Some Phenomenological
Implications of String Loop Effects

JAN LOUIS
Sektion Physik, Universität München
Theresienstr. 37, D-80333 München, Germany

and

YOSEF NIR
Department of Particle Physics
Weizmann Institute of Science, Rehovot 76100, Israel

Abstract

We investigate low energy implications of string loop corrections to supergravity couplings which break a possible flavor universality of the tree level. If Supersymmetry is broken by the dilaton $F$-term, universal soft scalar masses arise at the leading order but string loop corrections generically induce flavor-non-diagonal soft terms. Constraints from flavor changing neutral currents (FCNC) and CP violation then require a large supersymmetry breaking scale and thus heavy gluinos and squarks. If Supersymmetry is broken by moduli $F$-terms, universality at the string tree level can only be guaranteed by extra conditions on the Kähler potential. A large hierarchy between the gluino and squark masses ensures that FCNC and CP violation constraints are satisfied. If the soft scalar masses vanish at the string tree level, the cosmological problems related to light moduli can be evaded. However, generic string loop corrections violate FCNC bounds and require very heavy squark masses ($\approx 100\, \text{TeV}$).

1. Introduction

Supersymmetric extensions of the Standard Model (SM) are among the most promising candidates for new physics above the weak scale $M_Z$. In the minimal supersymmetric SM (MSSM) all particles of the SM are promoted to chiral $N = 1$ supermultiplets with one additional Higgs doublet added [1]. Supersymmetry is assumed to be broken by explicit but soft breaking terms which appear naturally in the low energy limit of spontaneously broken supergravity theories. This soft supersymmetry breaking introduces a number of new parameters into the Lagrangian which control the mass spectrum of the fermion’s superpartners. The parameter space spanned by the soft terms has been the subject of numerous investigations, mostly under some simplifying assumptions [1]. It has also been studied within the context of further extensions such as superstring theory [2–8].

Without specifying the precise origin of the non-perturbative effects in superstring theory but with a number of plausible assumptions, it has been possible to observe interesting features of the soft terms [4, 7, 8]. In particular, many signatures at $M_Z$ are entirely controlled by perturbative couplings in string theory and (almost) independent of the assumption about the unknown non-perturbative properties. The relevant perturbative couplings have indeed been computed for many string vacua at the leading order (tree level) in string perturbation theory. One finds that in most cases the standard assumptions of the MSSM are not fulfilled. For example, non-universal scalar masses as well as non-proportional $A$-terms with arbitrary CP-violating phases are easily generated. Non-universal scalar masses are particularly dangerous since they generically induce unacceptably large contributions to rare processes such as flavor changing neutral currents (FCNC) [9–12]. One way out is to look for possible mechanisms which naturally suppress non-universal scalar masses. This could be natural if SUSY breaking is communicated to the light particles by gauge interactions [13] or in models with a nonabelian horizontal symmetry [14]. (Abelian horizontal symmetries could
align quark and squark mass matrices, thus suppressing FCNC without squark degeneracy \[15,16].\)

In the context of string theory, universality of squark masses is achieved if the dominant source for supersymmetry breaking is the dilaton $F$-term \[7\] or in ‘no-scale’ type of models \[19\]. However, in both cases non-universal scalar masses might arise through couplings generated at the 1-loop level of string perturbation theory. Although for such couplings much less (string) information is currently available, it is possible to estimate the typical size of these corrections and hence estimate the physical implications for weak scale phenomenology. Furthermore, the generic CP-violating phases of the $A$ and $B$-terms are constrained in both scenarios and can be confronted with the bounds for the electric dipole moment of the neutron (EDMN)\[20\].

This paper is organized as follows: In Section 2 we present the general form of string loop corrections. Their implications on FCNC and CP violation when supersymmetry is broken by the dilaton $F$-term are studied in section 3. A similar analysis for supersymmetry breaking by moduli $F$-terms (including some cosmological consequences) are studied in section 4. A summary is given in Section 5.

\[\star\] Other recent investigations of the problem include refs. \[17, 18\].

2. String Loop Corrections

We first summarize the generic structure of the couplings in the low energy effective Lagrangian of string theory. In addition to the gravitational and gauge multiplets, the massless spectrum contains two types of chiral supermultiplets. First, matter fields $Q^I$ which are charged under the low energy gauge group $G$ and which contain the quark and lepton multiplets of the SM. Second, there are the gauge neutral supermultiplets $S$ (dilaton) and $M^I$ (moduli) which are flat directions of the perturbative effective potential and whose VEVs parameterize the perturbative degeneracy of the string vacuum.\[\star\] The couplings of the low energy effective Lagrangian for the massless multiplets are encoded in three scalar functions: the real Kähler potential $K$, the holomorphic superpotential $W$ and the holomorphic gauge kinetic function $f$.

$K$ summarizes the kinetic energy terms and at low energies can be expanded in the matter fields

\[
K = \kappa^{-2} \bar{K} + Z_{IJ} \bar{Q}^I e^{2\nu} Q^J + \left( \frac{1}{2} H_{IJ} Q^I Q^J + \text{h.c.} \right) + \cdots ,
\]

where the ‘\cdots’ in eq. (2.1) correspond to terms which are irrelevant for the present investigation. The matter fields $Q^I$ carry canonical dimension one whereas $S$ and $M^I$ are expected to receive Planck-sized VEVs and therefore are chosen to be dimensionless. The couplings $\bar{K}$, $Z_{IJ}$ and $H_{IJ}$ are dimensionless functions of $S$ and $M^I$ and only further constrained by the fact that the dilaton Re $S$ serves as the string-loop counting parameter. At the string tree level the dilaton couples universally in all string vacua; this universality is lost at the loop level but all

\[\star\] Strictly speaking there can also be singlet supermultiplets which are not moduli, i.e. which are not a flat direction of the effective potential. For the purpose of this article we include them among the matter fields $Q^I$. 

* Other recent investigations of the problem include refs. \[17, 18\].
The details of the mechanism responsible for the vanishing of the cosmological constant. Recently various mechanisms have been studied which also include low energy quantum corrections to the cosmological constant [22,23]. Most of our analysis here is insensitive to the details of the mechanism responsible for the vanishing of the cosmological constant.

\[ \hat{K} = -\ln (S + \bar{S}) + \sum_{n=0}^{\infty} \frac{\hat{K}^{(n)}(M, \bar{M})}{[8\pi^2(S + \bar{S})]^n}, \]

\[ Z_{IJ} = \sum_{n=0}^{\infty} \frac{Z_{IJ}^{(n)}(M, \bar{M})}{[8\pi^2(S + \bar{S})]^n}, \tag{2.2} \]

\[ H_{IJ} = \sum_{n=0}^{\infty} \frac{H_{IJ}^{(n)}(M, \bar{M})}{[8\pi^2(S + \bar{S})]^n}, \]

where \( \hat{K}^{(n)}, Z_{IJ}^{(n)} \) and \( H_{IJ}^{(n)} \) do not depend on the dilaton and their moduli dependence in general cannot be further constrained (they do depend on the details of the internal superconformal field theory).\(^\dagger\)

