SUSY SOFT BREAKING TERMS FROM STRING SCENARIOS

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ABSTRACT

The general SUSY soft breaking terms for a large class of phenomenologically relevant string scenarios (symmetric orbifolds) are given. They show a certain lack of universality, but not dangerous for flavor changing neutral currents. To get more quantitative results a specific SUSY breaking mechanism has to be considered, namely gaugino condensation in the hidden sector. Then, it turns out that squark and slepton masses tend to be much larger than scalar masses ($m_\phi \gtrsim 10 M_a$), which probably is a quite general fact. Experimental bounds and the requirement of a successful electroweak breaking without fine tuning impose further restrictions on the soft breaking terms. As a consequence the gluino and chargino masses should be quite close to their present experimental limits, whereas squark and slepton masses should be much higher ($\gtrsim 1 \text{ TeV}$).

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1. Introduction

In the last three years there has been substantial progress in our understanding of supersymmetry (SUSY) breaking in string theories\textsuperscript{1−7}. In particular, it has been learned that gaugino condensation effects in the hidden sector are able to break SUSY at a hierarchically small scale, at the same time as the dilaton $S$ and the moduli $T_i$ acquire reasonable vacuum expectation values (VEVs). (The VEV of $S$ is related to the value of the gauge coupling constant, while those of $T_i$ define the size and shape of the compactified space.) This has been realized in the context of symmetric orbifold constructions. The modular invariance in the target space\textsuperscript{3−5} as well as the presence of matter in the hidden sector play an important role in this process\textsuperscript{2,6,7}. On the other hand, the SUSY soft breaking terms in the observable sector are the phenomenological signature of any SUSY breaking mechanism, since they are responsible of the supersymmetric particle masses and the electroweak breaking process. Therefore, it is essential to completely calculate the soft breaking terms of these scenarios and study their phenomenological viability. The aim of this talk is to present some recent developments on this line of research\textsuperscript{8}.

In section 2 we give the general form of the soft SUSY breaking terms for symmetric orbifold constructions. Interestingly enough, they show a certain lack of universality, unlike the usual assumptions of the minimal supersymmetric standard model. In order to get more quantitative results a specific SUSY breaking mechanism has to be considered. This is done in section 3 assuming the above mentioned gaugino condensation in the hidden sector\textsuperscript{‡} (the only mechanism so far analyzed capable to generate a hierarchical SUSY breakdown in string constructions). Then the lack of universality of the soft breaking terms can be conveniently quantified. Furthermore, we observe a quite general fact: gaugino masses tend to be much smaller than scalar masses. In section 4 we investigate the phenomenological viability of the soft breaking terms. There are two types of tests that the soft breaking terms should pass. First, they have to be consistent with the experimental (lower) bounds on gaugino masses, squark masses, etc. Second, they should be small enough not to spoil the SUSY solution to the gauge hierarchy problem, guaranteeing a successful electroweak breaking. The latter is a naturalness requirement. We will show that the stringy scenarios considered are consistent with both tests, and that this imposes very interesting constraints on the parameters defining the SUSY breakdown. Finally, in section 5 we present our conclusions.

2. General characteristics

Following the standard notation we define the soft breaking terms in the observable sector by

$$\mathcal{L} = \mathcal{L}_{\text{SUSY}} + \mathcal{L}_{\text{soft}}.$$  \hspace{1cm} (1)

Here $\mathcal{L}_{\text{SUSY}}$ is the supersymmetric Lagrangian derived from the observable superpo-

\footnotesize
\textsuperscript{‡}A recent different analysis can be found in refs.\textsuperscript{[9,10]}
potential $W_{\text{obs}}$, which includes the usual Yukawa terms $W_Y$ and a mass coupling $\mu H_1 H_2$ between the two Higgs doublets $H_1$, $H_2$. Assuming canonically normalized fields, $L_{\text{soft}}$ is given by

$$
L_{\text{soft}} = - \sum_{\alpha} m_{\alpha}^2 |\phi_\alpha|^2 - \frac{1}{2} \sum_{a=1}^{3} M_a \bar{\lambda}_a \lambda_a - (Am_{3/2} W_Y + Bm_{3/2} \mu H_1 H_2 + \text{h.c.})
$$

(2)

where $m_{3/2}$ is the gravitino mass, $\phi_\alpha$ represent the scalar components of the supersymmetric particles, and $\lambda_a$ are the $U(1)_Y$, $SU(2)$, $SU(3)$ gauginos.

