Problems for Supersymmetry Breaking by the Dilaton in Strings from Charge and Color Breaking

J.A. Casas†,§, A. Lleyda¶, and C. Muñoz¶

§ Santa Cruz Institute for Particle Physics
University of California, Santa Cruz, CA 95064, USA
casas@cc.csic.es

¶ Departamento de Física Teórica C–XI
Universidad Autónoma de Madrid, 28049 Madrid, Spain
amanda@delta.ft.uam.es cmunoz@ccuam3.sdi.uam.es

Abstract

The general constraints on the parameter space of soft-breaking terms, in order to avoid dangerous charge and color breaking minima, are applied to the four-dimensional string scenario where the dilaton is the source of supersymmetry breaking (dilaton-dominated limit). The results indicate that the whole parameter space is excluded on these grounds after imposing the present experimental data on the top mass. The inclusion of a non-vanishing cosmological constant does not improve essentially the prospects. Possible way-outs to this situation are briefly discussed.
1 Introduction

The presence of scalar fields with color and electric charge in supersymmetric (SUSY) theories induces the possible existence of dangerous charge and color breaking minima, which would make the standard vacuum unstable [1-8]. This is not necessarily a shortcoming since many SUSY models can be discarded on these grounds, thus improving the predictive power of the theory. A complete analysis of all the potentially dangerous directions in the field space of the minimal supersymmetric standard model (MSSM) was carried out in ref.[6]. It was shown there that the corresponding constraints on the soft parameter space \((m, M, A, B)\) are very strong. As a matter of fact, there are extensive regions of this space that become forbidden producing important bounds, not only on the value of the trilinear scalar term \((A)\), but also on the values of the bilinear scalar term \((B)\) and the scalar and gaugino masses \((m, M)\) respectively.

On the other hand, in four-dimensional strings, working at the perturbative level, it is possible to obtain information about the structure of soft SUSY-breaking terms [9-14]. The basic idea is to identify some chiral fields whose auxiliary components could break SUSY by acquiring a vacuum expectation value (VEV). This is the case of the dilaton and the moduli fields. The important point in this assumption of locating the seed of SUSY breaking in the dilaton/moduli sector, is that it leads to some interesting relationships among different soft terms which could perhaps be experimentally tested. This general analysis was applied in particular to the gaugino condensation scenario in ref.[9], whereas in refs.[10-14] no special assumption was made about the possible origin of SUSY breaking.

The dilaton-dominated limit [11, 12], where only the dilaton field, \(S\), contributes to SUSY breaking is specially interesting. The dilaton field, whose VEV determines the tree-level gauge coupling, is present in any four-dimensional string and couples at tree-level in a universal manner to all particles. Therefore, this limit is model independent and, as a consequence, the soft terms are independent of the four-dimensional string considered. Besides, their expressions are quite simple since they are universal and essentially depend on only two parameters, \(m_{3/2}\) and \(B\), where \(m_{3/2}\) is the gravitino mass. Actually, universality is a desirable property, not only to reduce the number of independent parameters in the MSSM, but also for phenomenological reasons, particularly to avoid flavour changing neutral currents. Because of the simplicity of this scenario, the corresponding low-energy predictions are quite precise [16, 12, 17]. For example, the first and second generation squarks are almost degenerate with the gluino and are much heavier than sleptons.

From all the above reasons it is clearly of the utmost importance to study the consistency of the dilaton-dominated scenario with the possible existence of dangerous charge and color breaking minima. This is the aim of this paper. In fact, we will show that charge and color breaking constraints are so important that the whole parameter space is forbidden and, as a consequence, the dilaton-dominated limit is excluded on these grounds.

