NON–UNIVERSAL SOFT SCALAR MASSES IN SUPERSYMMETRIC THEORIES

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Abstract

The existence of non–universal soft masses is the most general situation in supersymmetric theories. We study the consequences that this situation has for the low–energy sparticle spectrum. In particular, we analyze in detail the contribution to the scalar mass renormalization group equations of the $U(1)_Y$ D–term. We obtain analytic expressions for the evolution of masses of the three generations and these allow us to show that such a contribution can produce important modifications on the spectrum. The necessity to avoid flavour changing neutral currents does not constrain this result. Finally, we discuss a realistic example in the context of string theory where the departure from universality is large.

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Supersymmetry (SUSY) is very important in order to solve the gauge hierarchy problem and to unify all fundamental interactions. The simplest supersymmetric extension of the Standard Model is the so-called Minimal Supersymmetric Standard Model (MSSM). It is an effective low–energy supergravity (SUGRA) theory with the Standard Model gauge group, global $N = 1$ SUSY broken softly and the minimal particle content (i.e. three generations of quark and lepton superfields plus two Higgs doublets, $H_1$ and $H_2$, of opposite hypercharge as well as the gauge superfields) $[1]$. In particular, the soft scalar masses generated after SUSY–breaking in the effective observable scalar potential are of the type:

$$V_{eff} = \sum_{\phi} m_\phi^2 |\phi|^2$$

(1)

The rich spectrum of sparticles of the MSSM has been studied almost exclusively under the assumption $[1] m_\phi = m$, i.e. the scalars have universal masses at the unification scale ($M_{\text{GUT}}$). In this way, after renormalizing the theory, all the physical masses can be calculated in terms of a small number of soft parameters and therefore the analysis is simplified. If, in the SUGRA theory under consideration, the kinetic terms of the fields are canonical then the above assumption is true and $m$ is the gravitino mass. However, for a general Kähler manifold, the relation $m_\phi = m$ is not fulfilled and non–universal scalar masses can emerge:

$$m_\phi = c_\phi m$$

(2)

where $c_\phi$ are constants whose values will depend on the specific mechanism of SUSY–breaking. Now the analysis of the sparticle spectrum is more difficult since new soft parameters must be taken into account. In fact, this is what happens in the context of SUGRA theories coming from superstrings, where the soft masses of the scalar fields can be explicitly calculated and they show a lack of universality $[3-5]$ unlike the usual assumption of the MSSM.

$^1$Common gaugino masses are also usually considered in the analyses of the MSSM. See ref.$[2]$ for exceptions, where the effect of relaxing this assumption on neutralinos is studied.
In this paper we study the consequences that the departure from the common assumption of universal soft scalar masses has for the low–energy SUSY spectrum. In particular, we analyze in detail the contribution to the scalar mass renormalization group equations (RGEs) of the \( U(1)_Y \) D–term (see eq.\((3)\)). We give analytic expressions for the evolution of masses of the three generations and these allow us to show that such a contribution can produce important modifications on the spectrum. The necessity to avoid flavour changing neutral currents (FCNC) does not constrain this result. Finally, we discuss a realistic example in the context of string theory where the departure from universality is large.

Let us consider for the moment the case of first– and second–generation sparticles. To calculate the soft SUSY breaking masses at low energies from the parameters at the unification scale, we have to use the RGEs \([6-9]\). Although Yukawa couplings contribute in the RGEs, their small effects can be safely neglected unlike the case of the third generation, where top, bottom, and tau Yukawa couplings can be large. Then

\[
\frac{dm^2_\phi}{dt} = 3 \sum_{i=1}^{3} \frac{\alpha_i}{\pi} C^\phi_i M_i^2 - \frac{\alpha_1}{4\pi} Y^\phi S ,
\]  

