IASSNS-HEP-97-137
CERN-TH/98-75
RU-97-97
hep-ph/9803310
March 1998

Stabilized Singlets in Supergravity as a Source of the μ-parameter

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Abstract

Within the context of supergravity-coupled supersymmetry, fields which are gauge and global singlets are usually considered anathema. Their vacuum expectation values are shifted by quadratically divergent tadpole diagrams which are cutoff at the Planck scale, destabilizing the classical potential and driving the singlet field to large values. We demonstrate a new and generic mechanism which stabilizes the singlet in the presence of an extended gauge symmetry. Such a symmetry will be broken down to the Standard Model by the supergravity interactions near the scale of spontaneous supersymmetry-breaking in the hidden-sector ($\sim 10^{10-11}$ GeV). The resulting singlet expectation value is stabilized and naturally of order the gravitino mass, providing therefore a weak-scale mass for the Higgs fields of the supersymmetric Standard Model (a μ-parameter). The resulting low-energy theory is the minimal supersymmetric Standard Model, with all new fields decoupling at the intermediate scale.

One of the primary motivations for low-energy supersymmetry (SUSY) is the stabilization of the weak scale (m_W) against quadratic divergences which appear at the quantum level and can drive m_W up to the Planck scale (M_P). To accomplish such a feat, one typically postulates that there are two separate sectors in the Lagrangian describing all of physics: a visible sector in which exist the Standard Model (SM) and its SUSY extension, the Minimal Supersymmetric Standard Model (MSSM); and a hidden sector in which exist the fields and interactions necessary to spontaneously break SUSY (see *e.g.* [1]). Consistency of the model requires that the hidden and visible sectors must not communicate with one another directly via renormalizable, tree-level interactions, and so one must also postulate a mediation mechanism which communicates SUSY-breaking from the hidden to the visible sector.

The simplest and most elegant means for communicating SUSY-breaking is via supergravity interactions; in fact, this mechanism is inherited directly from the gauging of SUSY and requires no additional sectors or interactions. In the picture of supergravity mediation, SUSY is spontaneously broken in the hidden sector at a scale $M_{\rm SUSY}$. This scale is then communicated to the visible sector through Planck-scale suppressed operators so that the visible SUSY-breaking scale is $M_{\rm SUSY}^2/M_P$. Setting this scale to m_W (in order to have weak-scale SUSY), one finds that the hidden sector scale is then $M_{\rm SUSY} \simeq \sqrt{m_W M_P} \equiv M_{\rm int} \simeq 10^{10-11}\,{\rm GeV}$. The resulting picture has three scales (though they are not all independent): M_P, $M_{\rm int}$ and m_W. The physics in each sector, however, is only sensitive to two of these scales: M_P and $M_{\rm int}$ in the hidden sector, and M_P and m_W in the visible sector. In both cases, one expects that SUSY-preserving masses should be near M_P, and that SUSY-breaking masses in each sector should be near the scale which controls SUSY-breaking in that sector.

However, there is one necessary, and another possible, violation of these expectations in low-energy SUSY. The former is often referred to as the μ-problem [2]. The superpotential of the MSSM must for phenomenological reasons contain a term $\mu H\overline{H}$ which mixes the two Higgs doublets responsible for electroweak symmetry breaking, with $\mu \sim m_W$. However, μ is both SUSY- and $SU(2)\times U(1)$-invariant and so we expect that its natural scale should be the Planck scale. The other possible violation arises whenever one wishes to extend the gauge structure of the MSSM to include new interactions broken at some scale intermediate to m_W and M_P. In particular, there exist several arguments that there should be new physics in the visible sector at a scale close to $M_{\rm int} \sim 10^{10-11}\,{\rm GeV}$ (for example, a right-handed neutrino at such a scale could provide a suppressed neutrino mass through the seesaw mechanism consistent with current indications). It is remarkable that this new scale for visible sector physics is also roughly the scale of SUSY-breaking in the hidden sector, though there is no obvious connection. In fact we will show in this letter that there may indeed be a deep underlying connection. Specifically, we will demonstrate a mechanism which simultaneously solves the μ-problem while generating a scale for new visible sector physics close to the scale $M_{\rm SUSY}$ of hidden sector physics.