---

1 For possible explicit SUSY-breaking mechanisms where this limit might be obtained see ref.[15].
2 The phenomenology of SUSY breaking by the dilaton in the context of a flipped SU(5) model was also studied in ref.[18].
2 Basic Ingredients

Let us briefly review the basic ingredients required for this analysis. First we will concentrate on the form of soft breaking terms. The general form of the soft SUSY-breaking Lagrangian in the context of the MSSM is given by

\[ \mathcal{L}_{soft} = \frac{1}{2} \sum_{a=1}^{3} M_a \lambda_a \lambda_a - \sum_i m_i^2 |\phi_i|^2 - (A_{ijk} W_{ijk} + B \mu H_1 H_2 + \text{h.c.}), \] (1)

where \( W_{ijk} \) are the usual terms of the Yukawa superpotential of the MSSM with \( i = Q_L, u_R, d_R, L, e_R, H_1, H_2 \), and \( \phi_i, \lambda_a \) are the canonically normalized scalar and gaugino fields respectively. In the dilaton-dominated scenario \[11, 12\], neglecting string loop corrections, one obtains the following results for the above scalar masses, gaugino masses and soft trilinear terms

\[ m_i^2 = m_{3/2}^2 + V_0, \]
\[ M_a = \sqrt{3 m_{3/2}^2 + V_0} \ e^{-i\alpha}, \]
\[ A_{ijk} = -M_a, \]

(2)

where, for the sake of completeness, we have included the VEV of the scalar potential (i.e. the cosmological constant) \( V_0 \), and a possible phase \( \alpha \) of the dilaton F-term \[12\]. Notice that we are using the standard supergravity mass units where \( M_{\text{Planck}}/\sqrt{8\pi} = 1 \).

The value of the bilinear term \( B \) is more model dependent and deserves some additional comments. Indeed, \( B \) depends not only on the dilaton-dominance assumption but also on the particular mechanism which could generate the associated (electroweak size) \( \mu \) term \[19\]. For example, the interesting possibilities of generating it through the Kähler potential \[20, 21, 11, 22, 23\] or the superpotential \[21, 23\] give rise, in the dilaton-dominated limit, to the following value of \( B \) \[12, 11, 24, 19, 14\]:

\[ B = 2m_{3/2} + \frac{V_0}{m_{3/2}}. \] (3)

The previous expressions for the soft terms can be simplified taking into account several experimental restrictions. From the limits on the electric dipole moment of the neutron it seems reasonable to impose in what follows \( \alpha = 0 \mod \pi \). On the other hand, experimental constraints in present cosmology allow us to assume vanishing cosmological constant \( V_0 = 0 \) (we will see later on that our conclusions will not be modified if we give up this assumption). Then

\[ m_i^2 = m_{3/2}^2, \]
\[ M_a = \pm \sqrt{3} m_{3/2}, \]
\[ A_{ijk} = -M_a, \]

(4)

and the \( B \) term associated with the mechanisms explained above in order to solve the \( \mu \) problem is\(^3\)

\[ B = 2m_{3/2}. \] (5)

\(^3\)Phenomenological aspects of the \( B=2m_{3/2} \) scenario have been studied in refs. [16, 22, 23].
Expressions (2-5) are to be understood at the string scale $M_{\text{string}} \simeq 0.5 \times g_{\text{string}} \times 10^{18}$ GeV [20], where $g_{\text{string}} = (\Re S)^{-1/2} \simeq 0.7$. In the following we will assume that the discrepancy between the unification scale of the gauge couplings, $M_X$, and the string unification scale, $M_{\text{string}}$, can be explained by the effect of string threshold corrections [27].

In the present paper we have taken the expressions of the soft terms given by eqs.(4, 5) as our starting point to work, considering the value of $B$ given by eq.(5) as a guiding example. Then, by varying the value of $B$ (and also $V_0$) we will eventually obtain the most general results. Concerning the value of the $\mu$ parameter, we will fix it as usual from the requirement of correct electroweak breaking [4].

The second basic ingredient of our analysis concerns the constraints associated with the existence of dangerous directions in the field space. As was mentioned in the introduction, a complete analysis of this issue, including in a proper way the radiative corrections to the scalar potential, was carried out in ref.[6]. The most relevant results obtained there for our present task are the following.

There are two types of constraints: the ones arising from directions in the field-space along which the (tree-level) potential can become unbounded from below (UFB), and those arising from the existence of charge and color breaking (CCB) minima in the potential deeper than the standard minimum.

