Disoriented and Plastic Soft Terms: 
A Dynamical Solution to the Problem 
of Supersymmetric Flavor Violations

S. Dimopoulos,\textsuperscript{a} G.F. Giudice\textsuperscript{a} and N. Tetradis\textsuperscript{b}

\textsuperscript{a}Theoretical Physics Division, CERN  
CH-1211 Geneva 23, Switzerland  
\textsuperscript{b}Theoretical Physics, University of Oxford  
1 Keble Road, Oxford OX1 3NP, U.K.

Abstract

We postulate that the orientation of the soft supersymmetry-breaking terms in flavor space is not fixed by tree level physics at the Planck scale; it is a dynamical variable which depends on fields that have no tree level potential. These fields can be thought of as either moduli or as the Nambu-Goldstone bosons of the spontaneously broken flavor symmetry which is non-linearly realized by the soft terms. We show that the soft terms align with the quark and lepton Yukawa couplings, just as spins align with an external magnetic field. As a result, the soft terms conserve individual lepton numbers and do not cause large flavor or CP violations. The vacuum adjusts so as to allow large sparticle splittings to naturally coexist with flavor conservation. Consequently, the resulting phenomenology is different from that of minimal supersymmetric theories. We also propose theories in which the shape of the soft terms in flavor space is a dynamical variable which depends on fields that have no tree level potential. This dynamically leads to partial degeneracy among sparticles and further supression of flavor violations. The ideas of this paper suggest a connection between the space of moduli and the spontaneously broken flavor group.

\textsuperscript{1}On leave of absence from the Physics Department, Stanford University, Stanford CA 94305, USA. 
\textsuperscript{2}On leave of absence from INFN, Sezione di Padova, Padua, Italy.
1. Universal versus Disoriented Soft Terms

The soft supersymmetry (SUSY)-breaking terms \([1, 2]\) are important for at least two reasons. First, they are the key ingredient which made the construction of realistic supersymmetric theories possible \([1]\). Second, they are experimentally measurable quantities since they determine the masses of sparticles. In early works, motivated by the need to avoid large flavor violations, it was postulated that soft terms satisfy universality \([1]\). Universality states that the squarks and sleptons of the three families are all degenerate in mass at some scale \(\sim M_{\text{GUT}}\).

Universality has a geometric interpretation which is useful to appreciate. To do this, consider the limit in which all but the gauge couplings of the supersymmetric standard model are set to zero. The resulting theory possesses a \(U(3)^5\) global symmetry which is called flavor symmetry. The 3 stands for the number of families and the 5 for the number of \(SU(3) \times SU(2) \times U(1)\) superfield members in a family, which will be labelled by \(A = Q, \bar{U}, \bar{D}, L, \bar{E}\). The flavor symmetry is simply a manifestation of the fact that gauge forces do not distinguish particles with identical gauge quantum numbers. Universality states that the five \(3 \times 3\) sparticle squared mass matrices \(m^2_A\) are flavor singlets, \(i.e.\) proportional to the identity. They are spheres in flavor space and they realize the flavor symmetry in the Wigner mode. In this paper we wish to suggest an alternative mechanism to universality for avoiding large flavor violations.

Let \(\Lambda\) be a high energy scale at which supersymmetry breaking occurs and the soft terms are determined. \(\Lambda\) can be of the order of the Planck mass \(M_{\text{Pl}}\) – as in supergravity – or smaller, equal to the mass of the messengers that communicate supersymmetry breaking to the ordinary particles. Our fundamental hypothesis is that physics at the scale \(\Lambda\) fixes the eigenvalues of the soft terms \(m^2_A\) but leaves their direction in flavor \(U(3)^5\) space undetermined. In other words, the potential energy \(V_\Lambda\) of the sector which determines the soft terms at the scale \(\Lambda\) is flavor \(U(3)^5\) invariant. \(V_\Lambda\) does not depend on the \(U(3)^5\) angles which are flat directions of the potential and which will be called here “moduli”. The moduli determine the direction in which the soft terms point in flavor space. They can be thought of as the Goldstone bosons of the flavor group which is spontaneously broken by the soft terms \(m^2_A\) themselves and are therefore “disoriented” in flavor space. Therefore, the simplest way to state our hypothesis is: the soft terms realize the flavor symmetry in the Goldstone mode. In contrast, universality states that the soft terms realize the flavor symmetry in the Wigner mode.

Our next assumption is that at energies below \(\Lambda\) we have the minimal supersymmetric particle content \([3]\) (along with the decoupled gauge singlet Goldstones/moduli). We will show that the orientation of the soft terms is determined by physics at lower energies – in particular the flavor-breaking fermion masses – in a calculable way.

A simple analogy is to think of the soft terms \(m^2_A\) as a spin \(\vec{s}\) in space and \(U(3)_A\) as ordinary rotational invariance. The magnitude of \(\vec{s}\) is determined by some unspecified “high energy” dynamics to be non-zero. This forces rotational invariance to break sponta-
neously. $\vec{s}$ can point in any direction until we turn on an external magnetic field $\vec{B}$ which explicitly breaks the rotational invariance and forces $\vec{s}$ to align parallel to $\vec{B}$. Notice that alignment (or anti-alignment) is preferred and the maximal subgroup possible, $SO(2)$, is preserved. This completes the analogy between $\vec{s}$ and the soft terms on one hand and between $\vec{B}$ and the fermion masses on the other. Perfect alignment would mean that the maximal subgroup consisting of the product of all vectorial $U(1)$ quantum numbers is preserved and consequently there is no flavor violation. In the quark sector since the Kobayashi-Maskawa matrix $K \neq 1$ this is not possible, but the dynamics will adjust as to reduce flavor violations.

2. Alignment

Consider the supersymmetric $SU(3) \times SU(2) \times U(1)$ theory with minimal particle content, whose gauge interactions possess an $U(3)^5$ global flavor symmetry. As in the standard model, the Yukawa couplings break the symmetry. In addition, flavor symmetry is violated here also by the soft SUSY-breaking terms which in general lead to phenomenologically unacceptable contributions to flavor-changing neutral current (FCNC) processes. Let us concentrate first on the soft SUSY-breaking masses $m^2_A$. Our hypothesis is that the $m^2_A$ are general Hermitian matrices whose eigenvalues are fixed at the high scale $\Lambda$ where supersymmetry is broken, but whose orientation is a dynamical variable determined by physics below $\Lambda$. The soft SUSY-breaking masses $m^2_A$ are thus promoted to fields:

\[ m^2_A \rightarrow \Sigma_A \equiv U_A^\dagger \Sigma_A U_A \quad A = Q, \bar{U}, \bar{D}, L, \bar{E}. \]  

(2.1)

$\Sigma_A$ are diagonal matrices with real, positive eigenvalues ordered according to increasing magnitude and $U_A$ are $3 \times 3$ unitary matrices.

