Gravitational divergences as a mediator of supersymmetry breaking

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Abstract

Gravitational divergences associated with singlet fields in supersymmetric theories are reexamined, and their possible contributions to the low-energy effective theory are pointed out. We demonstrate that such divergences are not necessarily harmful and that Planck-scale physics could play an important role in models of low-energy supersymmetry breaking via a radiatively induced tadpole term in the scalar potential. In this case, gravitational divergences play the role of the supersymmetry breaking mediator, leading to a simple realization of the so-called messenger model. We also point out a new mechanism for the generation of mass terms for the Higgs fields in models of low-energy supersymmetry breaking, as well as a horizontal messenger mechanism in which the horizontally charged scalars are rendered heavy. Implications to the flavor problem in supersymmetric models are also discussed.
Softly broken $N = 1$ supersymmetric theories are only logarithmically divergent and can naturally accommodate weak-scale scalar fields. In particular, supersymmetry offers a well-defined framework for embedding and for extrapolating to high energies the Standard Model of electroweak and strong interactions (SM), which contains a Higgs boson.

Supersymmetric extensions of the standard model postulate an equal number of fermionic and bosonic degrees of freedom in the low-energy theory, leading to the exact cancelation of quadratic divergences (which are not gravitational) in both supersymmetric and softly broken theories. The minimal supersymmetric extension of the standard model (MSSM) essentially compliments each SM fermion (boson) with a boson (fermion) superpartner, and a plethora of new soft mass (and possibly mixing and phase) parameters are needed to describe the superpartner potential. It must also contain a pair of SU(2) Higgs doublets with an opposite hypercharge (due to gauge anomalies involving the Higgs fermions), as well as the corresponding supersymmetry conserving mass term, $\mu H_1 H_2$, which renders the Higgs fermions sufficiently massive.

The MSSM provides a well-motivated and a well-defined framework for the study and search of physics beyond the SM and has attracted great attention in recent years. Nevertheless, the MSSM, or any other supersymmetric extension of the SM, is most probably only the low-energy effective approximation of a more fundamental theory which may explain the origin of its many parameters. Since it contains only logarithmic divergences, it is rather insensitive to the details of the high-energy theory. This situation is a consequence of supersymmetry, as explained above, but also of the fact that the MSSM does not contain a field which is a singlet under all global and local symmetries (i.e., a universal singlet).

From the model-building point of view, singlet fields — whether fundamental or composite objects — are a useful tool because of their trivial transformation properties, and can be easily incorporated into extended frameworks. Their introduction, however, reintroduces quadratic divergences to the low-energy theory, i.e., quadratically divergent terms associated with singlet fields appear in the softly broken theory $[1-8]$. The divergences are induced by the gravitational interactions of the singlet. (Similar divergences also appear in the vacuum energy $[9,7]$.) The additional divergences render the singlet sensitive to the physics at very high energies. This is similar to the situation with a fundamental Higgs boson field in the SM (in the absence of supersymmetry).

The singlet could become heavy, in which case, it either decouples and is irrelevant and harmless or it destabilizes the MSSM parameters (if it is coupled to other low-energy fields and renders those fields heavy). Of course, in the latter case one has to require the absence of such a harmful singlet. On the other hand, a universal singlet may be safely incorporated into the low-energy theory in certain situations. Furthermore, it may even be the agent that parameterizes the high-energy physics in its low-energy effective approximation. In this paper we will attempt to examine this possibility. We will investigate whether such a situation exists and if so, its implications for the low-energy effective theory. We will do so by examining the divergent scalar potential of the universal singlet focusing on the possible role of the quadratically divergent terms in the low-energy theory. In particular, we will reexamine the relations between the potential and the scale of supersymmetry breaking, $M_{SUSY}$.

