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Metastable supersymmetry breaking in terms of Peccei–Quinn mechanism

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ABSTRACT

Gauge-mediated supersymmetry breaking at metastable vacuum is reconsidered in terms of the Peccei–Quinn (PQ) mechanism. We suggest that for acceptable μ-value generation, such a class of model should involve messenger mass generation via the PQ-breaking process. We then construct a model in context of the Izawa–Yanagida–Intriligator–Thomas superpotential with the aid of the effective *Kähler* coupling induced by an additional super Yang–Mills sector. Therein, the PQ-charged fundamental singlet plays a crucial role in accommodating a sizable μ-value.

1. Introduction

Metastable supersymmetry (SUSY) breaking models such as the Intriligator–Seiberg–Shih (ISS) model [1] provide concise descriptions for gauge mediation [2], [3,4]. However, such a class of model might less definitely determine the hidden–structure, because of no continuous R-symmetry.

In this study, we address whether the Peccei–Quinn (PQ) mechanism [8–11], instead, can reveal the hidden structure. It should be noted that a fairly convincing argument regarding this prospect has already been proposed in that the dynamical SUSY breaking [5–7, 24,25] may coincide with the spontaneous breakdown of PQ symmetry; the authors of [12,13] discussed $Sp(N)$ gauge theory based on pure gravity mediation (PGM) [14–17]. Meanwhile, in [18], the gauge-mediated SUSY breakdown was argued in context of the Kim–Shifman–Vainstein–Zakharov (KSVZ) mechanism [20,21]. Furthermore, the viability of nonthermally produced dark matter was discussed in context of the $Sp(N) \times Sp(M)$ model that might eventually be unified into a larger $Sp(N_C)$ gauge group [19].

Notice that the PQ symmetry, even if only approximate, should adequately constrain the possible form of the superpotential to maintain the so-called theta angle $|\theta| < 10^{-10}$ [26–28]. Throughout this study, we accordingly assume that renormalizable terms respect $U(1)_{PQ}$, and the explicitly PQ-breaking interactions, while maintaining an acceptable $|\theta|$-value, only arise as Planck-suppressed terms. In addition, the discrete R-symmetry is postulated to remain unbroken below the Planck scale.

We then observe that if the mass generation of messengers or (dual) singlets is ruled by the PQ-conserving process as in [4,22], the μ-value (higgsino mass parameter) may be rendered unacceptably small in light of the QCD anomaly. Based on this consideration, we will construct a model based on the Izawa–Yanagida–Intriligator–Thomas (IYIT) superpotential, wherein the PQ-charged (fundamental) singlets realize the relevant mass, consequently making a sizable μ-value emerge in a consistent way.

This paper is organized as follows. In the next section, we argue that for significant μ-generation, the mass generation mechanism for messengers apparently involves a PQ-breaking process. In Sec. 3, we then address an IYIT-type model (assuming $SU(3)$) possessing a super Yang–Mills sector (assuming $SU(2)$) and examine the role of fundamental singlets. We argue that $U(1)_{PQ}$ and the discrete R-symmetry (assuming Z_{6R}) should reduce the superpotential to the O'Raifeartaigh form up to higher dimensional operators. At this stage, it is crucial that an effective *Kähler* coupling, induced by the $SU(2)$ gaugino condensation, plays an important role in accommodating the gauge mediation successfully, while ensuring the (meta)stability of the vacuum.[1] Then, we deduce that for f_a (axion decay constant [33–36]) $\simeq 10^{11-12}$ GeV, the effective $|\mu| \simeq \mathcal{O}(1)$ TeV and minimal supersymmetric standard model (MSSM) soft mass of $\mathcal{O}(1)$ TeV can be generated from the messenger loops, where $m_{3/2} \simeq$ 1–100 MeV. Furthermore, we address

[1] At this point, it should be emphasized that the author in [37] discussed the mechanism of stabilizing the messenger fields by the supergravity (SUGRA) effects, which will be applied to our model.

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the compatibility of such a class of model with the electroweak symmetry breaking (EWSB) in Sec. 4. We there find that the introduction of extra messengers serves as a prescription to evade the μ-B_μ problem [40] (which is encountered by the messenger coupling in Sec. 3), possibly leading to the desirable low-scale phenomenology. In the Appendix, we discuss the Planck-suppressed operators that may explicitly violate $U(1)_{PQ}$. In the final section, the conclusions are drawn.

2. Necessity of the PQ-breaking process for messenger and singlet mass generation

Let us address how the QCD anomaly associated with $U(1)_{PQ}$ influences the ISS gauge-mediation model.