Our fundamental hypothesis can now be restated: $\Sigma_A$ are fixed by physics at some very high scale $\Lambda$ – say $\Lambda \sim M_{PL}$, for concreteness – whereas $U_A$ are determined only by lower energy physics, namely the energetics of the supersymmetric $SU(3) \times SU(2) \times U(1)$ theory. For any given $A$, let us write

\[ U_A = \exp \left( i \sum_\alpha \lambda^\alpha \sigma^\alpha_A \right), \]

(2.2)

where $\lambda^\alpha/2$ are the generators of the flavor group broken by $\Sigma_A$, in short the six generators of $SU(3)/U(1)^2$. Thus $\sigma^\alpha_A$ can be thought of as the Goldstone bosons of the flavor $U(3)$ group that has been spontaneously broken by the $\Sigma_A$ VEV. In reality, the $\sigma^\alpha$ are pseudo-Goldstone bosons, because quark and lepton masses explicitly break flavor invariance. According to our fundamental hypothesis the potential $V_\Lambda(\sigma^\alpha_A)$ of the soft terms at the scale $\Lambda$ is flat, so that the expectation value of $\sigma^\alpha_A$ is undetermined. However, the effective potential at a lower scales (such as the supersymmetry scale $m_s$ or the weak scale) receives

\[ \text{The possibility that the third generation Yukawa couplings depend on dynamical variables was considered in ref. \[3\]; similar suggestions have also been proposed in ref. \[4\].} \]
quantum corrections from the integration of fluctuations with characteristic momenta between $\Lambda$ and $m_s$. It is the dynamics of these fluctuations that fixes the value of $\sigma^A_\alpha$ in such a way that the soft SUSY-breaking mass terms are aligned with the Yukawa couplings.

The natural setting in which to carry our discussion is provided by the approach to the renormalization group introduced by Wilson [5]. In his formalism the effects of quantum fluctuations with characteristic momenta $q^2$ larger than a given cutoff $k^2$ are included in a $k$-dependent effective action $\Gamma_k$. The scale $k$ can be viewed as the coarse-graining scale, beyond which the details of the system are not probed. As a result fluctuations with characteristic wavelengths smaller than $2\pi/k$ are integrated out and their effects are incorporated in the couplings in $\Gamma_k$. An exact renormalization group equation describes how the effective action $\Gamma_k$ changes as the scale $k$ is lowered and the effects of fluctuations with larger wavelengths are taken into account.

In our problem $k$ can be identified initially with the high scale $\Lambda$ where supersymmetry is broken. We are interested in the effect of fluctuations on the shape of the potential $V_k(\sigma_A)$ as the scale $k$ is lowered from $k = \Lambda$ to $k = m_s$. In appendix A we derive the equation which describes the evolution of $V_k(\sigma_A)$. It is

$$\frac{\partial V_k(\sigma_A)}{\partial t} = -\frac{k^4}{16\pi^2} \text{Str} \log \left[ 1 + \frac{\mathcal{M}^2(\sigma_A, k)}{k^2} \right], \tag{2.3}$$

where $t = \log(k/\Lambda)$ and $\mathcal{M}^2(\sigma_A, k)$ is the running mass matrix of the theory. This equation, when combined with the evolution equation for the running mass matrix

$$\frac{\partial \mathcal{M}^2(\sigma_A, k)}{\partial t} = \beta_{\mathcal{M}}(\xi^i(k)), \tag{2.4}$$

describes how the potential evolves as the coarse-graining scale is lowered and fluctuations with smaller characteristic momenta are incorporated in it. The $\beta$-function for the mass matrix can be obtained from the $\beta$-functions $\beta_{\xi^i}$ of the running couplings of the theory $\xi^i(k)$. The boundary conditions at the scale $k = \Lambda$ are given by the assumed (flat) form of $V_\Lambda(\sigma_A)$ and the tree level form of the mass matrix $\mathcal{M}^2(\sigma_A, \Lambda) = \mathcal{M}^2(\sigma_A)$. We have not taken into account the fluctuations of the Goldstone fields $\sigma^A_\alpha$, despite the fact that they are massless at tree level. The reason is that their contributions to the effective potential which introduce a non-trivial $\sigma^A_\alpha$ dependence are suppressed by powers of $\Lambda$ relative to the ones we have included. This can be checked in perturbation theory if we use the fields $\sigma^A_\alpha = \sigma^A_\alpha \Lambda$ which have appropriate mass dimensions and consider general kinetic and potential terms, invariant under non-linear realizations of the $SU(3)/U(1)^2$ symmetry. The $\beta$-functions $\beta_{\xi^i}$, which are needed for the calculation of $\beta_{\mathcal{M}}$, must be consistently calculated within the scheme that we have introduced. However, in an expansion in powers of the couplings they can be obtained from standard perturbative calculations [8] (at least to one loop, where no scheme dependence is expected).

Let us first consider eq. (2.3) with constant $\mathcal{M}^2(\sigma_A, k) = \mathcal{M}^2(\sigma_A, \Lambda) = \mathcal{M}^2(\sigma_A)$, for
which it can be easily integrated. Keeping the leading terms in $\Lambda$ for $k = 0$ we find

$$V_0(\sigma_A) = V_\Lambda(\sigma_A) + \frac{1}{32\pi^2} \Lambda^2 \text{Str} \mathcal{M}^2(\sigma_A) + \frac{1}{64\pi^2} \text{Str} \mathcal{M}^4(\sigma_A) \log \left( \frac{\mathcal{M}^2(\sigma_A)}{\Lambda^2} \right). \quad (2.5)$$

This the standard one loop result for the effective potential. For most of the evolution from $k = \Lambda$ to $k = 0$ described by eq. (2.3), we have $\tilde{\mathcal{M}}^2/k^2 \ll 1$ and the logarithm can be expanded around one, so that

$$\frac{\partial V_k(\sigma_A)}{\partial t} = -\frac{k^2}{16\pi^2} \text{Str} \tilde{\mathcal{M}}^2(\sigma_A, k) + \frac{1}{32\pi^2} \text{Str} \tilde{\mathcal{M}}^4(\sigma_A, k) \ldots \quad (2.6)$$

For constant $\tilde{\mathcal{M}}^2$ this approximation leads to

$$V_k(\sigma_A) = V_\Lambda(\sigma_A) + \frac{1}{32\pi^2} (\Lambda^2 - k^2) \text{Str} \mathcal{M}^2(\sigma_A) + \frac{1}{64\pi^2} \text{Str} \mathcal{M}^4(\sigma_A) \log \left( \frac{k^2}{\Lambda^2} \right). \quad (2.7)$$

Comparison with eq. (2.5) indicates that the part of the evolution with $k^2 \lesssim \tilde{\mathcal{M}}^2$ simply takes into account threshold effects which lead to the decoupling of the massive modes. As a result the masses replace $k^2$ as an effective infrared cutoff in the logarithm. The quadratic contribution, on which our discussion is based, is unaffected by the approximation. We replace, therefore, eq. (2.3) by eq. (2.6) and neglect the second term in the r.h.s.

