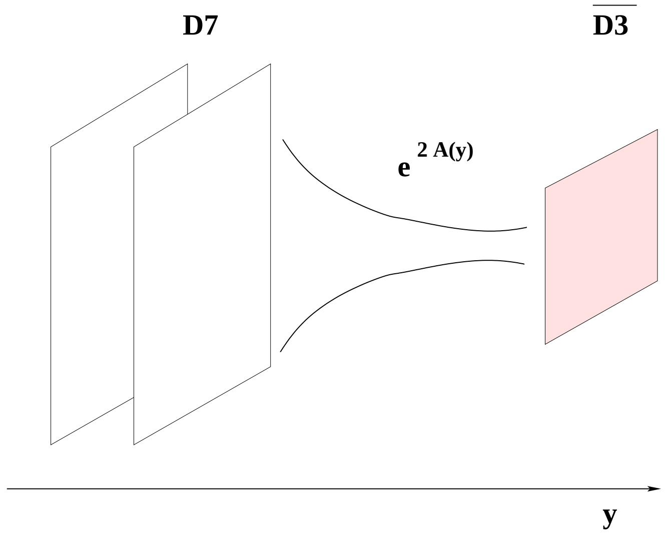

Figure 1: The KKLT setup with the MSSM fields on D7 branes. The factor $e^{2A(y)}$ represents warping along the compact dimension y (Klebanov-Strassler throat).

Concerning localization of the MSSM fields, there are a few options: they can live on D7 branes, D3 branes or on D7 and D3 branes.[2] There are certain advantages and disadvantages to each of these choices. If the observable fields are localized on the D3 branes, the MSSM spectrum is plagued by negative slepton masses squared — the usual problem of anomaly mediated SUSY breaking. Furthermore, it is difficult to get (semi-) realistic quark/lepton flavour structures. In the case of D7 branes, such problems do not arise: SUSY breaking is communicated by both the anomaly and the modulus F-term such that all masses can be made positive. For the flavour structures, in principle one can use the successful technology of intersecting branes [14]. On the other hand, the theoretical calculations are not well under control since

$$\mathrm{Re}\, T = \frac{1}{g_{\mathrm{GUT}}^2} \simeq 2 \tag{2.8}$$

requires a non-perturbative string coupling (cf. eqs.(2),(6) of ref. [4]). In any case, there are still some outstanding theoretical issues in this setup which have to do, for instance, with effects of explicit SUSY breaking contributions. We will not attempt to resolve these problems here. Instead, we will use the KKLT scenario as motivation to study certain patterns of soft SUSY breaking terms. As argued above, the setup with the MSSM on D7 branes is phenomenologically more appealing and we will take it as an assumption.

3. The MSSM soft terms

In this section we discuss the soft supersymmetry breaking terms for the MSSM fields living on D7 branes in the KKLT setup. Their main feature is that they interpolate

[2] In the original KKLT proposal, the MSSM fields were implicitly assumed to be localized on D3 branes. As we mention below, this choice is problematic due to the presence of tachyonic sleptons.

between the soft terms of anomaly-mediated SUSY breaking[3] and those of gravity-mediated SUSY breaking. This pattern appears generically whenever moduli are stabilized close to a supersymmetric point and leads to distinct phenomenology.

The Kähler potential and the kinetic function for the MSSM gauge fields are given by [13]

$$K = -3\ln(T+\overline{T}) + \sum_i \frac{|Q_i|^2}{(T+\bar{T})^{n_i}}\,, \qquad f_a = T\,. \tag{3.1}$$

Here Q_i are the MSSM matter fields and $a = 1, 2, 3$ runs over the GUT normalized U(1), SU(2) and SU(3) group factors. $n_i = \{0, 1/2, 1\}$ are constants depending on the origin and localization of the matter fields. For definiteness, we will set $n_i = 0$ in what follows. The choice $n_i = 1$ would lead to tachyons, whereas phenomenology of the $n_i = 1/2$ model would be quite similar to that of the $n_i = 0$ case (with larger tachyonic areas in parameter space).

The gaugino and soft scalar masses are generated by the auxiliary component F_T of the modulus superfield. Their magnitude is controlled by $F_T \sim m_{3/2}/a \ll m_{3/2}$. These tree level terms are much smaller than the gravitino mass and are comparable to the loop-suppressed anomaly-mediated contributions. The scale of the anomaly-mediated contributions is set by $F_\Phi/16\pi^2$, where $F_\Phi \equiv m_{3/2} + \frac{1}{3}F_T\partial_T K \simeq m_{3/2}$ and Φ is the conformal compensator. It is convenient to parametrize our F-terms in terms of a new scale M_s defined by

$$F_\Phi \simeq m_{3/2} = 16\pi^2 M_s\,, \qquad \frac{F_T}{T+\overline{T}} = \alpha\, M_s\,. \tag{3.2}$$

Here α depends on the shape of the lifting potential and is given by $\alpha \approx 16\pi^2 \frac{n}{2a\,\mathrm{Re}\,T}$ (note the difference from α defined by Choi *et al.* in ref. [9]!). Its precise value depends on further details of the model such as the string scale, the gravitino mass, etc.[4] For the original KKLT lifting potential $n = 2$ and α lies in the range $4.8 \div 6$. With other choices of the lifting potential, different values of α can be obtained, e.g. $\alpha = 7 \div 9$ for $n = 3$. In the limit $\alpha \to 0$ we recover pure anomaly mediation, while $\alpha \gg 5$ corresponds to gravity (modulus) mediation.

The soft terms in the mixed anomaly-modulus mediation scenario are controlled by the scale M_s and given by [4]:

$$\begin{aligned} M_a &= M_s\left[\alpha + b_a g_a^2\right]\,,\\ m_i^2 &= M_s^2\left[\alpha^2 - \dot{\gamma}_i + 2\alpha(T+\overline{T})\partial_T\gamma_i\right]\,,\\ A_{ijk} &= M_s\left[3\alpha - \gamma_i - \gamma_j - \gamma_k\right] + \Delta A_{ijk}\,. \end{aligned} \tag{3.3}$$

Here b_a are the beta function coefficients for the gauge couplings g_a, γ_i is the anomalous

[3] The anomaly-mediated contribution is usually present in string models [15], but may be absent in certain cases [16].

[4] This dependence appears since $a\,\mathrm{Re}\,T$ which solves the equation $e^{-a\,\mathrm{Re}\,T} = \frac{3}{2a\,\mathrm{Re}\,T}\frac{w_0}{C}$ depends on C and the gravitino mass.

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dimension and $\dot{\gamma}_i = 8\pi^2 \frac{\partial \gamma_i}{\partial \log \mu}$. In supersymmetric models,

$$\begin{aligned}\gamma_i &= 2\sum_a g_a^2 C_2^a(Q_i) - \sum_{y_i} |y_i|^2 \,, \\ \dot{\gamma}_i &= 2\sum_a g_a^4 b_a C_2^a(Q_i) - \sum_{y_i} |y_i|^2 b_{y_i} \,, \end{aligned} \tag{3.4}$$

where $C_2(Q_i)$ is the quadratic Casimir corresponding to the gauge representation of Q_i. In the second term, the sum runs over all physical Yukawa couplings y_i involving Q_i. The coefficient b_{y_i} describes the running of the Yukawa couplings, $\frac{\partial y_i}{\partial \log \mu} = \frac{1}{16\pi^2} y_i b_{y_i}$. Finally, the scalar soft masses contain a mixed anomaly-modulus contribution proportional to $\partial_T \gamma_i$ which appears due to the T-dependence of the gauge and physical Yukawa couplings,

$$(T + \overline{T})\partial_T \gamma_i = -2\sum_a g_a^2 C_2^a(Q_i) + 3\sum_i |y_i|^2 + \delta_i \,. \tag{3.5}$$

All relevant RG parameters are listed in appendix A. Numerically, the anomaly and gravity mediated pieces in eq. (3.3) are roughly the same at $\alpha \lesssim 3$.

δ_i in eq. (3.5) and ΔA_{ijk} in the expression for the A-terms account for a potential T-dependence of the Yukawa couplings,[5] $\Delta A_{ijk} \propto \partial_T \ln Y_{ijk}$. The presence of this term as well as its specific form depend on the theory of flavour and cannot be analyzed in full generality. For simplicity, we will omit this term in most of our phenomenological analyses, yet we will comment on some of the effects it can generate.

The remaining two parameters important for SUSY phenomenology are the μ and the $B\mu$ terms. Since these are responsible for electroweak symmetry breaking, their magnitude is bounded by the scale of the soft masses, that is around 1 TeV. This is rather difficult to achieve in models similar to the anomaly mediation scenario since the natural value for the B-term would be $F_\Phi \simeq m_{3/2} \gg M_s$. Nevertheless, the desired values can be obtained with some finetuning given the μ-term is generated in the superpotential, $\Delta W = \kappa H_1 H_2$, as well as the Kähler potential, $\Delta K = \kappa' H_1 H_2 + \text{h.c.}$ Then, parametrizing μ as

$$\mu = \mu_W + \mu_K \,, \tag{3.6}$$

one has

$$B\mu = c_1 m_{3/2}\mu_W + c_2 m_{3/2}\mu_K \,, \tag{3.7}$$

with $c_{1,2}$ being order one constants which depend on κ and κ'. Adjusting μ_W and μ_K appropriately, μ and B of order M_s can be obtained [9]. The practical conclusion is that, lacking a compelling model of generating μ and B, they should be treated as adjustable parameters so as to produce correct electroweak symmetry breaking.