The characteristics of soft SUSY breaking terms in the observable sector are, to a great extent, determined by the type of $N = 1$ SUGRA theory which appears in four dimensions. As it is known, this is characterized by the gauge kinetic function $f$, the Kähler potential $K$ and the superpotential $W$. It is also customary to define $G = K + \log |W|^2$. These functions are determined, in principle, in a given compactification scheme, although in practice they are sufficiently well known only for orbifold compactification schemes, which on the other hand have proved to possess very attractive features from the phenomenological point of view.

More precisely, in the general case when the gauge group contains several factors $G = \prod_a G_a$, the exact gauge kinetic functions in string perturbation theory, up to small field–independent contributions, are\textsuperscript{12,13}

$$
f^a = k^a S + \frac{1}{4\pi^2} \sum_{i=1}^{3} \left( \frac{1}{2} b_i^a - k^a \delta_i^{GS} \right) \log(\eta(T_i))^2,
$$

(3)

with

$$
b_i^a = C(G^a) - \sum_{\Phi} T(R_{\Phi}^a)(1 + 2n_{\Phi}^a),
$$

(4)

where the meaning of the various quantities appearing in Eqs.\textsuperscript{(3,4)} is the following: $k^a$ is the Kac–Moody level of the $G^a$ group ($k^a = 1$ is a very common possibility), $S$ is the dilaton field, $T_i$ ($i = 1, 2, 3$) are untwisted moduli whose real parts give the radii of the three compact complex dimensions of the orbifold (Re $T_i \propto R_i^2, i = 1, 2, 3$), $\delta_i^{GS}$ are 1–loop contributions coming from the Green–Schwarz mechanism, which have been determined for the simplest (2,2) $Z_N$ orbifolds\textsuperscript{13}; $\eta(T_i)$ is the Dedekind function, $\Phi$ labels the matter fields transforming as $R_{\Phi}^a$ representations under $G^a$, and $n_{\Phi}^a$ are the corresponding modular weights. $C(G)$ denotes the Casimir operator in the adjoint representation of $G$ and $T(R)$ is defined by $\text{Tr}(T^i T^j) = T(R) \delta^{ij}$. The Kähler potential $K$ is given by\textsuperscript{14,13}

$$
K = - \log Y - \sum_{i=1}^{3} \log(T_i + \bar{T}_i) + \sum_{\alpha} \prod_i (T_i + \bar{T}_i)^{n_{\alpha}^i} |\phi_\alpha|^2 + O(|\phi_\alpha|^4),
$$

(5)

where $\phi_\alpha$ are the matter fields and

$$
Y = S + \bar{S} + \frac{1}{4\pi^2} \sum_{i=1}^{3} \delta_i^{GS} \log(T_i + \bar{T}_i),
$$

(6)

2
which can be considered\(^{13}\) as the redefined gauge coupling constant (up to threshold corrections) at the unifying string scale \((Y = 2g_{#text{str}}^2)\). Finally, the perturbative superpotential \(W^p\) at the renormalizable level has the form

\[
W^p = h_{IJK} \Phi_I^a \Phi_J^b \Phi_K^c + h'_IJK(T_i) \Phi_I^a \Phi_J^b \Phi_K^c + h''_{IJK} \Phi_I^a \Phi_J^b \Phi_K^c ,
\]

