An Extension for Direct Gauge Mediation of Metastable Supersymmetry Breaking (Postprint)

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Abstract

We study the direct mediation of metastable supersymmetry breaking by a Φ2Φ2Φ^2 (ΦΦΦ is the meson field) deformation to the ISS model and extend it by splitting ΦΦΦ into two parts and gauging the flavor symmetry. We find that with such an extensi

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Preamble

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Abstract

We study direct mediation of metastable supersymmetry breaking through a Φ^2-deformation of the ISS model, extending it by splitting both $\mathrm{Tr}\Phi$ and $\mathrm{Tr}\Phi^2$ terms in the superpotential and gauging the flavor symmetry. We find that this extension yields sufficiently long-lived metastable vacua and generates proper gaugino masses. Furthermore, it allows for constructing models that avoid the Landau pole problem. In particular, our metastable vacua exhibit a large region for the parameter m_3 that satisfies phenomenological requirements and permits a low SUSY breaking scale ($h\mu^2 \sim 100$ TeV).

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Introduction

Dynamical supersymmetry (SUSY) breaking provides a compelling solution to the gauge hierarchy problem, though realistic models must generally satisfy numerous theoretical constraints. Moreover, phenomenological considerations become complex when these dynamical SUSY breaking effects are mediated to the visible sector.

Recently, Intriligator, Seiberg and Shih (ISS) [1] discovered metastable supersymmetry breaking in the surprising context of a vector-like theory, offering a natural framework for dynamical SUSY breaking and its mediation. Their model (the ISS model) has attracted considerable attention, but suffers from several problems. One is the Landau pole problem, which implies that unification cannot be simply realized. Another is that an accidental R-symmetry (a generic feature of SUSY breaking models) forbids gaugino masses. Numerous approaches have been proposed to address these issues [2–6].

In the ISS model, SUSY breaking mediation can proceed via gauge interactions. Indeed, gauge mediation of dynamical SUSY breaking was previously proposed in [7, 8], attempting to use a QCD-like strong interaction to break supersymmetry dynamically while identifying the Standard Model gauge group as a subgroup of the flavor symmetry. However, these early models suffered from phenomenological problems such as the Landau pole problem and the gaugino mass problem, causing the idea to be abandoned for some time. The advent of the ISS model revived this approach [3, 4].

The ISS model can yield the required phenomenology through certain deformations. In [2], a $\mathrm{Tr}\Phi^2$ term was introduced to the superpotential (Φ^2-deformation), and in [3] the $\mathrm{Tr}\Phi$ term was split. In this work, we consider a more general deformation of the ISS model by splitting both $\mathrm{Tr}\Phi$ and $\mathrm{Tr}\Phi^2$ terms in the superpotential. We find that this general deformation satisfies phenomenological constraints, including generating appropriate gaugino masses and avoiding the Landau pole.

This paper is organized as follows. In Section II we briefly review the ISS model and its Φ^2-deformation. In Section III we present our general deformation and discuss its phenomenology. Finally, in Section IV we present our conclusions.

II. The ISS Model and Its Φ^2-Deformation

In the ISS model, the hidden sector consists of $N = 1$ supersymmetric QCD with massive quarks satisfying $N_c < N_f < \frac{3}{2}N_c$, where N_f is the flavor number and N_c the color number. The superpotential in the dual magnetic theory is $W = hq_i\Phi_{ij}\tilde{q}_j - h\mu^2\mathrm{Tr}\Phi$, where q and $\tilde{q}$ are the quark and anti-quark fields, Φ is the meson field, and $i, j = 1, ..., N_f$ are flavor indices.