Concerning the UFB directions (and corresponding constraints), there are three of them, labelled as UFB-1, UFB-2, UFB-3 in [3]. It is worth mentioning here that in general the unboundedness is only true at tree-level since radiative corrections eventually raise the potential for large enough values of the fields, but still these minima can be deeper than the realistic one (i.e. the SUSY standard model vacuum) and thus dangerous. The UFB-3 direction, which involves the fields $\{H_2, \nu_{L_i}, e_{L_j}, e_{R_j}\}$ with $i \neq j$ and thus leads also to electric charge breaking, yields the strongest bound among all the UFB and CCB constraints. The explicit form of this bound is as follows. By simple analytical minimization it is possible to write the value of all the relevant fields along the UFB-3 direction in terms of the $H_2$ one. Then, for any value of $|H_2| < M_{\text{string}}$ satisfying

$$|H_2| > \frac{\mu^2}{4\lambda_{e_j}^2} + \frac{4m_{L_i}^2}{g'^2 + g_2^2} - \frac{|\mu|}{2\lambda_{e_j}},$$

the value of the potential along the UFB-3 direction is simply given by

$$V_{\text{UFB-3}} = (m_2^2 - \mu^2 + m_{L_i}^2)|H_2|^2 + \frac{|\mu|}{\lambda_{e_j}}(m_{L_j}^2 + m_{e_j}^2 + m_{L_i}^2)|H_2| - \frac{2m_{L_i}^4}{g'^2 + g_2^2}. \quad (7)$$

Otherwise

$$V_{\text{UFB-3}} = (m_2^2 - \mu^2)|H_2|^2 + \frac{|\mu|}{\lambda_{e_j}}(m_{L_j}^2 + m_{e_j}^2)|H_2| + \frac{1}{8}(g'^2 + g_2^2) \left[ |H_2|^2 + \frac{|\mu|}{\lambda_{e_j}}|H_2|^2 \right]^2. \quad (8)$$

4The value of $\mu$ can also be fixed once we choose a particular mechanism for solving the $\mu$ problem, e.g. if $\mu$ is generated through the Kähler potential in dilaton-dominated orbifold models, then $\mu = m_3/2$ [14, 24]. However, we prefer to eliminate $\mu$ in terms of the other parameters by imposing appropriate electroweak breaking, as mentioned above, the reason being that this provides the most general analysis.
In eqs.(7,8) $\lambda_{e_j}$ is the leptonic Yukawa coupling of the $j-$generation and $m_2^2$ is the sum of the $H_2$ squared soft mass, $m_{H_2}^2$, plus $\mu^2$. Then, the UFB-3 condition reads

$$V_{\text{UFB-3}}(Q = \hat{Q}) > V_{\text{real min}} \ ,$$

where $V_{\text{real min}} = -\frac{1}{8} (g'^2 + g_2^2) (v_2^2 - v_1^2)^2$, with $v_{1,2}$ the VEVs of the Higgses $H_{1,2}$, is the realistic minimum evaluated at $M_S$ (see below) and the $\hat{Q}$ scale is given by

$$\hat{Q} \sim \text{Max}(g_2 |e|, \lambda_{top} |H_2|, g_2 |H_2|, g_2 |L_i|, M_S) \quad \text{with} \quad |e| = \frac{|\mu|}{\lambda_{e_j} |H_2|} \quad \text{and} \quad |L_i|^2 = -\frac{4m_2^2}{g'^2 + g_2^2} + (|H_2|^2 + |e|^2).$$

Finally, $M_S$ is the typical scale of SUSY masses (normally a good choice for $M_S$ is an average of the stop masses, for more details see refs.[4, 29, 6]). Notice from (4,8) that the negative contribution to $V_{\text{UFB-3}}$ is essentially given by the $m_2^2 - \mu^2$ term, which can be very sizeable in many instances. On the other hand, the positive contribution is dominated by the term $\propto 1/\lambda_{e_j}$, thus the larger $\lambda_{e_j}$ the more restrictive the constraint becomes. Consequently, the optimum choice of the $e-$type slepton is the third generation one, i.e. $e_j = \text{stau}$.

Concerning the CCB constraints, let us mention that the “traditional” CCB bounds [1], when correctly evaluated (i.e. including the radiative corrections in a proper way), turn out to be extremely weak. However, the “improved” set of analytic constraints obtained in ref.[3], which represent the necessary and sufficient conditions to avoid dangerous CCB minima, is much stronger. It is not possible to give here an account of the explicit form of the CCB constraints used in the present paper. This can be found in section 5 of ref.[3], to which we refer the interested reader.