Let us now overview main features of the resulting SUSY spectrum.

(i) Moduli/gravitino problem. A characteristic feature of the spectrum is a moderate hierarchy (a factor of 30 or so) between the MSSM soft masses and the gravitino mass as well as between the gravitino mass and the moduli masses. As discussed in

[5] In simple cases, the holomorphic Yukawa couplings are T-independent [17].

ref. [10], this is advantageous from the cosmological perspective since the gravitino and moduli are heavy enough to decay before the nucleosynthesis and not to affect abundances of light elements.

(ii) Tachyons. Pure anomaly mediation is notorious for its negative slepton mass squared problem. In the KKLT setup, there is an additional gravity mediated contribution which rectifies the problem. The absence of tachyons imposes a lower bound on the parameter α. Indeed, the GUT scale boundary condition for the slepton masses of the first two generations reads

$$\begin{aligned} m_L^2 &\approx \left(-1-2\alpha+\alpha^2\right) M_s^2\,, \\ m_E^2 &\approx \left(-2-\alpha+\alpha^2\right) M_s^2\,. \end{aligned} \tag{3.8}$$

To avoid tachyonic sleptons, $\alpha > 2$ is required. For the squarks,

$$\begin{aligned} m_Q^2 &\approx \left(2-4\alpha+\alpha^2\right) M_s^2\,, \\ m_U^2 &\approx \left(1-3\alpha+\alpha^2\right) M_s^2\,, \\ m_D^2 &\approx \left(2-3\alpha+\alpha^2\right) M_s^2\,. \end{aligned} \tag{3.9}$$

Although the squark masses are positive in pure anomaly mediation, they become tachyonic[6] for $0.5 < \alpha < 4$ due to the mixed anomaly-modulus contribution proportional to α. In conclusion, the tachyons which signify color or charge breaking minima are absent for $\alpha > 4$. This bound has important implications for phenomenology. In particular, most of the parameter space with characteristic signals of anomaly mediation such as a wino LSP is excluded. Curiously, $\alpha \sim 5$ predicted by the original KKLT model is on the safe side.

(iii) LSP. In the non-tachyonic region, the bino is the lightest gaugino. Our numerical analysis shows however that for $4 \lesssim \alpha \lesssim 8$ the LSP is dominated by the Higgsino component. This can be explained as follows. The anomaly and modulus mediated contributions add up in the GUT scale bino mass, $M_1 \approx M_s(\alpha + 33/10)$ but partially cancel in the gluino mass $M_3 \approx M_s(\alpha - 3/2)$. It is well known that the low energy value of the Higgs mass parameter $m_{H_2}^2$ and, consequently, the μ-term, is typically controlled by the GUT scale gluino mass, $\mu^2 \approx (2 \div 3)M_3^2$. Thus, for intermediate α where the suppression of the gluino mass is effective, we get $|\mu| < M_1(\,\mathrm{TeV}) \approx 0.4M_1$ and the LSP is higgsino-like. This is certainly desired from the SUSY dark matter perspective. We also note that the stau LSP is not possible in this scenario since unlike in mSUGRA the scalar and the gaugino masses cannot be varied independently.

(iv) Mirage unification. An interesting feature of the scenario is the occurence of mirage unification [9]. That is, even though the gaugino and the scalar masses do not

[6] For $\alpha > 2$, the squark masses squared are positive at the EW scale due to the RG running. However, $2 < \alpha < 4$ lead to tachyonic squarks at the GUT scale which signifies existence of color breaking minima in the effective potential.

unify at the GUT scale, RG running of these quantities makes them unify at some intermediate scale. Indeed, the solutions to the 1-loop RG equations (neglecting Yukawa contributions) read

$$\begin{aligned} M_a(\mu) &= M_s \frac{\alpha + b_a g_{GUT}}{1 - \frac{b_a g_{GUT}}{8\pi^2} \log \frac{\mu}{M_{GUT}}}\,, \\ m_i^2(\mu) &= M_s^2 \alpha^2 \left(1 + 2\frac{C_a(Q_i)}{b_a}\right) - 2\frac{C_a(Q_i)}{b_a} M_a^2(\mu)\,. \end{aligned} \tag{3.10}$$

At the mirage scale $\mu_{\rm mir}$,

$$\mu_{\rm mir} = M_{GUT} \mathrm{e}^{-8\pi^2/\alpha}\,, \tag{3.11}$$

all gaugino and scalar masses of the first two generations unify,

$$M_a^2(\mu_{\rm mir}) = m_i^2(\mu_{\rm mir}) = (\alpha M_s)^2\,. \tag{3.12}$$

This is truly a mirage scale as there is no physical threshold associated with it. Furthermore, the third generation scalar and the Higgs mass parameters do not unify at that scale. We note that for $\alpha \approx 5$ the mirage unification occurs at an intermediate scale, $\mu_{\rm mir} \sim 10^{11}$ GeV. In this case, the low energy spectroscopy is in some respects similar to that of gravity mediation with an intermediate string scale. In particular, the hierarchy between the squark and slepton masses is reduced, as compared to mSUGRA models.

(v) FCNC problem. The FCNC problem can only be addressed after realistic Yukawa flavour structures have been obtained. The problem appears when the n_i parameters of eq. (3.1) are generation-dependent. n_i are generally correlated with the Yukawa structures [18] such that the problem might actually be absent in realistic models. In our analysis, we simply assume that all $n_i = 0$ in which case the FCNC are suppressed.

In any case, as we argue in the next section, consistency with accelerator constraints requires a heavy SUSY spectrum, 1-5 TeV. Since the scalar mass matrix is diagonal, even generation-dependent choices for n_i would not lead to any significant FCNC problem [19].

(vi) CP problem. The equations of motion require $\mathrm{Arg}(F_\Phi) \simeq \mathrm{Arg}(F_T)$. As a result, the CP phase in the gaugino masses is aligned with the universal CP phase of the A-terms. This means that the physical phases $\mathrm{Arg}(M_a^* A)$ vanish. Yet, there remain two sources of dangerous physical phases. First, it is a phase of the type $\mathrm{Arg}(M_a^* B)= \mathrm{Arg}\Big(M_a^*(B\mu)\mu^*\Big)$. From eqs. (3.6), (3.7) it is clear that this phase is proportional to

$$\mathrm{Arg}(\mu_W^* \mu_K)\,.$$

It can be associated with the phase of the μ-term since M_a and $B\mu$ can be made real by $\mathrm{U}(1)_R$ and $\mathrm{U}(1)_{\rm PQ}$ rotations. This phase is strongly constrained by EDM experiments, $\phi_\mu < 10^{-2}$ (see e.g. [20]). Since there is no reason for μ_W and μ_K to

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be aligned and the presence of both is required by correct electroweak symmetry breaking, EDMs are overproduced unless the SUSY spectrum is heavy. We note that the same problem appears in the well known dilaton-domination scenario.

The second source of EDMs is associated with A-term non-universality [21], namely the term ΔA_{ijk} in eq. (3.3). Even if the A-terms could be made real by a $\mathrm{U}(1)_R$ rotation, they would be flavour-dependent. That is, they would not be aligned with the Yukawa matrices. The latter are necessarily complex and require diagonalization involving complex rotation matrices. Specifically, defining $\hat{A}_{ij} \equiv A_{ij} Y_{ij}$ with i, j being the flavour indices, the Yukawas and the A-terms transform under a basis change as

$$\begin{aligned} Y &\to V_L^\dagger \, Y \, V_R \,, \\ \hat{A} &\to V_L^\dagger \, \hat{A} \, V_R \,, \end{aligned} \tag{3.13}$$

where $V_{L,R}$ diagonalize the Yukawa matrices in the up– and down-sectors. In the basis where the Yukawa matrices are diagonal, the A-terms have a general form and their diagonal entries involve CP phases. The resulting EDMs usually exceed the experimental bounds by orders of magnitude [21]. In our analysis, we will simply assume that the dangerous term ΔA_{ijk} is absent.

We find however that although the CP phases are present generically, the induced EDMs are suppressed due to the heavy SUSY spectrum (1-5 TeV) such that no significant CP problem exists. In what follows, we set the CP phases to zero for simplicity.

To conclude this section, we find that the KKLT setup leads to an interesting pattern of the soft masses. Although it may not solve all the problems, it has a number of positive features, in particular with regard to cosmology. In the next section, we present our detailed numerical study of the spectrum and low energy observables.

4. Phenomenology

As discussed in the previous section, the model contains four free parameters at the GUT scale: the gravitino mass $m_{3/2} = 16\pi^2 M_s$, the modulus to anomaly mediation ratio parametrized by α, the μ-term and the B term. The absolute value of μ is determined by requiring correct electroweak symmetry breaking, whereas its sign remains free. Further, it is conventional to trade B for a low energy parameter $\tan\beta = \langle H_2^0 \rangle / \langle H_1^0 \rangle$, which is a function of B and other GUT scale parameters. Thus, the parameter space for phenomenological studies is

$$m_{3/2}\,, \qquad \alpha\,, \qquad \tan\beta\,, \qquad \mathrm{sgn}(\mu)\,. \tag{4.1}$$

These are our input parameters at the GUT scale, $\sim 2 \times 10^{16}$ GeV. We assume that effective field theory is valid below this scale and use RG equations to derive the low energy SUSY spectrum. This is really an assumption since the string coupling is large in the regime considered and the effective field theory approach may not be valid. To this end, we use the bottom-up perspective and study the pattern of the soft terms hinted by the KKLT model.