In the low-energy region, the F-terms of Φ cannot be simultaneously set to zero because the rank of $q\tilde{q}$ is $N_f - N_c$, which is smaller than the rank of Φ (which equals N_f). This yields SUSY-breaking vacuum energy $V = N_c|h\mu^2|^2$. However, at high energies below the scale $\langle h\Phi\rangle$, the quarks are integrated out and the effective theory becomes $SU(N_f - N_c)$ pure Yang-Mills, where non-perturbative corrections to the superpotential restore the SUSY vacua, rendering the low-energy SUSY-breaking vacua only metastable.

The Φ^2-deformation of the ISS model was proposed by Giveon and Kutasov (GK) [2]. Its superpotential can be derived from brane configurations [9–11] and takes the form $W = hq_i\Phi_{ij}\tilde{q}_j - h\mu^2\mathrm{Tr}\Phi + h^2\mu_\phi\mathrm{Tr}\Phi^2$, where μ_ϕ is a new energy scale. Note that there is no non-perturbative superpotential from gaugino condensation effects. We assume $\ll$ to ensure the Φ^2-deformation remains a perturbation to the ISS model in the low-energy region.

This deformed supersymmetric QCD has a rich landscape of supersymmetric and non-supersymmetric vacua. The field expectation values are given by $\langle h\Phi\rangle = I_{N_f-k}\langle q\tilde{q}\rangle = \mu^2 I_k 0$ in supersymmetric vacua (where I_n denotes an $n\times n$ unit matrix), and $\langle h\Phi\rangle = 0h\Phi_n I_n I_{N_f-k-n}\langle q\tilde{q}\rangle = \mu^2 I_k 00$ in metastable non-supersymmetric vacua. In the latter case, we must consider the one-loop contribution to the potential, which to leading order is given by [1]

$$V_{1-loop} = b|h^2\mu|^2\mathrm{Tr}\Phi^\dagger\Phi\ln(\Phi^\dagger\Phi/\mu^2) - 1$$

with b being a constant given by $b = (N_f - N_c)/(16\pi^2)$. In our following analysis we fix $b = 0.01$ for convenience. The full potential for $\Phi_n, q, \tilde{q}$ then takes the form

$$V = |h|^2[|\Phi_n q|^2 + |\Phi_n\tilde{q}|^2 + |q\tilde{q} - \mu^2 I_n + h\mu_\phi\Phi_n|^2] + b|h\mu|^2\mathrm{Tr}\Phi^\dagger\Phi.$$

To be free of tachyons in the dual quark direction and to ensure the supersymmetric vacua are far enough from metastable vacua, we have (neglecting phase factors in energy scales and coupling constants)

$$\mu_\phi/b \ll h \gg \mu^2/(hb) \ll \mu_\phi.$$

Here the first constraint ensures sufficiently long-lived metastable vacua (as discussed below), while the second arises from special properties of this deformation (also applicable to our more general deformation in the following section) and can be obtained from analyzing the potential in Eq. (9) via $\partial V/\partial\Phi$. These constraints favor small h.

To check if the GK metastable vacua are sufficiently long-lived, we estimate the decay rate by evaluating the Euclidean action S_1 from the GK metastable vacua to their corresponding true vacua and the action S_2 from the GK metastable vacua to the ISS metastable vacua. Using the triangle approximation [14, 15], we estimate the bounce action as

$$S \sim (\Delta\Phi)^4/\Delta V \sim (\mu_\phi/(bh))^4 \left(\frac{\mu_\phi/(bh)}{\mu}\right)^4 \left(\frac{\mu}{h\mu_\phi}\right)^4$$

where $\Delta\Phi$ is the interval between the two vacua. Therefore, under the conditions in Eq. (10), we can obtain a sufficiently long lifetime for the universe even if h is not too small.

After constructing a dynamical supersymmetry breaking model, the next task is to mediate the breaking effects to the visible sector. Although constraints from metastable vacuum lifetime are weak, gaugino masses may be much smaller than sfermion masses in direct gauge mediation models where a flavor subgroup in the ISS sector is gauged. In the following section, we propose a more general deformation by splitting both $\mathrm{Tr}\Phi$ and $\mathrm{Tr}\Phi^2$ terms in the superpotential. With such a deformation, we can avoid the hierarchy between gaugino and sfermion masses and, furthermore, evade the Landau pole problem.