2.1. Implication of QCD anomaly

As discussed in [3], the magnetic $SU(N_F - N)$ theory ($3N/2 > N_F > N \geq 3$) should involve the superpotential terms (up to the dimensionless coefficients):

$$W \supset \frac{\Lambda_{ISS}}{M_{pl}} S^{ij} \bar{f} f + m\Lambda_{ISS} S^{ij} + M\bar{f} f, \tag{1}$$

where S^{ij}, $\bar{f}, f$, and m, M denote the dual singlet, messenger superfields, and mass parameters, respectively. In addition, Λ_{ISS} is the dynamical scale and $M_{pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck scale. Provided that both m and M originate from the $U(1)_{PQ}$-conserving process such as $\langle \text{Tr } W_\alpha^2 \rangle / M_{pl}^2$ (W_α is the gauge supermultiplet of the Yang–Mills sector [4]) or $\langle \Phi\Phi \rangle / M_{pl}$ via the $Sp(N)$-quantum-modified constraint $Pf(\Phi^k \Phi^l) = (\Phi\Phi)_a + \Phi\Phi = \Lambda_{Sp(N)}^{2(N+1)}$ (Φ^k is the $Sp(N)$-fundamental, while a represents the index for $SO(5)$-flavor symmetry)[2] [22,23], $f\bar{f}$ as well as S^{ij} can apparently carry no PQ-charge (see Appendix for a more rigid argument), which implies that the messenger fermion loops do not contribute to the $U(1)_{PQ}[SU(3)_{color}]^2$ anomaly. Hence, such a class of model might be subject to the Dine–Fischler–Srednicki–Zhitnitski (DFSZ) model [29,30]. In what follows, we evaluate the μ-generation under such a circumstance.

2.2. μ-generation

- **Higgs-messenger coupling**

Given that the up/down type of the Higgs superfields H_u, H_d couple with the messengers (denoted $\bar{f}_1, f_1, \bar{f}_2, f_2$) as follows:

$$W \supset H_u f_1 \bar{f}_2 + H_d \bar{f}_1 f_2, \tag{2}$$

the effective μ-term is expected to appear through the one-loop corrected *Kähler* term that takes the form of[3]

$$\Delta K \simeq \frac{1}{16\pi^2} \frac{S^{ij\dagger} H_u H_d}{M} + h.c., \tag{4}$$

after the messengers have been integrated out [38–40]. However, we point out that the PQ invariance forbids the appearance of ΔK,[4] as it is supposed that

$$q_{S^{ij}} = 0, \;\; q_{H_u} + q_{H_d} \neq 0, \tag{5}$$

where $q_{S^{ij}}$ is the PQ-charge of S^{ij}, while q_{H_u}, q_{H_d} are those of H_u, H_d, respectively, that add up to a non-vanishing value via the DFSZ mechanism. Hence, it may be difficult for this sort of model to admit such couplings.

- **Giudice–Masiero mechanism**

It is readily found from Eq. (5) that the *Kähler* potential cannot include $H_u H_d + h.c.$ and, therefore, the GM mechanism [31,32] does not work here.

- **PQ-charged magnetic quark coupling with Higgs bilinear**

The third candidate of μ-generation can be realized from the following interaction (the other couplings should be forbidden, as mentioned in the Appendix):

$$W \supset \frac{Q_{ISS}^1 \cdots Q_{ISS}^N}{M_{pl}^{N-1}} H_u H_d$$
$$\longrightarrow \frac{\Lambda_{ISS}^N}{M_{pl}^{N-1}} \left(\frac{q}{\Lambda_{ISS}} \right)^{N_f - N} H_u H_d. \tag{6}$$

Here, Q_{ISS} is the electric quark. Meanwhile, q is the magnetic quark that is allowed to be PQ-charged,[5] which might develop the VEV as follows:

$$\langle q \rangle \simeq \sqrt{m\Lambda_{ISS}}. \tag{8}$$

Then, for $N_F > N \geq 3$,

$$|\mu| \lesssim 1 \text{ MeV}, \tag{9}$$

because it is supposed that $m \simeq 10^{7-8}$ GeV and $\Lambda_{ISS} \simeq 10^{11-12}$ GeV [3,4]. To summarize, the representative μ-generating mechanism is much unlikely to accommodate the sizable μ-value in a consistent manner with the PQ mechanism.[6]

3. $SU(3) \times SU(2) \times U(1)_{PQ} \times Z_{6R}$ model

Here, we attempt to construct the $SU(3) \times U(1)_{PQ}$ with a $SU(2)$ Yang–Mills sector model, using a rather minimal content of messengers for the time being. Z_{6R} symmetry is also imposed. The main distinction from the ISS-type is that several fundamental singlets are present in this model.

3.1. Outline

Introducing two kinds of messenger superfields, we present the content of the fields in Table 1 ($i, j = 1$–4):

Let the $U(1)_{PQ}$- and Z_{6R}-charges be assigned as shown in Table 2.

---

[2] It is understood that $Pf(\Phi^k\Phi^l)$ (and hence $\langle\Phi\Phi\rangle$), as a whole, does not carry the PQ-charge, so the unwanted $U(1)_{PQ}[Sp(N)]^2$ anomaly can be eliminated.

[3] If multiple messengers are present, S could be [38]

$$\Delta K \simeq \frac{1}{16\pi^2} \left[\frac{S^{ij\dagger} S^{kl\dagger}}{|S^{ij}|^2 - |S^{kl}|^2} \cdot \log \left| \frac{S^{ij}}{S^{kl}} \right|^2 \right] H_u H_d + h.c. \tag{3}$$

[4] This is initially because either or both the terms in Eq. (2) cannot be allowed.

