Abstract

The stability of the standard vacuum imposes constraints on flavor violating trilinear soft terms which are stronger than the laboratory bounds coming from the absence of neutral flavor violations (FCNC). Furthermore, contrary to the FCNC bounds, these constraints persist even if the scale of supersymmetry breaking is arbitrarily large.
1 Introduction

Supersymmetry has sources of flavor violation which are not present in the Standard Model [1]. These arise from the possible presence of non-diagonal terms in the squark and slepton mass matrices, which come from the soft breaking Lagrangian, $\mathcal{L}_{\text{soft}}$: 

\[
- \mathcal{L}_{\text{soft}} = \left( m_{\tilde{t}}^2 \right)_{ij} \bar{L}_i L_j + \left( m_{\tilde{\tau}}^2 \right)_{ij} \bar{\tau}_i \tau_j \\
+ \left( m_{\tilde{Q}}^2 \right)_{ij} \bar{Q}_i Q_j + \left( m_{\tilde{u}}^2 \right)_{ij} \bar{u}_i u_j + \left( m_{\tilde{d}}^2 \right)_{ij} \bar{d}_i d_j \\
+ \left[ A_{ij} \bar{L}_i H_1 e_j + A_{ij} \bar{Q}_i H_2 u_j + A_{ij} \bar{Q}_i H_1 d_j + \text{h.c.} \right] + \cdots, \\
\]

where \(i, j = 1, 2, 3\) are generation indices. A usual simplifying assumption of the MSSM is that \(m_{\tilde{t}}^2\) is diagonal and universal and \(A_{ij}\) is proportional to the corresponding Yukawa matrix. However, there is no compelling theoretical argument for these hypotheses.

The size of the off-diagonal entries in \(m_{\tilde{t}}^2, A_{ij}\) is strongly restricted by FCNC experimental data [1, 2, 3, 4, 5, 6, 7, 8]. In this paper we will focus our attention on the \(A_{ij}(f)\) terms; a summary of the corresponding FCNC bounds is given in the second column of Table 1 [4, 8]. The \((\delta_{LR}^{(f)})_{ij}\) parameters used in the table are defined as

\[
(\delta_{LR}^{(f)})_{ij} \equiv \frac{\left( \Delta M_{LR}^2 (f) \right)_{ij}}{M_{av}^{(f)}} ,
\]

where \(f = u, d, l; M_{av}^{(f)}\) is the average of the squared sfermion (\(\tilde{f}_L\) and \(\tilde{f}_R\)) masses and \(\left( \Delta M_{LR}^2 (f) \right)_{ij}\) are the off-diagonal entries in the sfermion mass matrices

\[
\left( \Delta M_{LR}^2 (u) \right)_{ij} = A_{ij}^{(u)} \langle H_2^0 \rangle, \quad \left( \Delta M_{LR}^2 (d) \right)_{ij} = A_{ij}^{(d)} \langle H_1^0 \rangle, \quad \left( \Delta M_{LR}^2 (l) \right)_{ij} = A_{ij}^{(l)} \langle H_1^0 \rangle. \\
\]

In this paper we show that the \(A_{ij}(f)\) terms are also restricted on completely different grounds, namely from the requirement of the absence of dangerous charge and color breaking (CCB) minima or unbounded from below (UFB) directions. As we will see, these bounds are in general stronger than the FCNC ones. Other properties of these bounds are the following:

\(i)\) Some of the bounds, particularly the UFB ones, are genuine effects of the non-diagonal \(A_{ij}(f)\) structure, i.e. they do not have a “diagonal counterpart”.

\(ii)\) Contrary to the FCNC bounds, the strength of the CCB and UFB bounds does not decrease as the scale of supersymmetry breaking increases.

In sections 2 and 3 we derive the bounds. In section 4 we discuss their implication for various theories.