Of the existing suggestions for solving the μ-problem of the MSSM, two bear some kinship to our proposal. The first such suggestion is to extend the matter content of the MSSM by a singlet, S, which couples to $H\overline{H}$. If S receives a weak-scale vacuum expectation value (vev), it dynamically provides a μ-parameter, $\mu \sim \langle S \rangle$. This is known as the Next-to-Minimal SUSY Standard Model (NMSSM) [3]. The second solution is the Giudice-Masiero mechanism [4] in which non-minimal, non-renormalizable Kähler couplings mix the hidden and visible sectors. Thus what appears to be a SUSY-preserving μ-parameter originates in fact from SUSY-breaking in the hidden sector. Our mechanism has ingredients from both proposals in that we add a gauge singlet to the spectrum along with a non-minimal Kähler potential. But the singlet will decouple from the low-energy theory, and no potentially dangerous mixing of hidden and visible sectors is necessary. Our proposal is more closely related to that of Sen (and as updated by Barr) [5] for reviving the "sliding singlet" solution for the doublet-triplet splitting problem of $SU(5)$. There, as here, dangerous destabilizing operators in the effective potential need to be suppressed; in Ref. [5], this is accomplished by extending $SU(5)$ in a very particular way. Their models appear, in fact, to be special cases of a more general principle which we will identify and explore here.

Begin by considering the NMSSM, which has two well-known problems. At the renormalizable level, the NMSSM has a Z_3 symmetry. If that symmetry is preserved to all orders, then the vev of S will break the symmetry at the weak scale and produce cosmologically dangerous domain walls. If on the other hand the Z_3 symmetry is not preserved by higher-order terms in the Lagrangian, then S carries no conserved quantum numbers. In this latter case, S will generically develop tadpoles, in the presence of spontaneously-broken SUSY, whose quadratic divergences are cut off by the Planck scale [6, 7]. The resulting shift in the potential for S causes it to slide to large values far above the weak scale. If it were to couple to the MSSM Higgs fields, they would receive unacceptably large masses, destabilizing the weak scale. Therefore one concludes that not only do singlets fail to provide a viable μ-parameter, but they cannot even be allowed to couple to light fields.

We now present a toy model which solves this destabilization problem of the NMSSM, while at the same time introducing a new visible sector interaction whose scale will naturally fall at $M_{\rm int}$. Within this model, singlets can couple to MSSM fields, and in particular, can provide dynamical μ-terms at the weak scale. The model itself demonstrates a very general mechanism, though it already contains all of the ingredients necessary to be phenomenologically viable.

The Model. Consider a superpotential

$$W = \lambda_H S H\overline{H} + \lambda_\Sigma S \Sigma\overline{\Sigma} \tag{1}$$

where $H, \overline{H}$ carry charges ± 1 under a gauge symmetry $U(1)_H$, $\Sigma, \overline{\Sigma}$ are charged ± 1 under another gauge symmetry $U(1)_\Sigma$, and S is a gauge singlet. We require $\Sigma, \overline{\Sigma}$ to be neutral under $U(1)_H$ and, for simplicity, assume that $H, \overline{H}$ are also neutral under

$U(1)_\Sigma$, though they need not be. The D-term for $U(1)_\Sigma$ is then simply

$$D_\Sigma = g_\Sigma \left(|\Sigma|^2 - |\overline{\Sigma}|^2 \right), \tag{2}$$

and similarly for D_H. To apply this toy model to the MSSM, we identify $H, \overline{H}$ as the usual Higgs doublets, and extend $U(1)_H$ to the Standard Model gauge group; Σ and $\overline{\Sigma}$ are new fields charged under a new gauge symmetry $U(1)_\Sigma$.