[5] The PQ-charge must be assigned as follows:

$$\bar{q}_i = -q_j = r \neq 0, \text{ for } \forall\, i, j, \tag{7}$$

as all of S^{ij} are PQ-neutral.

[6] The gravity-mediation scheme leaves the ISS model compatible with the PQ mechanism. Particularly, the larger f_a [41] ($\simeq \langle q \rangle$) can embody the nonthermally produced sparticle mass based on WIMPzillas [42–45].

Table 1
Representation of matter.

	$SU(3)$	$SU(2)$	$SU(5)_{\rm GUT}$
$Q^i,\ \bar{Q}^j$	$3,\ \bar{3}$	$1,\ 1$	$1,\ 1$
$Y_{ij},\ Z,\ \bar{Z}$	$1,\ 1,\ 1$	$1,\ 1,\ 1$	$1,\ 1,\ 1$
$f_1,\ \bar{f}_1$	$1,\ 1$	$1,\ 1$	$5,\ \bar{5}$
f_2	1	1	1

Table 2
Charge assignment.

	$Z,\ \bar{Z},$	$f_1\ \bar{f}_1\ f_2$
$U(1)_{PQ}$	$1,\ -1$	$-\frac{1}{2},\ -\frac{1}{2},\ \frac{1}{2}$
Z_{6R}	$2,\ 2$	$0,\ 0,\ 0$
	$Q^1Q^2Q^3,$	$\bar{Q}^1\bar{Q}^2\bar{Q}^3$
$U(1)_{PQ}$	$-1,$	1
Z_{6R}	$0,$	0

For the other superfields, given that the relations hold as below:

$$1 < m^{11} < m^{12} < m^{13} < m^{21} < \frac{3}{2},$$
$$m^{14} = -1,\ m^{44} = 0,\ y_{kl} = -m^{kl},\ (k,l = 1-3) \tag{10}$$

it follows that[7]

$$m^{34} < m^{31} < m^{32} < m^{33} < -2,\ -1 < m^{24} < -\frac{1}{2},$$
$$(m^{21}) < m^{22} < m^{23} < 2 < m^{41} < m^{42} < m^{43}. \tag{12}$$

Each small letter corresponds to the PQ-charge of $Q^i\bar{Q}^j, Y_{ij}$, respectively. In addition, the R-charge is set to $\mathcal{R}(Y_{ij}) = 2, \mathcal{R}(Q^i) = \mathcal{R}(\bar{Q}^j) = 0$.[8] Notice that no $U(1)_{PQ}[SU(3)]^2$ anomaly appears. Meanwhile, the messenger fermion loops contribute to the QCD anomaly. (Later, we mention the total QCD anomaly coefficient combined with the PQ-charge of standard model (SM) quarks.)

The most general superpotential consistent with all the symmetries is then given by

$$W = Y_{kl}\bar{Q}^kQ^l + \frac{1}{M_{pl}}ZQ^1Q^2Q^3 + \frac{1}{M_{pl}}\bar{Z}\bar{Q}^1\bar{Q}^2\bar{Q}^3 + Z\bar{Q}^4Q^1 + \frac{Tr\left(W_\alpha^2\right)}{M_{pl}^2}\bar{Q}^4Q^4 + \frac{Tr(W_\alpha^2)Q^2Q^3Q^4}{M_{pl}^3} - \lambda_1 Z\bar{f}_1f_1 - \lambda_2\bar{Z}f_2f_2 + \cdots, \tag{13}$$

where W_α is the $SU(2)$-gauge supermultiplet, that carries the vanishing PQ-charge and $\mathcal{R}(W_\alpha) = 1$. Moreover, the dimensionless coefficients of $\mathcal{O}(0.1–1)$ are omitted, except for the messenger coupling, which will be discussed later. The ellipse denotes the Planck-suppressed operators–$\left(Z\bar{Z}\right)^2/M_{pl}$ etc.. In the presence of these terms, supersymmetric vacua potentially exist, whereas the metastable SUSY breaking cannot be destabilized, because of their small effects (see Sec. 3.3 for details).

Through the gaugino condensation (whose scale is denoted Λ'), $\bar{Q}^4, Q^4$ become massive, and thus, after integrating them out, one obtains the one-loop corrected *Kähler* term:

$$\Delta K \simeq -\frac{1}{16\pi^2}\frac{|Z|^4}{M_*^2}, \tag{14}$$

where $M_* = \Lambda'^3/M_{pl}^2$. Then, below the $SU(3)$ dynamical scale (denoted Λ, and we suppose $\Lambda \ll M_*$), composite mesons and baryon/anti-baryon should be formed:

$$\frac{Q^k\bar{Q}^l}{\Lambda} \to M^{kl},\ \frac{Q^1Q^2Q^3}{\Lambda^2} \to B,\ \frac{\bar{Q}^1\bar{Q}^2\bar{Q}^3}{\Lambda^2} \to \bar{B}. \tag{15}$$

Accordingly, the effective superpotential of $SU(3)$ dynamics is as follows[9,10]:

$$W = \Lambda Y_{kl}M^{kl} + m_1ZB + m_2\bar{Z}\bar{B},\ (k,l = 1–3), \tag{16}$$

with the quantum modified constraints [6]:

$$\frac{\det M^{kl}}{\Lambda} - B\bar{B} = \Lambda^2, \tag{17}$$

where

$$m_1 \simeq m_2 \simeq \frac{\Lambda^2}{M_{pl}}. \tag{18}$$

In what follows, we assume $m_1 = m_2 \equiv \bar{m}$, for simplicity.