\(^1\)We work in a basis for the superfields where the Yukawa coupling matrices are diagonal.
2 Constraints on $A_{ij}$ from CCB and UFB

Let us start with a CCB bound. Consider the off-diagonal trilinear scalar coupling

$$A_{ij}^{(l)} \bar{e}_L H_1^0 \mu_R + \text{h.c.} \quad (4)$$

Along the field-space direction $|e_L|^2 = |H_1^0|^2 = |\mu_R|^2 = a$, the $SU(3) \times SU(2) \times U(1)$ D-terms are vanishing. Since the phases of the fields can be chosen so that (4) becomes negative, the relevant terms of the potential are

$$V = m_{\mu_R}^2 |e_L|^2 + m_{\mu_R}^2 |\mu_R|^2 + m_1^2 |H_1|^2 + |\lambda_e \bar{e}_L H_1^0|^2 + |\lambda_{\mu_R} H_1^0 |\mu_R|^2 - 2 |A_{ij}^{(l)} \bar{e}_L H_1^0 \mu_R|^2 = (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2) a^2 + (\lambda_e^2 + \lambda_{\mu_R}^2) a^4 - 2 |A_{ij}^{(l)} a^3 . \quad (5)$$

Neglecting the $\lambda_e^2$ term, it is straightforward to check that a deep CCB minimum appears at $a \sim 2 |A_{ij}^{(l)}|/\lambda_{\mu_R}^2$ unless

$$|A_{ij}^{(l)}|^2 \leq \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2) \quad (6)$$

This bound is analogous to the “traditional” CCB bounds for $|A_{ij}^{(d)}| \leq 3 \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2), |A_{ij}^{(l)}| \leq 3 \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2)$. The bound (6) is easily generalized to other couplings

$$|A_{ij}^{(b)}|^2 \leq \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2), \quad k = \text{Max} (i, j)$$

$$|A_{ij}^{(d)}|^2 \leq \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2), \quad k = \text{Max} (i, j)$$

$$|A_{ij}^{(l)}|^2 \leq \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2), \quad k = \text{Max} (i, j) \quad (7)$$

These bounds are, in general, stronger than the corresponding FCNC ones. Actually, they can be made more restrictive by considering extra scalar fields in the potential. This works in the same way as for the “traditional” CCB bounds for diagonal $A$–terms, namely $|A_{11}^{(d)}| \leq 3 \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2), |A_{22}^{(l)}| \leq 3 \lambda_{\mu_R}^2 (m_{\mu_R}^2 + m_{\mu_R}^2 + m_1^2)$. Other possible improvements are more model-dependent, but do not change the order of magnitude of the bounds.

Let us now derive a simple UFB bound. Consider again the trilinear term of eq.(4) and the following direction in the (scalar) field–space

$$|e_L|^2 = |\mu_R|^2 = |\nu_L|^2 = |H_1^0|^2 \equiv a^2 , \quad (8)$$

which of course requires $|H_1|^2 < a^2$. Then, the $SU(2) \times U(1)$ D–terms get cancelled,

$$V_{D-terms} = \frac{1}{8} g^2 \left[ |H_1^0|^2 + |\nu_L|^2 - |e_L|^2 \right]^2 + \frac{1}{8} g^2 \left[ |H_1^0|^2 + |\nu_L|^2 + |e_L|^2 - 2 |\mu_R|^2 \right]^2 = 0 , \quad (9)$$
so that the scalar potential is given by
\[
V = a^2 \left[ m_{eL}^2 + m_{\mu R}^2 + m_{\nu_r}^2 - 2 |A_{12}^{(l)}| H_1^0 + |\lambda_\mu H_1^0|^2 \right] + \left( m_1^2 - m_{\nu_r}^2 \right) |H_1^0|^2 .
\] (10)