(3)

where \( t \equiv \ln \frac{M_\text{GUT}}{Q^2} \), \( M_i \) are the gaugino masses, \( Y^\phi \) are the hypercharges, \( C^\phi_i \) is the quadratic Casimir corresponding to each scalar (\( C = \frac{N^2-1}{2N} \) for \( SU(N) \), \( C = Y^2 \) for \( U(1)_Y \)) and the gauge coupling constants at \( M_\text{GUT} \) verify \( \alpha_3(0) = \alpha_2(0) = \frac{5}{3} \alpha_1(0) = \alpha_\text{GUT} \). Notice that in these equations there is a contribution to the scalar masses coming from the \( U(1)_Y \) D–term\(^2\) parametrized in the general case by:

\[
S = \sum_\phi d(R_\phi) Y^\phi m^2_\phi ,
\]  

(4)

where \( d(R_\phi) \) is the dimension of the representation \( R \) associated to the \( \phi \) scalar. The

\(^2\)This term was also included in a different context in the analysis of ref.\([10]\). The authors studied the phenomenological consequences of spontaneous breaking of intermediate scale gauge symmetries induced by universal soft terms.
evolution of $S$ is given by \[9\]:

$$
\frac{dS}{dt} = -\frac{\alpha_1}{4\pi} \left( \sum_{\phi} d(R_\phi)Y^2_{\phi} \right) S .
$$

(5)

For the MSSM eqs.\(\text{(4,5)}\) read

$$
S = m^2_2 - m^2_1 + \sum_{\text{generations}} \left( m^2_{Q_L} + m^2_{D_R} + m^2_{E_R} - m^2_{L_L} - 2m^2_{U_R} \right) ,
$$

(6)

$$
\frac{dS}{dt} = -\frac{\alpha_1}{4\pi} b_1 S ,
$$

(7)

where $U_R$, $D_R$ ($Q_L = (U_L, D_L)$) stand for the right–handed (left–handed) squarks, $E_R$ ($L_L = (V_L, E_L)$) for the right–handed (left–handed) sleptons and $b_1 = 11$ is the one–loop coefficient of the beta function calculated using the following hypercharges: $6Y_{Q_L} = 3Y_{D_R} = Y_{E_R} = -2Y_{L_L} = -\frac{3}{2}Y_{U_R} = 1$. Finally, $m^2_{1,2} \equiv m^2_{H_{1,2}} + \mu^2$, where $\mu$ is the supersymmetric mass term in the superpotential $\mu H_{1}H_{2}$. The solution of eq.\(\text{(5)}\)

$$
S(t) = \frac{S(0)}{1 + \frac{\alpha_1(0)b_1}{4\pi} t} ,
$$

(8)

implies that if $S$ vanishes at some scale, it vanishes everywhere. For the case of universal boundary conditions for the scalar masses $m_\phi(0) = m$, $S(0) = 0$ (see eq.\(\text{(3)}\)). However in the most general case of non–universality, $S(0) \neq 0$ and therefore $S$ must be included in the RGEs \(\text{(3)}\). We will see below that although this term in eq.\(\text{(3)}\) is multiplied by $\alpha_1$, its influence can still be sufficiently large as to modify substantially the low–energy scalar masses. It is worth noticing here that although the necessity to avoid large FCNC constrains the amount of non–universality and suggests that, in a viable model, d– and s–squarks are nearly degenerate at $M_{GUT}$, and likewise for the u– and c–squarks and the selectron and smuon, different scalars within the same generation can have different masses. This allows the second term in eq.\(\text{(3)}\) to be large. Besides, the first term depends on the difference between the masses of the Higgses and therefore is not constrain by FCNC either.
Eq. (3) can be analytically solved and for the MSSM we obtain:

\[ m_\phi^2(t) = m_\phi^2(0) + \sum_{i=1}^{3} M_i^2(0) f_i^\phi(t) - Y_\phi S(0) f_1'(t) + M_Z^2(T^\phi_{3L} \cos^2 \theta_W - Y_\phi \sin^2 \theta_W) \cos 2\beta \]  

(9)

where

\[ f_i^\phi(t) = \frac{2 C_i^\phi}{b_i} \left[ 1 - \frac{1}{(1 + \frac{\alpha_i(0)}{4\pi} b_i t)^2} \right], \quad f_1'(t) = \frac{1}{b_1} \left[ 1 - \frac{1}{(1 + \frac{\alpha_1(0)}{4\pi} b_1 t)^2} \right], \]

\[ \tan \beta = \frac{\langle H_2 \rangle}{\langle H_1 \rangle}; \quad b_3 = -3, \quad b_2 = 1, \quad b_1 = 11. \]  