Once $\tan\beta$ and $\mathrm{sgn}(\mu)$ have been fixed, we scan over the gravitino mass $0 < m_{3/2} < 150\,\mathrm{TeV}$ and $0 < \alpha < 10$. The low energy mass spectrum is calculated using the Fortran package `SUSPECT` [22] and its routines described in detail in ref. [23]. Evaluation of the $b \to s\gamma$ branching ratio, the anomalous magnetic moment of the muon and the relic neutralino density is carried out using the routines of `micrOMEGAs1.3.1` [24, 25].

In what follows, we divide constraints on the model into two classes which we call "theoretical" and "accelerator". The theoretical constraints include correct electroweak symmetry breaking, absence of color and charge breaking minima, as well as the dark matter abundance consistent with the WMAP limits. The accelerator constraints include bounds on the Higgs and sparticle masses, the $b \to s\gamma$ branching ratio and similar observables.

4.1 Theoretical constraints

4.1.1 Electroweak symmetry breaking

Minimizing the MSSM Higgs potential leads to the standard relation

$$\mu^2 = \frac{-m_{H_2}^2 \tan^2\beta + m_{H_1}^2}{\tan^2\beta - 1} - \frac{1}{2} M_Z^2 \,, \tag{4.2}$$

imposed at the SUSY breaking scale defined by the average stop mass, $M_{\mathrm{SUSY}} = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$. In most cases, it is well approximated by

$$\mu^2 \approx -m_{H_2}^2 - \frac{1}{2} M_Z^2 \,. \tag{4.3}$$

When the right hand side is negative, electroweak breaking cannot occur. $m_{H_2}^2$ at M_{SUSY} is computed by using its RG evolution from the GUT scale, $\frac{\partial m_{H_2}^2}{\partial \log\mu} \approx 6 y_t^2 (m_{H_2}^2 + m_{U_3}^2 + m_{Q_3}^2 + A_t^2)$ with μ being the scale parameter. The result depends most sensitively on the gluino mass M_3 at the GUT scale which increases $m_{U_3}^2$ and $m_{Q_3}^2$. Typically, one finds $m_{H_2}^2(M_{\mathrm{SUSY}}) \approx -(2 \div 3) M_3^2$.

In the model under consideration, the anomaly and the gravity contributions appear in M_3 with opposite signs, $M_3 \approx M_s(\alpha - 3/2)$. For low α, the effect of M_3 on $m_{H_2}^2(M_{\mathrm{SUSY}})$ is suppressed such that other RG contributions become more important and a negative $m_{H_2}^2(M_{\mathrm{SUSY}})$ cannot be obtained. Thus, the requirement of correct electroweak symmetry breaking imposes a lower bound on α. Taken together with the constraint from the absence of tachyons, this bounds requires typically $\alpha > 4 \div 6$.

4.1.2 Colour and charge breaking minima

Generically, supersymmetric models have many flat directions in the field space. SUSY breaking terms usually lift these directions, but may also induce global or deep minima which break the electric charge and colour symmetries (CCB minima) [26]. It is important to verify that such minima do not develop.

Some of the dangerous CCB minima appear along D-flat directions when the trilinear A-terms are sufficiently large. Absence of such minima imposes constraints on the magnitude of the A-terms. In particular,

$$A_t^2 \lesssim 3(m_{H_2}^2 + m_{t_R}^2 + m_{t_L}^2)\,. \tag{4.4}$$

Eq. (3.3) implies that this constraint is usually respected. We have also checked this statement numerically.

Another type of constraints comes from $F-$ and $D-$ flat directions. Among the dangerous flat directions are those corresponding to the gauge invariants LH_2 and LLE, LQD. Absence of CCB minima along these directions usually guarantees their absence along the remaining directions (see e.g. [27]). A CCB minimum develops for a negative $m_{H_2}^2 + m_L^2$ due to the negative and large in magnitude $m_{H_2}^2$ at low energies. We find however that this does not occur in viable regions of the parameter space, mainly due to the negative anomaly mediated contribution to M_3 which reduces the magnitude of $m_{H_2}^2$. Altogether, absence of CCB minima does not constrain the model significantly.

4.1.3 Neutralino dark matter

The 2σ WMAP limit on the dark matter relic abundance is [28]

$$0.094 \lesssim \Omega_{DM} h^2 \lesssim 0.129\,. \tag{4.5}$$

In SUSY models, the typical dark matter candidate is the lightest neutralino and it is the case here. In most of the parameter space, the lightest neutralino χ_1^0 is the LSP. Assuming R-parity conservation it is stable. Then to get the consistent dark matter abundance one has to make sure that the neutralinos annihilate efficiently enough to satisfy the bound (4.5). In this computation we will assume that the LSP abundance is thermal. Further, we will treat regions of the parameter space violating the upper bound in (4.5) as “ruled out”, those within the bounds as “favoured” and those below the lower bound as “allowed”. The last case implies that there are additional ingredients to dark matter, beyond the MSSM, or that dark matter production is non-thermal.

The four neutralinos $\chi^0_{i=1,2,3,4}$ are superpositions of the neutral fermionic partners of the electroweak gauge bosons $\tilde{B}^0$ and $\tilde{W}_3^0$, and the superpartners of the neutral Higgs bosons $\tilde{H}_u^0$, $\tilde{H}_d^0$. In the basis ($\tilde{B}^0$, $\tilde{W}_3^0$, $\tilde{H}_u^0$, $\tilde{H}_d^0$), the neutralino mass matrix is given by

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -m_Z\cos\beta\sin\theta_W & m_Z\sin\beta\sin\theta_W \\ 0 & M_2 & m_Z\cos\beta\cos\theta_W & -m_Z\sin\beta\cos\theta_W \\ -m_Z\cos\beta\sin\theta_W & m_Z\cos\beta\cos\theta_W & 0 & -\mu \\ m_Z\sin\beta\sin\theta_W & -m_Z\sin\beta\cos\theta_W & -\mu & 0 \end{pmatrix}.$$

This is diagonalized by an orthogonal matrix Z such that the lightest neutralino is given by

$$\tilde{\chi}_1^0 = Z_{11}\tilde{B}^0 + Z_{12}\tilde{W}_3^0 + Z_{13}\tilde{H}_d^0 + Z_{14}\tilde{H}_u^0\,. \tag{4.6}$$

$\tilde{\chi}_1^0$ is usually called “gaugino-like” if $P \equiv |Z_{11}|^2 + |Z_{12}|^2 > 0.9$, “Higgsino-like” if $P < 0.1$, and “mixed” otherwise.

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It is instructive first to recall the situation with dark matter in the minimal supergravity model (mSUGRA). In most of the parameter space, the lightest neutralino is mainly the bino and, as a consequence, the annihilation cross section is small producing too large relic abundance. Nevertheless, there are three corridors in the parameter space where the cross section is enhanced. First, there is the stau-neutralino coannihilation branch, i.e. the region where the stau mass is almost degenerate with that of the LSP. Second, there is the A-pole region where $4(m_{\tilde{\chi}_1^0})^2 \sim m_A^2 \approx m_{H_1}^2 - m_{H_2}^2 - M_Z^2$ and the dominant neutralino annihilation process is due to the s-channel pseudo-scalar Higgs exchange. Finally, the annihilation cross section is enhanced if the LSP is of the Higgsino type, which occurs for small μ. In that case the neutralinos annihilate efficiently through the Z boson exchange and also coannihilate with the charginos.

We find that the first option cannot be realized in the model under consideration. The reason is that the stau is always much heavier than the LSP since, unlike in mSUGRA, the gaugino and the scalar masses cannot be varied independently. However, the A-pole and the Higgsino LSP corridors are indeed present and the WMAP bounds are respected in considerable regions of the parameter space.

4.2 Accelerator constraints

4.2.1 Direct search constraints

An important constraint on the parameters of the model comes from lower bounds on the sparticle and Higgs masses due to direct collider searches. We implement these bounds by first ensuring absence of tachyons in the squark and slepton sectors and then applying the LEP2 constraints. The most restrictive bounds are due the chargino mass constraint, $m_{\chi^+} > 103.5\,\mathrm{GeV}$, and, particularly, due to the lightest Higgs mass constraint. In the decoupling limit $M_A \gg M_Z$ which is applicable in all of the viable parameter space, the latter bound is $m_h > 114\,\mathrm{GeV}$ at 3σ. It is well known that this bound is sensitive to the value of the top mass. In most of our analysis, we have used the central value $m_t = 178\,\mathrm{GeV}$. We have subsequently studied sensitivity of the results to the precise value of the top mass by considering the 2σ limiting cases, $m_t = 174\,\mathrm{GeV}$ and $m_t = 182\,\mathrm{GeV}$.