3 Results

In Fig.1 we have presented in detail the interesting case (and guiding example) $B=2m_{3/2}$ for the two possible values of gaugino masses, $M \equiv M_a = \pm \sqrt{3} m_{3/2}$, see eqs.(4, 5). In the plots of the figure we have shown how the parameter space defined by $m_{3/2}$ (the only soft parameter left independent) and $M_{\text{top}}$ (which we let vary for completeness) is totally excluded by the different constraints in the game. Some comments are in order here.

First, we have taken $m_{3/2} \leq 500$ GeV since larger values of $m_{3/2}$ would induce too large SUSY-mass spectra; e.g. $m_{3/2}=500$ GeV implies gluino and first and second generation squark masses of order 2.5 TeV, conflicting the absence-of-fine-tuning requirements [30, 29]. On the other hand, rather than assuming a particular value of the top mass, we have preferred to vary the physical (pole) top mass, $M_{\text{top}}^{\text{phys}}$, between 150 and 200 GeV. Actually, it is not always possible to choose the boundary condition of the top Yukawa coupling $\lambda_{\text{top}}$ so that the physical (pole) mass is reproduced because the renormalization group (RG) infrared fixed point of $\lambda_{\text{top}}$ puts an upper bound on the running top mass $M_{\text{top}}$, namely $M_{\text{top}} \leq 197 \sin \beta$ GeV [31], where $\tan \beta = v_2/v_1$. The corresponding restriction in the parameter space (black regions in Fig.1) is certainly substantial in the case $M=\sqrt{3} m_{3/2}$, yielding $M_{\text{top}}^{\text{phys}} < 167$ GeV, which is by itself a remarkable result.[4]

However, in the case $M=-\sqrt{3} m_{3/2}$ the restriction is small allowing in principle large top masses.
The region excluded by the CCB bounds is denoted by circles in the figure. The sign of the trilinear soft term $A_{ijk}$ is important in these type of constraints as can be seen in the figure (recall that $A \equiv A_{ijk} = -M$). Whereas the case $M = \sqrt{3}m_{3/2}$ is not constrained at all, in the case $M = -\sqrt{3}m_{3/2}$ the whole parameter space left allowed by the previous “top-fixed-point constraint” is excluded by the CCB bounds.

Anyway, it is apparent from Fig.1 that the the restrictions coming from the UFB constraints (small filled squares) are very strong in both cases. Most of the parameter space is in fact excluded by the UFB-3 constraint, which has been explained in the previous section.

Finally, we have also plotted in Fig.1 the region excluded by the experimental bounds on SUSY particle masses (filled diamonds). Conservatively enough, we have imposed

$$
M_{\tilde{g}} \geq 120 \text{ GeV} \ , \ M_{\tilde{\chi}^\pm} \geq 45 \text{ GeV} \ , \ M_{\tilde{\chi}^0} \geq 18 \text{ GeV} \ , \\
M_{\tilde{q}} \geq 100 \text{ GeV} \ , \ M_\ell \geq 45 \text{ GeV} \ , \ M_t \geq 45 \text{ GeV} , \quad (10)
$$

in an obvious notation. The corresponding forbidden area comes mainly from too small masses for neutralinos and charginos in the case $M = \sqrt{3}m_{3/2}$, and for the sleptons when $M = -\sqrt{3}m_{3/2}$. The ants indicate regions which are excluded by negative squared mass eigenvalues, in this case the sneutrinos.

Notice from Fig.1 that there are areas that are simultaneously constrained by different types of bounds.

At the end of the day, the allowed region (white) left is very small. Only in the case $M = -\sqrt{3}m_{3/2}$ and for $m_{3/2}$ larger than 320 GeV, which on the other hand corresponds to gluino and first and second generation squark masses heavier than 1.5 TeV, the dangerous minima are not present. However, this occurs for $M_{\text{top}}^{\text{phys}} < 157$ GeV. Thus, using the present experimental data on the top mass, $M_{\text{top}}^{\text{exp}} = 180 \pm 12$ GeV [13], we conclude that the dilaton dominance limit for SUSY breaking in strings with $B = 2m_{3/2}$ is excluded on charge and color breaking (CCB and UFB) grounds. It is worth mentioning here that, even if we relax the previous fine-tuning requirement (i.e. $m_{3/2} \leq 500$ GeV) by admitting higher values of $m_{3/2}$, we have checked that the growing of the allowed region is anyway remarkably slow, so one has to go to extremely high values of $m_{3/2}$ to get acceptable top masses.