Minimizing with respect to $H_1$, we find
\[
H_1^0 = \frac{|A_{12}^{(l)}| a^2}{\lambda_\mu^2 a^2 + (m_1^2 - m_{\nu_r}^2)} ,
\] (11)

which satisfies $|H_1|^2 < a^2$ for large enough values of $a$. Then, the potential of eq.(11) becomes
\[
V = a^2 \left[ m_{eL}^2 + m_{\mu R}^2 + m_{\nu_r}^2 - |A_{12}^{(l)}|^2 \frac{a^2}{\lambda_\mu^2 a^2 + (m_1^2 - m_{\nu_r}^2)} \right] .
\] (12)

So, the potential becomes deeply negative unless
\[
|A_{12}^{(l)}|^2 \frac{a^2}{\lambda_\mu^2 a^2 + (m_1^2 - m_{\nu_r}^2)} \leq \left( m_{eL}^2 + m_{\mu R}^2 + m_{\nu_r}^2 \right) \] (13)

The above equation should be satisfied for any value of $a$ such that $|H_1^0|$, as given by (11), satisfies $|H_1|^2 < a^2$. An interesting limit occurs for $a \gg \frac{(m_1^2 - m_{\nu_r}^2)}{\lambda_\mu^2}$. Then
\[
H_1^0 = \frac{|A_{12}^{(l)}|}{\lambda_\mu^2} ,
\] (14)

and, provided $a^2 > |H_1|^2$, the potential (12) reads
\[
V = a^2 \left[ m_{eL}^2 + m_{\mu R}^2 + m_{\nu_r}^2 - |A_{12}^{(l)}|^2 \frac{a^2}{\lambda_\mu^2} \right] .
\] (15)

and the previous bound (14) simply becomes
\[
|A_{12}^{(l)}|^2 \leq \lambda_\mu^2 \left( m_{eL}^2 + m_{\mu R}^2 + m_{\nu_r}^2 \right) \] (16)

The previous UFB example is useful to check the property (i) mentioned in the introduction. Indeed, it is easy to verify that this kind of UFB direction cannot take place if $A_{ij}$ is diagonal (property (i)) since in that instance the quartic part of the $H_1$ F-term would not be vanishing. Also, notice that property (ii) is a consequence of the fact that both members of eq.(16) scale in the same way as the typical supersymmetry breaking mass increases. Of course, this also holds for the CCB constraints summarized in eqs. (7).

Let us now extend this UFB constraint to the other trilinear terms. For the leptonic ones eq.(16) is immediately generalized to
\[
|A_{ij}^{(l)}|^2 \leq \lambda_{ek}^2 \left( m_{e_{ki}}^2 + m_{e_{Rj}}^2 + m_{\nu_m}^2 \right) , \quad k = \text{Max} \ (i, j) , \ m \neq i, j .
\] (17)
For the $A_{ij}^{(d)}$ terms things work in a very similar way, interchanging $e \leftrightarrow d$. More precisely, taking the following direction in the (scalar) field–space

$$d_{Li}^2 = |d_{Rj}|^2 = |\nu_m|^2 + |H_1^0|^2 \equiv a^2, \quad m \neq i, j$$  

(18)

the $SU(3) \times SU(2) \times U(1)$ D–terms get cancelled,

$$V_{D-terms} = \frac{1}{6} g_3^2 \left[ |d_{Li}|^2 - |d_{Rj}|^2 \right]^2$$

$$+ \frac{1}{8} g_2^2 \left[ |H_1^0|^2 + |\nu_m|^2 - |d_{Li}|^2 \right]^2$$

$$+ \frac{1}{8} g^2 \left[ |H_1^0|^2 + |\nu_m|^2 - \frac{1}{3} |d_{Li}|^2 - \frac{2}{3} |d_{Rj}|^2 \right]^2 = 0 \quad (19)$$

Therefore, following the same steps as in the leptonic case, see eqs. (10)–(17), we arrive at the corresponding UFB bound for $A_{ij}^{(d)}$ terms

$$|A_{ij}^{(d)}|^2 \leq \lambda_{d_k}^{2} \left( m_{d_{Li}}^2 + m_{d_{Rj}}^2 + m_{\nu_m}^2 \right), \quad k = \text{Max} \ (i, j) \quad (20)$$