4.2.2 BR($b \to s\gamma$)

The supersymmetric spectrum is constrained indirectly by the branching ratio of the $b \to s\gamma$ decay. The most important SUSY contributions involve the chargino-stop loops as well as the top-charged Higgs loops. We impose the 3σ bound from CLEO [29] and BELLE [30], $2.33 \times 10^{-4} \leq BR(b \to s\gamma) \leq 4.15 \times 10^{-4}$. We find that, typically, the $b \to s\gamma$ bound is more important for $\mu < 0$, but can also be relevant for $\mu > 0$, particularly at large tanβ.

4.2.3 Muon $g - 2$

The 2.7σ deviation of the experimental value of the muon anomalous magnetic moment [31] from the SM prediction [32] may be interpreted as indirect evidence for physics beyond the Standard Model and, in particular, supersymmetry. This deviation favours a relatively light SUSY spectrum and a specific set of SUSY parameters, e.g. a positive μ. We find

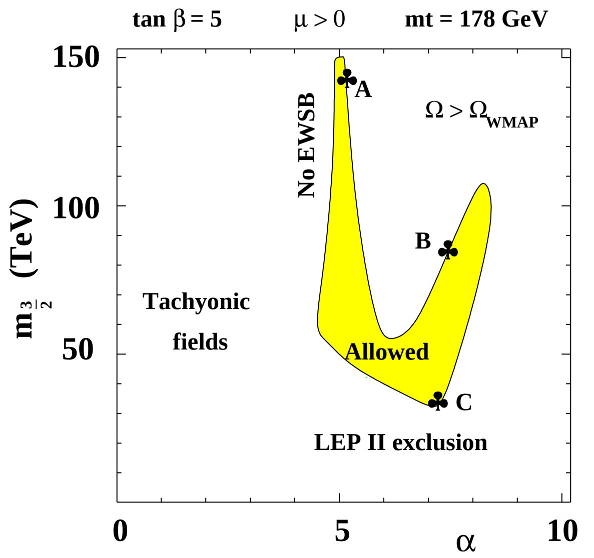

Figure 2: Example of the parameter region consistent with the theoretical and accelerator constraints.

however that consistency with other data requires a rather heavy spectrum in our model such that the muon $g-2$ deviation cannot be explained (unless $\tan\beta$ is large). Thus we will treat $g-2$ simply as a 3σ constraint on the model and will display the 2σ bands where relevant.

The discrepancy $\delta a_\mu \equiv a_\mu^{\rm exp} - a_\mu^{\rm SM}$ is measured to be $\delta a_\mu = (27.1 \pm 10) \times 10^{-10}$ if e^+e^- annihilation data for the calculation of $a_\mu^{\rm SM}$ are used. When the tau data are used instead, a smaller discrepancy is found. In this case, the 3σ bound is $\delta a_\mu > -6 \times 10^{-10}$, which we use in our analysis as a bound on the SUSY contribution $\delta a_\mu^{\rm SUSY}$. For $\mu < 0$, this excludes part of the parameter space with a relatively light spectrum. For $\mu > 0$, it imposes no constraint. In that case, we display in our figures the 2σ band $\delta a_\mu^{\rm SUSY} = 7.1 \times 10^{-10}$ for reference.

4.2.4 BR($B_s \to \mu^+\mu^-$)

For completeness, we include the bound on the $B_s \to \mu^+\mu^-$ branching ratio [33] $\mathrm{BR}(B_s \to \mu^+\mu^-) < 2.9 \times 10^{-7}$. It is known that it does not impose any significant constraints on the parameter space of mSUGRA. However, for non-universal soft terms which we are dealing with, the constraint may be significant [34], especially for large $\tan\beta$ and low Higgs masses. In practice, we find that the $\mathrm{BR}(B_s \to \mu^+\mu^-)$ constraint is satisfied automatically in regions of parameter space allowed by other considerations.

4.3 Example

Before going into a detailed discussion of our results let us present an example of the parameter space allowed by all the constraints. Figure 2 displays the surviving region in the plane $(\alpha,\ m_{3/2})$ for $\tan\beta = 5$ and $\mu > 0$. The area with $\alpha < 5$ or so is excluded by the presence of tachyons and absence of electroweak symmetry breaking. On the other

	A	**B**	**C**
$\tan\beta$	5	5	5
α	4.75	7	7.1
$m_{3/2}$ (TeV)	140	75	40
M_1	3308	2248	1179
M_2	3780	2877	1538
M_3	5616	5148	2903
$m_{\chi_1^0}$	974	2176	1163
$m_{\chi_2^0}$	976	2208	1329
$m_{\chi_1^+}$	975	2202	1321
$m_{\tilde{g}}$	5891	5391	3047
m_h	118	118	116
m_A	5115	4597	2573
m_H	5137	4616	2581
μ	955	2186	1327
$m_{\tilde{t}_1}$	4483	4000	2266
$m_{\tilde{t}_2}$	5477	4952	2798
$m_{\tilde{c}_1},\ m_{\tilde{u}_1}$	5792	5268	2972
$m_{\tilde{c}_2},\ m_{\tilde{u}_2}$	5951	5452	3075
$m_{\tilde{b}_1}$	5466	4946	2792
$m_{\tilde{b}_2}$	5902	5303	2988
$m_{\tilde{s}_1},\ m_{\tilde{d}_1}$	5761	5237	2955
$m_{\tilde{s}_2},\ m_{\tilde{d}_2}$	5951	5453	3076
$m_{\tilde{\tau}_1}$	4662	3644	1974
$m_{\tilde{\tau}_2}$	4669	3784	2061
$m_{\tilde{\mu}_1},\ m_{\tilde{e}_1}$	4332	3470	1881
$m_{\tilde{\mu}_2},\ m_{\tilde{e}_2}$	4507	3701	2017
$m_{\tilde{\nu}_3}$	4506	3700	2016
Ωh^2	0.099	0.105	0.094

Table 1: Sample spectra. All masses are in GeV, except for $m_{3/2}$ which is given in TeV.

hand, a large α region corresponding to the modulus dominated SUSY breaking is excluded by excessive dark matter abundance. The accelerator constraints yield a lower bound on the gravitino mass, $m_{3/2} \gtrsim 30\,\text{TeV}$. Very large values of $m_{3/2}$, except perhaps for a very thin strip, are excluded by a combination of the dark matter and electroweak symmetry breaking constraints. In table 1, we provide the SUSY spectrum for 3 representative points A,B,C in the surviving parameter space. These points are chosen such that the resulting dark matter abundance is consistent with the upper *and* lower WMAP bounds ("favoured" neutralino abundance).

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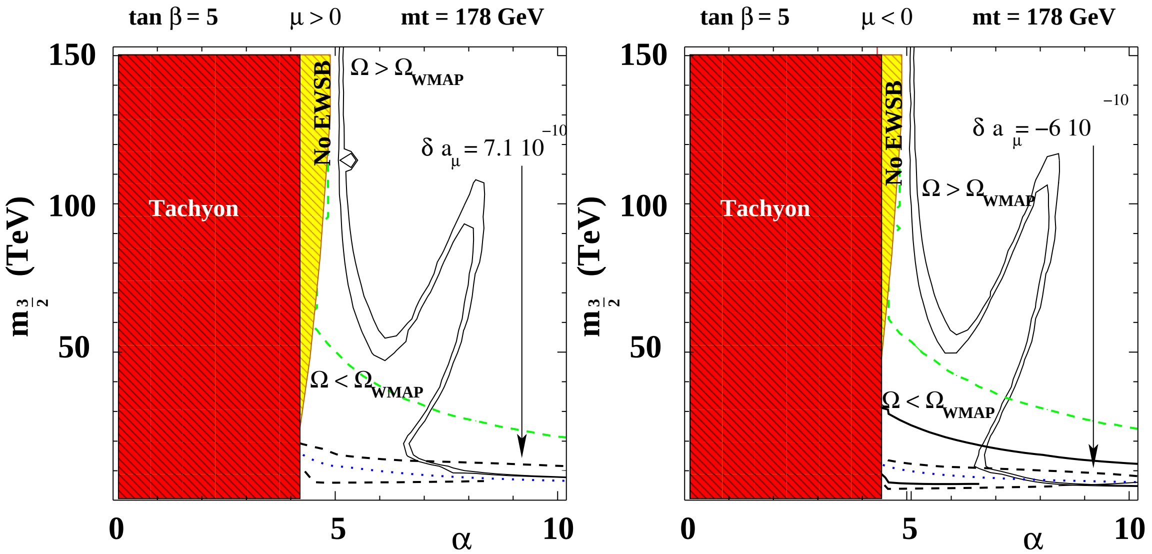

Figure 3: Constraints on the parameter space (α, $m_{3/2}$) of the KKLT scenario at $\tan\beta = 5$, $\mu > 0$ (left) and $\mu < 0$ (right). The region below the light grey (green) dashed line is excluded by the bound on the Higgs mass; below the the dotted line — by the bound on the chargino mass; below the solid line — by $b \to s\gamma$ (not shown on the left plot). The area between the black contours satisfies the WMAP constraint $0.094 \leq \Omega_{\tilde{\chi}_1^0}h^2 \leq 0.129$, whereas the region above it is excluded due to excessive LSP relic abundance. The black dashed line corresponds to a 2σ limit $\delta a_\mu^{\rm SUSY} > 7 \times 10^{-10}$ (left) and a 3σ limit $\delta a_\mu^{\rm SUSY} > -6 \times 10^{-10}$ (right). The lower black dashed line gives a 2σ upper bound $\delta a_\mu^{\rm SUSY} < 47 \times 10^{-10}$. BR($B \to \mu^+\mu^-$) imposes only a weak constraint and is not shown.