The scalar potential and the Yukawa couplings \( \hat{Y}_{I JL} \) are determined by the superpotential \( W \) which is not renormalized at any order in string perturbation theory. The perturbative \( W \) is completely independent of the dilaton \( S \) but non-perturbative corrections can introduce further dilaton (and moduli) dependence into \( W \). Expanding in \( Q^I \) we have

\[ W = \hat{W}(S, M^i) + \frac{1}{2} \hat{\mu}_{IJ}(S, M^i)Q^I Q^J + \frac{1}{3} \hat{\kappa}_{I J L}(M^i)Q^I Q^J Q^L + \cdots. \tag{2.3} \]

where the ‘\( \cdots \)’ stand for non-renormalizable interactions. \( \hat{W} \) is identically zero at any order in string perturbation theory and arises only from non-perturbative physics. (Similarly, the \( S \)-dependence in \( \hat{\mu} \) is induced at the non-perturbative level.) Without specifying the precise nature of such non-perturbative effects they can be parameterized by \( \hat{W} \). We assume that \( \hat{W} \) is such that it breaks supersymmetry by generating non-vanishing moduli \( F \)-terms \( \langle F^i \rangle \) and/or a dilaton

\[ F \text{-term } \langle F^\phi \rangle. \] 

To simplify our notation let us introduce an index \( \phi \) which runs over both the moduli and dilaton direction, \( i.e. \phi = (i, S) \). Using this notation the \( F \)-terms are given by

\[ \mathcal{F}^{\phi} = k^2 e^{\hat{K}/2} \hat{K}^{\phi\phi}(\partial_\phi \hat{W} + \hat{W} \partial_\phi \hat{K}), \tag{2.4} \]

while the scale of supersymmetry breaking is parameterized by the (complex) gravitino mass

\[ m_{3/2} = k^2 e^{\hat{K}/2} \hat{W}. \tag{2.5} \]

We further assume that, at the minimum of the potential, a dilaton VEV \( \langle S \rangle \) and moduli VEVs \( \langle M^i \rangle \) are generated and hence the perturbative vacuum degeneracy is (partially) lifted. Finally, the cosmological constant is assumed to be zero which implies\(^\ddagger\)

\[ |m_{3/2}|^2 = \frac{1}{3} \hat{K}_{\phi\phi} F^{\phi\phi} \mathcal{F}^{\phi}. \tag{2.6} \]

Under these assumptions (spelled out in more detail in ref. [7]) soft supersymmetry breaking terms are generated in the observable sector. In particular, the potential for the observable matter scalars (which we also call \( Q^I \)) contains the following soft supersymmetry breaking terms:

\[ V^{(SSB)} = M_{Q}^{2} Q^{I} \bar{Q}^{J} + (\frac{1}{3} A_{I J L} Q^{I} Q^{J} Q^{L} + \frac{1}{2} B_{I J} Q^{I} Q^{J} + h.c.), \tag{2.7} \]

where the parameters \( m^2, A, B \) are moduli and dilaton dependent and not necessarily flavour diagonal [4, 7, 8]. Specifically,

\[ m_{IJ}^{2} = |m_{3/2}|^{2} Z_{IJ} - \hat{F}^{\phi} \mathcal{F}^{\phi} R_{\phi\phi I J}, \tag{2.8} \]

where the flavour dependence can arise through the (perturbative) curvature

\(^\dagger\) \( \hat{K}^{(1)} \) and \( Z_{IJ}^{(1)} \) are the four-dimensional analogue of the Green-Schwarz term and have recently been computed in some orbifold vacua [21].

\(^\ddagger\) Recently various mechanisms have been studied which also include low energy quantum corrections to the cosmological constant [22,23]. Most of our analysis here is insensitive to the details of the mechanism responsible for the vanishing of the cosmological constant.
The gauge couplings

\[ R_{\phi i} = \partial_\phi \tilde{\phi} Z_{I J} - \Gamma^N_{\phi i} Z_{N L} \tilde{\Gamma}^L_{\phi J}, \quad \Gamma^N_{\phi i} = Z^{N J} \partial_\phi Z_{J I}, \tag{2.9} \]

and hence the standard assumption of universal (flavour independent) soft masses might not hold. Furthermore,

\[ A_{I J L} = F^\phi (\partial_\phi Y_{I J L} - \Gamma^N_{\phi i} Y_{I J L}) + \frac{1}{2} \tilde{K}_\phi Y_{I J L}, \quad Y_{I J K} = e^{K/2} \tilde{Y}_{I J K}, \tag{2.10} \]

where the first two terms are in general not proportional to the Yukawa couplings. Similarly,

\[ B_{I J} = 2|m_3/2|^2 H_{I J} + m_3/2 F^\phi D_\phi H_{I J} - \tilde{m}_3/2 \tilde{F}^\phi \tilde{D}_\phi H_{I J} + e^{K/2} \left[ F^\phi (\partial_\phi \mu_{I J} + \tilde{K}_\phi \tilde{\mu}_{I J} - 2\tilde{K}^\phi_{I J} \tilde{\mu}_{K J}) - \tilde{m}_3/2 \tilde{\mu}_{I J} \right], \tag{2.11} \]

(where \( D_\phi H_{I J} = \partial_\phi H_{I J} - 2\Gamma^K_{I J} H_{K J} \)) is not necessarily proportional to \( \phi \)

\[ \mu_{I J} = e^{K/2} \tilde{\mu}_{I J} + m_3/2 H_{I J} - \tilde{F}^\phi \tilde{\partial}_\phi H_{I J}. \tag{2.12} \]

Hence, the parameters of \( V^{(SSB)} \) in eq. (2.7) in general do not satisfy the property of flavour-independence which is commonly assumed in the MSSM.

Finally, there is one more soft term induced: the gauginos acquire a mass given by

\[ \tilde{m}_a = F^\phi \partial_\phi \ln g_a^{-2}, \tag{2.13} \]

where \( g_a^{-2} \) are the gauge couplings (\( a \) labels the simple factors in the gauge group). In string theory the gauge couplings are universal at the leading order and determined by the VEV of the dilaton. Non-universality and moduli dependence is only introduced via one-loop threshold corrections \( \Delta_a \) [24]

\[ g_a^{-2}(M_{\text{String}}) = \Re S + \frac{\Delta_a (M, M)}{16\pi^2}. \tag{2.14} \]

\( M_{\text{String}} \) denotes the characteristic scale of string theory; numerically \( M_{\text{String}} \approx 5 \times 10^{17} \) GeV which is close to the supersymmetric GUT-scale \( M_{\text{GUT}} \approx 3 \times 10^{16} \) GeV. In this paper we do not make any distinction between the two scales and denote them both by \( M_X \). As a consequence of the very special field dependence of the gauge couplings, the gaugino masses are universal at the leading order and obey

\[ \tilde{m}_a = \tilde{m}_{1/2} + \frac{\alpha_X}{4\pi} \tilde{m}_a^{(1)} + \cdots, \tag{2.15} \]

where

\[ \tilde{m}_{1/2} = \frac{FS}{(S + \bar{S})}, \quad \tilde{m}_a^{(1)} = F^{i} \partial_i \Delta_a - F^{S} \Delta_a, \tag{2.16} \]

\[ \alpha_X = \frac{g^2(M_X)}{4\pi} = \frac{1}{2\pi(S + \bar{S})}. \]

Note that the universal gaugino mass \( \tilde{m}_{1/2} \) is directly proportional to the dilaton \( F^{S} \) and that both \( \tilde{m}_{1/2} \) and \( \tilde{m}_a^{(1)} \) are of order \( O(m_3/2) \).