At the level of the superpotential, there exists the usual Z_3 which forbids explicit mass terms from appearing in W [3]. This symmetry is broken by $S \neq 0$, which could lead to creation of electroweak scale domain walls via the Kibble mechanism. The appearance of an S^3 term is forbidden by an R-symmetry under which $R(W) = 2$ and $R(S) = 0$.

However, we will assume that the Z_3 symmetry of the superpotential is only an accidental symmetry. This is a natural expectation since global symmetries are generally not preserved by quantum gravity effects (unless they are remnants of broken gauge symmetries) [8]. In particular, we expect that gravity-induced global symmetry-breaking will appear as non-renormalizable, explicit symmetry-breaking terms in the Kähler potential, K. The Z_3 symmetry can thus be a symmetry of the effective superpotential without being a symmetry of the entire action. This is equivalent to the statement that the S field is a true singlet, carrying no conserved quantum numbers.

After SUSY-breaking, non-zero tadpoles for S will generically arise with light chiral fields circulating in the loops [6, 7, 9, 10, 11]. These tadpoles appear at $\mathcal{O}(\hbar^n/M_P)$ due to supergravity corrections from Planck-suppressed operators. Because the exact source of the couplings which generate the tadpoles is highly model-dependent, we do not know *a priori* at what loop order non-zero contributions are generated. For example, it is known that for a flat Kähler metric, non-zero tadpoles do not arise until two-loops [9, 10]; however, for a non-flat metric they may arise at one-loop.

We will analyze these tadpole contributions following the formalism laid out in Refs. [7, 10]. Putting aside the question of loop-order at which the tadpoles arise, it is sufficient to consider these contributions as coming from an effective Kähler potential of the form

$$K = \left\{ 1 + \frac{1}{M_P}(S + S^\dagger) \right\} \Phi_i \Phi_i^\dagger. \tag{3}$$

The first term is just the canonical contribution to K of a chiral superfield Φ_i, while the second term explicitly breaks the Z_3-symmetry present in the superpotential. In general, an n-loop induced tadpole generates a term in the effective Lagrangian of the form:

$$\mathcal{L}_{\text{eff}} \sim \frac{N}{(16\pi^2)^n} \frac{\Lambda^2}{M_P} \int d^4\theta \, e^{K/M_P^2} \, (S + S^\dagger) \tag{4}$$

where N counts the number of light chiral superfields that appear in the loops and Λ is the cutoff for the quadratic divergence. (This expression results from a full supergravity calculation [10].) Henceforth we make the reasonable assumption that $\Lambda \simeq M_P$.

After SUSY-breaking, the superspace density, e^{K/M_P^2}, has the expansion:

$$e^{K/M_P^2} = 1 + \frac{1}{M_P^2}\left\{\theta^2 K_i F^i + \overline{\theta}^2 K_{\overline{\imath}} F^{\overline{\imath}} + \theta^2\overline{\theta}^2 \left(K_{i\overline{\jmath}} + \frac{K_i K_{\overline{\jmath}}}{M_P^2}\right) F^i F^{\overline{\jmath}}\right\} \tag{5}$$

where K_i ($K_{\overline{\imath}}$) is the derivative of the Kähler potential with respect to a superfield Φ_i ($\Phi_i^\dagger$). Even though i, j are summed over all fields in the theory, the vev of the superspace density is obviously dominated by the fields responsible for SUSY-breaking. In particular, we can take $F^i \sim M_{\rm SUSY}^2$ which is the scale of SUSY-breaking in the hidden sector and is related to the gravitino mass, $m_{3/2}$, via $M_{\rm SUSY}^2 \simeq m_{3/2} M_P$ (assuming cancellation of the cosmological constant). Throughout this article we will always assume that the communication of SUSY-breaking to the visible sector is via supergravity-induced terms, and thus $m_{3/2} \simeq m_W$. (See Ref. [11] for the case of tadpoles in gauge-mediated models for which $m_{3/2} \ll m_W$.)