Supersymmetry is broken in some sector of the theory at a scale $M_{SUSY}^2 \equiv \sqrt{3} m_{3/2} M_P$, where $m_{3/2}$ and $M_P = M_{\text{Planck}}/\sqrt{8\pi}$ are the gravitino and the reduced Planck mass, re-
spectively, and we impose hereafter the condition of a vanishing cosmological constant. If
supersymmetry breaking, parameterized here by $m_{3/2}$, is mediated to the SM (observable)
sector by the gravitational interactions of the SM fields, then the gravitino mass also pa-
parameterizes the (soft) mass scale of the superpartners of the ordinary fermions and of the
gauge and Higgs bosons (the sparticles). It leads to the constraint $m_{3/2} \simeq m_{\text{sparticle}} \simeq m_{\text{weak}}$,
which fixes the scale of the gravitino mass. The one-loop divergences of the cosmological
constant introduce a $\sim 100\%$ ambiguity in $m_{3/2}$ and in related low-energy parameters [9],
but do not affect the qualitative relation between the scales. If the only interactions between
the supersymmetry breaking sector and the observable sector are gravitational, as one often
assumes, then supersymmetry breaking is said to have taken place in a hidden sector.
Such models are sometimes classified as models of gravitationally mediated supersymmetry
breaking. It is impossible to realize in a satisfactory fashion a light universal singlet in
the above framework. Radiatively induced quadratic divergences, which are proportional to
$m_{3/2}^2$, destabilize the singlet field [3–4]. It slides at the quantum level to large field values
and decouples from the low-energy theory. As explained above, the singlet could be either
irrelevant or harmful in this case. The decoupling of the singlet is, however, a result of the
ad hoc relation $m_{3/2} \simeq m_{\text{weak}}$, and it is not a generic property of the theory.

Below, we will examine possible frameworks in which the constraint $m_{3/2} \simeq m_{\text{weak}}$ is
relaxed, i.e., we will require only $m_{3/2} \lesssim m_{\text{weak}}$. Indeed, we find that the divergent singlet
might not decouple from the low-energy theory. Furthermore, it may provide an important
model-building tool. The relation between the value of the singlet field and the supersymme-
try breaking scale, which are correlated by $m_{3/2}$, allows for new mechanisms of gravitational
mediation of supersymmetry breaking in which $m_{\text{weak}}$ is a complicated function of $m_{3/2}$. The
superpartners of the ordinary fermion and of the gauge and Higgs bosons are rendered mas-
vive by, e.g., gauge loops rather than by their gravitational interactions. The gravitational
interactions, which affect only the singlet field directly, are only a trigger for the genera-
tion of the MSSM soft supersymmetry breaking mass parameters. The presence of a light
universal singlet field in the theory may be difficult to understand from the effective field
theory point of view. By definition, the theory does not contain any local or global sym-
metry which would prevent $O(M_P)$ tree-level mass terms for the singlet. However, string
theory often contradicts this somewhat naive argument, and could contain in its low-energy
limit massless states that cannot be understood from the low-energy point of view. In fact,
such states may provide a unique low-energy signature of string theory. Note also that the
singlet may not be a fundamental object. This would be the case, for example, if it is a
bilinear composite given by duality transformations, in which case its mass could correspond
to a nonrenormalizable operator in the more fundamental strongly coupled theory [10] and
could be highly suppressed[1]. Here, we will simply assume that the theory contains a light
universal singlet, $S$, and examine its consequences.

The radiative emergence of dimensionful parameters for the singlet field in the low-energy

\footnote{However, in this case one has to forbid mass terms in the strongly coupled theory which would
correspond to a tree-level tadpole term in the weakly coupled theory. If the strongly coupled
theory contains also fundamental singlets, fundamental-composite mixing could arise from the
duality transformations of Yukawa operators.}
effective theory is forbidden by supersymmetry nonrenormalization theorems. Nevertheless, terms proportional to the supersymmetry breaking parameter $m_{3/2}$ are not forbidden by any mechanism and could still appear at low-energy. Their gravitational origin is imprinted in their proportionality to the gravitino mass. Indeed, it has been shown that gravitationally induced tadpole terms for the singlet superfield typically appear at one \cite{3-6} or two \cite{7} loop order. It is instructive to construct the divergent potential in a simple example, which can then be generalized. The tadpole terms could arise, e.g., due to Planck mass suppressed trilinear terms in the Kahler potential,

$$K = \sum_I \left[ 1 + \frac{\alpha_I}{M_P} (S + S^\dagger) \right] \Phi_I \Phi_I^\dagger + \ldots,$$