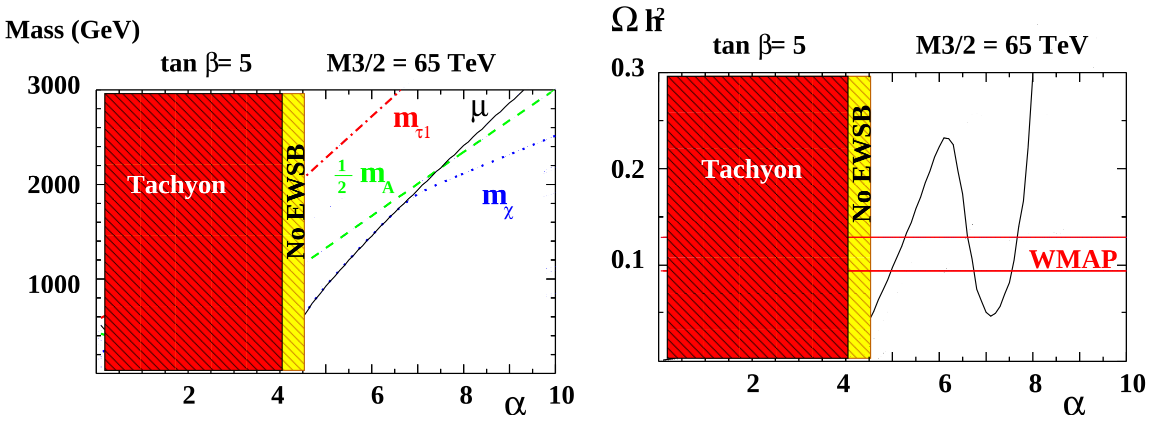

Figure 4: SUSY spectrum and the relic density as functions of α for $\tan\beta = 5$, $\mu > 0$ and $m_{3/2} = 65\,$TeV.

4.4 Numerical results

Our numerical results are summarized in figures 3– 8. These plots display contours corresponding to various constraints in the $(m_{3/2}, \alpha)$ plane for $\tan\beta = 5, 35$, a positive and

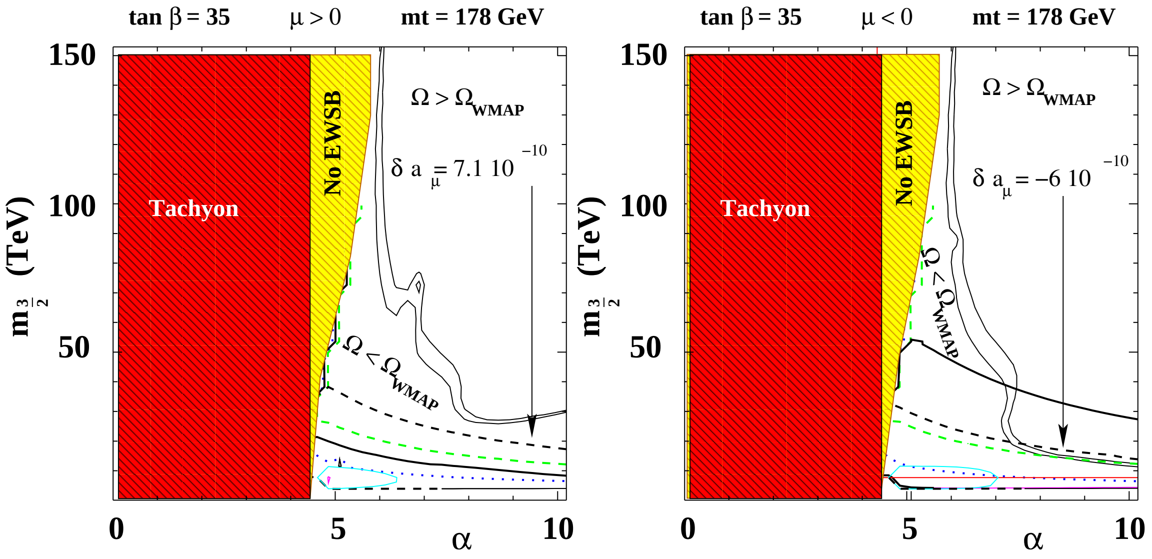

Figure 5: As in figure 3 but for $\tan\beta = 35$.

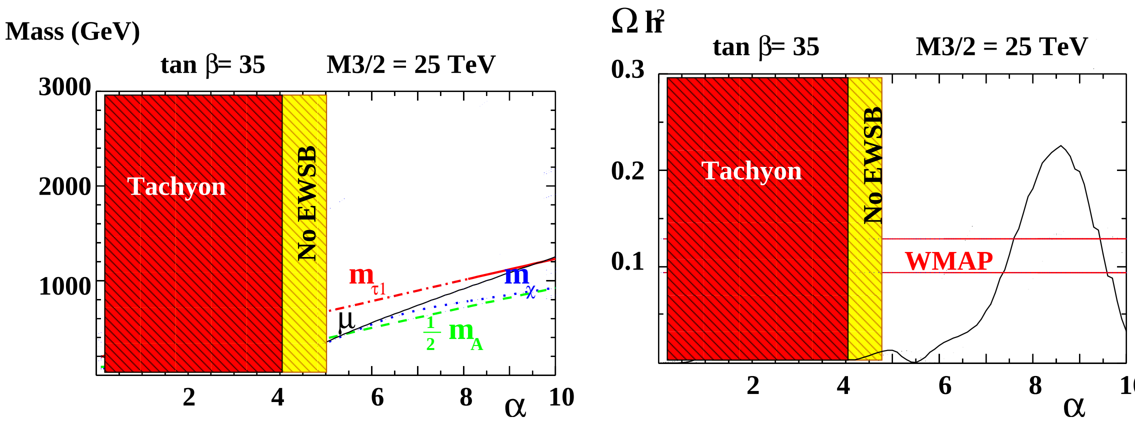

Figure 6: SUSY spectrum and the relic density as functions of α for $\tan\beta = 35$, $\mu > 0$ and $m_{3/2} = 25\,\text{TeV}$.

a negative μ-parameter and $m_t = 174, 178, 182\,\text{GeV}$. In addition, figures 4 and 6 show the evolution of the SUSY spectrum and the neutralino relic density with α when $m_{3/2}$ has been fixed.

There are several features of our analysis that are insensitive to $\tan\beta$ and $\operatorname{sgn}\mu$. In all of the considered cases, the parameter space with $\alpha \lesssim 4$ is excluded by requiring absence of tachyons. As explained in section 3, $\alpha \sim 0$ corresponds to pure anomaly mediation which predicts tachyonic sleptons. For $4 > \alpha > 2$, the sleptons have positive masses squared but the squarks turn tachyonic. This feature is specific to our scenario and appears due to the mixed modulus-anomaly contribution proportional to α in eq. (3.9).

Another robust feature is the presence of a "no electroweak symmetry breaking" (NO EWSB) region adjacent to the tachyonic area. As elaborated in subsection 4.1.1, it is related to the suppression of the gluino mass at low α. Electroweak symmetry breaking

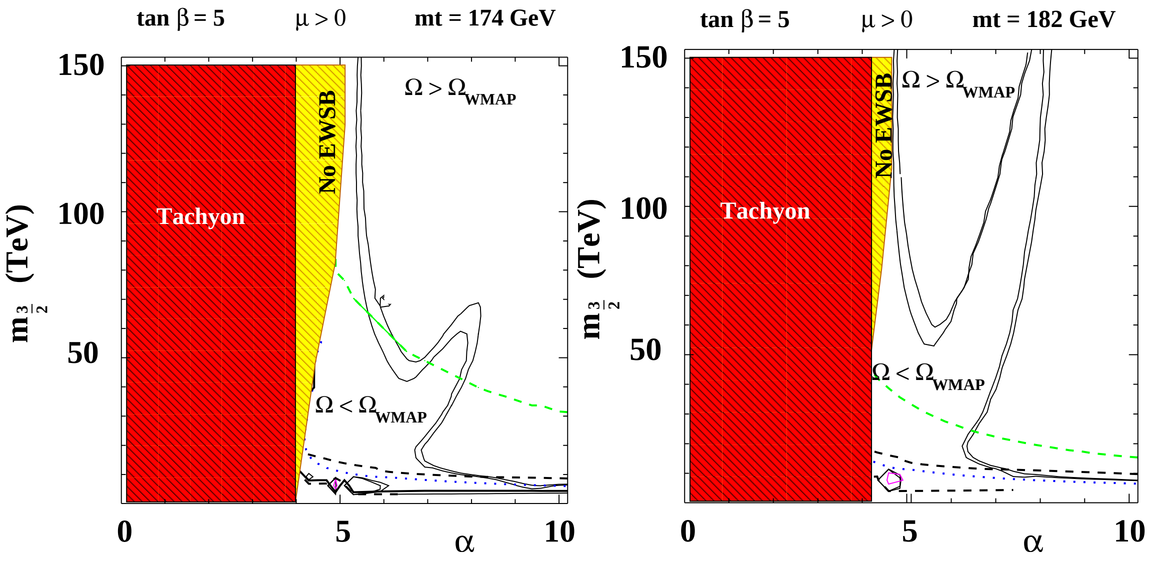

Figure 7: As in figure 3 but for tan$\beta = 5$, $\mu > 0$, $m_t = 174\,\mathrm{GeV}$ (left) and $182\,\mathrm{GeV}$ (right).