(10)

In eq.(9) we have already included the contribution of the scalar potential D–terms to the scalar masses after the Higgs bosons get vacuum expectation values (VEVs). One has to add to eq.(9) the \((mass)^2\) of the corresponding fermionic partner. The \(f's\) evaluated at \(M_Z\) are:

\[ f_1'(M_Z) = 0.054, \quad f_3^\phi(M_Z) = 7.093, \quad f_2^\phi(M_Z) = 0.496, \quad f_1^\phi(M_Z) = 0.152 \ Y_\phi^2 \]  

(11)

where the different \(\alpha_i(0)\) were computed by running up to \(M_{GUT}\) the three gauge coupling constants obtained from the following experimental inputs:

\[ M_Z = 91.175 \ \text{GeV}, \quad \alpha_3(M_Z) = 0.125, \quad \alpha_{em}(M_Z) = \frac{1}{127.9}, \quad \sin^2 \theta_W(M_Z) = 0.23 \]  

(12)

As discussed above, in general, \(m_\phi(0) = c_\phi m\) so it is more useful for us to write eq.(9) as

\[ m_\phi^2(t) = (c_\phi^2 - c^2 Y_\phi f_1'(t)) \ m^2 + \sum_{i=1}^{3} M_i^2(0) f_i^\phi(t) + M_Z^2(T^\phi_{3L} \cos^2 \theta_W - Y_\phi \sin^2 \theta_W) \cos 2\beta \]  

(13)

where the influence of the \(U(1)_Y\) D–term becomes apparent as a modification of the soft masses at \(M_{GUT}\). The value of \(c^2\) is defined as (see eq.(3)):

\[ c^2 \equiv c_2^2 - c_1^2 + \sum_{\text{generations}} \left( c_{Q_L}^2 + c_{D_R}^2 + c_{E_R}^2 - c_{L_L}^2 - 2 c_{U_R}^2 \right). \]  

(14)

Notice that \(c_\phi^2\) must be a positive number, otherwise the mass squared of the scalars would already become negative at \(M_{GUT}\). However, let us remark that \(c^2\) can be a negative number. Finally, the universal case is given by \(c^2 = 0.\)

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3 Notice to obtain the physical scalar masses in eq.(9) the RGEs must be integrated to \(t = \ln \frac{M_{GUT}^2}{m_\phi^2}.\) We will take into account this subtleness in the figures below.
Let us now discuss the importance that the $U(1)_Y$ D-term contribution to the RGEs has for the low-energy spectrum. Notice to neglect this contribution the condition $f'_1(t) \mid c^2 Y_\phi \mid \ll c^2 \phi$ must be fulfilled. But this is by no means the most general situation since $|c^2|$ can be sufficiently large as to compensate the small value of $f'_1$ (see eq.(11)). It is very easy to build examples of this type. For instance $m_{H_1}(0) = m_{H_2}(0) = m$ and $m_{L_L}(0) = m_{E_R}(0) = m_{U_R}(0) = \frac{1}{2} m_{D_R}(0) = \frac{1}{2} m_{Q_L}(0) = m$ for the three generations of particles give rise to $c^2 = 18$ and therefore $f'_1(t) \mid c^2 Y_\phi \mid \sim c^2 \phi$ for $\phi = L_L, E_R, U_R$. In order to see the numerical importance that the $U(1)_Y$ D-term contribution can have in the computation of low-energy scalar masses, let us consider the previous example for the $E_R$-type sleptons and $U_R$-type squarks taking as input soft parameters $m = 150$ GeV and $M_i(0) = 50$ GeV $i = 1, 2, 3$. Then, for $c^2 = 18$ and taking for simplicity a vanishing scalar potential D–term (i.e. $\tan\beta = 1$) we obtain the following masses: $m_{E_R}(M_Z) = 32$ GeV and $m_{U_R}(M_Z) = 234$ GeV. However, the universal case $c^2 = 0$ implies $m_{E_R}(M_Z) = 151$ GeV and $m_{U_R}(M_Z) = 201$ GeV. This particular example shows that the value of $c^2$, which determines the departure from universality of soft scalar masses, can be crucial to determine the correct supersymmetric spectrum. The larger is $|c^2|$ the more important will be the $U(1)_Y$ D–term.

