Gauge Mediation Simplified

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Gauge mediation of supersymmetry breaking is drastically simplified using generic superpotentials without $U(1)_R$ symmetry by allowing metastable vacua.

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Supersymmetry [1] is a manifestation of quantum spacetime dimensions that has been hypothesized as the leading candidate for physics beyond the standard model of particle physics. It is a necessary ingredient of string theory which promises to unify gravity and quantum physics. If present below the TeV energy scale, it stabilizes the 16 orders of magnitude of hierarchy between the electroweak and Planck scales against radiative corrections [2]. In addition, it provides a natural candidate for the cold dark matter of the Universe [3], achieves unification of gauge forces with the minimal particle content [4], and may play an important role in keeping the inflaton potential flat to account for the observed near scale-invariant density fluctuations. There is a strong anticipation of its discovery at the Large Hadron Collider (LHC) that starts its operation at the end of 2007.

However, there is a growing concern that its imprint on rare phenomena should have already been seen experimentally if the superparticles indeed exist below the TeV scale to stabilize the hierarchy. In particular, recent beautiful data from B factories [5] and Tevatron [6] have not shown hints of superparticles in their flavor-changing effects. Similarly, recent upper limits on electric dipole moments of neutron, electron, and mercury atom also contradict with sub-TeV superparticles if they have large CP violation [7]. Unless there is a good reason to believe that the superparticles do not contribute to flavor-changing or CP-violating effects, sub-TeV superparticles appear unlikely. Is there such a good reason? If so, does it require very special models or apply generically to many supersymmetric models? These are urgent questions associated with the mechanism of supersymmetry breaking and its mediation to the supersymmetric standard model.

Let us first consider the issue of supersymmetry breaking. Breaking supersymmetry has been a nontrivial task. A general argument by Nelson and Seiberg is that it requires a theory with an exact continuous $U(1)_R$ symmetry if we assume that the superpotential is generic [8]. In addition, an argument based on the Witten index [9] said that the theory must be chiral. This is because one can continuously deform a vectorlike theory by mass terms to a pure Yang-Mills theory, which is known to have a nonzero Witten index (dual Coxeter number) and hence supersymmetric vacua. Chirality and $U(1)_R$ invariance strongly limit the choice of possible theories that break supersymmetry. Later, vectorlike models were found [10]. They evade the Witten index argument because the mass terms can always be absorbed by shifting singlet fields in the theory. The required superpotential, however, is not generic unless one imposes an exact $U(1)_R$ symmetry.

The requirement of an exact $U(1)_R$ symmetry is unfortunate, because exact global symmetries are not expected to exist in quantum theory of gravity such as the field-theory limit of string theory. In addition, embedding a model of supersymmetry breaking into supergravity requires explicit breaking of $U(1)_R$ to allow for a constant term in the superpotential needed for canceling the cosmological constant. Once $U(1)_R$ is not an exact symmetry, it is not clear how one can justify the form of the superpotential required for supersymmetry breaking.

Even if supersymmetry is successfully broken, it is certainly not enough. One needs to mediate its effects to the supersymmetric standard model sector in such a way that it does not lead to large flavor-changing or CP-violating effects. An attractive idea for achieving this is to use standard model gauge interactions for the mediation, since they conserve both flavor and CP, and this possibility includes a class of models called gauge mediation [11,12]. Unfortunately, however, all known models constructed along these lines rely on very specific choices of particle content and/or gauge groups, and it is fair to say that they are far from generic among possible supersymmetric theories that may come out from a fundamental theory such as string theory. This is mainly because the mediation mechanism should be added very carefully to a model of supersymmetry breaking if we want to preserve its fragile structures, including the one associated with $U(1)_R$ symmetry. Moreover, once $U(1)_R$ is imposed, it needs to be spontaneously broken and its breaking effect mediated so that the gaugino masses are successfully generated. The spontaneously broken $U(1)_R$ symmetry also raises the issue of a potentially hazardous R axion.