For the $A_{ij}^{(u)}$ things change since we need the contribution of two sleptons of different generations $|e_{Li}^2|, |e_{Rj}^2| \ (p \neq q)$, rather than $|\nu_m|^2$, in order to cancel the D–terms. To see this, notice that in this case the D–terms have the form (we take $m \neq p, q$ for simplicity)

$$V_{D-terms} = \frac{1}{6} g_3^2 \left[ |u_{Li}|^2 - |u_{Rj}|^2 \right]^2$$

$$+ \frac{1}{8} g_2^2 \left[ |H_2^0|^2 - |e_p|^2 + |e_{Li}|^2 - |u_{Li}|^2 \right]^2$$

$$+ \frac{1}{8} g^2 \left[ |H_2^0|^2 - |\nu_m|^2 - 2 |e_{Rj}|^2 + \frac{1}{3} |u_{Li}|^2 - \frac{4}{3} |u_{Rj}|^2 \right]^2, \quad (21)$$

which are cancelled for

$$|\nu_m|^2 = 0, \quad |e_p|^2 = |e_{Rj}|^2,$$

$$|u_{Li}|^2 = |u_{Rj}|^2 = |e_{Li}|^2 + |H_2^0|^2 \equiv a^2.$$  

(22)

So, following again the steps of eqs. (10)–(17), we obtain the corresponding UFB bound for $A_{ij}^{(u)}$ terms

$$|A_{ij}^{(u)}|^2 \leq \lambda_{u_k}^{2} \left( m_{u_{Li}}^2 + m_{u_{Rj}}^2 + m_{e_{Li}}^2 + m_{e_{Rj}}^2 \right), \quad k = \text{Max} \ (i, j), \ p \neq q.$$  

(23)

The UFB bounds can also be slightly improved by considering extra scalar fields. However, the simplified limits of the bounds, i.e. eqs. (17, 20, 23), cannot be modified in a simple model-independent way.

The CCB and UFB bounds collected in eqs. (4, 17, 20, 23) must be imposed at a renormalization scale, $Q$, of the order of the VEVs of the relevant fields. This means
that the CCB bounds must be evaluated at a scale $Q \sim 2A_{ij}^{(f)}/\lambda_f^2$, while the UFB bounds must be imposed at any possible value of $Q \sim a$. This can be relevant in many instances. For example, for universal gaugino and scalar masses ($M_{1/2}$ and $m$ respectively) satisfying $M_{1/2} \gtrsim m$, the UFB bounds are more restrictive at $M_X$ than at low energies (especially the hadronic ones). This trend gets stronger as the ratio $M_{1/2}/m$ increases.

3 Numerical results

Let us express the previously obtained CCB and UFB bounds in terms of the $(\delta_{LR}^{(f)})_{ij}$ parameters defined in eqs.(2, 3). The CCB bounds, eqs.(7), read

$$
(\delta_{LR}^{(l)})_{ij} \leq \mathcal{M}_{e_k} \frac{[2M_{av}^{(l)} + m_1^2]^{1/2}}{M_{av}^{(l)}} \quad k = \text{Max} \ (i, j)
$$

$$
(\delta_{LR}^{(d)})_{ij} \leq \mathcal{M}_{d_k} \frac{[2M_{av}^{(d)} + m_1^2]^{1/2}}{M_{av}^{(d)}} \quad k = \text{Max} \ (i, j)
$$

$$
(\delta_{LR}^{(u)})_{ij} \leq \mathcal{M}_{u_k} \frac{[2M_{av}^{(u)} + m_2^2]^{1/2}}{M_{av}^{(u)}} \quad k = \text{Max} \ (i, j), \quad (24)
$$