A more complete analysis valid for any value of the soft parameters $m, M$ is shown in figs.1 and 2 for the right selectron and the right U-squark respectively. We take for simplicity $m_{E_R}(0) = m_{U_R}(0) = m$ and $M_i(0) = M_i = 1, 2, 3$. It is straightforward to extend this study for the other particles and non–universal gaugino masses using eq.(13). The solid (dotted) lines in the figures show the contours in the $m, M$ plane corresponding to different physical masses\(^4\) in the limit $\tan\beta = 1$ ($\tan\beta = \infty$). For squark masses both types of lines are almost indistinguishable since the contribution of the last term of eq.(13) is practically negligible compared to the other terms. Slepton masses are more sensitive to this contribution as can be seen in fig.1. Different possible values of $c^2$ are considered

\(^4\)We omit to plot the contours for negative values of $M$ since the figures are symmetrical as we will explicitly show below.
in order to study the modifications that they introduce in the spectrum in comparison with the universal case \( c^2 = 0 \). The curves in the figures are conics since eq.(13) can be written as

\[
\frac{M^2}{a^2} + \frac{m^2}{b^2} = 1
\]

with

\[
a^2 \equiv \frac{m_{\phi}^2(t) - M_Z^2(T_{3L}^\phi \cos^2 \theta_W - Y_\phi \sin^2 \theta_W) \cos 2\beta}{\sum_{i=1}^{3} f_i^\phi(t)}
\]

\[
b^2 \equiv \frac{m_{\phi}^2(t) - M_Z^2(T_{3L}^\phi \cos^2 \theta_W - Y_\phi \sin^2 \theta_W) \cos 2\beta}{c_\phi^2 - c^2 Y_\phi f_1^\phi(t)}
\]

For fixed values of the low–energy scalar masses \( a^2 \) is a constant whereas \( b^2 \) is a function of \( c^2 \) (as can be seen in the figures). When the condition \( c_\phi^2 - c^2 Y_\phi f_1^\phi(t) > 0 \) \( (< 0) \) is fulfilled the curves are ellipses (hyperbolas).

As mentioned above, it is obvious from the figures that the value of \( c^2 \) is very important in order to determine the correct supersymmetric spectrum at low–energies. A negative value of \( c^2 \) will increase the mass of the right selectron with respect to the universal case \( c^2 = 0 \). On the contrary, a positive value of \( c^2 \) will decrease its mass and for values \( c^2 \geq 20 \) the ellipses are transformed in hyperbolas. In the latter case the region of soft parameters \( m, M \) excluded by the present experimental bounds ( \( m_l > 45 \text{ GeV} \) ) may be enormously increased with respect to the universal case \( c^2 = 0 \), as can be seen in fig.1. For instance a bound \( |M| > 115 \text{ GeV} \) is obtained for \( \tan\beta = 1 \). For the right \( U \)–squark the analysis will be the opposite since its hypercharge \( Y_{UR} = -\frac{2}{3} \), which appears in the first term of eq.(13), has the opposite sign to the one of the right selectron \( Y_{ER} = 1 \). In fig.2 we see that a negative (positive) value of \( c^2 \) will decrease (increase) its mass with respect to the universal case \( c^2 = 0 \). Notice that the influence of the \( U(1)_Y \) \( D \)–term, although significant, is smaller than in the right selectron case (e.g. the ellipses are transformed in hyperbolas for \( |c^2| \simeq 30 \) due to the fact that \( |Y_{UR}| < |Y_{ER}| \) and the contribution of the second term of eq.(13) is more important for the right \( U \)–squark. Finally, let us remark
that a similar analysis can be carried out for the other particles of the first and second
generation. In ref.[11] a program where measurements of three sparticle masses at future
accelerators determine, in the universal case, the values of $\tan\beta$, $m$, $M$ was proposed. It
was also remark that, if additional sparticle masses are not consistent with the obtained
values of those parameters, that could be a signal of new physics. In our approach that
would be a signal of non--universality of soft terms and the above analysis would have to
be applied.