(1)

where the dots denote terms suppressed by higher orders of the Planck mass and include, in principle, mixing between hidden and observable fields (possibly leading to a non-trivial Kahler curvature). The summation is over the chiral superfields $\Phi_I$, $S$ is a singlet chiral superfield and $\alpha_I$ are $O(1)$ dimensionless couplings. The leading nonrenormalizable operators given in eq. (1) are allowed only if $S$ is a gauge and global symmetry (universal) singlet, which we assume hereafter.

The low-energy Lagrangian contains contributions from the Kahler potential $K$, the so-called $D$-terms,

$$\mathcal{L}_D = \int d^2 \theta d^2 \bar{\theta} E K,$$

(2)

where $\theta$ is the Grassman variable and $E$ is the superspace density $E = e^{K/M_P^2}$. In particular, eq. (2) includes the trilinear vertices $\alpha_I E S \Phi_I \Phi_I^\dagger / M_P$. A tadpole diagram, which is allowed by all symmetries, appears at one loop,

$$\mathcal{L}_D \sim \frac{N}{16\pi^2} \frac{\Lambda^2}{M_P^2} \int d^2 \theta d^2 \bar{\theta} E (S + S^\dagger),$$

(3)

where we set, for simplicity, all dimensionless couplings to unity, $N$ counts the number of massless fields $Q_I \in \Phi_I$, and $\Lambda$ is the momentum cutoff scale for the divergent loop diagram.

The superspace density can be expanded \cite{4}

$$e^{K/M_P^2} = 1 + \frac{1}{M_P^2} \left[ \theta^2 K_i F^i + \bar{\theta}^2 K_i^* F^i + \theta^2 \bar{\theta}^2 (K_{ij} + \frac{K_i K_j^*}{M_P^2}) F^i F^j + \ldots \right],$$

(4)

where $K_i$ ($K_{ij}$) denotes derivatives with respect to supersymmetry breaking field $Z^I = z^I + \theta^2 F^i \in \Phi_I$ ($Z^{I*}$) and $F^i \sim M_{SUSY}^2$. Substituting (4) into (3), taking $\Lambda \sim M_P$, as is appropriate in the case of gravitational interactions (2), and using $F^i \sim M_{SUSY}^2 \sim m_{3/2} M_P$, one derives the contribution to the low-energy potential

$$|\Delta V| \sim \frac{N}{16\pi^2} \frac{1}{M_P} \left\{ K_{i\alpha} m_{3/2} M_P F_s + (K_{ij} + \frac{K_i K_j^*}{M_P^2}) m_{3/2}^2 M_P^2 s + h.c. \right\},$$

(5)

where $S = s + \theta^2 F_s$.

One could distinguish two obvious limits: (1) Planckian values for supersymmetry breaking fields: $K_i \sim M_P$; (2) No Planckian values for supersymmetry breaking fields: $K_i \sim M_{SUSY} \ll M_P$. Taking the appropriate limits in eq. (3) one has in case (1)
\[ |\Delta V| \sim \frac{N}{16\pi^2} \left[ m_{3/2}^2 M_P s \pm m_{3/2} M_P F_s + h.c. \right]. \quad (6) \]

In case (2), only the tadpole term, \( m_{3/2}^2 M_P s \sim [M_{SUSY}^4/M_P] s \), is present. In either case, the low-energy potential could be written as \( \Delta V = M_P \sum_{n=1}^3 C_n m_{3/2}^n q_i^{3-n} \). In this form the mixing between the heavy and light sectors is seen more clearly. It is sufficiently suppressed if \( m_{3/2} \to 0 \), but otherwise it could destabilize the light scalars \( q_i \). Note also that each term in eq. \( (6) \) carries a degree of freedom associated with its phase. Our phase choice below will correspond to a non-trivial global minimum of the potential, \( i.e. \), we explicitly assume that the divergence is relevant to the low-energy theory. (This would be the case, for example, if the phase is determined dynamically.)