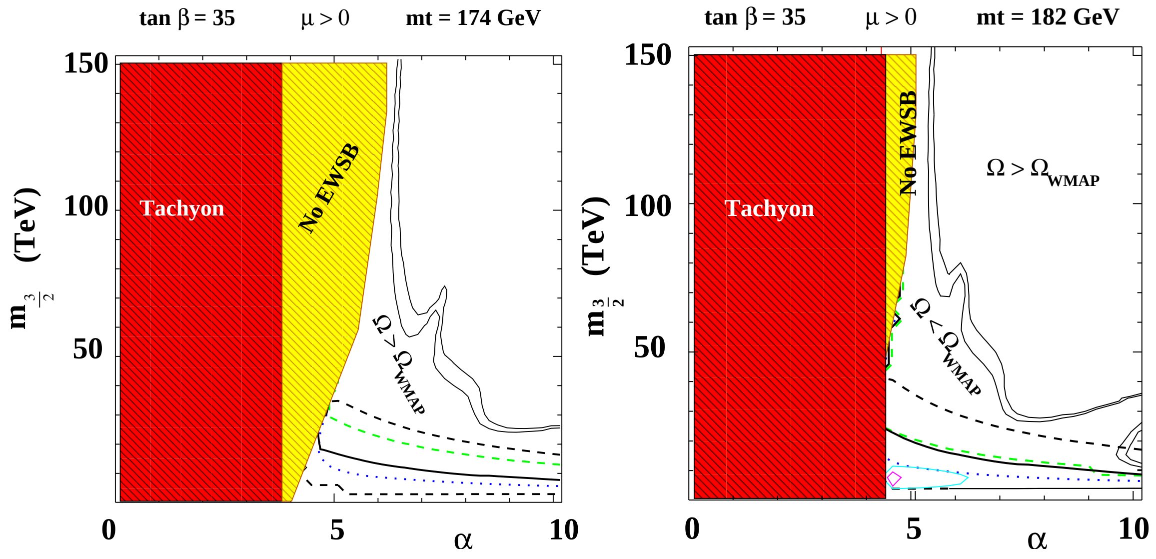

Figure 8: As in figure 3 but for tan$\beta = 35$, $\mu > 0$, $m_t = 174\,\mathrm{GeV}$ (left) and $182\,\mathrm{GeV}$ (right).

occurs when the Higgs mass parameter $m^2_{H_2}$ is negative and sufficiently large in magnitude (eq. (4.3)). The RG evolution of $m^2_{H_2}$ is controlled to a large extent by M_3 and for small gluino masses electroweak symmetry breaking is not possible. Appearance of the NO EWSB exclusion region at higher $m_{3/2}$ when α is fixed is associated with two loop effects.

In what follows, we study effects specific to certain regions of the parameter space.

4.4.1 Low tanβ regime

For small tan β, the most important accelerator bound is that on the lightest Higgs boson mass, see figure 3 (the green (light grey) dashed line). It sets a lower bound on the gravitino mass $m_{3/2} \gtrsim 30\,\text{TeV}$ which translates into a lower bound on the squark and slepton masses of order $2\,\text{TeV}$.

Clearly, such a heavy spectrum cannot explain the muon $g-2$ anomaly. For reference, we display in figure 3 the contour (black dashed) corresponding to $\delta a_\mu^{\text{SUSY}} = 7.1 \times 10^{-10}$. Above this line the SUSY contribution is too small to be relevant to the muon anomaly, yet it is allowed at a 3σ level.

The region between the two solid black contours satisfies the upper and lower WMAP bounds. The area above the contours corresponds to excessive neutralino abundance and is ruled out, whereas that below the contours is allowed given additional non-SUSY components of dark matter. Figure 4 explains the shape of the allowed region by tracking the composition of the LSP and the SUSY spectrum as a function of α at fixed $m_{3/2}$. For $\alpha \sim 5$ we are close to the NO EWSB region so that the μ term is small and the neutralino LSP is mainly a higgsino. Since $m_{\chi_1^0} \sim m_{\chi_1^+} \sim \mu$, the coannihilation with the chargino χ_1^+ is at work and, furthermore, the higgsino coupling to the Z allows for the efficient s-channel annihilation $\chi_1^0\chi_1^0 \to Z \to f\overline{f}$. This produces acceptable LSP relic abundance. As we increase α, the higgsino gets heavier and the LSP becomes more and more bino-like. The annihilation cross section decreases and the relic abundance becomes excessive already at $\alpha \sim 5.5$. As we go to even higher $\alpha \sim 7$, the neutralino mass approaches the value $M_A/2$, where the annihilation proceeds efficiently through the pseudo-scalar Higgs exchange. As a result, the relic density falls and the WMAP bounds are respected. At $\alpha > 8$ no efficient annihilation channel is available and dark matter is overproduced.

It is interesting to remark here that, in contrast with mSUGRA, the A-pole annihilation opens up for a *higgsino*-like neutralino, and not a *bino*-like one. Similar merging between the higgsino and the higgs annihilation branches has recently been observed in effective supergravity models with non-universal gaugino masses [36]. This effect allows for a large zone of the parameter space respecting the WMAP bounds at $\tan\beta = 5$ through efficient annihilation-coannihilation processes, which is not the case in mSUGRA.

4.4.2 Large $\tan\beta$ regime

At large $\tan\beta$, the Higgs mass constraint becomes less stringent. The main reason is that at tree-level $m_h \sim M_Z|\cos 2\beta|$, which increases with $\tan\beta$. It requires (figure 5) $m_{3/2} \gtrsim 20 - 30\,\text{TeV}$ as compared to $m_{3/2} \gtrsim 30 - 50\,\text{TeV}$ for low $\tan\beta$.

On the other hand, the $b \to s\gamma$ amplitude grows with $\tan\beta$ and becomes more important. In particular, it provides the most severe accelerator bound for $\mu < 0$ and excludes a large portion of the parameter space.

We observe that at large $\tan\beta$ even a heavy SUSY spectrum contributes significantly to the muon $g-2$ and ameliorates the anomaly for $\mu > 0$ (figure 5, region below the black dashed line).

$\text{BR}(B_s \to \mu^+\mu^-)$ does not provide any considerable constraint. Indeed, this observable is relevant for highly non-universal cases [34] which in our model are excluded by the presence of tachyons (low α regime).

The evolution of the relic density with α (figure 6) differs from the low $\tan\beta$ case because the pseudo-scalar Higgs exchange followed by a Higgs decay into $b\overline{b}$ pairs now

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dominates the relic abundance calculation. This process is efficient at large $\tan\beta$ for two reasons. First, the pseudo-scalar Higgs mass $m_A^2 \approx m_{H_1}^2 - m_{H_2}^2$ is much smaller due to the negative bottom quark Yukawa RG contribution to $m_{H_1}^2$. For example, $m_A \sim 2\mu \sim 2m_{\chi_1^0}$ already at $\alpha \sim 5$ (figure 6, left). Second, the $Ab\bar{b}$ coupling is proportional to $\tan\beta$ and the corresponding cross section $\sigma_{\chi\chi\to A\to b\bar{b}}$ grows as $\tan^2\beta$. As a result, the relic abundance is well below the WMAP range for $4 < \alpha < 7$. As the bino component of the neutralino increases, the relic density grows to its maximum value around $\alpha \sim 9$. Then it drops again for $\alpha \sim 10$ where $m_A \sim 2M_1 \sim 2m_{\chi_1^0}$ corresponding to an opening of a bino-like A-pole.

4.4.3 Influence of the sign of μ

It is well known that the $b \to s\gamma$ constraint is more important for $\mu < 0$ [37]. The reason is that in this case the SUSY contributions interfere constructively with those of the SM increasing the branching ratio, especially at large $\tan\beta$. This effect is clearly seen from figures 3 and 5 (right).

The SUSY contribution to the muon $g-2$ usually has the same sign as μ. Thus, a positive μ is preferred by the muon anomaly. In any case, the $g-2$ discrepancy is not conclusive and we treat it as a 3σ constraint.

Other observables are less sensitive to the sign of μ.

4.4.4 Uncertainties due to the top mass

We find that some of the results are very sensitive to the precise value of the top mass. To take this into account, we provide the exclusion plots for 3 values of the top mass: the central value $m_t = 178\,\mathrm{GeV}$ and the 2σ limits $m_t = 174, 182\,\mathrm{GeV}$[7] (figures 7, 8).

The top mass affects most of all the Higgs mass bound and the relic density. The former is sensitive to m_t through the one loop correction $\delta m_h^2 \propto \frac{m_t^4}{m_W^2}\log\left(\frac{m_{\tilde{t}}^2}{m_t^2}\right)$. For a heavier top, a larger portion of the parameter space is allowed by the Higgs mass constraint (figures 7 and 8, right).