Let us now consider the case of third--generation sparticles. The RGEs are more
complicated than in the previous case since Yukawa couplings are present. However, as it
is well known, in the interesting aproximation where all Yukawa couplings except the one
of the top $h_t$ are neglected, the formulae are simplified quite a lot and one can give some
analytic expressions for the evolution of masses. In this aproximation the above analysis
still applies for all the sleptons and squarks of the third generation except stops and left sbottom. The RGEs for the latter are:

\[
\begin{align*}
\frac{dm^2_{Q_L}}{dt} &= \left( \frac{16}{3} \frac{\alpha_3}{4\pi} M^2_3 + 3 \frac{\alpha_2}{4\pi} M^2_2 + \frac{1}{9} \frac{\alpha_1}{4\pi} M^2_1 \right) - \frac{h_t^2}{(4\pi)^2} (m^2_{Q_L} + \mu^2) - \frac{1}{6} \frac{\alpha_1}{4\pi} S \\
\frac{dm^2_{U_R}}{dt} &= \left( \frac{16}{3} \frac{\alpha_3}{4\pi} M^2_3 + \frac{16}{9} \frac{\alpha_1}{4\pi} M^2_t \right) - 2 \frac{h_t^2}{(4\pi)^2} (m^2_{Q_L} + \mu^2) + \frac{2}{3} \frac{\alpha_1}{4\pi} S \\
\frac{dm^2_{d}}{dt} &= \left( \frac{3}{4\pi} \frac{\alpha_2}{4\pi} M^2_2 + \frac{\alpha_1}{4\pi} M^2_1 \right) + \left( \frac{3}{4\pi} \frac{\alpha_2}{4\pi} + \frac{\alpha_1}{4\pi} \right) \mu^2 - 3 \frac{h_t^2}{(4\pi)^2} \mu^2 + \frac{1}{2} \frac{\alpha_1}{4\pi} S \\
\frac{dm^2_{b}}{dt} &= \left( \frac{3}{4\pi} \frac{\alpha_2}{4\pi} M^2_2 + \frac{\alpha_1}{4\pi} M^2_1 \right) + \left( \frac{3}{4\pi} \frac{\alpha_2}{4\pi} + \frac{\alpha_1}{4\pi} \right) \mu^2 - 3 \frac{h_t^2}{(4\pi)^2} (m^2_{Q_L} + \mu^2) - \frac{1}{2} \frac{\alpha_1}{4\pi} S
\end{align*}
\]

where we have also included the RGEs of the Higgs doublets $H_1$ and $H_2$ since they appear
in the squark formulae.

Eqs.(17) can be analytically solved and we obtain (taking into account the lack of
universality in the scalar masses):

\[
\begin{align*}
\end{align*}
\]
\[ m_{Q_L}^2(t) = m_{Q_L}^2(0) + M^2 \left( \frac{1}{3} e(t) + \frac{8}{3} \frac{\alpha_3(0)}{4\pi} f_3(t) + \frac{\alpha_2(0)}{4\pi} f_2(t) - \frac{1}{9} \frac{\alpha_1(0)}{4\pi} f_1(t) \right) + \frac{1}{3} AM f(t) - \frac{1}{3} k(t) A^2 - \frac{h_t^2(0) F(t)}{(4\pi)^2 D(t)} (m_{Q_L}^2(0) - \mu^2(0) + m_{Q_L}^2(0) + m_{U_R}^2(0)) - \frac{1}{6} S(0) f'_1(t) \]

\[ m_{U_R}^2(t) = m_{U_R}^2(0) + M^2 \left( \frac{2}{3} e(t) - H_s(t) + \frac{2}{3} AM f(t) - \frac{2}{3} k(t) A^2 - \frac{2}{3} h_t^2(0) F(t)}{(4\pi)^2 D(t)} (m_{Q_L}^2(0) - \mu^2(0) + m_{Q_L}^2(0) + m_{Q_L}^2(0)) + \frac{2}{3} S(0) f'_1(t) \]