On the other hand, the scalar masses given by eq. (2.8) are in general flavour-dependent (non-universal) already at the leading order of perturbation theory when \( Z_{I J} \) is approximated by its tree level contribution \( Z_{I J}^{(0)} \). However, there are scenarios where universal scalar masses and \( A \)-terms do appear at the leading order and non-universality is only introduced at the one-loop level. For those cases - which are the focus of this paper - we have

\[ m_{ij}^2 = m_0^2 Z_{ij}^{(0)} + \frac{\alpha_X}{4\pi} m_{ij}^{(1)} + \cdots, \tag{2.17} \]

\[ A_{i j} = A_0 Y_{i j} + \frac{\alpha_X}{4\pi} A_{i j}^{(1)} + \cdots. \]
3. Supersymmetry Breaking by the Dilaton

Under the assumption that only a dilaton F-term $\langle F^S \rangle$ is generated by the non-perturbative physics, the soft parameters simplify considerably at the leading order and the scalar masses and $A$-terms are indeed universal. This is a consequence of the universal couplings of the dilaton at the string tree level. Specifically one finds [7, 8]

$$m_0^2 = |m_{3/2}|^2 = \frac{1}{6} |F^S|^2 / (S + S), \quad A_0 = - \frac{F^S}{S + S}, \quad \tilde{m}_{1/2} = \frac{F^S}{S + S}, \quad (3.1)$$

while $B$ and $\mu$ are independent parameters. (If, in addition, $\tilde{\mu} = 0$ holds in eq. (2.12), $B$ and $\mu$ are related via $B = 2 \tilde{m}_{3/2} \mu$ but we do not assume this relation here.) Given the soft terms (3.1) generated at $M_X$, standard RG-analysis can be used to compute the supersymmetric mass spectrum at low energies [25].

One finds that all squark masses $m_{\tilde{q}}$ are essentially degenerate with the gluino mass $m_3$

$$m_{\tilde{q}} \simeq \tilde{m}_3 \simeq 5 m_{3/2}, \quad (3.2)$$

whereas the slepton masses obey

$$m_{\tilde{\ell}} \simeq 0.3 \tilde{m}_3 \simeq 1.5 m_{3/2}. \quad (3.3)$$

In order to evade the direct experimental bounds [27] on scalar and gaugino masses, eqs. (3.2), (3.3) imply

$$m_{3/2} > 30 \text{ GeV}. \quad (3.4)$$

In this section we study the physical properties of this scenario beyond the leading order approximation. In particular, we assume generic $\mathcal{O}(1)$ one-loop couplings $Z_{IJ}^{(1)}$ which induce flavour-dependent scalar masses and non-proportional $A$-terms at the next order. From eqs. (2.8)-(2.10) we learn (using (3.1))

$$m_{ij}^{2(1)} = -5 \frac{|m_{3/2}|^2 Z_{ij}^{(1)}}{2} \sim \mathcal{O}(m_{3/2}^2),$$

$$A_{ij}^{(1)} = \frac{F^S}{S + S} \left( -K^{(1)} Y_{IJL} + 3Z_{ij}^{(1)} Z_{JNL} + \partial_j \hat{K}^{(1)} \bar{B}_i \hat{D}_j Y_{IJL} \right) \sim \mathcal{O}(m_{3/2} Y), \quad (3.5)$$

where $Z^{(0)}$, $Z^{(1)}$ and $\hat{K}^{(1)}$ are all functions of $\mathcal{O}(1)$.

3.1. Constraints from Flavor Changing Neutral Currents

Let us first focus on the constraints implied by the smallness of FCNC. We use the notation of ref. [15] and the calculations of ref. [28]. The experimental bounds from FCNC constrain the sfermion masses at the weak scale $M_f^2$ ($f = u, d, \ell$) which are determined in terms of the soft input parameters (3.1) and (3.5) generated at the high energy scale $M_X$. In the basis where fermion mass-matrices are diagonal and gaugino couplings are diagonal, the sfermion masses appear in $3 \times 3$ submatrices,

$$M_f^2 = \begin{pmatrix} M_{LL}^2 & M_{LR}^2 \\ M_{RL}^2 & M_{RR}^2 \end{pmatrix}, \quad (3.6)$$

where the soft scalar masses $m_{ij}^2$ contribute to the diagonal blocks $M_{LL}^2, M_{RR}^2$ while the $A$-terms directly determine $M_{LR}^2 \sim 1$. As SUSY breaking by the dilaton leads to approximately degenerate sfermions in each sector, it is convenient to define the average sfermion mass-squared, $m_{\tilde{f}}^2$. FCNCs are then proportional to

$$(\delta_{MN}^f)_{ij} = \frac{(M_{MN}^2)_{ij}}{m_{\tilde{f}}^2}, \quad i \neq j \quad (3.7)$$

and the strongest constraints on non-universality arise from the light generations.

* See also ref. [26].

† See also refs. [29].

‡ The diagonal elements in $M_{LR}^2$ also depend on $\mu$. 
(For squarks the bounds are particularly strong on the combination \( \delta^f_{12} \) = \( \sqrt{(\delta^f_{12})_1^2 (\delta^f_{12})_2^2} \). One finds [28, 15]

\[
\begin{align*}
\text{Re} \left( \delta_{12}^d \right) & \leq 6 \times 10^{-3} \left( \frac{m_\tilde{d}}{1 \text{ TeV}} \right), \\
\text{Re} \left( \delta_{12}^q \right) & \leq 8 \times 10^{-3} \left( \frac{m_\tilde{d}}{1 \text{ TeV}} \right), \\
\text{Im} \left( \delta_{12}^d \right) & \leq 5 \times 10^{-4} \left( \frac{m_\tilde{d}}{1 \text{ TeV}} \right), \\
\text{Im} \left( \delta_{12}^q \right) & \leq 7 \times 10^{-4} \left( \frac{m_\tilde{d}}{1 \text{ TeV}} \right),
\end{align*}
\]

(3.8)

\[
(\delta_{12}^f)_{12} \leq 1.5 \times 10^{-2} \left( \frac{m_\tilde{f}}{0.3 \text{ TeV}} \right)^2, \\
(\delta_{12}^d)_{12} \leq 5 \times 10^{-6} \left( \frac{m_\tilde{d}}{0.3 \text{ TeV}} \right),
\]

(3.9)

from the bound on \( \text{BR}(\mu \rightarrow e\gamma) \). The bounds (3.8) and (3.9) have been evaluated under the assumption \( m_\tilde{q} \approx \tilde{m}_3 \) and \( m_\tilde{f} \approx 2 \tilde{m}_1 \) as appropriate for dilaton-induced SUSY breaking (cf. eq. (3.2)). In the slepton sector the bound on \( (\delta_{12}^f)_{12} \) also depends on \( (M^{LR}_{12})_2 \) and we have used the value \( (M^{LR}_{12})_2 = -3.8 m_\tilde{m} m_{3/2} \) in (3.9) as a characteristic value for the dilaton scenario.\(^\S\) All bounds quoted are only accurate up to factors of \( \mathcal{O}(1) \) due to hadronic uncertainties in the squark sector and the dependence on \( (M^{LR}_{12})_2 \) in the slepton sector. Finally, the bounds from \( \Delta m_B, \Delta m_D \) and radiative \( \tau \) decays are much milder than (3.8) and (3.9) and play no role in our analysis.