where \(\Phi_I^a\) (\(\Phi_I^b\)) are untwisted (twisted) charged matter fields. The value of \(h_{IJK}\), \(h'_IJK\) for the allowed couplings is simply a constant, while \(h'_IJK(T_i)\) is complicated but known\(^{15}\) functions of \(T_i\). Besides \(W^p\), there is a non-perturbative piece, \(W^{np}\), usually triggered by gaugino condensation effects in the hidden sector, which is crucial to break SUSY. In the following we will assume \(\langle W^p \rangle = 0\), \(\langle W^{np} \rangle \neq 0\), as it happens in all SUSY breaking scenarios so far analysed. \(W^{np}\) depends on the \(S\) and \(T\) fields and, sometimes, on certain matter fields \(A\), which are singlet under the relevant gauge groups\(^{6,7}\). As has been shown\(^7\), the condition \(\partial W / \partial A = 0\) is the correct one to integrate out these fields. In consequence, we can use

\[
W^{np} = W^{np}(S, T_i)
\]

without any loss of generality. Expressions \([3, 5]\) are to be understood at the string scale\(^{12}\) \(M_{#text{str}} = 0.527 \times g \times 10^{18}\) GeV, where \(g \approx 1 / \sqrt{2}\) is the corresponding value of the gauge coupling constant.

The next step is to calculate, once SUSY breaking is assumed, the form of the soft breaking terms based on Eqs.\([3, 5]\). This can be done by comparing the form of the corresponding \(L_{SU(2)R}A\)\(^{16}\) with \([4, 6]\). We give below the general expressions for all the soft breaking terms\(^8\). (Previous work on the subject can be found in refs.\([3, 17, 5]\).) These are given in the usual overall modulus simplification \(T = T_1 = T_2 = T_3\), the subindex \(\phi\) (\(\phi = S, T\)) denotes \(\partial / \partial \phi\) and \(\delta^GS = \sum_i \delta^GS_i\). Gaugino and scalar masses are obtained for the canonically normalized fields.

**Gaugino masses:**

\[
M_{a}(Q) = 2\pi \alpha_o(Q) m_{3/2} \left\{ \kappa^2 Y^2 \left( -\frac{1}{Y} + \frac{\bar{W}_S}{W} \right) + \frac{Y(T + T)^2}{3Y + \delta^GS_T} \frac{\bar{G}_2(T)}{\pi} \times \left( \sum_i \frac{h_i^a}{16\pi^2} - k^a \delta^GS W_S \right) \left[ \frac{1}{T + T} \left( 3 + \frac{\delta^GS W_S}{4\pi^2 W} - \frac{\bar{W}_S}{W} \right) \right] \right\} ,
\]

where \(m_{3/2} = e^{K/2} |W|\), \(Q\) is the renormalization group scale and \(\bar{G}_2(T) = - \left(\frac{T - T}{T + T} + 4\pi \eta^{-1} \frac{\partial}{\partial T}\right)\).

**Scalar masses:**

\[
m_{\phi^2} = V_o + m_{3/2}^2 \left[ 1 + \frac{n_o Y^2}{(3Y + \delta^GS_T)^2} \left( \frac{\delta^GS W_S}{4\pi^2 W} - \frac{\bar{W}_S}{W} \right) \right]^2 ,
\]

where \(V_o = \langle V \rangle\) is the value of the cosmological constant, which we will assume to vanish through the paper.

**Trilinear scalar terms:**
A term $h\Phi_1\Phi_2\Phi_3$ in the superpotential induces a trilinear scalar term of the form

\[
\mathcal{L}_{\text{tril}} = -m_{3/2} A \hat{h}\phi_1\phi_2\phi_3 + \text{h.c.}
\]

\[
= -m_{3/2} \left\{ \left( 1 - Y \frac{\bar{W}_S}{W} \right) + \left( 3 + \delta^{GS} \frac{\bar{W}_S}{4\pi^2} - (T + \bar{T}) \frac{\bar{W}_T}{W} \right) \right\} \hat{h}\phi_1\phi_2\phi_3 + \text{h.c.} ,
\]

(11)

where $\phi_\alpha$ are the (properly normalized) scalar components of the respective superfields, $n_\alpha$ are the corresponding modular weights, and $\hat{h} = e^{K/2} \prod_{\alpha=1}^{3} (K_{\alpha\bar{\alpha}})^{-1/2}$ is the effective Yukawa coupling between the physical fields. Notice that for the untwisted case $[n_\alpha = -1, h \neq h(T)]$ the previous expression is drastically simplified.