\[ m_1^2(t) = \mu^2(0) l(t) + (m_1^2(0) - \mu^2(0)) + M^2 g(t) + \frac{1}{2} S(0) f'_1(t) \]

\[ m_2^2(t) = \mu^2(0) l(t) + M^2 e(t) + AM f(t) - k(t) A^2 + (m_2^2(0) - \mu^2(0)) - \frac{3}{3} \frac{h_t^2(0) F(t)}{(4\pi)^2 D(t)} (m_{Q_L}^2(0) - \mu^2(0) + m_{Q_L}^2(0) + m_{Q_L}^2(0)) - \frac{1}{2} S(0) f'_1(t) \]

where \( f'_1 \) was defined in eq. (10) and the functions \( l, e, f, k, H_s, f_i, F, D \) are defined in eq. (17) and appendix B of ref. [7]. For the sake of simplicity we have taken \( M_i(0) = M \quad i = 1, 2, 3 \) in the previous equations but it is straightforward to extend this study for non–universal gaugino masses. One has to add to the squarks in eq. (18) the scalar potential D–term contribution as in eq. (11) and the \((mass)^2\) of the corresponding fermionic partner.

The analysis of the \( U(1) \_D \)–term contribution to the low–energy stop and left sbottom masses is more complicated than for the rest of particles since new parameters appear explicitly in the formulae, in particular, the top–quark mass \( m_t \), the \( \mu \) parameter and the trilinear soft term \( A_t(0) \equiv A \). One can eliminate one of them (e.g. \( \mu \), which is the one we know the least) in terms of the others by imposing appropriate symmetry breaking at the weak scale. We are not going to review here this well known technology which may be found e.g. in refs. [7,8] (see [1] for reviews). Let us just remark that with this constraint the vacuum expectation values of the Higgs bosons will be fixed and therefore the scalar potential D–term contribution will be as well. The value of the top mass is quite constrained by the LEP and CDF data so that \( m_t \) is not a source of big uncertainty. Now,

\[ ^5 \text{They depend not only on } t \text{ but also on } \alpha_i(0) \text{ and } h_t^2(0) \]
the appearance of the $U(1)_Y$ D–term in the above formulae is important not only because contributes to the physical masses but also because modifies the electroweak symmetry breaking conditions. Notice that $\tan\beta$ which appears in chargino and neutralino masses and in the ratio of U– and D–quark masses depends on $m_{1,2}$. In order to see the qualitative importance of this contribution to the third generation squark masses, let us choose $A = 0$ and for the sake of simplicity we are going to further assume that $B = A - m$. In fact, the latter assumption has in general a small influence on the squark spectra. In fig.3 we have plotted as an example different lightest stop physical mass contours in the $m, M$ plane for $m_t \simeq 145$ GeV after diagonalizing the mass matrix. The diagonalization is due to the usual $t_L - t_R$ mixing terms. They also originate a shift in the fig.3 with respect to the other generations. We compare the non–universal case $c^2 = -15$ (taking e.g. the simple case $m_{Q_L}(0) = m_{U_R}(0) = m_{H_{1,2}}(0) = m$) with the universal one $c^2 = 0$. The same as the other generations, it is obvious from the figure that the value of $c^2$ is important to determine the low–energy masses. In fact, its influence is even larger than for the two first generations (notice that the physical mass contours behave as hyperbolas for $c^2$ around $-15$ instead of $-30$ as in fig.2), due to the extra contributions proportional to $m^2$ which appear in eq.(18) with respect to eq.(9). For $c^2 = 0$, the region of small M cannot be reached because the symmetry breaking conditions have no solution. The analysis for the left sbottom can be carried out in the same way. However, the influence of $c^2$ is less important than for the lightest stop since the hypercharge $Y_{h_u} = \frac{1}{6}$ is small and there is no mixing.

Finally, let us show an explicit example where soft masses are not universal, contrary to the usual assumptions in the MSSM. In string theory there is no reason to expect soft scalar masses degeneracy to hold generically
7. In fact, in orbifold constructions the

6This is the simplest case in order to compare the third generation masses with those of the first and second generation. Notice that the physical mass contours in the $m, M$ plane will also be conics for fixed values of $h_t(0)$ (see eq.(18)).