We use the perturbative expressions of ref. [8] for the $\beta$-functions $\beta_{\xi^i}$. This is expected to be a good approximation for the small couplings relevant for our investigation. The general form of the perturbative $\beta$-function is

$$\beta_{\tilde{\mathcal{M}}} = \frac{1}{16\pi^2} \beta^{(1)}_{\tilde{\mathcal{M}}} + \frac{1}{(16\pi^2)^2} \beta^{(2)}_{\tilde{\mathcal{M}}} \ldots, \quad (2.8)$$

where only the first two loop contributions are considered. Starting from the above expression we can iteratively derive an approximate solution of eq. (2.4) for the small values of $|\Delta t|/16\pi^2 = |\log(m_s/\Lambda)|/16\pi^2 \simeq (32 - 37)/16\pi^2$ which are relevant for our problem. The leading terms are given by

$$\tilde{\mathcal{M}}^2(\sigma_A, k) = \mathcal{M}^2(\sigma_A) + \frac{1}{16\pi^2} \beta^{(1)}_{\mathcal{M}} \log \left( \frac{k}{\Lambda} \right) + \frac{1}{(16\pi^2)^2} \beta^{(2)}_{\mathcal{M}} \log \left( \frac{k}{\Lambda} \right) + \frac{1}{2(16\pi^2)^2} \left[ \beta^{(1)}_{\xi^i} \partial \beta^{(1)}_{\mathcal{M}} / \partial \xi^i \right] \log^2 \left( \frac{k}{\Lambda} \right) \ldots \quad (2.9)$$

The $\beta$-functions are evaluated at $k = \Lambda$ in terms of the tree-level values of the couplings and masses. The last term in the second line of the above expression results from the quadratic term in the Taylor expansion of $\tilde{\mathcal{M}}^2$. For the quantity in the square brackets summation over $i$ is assumed and $\beta_{\xi^i}$ (which is in general a matrix) must be substituted at the point where the derivative with respect to $\xi^i$ is taken in the expression for $\beta^{(1)}_{\mathcal{M}}$. The integration of eq. (2.6) is now straightforward. We find

$$V_{m_s}(\sigma_A) = V_\Lambda(\sigma_A) + \Lambda^2 \frac{1}{32\pi^2} \text{Str} \left[ \mathcal{M}^2 - \frac{1}{32\pi^2} \beta^{(1)}_{\mathcal{M}} + \frac{1}{1024\pi^4} \left[ \beta^{(1)}_{\xi^i} \partial \beta^{(1)}_{\mathcal{M}} / \partial \xi^i - 2 \beta^{(2)}_{\mathcal{M}} \right] \right]. \quad (2.10)$$
As we have already pointed out the r.h.s. of eq. (2.10) must be evaluated in terms of the tree-level values of the parameters $\xi^l$ at the scale $\Lambda$. We also need to express the mass matrix of the theory in terms of $\xi^l$. Our treatment is simplified by the fact that the Higgs field has not yet developed an expectation value at $k = \Lambda$. The complications arising from the non-zero Higgs field expectation value at low scales are neglected in our approximation. The induced error is small, as the most significant contributions in the integration of eq. (2.6) come from scales $k \sim \Lambda$.

We start by considering the slepton and Higgs mass matrices, which are given by

$$
\mathcal{M}_{e,u,d}^2 = \begin{pmatrix} m_{L,Q,Q}^2 & 0 \\ 0 & m_{E,\bar{U},\bar{D}}^2 \end{pmatrix}
$$

and

$$
\mathcal{M}_H^2 = \begin{pmatrix} m_{H_u}^2 + \mu^2 & B^\dagger \\ B & m_{H_d}^2 + \mu^2 \end{pmatrix}
$$

The term $\text{Str}\mathcal{M}^2$ in the r.h.s. of eq. (2.10) gives a $\sigma_3^0$-independent contribution. The one loop $\beta$-functions for $m_A$ and $m_{H_u,d}^2$ which are relevant for the discussion of the orientation of the $\sigma_3^0$ fields can be obtained from ref. [8]. We list only the parts which remain $\sigma_3^0$-dependent after the trace is taken:

$$
\begin{align}
\text{Tr}\beta^{(1)}_{m_L^2} &= \text{Tr} \left[ m_L^2 \lambda_e^\dagger \lambda_e + 2 \lambda_e^\dagger m_L^2 \lambda_e + \lambda_e^\dagger \lambda_e m_L^2 \right] + (\sigma_3^0 - \text{indep.}) \\
\text{Tr}\beta^{(1)}_{m_E^2} &= \text{Tr} \left[ 2 m_E^2 \lambda_e^\dagger \lambda_e + 4 \lambda_e^\dagger m_E^2 \lambda_e + 2 \lambda_e^\dagger \lambda_e m_E^2 \right] + (\sigma_3^0 - \text{indep.}) \\
\text{Tr}\beta^{(1)}_{m_Q^2} &= \text{Tr} \left[ m_Q^2 (\lambda_u^\dagger \lambda_u + \lambda_d^\dagger \lambda_d) + (\lambda_u^\dagger \lambda_u + \lambda_d^\dagger \lambda_d) m_Q^2 + 2 \lambda_u^\dagger m_Q^2 \lambda_u + 2 \lambda_d^\dagger m_Q^2 \lambda_d \right] \\
&\quad + (\sigma_3^0 - \text{indep.}) \\
\text{Tr}\beta^{(1)}_{m_H^2} &= \text{Tr} \left[ m_Q^2 \lambda_u^\dagger \lambda_u + \lambda_u^\dagger m_H^2 \lambda_u \right] + (\sigma_3^0 - \text{indep.}) \\
\beta^{(1)}_{m_H^2} &= 6 \text{Tr} \left[ m_Q^2 \lambda_u^\dagger \lambda_u + \lambda_u^\dagger m_H^2 \lambda_u \right] + (\sigma_3^0 - \text{indep.}) \\
\beta^{(1)}_{m_H^2} &= \text{Tr} \left[ 6 m_Q^2 \lambda_d^\dagger \lambda_d + 6 \lambda_d^\dagger m_Q^2 \lambda_d + 2 m_L^2 \lambda_e^\dagger \lambda_e + 2 \lambda_e^\dagger m_L^2 \lambda_e \right] + (\sigma_3^0 - \text{indep.}),
\end{align}
$$