**Bilinear scalar terms:**

It is not clear by now where a bilinear term of the form $\mu H_1 H_2$ in $W_{\text{obs}}$ has its origin, although it is well known that this term is necessary in order to break $SU(2) \times U(1)_Y$ successfully. This is the so–called $\mu$ problem\footnote{Possible solutions to the $\mu$ problem can be found in refs.[18,9].}. Assuming that this term is actually present, the corresponding bilinear term in the scalar Lagrangian turns out to be

\[
\mathcal{L}_{\text{bil}} = -m_{3/2} B \hat{\mu} H_1 H_2 + \text{h.c.}
\]

\[
= -m_{3/2} Y \left\{ \frac{\bar{W}_S}{W} + \left( 3 + \frac{\bar{W}_S \delta^{GS}}{4\pi^2} - \frac{\bar{W}_T}{W} (T + \bar{T}) \right) \right\} \hat{\mu} H_1 H_2 + \text{h.c.}
\]

(12)

with a notation similar to that of Eq.(11). Again, $\hat{\mu} = e^{K/2} \prod_{\alpha=1}^{2} (K_{H_\alpha H_{\bar{\alpha}}})^{-1/2} \mu$ is the effective parameter giving, for example, the coupling between the two higgsinos. For simplicity in the notation, we will drop the hat from here on. Because of the above mentioned ignorance about the origin of $\mu$, we prefer in the following to consider $B$ as an unknown parameter.

The main conclusion at this stage is that some of the common assumptions of the minimal supersymmetric standard model, in particular universality for all the gaugino masses, scalar masses and trilinear terms, do not hold in general (this was already noted in ref.[17]). For $m_{3/2}$ this lack of universality is related to the different values of the modular weights ($n_\alpha$). Since large classes of fields (e.g. for the $Z_3$ orbifold all the untwisted or all the twisted fields) share the same modular weights, this is not necessarily dangerous for flavor changing neutral currents. Nevertheless, in order to get more quantitative results we need to know $m_{3/2}$, $\langle W \rangle$, $\langle W_S \rangle$, etc. This can only be done in the framework of a SUSY breaking scenario. We turn to this

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3. Realization of SUSY breaking by gaugino condensation

As was mentioned above, the only mechanism so far analysed, capable of generating a hierarchical SUSY breakdown in superstring theories, is gaugino condensation in the hidden sector. An extensive study of the status and properties of this mechanism can be found in ref.[7]. Let us briefly summarize here the main characteristics. For a hidden sector gauge group $G = \prod_b G_b$, gaugino condensation induces a non-perturbative superpotential

$$W^{np} = \sum_b d_b \left[ e^{-3k_b S/2\beta_b} \right] \left[ \eta(T) \right]^{6-(3k_b\delta^{GS}/4\pi^2\beta_b)} ,$$

where $\beta_b$ are the corresponding beta functions and $d_b$ are constants. Notice that $W^{np}_T = -2\eta^{-1} \eta(T) \left( 3W^{np} + \delta^{GS} W^S_{np} \right)$, which can be used to eliminate $W_T$ in (13). There is a large class of scenarios for which SUSY is spontaneously broken (essentially through a non-vanishing $T$ F-term), yielding reasonable values of $m_{3/2}$ and $\langle Y \rangle$ without the need of fine tuning\(^7\). A common ingredient of these models is the existence of more than one condensing group in the hidden sector. The values of $m_{3/2}$ and $\langle Y \rangle$ turn out to depend almost exclusively on which the gauge group and matter content of the hidden sector are. Practically, any values of $m_{3/2}$ and $\langle Y \rangle$ are available by appropriately choosing the hidden sector. Consequently, we can consider $m_{3/2}$ and $Y$ as free parameters since no dynamical mechanism is already known to select a particular string vacuum. On the other hand, both $m_{3/2}$ and $\langle Y \rangle$ are almost independent of the value of $\delta^{GS}$. Likewise, the value of $\langle T \rangle$ turns out to be $\sim 1.23$ in all cases\(^3\,\,7\). This stability does not hold, however, for the other quantities ($\langle W \rangle$, $\langle W_S \rangle$, etc.) appearing in Eqs.(13–14). For example, as was pointed out in ref.[3], the combination $YW_S - W$ (which is proportional to the $S$ F-term and appears in several places in the previous equations) is vanishing in the minimum of the scalar potential for $\delta^{GS} = 0$, but this is no longer true for $\delta^{GS} \neq 0$. In this case, however, the scalar potential is much more involved, so a numerical analysis is in general necessary. Fortunately, the results admit quite simple and useful parameterizations describing them very well (within 1% of accuracy). Next, we give these parameterizations for the common $k^a = 1$ case.