The elevation of the spurion derivation used above to a supergravity derivation implies new divergent diagrams due to (divergent) renormalization of the kinetic terms. The two sets of one-loop diagrams cancel \([6]\), unless the Kahler curvature is non trivial in the appropriate direction \([6,7]\). However, the cancelation is accidental and, in general, is limited to one loop \([7]\). More generally one has

\[ |\Delta V| \sim \frac{N}{(16\pi^2)^n} \left[ m_{3/2}^2 M_P s \pm m_{3/2} M_P F_s + h.c. \right], \quad (7) \]

where \( n \) denotes the loop order of the first non-vanishing divergent contribution to the potential, and similarly in case (2). We typically expect \( n \leq 2 \).

Indeed, if the scalar potential \( V \sim m_{3/2}^2 [s^2 - M_P s] \) then \( s \sim M_P \) decouples from the low-energy theory. (The singlet may be stabilized in intermediate scales if the scalar potential contains a quartic term.) In the limit \( m_{3/2} \to 0 \) one can neglect any soft contributions \( \propto m_{3/2} \ll \sqrt{m_{3/2} M_P} \) to the low-energy potential other than eq. \( (7) \), provided that the superpotential is not trivial. The low-energy theory is given in this case by the effective superpotential and by the divergent corrections. Here, we will assume for simplicity the limit of case (2), \( i.e. \), \( K_i/M_P \to 0 \), and consider a simple-minded toy model. The toy model is given by the superpotential

\[ W = W(S) + \lambda S V \bar{V}, \quad (8) \]

where \( V = v + \theta^2 F_v \) and \( \bar{V} \) are a vector-like pair of charged chiral superfields, and we choose

\[ W(S) = \frac{\kappa}{3} S^3. \quad (9) \]

The scalar potential for \( s \) is given by eqs. \( (7) \) and \( (8) \),

\[ V(s) = -\frac{M_{SUSY}^4}{M_P} s + \left| \frac{\partial W(S)}{\partial S} \right|^2 = -\frac{M_{SUSY}^4}{M_P} s + \kappa^2 s^4. \quad (10) \]

We omit hereafter the dimensionless factor \( N/(16\pi^2)^n \sim O(1 - 10^{-4}) \) in eq. \( (7) \). It is absorbed in \( M_{SUSY}^4 \) and corresponds to an order of magnitude ambiguity in \( M_{SUSY} \). In general, we expect modification of the quartic coupling at the quantum level due the field dependent masses of \( S, V \) and \( \bar{V} \) in the one-loop effective potential. This is relevant only if \( \kappa^2 \lesssim 1/64\pi^2 \). More typically, however, one expects that all dimensionless couplings are \( O(1) \).
In deriving eq. (10) we implicitly assumed that in the global minimum $v, \bar{v}$ and $F_v, F_{\bar{v}}$ vanish. This is indeed the case if the charge breaking direction $F^*_s = \lambda v \bar{v} + \kappa s^2 = 0$ is only a local minimum of the model. The corresponding constraint reads

$$2\lambda > \kappa,$$

which is easily satisfied and which we will assume.