The experimental bounds (3.8) and (3.9) can now be compared with the theoretical ‘predictions’ of the dilaton scenario which follow from eqs. (3.1) and (3.5). Let us first note that even for the universal boundary conditions (3.1) renormalization effects induce small \( \delta \)’s at low energies which obey (3.8) and (3.9) [28, 29]. The point we want to study here is the implication of the non-universality implied by (3.5). The running of the off-diagonal mass-matrix elements of the first two generations is negligibly small [1] and hence we can estimate at the weak scale:

\[
\begin{align*}
(\delta_{MM}^f)_{12} & \approx \frac{\alpha_X m_{12}^{(1)}}{4\pi m_\tilde{q}} \approx 1.2 \times 10^{-4}, \\
(\delta_{LR}^f)_{12} & \approx 3 \frac{\alpha_X m_{3/2} m_{3/2}}{4\pi m_{\tilde{q}}^2} \approx 4 \times 10^{-7} \left( \frac{1 \text{ TeV}}{m_{\tilde{q}}} \right), \\
(\delta_{MM}^d)_{12} & \approx \frac{\alpha_X m_{12}^{(1)}}{4\pi m_\tilde{q}} \approx 1.5 \times 10^{-3}, \\
(\delta_{LR}^d)_{12} & \approx 1.5 \frac{\alpha_X m_{3/2} m_{3/2}}{4\pi m_\tilde{q}^2} \approx 1 \times 10^{-6} \left( \frac{0.3 \text{ TeV}}{m_\tilde{q}} \right),
\end{align*}
\]

(3.10)

where we used eqs. (3.2), (3.3), (3.5) and \( \alpha_X = 1/24 \). Also (3.10) are only order of magnitude estimates and factors of \( \mathcal{O}(1) \) are neglected.

Comparing (3.10) with (3.8), (3.9) we find that in the squark sector the only potentially interesting bound arises from \( \epsilon_K \). Assuming phases of \( \mathcal{O}(1) \) in the mass matrix (3.6) or equivalently \( \text{Im}(\delta^f_{12} M_{12}) \sim \text{Re}(\delta^f_{12} M_{12}) \), we find that (3.8) can be satisfied by slightly raising \( m_\tilde{q} \).

\[
m_\tilde{q} \geq 180 \text{ GeV} \quad \Rightarrow \quad \tilde{m}_3 \geq 180 \text{ GeV} \quad (m_{3/2} \geq 40 \text{ GeV}).
\]

(3.11)

In the slepton sector the constraint is stronger and (3.9) can only be satisfied for

\[
m_\tilde{f} \geq 135 \text{ GeV} \quad \Rightarrow \quad \tilde{m}_3 \geq 450 \text{ GeV} \quad (m_{3/2} \geq 90 \text{ GeV}).
\]

(3.12)

The fact that the stronger constraint arises in the slepton sector is a consequence of the large renormalization effect in the squark sector due to the gluino mass which enhances the average squark masses and therefore weakens the FCNC constraints [14, 8, 18].

To summarize, when SUSY is broken by the dilaton, universality and proportionality are violated at the string loop level. The effect on mass differences

\(^\S\) \((M^{LR}_{12})_2 \) depends on a phase \( \phi_B \), defined in the next section. Here we take \( \phi_B = 0 \) which gives the weakest constraint.
in the various neutral meson systems is small. If the phase in the universality violating terms is of $O(1)$, a lower bound on the down squark masses arises, $m_{\tilde{d}} \geq 180 \text{ GeV}$. The effects on the decay $\mu \to e\gamma$ due to violation of either universality or proportionality are more significant and give a lower bound on the charged slepton masses, $m_{\tilde{\ell}} \geq 135 \text{ GeV}$. As all sfermion masses are fixed by the gluino mass in this scenario, we conclude that the most stringent constraint is the one from the lepton sector and requires $\tilde{m}_3 \geq 450 \text{ GeV}$ (or equivalently $m_{3/2} \geq 90 \text{ GeV}$). However, it should be stressed that such estimates are only accurate up to factors of $O(1)$.

### 3.2. Constraints from CP Violation

In the previous section we investigated the effects of violation of universality and proportionality by string loop effects. In this section we study then CP violating effects that arise at the leading order and are implied by eqs. (3.1). Such effects are constrained by the upper bounds on the electric dipole moments of the neutron (EDMN) and of various atoms and molecules.

When both universality and proportionality hold, there are, in general, two new CP violating phases (in addition to the CKM phase $\delta_{KM}$ and the strong CP phase $\theta_{QCD}$) [30]:

$$\phi_A \equiv \arg \left( A_0 \tilde{m}_{1/2}^* \right), \quad \phi_B \equiv \arg \left( B_0 \tilde{m}_{1/2}^* \right),$$

(3.13)

(where $B_0 = B_{1J}/\mu_{1J}$). In eq. (2.5) we defined $m_{3/2}$ as a complex quantity, its complex conjugate is $\tilde{m}_{3/2} = \kappa^2 e^{iK/2} \bar{W}$. Using eqs. (3.1) we conclude that $\phi_A$ vanishes at tree level while there is no significant simplification for $\phi_B$ and we expect [8]

$$\phi_A = O \left( \frac{\alpha_X}{4\pi} \right), \quad \phi_B = O \left( 1 \right).$$

(3.14)

The contributions to the EDMs of the neutron and of various atoms from $\phi_A$ and $\phi_B$ were estimated in ref. [31]. The appropriate modifications of their estimates to our case read, in the limit $\phi_B \gg \phi_A$,

$$|d_N| \sim 1.4 \times 10^{-24} \text{ cm } \left( \frac{100 \text{ GeV}}{m_{3/2}} \right)^2 \left( \sin \phi_B \right)$$

(3.15)

(where the leading contributions come from the light quark EDMs and CDMs),

$$|d_{Tl}| \sim 1.6 \times 10^{-22} \text{ cm } \left( \frac{100 \text{ GeV}}{m_{3/2}} \right)^2 \left( \sin \phi_B \right)$$

(3.16)

(where the leading contribution comes from the EDM of the electron), and

$$|d_{He}| \sim 3 \times 10^{-26} \text{ cm } \left( \frac{100 \text{ GeV}}{m_{3/2}} \right)^2 \left( \sin \phi_B \right)$$

(3.17)

(where the leading contribution comes from the nonderivative nucleon-nucleon coupling). The experimental bounds [33–35],

$$|d_N| \leq 1.2 \times 10^{-25} \text{ cm},$$

$$|d_{Tl}| \leq 6.6 \times 10^{-24} \text{ cm},$$

$$|d_{He}| \leq 1.3 \times 10^{-27} \text{ cm},$$

require

$$m_{3/2} \geq 480 \text{ GeV} \sqrt{\sin \phi_B}.$$ 

(3.18)

(We have also checked the bounds from $d_{Cs}$, $d_{Xe}$ and $d_{TlF}$ and found that they are weaker.) Even for $\sin \phi_B \sim 0.1$, we need $m_{\tilde{g}} \geq 800 \text{ GeV}$ which is stronger than any of the FCNC bounds (3.11), (3.12). Finally, we note that if $\phi_B = O(\frac{\alpha_X}{4\pi})$, then (3.18) is satisfied for $m_{3/2} \geq 30 \text{ GeV}$, which coincides with the direct limit (3.4).