where $\lambda_{e,u,d}$ are the Yukawa matrices. We choose a basis in which the the matrices $\lambda_{e,u}$ are diagonal and related to the observable fermion mass matrices at low energies through $\lambda_e = m_e/v_1, \lambda_u = m_u/v_2$, where $v_{1,2}$ are the the Higgs field expectation values. Then $\lambda_d$ is given by $\lambda_d = m_d K^{\dagger}/v_1$, with $K$ the CKM matrix and $m_d$ diagonal. The trace of $\beta^{(1)}_M$ can now be easily evaluated, with the result

$$
\begin{align}
\text{Str}\beta^{(1)}_M &= \text{Tr} \left[ 8 \lambda_e^\dagger \lambda_e m_L^2 + 8 \lambda_e^\dagger \lambda_e m_E^2 + 14 (\lambda_u^\dagger \lambda_u + \lambda_d^\dagger \lambda_d) m_Q^2 + 14 \lambda_u^\dagger m_Q^2 \lambda_u + 14 \lambda_d^\dagger m_Q^2 \lambda_d \right].
\end{align}
$$
Let us consider the first term in the r.h.s. of the above expression. Starting from the definitions of eqs. (2.1), (2.2) we find for small $\sigma_\alpha^L$

$$\text{Tr}(\lambda^\dagger \lambda m^2_L) = -\sum_\alpha (\sigma_\alpha^L)^2 \sum_{i>j} |\lambda^\alpha_{ij}|^2 \frac{1}{e^2} \left( \tilde{\Sigma}_{Li} - \tilde{\Sigma}_{Lj} \right) \left( m^2_{\epsilon_i} - m^2_{\epsilon_j} \right) + \mathcal{O}(\sigma^3),$$  

(2.21)

with $m^2_i$ the charged lepton masses ordered according to increasing magnitude. Clearly the effective potential, as it is given by eqs. (2.10), (2.20), (2.21) has a minimum at $\sigma^L_\alpha = 0$. The same conclusion can be easily reached for $\Lambda = \tilde{E}$. The result $U_L = U_E = 1$ has the important consequence that the $e, \mu, \tau$ lepton numbers are separately conserved. Since slepton and lepton mass matrices are parallel, they both preserve the same $U(1)^3$ symmetry and individual lepton number violating processes like $\mu \rightarrow e\gamma$ do not occur in this theory. Complete alignment of the squark-quark sectors is not possible due to the presence of the Kobayashi-Maskawa matrix. The soft terms $m^2_U$, $m^2_D$ align with the quark mass matrices $m^2_u$, $m^2_d$ respectively, while $m^2_Q$ aligns with the linear combination $m^2_u + Kr^2m^2_dK^\dagger$. In this way FCNC processes are adequately suppressed.

The predicted masses of the pseudo-Goldstone bosons $\sigma^A_\alpha$ can be computed from the effective potential of eqs. (2.10), (2.20), (2.21). We find the approximate expression

$$m^2_{\sigma^A_\alpha} = \frac{m^2_s}{16\pi^2} \sum_{i>j} |\lambda^\alpha_{ij}|^2 \frac{m^2_{\epsilon_i} - m^2_{\epsilon_j}}{m^2_{\text{weak}}},$$  

(2.22)

where $m_{\epsilon_i}$ are the masses of fermions in the superfield $A$, and we have assumed that the sparticle splittings are of the order $m_s$ and approximated the Kobayashi-Maskawa matrix by the unit matrix. The pseudo-Goldstone masses are proportional to the scale of the soft terms $m_s$, and to the Yukawa couplings, which explicitly break the flavor symmetry. They range roughly between $(10^{-3} - 1)$ GeV. The couplings of the pseudo-Goldstone bosons to the particles of the minimal supersymmetric standard model are extremely weak as they are suppressed by $m_s/\Lambda$.

The quadratic momentum dependence in eq. (2.6) shows that the alignment is determined by the behavior of the theory at energies just below the scale $\Lambda$. This feature is not appealing since it introduces a sensitivity to the details of the high-energy physics. In view of this it is not appropriate to think of the Goldstones/moduli as determined by low-energy physics. Unfortunately this ultraviolet sensitivity is bound to frustrate all attempts to convert parameters of the supersymmetric theory – through their dependence on moduli – to dynamical variables of the low-energy theory. It originates in the quadratic dependence of the energy on the cut-off, a feature present in theories of softly-broken low-energy supersymmetry.

3. Alignment of the $A$-terms

The triscalar $A$-terms break both supersymmetry and chirality; thus they resemble the soft masses $m^2_A$ in one sense and the Yukawa couplings in another. Consequently there

---

5Similar observations were made in ref. [9].
are three possibilities:

a) The first is that the $A$-terms are disoriented and independent of $m^2_A$. They can be parametrized as

$$A_a = V_a \bar{\Delta}_a V_a,$$

with $a = e, u, d$. $\bar{\Delta}_a$ is the diagonal matrix of the eigenvalues of $A_a$, and $V_a, \bar{V}_a$ are unitary matrices, analogous to $U_A$, which include new Goldstone fields $\delta^\alpha_a, \bar{\delta}^\alpha_a$ whose values are postulated to be undetermined by the potential $V_\Lambda$ at the high scale $\Lambda$. Their expectation values at low scales are determined by the minimum of the potential of eq. (2.10), in which the additional fields $\delta^\alpha_a, \bar{\delta}^\alpha_a$ now appear. The term proportional to $\text{Str} \beta(1)_M$ in the r.h.s. of eq. (2.10) is $A_a$-independent and does not determine $\delta^\alpha_a, \bar{\delta}^\alpha_a$. One has to evaluate the last two terms which do depend on $A_a$ and can fix the values of $\delta^\alpha_a, \bar{\delta}^\alpha_a$. Making use of the results of ref. [8] for $\beta(1)_i, \beta(2)_i$ we conclude that all the terms which depend quadratically on $A_a$, e.g. $\text{Tr} A^\dagger A \lambda^\dagger \lambda$, come with a positive sign in the potential and anti-align the $A$-terms with the Yukawa couplings. This still implies the existence of a $U(1)^3$ symmetry – approximate for quarks, exact for leptons – which suppresses flavor violations. However, the terms linear in $A_a$, e.g. $\text{Tr} A \lambda^\dagger$, which are proportional to the gaugino masses, align the $A$-terms with the Yukawa couplings, irrespectively of their sign. Therefore, depending on whether the terms linear or quadratic in $A_a$ dominate we expect alignment or anti-alignment. Either possibility guarantees flavor conservation as either one implies an approximate $U(1)^3$ symmetry. The former situation occurs if the gaugino mass is much bigger than the $A$-terms; and the latter in the opposite case.