While $K_{ij} \sim \mathcal{O}(1)$, the value of K_i is more model-dependent. For example, in the Polonyi model (see [1]), the field responsible for SUSY-breaking has an F-term $\sim m_W M_P$, but its scalar vev is $\sim M_P$; thus $K_i \sim M_P$. However, this is not a necessary ingredient and $K_i \ll M_P$ is possible. To be general we will write

$$K_i = \epsilon_i M_P \quad \text{for} \quad |\epsilon_i| \leq 1 \tag{6}$$

to parameterize our ignorance of the value of K_i. The contribution to the effective Lagrangian coming from the tadpole is then:

$$\begin{aligned}\mathcal{L}_{\rm eff} &\sim \frac{N}{(16\pi^2)^n}\left\{K_{\overline{\imath}} m_{3/2} F_S + \left(K_{i\overline{\jmath}} + \frac{K_i K_{\overline{\jmath}}}{M_P^2}\right) m_{3/2}^2 M_P S + h.c.\right\} \\ &\equiv \beta\epsilon m_{3/2} M_P F_S + \gamma m_{3/2}^2 M_P S + h.c.\end{aligned} \tag{7}$$

where ϵ is the maximum ϵ_i among the SUSY-breaking fields, and γ, β are complex coefficients which include the loop suppression factors $(16\pi^2)^{-n}$ and counting factors N, and so whose magnitudes are roughly $\mathcal{O}(10^{-4} - 1)$.

We are now in a position to write down the full scalar potential after SUSY-breaking, including supergravity-mediated soft masses as well as the tadpole contributions. Begin by considering the contributions to the scalar potential involving the F_S auxiliary field:

$$V_{F_S} = (\beta\epsilon m_{3/2} M_P F_S + h.c.) - |F_S|^2 - \left(F_S \frac{\partial W}{\partial S} + h.c.\right) \tag{8}$$

where the first term is the contribution of the tadpole. On integrating out all auxiliary fields, one finds that F_S is shifted from its canonical form by the tadpole contribution:

$$F_S^\dagger = -\frac{\partial W}{\partial S} + \beta\epsilon m_{3/2} M_P, \tag{9}$$

while all other F-terms (*e.g.* F_Σ and F_H) are canonical. The D-terms associated with the gauge fields also take their canonical forms.

The full scalar potential after soft SUSY-breaking can then be written:

$$\begin{aligned} V &= \sum_i m_i^2|\varphi_i|^2 + |\lambda_\Sigma S|^2\left(|\Sigma|^2+|\overline{\Sigma}|^2\right) + |\lambda_H S|^2\left(|H|^2+|\overline{H}|^2\right) \\ &\quad + m_{3/2}^2 M_P(\gamma S+\gamma^\dagger S^\dagger) + \left|\lambda_\Sigma \Sigma\overline{\Sigma} + \lambda_H H\overline{H} - \beta\epsilon m_{3/2}M_P\right|^2 \\ &\quad + \frac{g_\Sigma^2}{2}\left(|\Sigma|^2-|\overline{\Sigma}|^2\right)^2 + \frac{g_H^2}{2}\left(|H|^2-|\overline{H}|^2\right)^2 + A, B\text{-terms}, \end{aligned} \tag{10}$$

where the first term represents the gravitationally-induced soft SUSY-breaking masses, $m_i^2 \sim m_{3/2}^2$, for the fields $\varphi_i = \{S, \Sigma, \overline{\Sigma}, H, \overline{H}\}$. The final terms are gravitationally-induced soft-breaking bilinear, B, and trilinear, A, terms which for simplicity we ignore hereafter; they do not change our results substantially. Note that the potential as written requires that $m_S^2 \geq 0$ in order to be bounded from below (this condition is modified in the presence of B-terms). Indeed, one expects $m_S^2 > 0$ at tree level and it will only be driven negative if its coupling to either of the two sets of Higgs fields is fairly large. Henceforth we will take all soft squared-masses to be equal to $m_{3/2}^2 > 0$.