**Scalar masses:**

$$m_{\phi^a}^2 = m_{3/2}^2 \left[ 1 + n_a (0.078 - 1.3 \times 10^{-4} \delta^{GS}) \right] .$$

**Gaugino masses:**

For the $Z_3$ and $Z_7$ orbifolds the threshold contributions to the $f$ function [see Eq.(3)] are known to vanish\(^1\,\,12\) and the following equality holds for all the gauge group factors

$$\sum_i b_i^a - 2k^a \delta^{GS} = 0 .$$
As a consequence, there is a cancellation of the second term of Eq.(9). The value of the first term at the minimum of the potential is easily parameterizable as

$$M_a(Q) = -\alpha_a(Q) \ m_{3/2} \ (0.0120 \ \delta^{GS} + 0.019) \ .$$ \hspace{1cm} (16)

Hence, for the $Z_3$ and $Z_7$ orbifolds the value of $M_a$ is completely given in terms of $m_{3/2}$ and $\delta^{GS}$. For the rest of the $Z_N$ orbifolds things are different since Eq.(15) is no longer true, allowing also for a lack of universality of the gaugino masses. Then, an equation similar to Eq.(16) can be written in terms of $m_{3/2}$, $\delta^{GS}$ and $\sum b_i^a$. Let us note, however, that for all the remaining $Z_N$ orbifolds (except for the $Z_6$–II) the cancellation (15) still takes place for two of the three compactified dimensions, i.e. $b_i^a - 2k^a \delta_i^{GS} = 0 \ (i = 1, 2)$. Roughly speaking, our numerical results indicate that parameterization (16) is still valid in these cases within a 30% error. As will be seen in the next section, this is enough for our purposes.

**Trilinear and bilinear scalar terms:**

If the three fields under consideration are untwisted, i.e. $n_1 = n_2 = n_3 = -1$, $A$ turns out to be very small in all the cases: $A^{untw} \lesssim 10^{-3} \ll 1$ (if $\delta^{GS} = 0$, $A^{untw} = 0$ exactly). For twisted fields a good parameterization is

$$A = A^{untw} + (0.28 - 2.3 \times 10^{-4} \delta^{GS}) \left(3 + \sum_{\alpha=1}^{3} n_\alpha \right) + (0.69 - 6.9 \times 10^{-4} \delta^{GS}) \frac{h_T}{h} \ ,$$ \hspace{1cm} (17)

where the precise value of $h_T/h$ depends on the specific Yukawa coupling considered (see ref.[15]). In general $h_T/h$ is negative and $|h_T/h| \lesssim O(1)$. Finally, as explained above, we choose to leave the value of the bilinear term coefficient, $B$, free owing to our ignorance of the origin of $\mu$ [see Eq.(12)].

Equations (14,16,17) summarize the input of soft breaking terms to be studied from the phenomenological point of view. Before entering in a more detailed analysis, let us point out their most outstanding characteristics. First the scalar masses are not universal (for fields with different modular weights $n_\alpha$), but the deviation from universality is very small (since $n_\alpha = O(1)$), essentially consistent with flavor changing neutral currents$^{20}$. Second, gaugino masses tend to be much smaller than scalar masses. In fact, for $\delta^{GS} = 0$ the first term of Eq.(9) (proportional to the $S$ F-term) vanishes at the minimum of the potential$^3$. Then, the resulting gaugino mass is very tiny, as is reflected in Eq.(16), since it is generated thanks to the one-loop threshold correction of Eq.(3). Actually, $\delta^{GS} \neq 0$ is itself a one-loop effect, so it is not surprising that $M_a \ll m_\phi$. Since typically$^{13} \delta^{GS} \lesssim 50$, this means $M_a \lesssim \frac{1}{10} m_\phi$. This is probably a quite general fact$^{14}$ for a wide class of non-perturbative

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$^\dagger$ For the $Z_6$–II case the cancellation holds in one complex dimension. On the other hand, the $Z_6$–II orbifold is one of the less interesting orbifolds from the phenomenological point of view$^{18}$.