3.2. Implementation of gauge mediation

Following [46], we define $Z_\pm$ and $B_\pm$ as:

$$Z_+ = \frac{Z+\bar{Z}}{\sqrt{2}},\ Z_- = \frac{Z-\bar{Z}}{\sqrt{2}},$$
$$B_+ = \frac{B+\bar{B}}{\sqrt{2}},\ B_- = \frac{B-\bar{B}}{\sqrt{2}}. \tag{19}$$

Eliminating B_- in terms of Eq. (17), one rewrites Eq. (16) as follows:

$$W = \Lambda Y_{kl}M^{kl} + \bar{m}Z_+B_+ + \bar{m}Z_-\left(\Lambda^2 - \frac{\det M}{\Lambda} + B_+^2\right)^{1/2}. \tag{20}$$

With Eq. (14), the scalar potential is then approximately given by[11]:

$$V \supset \bar{m}^2|B_+|^2 + \bar{m}^2|Z_+|^2 + \frac{\bar{m}^2\Lambda^2}{16\pi^2M_*^2}|Z_-|^2, \tag{21}$$

after integrating out Y_{kl}, M^{kl}. Hence the metastable SUSY-breaking vacuum should be at

$$B_+ \simeq 0,\ Z_\pm \simeq 0,\ \ f_1 = \bar{f}_1 = f_2 = 0. \tag{22}$$

---

[7] Each assignment is as follows:

$$m^{22} = m^{12} + m^{21} - m^{11},\ m^{23} = m^{13} + m^{21} - m^{11},$$
$$m^{31} = m^{11} - m^{12} - m^{13} - m^{21},\ m^{32} = -m^{13} - m^{21},$$
$$m^{33} = -m^{12} - m^{21},\ m^{24} = -m^{11} + m^{21} - 1,$$
$$m^{34} = -m^{12} - m^{13} - m^{21} - 1,$$
$$m^{41} = m^{11} + 1,\ m^{42} = m^{12} + 1,\ m^{43} = m^{13} + 1. \tag{11}$$

[8] This assignment yields $Z_{6R}[SU(3)]^2$ anomaly, which however disappears owing to the mass generation of $Q^4, \bar{Q}^4$.

[9] $Z\bar{B}, \bar{Z}B$ or $Y_{ij}B, Y_{ij}\bar{B}$, despite the Planck-suppressed terms, must be forbidden by the PQ mechanism, as mentioned in the Appendix.

[10] After integrating out $Q^4, \bar{Q}^4$, the effective superpotential of $\kappa ZQ^1Q^2Q^3/M_{pl}, (|\kappa| = \mathcal{O}(0.1-1))$ is generated, which does not significantly influence on our further analysis.

[11] We ignore the mixing term between Z_+ and Z_-, and omit the B_+-loop-induced mass term of Z_-, because these are small corrections.

The SUGRA correction, however, induces the term [37]:

$$-3m_{3/2}\bar{m}\Lambda\left(Z_- + Z_-^*\right), \tag{23}$$

and thus, Z_- develops the VEV[12,13,14]:

$$Z_- \simeq \frac{48\pi^2 m_{3/2} M_*^2}{\bar{m}\Lambda}. \tag{25}$$

In addition, the F-term condenses:

$$F_{Z_-} \simeq \bar{m}\Lambda, \tag{26}$$

where $m_{3/2} = \bar{m}\Lambda/\sqrt{3}M_{pl}$. Consequently, the messenger scale (denoted Λ_{mess}) and f_a can be obtained:

$$\Lambda_{mess} \simeq \frac{\bar{m}^2\Lambda^2}{48\pi^2 m_{3/2} M_*^2}, \quad f_a \simeq \Lambda, \tag{27}$$

and the SUSY-breaking scale is as follows:

$$\Lambda_{SUSY} \simeq \sqrt{\bar{m}\Lambda}. \tag{28}$$

3.3. Physical viability of the vacuum

Let us examine the viability of this vacuum. The tachyonic mass of the messenger can be avoided by

$$\sqrt{2}|F_{Z_-}| < \lambda_1 |Z_-|^2, \quad \sqrt{2}|F_{Z_-}| < \lambda_2 |Z_-|^2. \tag{29}$$

In addition, the messenger loop generates the extra scalar potential that contributes the negative squared mass to $Z_\pm$:

$$V \supset \frac{5\lambda_1^2|F_{Z_-}|^2}{64\pi^2}\log\frac{|Z_+ + Z_-|^2}{2\Lambda_{CUT}^2} + \frac{\lambda_2^2|F_{Z_-}|^2}{64\pi^2}\log\frac{|Z_+ - Z_-|^2}{2\Lambda_{CUT}^2}, \tag{30}$$

where Λ_{CUT} is some cutoff scale. Hence, λ_1, λ_2 place a limitation on

$$\frac{3\sqrt{2}}{\left(48\pi^2\right)^2}\cdot\frac{M_{pl}\Lambda^3}{M_*^4} < \lambda_1 < \frac{48\pi^2}{\sqrt{15}}\cdot\frac{M_*}{M_{pl}},$$
$$\frac{3\sqrt{2}}{\left(48\pi^2\right)^2}\cdot\frac{M_{pl}\Lambda^3}{M_*^4} < \lambda_2 < \frac{48\pi^2}{\sqrt{3}}\cdot\frac{M_*}{M_{pl}}. \tag{31}$$