b) A second possibility is that the orientation of the $A$-terms is given by the same matrices $U_A$ that occur in $m^2_A$, i.e.

$$A_e = U^\dagger L \bar{\Delta}_e L_E \quad A_u = U^\dagger Q \bar{\Delta}_u U_U \quad A_d = U^\dagger Q \bar{\Delta}_d U_D.$$

In this case the orientation of the $U_A$ is fixed by the dominant one loop effects of the previous section and one has to hope that the frozen parameters $\bar{\Delta}_a$ commute with the corresponding mass matrices $m_a$.

c) Finally, a third possibility is that the $A_a$ themselves are frozen parameters; this is identical to what happens in the minimal supersymmetric standard model and one has again to hope that the $A_a$ commute with the corresponding $\lambda_a$.

For the rest of the paper we shall assume that the $A$-terms commute with the corresponding Yukawa matrices and do not cause significant flavor violations.

### 4. Flavor Violating Processes

The first consequence of the results of the previous sections is that, as a result of alignment, all three lepton numbers are individually conserved. This is obviously not possible in the quark sector, since the up and down quarks themselves do not have parallel mass matrices.\footnote{Alignment as a solution of the flavor problem in supersymmetric theories was also considered, in a different context, in ref. [10].} The quark flavor violations are best discussed by going, via a superfield rota-
tion, to the quark Yukawa eigenbasis where both up and down masses are diagonal and the squark Yukawas have the form:

\[
\mathcal{M}_u^2 = \begin{pmatrix}
    (m_u^2 + S^\dagger S + D_{uL}) & \tilde{\Delta}_u + \frac{\mu}{\tan \beta} m_u \\
    \tilde{\Delta}_u + \frac{\mu}{\tan \beta} m_u & m_u^2 + \Sigma_Q + D_{uR}
\end{pmatrix},
\]

\[
\mathcal{M}_d^2 = \begin{pmatrix}
    (m_d^2 + K^\dagger S^\dagger S K + D_{dL}) & \tilde{\Delta}_d + \mu \tan \beta m_d \\
    \tilde{\Delta}_d + \mu \tan \beta m_d & m_d^2 + \Sigma_D + D_{dR}
\end{pmatrix}
\]

(4.1)

All flavor violation is contained in \( S \) and \( K \). The off-diagonal elements of \( S \) are much smaller than those of the Kobayashi-Maskawa matrix:

\[
S_{23} \simeq K_{cb} \frac{m_b^2}{m_t^2} \sim 2 \times 10^{-5}
\]

\[
S_{13} \simeq K_{ub} \frac{m_b^2}{m_t^2} \sim 2 \times 10^{-6}
\]

\[
S_{12} \simeq \frac{|K_{us}K_{cs}m_s^2 + K_{ub}K_{cb}m_b^2|}{m_c^2} \sim 5 \times 10^{-3}
\]

(4.2)

and therefore they do not significantly affect FCNC processes, although they may contribute to CP-violating processes. Then, in the approximation \( S = 1 \), all new flavor violations occur in the \( D_L \) sector, as can be seen from the squark mass matrices in Eq. (4.1).

The most stringent constraint comes from the contribution of squark-gluino loops to the real part of the \( K^0 - \bar{K}^0 \) mixing:

\[
\frac{\Delta m_K}{m_K} \simeq \frac{\alpha_s^2}{54} \frac{B_K}{M_{\tilde{g}}^2} \Re(X)
\]

(4.3)

\[
X \equiv \sum_{i,j} K_{is} K_{id}^* K_{js} K_{jd}^* f\left(\frac{m_{Q_i}^2}{M_{\tilde{g}}^2}, \frac{m_{Q_j}^2}{M_{\tilde{g}}^2}\right)
\]

(4.4)

\[
f(x, y) \equiv \frac{1}{x - y} \left[\frac{(11x + 4)x}{(x - 1)^2} \log x - \frac{15}{x} - (x \to y)\right],
\]

(4.5)

where \( f_K = 165 \text{ MeV} \) is the kaon decay constant, \( B_K \) parametrizes the hadronic matrix element, and \( M_{\tilde{g}} \) is the gluino mass. Assuming \( M_{\tilde{g}}^2 = m_Q^2 \) and keeping the leading contribution in the squark mass splitting, one finds

\[
\Re(X) = \frac{\sin^2 \theta_c}{6} D_{21}^2,
\]

(4.6)

where \( \theta_c \) is the Cabibbo angle and

\[
D_{ij} \equiv \frac{m_{Q_i}^2 - m_{Q_j}^2}{m_{Q_i}^2}.
\]

(4.7)
If we require that the gluino contribution in Eq. (4.3) does not exceed the experimental value of $\Delta m_K/m_K$, we obtain the constraint:

$$D_{21} < 0.1 \frac{m_Q}{300 \text{ GeV}}.$$  \hspace{1cm} (4.8)

The squark-gluino contribution to the imaginary part of $K^0 - \bar{K}^0$ mixing is given by:

$$\langle |\epsilon| \rangle_{\tilde{g}} = \frac{m_K}{\Delta m_K} \frac{f_K^2 B_K}{108 \sqrt{2}} \frac{\alpha_s^2}{M_{\tilde{g}}^2} \text{Im}(X)$$  \hspace{1cm} (4.9)

With the same approximation used before, we obtain

$$\text{Im}(X) = \frac{1}{3} |K_{ua}| |K_{ub}| |K_{cb}| \sin \delta D_{32} D_{21},$$  \hspace{1cm} (4.10)

where $\delta$ is the CP-violating phase in the Kobayashi–Maskawa matrix. This does not exceed the experimental value for $|\epsilon|$ if

$$\sqrt{D_{21} D_{31}} < \frac{m_Q}{300 \text{ GeV}}.$$  \hspace{1cm} (4.11)

There is no significant constraint coming from $B^0 - \bar{B}^0$ mixing and, in the limit $S = 1$, there is no new gluino-mediated contribution to $D^0 - \bar{D}^0$ mixing.

The constraints from FCNC on our model are much milder than those on a general supersymmetric $SU(3) \times SU(2) \times U(1)$ theory with minimal particle content and non-universal frozen soft-terms [11]. The reason is that in our theory, just as in the standard model, flavor violations are proportional to the Kobayashi-Maskawa angles; however, they are also suppressed by the large sparticle masses. Therefore, our contributions to rare processes can compete with the standard model contributions only if the latter have light quark suppressions, as in $\Delta m_K/m_K$ where $(\Delta m_K/m_K)_{SM} \sim G_F m^2$.