It is now straightforward to minimize the potential (10),

$$s = \left( \frac{M_{SUSY}^4}{4\kappa^2 M_P} \right)^{\frac{1}{3}}.$$  \hspace{1cm} (12)

Note that the result for $s$ is only weakly sensitive to $\kappa$. The order of magnitude ambiguity in $M_{SUSY}$ corresponds to a similar ambiguity in $s$. It could be compensated, if desired, by an appropriate choice of $\kappa$. The inclusion of all $O(m_3/2)$ soft terms would shift $s$ by $O(m_3/2)$. We will verify that the shifts are negligible when discussing specific examples.

We also have $F_s/s = \kappa s$. The nonvanishing $S$ components generate masses for the vector-like scalar pair $(v, \bar{v})$,

$$M_{v\bar{v}}^2 \sim \left( \frac{\lambda^2 s^2 \lambda F_s}{\lambda F_s^* \lambda^2 s^2} \right).$$  \hspace{1cm} (13)

The diagonal term is a supersymmetric mass term, i.e., the corresponding fermions have a similar Dirac mass term $\lambda s \bar{v}$. Similarly, field dependent masses are induced for $s$ and for its fermion partner $\tilde{s}$, and are given by replacing $\lambda$ with $\kappa/3 (\kappa)$ in the diagonal (off-diagonal) terms in eq. (13).

In the remainder of this paper we will discuss the application of our toy model to various (phenomenological) schemes of supersymmetry breaking. The most obvious framework that merits examination is that of low-energy supersymmetry breaking: $M_{SUSY}^2 \ll m_{\text{weak}} M_P$ and $m_{3/2} \to 0$. Typically, one assumes that if supersymmetry is broken at low energy then it is broken in a secluded (rather than hidden) sector. The secluded sector communicates via new gauge interactions with a messenger sector, which, in turn, communicates via the ordinary gauge interactions with the observable sector. The new gauge and messenger Yukawa interactions mediate the supersymmetry breaking to a (SM) singlet messenger $N = n + \theta^2 F_n$, which parameterizes the supersymmetry breaking in the messenger sector. The singlet $N$ interacts also with SM non-singlet messenger fields $M$ and $\bar{M}$. The Yukawa interaction $\lambda N M \bar{M}$ communicates the supersymmetry breaking to the messengers $M$ and $\bar{M}$ as in eq. (13). In turn, the vector-like pair $M$ and $\bar{M}$, which transforms under the SM gauge group, communicates the supersymmetry breaking to the ordinary MSSM fields via gauge loops. The gauge loops commute with flavor and, thus, the spectrum is charge dependent but flavor diagonal, if one ensures that all other possible contributions to the soft spectrum are absent or are strongly suppressed. Such models [11] are often referred to as “gauge mediation of supersymmetry breaking” or messenger models. The sparticle spectrum, and hence, the weak scale, are given in this framework by $m_{\text{weak}} \sim m_{\text{sparticle}} \sim (\alpha_i/4\pi)(F_n/n)$, where $\alpha_i$ is the relevant gauge coupling at the scale $\Lambda_N = F_n/n \sim 10^5$ GeV. One also assumes a similar scale for the supersymmetry breaking in the secluded sector, $M_{SUSY} \sim \Lambda_n$ (but $M_{SUSY}$ could be one or two orders of magnitude higher).
The universal singlet $S$ could couple in this framework to either the messenger fields, $VV \rightarrow MM$, or to the MSSM Higgs doublets, $VV \rightarrow H_1H_2$. By properly assigning, e.g., Peccei-Quinn or $Z_3$ charges, one could forbid one of the coupling while allowing the other. Alternatively, one could assume that one of the Yukawa terms is forbidden by the selection rules in a string theory (see discussion above and Ref. [12]).