In this Letter, we advocate to discard $U(1)_R$ symmetry altogether from the theory, and allow for completely generic superpotentials. According to the Nelson-Seiberg argument, such a theory would not break supersymmetry.

Yet, it may have a *local* supersymmetry breaking minimum. Supersymmetry is broken if the low-energy limit of the supersymmetry breaking sector has an *accidental* $U(1)_R$ symmetry, which nonetheless is broken by its coupling to messengers of supersymmetry breaking. Indeed, we show a very simple class of models of this type. The models do not have a fundamental singlet field, eliminating aesthetic and various fine-tuning problems in cosmology and preserving the hierarchy. The gauginos and scalars in the supersymmetric standard model sector obtain flavor universal masses by standard model gauge interactions through loops of the messengers, i.e., through gauge mediation. Given the absence of $U(1)_R$, there is no problem in generating gaugino masses, and no dangerous R axion arises.

An explicit model that realizes our general philosophy is a supersymmetric $SU(N_c)$ QCD with massive vectorlike quarks Q^i and $\bar{Q}^i$ $(i = 1, \cdots, N_f)$. In addition, we introduce massive messengers f and $\bar{f}$ and write the most general superpotential consistent with the gauge symmetry. This is the entire model. The important terms in the superpotential are given by

$$W_{\rm tree} = m_{ij}\bar{Q}^i Q^j + \frac{\lambda_{ij}}{M_{\rm Pl}}\bar{Q}^i Q^j \bar{f} f + M\bar{f}f, \tag{1}$$

where λ_{ij} are coupling constants [13]. (The effects of other terms will be discussed later.) For concreteness, we take the messengers f, $\bar{f}$ to be in $\mathbf{5} + \mathbf{5}^*$ representations of $SU(5)$ in which the standard model gauge group is embedded.

Intriligator, Seiberg, and Shih (ISS) pointed out that supersymmetric $SU(N_c)$ QCD in the free magnetic phase $(N_c + 1 \le N_f < \frac{3}{2}N_c)$ breaks supersymmetry on a metastable local minimum if the quark masses m_{ij} are much smaller than the dynamical scale Λ [14]. Note that in the ISS model a $U(1)_R$ symmetry is broken only down to Z_{2N_c} which prevents the gaugino masses. In the present model, however, the coupling to the messengers breaks it down to Z_2, so that the model does not have any R symmetry beyond R parity.

For the sake of concreteness, we discuss the case without the magnetic gauge group $N_f = N_c + 1$ below, although any $N_c + 1 \le N_f < \frac{3}{2}N_c$ works equally well. At energies below the dynamical scale, the nonperturbative low-energy effective superpotential is described as [15]

$$W_{\rm dyn} = \frac{1}{\Lambda^{2N_f-3}}(\bar{B}_i M^{ij} B_j - {\rm det}M^{ij}), \tag{2}$$

where $M^{ij} = \bar{Q}^i Q^j$, $B_i = \epsilon_{ii_1\cdots i_{N_c}} Q^{i_1}\cdots Q^{i_{N_c}}/N_c!$, and $\bar{B}_i = \epsilon_{ii_1\cdots i_{N_c}}\bar{Q}^{i_1}\cdots\bar{Q}^{i_{N_c}}/N_c!$ are meson, baryon, and antibaryon chiral superfields, respectively. In the following, we adopt the basis in which the quark mass matrix is diagonal, $m_{ij} = -m_i\delta_{ij}$, with m_i real and positive. We also assume that they are ordered as $m_1 > m_2 > \cdots > m_{N_f} > 0$ without loss of generality. Here, we have taken all masses different to avoid (potentially) unwanted Nambu-Goldstone bosons.