It is noteworthy that we do not obtain any constraints from either $\mu \to e\gamma$ or $\epsilon$. These provide by far the strongest constraints on general supersymmetric models. In our case, $\mu \to e\gamma$ vanishes whereas $\epsilon$ is small because it is proportional to the Jarlskog invariant $J$ of the standard model and is further suppressed by sparticle masses. The only significant constraint we have is from $\Delta m_K$ Eq. (4.8). It can be accounted for in several ways. One is by invoking heavy gluinos, which cause the squark masses to approach one another in the infrared. Furthermore, in Sect. 5, we will show how the dynamics of the moduli can adjust to render the squarks of the two heavy generations degenerate.

We end with a cosmological caveat. Because the moduli couple very weakly with strength $\sim M^{-1}_{\text{Pl}}$, they do not efficiently lose energy. As a result, they do not reach their minima in simple cosmologies [12], unless they happen to accidentally start out near their vacuum. Recently, there have been a revival of suggestions [13] on how to solve the problem and to allow the moduli to cosmologically relax to their ground state. Such a mechanism is clearly necessary to ensure flavor alignment. Even more, it is necessary to ensure that the Universe is not overclosed by coherent oscillations of the moduli.
5. Plastic Soft Terms

In previous sections we have conjectured that the potential \( V_\Lambda \) at the scale \( \Lambda \) where SUSY is broken leaves the orientation of the soft terms undetermined, but fixes their eigenvalues. In this section we wish to relax the latter hypothesis. We envisage that the supersymmetry-breaking dynamics at \( \Lambda \) provide the low-energy theory with a constraint which fixes the overall scale \( m_s \) but does not necessarily freeze all three eigenvalues. Some functions of the eigenvalues can correspond to flat directions which remain undetermined until we turn on the Yukawa couplings. Of course, our postulate that the supersymmetry-breaking mechanism respects the flavor symmetry requires that the constraints that fix \( m_s \) have to be flavor singlets.

Let us consider the case of vanishing left-right mixings in the squark and slepton mass matrices and focus on the fields \( \Sigma \) defined in eq. (2.1). (For this section we drop the subscript \( A \).) Suppose that the dynamics at the scale \( \Lambda \) fixes the two lowest-dimension flavor-singlet operators:

\[
\text{Tr} \Sigma = T, \quad \text{Tr} \Sigma^2 = T^2,
\]

where \( T \) and \( T^2 \) are numbers of order \( m_s^4 \).

These are two constraints on three eigenvalues, thus one combination of eigenvalues remains a flat direction whose VEV will be determined by low-energy physics in a calculable way. It is easy to identify the flat direction. The above constraints are not just \( SU(3) \) invariant, but are \( SO(8) \) invariant, and they force the spontaneous breakdown \( SO(8) \rightarrow SO(7) \), giving rise to seven Goldstone bosons. Six of them are a consequence of the breaking \( SU(3) \rightarrow U(1)^2 \) and can be identified with the fields \( \sigma \). The seventh is the new flat direction \( \theta \) which allows the eigenvalues of \( \Sigma \) to slide along a valley which preserves the above constraints.

The field \( \Sigma \) satisfying Eq. (5.1) can be expressed as

\[
\Sigma = TU^\dagger \left[ \frac{1}{3} - x(\cos \theta \lambda_8 + \sin \theta \lambda_3) \right] U,
\]

where \( \lambda_{3,8} \) are the two diagonal Gell-Mann matrices, \( U \) denotes an \( SU(3)/U(1)^2 \) rotation, and

\[
x \equiv \sqrt{\frac{3T^2 - T^2}{6T^2}}, \quad 0 \leq x \leq \frac{1}{\sqrt{3}}.
\]

Our assumption is that the six parameters contained in \( U \) and the angle \( \theta \) are dynamical variables, related to flat directions of the moduli fields. The soft term \( \Sigma \) is not only “disoriented” in flavor space, but is also “plastic”, since the pattern of eigenvalues can be deformed. Plasticity is disorientation in \( SO(8) \) space. In contrast to \( SU(3) \), \( SO(8) \) allows rotations in the \( \lambda_3 - \lambda_8 \) plane.

The effective potential for \( \Sigma \) is given by eqs. (2.10), (2.20), and its minimum occurs for \( \cos \theta \simeq 1 \). This implies that the vacuum has an approximate \( SU(2) \times U(1) \) symmetry which insures the degeneracy of the soft masses of same charge sparticles belonging to the two lightest generations. Consequently, \( m^2_{Q_1} \) and \( m^2_{Q_2} \) are approximately equal, and this
provides for the desired suppression of the real part of $K_0 - \bar{K}_0$ mixing. Plasticity can be extended to the $A$-terms and, as discussed in section 3, they can align or anti-align with the Yukawas, depending on whether the gaugino mass is much larger than the $A$-terms or vice versa.

6. Minimal Unification

Until now we have been working under the hypothesis that below the scale $\Lambda$, where the supersymmetry breakdown occurs, we have the minimal supersymmetric $SU(3) \times SU(2) \times U(1)$ particle content. We now consider the possibility that the theory below $\Lambda$ is some minimal supersymmetric GUT.

In minimal supersymmetric GUTs the gauge symmetry is increased to $SU(5)$ or $SO(10)$ and the number of chiral multiplets decreases. This means that the flavor group is no longer $U(3)^5$, but it is smaller: $U(3)_5 \times U(3)_10$ in the case of $SU(5)$, and just $U(3)_16$ for $SO(10)$. If we also assume that the soft terms are as minimal as possible, namely singlets under the GUT group, then we have a very constrained system with a small flavor group and a small number of parameters in the soft terms. Are there enough moduli/Goldstones available to align sufficiently and avoid problems with flavor violations?

For simplicity, let us discuss the minimal $SO(10)$ model in which the Yukawa coupling superpotential between the ordinary fermions in the $16$ representation and the Higgs fields $H_{u,d}$ is

$$W_Y = 16\lambda_u 16H_u + 16K\lambda_d K^T 16H_d.$$  \hspace{1cm} (6.1)

For simplicity, we will ignore the $A$ trilinear terms and write the soft supersymmetry-breaking Lagrangian as:

$$\mathcal{L}_{\text{Soft}} = \frac{m_s^2}{F} 16^\dagger U^\dagger \Sigma U 16.$$ \hspace{1cm} (6.2)

The crucial difference between this minimal-GUT case, with gauge-singlet $\Sigma$, and the previous $SU(3) \times SU(2) \times U(1)$ analysis is apparent from Eq. (6.2). Now there is just one $U$ available, instead of 5, to do all the alignments necessary to reduce flavor violations. It is clear that $U$ will align $\Sigma$ parallel to $m_u$, since $m_u$ gives the largest contribution to the energy. This implies that all sparticle mass matrices will be parallel to $m_u$, whereas the down-quark and charged-lepton mass matrices will be misaligned from $m_u$ by angles of the order of the Kobayashi-Maskawa angles.