while the UFB bounds, eqs.(17, 20, 23), can be essentially expressed as

$$
(\delta_{LR}^{(l)})_{ij} \leq \frac{\mathcal{M}_{e_k} \sqrt{3}}{M_{av}^{(l)}} \quad k = \text{Max} \ (i, j)
$$

$$
(\delta_{LR}^{(d)})_{ij} \leq \frac{\mathcal{M}_{d_k}}{M_{av}^{(d)}} \quad k = \text{Max} \ (i, j)
$$

$$
(\delta_{LR}^{(u)})_{ij} \leq \frac{\mathcal{M}_{u_k} \left(2M_{av}^{(u)} + 2M_{av}^{(l)}\right)^{1/2}}{M_{av}^{(u)}} \quad k = \text{Max} \ (i, j). \quad (25)
$$

In eqs.(24, 25) $\mathcal{M}_{f_k}$ represents the mass of the fermion $f_k$.

These CCB and UFB bounds are almost always stronger than the corresponding FCNC bounds. This is illustrated in Table 1 for the particular case $\mathcal{M}_{f_k} = 500$ GeV. The only exception is $(\delta_{LR}^{(l)})_{12}$, which is experimentally constrained by the $\mu \rightarrow e, \gamma$ process. As the scale of supersymmetry breaking increases the FCNC bounds are easily satisfied whereas the CCB and UFB bounds continue to strongly constrain the theory. Another case in which the FCNC constraints are satisfied is when approximate “infrared universality” emerges from the RG equations [12, 4, 7]. Again, the CCB and UFB bounds continue to impose strong constraints on such theories. This is because, as argued before, these bounds have to be evaluated at different large scales and do not benefit from RG running.
Table 1: FCNC bounds versus CCB and UFB bounds on \( (\delta_{LR}^{(f)})_{ij} \) for \( M_{av}^{(f)} = 500 \text{ GeV} \). The bounds have been obtained from ref.\[8\] taking \( x = (m_{\text{gaugino}}/M_{av}^{(f)})^2 = 1 \).

4 Implications for supersymmetric and superstring models

4.1 Fritzsch models

In the Fritzsch ansatz \[13\], the Yukawa–coupling matrices of quarks in the interaction basis have the following texture

\[
\lambda^{(u)} = \begin{pmatrix}
0 & \sqrt{\lambda_u \lambda_c} & 0 \\
\sqrt{\lambda_u \lambda_c} & 0 & \sqrt{\lambda_c \lambda_t} \\
0 & \sqrt{\lambda_c \lambda_t} & \lambda_t
\end{pmatrix},
\lambda^{(d)} = \begin{pmatrix}
0 & \sqrt{\lambda_d \lambda_s} & 0 \\
\sqrt{\lambda_d \lambda_s} & 0 & \sqrt{\lambda_s \lambda_b} \\
0 & \sqrt{\lambda_s \lambda_b} & \lambda_b
\end{pmatrix},
\] (26)

where the magnitude of the entries is to be understood in an approximate way. The (1,3), (3,1) entries can be filled, if desired, following the same pattern (e.g. \( \lambda_{13}^{(d)} \sim \sqrt{\lambda_d \lambda_b} \)) with almost no effect in the results. By analogy, a similar texture can be assumed for the lepton matrix, \( \lambda^{(l)}_{ij} \).

Let us make the further assumption that the trilinear soft terms, \( A_{IJK} \phi_I \phi_J \phi_K \), are such that

\[
A_{IJK} \sim O(1) \times M_{\text{SUSY}} \times \lambda_{IJK},
\] (27)

\[\text{6}\]
where $\lambda_{IJK}$ is the corresponding Yukawa coupling in the superpotential. This occurs in simple SUGRA scenarios.