$^\parallel$ Work in progress.
4. Phenomenological viability of the soft breaking terms

4.1. Experimental constraints

As was mentioned in the introduction, there are two types of constraints on soft breaking terms: observational bounds and naturalness bounds. The first ones mainly come from direct production of supersymmetric particles in accelerators. This gives the following lower bounds on supersymmetric particle masses, as reported by the Particle Data Group\textsuperscript{21}

\begin{align*}
M_3 &> 79 \text{ GeV} , \quad M_{\chi^\pm} > 45 \text{ GeV} \\
m_{\tilde{q}} &> 74 \text{ GeV} , \quad m_{\tilde{t}} \gtrsim 45 \text{ GeV} ,
\end{align*}

(18)

where $M_3$ is the gluino mass, $\chi^\pm$ is the lightest chargino, and $\tilde{q}, \tilde{t}$ collectively denote squarks and sleptons respectively.\textsuperscript{**} Since in the stringy schemes we are analyzing gaugino masses are much smaller than scalar masses, the relevant experimental (lower) bounds for us are those corresponding to $M_3$ and $M_{\chi^\pm}$. In particular, for the gluino mass [(see Eq.(16)]

\begin{align*}
M_3 &\simeq \alpha_3(M_Z)m_{3/2} \left(0.0120 \delta^{GS} + 0.019\right) > 79 \text{ GeV} .
\end{align*}

(19)

This translates into lower bounds on $m_{3/2}$ and (using Eq.(14)) $m_{\tilde{q}}, m_{\tilde{t}}$:

\begin{align*}
m_{3/2} &\gtrsim \frac{79 \text{ GeV}}{\alpha_3(M_Z) \left(0.0120 \delta^{GS} + 0.019\right)} \\
m_{\tilde{q}}(M_{Str}), m_{\tilde{t}}(M_{Str}) &\gtrsim (79 \text{ GeV}) \frac{\sqrt{1 + n_\alpha(0.078 - 1.3 \times 10^{-4} \delta^{GS})}}{\alpha_3(M_Z)(0.0120 \delta^{GS} + 0.019)},
\end{align*}

(20)

Similar bounds, but involving also $\mu$ and $\sin 2\beta$, are obtained from the lightest chargino mass. However, these also involve the values of $\mu$ (whose origin is not clear yet) and $\sin 2\beta$ (where $\tan \beta = \langle H_2 \rangle / \langle H_1 \rangle$) and, hence, they are not as clean as those of Eq.(20).

4.2. Naturalness constraints

One of the most interesting features of low-energy supergravity is that the electroweak breaking can arise as a direct consequence of SUSY breaking\textsuperscript{23}. More precisely, the corresponding scalar potential for the Higgs fields is given by

\begin{align*}
V(H_1, H_2) = \frac{1}{8}(g^2 + g'^2) \left(|H_1|^2 - |H_2|^2\right)^2 + \mu_1^2|H_1|^2 + \mu_2^2|H_2|^2 - \mu_3^2(H_1H_2 + \text{h.c.}) ,
\end{align*}

(21)

where $\mu_1, \mu_2, \mu_3$ are related to the soft breaking parameters via well–known\textsuperscript{23} expressions. The minimization of $V(H_1, H_2)$ yields $\langle H_1 \rangle, \langle H_2 \rangle$ and thus $M_Z(a_i, h_t)$, where $a_i$ are the theoretical parameters defining the SUSY breaking and $h_t$ is the top Yukawa