Thus, unless sleptons are highly degenerate in mass, $\mu_{L,R} \rightarrow e_{R,L}\gamma$ transitions are proportional to a mixing angle $K_{e\mu} = K_{us} \simeq \sin \theta_c$ and occur at an unacceptable rate. In $SU(5)$ only the right-handed sleptons are misaligned from the lepton mass matrix, and the amplitude for $\mu_L \rightarrow e_R + \gamma$ is again proportional to the Cabibbo angle $\sin \theta_c \simeq \sqrt{d/s}$. Of course, minimal $SO(10)$ and $SU(5)$ theories have a problem: they predict $m_d = m_e$ and this is the reason why they give $\mu \rightarrow e\gamma$ proportional to $\sqrt{d/s}$. However even if we extend the theory à la Georgi–Jarlskog, the $\mu \rightarrow e\gamma$ amplitude is still problematic, being proportional to $\sqrt{e/\mu}$. 

11
The reason for this failure is that in minimal supersymmetric GUTs with minimal GUT-invariant soft terms, the few available soft terms just align with \( m_u \), leaving some mismatch between down quarks and squarks and more importantly between leptons and sleptons. This causes difficulties with individual lepton violating processes, which were not originally present in supersymmetric GUTs with universality at \( M_{GUT} \).

The problem could be cured in more complicated GUTs with a larger flavor structure, necessary perhaps to explain the fermion mass pattern, which would allow for more freedom in the low-energy alignment of the soft-breaking masses.

A strong degeneracy between the first two generations of sleptons and down squarks suppresses the most dangerous processes and could therefore represent an alternative solution. In the previous section we have shown that this occurs in the plastic soft-term scenario and the degeneracy of the sparticles of the first two generations is predicted. The dynamics of the plastic soft terms cures the disease in the dynamics of the disoriented soft terms: in GUTs the large up-type quark Yukawa couplings force the sleptons to misalign, but insure that the first two generations are almost degenerate in mass. The decay \( \mu \to e\gamma \) can still occur via virtual \( \tilde{\tau} \) exchange, and its rate is just below the present experimental limit. Interesting effects in lepton-number violating \( \tau \)-decays can be present. The plastic GUT scenario allows therefore the construction of phenomenologically viable models which are predictive and represent possible alternatives to the minimal supersymmetric standard model with universal boundary conditions at \( M_{GUT} \).

7. Conclusions

We proposed “disorientation” as an alternative to universality for suppressing flavor violation in supersymmetric theories. Universal soft terms realize the flavor symmetry in the Wigner mode. Disoriented soft terms realize it in the Nambu-Goldstone mode; this allows large sparticle splittings and has the appeal that the absence of flavor violations is a consequence of a dynamical calculation.

The Goldstone particles can be thought of as either the consequence of a spontaneously broken flavor symmetry or perhaps could be identified with some of the flat directions (moduli) that frequently occur in supersymmetric or superstring theories. In the latter case there would be an important connection between the space of the moduli and the flavor group.

Why did our mechanism work? Promoting some of the parameters of the low-energy theory to fields allowed us to exploit nature’s preference for states of maximal possible symmetry. This is the reason why: the spin aligns with an external magnetic field, preserving \( SO(2) \); sleptons align with leptons, preserving individual lepton number conservation \( U(1)^3 \); squarks align –as much as possible– with the quarks, preserving an approximate \( U(1)^3 \); the 7th goldstone boson of the plastic scenario chooses to relax at its special value where the symmetry is enhanced to \( SU(2) \times U(1) \) and pairs of sparticles are degenerate. Nature’s frequent preference for states of higher symmetry fully accounts for our mechanism for the suppression of flavor violation. More importantly, it leads us to new
supersymmetric phenomenology and the peaceful coexistence of split sparticles and flavor conservation.

We have found that the dynamics of alignment occurs at large scales and is sensitive to details of Planckian physics. In light of this, is disorientation better than universality? Both are strong hypotheses which rely on the existence of an approximate symmetry in a sub-sector of the full theory. Which is better can only be decided in the context of a complete theory which addresses the full flavor problem and explains fermion masses. Only then can we see how the soft terms avoid being directly infested with large flavor violations from Planckian physics.

Acknowledgements:

It is a pleasure to thank L. Hall, C. Kounnas, G. Veneziano and F. Zwirner for valuable discussions. S.D. thanks the Department of Theoretical Physics of the University of Oxford for hospitality during the course of the work.
Appendix A: The evolution equation for $V_k$

We are interested in the effect of fluctuations on the shape of the potential $V_k(\sigma_A)$ as the scale $k$ is lowered from $k = \Lambda$ to $k = m_s$. For this reason we introduce an effective infrared cutoff term $R_k(q^2)$ in the momentum integrations appearing in the loop contributions to the potential $V_k(\sigma_A)$. This term prevents the integration of modes with momenta $q^2 \lesssim k^2$. The effective potential at one loop is now given by

$$V_k(\sigma_A) = V_\Lambda(\sigma_A) + \frac{1}{2} \int_\Lambda d^4 q (2\pi)^4 \text{Str} \log \left[q^2 + R_k(q^2) + M^2(\sigma_A)\right]. \tag{A.1}$$

In the formulation by C. Wetterich the cutoff term is chosen as

$$R_k(q^2) = \frac{Z_k q^2 f_k^2(q^2)}{1 - f_k^2(q^2)}. \tag{A.2}$$

The function

$$f_k^2(x) = \exp\left\{-2a \left(\frac{q^2}{k^2}\right)^b\right\} \tag{A.3}$$

can be used for the implementation of a sharp or smooth cutoff through an appropriate choice of the two free parameters $a, b$. $Z_k$ is a $k$-dependent matrix in field space whose precise definition is given in the following. An ultraviolet cutoff $\Lambda$ is assumed for the momentum integration. Notice that for $k = \Lambda$ the one loop contribution automatically vanishes. In the limit $k \to 0$ the cutoff term $R_k(q^2)$ is removed and the integration reproduces the standard one loop result for the effective potential without a cutoff. Taking the partial derivative of $V_k$ with respect to $t = \log(k/\Lambda)$ results in the evolution equation

$$\frac{\partial V_k(\sigma_A)}{\partial t} = \frac{1}{2} \int_\Lambda d^4 q (2\pi)^4 \text{Str} \frac{\partial R_k}{\partial t} \left[q^2 + R_k(q^2) + M^2(\sigma_A)\right]^{-1}. \tag{A.4}$$