Minimization of the Potential. To continue further, we take $\epsilon \simeq 1$, *i.e.*, $K_i \simeq M_P$ which is the generic choice; small deviations of ϵ away from 1 can be absorbed into β. Writing down the minimization conditions for the potential is trivial, but as the potential is quite complicated, it has many local minima besides the true global one. However, there are two lowest-lying minima, both along directions that are D-flat up to weak-scale corrections, *i.e.*, $\Sigma \simeq \overline{\Sigma}$ and $H \simeq \overline{H}$.

At a first minimum, denoted V_1,

$$\begin{aligned} \Sigma = \overline{\Sigma} = H = \overline{H} &= 0, \\ S &\simeq -\gamma^\dagger M_P, \\ |F_S| &\simeq |\beta m_{3/2} M_P|, \\ V_1 \equiv V_{\min} &\simeq (|\beta|^2 - |\gamma|^2) m_{3/2}^2 M_P^2. \end{aligned}$$

This minimum represents the case usually considered in the literature for singlets with non-zero tadpoles — their vevs are pulled up to the Planck scale, taking with them any matter to which they couple. This is precisely the reason it was argued in Refs. [7, 10] that the vev of a true singlet cannot be responsible for the μ-term in the MSSM.

At a second minimum, V_2,

$$\begin{aligned} \Sigma\overline{\Sigma} &= \frac{\beta m_{3/2} M_P}{\lambda_\Sigma}, \\ H = \overline{H} &= 0, \end{aligned}$$

$$
\begin{aligned}
S &\simeq -\frac{\gamma^\dagger}{2|\lambda_\Sigma \beta|} m_{3/2}, \\
F_S &\sim m_{3/2}^2, \\
F_{\Sigma,\overline{\Sigma}} &\simeq \lambda_\Sigma m_{3/2}^{3/2} M_P^{1/2}, \\
V_2 \equiv V_{\rm min} &\simeq \frac{1}{|\lambda_\Sigma|} \left(|\beta| - \left| \frac{\gamma^2}{2\beta} \right| \right) m_{3/2}^3 M_P.
\end{aligned}
$$

The Σ-fields receive vev's of $\sim \sqrt{m_{3/2} M_P}$ to cancel off the F_S contribution to the potential. These large Σ-vev's then produce masses for the S-field (through the F_Σ terms) which stabilizes the S-vev against the tadpole-induced linear potential. The resulting vev of S is then only $\langle S \rangle \sim m_{3/2} \simeq m_W$!

Any gauge symmetry carried by the Σ-fields will be broken at the scale of their vev's. Up to the loop factors buried in β, this is the intermediate scale, $M_{\rm int}$. It is also, not coincidentally, the scale of SUSY-breaking in the hidden sector. In fact, one may interpret the physics at this minimum as the tadpoles communicating to the Σ-fields the true scale of SUSY-breaking, up to the loop factors.

There is also a third minimum, V_3, which is identical to V_2 except that the would-be MSSM Higgs fields, H and $\overline{H}$, play the role of Σ and $\overline{\Sigma}$ and receive vev's $\sim M_{\rm int}$, with λ_H replacing λ_Σ in all expressions. This is clearly *not* the desired minimum but is instead another example of how the tadpole can destabilize the weak scale. (Note that points at which $H, \overline{H}, \Sigma, \overline{\Sigma}$ all get vev's simultaneously are not even local minima of the potential.)

One still needs to resolve which of the three minima is the global one. Clearly $|V_{2,3}| \ll |V_1|$, but if $V_1 < 0$ then V_1 will be the global minimum. It is easy to show that $V_1 < 0$ if and only if $|\gamma|^2 > |\beta|^2$. As we know of no general argument that can set the relative sizes of γ and β, we simply regard the inequality $|\gamma|^2 < |\beta|^2$ to be a condition for the universe to lie in the interesting minimum, which, given the arbitrariness of β and γ, is not a restrictive requirement.