In terms of fields with canonical dimensions $S^{ij} = M^{ij}/\Lambda$, $b_i = B_i/\Lambda^{N_f-2}$ and $\bar{b}_i = \bar{B}_i/\Lambda^{N_f-2}$, the dynamical superpotential of Eq. (2) together with the quark mass terms [the first term of Eq. (1)] can be written as [16]

$$W_{\rm ISS} = \bar{b}_i S^{ij} b_j - \frac{{\rm det}S^{ij}}{\Lambda^{N_f-3}} - m_i\Lambda S^{ii}. \tag{3}$$

For $N_f > 3$, the superpotential term ${\rm det}S^{ij}$ is irrelevant and can be ignored to discuss physics around the origin $S^{ij} = 0$ [17]. The superpotential of Eq. (3) then leads to a local minimum at

$$b = \bar{b} = \begin{pmatrix} \sqrt{m_1\Lambda} \\ 0 \\ \vdots \\ 0 \end{pmatrix}, \qquad S^{ij} = 0, \tag{4}$$

where supersymmetry is broken because $F_{S^{ij}} = -(\partial_{S^{ij}}W)^* = m_i\delta_{ij}\Lambda \neq 0$ for $i, j \neq 1$. Even though S^{ij} $(i, j \neq 1)$ are classically flat directions, they are lifted by the one-loop Coleman-Weinberg potential. As a result, the origin $S^{ij} = 0$ is a local minimum, with curvature $m^2_{S^{ij}} \sim m\Lambda/16\pi^2$ for all $m_i \sim m$. It is long-lived as long as $m_i \ll \Lambda$, where the weakly coupled analysis of the low-energy theory is valid.

The existence of a supersymmetry breaking minimum of Eq. (4) can be viewed as a result of an accidental (and approximate) $U(1)_R$ symmetry possessed by the superpotential of Eq. (3) with the R-charge assignments $R(S^{ij}) = 2$, $R(b_i) = R(\bar{b}_i) = 0$, in the limit of neglecting the irrelevant term of ${\rm det}S^{ij}/\Lambda^{N_f-3}$. In fact, this accidental $U(1)_R$ symmetry is also a reason for the origin $S^{ij} = 0$ being the minimum of the effective potential as a symmetry enhanced point. This picture is corrected by the coupling of Q^i and $\bar{Q}^i$ to the messengers and by higher dimension terms in the superpotential omitted in Eq. (1), which introduce $U(1)_R$ violating effects to the supersymmetry breaking sector. These effects, however, can be easily suppressed as we will see later, and the basic picture described above can be a good approximation of the dynamics.

At the supersymmetry breaking minimum of Eq. (4) [with S^{ij} slightly shifted due to $U(1)_R$ violating effects], the messenger fields have both supersymmetric and holomorphic supersymmetry breaking masses:

$$M_{\rm mess} = M + \frac{\lambda_{ij}\Lambda}{M_{\rm Pl}}\langle S^{ij}\rangle \simeq M, \tag{5}$$

and

$$F_{\rm mess} = \frac{\lambda_{ij}\Lambda}{M_{\rm Pl}}F_{S^{ij}} = \frac{\bar{m}\Lambda^2}{M_{\rm Pl}}, \tag{6}$$

where

$$\bar{m} \equiv \sum_{i \neq 1} \lambda_{ii} m_i. \tag{7}$$

The usual loop diagrams of the messenger fields then induce gauge-mediated scalar and gaugino masses in the supersymmetric standard model sector, of the magnitude [11,12]

$$m_{\rm SUSY} \simeq \frac{g^2}{16\pi^2} \frac{\bar{m}\Lambda^2}{M M_{\rm Pl}}, \tag{8}$$

where g represents generic standard model gauge coupling constants. Since these masses are generated by standard model gauge interactions, they are flavor universal. Once we introduce the term $\mu H_u H_d$ in the superpotential with $\mu \sim m_{\rm SUSY}$, where H_u and H_d are the two Higgs doublets of the supersymmetric standard model, the electroweak symmetry can be successfully broken. This does not induce any new CP violation in the supersymmetric standard model sector beyond that in the standard model. Therefore, sub-TeV superparticles are phenomenologically allowed in this class of models.