* Strictly speaking one should also take the renormalization of $\phi_B$ into account. However, in ref. [32] it was shown that $\phi_B$ does not renormalize and therefore we can use the boundary values at $M_X$. 
To summarize, of the two new CP violating phases, one vanishes at string tree level and poses no phenomenological problems. The other is expected, in general, to be of $\mathcal{O}(1)$, in which case it would require gluino mass above 800 GeV. Under special circumstances it could be suppressed\(^\dagger\) but we see no simple mechanism to guarantee its vanishing.

4. Supersymmetry Breaking induced by the Moduli

4.1. Non-Universal Soft Terms

If the dominant source of supersymmetry breaking are moduli $F$-terms $\langle F^i \rangle$, the soft scalar masses are generically non-universal at the string tree level and of $\mathcal{O}(m_{3/2})$. The $A$-terms are not proportional to the Yukawa couplings and of $\mathcal{O}(m_{3/2})$. The gaugino masses are also non-universal but, more importantly, they are suppressed since $F^S \approx 0$ implies $\tilde{m}_{1/2} \approx 0$ via eq. (2.16). Instead we have

\[
\tilde{m}_a = \frac{\alpha_X}{4\pi} \tilde{m}_a^{(1)},
\]

where $\tilde{m}_a^{(1)} = \mathcal{O}(m_{3/2})$. The current lower bound on the gluino mass [27] implies

\[
150 \text{ GeV} < \tilde{m}_3(M_Z) = 3 \tilde{m}_3(M_X) \approx 3 \frac{\alpha_X}{4\pi} m_{3/2},
\]

or equivalently $m_{3/2} \gtrsim 15 \text{ TeV}\(^\ddagger\) Generically, this leads to squark and slepton masses of the same order of magnitude and hence radiative electroweak symmetry breaking requires a major fine-tuning in order to keep $M_X$ at 90 GeV [38]. However, large scalar masses also suppress the contributions to FCNC processes.

\(^\dagger\) For example, when $\tilde{\mu} = 0$ and $\partial_5 \tilde{W}/\tilde{W}$ is real. Different mechanisms are proposed in refs. [36, 37].

\(^\ddagger\) A similar bound follows from the charginos. Again, this bound is correct only up to factors of $\mathcal{O}(1)$ and in specific models a smaller $m_{3/2}$ might appear [5].

For a small ratio $\tilde{m}_3^2/m_f^2 \simeq 10^{-4}$ (which follows from eq. (4.2)), the experimental constraints slightly change compared to (3.8) and (3.9) and now read

\[
\begin{align*}
\text{Re} \left( \langle \delta_{12}^d \rangle \right) &\leq 1.3 \times 10^{-2} \left( \frac{m_f}{1 \text{ TeV}} \right)^2, \quad \text{Re}(\delta_{LR}^d)_{12} \leq 5 \times 10^{-3} \left( \frac{m_f}{1 \text{ TeV}} \right); \\
\text{Im} \left( \langle \delta_{12}^d \rangle \right) &\leq 1.1 \times 10^{-3} \left( \frac{m_f}{1 \text{ TeV}} \right), \quad \text{Im}(\delta_{LR}^d)_{12} \leq 4 \times 10^{-4} \left( \frac{m_f}{1 \text{ TeV}} \right),
\end{align*}
\]

from $\Delta m_K$ and $\epsilon_K$, and

\[
(\delta_{MM}^d)_{12} \leq 2 \times 10^{-2} \left( \frac{m_f}{0.3 \text{ TeV}} \right)^2, \quad (\delta_{LR}^d)_{12} \leq 1 \times 10^{-4} \left( \frac{m_f}{0.3 \text{ TeV}} \right),
\]

from the bound on BR($\mu \to e\gamma$). For large scalar masses, the strongest constraint arises in the down-squark sector from $K - \bar{K}$ mixing. (The slepton constraint becomes weaker due to its scaling behaviour.) For $m_3^2 \ll m_f^2$ the $\delta$'s (defined in eq. (3.7)) do not renormalize and are given directly by their boundary values at $M_X\(^\S\)) Off-diagonal scalar mass matrix elements of $\mathcal{O}(m_{3/2})$ then imply $\langle \delta_{12}^d \rangle \simeq 1$ and hence $m_{3/2} > 75 \text{ TeV}$ (or even $m_{3/2} > 650 \text{ TeV}$ if $\text{Im} \langle \delta_{12}^d \rangle \sim \text{Re} \langle \delta_{12}^d \rangle$) is required in order to satisfy eqs. (4.3). This bound is much stronger than the direct bound (4.2).

To summarize, for supersymmetry breaking induced by moduli the gaugino masses are suppressed and the experimental bound on the gluino implies rather large squark and slepton masses. At the same time flavor non-diagonal soft terms are present already at the string tree level and despite the large scalar masses they violate the FCNC bounds. Thus, one needs at least an approximate universality at leading order.

\(^\S\) Non-proportional $A$-terms can renormalize the $\delta$'s and weaken the constraints but this mechanism is not available for $\tilde{m}_3^2 \ll m_f^2$ [18].
4.2. Universal Soft Terms

In the moduli dominated scenario universal soft terms appear at leading order whenever the couplings $Z_{ij}^{(0)}$ satisfy

$$Z_{ij}^{(0)} = h(M, M) \delta_{ij}. \quad (4.5)$$

The unit matrix $\delta_{ij}$ in eq. (4.5) is not the only solution which guarantees universal soft terms. Rather, there could be an arbitrary matrix which only depends on moduli whose $F$-terms vanish but which is independent on all moduli whose $F$-terms break supersymmetry. Indeed, $Z_{ij}^{(0)}$ obeys such a 'split' in string vacua based on $(2,2)$ compactifications where the metric for the $\mathcal{Z}$ (of $E_6$) only depends on the $(1,2)$ moduli through an overall scale factor $\hat{h}$ [39]. (Similarly, the metric for the $\mathcal{Z}$ only depends on the $(1,1)$ moduli through an overall scale factor.)