Then the $A_{ij}^{(f)}$ matrices, in the basis where the fermion matrices are diagonalized, have essentially the same texture as the Fritzsch matrices, i.e.

$$A_{ij}^{(f)} \sim O(1) \times M_{\text{SUSY}} \times \lambda_{ij}^{(f)}.$$  \tag{28}

Now it is easy to check that the CCB and UFB conditions obtained in the previous section (see e.g. eqs.(24, 25)) are in general automatically fulfilled since typically $|A_{ij}^{(f)}| \propto \sqrt{\lambda_i \lambda_j} < \lambda_k$ with $k = \text{Max}(i, j)$.

These models, therefore, are safe with respect to CCB and UFB bounds. This, however, is the exception rather than the rule and in general the CCB and UFB bounds strongly constraint theories. Consider, for example, the so called Democratic scenarios \[14\] in which all the elements of the fermion mass matrices are very close to 1. In these theories the approximate proportionality of equation 27 is inadequate and the CCB and UFB constraints impose severe constraints on the models.

### 4.2 Superstring Scenarios

The most interesting application of the CCB and UFB constraints obtained here is to generic SUGRA frameworks, particularly superstring scenarios.

The general SUGRA expression for $A_{IJK}$, as defined in eq.(1), is given by \[15, 16\]

$$A_{IJK} = \frac{1}{3} F^\phi \left[ \partial_\phi \lambda_{IJK} - \Gamma_N^N (\lambda_{IJK})_N + \frac{1}{2} (\partial_\phi K) \lambda_{IJK} \right],$$  \tag{29}

with

$$F^\phi = e^{K/2} \frac{\partial}{\partial \Phi} (\partial_\phi W + (\partial_\phi K) W), \quad g_{IJ} \equiv \partial^2 K/\partial \Phi_I \partial \Phi_J, \quad \Gamma_N^N \equiv g_{N J} \partial_\phi g_{J 1}. \tag{30}$$

$K$ and $W$ are the Kähler potential and the original SUGRA superpotential respectively, while $I, J, L, N$, as well as $\phi$ run over all the chiral fields. $\lambda_{IJK}$ are the Yukawa couplings in the effective superpotential (they are a factor $e^{K/2}$ the original ones).

It is clear from (29) that only the third term in the right hand side satisfies the proportionality relation (27). The second term is particularly relevant for our concern since it mixes different Yukawa couplings through the non-trivial structure of the Kähler manifold. Thus, the constraints summarized in eqs.(24, 25) put non-trivial constraints on geometric properties of the SUGRA structure.

In superstring theories the Kähler metric, $g_{IJ}$, for the observable fields depends at tree level on the moduli, $T_i$, and at higher orders also on the dilaton, $S$ \[17\]. The Yukawa couplings depend at tree level on the moduli, with some exceptions, as for the untwisted fields in orbifold constructions. As a consequence, the general situation is $A_{IJK} = O(m_{3/2})$ \[16, 18\]. A exception to this rule occurs for dilaton-dominance SUSY breaking, i.e. when only $F_S \neq 0$. In that case, working at tree level, one gets exact universality of the soft breaking terms (although this universality is spoiled at higher orders). However, for moduli-dominated SUSY breaking this will not be the general case.
Here, rather than making a detailed analysis of all the possibilities, we will illustrate the relevance of the CCB and UFB bounds by applying them to an interesting superstring inspired model that has been discussed in the literature [4]. Namely, consider the following ansatz for the $A^{(f)}$ matrices

$$A^{(f)} = \begin{pmatrix} 0 & 0 & A^f \\ 0 & 0 & A^f \\ A^f & A^f & A^f \end{pmatrix}, \quad (31)$$

where $f = u, d, l$. The off-diagonal entries in this matrix are useful in order to decrease the low–energy magnitude of the additional sources of flavor violation $\left(\Delta m_{LL}^{(f)}\right)^2$, $\left(\Delta m_{RR}^{(f)}\right)^2$, through the RG running provided $A^f$ is large enough.