Altogether, if $\Lambda = \mathcal{O}(10^{11\text{–}12})$ GeV and $\Lambda' = \mathcal{O}(10^{16})$ GeV, one finds that

$$m_{soft} \simeq \frac{\alpha_{\rm MSSM}}{4\pi}\cdot\Lambda_{mess} = \mathcal{O}(1)\ \text{TeV},$$
$$m_{3/2} = \mathcal{O}(1\text{–}100)\ \text{MeV}, \ f_a = \mathcal{O}(10^{11\text{–}12})\ \text{GeV}, \tag{32}$$

and

$$10^{-6} \lesssim \lambda_1, \lambda_2 \lesssim 10^{-3}. \tag{33}$$

This simultaneously implies that the flavor-changing processes can be significantly suppressed because

$$m_{soft} \gg 100\cdot m_{3/2}. \tag{34}$$

Next, to discuss the longevity of the metastable vacuum, we should focus on some Planck-suppressed interactions. Among them is the leading term:

$$\Delta W = \frac{\left(Z\bar{Z}\right)^2}{4M_{pl}}, \tag{35}$$

which gives rise to the appearance of several supersymmetric vacua. Noting that each stationary point is at

$$Z_{SUSY} \simeq \bar{Z}_{SUSY} \simeq -\Lambda \tag{36}$$

and both of them are the furthest from metastable vacuum,[15] the following conditions are necessary to suppress the tunneling into the supersymmetric vacuum [47]:

$$\frac{Z_{SUSY}^4}{\Lambda_{SUSY}^4}, \ \frac{\bar{Z}_{SUSY}^4}{\Lambda_{SUSY}^4} > 400, \tag{38}$$

which is easily satisfied.

In addition to that, there may exist another non-supersymmetric vacuum at $f_1\bar{f}_1 \simeq \bar{m}\Lambda/\lambda_1$, $f_2 \simeq \sqrt{\bar{m}\Lambda/\lambda_2}$. Even if such a state has a lower energy than our viable vacuum, the tunneling rate (derived from the thin-wall approximation) is significantly suppressed as it holds that

$$\bar{m}\Lambda\cdot\langle Z_-\rangle > F_{Z_-}^{3/2} \tag{39}$$

in our model. Here $\langle Z_-\rangle$ is the VEV of Eq. (25).

3.4. Net QCD anomaly and effective μ

Let us verify that the PQ mechanism normally works. The Higgs superfields, which belong to $5, \bar{5}$ of $SU(5)_{GUT}$, are assigned the $U(1)_{PQ}$- and Z_{6R}-charges as $H_u(0,2)$, $H_d(0,2)$ (the Z_{6R}-invariance forbids the bare μ-term). Provided that the SM quarks should be correspondingly PQ-charged, the net QCD anomaly can be written as follows:

$$-\frac{1}{2}\cdot\frac{g_s^2}{32\pi^2}\cdot\frac{a}{f_a}G_{\mu\nu}\tilde{G}_{\mu\nu}, \tag{40}$$

where g_s, a, and $G_{\mu\nu}$ are the $SU(3)_{color}$ coupling constant, axion, and SM gluon, respectively. The factor $-1/2$ results from the messenger multiplets, while the SM quarks give no contribution [12]. Then, given that the following terms can be among the superpotential:

$$W \supset \lambda_u H_u \bar{f}_1 f_2 + \lambda_d H_d f_1 f_2, \tag{41}$$

the effective μ-term is generated:

$$\mathscr{L} \supset \int d^2\theta\, \frac{\lambda_u\lambda_d}{16\pi^2}\Lambda_{mess}H_uH_d + h.c., \tag{42}$$

---

[12] Note that without the effective *Kähler* term of Eq. (14), Z_- would exhibit the VEV of $\mathcal{O}(M_{pl})$ via the SUGRA correction.

[13] Its VEV as well as the $SU(2)$-gaugino condensation break Z_{6R} down to Z_{2R} symmetry.

[14] Accordingly, it follows that:

$$Z \simeq \frac{48\pi^2 m_{3/2} M_*^2}{\sqrt{2}\bar{m}\Lambda}, \ \bar{Z} \simeq -\frac{48\pi^2 m_{3/2} M_*^2}{\sqrt{2}\bar{m}\Lambda}. \tag{24}$$

[15] Aside from the SUGRA correction, near the origin of $Z_\pm$, the scalar potential is dominated by the following terms:

$$V \supset \bar{m}^2|Z_+|^2 + \frac{\bar{m}^2\Lambda^2}{16\pi^2 M_*^2}|Z_-|^2, \tag{37}$$

which is (meta)stabilized at $Z_\pm = 0$. Z_- is then shifted by the SUGRA–induced potential.