7However, it is worthy of remark that gaugino masses are degenerated at $M_{GUT}$ due to the universal contribution of $S$ to the gauge kinetic function up to small one–loop effects [3-5].
masses depend on the modular weights of the particles and therefore they show a lack of universality [3-5]. In particular, assuming vanishing cosmological constant, the following result was obtained [5]:

\[ m^2_\phi = m^2 (1 + n_\phi \cos^2 \theta) \] (19)

where the modular weights of the matter fields \( n_\phi \) are normally negative integers. For example, in the case of \( Z_N \) orbifolds the possible modular weights of matter fields are \(-1,-2,-3,-4,-5\). Fields belonging to the untwisted sector have \( n_\phi = -1 \). Fields in twisted sectors of the orbifold but without oscillators have usually modular weight -2 and those with oscillators have \( n_\phi \leq -3 \) (we direct the reader to ref.[3] for an explanation of these points). The angle parameter \( \theta \) says where the source of SUSY–breaking resides, either predominantly in the dilaton \( S \) sector (\( \sin \theta = 1 \) limit) or in the modulus \( T (\sin \theta = 0 \) limit). In particular: \( \tan \theta = \frac{\langle F_S \rangle}{\langle F_T \rangle} \), where \( F_S \) is the dilaton auxiliary field (the VEV of \( S \) gives the inverse square of the gauge coupling constant) and \( F_T \) is the modulus auxiliary field (the VEV of \( T \) parametrize the size and shape of the compactified space). Again we direct the reader to ref.[5] for going into details.

Notice that only the \( \sin \theta = 1 \) limit is universal⁸. As \( \sin \theta \) decreases, the modular weight dependence increases and the resulting soft terms are not universal [3]. In particular, fields with higher (negative) modular weight have smaller soft masses than those with smaller weight. For the masses of scalar particles with modular weights -1,-2 and -3 eq. (19) gives respectively

\[ m_{-1}^2 = m^2 \sin^2 \theta ; \quad m_{-2}^2 = m^2 (1 - 2 \cos^2 \theta) ; \quad m_{-3}^2 = m^2 (1 - 3 \cos^2 \theta) \] (20)

This was in fact the case of ref.[14] where the following modular weights for the MSSM fields were obtained (assuming flavour independence) by imposing the appropriate joining

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⁸This limit was obtained in ref.[13]. However, it is worth noticing here that in gaugino condensation, the only mechanism so far analysed capable of generating a hierarchical SUSY breakdown, such limit is not fulfilled.
of coupling constants at $M_{GUT}$:

$$n_{QL} = n_{DR} = -1 ; \quad n_{UR} = -2 ; \quad n_{LL} = n_{ER} = -3 ; \quad n_{H_1} + n_{H_2} = -4, -5$$  \hspace{1cm} (21)

This case shows that, even for large $sin\theta$, the deviation from universality can be important. Let us consider for example $sin\theta = 0.82$. Then, from eqs.(20) and (21) we obtain

$$m_{QL,DR} = 0.82 \text{ m} ; \quad m_{UR} = 0.59 \text{ m} ; \quad m_{LL,ER} = 0.13 \text{ m}$$  \hspace{1cm} (22)

Thus, the non-universality of soft masses is large and the influence of the $U(1)_Y$ D–term in the spectrum is crucial. Notice that, taking for the sake of definiteness $n_{H_2} = -1, n_{H_1} = -3, c^2 = 2.62$ whereas e.g. $c^2_{UR} = 0.34$ and $c^2_{LL,ER} = 0.02$. This produces a large effective modification of $U_R, L_L, E_R$ masses at $M_{GUT}$ (see the first term of eq.(13)).

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Figure captions

FIGURE 1: Allowed values of the soft parameters m, M for different right selectron ($E_R$) masses. The solid (dotted) lines correspond to the limit $\tan\beta = 1$ ($\tan\beta = \infty$). $c^2$ is the parameter defined in eq. (14) and represents the lack of universality of the soft scalar masses.

FIGURE 2: The same as fig. 1 but for the right U–squark ($U_R$).

FIGURE 3: Allowed values of the soft parameters m, M for different lightest stop masses.