The authors of ref.[4] showed that for $A^f = O(1) \times m_{3/2}$ the model is safe with respect to all the FCNC constraints, even if the initial values of $\left(\Delta m_{LL}^{(f)}\right)^2$ are $O(m_{3/2})$, provided that a moderate hierarchy $M_{1/2}/m_{3/2} = O(10)$ takes place. A slighter, but still appreciable, improvement is obtained by setting $A^u = O(\lambda_t) \times m_{3/2}$, $A^d = O(\lambda_b) \times m_{3/2}$, $A^l = O(\lambda_t) \times m_{3/2}$. It is interesting to stress that, from the previous discussion on $A$-terms in superstring constructions, this scenario may occur in the framework of superstrings if the $A$-terms associated to the lighter generations are small.

Nevertheless, it is easy to check that this model does not survive the CCB and UFB bounds (see e.g. eqs.(24, 25). Even in the mentioned (less strong) case in which $A^f$ is proportional to the largest Yukawa coupling of the $f$-type, the CCB and UFB bounds can only be satisfied by proper choice of the values of the various $O(1)$ constants.

## 5 Remarks

The CCB and UFB bounds presented here are conservative; they correspond to sufficient, but not necessary, conditions for the stability of the standard vacuum. It is possible that we live in a metastable vacuum [19, 20], whose lifetime is longer than the age of the universe. This softens the constraints obtained here. However, it is conceptually difficult to understand how the cosmological constant is vanishing precisely in a local “interim” vacuum. It is also interesting that many of the UFB directions found here are really unbounded from below and, if present, make the theory ill defined until Planckean physics comes to the rescue.

To conclude, the stability bounds presented here are one more manifestation of the supersymmetric flavor problem [3, 4]. They have the unique feature that they cannot be satisfied by simply increasing the scale of supersymmetry breaking. The simplest cure to all the flavor problems is found in theories where supersymmetry breaking originates at low energies and is communicated to the ordinary sparticles via gauge interactions.

## References
[1] S. Dimopoulos and H. Georgi, *Nucl. Phys.* **B193** (1981) 150

[2] F. Gabbiani and A. Masiero *Nucl. Phys.* **B322** (1989) 235

[3] J.S. Hagelin, S. Kelley and T. Tanaka *Nucl. Phys.* **B415** (1994) 293

[4] D. Choudhury, F. Eberlein, A. Konig, J. Louis and S. Pokorski, *Phys. Lett.* **B342** (1995) 180

[5] R. Barbieri and L.J. Hall, *Phys. Lett.* **B338** (1994) 212

[6] S. Dimopoulos and D. Sutter *Nucl. Phys.* **B452** (1995) 496

[7] B. de Carlos, J.A. Casas and J.M. Moreno, *Phys. Rev.* **D53** (1996) 6398.

[8] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, [hep-ph/9604387](http://arxiv.org/abs/hep-ph/9604387)

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

[10] M. Drees, M. Glück and K. Grassie, *Phys. Lett.* **B157** (1985) 164; J.F. Gunion, H.E. Haber and M. Sher, *Nucl. Phys.* **B306** (1988) 1; H. Komatsu, *Phys. Lett.* **B215** (1988) 323, P. Langacker and N. Polonsky, *Phys. Rev.* **D50** (1994) 2199

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

[12] M. Dine, A. Kagan and S. Samuel, *Phys. Lett.* **B243** (1990) 250

[13] H. Fritzsch, *Phys. Lett.* **B70** (1977) 418; P. Ramond, R.G. Roberts and G.G. Ross, *Nucl. Phys.* **B406** (1993) 19

[14] H. Fritzsch, D. Holtmannspotter Phys.Lett.B338:290-294,1994.

[15] L. Hall, J. Lykken and S. Weinberg, *Phys. Rev.* **D27** (1983) 2359

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

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

[18] J. Louis and Y. Nir, *Nucl. Phys.* **B447** (1995) 18.

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

[20] A. Kusenko, P. Langacker and G. Segre, UPR-0677-T, [hep-ph/9602414](http://arxiv.org/abs/hep-ph/9602414).