III. A More General Deformation to the ISS Model

In [3] the $\mathrm{Tr}\Phi$ term in the superpotential was split. Here we propose a more general deformation with superpotential

$$W = h\mathrm{Tr}(q\tilde{q}\Phi) - h\mu_1^2\mathrm{Tr}Y - h\mu_2^2\mathrm{Tr}\hat{\Phi} + h^2m_1^2\mathrm{Tr}Y^2 + h^2m_2^2\mathrm{Tr}\hat{\Phi}^2 + h^2m_3\mathrm{Tr}Z\tilde{Z} + h^2m_3\mathrm{Tr}\tilde{Z}\hat{\Phi}Z$$

where Y is an $N_1 \times N_1$ matrix, $\hat{\Phi}$ is an $N_2 \times N_2$ matrix, and m_1, m_2, m_3, μ_1 and μ_2 are mass scales. For $m_1 = m_2 = m_3 = \mu_\phi$ and $\mu_1 = \mu_2 = \mu$, our deformation reduces to the GK Φ^2-deformation [2]. The above superpotential has $SU(N_1) \times SU(N_2)$ flavor symmetry.

Our deformation not only exhibits a rich landscape of supersymmetric and non-supersymmetric vacua like the GK Φ^2-deformation, but also yields more appropriate phenomenology. We discuss several features below.

First, we examine the vacua in our deformation. The supersymmetric vacua are given by $\langle h\Phi\rangle = I_{N_1-k}\langle q\tilde{q}\rangle = \mu_1^2 I_k 0$, where k can run from 0 to N_1. Including the one-loop potential and following the procedure in [2], we obtain the metastable vacua as $\langle h\Phi\rangle = (m_1/b)I_{N_1}\langle q\tilde{q}\rangle = bI_{N_2-n}$, where we take $k = N_1$ and the vacuum energy to leading order is $V \simeq (N_2-n)|h\mu_2^2|^2$. Here the last component of Φ gives non-zero F-terms. As discussed in the preceding section, we have the condition $m_2^2/b \ll \mu_2^2 \leq m_2^2/(hb)$ for long-lived metastable vacua without tachyonic quarks. Note that above we took $k = N_1$; our following analysis also discusses the case $k = 0$ without presenting explicit structure. Other metastable vacua exist in our model, but those above suffice for constructing an appropriate phenomenological model.

In our metastable vacua in Eq. (16), the flavor symmetry $SU(N_2)$ breaks to $SU(n) \times SU(N_2 - n)$. We can gauge the flavor symmetry $SU(n)$ or $SU(N_2 - n)$ and embed the Standard Model gauge group into the gauged flavor symmetry to realize gauge mediation of SUSY breaking.

Now we examine the gaugino masses in our deformation. In gauge mediation with a superpotential that explicitly breaks R-symmetry, gaugino masses are given by [16]

$$m_\lambda \simeq \frac{g^2}{(4\pi)^2} \frac{F_{X_i}}{X_i} \log(\det M)$$

where $\bar{N}$ is a constant, M is the messenger field mass matrix, and X_i denotes a superfield in the hidden sector with $-F^* = \partial W / \partial X_i$. In our deformation, the form of M depends on which flavor symmetry, $SU(n)$ or $SU(N_2 - n)$, is gauged.