Let us generalize now the previous analysis by varying the value of $B$. In Fig.2 we have shown the representative examples $B=0, 3m_{3/2}$ with $M = -\sqrt{3}m_{3/2}$. Whereas for large values of $B$, $B \geq 3m_{3/2}$, the whole parameter space is excluded, for $B=0$ there is still a very small allowed region. However, this region is in fact excluded by the value of $M_{\text{top}}^{\text{exp}}$ in a similar fashion as it happened in the $B = 2m_{3/2}$ case (see above). Intermediate values of $B$ do not improve the situation. For negative values of $B$ the this possibility was not taken into account in ref. [16] where an upper bound of 180 GeV on the top mass was obtained for $M = \sqrt{3}m_{3/2}$, $A = -\sqrt{3}m_{3/2}$. The discrepancy between that upper bound and the one that we obtain here for the same case, 167 GeV, is due to the following: the RGEs [32] used in ref. [13] correspond to the soft Lagrangian of eq.(1) with a minus sign in front of the gaugino masses and therefore the associated boundary conditions that should have been used in order to get the correct result are $M = -\sqrt{3}m_{3/2}$, $A = -\sqrt{3}m_{3/2}$.
corresponding figures are the same, since they are invariant under the transformation 
\( B, A, M \rightarrow -B, -A, -M \). The same conclusion is obtained for the case \( M = \sqrt{3}m_{3/2} \).
Examples of this scenario are given in Fig.1 for \( B = 2m_{3/2} \) and in Fig.2 for \( B = 0 \) (notice that from the previous invariance the \( B = 0 \) figure is valid for \( M = -\sqrt{3}m_{3/2} \) as well as for \( M = \sqrt{3}m_{3/2} \)).

From the various figures it is clear that the CCB and UFB constraints do not allow the possibility of SUSY breaking in the dilaton-dominance limit in strings.

Let us now consider for the sake of completeness the possibility of a non-vanishing cosmological constant \( V_0 \). This may correspond to different attitudes concerning this problem. For example, one might think that the experimentally constrained cosmological term in present cosmology is not directly connected to the particle physics vacuum energy \( V_0 \). Another possibility is to admit a non-vanishing tree-level cosmological constant, requiring a vanishing fully renormalized one \cite{14}. Anyway, whatever the assumption is, it is worth relaxing the condition \( V_0 = 0 \) since, from eq.(2), this might be the only possibility to avoid the previous dramatical conclusions for the dilaton-dominated limit. Notice, however, that \( V_0 \) is constrained to be \( V_0 \geq -m_{3/2}^2 \) in order to avoid negative squared scalar masses (see eq.(2)) and, on the other hand, it should not be too large in order to keep the SUSY spectrum in the range of 1 TeV. In addition, to simplify the analysis we have initially taken the (theoretically well-motivated) value of \( B \) given in eq.(3).

In Fig.3 we have presented three representative examples of a non-vanishing cosmological constant \( (V_0 = -0.5m_{3/2}, 3m_{3/2}, 10m_{3/2}) \) with \( M = -\sqrt{3}m_{3/2} \). We see that for \( V_0 = -0.5m_{3/2} \) the whole parameter space is excluded. Larger values of \( V_0 \) do improve the situation, the best case being \( V_0 = 3m_{3/2} \). However, once more, the present value of \( M_{\text{phys}}^\text{top} \) essentially excludes this possibility for any reasonable SUSY spectrum. More precisely, notice that for \( m_{3/2} \geq 300 \) GeV we are already in the range of 2 TeV for the gluino and first and second generation squark masses, which are far too large from fine-tuning considerations. Finally, for the case \( M = \sqrt{3}m_{3/2} \) (not represented in the figures) the constraints are even stronger, e.g. for \( V_0 = 3m_{3/2} \) the whole parameter space is excluded.