In the case $W = SH_1H_2 + \frac{c}{3}S^3$, the mass term (13) corresponds to the supersymmetry preserving (usually denoted by $\mu$) and breaking (usually denoted by $m_{12}^2$, $m_3^2$, or $B$) mass terms in the MSSM Higgs potential. Substituting $M_{SUSY} \sim 10^6$ GeV in eqs. (12) and (13) one has $\mu \sim m_{12} \sim m_{\text{weak}}$, as required. In comparison, if $\mu$ and $m_{12}$ are generated by (Yukawa) loops, which would be the most natural mechanism in this framework, they would typically be generated at the same loop order and one has $m_{12}/\mu \sim \Lambda_n$, which is phenomenologically unacceptable (but see Ref. [13]). The universal singlet mechanism offers a simple resolution of this situation. (For other proposals, see Ref. [14].)

Alternatively, one could couple the universal singlet to the messenger fields, mediating supersymmetry breaking to the messenger sector via the gravitationally induced tadpole term in the scalar potential. Gravitational divergences play now the role of the supersymmetry breaking mediator, leading to a simple realization of the a messenger model. One need not introduce any new gauge interactions between the secluded and messenger sectors in order to mediate supersymmetry breaking to the messenger sector. The secluded sector is promoted to a truly hidden sector while the messenger sector, which now has a minimal content $\{S = N, V = M, \bar{V} = \bar{M}\}$, can be embedded in the observable sector. This significantly simplifies the theory and avoids potentially dangerous directions in field space [15]. (Such directions could appear if there are additional Yukawa interactions which are introduced in order to communicate the supersymmetry breaking from the subset of messengers which transform under the new gauge interactions to the singlet $N$.) Requiring $\Lambda_s = F_s/s > \sim 10^7$ GeV we find $M_{SUSY} \sim 10^8$ GeV. The gravitino mass is in the MeV range, and all terms proportional to $m_{3/2}^2$ (aside from the tadpole) can be safely neglected. We did not need to make any assumptions regarding the nature of supersymmetry breaking in the hidden sector, and we did not alter the generic phenomenology of the models. The embedding of only the messenger fields in the SM sector offers a simple and essentially model-independent framework in comparison to interesting but complicated proposals made recently for embedding both the SM and messenger sectors in the supersymmetry breaking sector [17].

If one raises $M_{SUSY}$, there is a possible interplay between $m_{3/2}$ and $(\alpha/4\pi)\Lambda_s$ mass terms, which allows us to consider a new class of messenger models. The messenger fields in this framework, $T$ and $\bar{T}$, are charged only under a gauged horizontal symmetry and are SM singlets. For concreteness, we will consider an SU(2) horizontal symmetry under which the third family and Higgs fields are singlets. The first and second family fields transform as doublets ($U_1$, $U_2$, etc., under the symmetry. For $M_{SUSY} \sim 10^{10-11}$ GeV one has $\Lambda_s = F_s/s \gtrsim 10^7$ GeV, leading to multi-TeV masses for all horizontally charged MSSM

\footnote{In the limit (1), i.e., $K_i/M_P \sim 1$, we find $M_{SUSY} \sim 10^5$ GeV.}

\footnote{The MeV gravitino may significantly constrain the cosmology [16], however, we do not discuss the cosmological implications in this paper.}
TABLE I. Mass scales in the horizontal messenger model.

| Particle                      | Mass scale |
|-------------------------------|------------|
| First and second family sfermions | $\frac{\Lambda}{4\pi}$ |
| Third family sfermions        | $m_{3/2}$  |
| Higgs fields                  | $m_{3/2}$  |
| SM gauginos                   | $m_{3/2}$  |
| $S, T, \bar{T}$               | $\Lambda_s$|
| Horizontal gauge fields       | $\lesssim \Lambda_s$|

scalars (assuming that the horizontal gauge coupling $\alpha_H \lesssim \alpha_{\text{weak}}$). As a result, generic MSSM contributions to flavor changing neutral currents involving the first and second family fermions (i.e., the flavor problem) are sufficiently suppressed, and there is no need to impose mass universality (as in the previous case) or fermion-sfermion mass alignment. The scalar components of the Higgs and third family fields are singlets of the symmetry and are lighter with only gravitational masses $\sim m_{3/2} \sim m_{\text{weak}}$. Thus, one avoids excessive fine-tuning in the Higgs potential (which radiatively couples to the third family fields through their Yukawa interactions). The spectrum is described in Table I. Although $m_{3/2}$ is not negligible in this case, the corresponding shift in the intermediate-scale $s$ is negligible. The horizontal messenger model offers a simple realization of a model with heavy first two families of sfermions [18], which is motivated by the resulting decoupling of the superpartners from sensitive low-energy flavor (and CP) observables.