The neutralino relic density is affected by m_t mainly in the A-pole region. There the neutralino is typically higgsino-like, $m_{\chi_1^0} \sim \mu$. The value of $\mu^2 \sim -m_{H_2}^2$ depends strongly on m_t via the top Yukawa contributions to $m_{H_2}^2$. For the same reason, the pseudo-scalar Higgs mass is sensitive to m_t, $m_A^2 \sim m_{H_1}^2 - m_{H_2}^2$. The net result is that for larger m_t, a broader A-pole region is available (figure 7). This effect disappears at large $\tan\beta$ in which case the bottom quark Yukawa decreases $m_{H_1}^2$ and, consequently, m_A^2.

We note that at large $\tan\beta$ and m_t an "mSUGRA-like" A-pole regime becomes available. This is seen in figure 8 (right) at $\alpha \sim 10$ and $m_{3/2} \sim 25\,\mathrm{TeV}$. The pole corresponds to annihilation of bino-like neutralinos through the pseudo-scalar Higgs.

4.4.5 Summary

The above analysis shows that there are considerable regions of parameter space where the model is consistent with all the constraints. The most restrictive accelerator bounds are

[7] The preliminary CDF/D0 update [38] yields $m_t = 172.7 \pm 2.9\,\mathrm{GeV}$. The central value $m_t = 178\,\mathrm{GeV}$ used in our analysis is within the 2σ interval. For smaller m_t the Higgs mass bound is more constraining, see figures 7, 8.

due to the Higgs mass constraint and BR($b \to s\gamma$). In some parts of parameter space, the muon $g-2$ anomaly is ameliorated. A positive μ-parameter is preferred, whereas both low and high values of $\tan\beta$ are allowed. The results are sensitive to the top mass such that its higher values lead to larger allowed regions of parameter space.

In most of the cases considered, the resulting SUSY spectrum is rather heavy. This can be understood as follows. The Higgs mass constraint yields a lower bound on the stop masses of order 1 TeV. Since all the SUSY masses are controlled by $m_{3/2}$, this bound implies a large $m_{3/2}$ and thus a heavy spectrum. This is different from the mSUGRA case where the scalar masses, the gaugino masses and the A-term can be varied independently. In spite of the heavy spectrum, the degree of fine-tuning to get the right EW breaking scale is similar to that of mSUGRA ($< 1\%$) as it is mainly sensitive to the 3rd generation scalar masses.

We note that if we do not insist on the neutralino being the dominant component of dark matter, for large $\tan\beta$ the spectrum is allowed to be lighter, 300 GeV — 1 TeV.

5. Conclusions

In this work, we have studied SUSY phenomenology of the KKLT-type flux compactification scenario with the MSSM on D7 branes. This setup leads to a specific pattern of the soft masses, with modulus and anomaly mediated contributions being comparable, and avoids the cosmological gravitino/moduli problems.

The parameter space includes 3 continuous variables $m_{3/2}, \alpha, \tan\beta$ and a discrete parameter sgnμ. The resulting SUSY spectrum is non-universal which distinguishes the model from mSUGRA and leads to distinct phenomenology. In particular, the neutralino LSP is often higgsino-like such that low $\tan\beta$ is allowed by dark matter considerations, in addition to the usual large $\tan\beta$ regime. We find that all experimental constraints can be satisfied simultaneously in large portions of parameter space. Curiously, $\alpha \sim 5$ required by the shape of the original KKLT lifting potential is consistent with the constraints.

We find that the SUSY spectrum is required to be quite heavy, typically between 1 and 5 TeV. Although this has certain merits in relation to the CP and flavour problems, it may be challenging to discover the superpartners at the LHC. Yet, at least part of the parameter space with the squarks and gluinos below roughly 3 TeV will be explored. We note also that in some cases the charginos may be long lived due to their near degeneracy with the LSP (table 1), which represents a typical anomaly mediation signature [39].

Finally, it is encouraging that a theoretical model conceived to address the moduli stabilization problem turned out to have remarkably healthy phenomenological properties.

Acknowledgments

We are grateful to Tilman Plehn for discussions and to Wilfried Buchmüller, Kiwoon Choi and Koichi Hamaguchi for comments on the manuscript. We would also like to thank Tania Robens for inspiring discussions, support and endurance.

A.F. was partially supported by the Polish KBN grant 2 P03B 129 24 for years 2003-2005 and by the EC Contract MRTN-CT-2004-503369 - network "The Quest for Unification: Theory Confronts Experiment" (2004-2008). The stay of A.F. at DESY is possible owing to the Research Fellowship granted by Alexander von Humboldt Foundation.

A. MSSM RG parameters

In this appendix, we list the MSSM renormalization group parameters which appear in the soft terms formulae (3.3). The U(1), SU(2) and SU(3) gauge couplings are denoted by g_a. Here U(1) is GUT normalized and related as $g_1 = \sqrt{5/3}g_Y$ to the hypercharge coupling.

The beta function coefficients b_a are defined as

$$\frac{\partial g_a}{\partial \log \mu} = \frac{1}{16\pi^2} b_a g_a^3 \,. \tag{A.1}$$

In the MSSM,

$$b_3 = -3 \,, \qquad b_2 = 1 \,, \qquad b_1 = \frac{33}{5} \,. \tag{A.2}$$

The anomalous dimension γ_i describes the RG dependence of the wave function renormalization Z_i,

$$\frac{\partial \log Z_i}{\partial \log \mu} = \frac{1}{8\pi^2} \gamma_i \,. \tag{A.3}$$

In supersymmetric theories, the following general formula holds:

$$\gamma_i = 2 \sum_a g_a^2 C_2^a(Q_i) - \sum_{y_i} |y_i|^2 \,. \tag{A.4}$$

In the second term, the sum runs over all Yukawa couplings y_i involving Q_i with appropriate color factors included. The quadratic Casimir $C_2(Q_i)$ takes the following values: $C_2^3 = 4/3$ for the SU(3) fundamental or anti-fundamental representation, $C_2^2 = 3/4$ for the SU(2) fundamentals, $C_2^1 = q_i^2$, where q_i is the U(1) charge of Q_i. The anomalous dimensions of the MSSM fields read:

$$\begin{aligned}
\gamma_{Q_p} &= \frac{8}{3} g_3^2 + \frac{3}{2} g_2^2 + \frac{1}{30} g_1^2 - (y_t^2 + y_b^2)\delta_{3p} \,, \\
\gamma_{U_p} &= \frac{8}{3} g_3^2 + \frac{8}{15} g_1^2 - 2y_t^2 \delta_{3p} \,, \\
\gamma_{D_p} &= \frac{8}{3} g_3^2 + \frac{2}{15} g_1^2 - 2y_b^2 \delta_{3p} \,, \\
\gamma_{L_p} &= \frac{3}{2} g_2^2 + \frac{3}{10} g_1^2 - y_\tau^2 \delta_{3p} \,, \\
\gamma_{E_p} &= \frac{6}{5} g_1^2 - 2y_\tau^2 \delta_{3p} \,, \\
\gamma_{H_1} &= \frac{3}{2} g_2^2 + \frac{3}{10} g_1^2 - 3y_b^2 - y_\tau^2 \,, \\
\gamma_{H_2} &= \frac{3}{2} g_2^2 + \frac{3}{10} g_1^2 - 3y_t^2 \,.
\end{aligned} \tag{A.5}$$

We have neglected all the Yukawa couplings except for the diagonal ones involving the third generation.

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The soft term formulae also involve $\dot{\gamma}_i = 8\pi^2 \frac{\partial \gamma_i}{\partial \log \mu}$. From eq. (A.4),

$$\dot{\gamma}_i = 2 \sum_a g_a^4 b_a C_2^a(Q_i) - \sum_{y_i} |y_i|^2 b_{y_i} \,. \tag{A.6}$$

Here b_{y_i} describes the running of the Yukawa couplings, $\frac{\partial y_i}{\partial \log \mu} = \frac{1}{16\pi^2} y_i b_{y_i}$. In the MSSM,

$$\begin{aligned}
\dot{\gamma}_{Q_p} &= -8g_3^4 + \frac{3}{2}g_2^4 + \frac{11}{50}g_1^4 - (y_t^2 b_{y_t} + y_b^2 b_{y_b})\delta_{3p} \,, \\
\dot{\gamma}_{U_p} &= -8g_3^4 + \frac{88}{25}g_1^4 - 2y_t^2 b_{y_t}\delta_{3p} \,, \\
\dot{\gamma}_{D_p} &= -8g_3^4 + \frac{22}{25}g_1^4 - 2y_b^2 b_{y_b}\delta_{3p} \,, \\
\dot{\gamma}_{L_p} &= \frac{3}{2}g_2^4 + \frac{99}{50}g_1^4 - y_\tau^2 b_{y_\tau}\delta_{3p} \,, \\
\dot{\gamma}_{E_p} &= \frac{198}{25}g_1^4 - 2y_\tau^2 b_{y_\tau}\delta_{3p} \,, \\
\dot{\gamma}_{H_1} &= \frac{3}{2}g_2^4 + \frac{99}{50}g_1^4 - 3y_b^2 b_{y_b} - y_\tau^2 b_{y_\tau} \,, \\
\dot{\gamma}_{H_2} &= \frac{3}{2}g_2^4 + \frac{99}{50}g_1^4 - 3y_t^2 b_{y_t} \,,
\end{aligned} \tag{A.7}$$

where

$$\begin{aligned}
b_{y_t} &= 6y_t^2 + y_b^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{13}{15}g_1^2 \,, \\
b_{y_b} &= y_t^2 + 6y_b^2 + y_\tau^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{7}{15}g_1^2 \,, \\
b_{y_\tau} &= 3y_b^2 + 4y_\tau^2 - 3g_2^2 - \frac{9}{5}g_1^2 \,.
\end{aligned} \tag{A.8}$$