The momentum integration is infrared and ultraviolet finite as the integrand deviates significantly from zero only for $q^2 \simeq k^2$. The renormalization group improvement consists in substituting the running mass matrix $M^2(\sigma_A, k)$ for the classical one, and multiplying $q^2$ by the wavefunction renormalization of the various fields $Z_k$. This takes into account the fact that the change in the effective potential when fluctuations with momenta $q^2 \simeq k^2$ are incorporated in it involves the full propagator of the theory at the scale $k$. We can now identify the matrix $Z_k$ appearing in the definition of eq. (A.2) with the wavefunction renormalization. Notice that the $t$-derivative of $R_k$ includes a contribution proportional to the anomalous dimension of the fields $\eta = -\partial(\log Z_k)/\partial t$. It can be checked that the explicit $Z_k$-dependence can be incorporated in the definition of the renormalized mass matrix $\tilde{M}^2(\sigma_A, k) = Z_k^{-1} M^2(\sigma_A, k)$. The integral in the r.h.s. of eq. (A.4) cannot be easily computed for general values of the parameters $a, b$ appearing in eq. (A.3).

However, in the limit of a sharp cutoff $b \to \infty$ the momentum integration can be carried out explicitly. Moreover, the contribution proportional to $\eta$ can be neglected, as it is suppressed by $1/b$. As a result, the effect of the wavefunction renormalization is completely
absorbed in the running of the renormalized mass matrix $\tilde{M}(\sigma_A, k)$. The final expression for the running of the potential is

$$\frac{\partial V_k(\sigma_A)}{\partial t} = -\frac{k^4}{16\pi^2} \text{Str} \log \left[ 1 + \frac{\tilde{M}^2(\sigma_A, k)}{k^2} \right].$$  (A.5)

A few remarks are due in order to clarify some steps in our derivation:

1) Even though we were led to eq. (A.5) through an intuitive way a more formal derivation is possible [6]. The $k$-dependent effective action $\Gamma_k$ for scalar fields can be obtained from the partition function through the usual Legendre transformation, if the infrared cutoff term of eq. (A.2) is added to the classical action so that low momentum modes do not propagate. An exact renormalization group equation describes the evolution of $\Gamma_k$ with $k$ [6]. This equation leads to eq. (A.5) for the potential. For fermions the discussion proceeds along parallel lines. A modified fermion propagator is used, so that the momentum integrations are cut off in the infrared [7]. For the discussion of supersymmetric theories the choice of cutoffs terms for scalars and fermions must preserve the supersymmetry at all scales. This is accomplished in the limit $b \to \infty$ that we have considered [7].

2) The matrix $Z_k$ appearing in eq. (A.2) includes the wavefunction renormalization for scalar and fermion fields. We have implicitly assumed that the fermionic part of the matrix involves the square of the term which renormalizes the fermion field. This is apparent from the way the renormalized masses $\tilde{M}$ are defined.
References

[1] S. Dimopoulos and H. Georgi, ”Supersymmetric GUTs”, p. 285, Second Workshop on Grand Unification, University of Michigan, Ann Arbor, April 24-26, 1981, eds. J. Leveille, L. Sulak, D. Unger; Birkhauser, 1981;
S. Dimopoulos and H. Georgi, *Nucl. Phys. B* **193** (1981) 150.

[2] L. Girardello and M.T. Grisaru, *Nucl. Phys. B* **194** (1982) 65.

[3] C. Kounnas, F. Zwirner, and I. Pavel, *Phys. Lett. B* **335** (1994) 403;
P. Binetruy and E. Dudas, *Phys. Lett. B* **338** (1994) 23 and preprint LPTHE Orsay 94/73, SPhT Saclay T94/145;
C. Kounnas, I. Pavel, G. Ridolfi, and F. Zwirner, preprint CERN-TH/95-11.

[4] Y. Nambu, preprint EFI 92-37;
V.I. Zakharov, in ”Properties of SUSY particles” (L. Cifarelli and V.A. Khoze eds., World Scientific, Singapore, 1993) p.65.

[5] For a review see K. Wilson and I. Kogut, Phys. Rep. **12**, 75 (1974).

[6] C. Wetterich, Phys. Lett. **B 301**, 90 (1993).

[7] C. Wetterich, Z. Phys. **C 48**, 693 (1990).

[8] I. Jack and D. Jones, Phys. Lett. **B 333**, 372 (1994);
S. Martin and M. Vaughn, Phys. Rev. **D 50**, 2282 (1994);
Y. Yamada, Phys. Rev. **D 50**, 3537 (1994);
I. Jack, D. Jones, S. Martin, M. Vaughn and Y. Yamada, Phys. Rev. **D 50**, 5481 (1994).

[9] J. Bagger, E. Poppitz, and L. Randall, preprint EFI-95-21 (1995);
H.-C. Cheng and N. Arkani-Hamed, unpublished.

[10] Y. Nir and N. Seiberg, *Phys. Lett. B* **309** (1993) 337.

[11] F. Gabbiani and A. Masiero, *Nucl. Phys. B* **322** (1989) 235;
J.S. Hagelin, S. Kelley and T. Tanaka, *Nucl. Phys. B* **415** (1994) 293.

[12] G. Coughlan, W. Fischler, E. Kolb, S. Raby and G. Ross, *Phys.Lett. B131* (1983) 59;
J. Ellis, D.V. Nanopoulos and M. Quirós, *Phys.Lett. B174* (1986) 176;
G. German and G.G. Ross, *Phys.Lett.*, **B172** (1986) 305;
O. Bertolami, *Phys.Lett B209* (1988) 277;
R. de Carlos, J.A. Casas, F. Quevedo and E. Roulet, *Phys.Lett. B318* (1993) 447;
T. Banks, D. Kaplan and A. Nelson *Phys.Rev. D49* (1994) 779;
T. Banks, M. Berkooz and P.J. Steinhardt, preprint RU-94-92.

[13] S. Dimopoulos and L. J. Hall, *Phys. Rev. Lett.* **60** (1988) 1899;
L. Randall and S. Thomas, MIT preprint, LMU-TPW-94-17;
G. Dvali, preprint IFUP-TH 09/95;
M. Dine, L. Randall and S. Thomas,preprint SLAC-PUB-95-6776;
T. Damour and A. Vilenkin, CNRS preprint.