Table 3
Charge assignment.

	$f_1,\ \bar{f}_1,\ f_2,\ f',\ \bar{f}'$	$H_{u,d}$	$Z,\ \bar{Z}$
$U(1)_{PQ}$	$-\frac{1}{2},\ -\frac{1}{2},\ 0,\ \frac{1}{2},\ \frac{1}{2}$	$\frac{1}{2}$	$1, -1$
Z_{6R}	$0,\ 0,\ -1,\ -2,\ -2$	3	$2,\ 2$

where λ_u, λ_d are of $\mathcal{O}(0.1-1)$, and it follows that

$$|\mu| = \mathcal{O}(1)\ \text{TeV}. \tag{43}$$

4. The μ-B_μ problem addressed; aiming at successful EWSB

For more comprehensive analysis, we here address whether this sort of model could be favored by the EWSB at weak scale. While the messenger content of $(5+\bar{5})$ and 1 is sufficient to provide $|\mu|$ as well as the MSSM soft mass at TeV scale, Eq. (41) generally exhibits a much larger $|B_\mu|$ than $|\mu|^2$, whereas $m^2_{H_{u,d}}$ should vanish up to 1-loop order of the messengers. This implies that the electroweak symmetry spontaneously breaks down at the messenger scale, and what is worse, the Higgs potential is destabilized along the D-flat direction (called the "μ-B_μ problem" [40]). We then follow [38] to present an approach for resolving this problem.

4.1. Extended model to suppress the B_μ-term

Let us suppose that an additional pair of messengers (denoted $f', \bar{f}'$ which belong to 5 and $\bar{5}$ of $SU(5)_{GUT}$, respectively) is introduced and the relevant superpotential takes the form of:

$$\begin{aligned} W \supset & \lambda_u H_u \bar{f}_1 f_2 + \lambda_d H_d f_1 f_2 \\ & - \lambda Z \bar{f}_1 f_1 + \lambda' \frac{\left[Tr\left(W_\alpha^2\right)\right]^2}{M_{pl}^5} f_2^2 \\ & + \lambda'' \frac{\left[Tr\left(W_\alpha^2\right)\right]^2}{M_{pl}^5} \bar{f}' f_1 + \lambda''' \frac{\left[Tr\left(W_\alpha^2\right)\right]^2}{M_{pl}^5} \bar{f}_1 f', \end{aligned} \tag{44}$$

which is the most general form consistent with the charge assignment up to higher-order terms shown in Table 3.[16]

Accordingly, at 1-loop level of the messengers, the soft mass parameters could be generated[17]:

$$\begin{aligned} & \mu \simeq \frac{\lambda_u \lambda_d}{16\pi^2} \cdot \frac{\lambda F_{Z_-}}{M'}, \quad m^2_{H_{u,d}} \simeq \frac{\lambda^2_{u,d}}{16\pi^2} \left(\frac{\lambda F_{Z_-}}{M'}\right)^2, \\ & B_\mu \simeq \frac{\lambda_u \lambda_d}{16\pi^2} \cdot \frac{\lambda \langle Z \rangle}{M'} \left(\frac{\lambda F_{Z_-}}{M'}\right)^2 \simeq \lambda \lambda_u \lambda_d \left(\frac{\lambda F_{Z_-}}{M'}\right)^2, \end{aligned} \tag{46}$$

where $M' \equiv \bar{\lambda}[Tr\left(W_\alpha^2\right)]^2/M_{pl}^5$ with $\bar{\lambda} \equiv \lambda' \simeq \lambda'' \simeq \lambda'''$. Then, for $\lambda \ll \mathcal{O}(1/16\pi^2)$, one obtains the desirable relations:

$$\begin{aligned} & 2\,|\mu|^2 + m^2_{H_u} + m^2_{H_d} \gg 2\left|B_\mu\right|, \\ & \left(m^2_{H_u} + |\mu|^2\right)\left(m^2_{H_d} + |\mu|^2\right) \gg \left|B_\mu\right|^2. \end{aligned} \tag{47}$$

Hence, we may deduce that the electroweak symmetry is left unbroken at the messenger scale.

[16] For this case, the QCD anomaly coefficient adds up to $-3/2$.

[17] Here, it is expected that the B_μ-term should be dominantly generated by:

$$K_{eff} \supset \mathcal{O}(1) \cdot \frac{\lambda_u \lambda_d}{16\pi^2} \cdot \frac{\lambda Z^\dagger}{M'} \frac{|\lambda Z|^2}{M'^2} H_u H_d + h.c.. \tag{45}$$