All mechanisms described above lead to simple realizations of theoretical frameworks which attempt to solve the flavor problem in supersymmetric models by either low-energy sfermion mass universality (gauge-mediation framework) or decoupling (heavy sfermions - light Higgs framework). Indeed, a priori there is no reason to expect sfermion mass universality to hold. In the horizontal messenger framework the universality assumption can be significantly relaxed (as in the models of Ref. [18]). In the gauge-mediation framework universality is achieved by ensuring the exclusiveness of the gauge-loop contributions to the soft spectrum. A similar situation arises in the gravitational mediation framework if the sfermion spectrum is determined primarily by gaugino loops. Nevertheless, generic non-universalities in the scalar spectrum might appear in the gravitational mediation framework, in particular, if this framework is the low-energy limit of a string theory [20]. Even if universality, i.e., $K_{q_i q'_j} = \delta_{i j}$, is imposed by hand, it could be spoiled by either large radiative corrections (e.g., if the model is embedded in a grand unified theory [21]), or more generically, by the tower of nonrenormalizable operators in the Kahler potential which mix heavy and

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4 We do not specify the mechanism for the breaking of the horizontal symmetry. It could involve the $T$ fields if the constraint (11) is not satisfied. In a given model one would have to ensure that the horizontal scale is sufficiently high to suppress contributions to flavor violation from the horizontal sector.

5 Note that such frameworks often involve very heavy sparticl es, probes of which were recently discussed in Ref. [19].
light fields \[22\] (recall that supergravity is a nonrenormalizable theory). Attempts to impose universality in (perturbative) string theory by assuming the dominance of supersymmetry breaking in the universal dilaton field \[20\] are also subject to radiative corrections \[23\]. (In addition, no convincing example of such models exist.) Clearly, universality may be realized if the coupling of the various scalar fields to the Goldstino component of the gravitino is dictated by a symmetry. In string theory this would be the case if the gauge quantum numbers are linked to the conformal weights, in which case, the moduli contributions to the soft spectrum are charge dependent but flavor diagonal, as in the gauge-mediation framework.

Here, we would like to show that even if heavy-light mixing is eliminated order by order in \(M_P^{-1}\), radiatively generated mixing among the light fields themselves could still spoil the universality assumption. The singlet field is replaced by a gauge invariant composite operator of the light fields, \(e.g., \ QQ^\dagger/M_P\). The Kahler potential \([1]\) is rewritten in this case as

\[
K = \sum_I \left[ 1 + \frac{\alpha_I}{M_P^2} QQ^\dagger \right] \Phi_I \Phi_I^\dagger + \ldots, \tag{14}
\]

and includes mixing among the light fields \(\{Q_I\}\). The \(D\)-terms \([2]\) now contain the soft mass parameters

\[
\Delta_Q^2 \sim \sum_I \frac{\alpha_I}{(16\pi)^n} m_Q^2/2, \tag{15}
\]

where the summation is over the light fields. Note that flavor indices in eq. \(15\) need not be diagonal, leading to both non-universalities and explicit flavor off-diagonal masses. For \(m_Q^2 \sim m_{3/2}\), the size of the effect depends on the size of the counting parameter \(N = \sum_I \alpha_I\) and on the loop-order, \(i.e., \Delta_Q^2/m_Q^2 \sim N/(16\pi)^n\). The effect could be significant in comparison to the (rough) phenomenological upper bound of \(\Delta_Q^2/m_Q^2 \lesssim 10^{-3} [24]\).