Even if we let \( B \) vary as an independent parameter (at the same time as \( V_0 \)) the results do not improve much. More precisely, the best case is found for \( V_0 = 3m_{3/2} \), \( B = 3m_{3/2} \). But, even in this extreme possibility, if one demands a reasonable SUSY spectrum (e.g. masses of order 1 TeV, which in this case requires \( m_{3/2} \sim 150 \) GeV) the top mass becomes too small (\( M_{\text{phys}}^\text{top} \leq 172 \) GeV), almost inconsistent with the present experimental lower bound (\( M_{\text{phys}}^\text{top} \geq 168 \) GeV).

To summarize the results, the dilaton-dominated limit is essentially excluded on charge and color breaking (CCB and UFB) grounds. Even in the few extreme cases consistent with the charge and color breaking constraints, the spectrum is inviable since either the top mass is too small or the SUSY spectrum is far too heavy from any fine-tuning consideration. The addition of a non-vanishing cosmological constant in the game does not improve this situation.
4 Concluding Remarks

The dilaton-dominated limit in strings, defined as the scenario where the only source of SUSY breaking is the dilaton, is highly interesting since at the string tree-level approximation its formulation is model independent (i.e. holds for any four-dimensional string) and yields quite precise low-energy predictions (SUSY spectrum). However, we have seen along this paper that, after imposing the present experimental data on the top mass, the whole parameter space of this scenario \((m_3/2, B)\) is excluded on charge and color breaking grounds \([6]\), i.e. by the existence of charge and color breaking minima deeper than the standard vacuum. Even allowing a non-vanishing cosmological constant does not improve essentially the situation. Due to the attractiveness of the dilaton-dominated limit, let us briefly discuss some possible way-outs to the previous dramatical conclusions.

One possibility is to accept that we live in a metastable vacuum, provided that its lifetime is longer than the present age of the universe \([4, 8]\), thus rescuing some points in the parameter space. This possibility, however, poses the cosmological problem of why our universe does not correspond to the global minimum of the potential (without invoking an anthropic principle). Even if a solution to that problem is found we would still have to face the rather bizarre situation of a future cosmological catastrophe, which does not seem very attractive. In addition, it is hard to understand how the cosmological constant is vanishing precisely in such “interim” vacuum. Anyway, despite the previous shortcomings, this is still a possible scenario which would be worth analyzing \([35]\).

A different possibility is to assume that also the moduli fields contribute to SUSY breaking. Then the soft terms are modified and possibly some regions in the parameter space would be allowed. This more general situation deserves further analysis \([35]\). Of course, this amounts to a departure of the pure dilaton-dominated scenario. On the other hand, it is interesting to note that explicit possible scenarios of SUSY breaking by gaugino condensation in strings, when analyzed at the one-loop level, lead to the mandatory inclusion of the moduli in the game (in fact the moduli are the main source of SUSY breaking in these cases) \([9, 36]\).

Finally, one may think that the perturbative and non-perturbative corrections to the “standard” tree-level dilaton-dominated scenario are important and can modify the previous conclusions. Due to analyticity and non-renormalization theorems these contributions are likely to affect in a substantial way only the Kähler potential \([37]\) (in fact, the gauge kinetic function does also receive perturbative corrections at one loop level, but not beyond). Actually, the one-loop string corrections to the Kähler potential (and the gauge kinetic function) have been calculated for orbifold models \([38]\) and they are rather small for sensible values of the moduli. Thus, it is reasonable to expect that further perturbative corrections will be even smaller. However, this is not the case for the string non-perturbative corrections, whose size could be much larger (see e.g. ref.\([39]\)). These corrections could be crucial to understand both the SUSY-breaking mechanism and the vanishing of the cosmological constant \([37]\); actually, it is possible to show \([40]\) that a tree-level dilaton-dominated scenario cannot have a global minimum of the dilaton potential at vanishing cosmological constant. Unfortunately,
the form of this type of corrections is very poorly known, which introduces additional sources of uncertainty in the analysis [37, 40].

Acknowledgements

We thank L.E. Ibáñez for useful comments.

References

[1] J.M. Frere, D.R.T. Jones and S. Raby, *Nucl. Phys.* **B222** (1983) 11; L. Alvarez-Gaumé, J. Polchinski and M. Wise, *Nucl. Phys.* **B221** (1983) 495; J.P. Derendinger and C.A. Savoy, *Nucl. Phys.* **B237** (1984) 307; C. Kounnas, A.B. Lahanas, D.V. Nanopoulos and M. Quirós, *Nucl. Phys.* **B236** (1984) 438.