Let us roughly examine the possibility of the successful EWSB. To this end, it must hold that at weak scale:

$$\left(m^2_{H_u} - \Delta m^2_{H_u} + |\mu|^2\right)\left(m^2_{H_d} + |\mu|^2\right) \simeq \left|B_\mu\right|^2, \tag{48}$$

where

$$\Delta m^2_{H_u} \simeq \frac{3h_t^2}{4\pi^2} m^2_{\tilde{t}} \log \frac{M'}{m_{\tilde{t}}} \tag{49}$$

is the loop correction from the top squark (whose mass is denoted $m_{\tilde{t}}$), and h_t is the top Yukawa coupling constant. By using $|\mu|^2 \ll m^2_{H_{u,d}}$ expected from Eq. (46) and $h_t \simeq 1/\sin\beta$, Eq. (48) is rewritten as:

$$16\pi^2 \lambda \simeq \left(1 - \frac{\mathcal{O}(10^{-2})}{\lambda_u^2 \sin^2\beta}\right)^{1/2}, \tag{50}$$

which can be made possible for $\lambda \ll 1/16\pi^2$ and $\lambda_u = \mathcal{O}(0.1-1)$, depending on $\tan\beta$.[18]

- **(Meta)stability of the vacuum**

Before ending this section, we verify the (meta)stability of the vacuum. To avoid the tachyonic mass of $f_1, \bar{f}_1, f', \bar{f}'$, their squared mass matrix (denoted $\mathcal{M}^2$) needs to be positive-definite:

$$\det \mathcal{M}^2 > 0. \tag{51}$$

Because of $\lambda \langle Z \rangle \ll M'$ (which follows from $\lambda \ll 1/16\pi^2$), we eventually require that:

$$\lambda F_{Z_-} < M'^2, \tag{52}$$

which is expected to hold, where $\lambda \lesssim 10^{-4}$ and $\Lambda = 10^{11\text{–}12}$ GeV, $\Lambda' = \mathcal{O}(10^{16})$ GeV.[19] The messenger loop yields the quadratic term of Z in the scalar potential:

$$\Delta V = \frac{5\,|\lambda|^4}{192\pi^2} \cdot \frac{\left|F_{Z_-}\right|^2}{M'^2} |Z|^2, \tag{54}$$

which only causes a slight shift of $\langle Z_- \rangle$ from the stationary point shown in Eq. (25). Besides, we evaluate the tunneling rate into the nonsupersymmetric vacuum at $\bar{f}_1 f_1 \simeq F_{Z_-}/\lambda$, $f_2 = 0$ and $f' = \lambda \langle Z \rangle f_1/M'$, $\bar{f}' = \lambda \langle Z \rangle \bar{f}_1/M'$. In case this state lowers the vacuum energy, it should be required that:

$$\frac{M' F_{Z_-}}{\lambda} > F_{Z_-}^{3/2}, \tag{55}$$

which is also satisfied in this model.

[18] B_μ evolves according to the renormalization group equation [49,50], whose effect is, however, less significant because of the larger $B(\equiv B_\mu/\mu)$ at the messenger scale than the wino mass or the A-coefficient involved.

[19] Note that the MSSM soft mass is approximately given by:

$$m_{soft} \simeq \frac{\alpha_{MSSM}}{4\pi} \cdot \frac{\lambda F_{Z_-}}{M'}. \tag{53}$$

Hence, μ, m_{soft} of $\mathcal{O}(1)$ TeV can be embodied as well in the parameter region of our interest.

5. Conclusions

The plausibility of SUSY breaking beyond the weak scale has recently been argued based on the prospect of its coincidence with the spontaneous breakdown of $U(1)_{PQ}$. Following this, we discussed the influence of the PQ mechanism on the metastable gauge-mediation model. First, we confirmed that if the $SU(N)$ hyper quark and messenger acquire their masses via the $U(1)_{PQ}$-conserving process, as in the ISS model [4,22], then the PQ mechanism might be described in context of the DFSZ model (see Appendix for detailed discussion). This implies that the Higgs-messenger interaction may be forbidden and the GM mechanism may not be allowed. Eventually, we found that for the ISS model, the vacuum expectation value (VEV) of the magnetic quark only manages to yield $|\mu| \lesssim 1$ MeV. Subsequently, we addressed an IYIT-type model, which is representative of the fundamental singlet-possessing model. We then showed that via $U(1)_{PQ}$ and Z_{6R}, the "$U(1)_R$*-violating*" operators must be Planck-suppressed, enabling the metastable SUSY-breaking vacuum to emerge. At this stage, the effective *Kähler* coupling, induced by the $SU(2)$-gaugino condensation, provided the appropriate mass for messengers via the DSB and the PQ-breaking process. In consequence, we obtained the desirable MSSM mass spectra as well as $|\mu| \simeq 1$ TeV, where $m_{3/2} = \mathcal{O}(1\text{–}100)$ MeV and $f_a = \mathcal{O}(10^{11-12})$ GeV. It should be emphasized that our deduction should be valid even when only the approximate $U(1)_{PQ}$ remains unbroken below the Planck scale.

In the study, we did not assume global symmetries except for the approximate $U(1)_{PQ}$, and therefore, no unwanted Nambu–Goldstone boson (NGB) appears.

Finally, we mentioned the possibility of weak scale EWSB in context of this sort of model. For its achievement, the additional content of messengers might serve a crucial role, even though the tuning problem among relevant parameters is still left unresolved.

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Appendix A

The global $U(1)_{PQ}$ symmetry, even though not exact, should adequately constrain the presence of higher-order couplings so that $|\theta|$ can settle down to a small value less than 10^{-10} [26–28]. Based on this consideration, one might even assign the correct PQ-charge to the Planck-suppressed terms.