(1) If we gauge the $SU(N_2 - n)$ flavor symmetry and embed the Standard Model group $SU(3) \times SU(2) \times U(1)$ into it, the messenger fields are ρ_2, R, Z_2. Using the notation

$$\Phi = \begin{pmatrix} Y & Z_1 & Z_2 \\ \tilde{Z}_1 & \hat{\Phi} & R \\ \tilde{Z}_2 & \tilde{R} & \Phi_2 \end{pmatrix},$$

where Φ_1 is an $n \times n$ matrix and Φ_2 is an $(N_2 - n) \times (N_2 - n)$ matrix, the mass matrix M is given in the basis $(\rho_2', Z_2, \rho_2'', R)$ by

$$M/h = \begin{pmatrix} \Phi_2 & \mu_1 & hm_3 & 0 \\ \mu_1 & 0 & 0 & 0 \\ hm_3 & 0 & \Phi_2 & \mu_2 \\ 0 & 0 & \mu_2 & 0 \end{pmatrix}.$$

Assuming $m_2 m_3 \gg b\mu_2^2$ and $\mu_1 \gg \langle \Phi_2 \rangle$, we have

$$m_\lambda \simeq \frac{g^2}{(4\pi)^2} \frac{F_{\Phi_2}}{\langle \Phi_2 \rangle} \log(\det M) \simeq \frac{g^2}{(4\pi)^2} \frac{F_{\Phi_2}}{\langle \Phi_2 \rangle} \log \left(\frac{hm_3 \langle \Phi_2 \rangle - \mu_2^2}{hm_2 \langle \Phi_2 \rangle - \mu_2^2} \right)^2$$

where $F_{\Phi_2} = h\mu_2^2$ and $\langle \Phi_2 \rangle = m_2/(hb)$ denotes the expectation value in the metastable vacuum shown in Eq. (16). The squark masses are

$$m_s \simeq \frac{g^2}{(4\pi)^2} \frac{F_{\Phi_2}}{\langle \Phi_2 \rangle} \left(\frac{\langle \Phi_2 \rangle}{\mu_1} \right)$$

where we assumed $\mu_1 \gg \langle\Phi_2\rangle$ and considered $\langle\Phi_2\rangle = m_2/(hb) \gg \mu_2$ as required by the absence of tachyonic messenger fields. This yields gaugino and squark masses of the same order.

Note that the assumptions $m_2 m_3 \gg b\mu_2^2$ and $\mu_1 \gg \langle\Phi_2\rangle$ are easily satisfied for sufficiently large m_3. We verified that these conditions do not affect our vacuum structure in Eqs. (15, 16) nor our vacuum lifetime calculations (where we use Eq. (11) with μ_ϕ and μ replaced by m_2 and μ_2, respectively).

Compared with [3], where the two independent scales must be nearly equal ($m_3 \sim \mu_1$), our study allows a larger viable region for the parameter m_3. As shown in our Landau pole analysis below, large m_3 is favored, enabling a SUSY breaking scale ($h\mu_2^2 \sim 100$ TeV) lower than that obtained in [3].

(2) If we gauge the $SU(n)$ flavor symmetry and embed the Standard Model group into it, the messenger fields are ρ_2, χ, R and Z_1. The mass matrix M is given in the basis (ρ_2, R, χ, Z_1) by

$$M/h = \begin{pmatrix} \Phi_2 & \mu_2 & hm_2 & 0 \\ \mu_2 & 0 & 0 & 0 \\ hm_2 & 0 & Y & \mu_1 \\ 0 & 0 & \mu_1 & 0 \end{pmatrix}.$$

If we assume the F-term of Y is non-zero (for example, taking $k = 0$ with non-zero N_1, so that in Eq. (16) the first diagonal element is $(m_1/b)I_{N_1}$ for $\langle h\Phi\rangle$ and $0I_{N_1}$ for $\langle q\tilde{q}\rangle$) and further assume $m_1 m_3 \gg b\mu_2^2$, we obtain

$$m_\lambda \simeq \frac{g^2}{(4\pi)^2}\frac{F_Y}{\langle Y\rangle}\log(\det M) \simeq \frac{g^2}{(4\pi)^2}\frac{F_Y}{\langle Y\rangle}\log\left(\frac{h^2 m_2 + (hm_3\langle Y\rangle - \mu_2^2)^2}{(hm_3\langle Y\rangle - \mu_2^2)^2}\right)$$

where $F_Y = h\mu_1^2$ and $\langle Y\rangle = m_1/(hb)$ denotes the expectation value in the metastable vacuum. The squark masses are

$$m_s \simeq \frac{g^2}{(4\pi)^2}\frac{F_{\Phi_2}}{\langle\Phi_2\rangle}\left(\frac{\langle Y\rangle}{\mu_1}\right).$$