[2] M. Claudson, L.J. Hall and I. Hinchliffe, *Nucl. Phys.* **B228** (1983) 501.

[3] M. Drees, M. Glück and K. Grassie, *Phys. Lett.* **B157** (1985) 164; J.F. Gunion, H.E. Haber and M. Sher, *Nucl. Phys.* **B306** (1988) 1; H. Komatsu, *Phys. Lett.* **B215** (1988) 323.

[4] G. Gamberini, G. Ridolfi and F. Zwirner, *Nucl. Phys.* **B331** (1990) 331.

[5] P. Langacker and N. Polonsky, *Phys. Rev.* **D50** (1994) 2199; A. Bordner, *KUNS-1351*, [hep-ph/9506409](http://arxiv.org/abs/hep-ph/9506409).

[6] J.A. Casas, A. Lleyda and C. Muñoz, *FTUAM 95/11*, [hep-th/9507294](http://arxiv.org/abs/hep-th/9507294), to be published in *Nuclear Physics B*.

[7] T. Falk, K. Olive, L. Roszkowski and M. Srednicki, *UMN-TH-1411/95*, [hep-ph/9510308](http://arxiv.org/abs/hep-ph/9510308).

[8] A. Riotto and E. Roulet, *SISSA-163/95/EP*, [hep-ph/9512401](http://arxiv.org/abs/hep-ph/9512401).

[9] A. Font, L.E. Ibáñez, D. Lüst and F. Quevedo, *Phys. Lett.* **B245** (1990) 401; M. Cvetič, A. Font, L.E. Ibáñez, D. Lüst and F. Quevedo, *Nucl. Phys.* **B361** (1991) 194; B. de Carlos, J.A. Casas and C. Muñoz, *Phys. Lett.* **B299** (1993) 234; *Nucl. Phys.* **B399** (1993) 623; A. de la Macorra and G.G. Ross, *Phys. Lett.* **B325** (1994) 85.

[10] L.E. Ibáñez and D. Lüst, *Nucl. Phys.* **B382** (1992) 305.

[11] V.S. Kaplunovsky and J. Louis, *Phys. Lett.* **B306** (1993) 269.

[12] A. Brignole, L.E. Ibáñez and C. Muñoz, *Nucl. Phys.* **B422** (1994) 125 [Erratum: *B436* (1995) 747].
[13] S. Ferrara, C. Kounnas and F. Zwirner, *Nucl. Phys.* **B429** (1994) 589 [Erratum: **B433** (1995) 255].

[14] A. Brignole, L.E. Ibáñez, C. Muñoz and C. Scheich, *FTUAM 95/26*, [hep-ph/9508253](https://arxiv.org/abs/hep-ph/9508253).

[15] A. de la Macorra and G.G. Ross, *Nucl. Phys.* **B404** (1993) 321; V. Halyo and E. Halyo, *SU-ITP-96-4*, [hep-ph/9601328](https://arxiv.org/abs/hep-ph/9601328).

[16] R. Barbieri, J. Louis and M. Moretti, *Phys. Lett.* **B312** (1993) 451 [Erratum: **B316** (1993) 632].

[17] S. Khalil, A. Masiero and F. Vissani, *SHEP-95-30*, [hep-ph/9511284](https://arxiv.org/abs/hep-ph/9511284).

[18] J.L. Lopez, D.V. Nanopoulos and A. Zichichi, *Phys. Lett.* **B319** (1993) 451.

[19] For a recent review, see: C. Muñoz, ”Soft supersymmetry-breaking terms and the $\mu$ problem”, *FTUAM 95/20*, [hep-th/9507108](https://arxiv.org/abs/hep-th/9507108).

[20] G.F. Giudice and A. Masiero, *Phys. Lett.* **B206** (1988) 480.

[21] J.A. Casas and C. Muñoz, *Phys. Lett.* **B306** (1993) 288.

[22] G. Lopes-Cardoso, D. Lüst and T. Mohaupt, *Nucl. Phys.* **B432** (1994) 68.