- **For ISS-type**

First, $\bar{Q}^i_{ISS} Q^j_{ISS}$ and $\bar{f} f$ are addressed. The former is associated with an interaction in the Lagrangian density:

$$\mathcal{L} \supset m \bar{\psi}^i_{\bar{Q}_{ISS}} \psi^j_{Q_{ISS}}, \tag{56}$$

which develops a VEV of $m\Lambda^3_{ISS}$. (Here, $\bar{\psi}^i_{\bar{Q}_{ISS}}$ and $\psi^j_{Q_{ISS}}$ are the fermion components of $\bar{Q}^i_{ISS}$ and Q^j_{ISS}, respectively.) Hence, for $m \sim 10^{7-8}$ GeV and $\Lambda_{ISS} \sim 10^{11-12}$ GeV, such a term (S^{ij} as well) cannot violate $U(1)_{PQ}$, otherwise, $|\theta|$ would be too large.

For the latter, if $\bar{f} f$ carries a nonvanishing PQ-charge r_{pq}, the axion would be present in the following scalar potential term:

$$V \supset \left| \frac{\Lambda_{ISS}}{M^{pl}} \bar{f} f \exp \frac{i r_{pq} a}{f_a} + m \Lambda_{ISS} \right|^2$$

$$\longrightarrow F \left[\left(-f_1^* f_1 + f_2^* f_2 \right) \cos \frac{r_{pq} a}{f_a} \right] + iF \left[\left(f_1 f_2 - f_1^* f_2^* \right) \sin \frac{r_{pq} a}{f_a} \right], \tag{57}$$

where $F = m\Lambda^2_{ISS}/M_{pl}$ and f_1, f_2 denote the mass eigenstates of $f, \bar{f}$, whose squared masses are $M^2 - F, M^2 + F$, respectively. After integrating the messengers up to the 1-loop, the following is obtained:

$$V \supset \frac{F^2}{16\pi^2} \cos \frac{r_{pq} a}{f_a} + \frac{F^2}{64\pi^2} \cos \frac{2 r_{pq} a}{f_a} + (\text{subdominant terms}), \tag{58}$$

which would lead to too large a $|\theta|$ value.

Through the similar discussion, the explicitly PQ-breaking coupling such as $\frac{\Lambda_{ISS} S^{ij}}{M_{pl}} H_u H_d$ or $\frac{\bar{q}_i q_j}{M_{pl}} H_u H_d$ must be forbidden because the tree-level B_μ generates the corrected term via the Higgs boson loops:

$$V \supset \frac{B_\mu^2}{16\pi^2} \cos \frac{r_{pq} a}{f_a} - \frac{B_\mu^2}{64\pi^2} \cos \frac{2 r_{pq} a}{f_a} + (\text{subdominant terms}), \tag{59}$$

where

$$B_\mu = \frac{m \Lambda^2_{ISS}}{M_{pl}} \quad \text{or} \quad \frac{m \Lambda_{ISS} \left\langle S^{ij} \right\rangle}{M_{pl}} \tag{60}$$

and [3,48]

$$\left\langle S^{ij} \right\rangle \simeq \frac{\Lambda^2_{ISS}}{M_{pl}}. \tag{61}$$

Here, r_{pq} is the PQ-charge carried by the $H_u H_d$ term.

- **For IYIT type**

For this case, the Planck-suppressed couplings $Z\bar{B}, YB$ (or $\bar{Z}B, Y\bar{B}$), which might violate $U(1)_{PQ}$, should be addressed. Considering the following set of superpotential terms:

$$W \supset \frac{\Lambda^2}{M_{pl}} ZB + \frac{\Lambda^2}{M_{pl}} Z\bar{B}, \tag{62}$$

the following scalar couplings can be obtained:

$$\mathcal{L} \supset \frac{\Lambda^4}{M^2_{pl}} B^* \bar{B} + h.c., \tag{63}$$

which acquires the VEV of $\mathcal{O}(\Lambda^6/M^2_{pl})$.

For YB, if the following term is among the superpotential:

$$W \supset \frac{\Lambda^2}{M_{pl}} YB. \tag{64}$$

Thus, the scalar potential, receiving the SUGRA correction

$$V \supset \Lambda^2 |Y|^2 - 3 \frac{m_{3/2} \Lambda^3}{M_{pl}} \left(Y + Y^* \right), \tag{65}$$

yields $Y \simeq \Lambda^4/M^3_{pl}$. Combined with $W \supset \dfrac{\Lambda^2}{M_{pl}} ZB$, it is deduced that

$$V \supset \frac{\Lambda^4}{\sqrt{2} M^2_{pl}} Y Z^* + h.c., \tag{66}$$

whose VEV is $\mathcal{O}(100 \cdot \Lambda^8 \Lambda'^6 / M_{pl}^{10})$. Note that YB has the same PQ-charge as YZ^*. We thus conclude that for $\Lambda = \mathcal{O}(10^{11\text{–}12})$ GeV and $\mathcal{O}(10^{16})$ GeV, the Planck-suppressed couplings involved should not be PQ-charged.

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