\textsuperscript{**}Less conservative bounds have been reported elsewhere\textsuperscript{22}. The corresponding modification of our final results is completely straightforward.
coupling, which is taken as an independent parameter. For this process to be natural one must impose the absence of fine tuning on the values of the \( a_i \) parameters. This is usually accomplished through the standard criterion

\[
\left| \frac{a_i \cdot \partial M_Z^2(a_i, h_t)}{M_Z^2} \right| < \Delta, \tag{22}
\]

where \( \Delta \) measures the allowed degree of fine tuning\(^{††} \). In ref.[24], in the context of the minimal supersymmetric standard model (MSSM), the soft breaking parameters \( M, m, \mu, A, B \) were taken as the \( a_i \) parameters of Eq.(22). In fact, in the MSSM, the former are quantities whose origin is unknown rather than independent parameters. On the contrary, in our stringy context assuming gaugino condensation, these parameters (except \( \mu \) and \( B \)) are related each other via Eqs.(14–17). Consequently, in our case the role of the independent \( a_i \) parameters in (22) has to be played by \( m_{3/2}, \mu \) and \( B \). After the corresponding analysis (for details see ref.[8]) we arrive from (22) to upper bounds on the the values of \( m_{3/2}, \mu \) and \( B \), and thus on all the supersymmetric masses. Of course these bounds depend on the values of \( \Delta \) and \( \delta^{GS} \).

(Here we take conservatively\(^{25} \Delta = 50 \). On the other hand, typically

\[ \delta^{GS} < \sim 50. \]

Besides these upper bounds there are the experimental lower bounds obtained in the previous subsection. This gives the allowed ranges of variation for all the masses. For example, for a typical case (\( \delta^{GS} = 45 \)) we find

\[
\begin{align*}
1100 \text{ GeV} \leq m_{3/2} &\leq 4200 \text{ GeV}, & 350 \text{ GeV} \leq \mu \leq 450 \text{ GeV} \\
79 \text{ GeV} \leq M_3 &\leq 285 \text{ GeV}, & 45 \text{ GeV} \leq M_{\chi^\pm} \leq 80 \text{ GeV} \\
110 \text{ GeV} \leq m_{\text{top}} &\leq 165 \text{ GeV}, & 1070 \text{ GeV} \leq m_{\tilde{q}}, m_{\tilde{\ell}} \leq 4080 \text{ GeV},
\end{align*}
\]

(remarkably, the above upper bound on \( m_{\text{top}} \) arises naturally, without using the current theoretical upper bound). As a general fact we can say that the gluino and chargino masses should be quite close to their present experimental limits, whereas \( m_{\tilde{q}}, m_{\tilde{\ell}} \) should be much higher (\( \sim 1 \text{ TeV} \)).

5. Conclusions

The general form of the SUSY soft breaking terms is derivable for a large class of phenomenologically interesting string constructions, more precisely symmetric orbifolds. As a general characteristic the soft breaking terms show a certain lack of universality (unlike the usual assumptions of the minimal supersymmetric standard model), though this is not necessarily dangerous for flavor changing neutral currents. To get more quantitative results a specific SUSY breaking mechanism has to be considered, gaugino condensation being the most attractive on so far proposed. The most striking property of the soft breaking terms derived from gaugino condensation is that squark and slepton masses tend to be much larger than gaugino masses \( (m_\phi > \sim 10M_a) \), which probably is a quite general fact. Experimental bounds and the requirement of a successful electroweak breaking without fine tuning impose further

\(^{††} \) A discussion on the validity of this criterion can be found in ref.[25]
restrictions on the soft breaking terms. As a consequence the gluino and chargino masses should be quite close to their present experimental limits, whereas squark and slepton masses should be much higher (\sim 1 \text{ TeV}).

7. References

1. L. Dixon, talk presented at the A.P.S. D.P.F. Meeting at Houston (1990)


20. See e.g. J.S. Hagelin, S. Kelley and T. Tanaka, MIU–THP–92/59, and M. Dine talk in this Workshop


22. See e.g. T. Medcalf talk in this Workshop

23. For a recent review see: L.E. Ibáñez and G.G. Ross, CERN–TH.6412/92 (1992), to appear in Perspectives in Higgs Physics, ed. G. Kane, and references therein