Using eqs. (2.9) and (4.5) we find

$$\Gamma_{il}^j = \delta_{il} \partial_i \ln h , \quad R_{ijlj} = Z_{ij}^{(0)} \partial_i \partial_j \ln h. \quad (4.6)$$

Inserting into (2.8) and (2.10) results in

$$m_0^2 = |m_{3/2}|^2 - \frac{F^i}{2} \partial_i \partial_j \ln h \sim O(m_{3/2}^2),$$

$$A_{ijj} = F^i \left( \partial_i Y_{ijj} + Y_{ijj} \left( \frac{1}{3} \partial_i \hat{K} - 3 \delta_{ij} \ln h \right) \right) \sim O(m_{3/2} Y) \quad (4.7)$$

at the leading order (string tree level). For Yukawa couplings which only depend weakly on the supersymmetry breaking moduli (i.e. $\partial_i \tilde{Y}_{ijj} \approx 0$) the $A$-terms are strictly proportional to the Yukawa couplings ($A_0 = e^{\hat{K}/2} F^i (\partial_i \hat{K} - 3 \partial_i \ln h)$). However, similar to the dilaton case this universality might be lost at the next order for generic $Z_{ij}^{(1)}$ couplings which do not obey (4.5) and we can estimate the physical consequences implied by such non-universality. The main difference is that now the gaugino masses are much smaller than the scalar masses $\tilde{m}_a^2 \ll m_0^2$ and therefore no renormalization effects enter into the low energy scalar masses; they are directly determined by their boundary value $m_0^2$. Similarly, the $\delta$'s do not renormalize and for both sleptons and squarks we have

$$(\delta^f_{MM})_{12} \simeq \frac{\alpha \chi m_{1/2}^2}{4\pi m_0^2} \simeq 3.3 \times 10^{-3}, \quad (4.8)$$

where we used $m_{1/2}^2 \approx m_0^2$. (The $(\delta^f_{LR})_{12}$ are suppressed by an additional factor of the appropriate fermion mass divided by $m_f$ and thus provide no additional constraint.) Comparing the theoretical prediction (eq. (4.8)) with the experimental bounds (4.3), (4.4), we see that due to the large squark and slepton masses implied by the direct limits (4.2) all FCNC constraints are automatically satisfied.

Finally, let us discuss a specific example of the moduli dominated scenarios which is closely related to no-scale models [19]. For the special case of

$$h = e^{\hat{K}/3} \quad (4.9)$$

in eq. (4.5) (which can also be found in $(2,2)$ vacua), (4.6) and (4.7) imply*

$$\Gamma_{il}^N = \frac{1}{3} \delta_{il} \hat{K}_i , \quad R_{ijlj} = \frac{1}{3} \hat{K}_i Z_{ijlj}, \quad (4.10)$$

$$m_0^2 = 0, \quad A_{ijj} = e^{\hat{K}/2} F^i \partial_i \tilde{Y}_{ijj}. \quad (4.11)$$

If, in addition, the moduli dependence of the Yukawa couplings is weak, $\partial_i \tilde{Y}_{ijj} \approx 0$, the $A_{ijj}$ terms also vanish at tree level and we have instead

$$A_{ijj} = O\left( \frac{\alpha \chi m_{3/2}}{4\pi} \right).$$

Inserting eqs. (4.9) and (4.10) into eq. (2.11) gives, in general, no special cancellations for $B_{ijj}$. Note, however, that if $H_{ij} \approx 0$, then the scale of $B_{ijj}$ is

\[\text{\footnotesize\textbullet} \quad \text{In the dilaton-dominated scenario the renormalization of the squark and slepton masses are driven by the gaugino masses which is the reason for eqs. (3.2), (3.3).}\]

\[\text{\footnotesize\textbullet} \quad \text{Note that this does not require any constraint on $\hat{K}$ itself.}\]
set by $\mu_{IJ}$ which is independent of $m_{3/2}$. In such a case, $B_{IJ}$ could be much smaller than $m_{3/2}^2$ independently of the SUSY breaking mechanism. Therefore, in our analysis below, we allow $B_{IJ}$ to take arbitrary values (as long as they are phenomenologically acceptable).\footnote{\ifnum\value{footnote} = \currentheappage\addtocounter{footnote}{-1}
\ifnum\value{footnote} = \currentheappage\addtocounter{footnote}{-1}
$B_{IJ} = 0$ can also be arranged by choosing $\bar{K}$ appropriately \cite{19}.}

The bound (4.2) still holds but now the scalar masses also vanish at leading order and one expects

$$m_{IJ}^2 \simeq \frac{\alpha_X}{4\pi} m_{3/2}^2 > (850 \text{GeV})^2.$$\hspace{1cm} (4.12)

Hence, most parameters in the observable sector are decoupled from $m_{3/2}$ at the leading order and only arise from string loop effects and with the appropriate suppression. However, the contributions to FCNC processes are generically too large in this scenario. The bounds are somewhat different from (4.3) and (4.4) because in this case $m_{3/2}^2/m_f^2 = 10^{-2}$:

$$\text{Re}\left(\delta_{12}^d\right) \leq 1 \times 10^{-2} \left(\frac{m_{\tilde{d}}}{1 \text{ TeV}}\right), \quad \text{Re}(\delta_{LR}^d)_{12} \leq 6 \times 10^{-3} \left(\frac{m_{\tilde{d}}}{1 \text{ TeV}}\right);$$

$$\text{Im}\left(\delta_{12}^d\right) \leq 8 \times 10^{-4} \left(\frac{m_{\tilde{d}}}{1 \text{ TeV}}\right), \quad \text{Im}(\delta_{LR}^d)_{12} \leq 5 \times 10^{-4} \left(\frac{m_{\tilde{d}}}{1 \text{ TeV}}\right),$$\hspace{1cm} (4.13)

$$\text{Re}\left(\delta_{M(\bar{M})}^d\right)_{12} \leq 1 \times 10^{-1} \left(\frac{m_{\tilde{d}}}{0.3 \text{ TeV}}\right)^2; \quad (\delta_{LR}^d)_{12} \leq 1 \times 10^{-5} \left(\frac{m_{\tilde{d}}}{0.3 \text{ TeV}}\right).$$\hspace{1cm} (4.14)

The strongest constraint again arises in the down-squark sector from $K - \bar{K}$ mixing. For off-diagonal mass matrix elements of the same order as the average scalar masses, one has $(\delta_{M(\bar{M})})_{12} = O(1)$ which implies $m_{\tilde{d}} > 100 \text{ TeV}$ (or even $m_{\tilde{d}} > 1000 \text{ TeV}$ if Im $(\delta_{12}^d) \sim \text{Re}(\delta_{12}^d)$).

The bounds on $m_{3/2}$ from electric dipole moments are of $O(1 \text{ TeV})$ for phases of $O(1)$. Thus, with $m_{3/2} > O(10 \text{ TeV})$ these bounds are always satisfied.

4.3. Cosmological implications

The existence of light moduli, $M_i \sim M_Z$, with couplings to observable particles of order $1/m_F$, poses severe cosmological problems \cite{40–42}. Such moduli are likely to dominate the matter density of the universe until their decay. When they decay, at time $\tau_i \sim m_F^2/M_i^2$, they give a reheat temperature $T_R \sim \sqrt{m_F/\tau_i} \sim 10^{-6} \text{ GeV}$, too low for successful nucleosynthesis ($T_{NS} \sim 10^{-3} \text{ GeV}$).