[23] I. Antoniadis, E. Gava, K.S. Narain and T.R. Taylor, *Nucl. Phys.* **B432** (1994) 187.

[24] C. Muñoz, Proceedings of the International Workshop ”Beyond the Standard Model IV”, *World Scientific* (1995) 200, [hep-ph/9503314](https://arxiv.org/abs/hep-ph/9503314).

[25] B. de Carlos and J.A. Casas, *Phys. Lett.* **B349** (1995) 300 [Erratum: **B351** (1995) 604]; M.A. Diaz and S.F. King, *SHEP-95-30*, [hep-ph/9601230](https://arxiv.org/abs/hep-ph/9601230).

[26] V. Kaplunovsky, *Nucl. Phys.* **B307** (1988) 145 [Erratum: **B382** (1992) 436].

[27] K. Choi, *Phys. Rev.* **D37** (1988) 1564; L.E. Ibáñez, D. Lüst and G.G. Ross, *Phys. Lett.* **B272** (1991) 251; L.E. Ibáñez and D. Lüst, *Nucl. Phys.* **B382** (1992) 305; H.P. Nilles and S. Stieberger, *TUM-HEP-225/95*, [hep-th/9510009](https://arxiv.org/abs/hep-th/9510009).

[28] C. Muñoz, Proceedings of the 5th Hellenic School and Workshops on Elementary Particle Physics, Corfu (1995), *FTUAM 96/04*, [hep-ph/9601323](https://arxiv.org/abs/hep-ph/9601323).

[29] B. de Carlos and J.A. Casas, *Phys. Lett.* **B309** (1993) 320.

[30] R. Barbieri and G.F. Giudice, *Nucl. Phys.* **B306** (1988) 63.
[31] K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, *Prog. Theor. Phys.* **67** (1982) 1889;
L. Ibáñez and C. López, *Phys. Lett.* **B126** (1983) 54;
L. Alvarez-Gaumé, J. Polchinski and M. Wise, in ref.[1].

[32] L. Ibáñez and C. López, *Nucl. Phys.* **B233** (1984) 511;
L. Ibáñez, C. López and C. Muñoz, *Nucl. Phys.* **B256** (1985) 218.

[33] L. Montanet et al., *Phys. Rev.* **D50** (1994) 1173 and 1995 off-year partial update for the 1996 edition available on PDG WWW pages (URL: http://pdg.lbl.gov/).

[34] K. Choi, J.E. Kim and H.P. Nilles, *Phys. Rev. Lett.* **73** (1994) 1758.

[35] J.A. Casas, A. Lleyda and C. Muñoz, in preparation.

[36] S. Ferrara, N. Magnoli, T.R. Taylor and G. Veneziano, *Phys. Lett.* **B245** (1990) 409.

[37] T. Banks and M. Dine, *Phys. Rev.* **D50** (1994) 7454.

[38] L.J. Dixon, V. Kaplunovsky and J. Louis, *Nucl. Phys.* **B355** (1991) 649;
J.P. Derendinger, S. Ferrara, C. Kounnas and F. Zwirner, *Nucl. Phys.* **B372** (1992) 145.

[39] S.H. Shenker, Proceedings of the Cargese Workshop on Random Surfaces, Quantum Gravity and Strings, Cargese (France) 1990.

[40] Work in progress.
Figure Captions

**Fig.1** Excluded regions in the parameter space of the MSSM assuming SUSY breaking by the dilaton, with $B=2m_{3/2}$. The black region is excluded because it is not possible to reproduce the experimental mass of the top. The small filled squares indicate regions excluded by Unbounded From Below constraints. The circles indicate regions excluded by Charge and Color Breaking constraints. The filled diamonds correspond to regions excluded by the experimental lower bounds on SUSY-particle masses. The ants indicate regions excluded by negative scalar squared mass eigenvalues.

**Fig.2** The same as Fig.1 but with $B=0$, $3m_{3/2}$ and $M=-\sqrt{3}m_{3/2}$.

**Fig.3** The same as Fig.1 but with $M=-\sqrt{3}m_{3/2}$ and $V_0=-0.5m_{3/2}^2$, $3m_{3/2}^2$, $10m_{3/2}^2$.
Fig. 1
Fig. 3