But assuming $V_1 > V_{2,3}$, which of either pair of fields $\Sigma, \overline{\Sigma}$ or $H, \overline{H}$ gets a vev? Though the form of the potential for the Σ- and H-fields are the same, they have differing quadratic pieces (effective masses) coming from their soft masses and their couplings to S. Assuming equality of soft masses, V at its minimum is positive (given $|\gamma|^2 < |\beta|^2$) and scales as either $|\lambda_\Sigma|^{-1}$ or $|\lambda_H|^{-1}$. Thus the global minimum of the potential occurs where only the pair of fields with largest absolute coupling to S gets vevs $\sim M_{\rm int}$. That is, if $|\lambda_H| < |\lambda_\Sigma|$, it is energetically unfavorable for $H, \overline{H}$ to receive vevs $\sim M_{\rm int}$ and so only the $\Sigma, \overline{\Sigma}$ fields do. Note in particular that simple inequality is all that is needed to ensure that a gauge hierarchy develops; no large hierarchy is needed between the two couplings themselves. Unequal soft masses shift the condition slightly, but the same basic result will always hold. Thus we conclude that the S-vev can in fact provide the μ-term of the MSSM as long as $|\lambda_H| < |\lambda_\Sigma|$. The whole question of which gauge group is broken at the scale $m_{3/2}$ and which at

$M_{\rm int}$ may rest entirely on the relative size of two couplings (λ_H and λ_Σ) whose ratio is generically $\mathcal{O}(1)$!

Recapping, we have found that if $|\beta|^2 > |\gamma|^2$ and if $|\lambda_H| < |\lambda_\Sigma|$ then S develops a vev at the weak scale and provides a μ-term of the correct size for light MSSM Higgs fields. Henceforth we will assume that these conditions hold and examine the solution in minimum V_2. It is well-known that models with a dynamical μ-term can contain a Peccei-Quinn (PQ) symmetry which would be spontaneously broken and thus create an unwanted axion at the weak scale. To examine this possibility, promote the PQ-symmetry to the previously discussed R-symmetry under which S is neutral and all other superfields are singly charged. However, $R(\beta m_{3/2} M_P) = 2$ explicitly breaks the symmetry and the would-be axions are all given masses near the intermediate scale, rendering them harmless. But there still remains a residual PQ symmetry in the MSSM Lagrangian. This too is explicitly broken, this time by $F_S \sim m_W^2$, which generates a B-term in the Higgs sector ($\sim F_S^\dagger H\overline{H}$) and gives mass to the pseudoscalar Higgs/would-be axion.

In detail, the heavy spectrum contains two massive pseudoscalars (with masses near $\langle\Sigma\rangle$), a massless Goldstone which is eaten by the gauge field, and three massive scalars. In particular, the S-scalar itself is not in the light spectrum. It receives a mass via its coupling to Σ and $\overline{\Sigma}$ near $M_{\rm int}$. Analysis of the heavy fermion mass spectrum is more complicated since ψ_S mixes with the ψ_Σ and $\psi_{\overline{\Sigma}}$ fields and the gaugino, χ, of the broken gauge group. The mass matrix in the $\{\psi_S, \psi_\Sigma, \psi_{\overline{\Sigma}}, \chi\}$ basis is given by

$$M = \begin{pmatrix} 0 & \lambda\overline{\Sigma} & \lambda\Sigma & 0 \\ \lambda\overline{\Sigma} & 0 & \lambda S & g\Sigma \\ \lambda\Sigma & \lambda S & 0 & -g\overline{\Sigma} \\ 0 & g\Sigma & -g\overline{\Sigma} & m_\chi \end{pmatrix} \tag{11}$$

where the (4,4) element is the soft SUSY-breaking mass of the χ: $m_\chi \simeq m_{3/2}$. This mass matrix has all four eigenvalues of order $M_{\rm int}$, so there is no remnant of the dynamical S field at the weak scale, only its weak scale vev.