The cosmological implications of the moduli are drastically different if their masses are much higher than $M_Z$. A particularly interesting range is $M_i \sim 10^3 \text{ TeV}$. If this it the typical mass scale of moduli then \cite{43,44}

1. The universe becomes matter dominated by the heavy moduli long before they decay if the Hubble constant during inflation is larger than the moduli masses.

2. The moduli would decay at time $\tau_i \sim m_F^2/M_i \sim 1 \text{ sec}$. They will give a reheat temperature of $T_R \sim \sqrt{m_F/\tau_i} \sim \text{ a few MeV}$, just right for nucleosynthesis.

3. All decay products will thermalize very fast: with typical number density $n_F \sim T_{R}^4/M_i \sim 10^{-16} \text{ GeV}^3$, hadronic-interaction cross section $\sigma \sim 1/f_{\pi}^2$ and initial velocity $v \sim 1$, the thermalization rate $\sigma v n_F \sim 10^{-14} \text{ GeV}$ is much faster than the expansion rate.

4. Upon thermalization, the number of photons increases to $n_\gamma \sim 10^{-9} \text{ GeV}^3$, but (as baryon multiplicity in hadron scattering is $O(1-10)$) the number of baryons remains essentially unchanged, $n_{B+B} \sim n_F \sim 10^{-7} n_\gamma$. (If CP- and B-violating interactions – either directly in moduli couplings or indirectly in SUSY interactions – induce an asymmetry $\frac{n_B+n_{\bar{B}}}{n_B+n_{\bar{B}}} \sim 10^{-3}$, it would lead to the required baryon symmetry. However, such a large asymmetry is unlikely, as a suppression factor $\leq O(M_F)$ is unavoidable.)

Thus, while light moduli pose serious problems to nucleosynthesis, heavy
moduli \( \langle M_i \sim 100 \text{ TeV} \rangle \) could actually be responsible to nucleosynthesis. From eqs. (4.1), (4.11) and (4.12), we learn that when (a) SUSY is broken by the moduli, (b) \( Z_{I,J}^{(0)} = e^{K/3} \delta_{I,J} \), and (c) the moduli dependence of the Yukawa couplings is weak, then \( \tilde{m}_a, A_{IJJ} \ll m_f \ll m_{3/2} \) while the moduli masses are \( M_i = \mathcal{O}(m_{3/2}) \). The direct experimental bound on \( \tilde{m}_3 \) implies then

\[
\tilde{m}_3 \gtrsim 150 \text{ GeV}, \quad m_f \gtrsim 900 \text{ GeV}, \quad M_i \gtrsim 15 \text{ TeV},
\]

(4.15)

In this scenario, the moduli masses are necessarily heavy and consequently the cosmological problems related to light moduli can be evaded. However, the model faces two problems. First, in the previous section we found that if universality is violated at the string one-loop level, then \( m_f \) (and consequently all other scales) should be at least two orders of magnitude above the bound (4.15). In this case, a major fine-tuning (of order \( \frac{m_f^2}{m_f} \sim 10^{-6} \)) is required to produce the correct electroweak breaking scale, making this scenario very unattractive. In order that it remains viable, there should exist a mechanism that would guarantee universality to high enough string loop level that the various scales actually reside not far above the lower bounds (4.15). Second, even if such a mechanism does exist, the natural scale for \( M_Z \) would still be of \( \mathcal{O}(m_f) \). We were able to show, however, that with fine-tuning of \( \mathcal{O} \left( \frac{1}{15} m_{3/2} \right) \) (of either \( m_t \) or \( m_{\tilde{t}} \)) and \( \mu = \mathcal{O} \left( \frac{1}{15} m_{3/2} \right) \), \( B = \mathcal{O} \left( \frac{1}{15} m_{3/2}^2 \right) \) we get the correct scale for \( M_Z \).

5. Conclusion

In this paper we analyzed the effects of string loop corrections on rare processes at the weak scale. Since only limited information about these corrections is currently available, we estimated their typical order of magnitude and compared them with the stringent bounds implied by the small FCNC.

We find that in the dilaton scenario the experimental bounds can only be satisfied by raising the supersymmetry breaking scale,

\[
\tilde{m}_3 \gtrsim 450 \text{ GeV},
\]

which is a factor of 3 above the scale required by the direct experimental limits. For CP-violating phases of \( \mathcal{O}(1) \), constraints from EDMN require an even larger scale,

\[
\tilde{m}_3 \gtrsim 2.4 \text{ TeV} \sqrt{\sin \phi_B}.
\]

However, both estimates neglect factors of \( \mathcal{O}(1) \). In the moduli scenario the gaugino masses always only appear as string loop corrections and therefore are hierarchically smaller than the scalar masses,

\[
\tilde{m}_3 \gtrsim 150 \text{ GeV}, \quad m_f \gtrsim 15 \text{ TeV}.
\]

Even with this hierarchy, generic tree level soft terms violate the bounds from rare processes. For squarks to have their masses at the lower bound, \( m_{\tilde{q}} \sim 15 \text{ TeV} \), the soft scalar masses that appear at the string tree level have to be universal. Such universality does occur with extra conditions on the metric \( Z_{I,J} \).

We have not explicitly considered the case where moduli and dilaton \( F \)-terms are of the same order of magnitude \( F^S \sim F^a \). If eq. (4.5) holds, the soft parameters are essentially equivalent to the standard MSSM parameters at

\* For another solution of the cosmological moduli problem, that does not require heavy moduli, see ref. [44].

\† The moduli masses can only be much lower than \( m_{3/2} \) for special \( \hat{K} \) and \( \hat{W} \) [19].
leading order with independent $\tilde{m}_{1/2}, m_0^2, \mu, A_0, B$. The gaugino masses are not suppressed and therefore they drive the renormalization of the scalar masses. Without repeating the entire analysis we may conclude that, within the accuracy of our estimates, this leads to similar constraints as were found in the dilaton scenario. That is, the scale of supersymmetry breaking has to be raised compared to the scale required by the direct experimental limits.

In no-scale type scenarios also the scalar masses only appear at the loop level and

$$\tilde{m}_3 \gtrsim 150 \text{ GeV}, \ m_f \gtrsim 900 \text{ GeV}, \ m_{3/2} \gtrsim 15 \text{ TeV}. $$

This leads to the interesting possibility of a large hierarchy between the observable sparticle masses and the moduli masses with interesting cosmological consequences. However, FCNC constraints push the scale to at least two orders of magnitude above the lower bounds. For this scenario to be realistic, universality has to hold well beyond the string one-loop level.

Acknowledgements:

We thank M. Dine and N. Seiberg for initiating this investigation and T. Banks, L. Dixon, F. Eberlein, L. Ibáñez, A. König, S. Pokorski and S. Thomas for useful discussions. Y.N. is an incumbent of the Ruth E. Recu Career Development chair, and is supported in part by the Israel Commission for Basic Research, by the United States – Israel Binational Science Foundation (BSF), and by the Minerva Foundation. J.L. is supported by a Heisenberg fellowship of the DFG and would like to thank the Weizmann Institute and Einstein Center for hospitality and financial support.