Several conditions for the parameters need to be met for the model to be phenomenologically successful. Even though not necessary, we regard all the quark masses (and the couplings λ_{ij}) to be comparable, $m_i \sim m$ ($\lambda_{ij} \sim \lambda$), in the numerical estimates below.

First, we would like $m_{\rm SUSY}$ to stabilize the electroweak scale, and hence $m_{\rm SUSY} = O(100~{\rm GeV}–1~{\rm TeV})$. This corresponds to

$$\frac{\bar{m}\Lambda^2}{M M_{\rm Pl}} \approx 100~{\rm TeV}. \tag{9}$$

On the other hand, we would like the gauge-mediated contribution to the scalar masses dominate over the gravity-mediated piece to avoid excessive flavor-changing processes, leading to $m_{3/2} \approx m\Lambda/M_{\rm Pl} \lesssim 10^{-2} m_{\rm SUSY}$. Therefore,

$$mM \lesssim 10^{-4} \bar{m}\Lambda. \tag{10}$$

We also need the messengers to be nontachyonic,

$$M^2 > \frac{\bar{m}\Lambda^2}{M_{\rm Pl}}. \tag{11}$$

In addition, the analysis of supersymmetry breaking is valid only if m is sufficiently smaller than Λ:

$$m \lesssim 0.1\Lambda. \tag{12}$$

We now discuss the effects of $U(1)_R$ violation. These effects cause shifts of S^{ij} from the origin, which must be smaller than $\approx 4\pi\sqrt{m\Lambda}$ for the ISS analysis to be valid, and than $\approx M M_{\rm Pl}/\lambda\Lambda$ to avoid tachyonic messengers. One origin of $U(1)_R$ violation comes from higher dimension terms in the superpotential, omitted in Eq. (1). The dominant effect comes from

$$\Delta W = \frac{\lambda_{ijkl}}{M_{\rm Pl}} \bar{Q}^i Q^j \bar{Q}^k Q^l = \frac{\lambda_{ijkl}\Lambda^2}{M_{\rm Pl}} S^{ij} S^{kl}. \tag{13}$$

These terms may destabilize the minimum, since they lead to linear terms of S^{ij} in the potential through $F_{S^{ij}} = m_i \delta_{ij} \Lambda$ [18]. The squared masses of S^{ij} from the one-loop effective potential are $m^2_{S^{ij}} \sim m\Lambda/16\pi^2$, while the linear terms are $\sim(\lambda_{ijkk} m_k \Lambda^3/M_{\rm Pl}) S^{ij}$. Therefore, the shifts of the fields are $\Delta S^{ij} \sim 16\pi^2 \lambda_{ijkk}\Lambda^2/M_{\rm Pl}$. Requiring this to be sufficiently small, we obtain the condition

$$\frac{\lambda_{ijkk}\Lambda^2}{M_{\rm Pl}} \lesssim \min\left\{0.1(m\Lambda)^{1/2}, 10^{-2}\frac{M M_{\rm Pl}}{\lambda\Lambda}\right\}. \tag{14}$$

Similar conditions can be worked out for even higher order terms, but they are rather mild.

Another source of $U(1)_R$ violation comes from the coupling of Q^i and $\bar{Q}^i$ to the messengers, which shifts the minimum of S^{ij} at the loop level. The effect of the messengers on the S^{ij} effective potential can be calculated by computing the one-loop Coleman-Weinberg potential arising from the last two terms of Eq. (1):

$$W_{\rm mess} = \frac{\lambda_{ij}\Lambda}{M_{\rm Pl}} S^{ij} \bar{f} f + M \bar{f} f. \tag{15}$$