While the application of S to the MSSM μ-term does not require that any of the MSSM fields be charged under the $U(1)_\Sigma$ gauge group, it is easy to see that our mechanism still works even if this charge were carried by some or all of the MSSM fields. Difficulties could arise if the D_Σ-term were large after symmetry-breaking, $D_\Sigma \gg m_{3/2}^2$. Then all fields charged under that symmetry would be pulled up to the scale of the (large) D_Σ-term. However, the Σ-fields receive their vevs along an approximately D-flat direction $\Sigma \simeq \overline{\Sigma}$ where:

$$\Sigma^2 - \overline{\Sigma}^2 = \frac{1}{4g_\Sigma^2}\left(m_\Sigma^2 - m_{\overline{\Sigma}}^2\right). \tag{12}$$

Thus, just as in the case without tadpoles, the mass corrections to the MSSM fields are not dangerous, though perhaps observable [12].

It is clear from our analysis why it was necessary to forbid S^3 terms from the superpotential (such a term is automatically forbidden if S carries zero R-charge).

Generically, these terms disrupt minimization of the potential since S can take on a vev itself to cancel off the large F_S term in V. While this may provide a natural way of giving intermediate scale masses to vector-like matter, it does not give a μ-term. (The situation is again somewhat different in the case of gauge-mediated SUSY breaking [11].)

The Lagrangian of Eq. (7) receives corrections after SUSY-breaking from the one-loop effective potential. The dominant contribution is an S^4 term:

$$\Delta V = \frac{1}{64\pi^2}\mathrm{STr}\, M^4 \log(M^2/Q^2) \simeq \frac{\lambda_\Sigma^4}{64\pi^2}\left(\frac{m_{3/2}}{M_P}\right) S^4 + \cdots . \tag{13}$$

However, this correction is too small to affect the minimization of the potential or our results. As well, if mixing in K between the hidden and visible sectors is allowed, terms of the form $K = Z^\dagger Z S/M_P + h.c.$ generically arise. After carrying out the $d^4\theta$ integral, these echo the form of the tadpole-induced contributions, only with $\mathcal{O}(1)$ coefficients. Thus their entire effect can be absorbed into redefining $\beta, \gamma \sim \mathcal{O}(1)$, which leaves our central result unchanged.

Conclusions. In this letter, we have demonstrated a new mechanism for obtaining a weak-scale μ-term in the MSSM by adding a total singlet in conjunction with a new gauge interaction and its accompanying Higgs sector. We have used the tadpoles endemic to models with singlets to drive the breaking of the new symmetry at the intermediate scale $\sqrt{m_W M_P}$, to remove all vestiges of the singlet from the low-energy theory, and to render all would-be axions harmless. In this way we have pointed out a natural way to tie the introduction of intermediate-scale operators, such as those that might give mass to a right-handed neutrino, to the solution of the weak-scale μ-problem. The model itself is highly generic and can naturally explain the hierarchy between the weak and intermediate scales in terms of the ratio of two couplings which is $\mathcal{O}(1)$. It seems straightforward to generalize our discussion to more complete models of physics at intermediate scales [13].

Acknowledgements

We would like to thank K.S. Babu, G.R. Farrar, S.F. King and H.-P. Nilles for conversations. The work of CK is supported by Department of Energy contract DE-FG02-90ER40542 and through the generosity of Helen and Martin Chooljian. The work of SP is supported by the Polish Committee for Scientific Research, grant 2P03B0412 (1997-8) and by the US-Poland Maria Sklodowska-Curie Fund MEN/DOE-96-264. The work of NP is supported by National Science Foundation grant PHY-94-23002. We would all like to thank the Aspen Center for Physics, where this work was initiated, for its hospitality.

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[13] *Work in progress.*