REFERENCES

1. For a review see for example, H.-P. Nilles, Phys. Rep. C110 (1984) 1; H.E. Haber and G. Kane, Phys. Rep. C117 (1985) 75; R. Barbieri, Riv. Nuovo Cimento 11 (1988) 1; L.E. Ibáñez and G.G. Ross, in Perspectives in Higgs Physics, ed. G. Kane; F. Zwirner, preprint CERN-TH.6357/91, Talk at the Workshop on Physics and Experiments with Linear Colliders, Saarislka, Finland, Sep. 1991; and references therein.

2. P. Binétruy and M.K. Gaillard, Nucl. Phys. B358 (1991) 121.

3. M. Cvetič, A. Font, L.E. Ibáñez, D. Lüst and F. Quevedo, Nucl. Phys. B361 (1991) 194.

4. L. Ibáñez and D. Lüst, Nucl. Phys. B382 (1992) 305.

5. B. de Carlos, J.A. Casas and C. Muñoz, Phys. Lett. B299 (1993) 234.

6. A. de la Macorra and G.G. Ross, Nucl. Phys. B404 (1993) 321.

7. V.S. Kaplunovsky and J. Louis, Phys. Lett. B306 (1993) 269.

8. A. Brignole, L.E. Ibáñez and C. Muñoz, Nucl. Phys. B422 (1994) 125.

9. For early works see for example, S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150; J. Ellis and D.V. Nanopoulos, Phys. Lett. 110B (1982) 44; R. Barbieri and R. Gatto, Phys. Lett. 110B (1982) 211; M. Duncan, Nucl. Phys. B221 (1983) 285; J. Donoghue, H.-P. Nilles, and D. Wyler, Phys. Lett. B128 (1983) 55; A. Bouquet, J. Kaplan and C.A. Savoy, Phys. Lett. B148 (1984) 69.

10. L.J. Hall, V.A. Kostelecky and S. Raby, Nucl. Phys. B267 (1986) 415; H. Georgi, Phys. Lett. B169 (1986) 231.
11. L.J. Hall, J. Lykken and S. Weinberg, Phys. Rev. D27 (1983) 2359.
12. M. Dine, A. Kagan and S. Samuel, Phys. Lett. B243 (1990) 250.
13. M. Dine and W. Fischler, Phys. Lett. 110B (1982) 227;
   M. Dine and A.E. Nelson, Phys. Rev. D48 (1993) 1277.
14. M. Dine, A. Kagan and R. Leigh, Phys. Rev. D48 (1993) 4269.
15. Y. Nir and N. Seiberg, Phys. Lett. B309 (1993) 337.
16. M. Leurer, Y. Nir and N. Seiberg, Nucl. Phys. B420 (1994) 468.
17. A. Lleyda and C. Muñoz, Phys. Lett. B317 (1993) 82;
   N. Polonsky and A. Pomarol, Phys. Rev. Lett. 73 (1994) 2292;
   T. Kobayashi, D. Suematsu and Y. Yamagishi, Phys. Lett. B329 (1994) 27;
   D. Matalliotakis and H.P. Nilles, Munich preprint TUM-HEP-201/94;
   M. Olechowski and S. Pokorski, Munich preprint MPI-PHT/94-40;
   T. Kobayashi, D. Suematsu, K. Yamada and Y. Yamagishi, preprint KANAZAWA-94-16;
   P. Brax and M. Chemtob, Saclay preprint SACLAY-SPHT-94-128.
18. D. Choudhury, F. Eberlein, A. König, J. Louis and S. Pokorski, Munich preprint LMU–TPW 94-12.
19. E. Cremmer, S. Ferrara, C. Kounnas and D. Nanopoulos, Phys. Lett. 133B (1983) 61;
    for a review see, A.B. Lahanas and D.V. Nanopoulos, Phys. Rep. C145 (1987) 1.
20. For early works see for example,
    J. Ellis, S. Ferrara and D.V. Nanopoulos, Phys. Lett. 114B (1982) 231;
    W. Buchmüler and D. Wyler, Phys. Lett. 121B (1983) 321;
    J. Polchinski and M.B. Wise, Phys. Lett. 125B (1983) 393.
21. I. Antoniadis, E. Gava and K. Narain, Phys. Lett. B283 (1992) 209,
    Nucl. Phys. B383 (1992) 93;
    I. Antoniadis, E. Gava, K. Narain and T. Taylor, Nucl. Phys. B407 (1993) 706.
22. J.K. Kim and H.-P. Nilles, Phys. Rev. Lett. 73 (1994) 1758.
23. S. Ferrara, C. Kounnas and F. Zwirner, CERN preprint CERN-TH-7192-94.
24. L. Dixon, V. Kaplunovsky and J. Louis, Nucl. Phys. B355 (1991) 649.
25. R. Barbieri, J. Louis and M. Moretti, Phys. Lett. B312 (1993) 451,
    erratum ibid. B316 (1993) 632.
26. J. Lopez, D. Nanopoulos and A. Zichichi, Phys. Lett. B319 (1993) 451.
27. Particle Data Group, Phys. Rev. D50 (1994) 1173.
28. F. Gabbiani and A. Masiero, Nucl. Phys. B322 (1989) 235.
29. J. Hagelin, S. Kelley, T. Tanaka, Nucl. Phys. B415 (1994) 293;
    Mod. Phys. Lett. A8 (1993) 2737;
    T. Kosmas, G.K. Leontaris and J.D. Vergados, Phys. Lett. 219B (1989) 457,
    Prog. Part. Nucl. Phys. 33 (1994) 397;
    G.K. Leontaris, Z. Phys. C62 (1994) 91.
30. M. Dugan, B. Grinstein and L. Hall, Nucl. Phys. B255 (1985) 413.
31. W. Fischler, S. Paban and S. Thomas, Phys. Lett. B289 (1992) 373.
32. S. Bertolini and F. Vissiani, Phys. Lett. B324 (1994) 164.
33. I. Altarev et al., JETP Lett. 44 (1986) 460; K. Smith et al., Phys. Lett. B234 (1990) 191.
34. K. Abdullah et al., Phys. Rev. Lett. 65 (1990) 2347.
35. J.P. Jacobs et al., Phys. Rev. Lett. 71 (1993) 3782.
36. T. Kobayashi, M. Konmura, D. Suematsu, K. Yamada and Y. Yamagishi, 
   preprint KANAZAWA-94-17.
37. K. Choi, Phys. Rev. Lett. 72 (1994) 1592.
38. R. Barbieri and G.F. Giudice, Nucl. Phys. B306 (1988) 63.
39. L.J. Dixon, V.S. Kaplunovsky and J. Louis, Nucl. Phys. B329 (1990) 27.
40. G.D. Coughlan, W. Fischler, E.W. Kolb, S. Raby and G.G. Ross, 
   Phys. Lett. 131B (1983) 59.
41. T. Banks, D. Kaplan and A. Nelson, Phys. Rev. D49 (1994) 779.
42. B. de Carlos, A. Casas, F. Quevedo and E. Roulet, Phys. Lett. B318 (1993) 
   447.
43. T. Banks, A. Cohen, G. Moore and Y. Nir, unpublished.
44. L. Randall and S. Thomas, preprint MIT-CTP-2331.