Some comments are in place. Indeed, one expects the emergence of radiative corrections if certain terms which are allowed by all symmetries are eliminated by hand from the Kahler potential. (For example, this is the case of the universal singlet tadpole.) However, here we have a somewhat special case. The terms \(Q_1Q_1^\dagger Q_2Q_2^\dagger\) cannot be forbidden by any typical field-theory symmetry and hence, there are typically divergent radiative corrections to \(Q_1Q_1^\dagger\). The above terms could presumably be absent in a string theory framework in the case of T-duality. Otherwise, the radiatively induced tadpole (for the bilinears) aggravates the flavor problem in models of gravity mediated supersymmetry breaking. On the other hand, we saw above that a universal singlet tadpole term could offer alternative frameworks in which gravity is only a trigger for the generation of the MSSM parameters, and in which the flavor problem is solved by low-energy universality or decoupling.

We note in passing that the Kahler potential \([14]\) also provides a new mechanism for the generation of the Higgs mixing parameter \(m_{12}^2\). The bilinear \(QQ^\dagger\) is substituted in this case with the gauge invariant bilinear \(H_1H_2\). If lepton number is not a symmetry of the Kahler potential, the \(LH_1^\dagger\) and \(LH_2\) bilinears (\(L\) denotes a lepton doublet) would similarly generate Higgs-lepton mixing in the scalar potential.

It should be noted that a given theory contains only one scale for all tadpole operators which is given by \(M_{SUSY}\). Only one of the above mechanisms could be realized in a given
model. For example, our realization of the (ordinary) messenger model cannot benefit from our proposal for the generation of the Higgs mass parameters in models of low-energy supersymmetry breaking. However, a new source of tadpole diagrams could arise if there is a superpotential mixing between heavy and light singlets \[1\] (which are not necessarily universal singlets). The scale of the tadpole in this case depends on the heavy singlet mass \(M_H\) rather than on \(M_P\). One could couple the light singlet \(S_l\) to a vector-like multiplet by extending the superpotential

\[
W = M_H S_h^2 + m_l S_l^2 + \lambda_1 S_h^3 + \lambda_2 S_h^2 S_l + \lambda_3 S_l^3 + \lambda_4 S_h S_l^2
\]

of Ref. \[2\] with the appropriate Yukawa term \(W \to W + \lambda_5 S_l V \bar{V}\). One could construct, in principle, models with a few singlet fields and with both gravitational and heavy singlet tadpoles which are effectively described by

\[
|\Delta V| \sim m_{3/2}^2 M_P s \pm m_{3/2} M_H F_{s_l} + h.c.
\]

(Note that the assumptions in Ref. \[2\] correspond to our limit (2) above.) The singlet \(s\) is as before and its value is as in eq. (12), while \(s_l \sim \sqrt{m_{3/2} M_H}\). In this hybrid scenario for the messenger models, \(S\) couples to \(H_1 H_2\) while \(S_l\) couples to \(\bar{M} M\).

In conclusion, gravitationally induced tadpole terms in the low-energy theory aggravate the flavor problem if \(m_{\text{sparticle}} \approx m_{3/2}\), but may be desirable if \(m_{3/2} \ll m_{\text{sparticle}}\). In the latter case the universal singlet does not decouple and it parameterizes the effects of the heavy sector in the effective low-energy theory. In specific examples we were able to realize:

1. Mass terms for the Higgs fields if \(M_{SUSY} \sim 10^6\) GeV;
2. Gravitational triggering of a gauge mediation framework if \(M_{SUSY} \sim 10^8\) GeV;
3. A horizontal messenger model if \(M_{SUSY} \sim 10^{10}\) GeV.

All of the above mechanisms stem from a simple-minded toy model and are characterized by their simplicity. They merit further investigation and one may attempt to incorporate them individually (or simultaneously, if including heavy-light mixing) into more complete models.

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