The resulting effective potential takes the following generic form

$$\Delta V \approx \frac{\bar{m}^2\Lambda^4}{16\pi^2 M_{\rm Pl}^2} \mathcal{F}\left(\frac{\lambda_{ij}\Lambda S^{ij}}{M M_{\rm Pl}}\right), \tag{16}$$

where $\mathcal{F}(x)$ is a real polynomial function with the coefficients of $O(1)$ up to symmetry factors. The resulting shifts of S^{ij} are of order $\lambda^3 m\Lambda^4/M M_{\rm Pl}^3$, which are sufficiently small if

$$M \gtrsim \frac{\lambda^2 m^{1/2} \Lambda^{5/2}}{M_{\rm Pl}^2}. \tag{17}$$

Note that the coupling to the messengers in Eq. (15) does not generate a new supersymmetric minimum. However, turning on the expectation values for the messengers may allow for lowering the vacuum energy, depending on the combinations of m_{ij} and $\lambda_{ij}\bar{f}f$. Even if this is the case, the tunneling to a lower minimum at $\bar{f}f \approx m M_{\rm Pl}/\lambda$ can easily be made suppressed to the level consistent with the longevity of our Universe, if $M M_{\rm Pl}/\lambda \gtrsim m^{1/2}\Lambda^{3/2}$.

It is now easy to see that there is a wide range of parameters that satisfy the conditions Eqs. (9)–(12), (14), and (17). For instance, if we take $\lambda_{ij} \sim \lambda_{ijkl} \sim 1$, $\Lambda \sim 10^{11}$ GeV, $m \sim \bar{m} \sim 10^8$ GeV, and $M \sim 10^7$ GeV, then all the requirements are easily satisfied. Note that the conditions of Eqs. (14) and (17) are generically rather weak, unless Λ is close to $M_{\rm Pl}$. This is because the relevant interactions in Eqs. (13) and (15) arise from higher dimension operators suppressed by $M_{\rm Pl}$.

Finally, we discuss if there are any unwanted light fields in the model. The fermionic fields in S^{ij} (i, $j \neq 1$) are massless in the ISS model, but they acquire masses here due to the generic terms in Eq. (13) [19]. They can decay to standard model particles through their coupling to the messengers and hence harmless. There is a Nambu-Goldstone boson (NGB) of a spontaneously broken $U(1)_B$ symmetry, $b^1 - \bar{b}^1$, and its fermionic partner. Exactly massless NGB and fermion would be a radiation component of the Universe. Their abundance is diluted by an order of magnitude due to the QCD phase transition and is in general consistent with the constraint from the big-bang nucleosynthesis, $\Delta N_\nu \lesssim 1.5$ [20]. Alternatively, they can be made massive by gauging $U(1)_B$, or avoided entirely by employing an $SO(N_c)$ or $Sp(N_c)$ gauge group for supersymmetry breaking, instead of $SU(N_c)$. The gravitino is the lightest supersymmetric particle and hence stable if R parity is unbroken. It places an upper limit on the reheating temperature [21], which is acceptable, e.g., in leptogenesis models by nonthermal production of right-handed scalar neutrinos [22].

A generic consequence of the proposed class of models is that superparticles produced in the future accelerator experiments would decay to the gravitino. In addition, given that there is no flavor or CP problem, the superparticles can be relatively light, allowing for their copious production at the LHC and International Linear Collider. For example, the lightest neutralino may decay into a gravitino and a photon, giving extra signatures to be searched for [23]. A more dramatic possibility is that there is a charged particle (e.g., stau) that is quasistable and leaves a track with anomalously large dE/dx. It may decay inside or outside the detector depending on the scale of fundamental supersymmetry breaking $\approx \sqrt{m\Lambda}$.

In summary, we advocated gauge mediation models of supersymmetry breaking with generic superpotentials without $U(1)_R$ symmetry. Using metastable minima, we find a class of phenomenologically successful models without any elementary gauge singlet fields. The simplicity and generality of the models are quite remarkable. Given the wide variety of models available, there is good reason to expect that flavor and CP are conserved in supersymmetry. We conclude that sub-TeV supersymmetry is still quite possible despite the lack of experimental hints so far.

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