Thus, in this case, if $\mu_1 \sim \mu_2$ the squark masses can also be of the same order as gaugino masses.

Finally, we check whether our deformation suffers from the Landau pole problem. The mass spectrum can be read from the metastable vacua and depends on which flavor symmetry, $SU(n)$ or $SU(N_2 - n)$, is gauged. We find that our model has a Landau pole problem if we gauge $SU(n)$ symmetry, but is free of this problem if we gauge $SU(N_2 - n)$ symmetry. Below we demonstrate how to avoid the Landau pole when gauging $SU(N_2 - n)$.

When gauging $SU(N_2-n)$ and embedding the Standard Model group into it, the fields ρ_2' and Z_2 have masses of order m_2/b and $O(hm_3)$, respectively. The fields R and ρ_2'' have masses near the scale $h\mu_2$, and the pseudo-moduli Φ_2 has mass similar to the gauginos. For simplicity, we take $m_2/b \sim h\mu_2$ in the following calculation. The beta function coefficients of the $SU(3)$ gauge coupling b_3 are

$$
\begin{aligned}
b_3(\mu_R < m_\lambda) &= b_3^{SM} = -7, \\
b_3(m_\lambda < \mu_R < h\mu_2) &= -3 + N_2 - n, \\
b_3(h\mu_2 < \mu_R < hm_3) &= -3 + N_2 + N_f - N_c, \\
b_3(hm_3 < \mu_R < \Lambda) &= -3 + 2N_f - N_c, \\
b_3(\mu_R > \Lambda) &= -3 + N_c,
\end{aligned}
$$

where μ_R is the renormalization scale. In our analysis we use the definition $\beta_3 \equiv 16\pi^2 d\alpha_3/dt = b_3\alpha_3^2$, and take input parameters $M_{GUT} = 10^{16}$ GeV, $M_Z \simeq 90$ GeV, $m_\lambda \simeq 10^3$ GeV, and $\alpha_3(M_Z) \sim 0.18$.

The $SU(3)$ coupling at the GUT scale is

$$\alpha_3^{-1}(M_{GUT}) \simeq 5.6 - \log(M_Z/m_\lambda) + (N_f - N_c + n)\log(h\mu_2/m_\lambda) + (N_f - N_2)\log(hm_3/(h\mu_2)) + (2N_c - 2N_f)\log(M_{GU}$$

where we take $\Lambda = M_{GUT}$. For example, taking $N_2 - n = 5$, $N_f = 11$, $N_c = 9$, $n = 1$ and $N_1 = N_f - N_2 = 5$, we find that for SUSY breaking scale $h\mu_2 \sim 10^5$ GeV and $hm_3 \geq 10^7$ GeV, the Landau pole can be avoided below the unification scale, i.e., $\alpha_3^{-1}(M_{GUT}) > 1$. Under this condition, proper gaugino masses can be obtained from Eq. (26), and we have verified this is consistent with our other assumptions.

IV. Conclusion

In this work we considered a more general deformation of the ISS model by splitting both TrΦ and TrΦ^2 terms. We found that the corresponding metastable vacua can be sufficiently long-lived and generate proper gaugino masses. In particular, models can avoid the Landau pole problem if we gauge the $SU(N_2-n)$ flavor group and embed the Standard Model group into it.

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Note: Figure translations are in progress. See original paper for figures.