Finally, the soft scalar masses contain a mixed anomaly-modulus contribution proportional to $\partial_T \gamma_i$ which appears due to the T-dependence of the gauge couplings. In our model,

$$(T + \overline{T})\partial_T \gamma_i = -2 \sum_a g_a^2 C_2^a(Q_i) + 3 \sum_i |y_i|^2 \,, \tag{A.9}$$

such that

$$\begin{aligned}
(T + \overline{T})\partial_T \gamma_{Q_p} &= -\frac{8}{3}g_3^2 - \frac{3}{2}g_2^2 - \frac{1}{30}g_1^2 + 3(y_t^2 + y_b^2)\delta_{3p} \,, \\
(T + \overline{T})\partial_T \gamma_{U_p} &= -\frac{8}{3}g_3^2 - \frac{8}{15}g_1^2 + 6y_t^2\delta_{3p} \,, \\
(T + \overline{T})\partial_T \gamma_{D_p} &= -\frac{8}{3}g_3^2 - \frac{2}{15}g_1^2 + 6y_b^2\delta_{3p} \,, \\
(T + \overline{T})\partial_T \gamma_{L_p} &= -\frac{3}{2}g_2^2 - \frac{3}{10}g_1^2 + 3y_\tau^2\delta_{3p} \,, \\
(T + \overline{T})\partial_T \gamma_{E_p} &= -\frac{6}{5}g_1^2 + 6y_\tau^2\delta_{3p} \,, \\
(T + \overline{T})\partial_T \gamma_{H_1} &= -\frac{3}{2}g_2^2 - \frac{3}{10}g_1^2 + 9y_b^2 + 3y_\tau^2 \,, \\
(T + \overline{T})\partial_T \gamma_{H_2} &= -\frac{3}{2}g_2^2 - \frac{3}{10}g_1^2 + 9y_t^2 \,.
\end{aligned} \tag{A.10}$$

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RECEIVED: *September 21, 2006*

Erratum

Confusion in various soft term conventions has lead to an unfortunate error in our numerical analysis. Derivation of the soft terms in eq. (3.3) assumes that the Yukawa couplings between gauginos and matter fields contain the i factor: $\mathcal{L} = -i\sqrt{2}\partial_i\partial_{\bar{j}}K\chi^a Q_j^\dagger T^a\psi_i + \text{h.c.}$ The numerical codes use a different convention (without the i factor) which, effectively, amounts to changing the relative sign between the gaugino masses and the A–terms.

Correcting this error leads to the following modifications:

1. The A-terms are typically large at the TeV scale.
2. Some of the parameter space is excluded by the presence of the stop and stau LSP.
3. The neutralino LSP is usually a bino.
4. The SUSY spectrum can be lighter (below 1 TeV) in certain cases.
5. The shape of the allowed parameter space somewhat changes (figure 9).

In spite of these changes many qualitative results remain the same. In particular, there is a considerable part of the parameter space, e.g. around $\alpha \sim 5$, which is allowed by all the constraints. Typically, the spectrum is heavy (above 1 TeV), see table 2, although some exceptions can be found. This is enforced by the Higgs mass bound and BR($b \to s\gamma$).

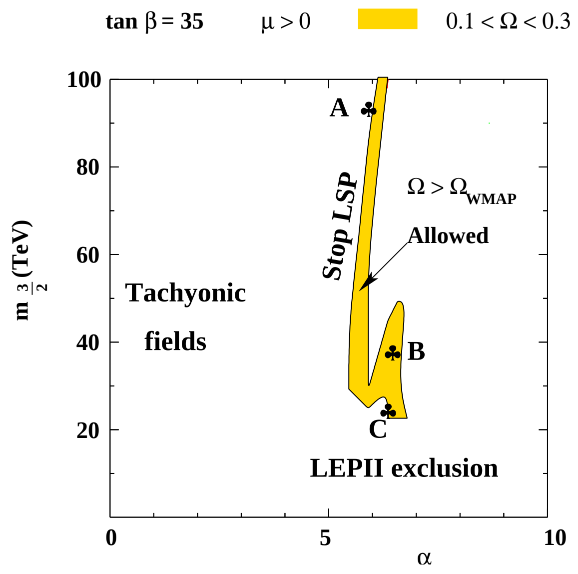

Figure 9: Example of the parameter region consistent with the theoretical and accelerator constraints.

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	A	**B**	**C**
$\tan\beta$	30	30	30
α	5.6	6.5	6.4
$m_{3/2}$ (TeV)	90	33	20
M_1	2342	912	1242
M_2	2838	1176	877
M_3	4711	2221	617
$m_{\chi_1^0}$	2341	911	533
$m_{\chi_2^0}$	2908	1206	709
$m_{\chi_1^+}$	2908	1206	709
$m_{\tilde{g}}$	4837	2286	1413
m_h	130	127	124
m_A	3961	1840	1121
m_H	3961	1839	1120
μ	3621	1768	1094
$m_{\tilde{t}_1}$	2462	1176	709
$m_{\tilde{t}_2}$	3676	1787	1144
$m_{\tilde{c}_1},\ m_{\tilde{u}_1}$	4862	2293	1414
$m_{\tilde{c}_2},\ m_{\tilde{u}_2}$	4738	2231	1376
$m_{\tilde{b}_1}$	3683	1756	1085
$m_{\tilde{b}_2}$	4198	1988	1235
$m_{\tilde{s}_1},\ m_{\tilde{d}_1}$	4863	2295	1416
$m_{\tilde{s}_2},\ m_{\tilde{d}_2}$	4728	2225	1373
$m_{\tilde{\tau}_1}$	2518	1031	601
$m_{\tilde{\tau}_2}$	3134	1355	816
$m_{\tilde{\mu}_1},\ m_{\tilde{e}_1}$	3461	1506	903
$m_{\tilde{\mu}_2},\ m_{\tilde{e}_2}$	3278	1407	841
$m_{\tilde{\nu}_3}$	3132	1348	804
Ωh^2	0.095	0.117	0.110

Table 2: Sample spectra. All masses are in GeV, except for $m_{3/2}$.

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Acceptable abundance of dark matter can be obtained either due to stop–coannihilation or the A-funnel. These results agree with [40, 41].

We thank the authors of [40], K. Choi, V. Löwen and H.P. Nilles for important communications.

Added references

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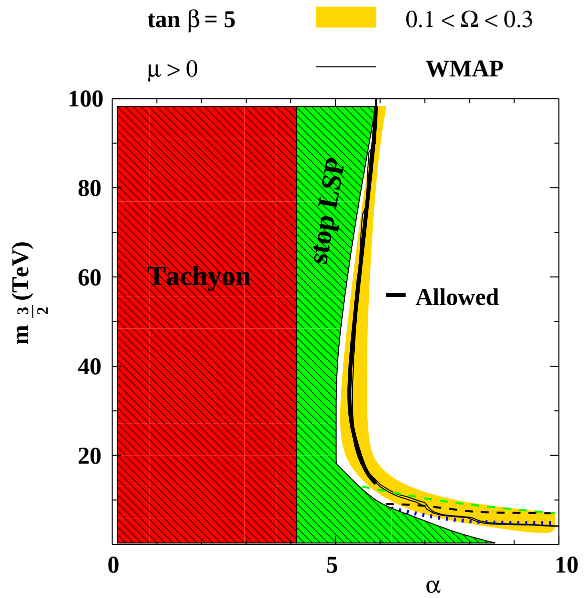

Figure 10: Constraints on the parameter space (see Fig.3 of the main paper).

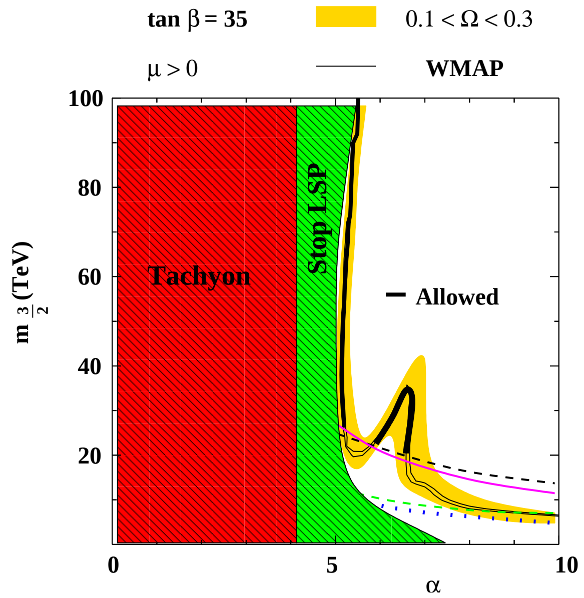

Figure 11: Constraints on the parameter space (see figure 3 of the main paper).