A new scheme for NMSSM in gauge mediation

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ABSTRACT: We propose a new framework for the next-to-minimal supersymmetric standard model (NMSSM) in gauge mediation, where in general the correct electroweak symmetry breaking (EWSB) is difficult to be explained. The difficulty is caused by the absence of a soft supersymmetry (SUSY) breaking mass for the NMSSM singlet S. In our framework, S is a meson in a hidden QCD. This QCD is responsible for the dynamical SUSY breaking, forming S, and the soft SUSY breaking mass for S, which is a key to explain the correct EWSB: all the ingredients for successful phenomenology originate from the common dynamics. From the requirement of the successful EWSB, the low-scale SUSY breaking around 100–1000 TeV is predicted. This is favored to avoid the large fine-tuning.

Contents

1 Introduction

Supersymmetry (SUSY) is a promising candidate beyond the standard model (SM). Dynamical SUSY breaking gives a beautiful explanation for the question why the electroweak scale is so small compared to the Planck scale or the unification scale [1] (for a review, see [2]) and the gauge mediation mechanism [3–5] (for early attempts, see [6–10]) (for reviews, see e.g. [11, 12]) can naturally explain why dangerous flavor-changing neutral currents are highly suppressed.[1]

However, the observed Higgs mass, 125 GeV, may conflict to a minimal realization of such a scenario [19]. To obtain the observed Higgs mass in the minimal supersymmetric standard model (MSSM), a significant radiative correction from top/stop loops [20–22] are required, then, in the gauge mediation scenario, sparticle masses should be very large because there is no large stop trilinear coupling.

The simplest modification would be to add a SM singlet chiral superfield S. In such a model, the SH_uH_d superpotential interaction provides an additional F-term contribution to the Higgs potential and it can push the lightest Higgs boson mass up without large stop contributions. This is known as the next-to-minimal supersymmetric standard model (NMSSM) and have been widely investigated so far (for reviews, see e.g. [23, 24]).

Although the additional F-term contribution to the Higgs potential is appealing, the viable parameter space is restricted because large singlet couplings, e.g., SH_uH_d tends to

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[1]In this sense, SUSY models explaining the muon $g - 2$ anomaly based on gauge mediation [13–18] are more convincing than those based on gravity mediation. This is because the light slepton and chargino/neutralino are always required.

blow up at high energy. Furthermore, such a singlet extension of the gauge mediation scenario is not straightforward: since the singlet does not have any SM gauge charge, the SUSY breaking hardly mediates to the singlet sector, which makes it difficult to achieve the correct electroweak symmetry breaking (EWSB); therefore, further extensions, e.g., introducing extra vector-like matter fields coupling to S [25–27],[2] and the coupling between S and messengers [28–31] have been studied.[3]

The solution may be "hidden". In this paper, we propose a possibility that the NMSSM singlet is a meson in the hidden sector. This meson is composed by particles which are charged under a strong gauge symmetry as in Fat Higgs models [33–35] (see also [36, 37]). The dynamics of the strong gauge symmetry provides a meta-stable SUSY breaking vacuum around the origin of the field space [38] and messengers of gauge mediation also couple to the SUSY breaking sector [39]. The SH_uH_d coupling is provided by integrating out some heavy particles charged under the strong gauge symmetry, then, the particle decoupling is also a trigger of the confining dynamics and creates a meta-stable SUSY breaking vacuum.

Since the singlet is also a member of the hidden sector, the SUSY breaking can mediate to the singlet directly. In fact, the singlet receives the SUSY breaking at two loop and it plays an important role to achieve the correct EWSB in our scenario. Although this is not a gauge-mediated SUSY breaking, such a soft mass of the singlet does not cause dangerous flavor-changing neutral currents either.

Supposing the TeV superparticle mass spectrum, a viable SUSY breaking scale will be about $\mathcal{O}(100)$–$\mathcal{O}(1000)$ TeV in our scenario. This would be also an appealing prediction. Such a $\mathcal{O}(100)$ TeV SUSY breaking provides a cosmologically safe gravitino mass of order eV. A large SH_uH_d coupling is possible because the singlet composite scale is not very far from the electroweak scale. Also by the naturalness discussion of the EWSB (for "natural supersymmetry" spectra, see e.g. [40–42]), such a low messenger scale is favored because radiative corrections to the Higgs potential can be relatively small due to the short running.

The rest of the paper is organized as follows. In the next section, we present our model and discuss the formation of the NMSSM singlet and the dynamical SUSY breaking in that model. The mediation of the SUSY breaking to ordinary superpartners is also discussed. In section 3, we show expected mass spectra at the electroweak scale and investigate the phenomenology of the model. In section 4, we conclude the discussion and comment on possible future directions.

2 The framework

In this section, we first present our model and describe the formation of the NMSSM singlet and dynamical SUSY breaking at a meta-stable vacuum. A 2-loop radiative correction from the SUSY breaking sector generates a negative soft mass-squared for the singlet. We also explain gauge-mediated SUSY breaking for ordinary superparticles.

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[2]In refs. [26, 27], it has been pointed out that the domain wall problem in the Z_3 invariant NMSSM is solved with these new vector-like matter fields, which make Z_3 anomalous.

[3]See also ref. [32].

	$\mathrm{SU}(N)_H$	$\mathrm{SU}(3)_C$	$\mathrm{SU}(2)_L$	$\mathrm{U}(1)_Y$
$Q_I\ (I=1,\cdots,N_f)$	**N**	**1**	**1**	0
$\bar{Q}_I\ (I=1,\cdots,N_f)$	**N̄**	**1**	**1**	0
Φ_c	**1**	**3**	**1**	$-1/3$
$\bar{\Phi}_c$	**1**	**3̄**	**1**	$1/3$
Φ_l	**1**	**1**	**2**	$1/2$
$\bar{\Phi}_l$	**1**	**1**	**2**	$-1/2$
f	**N**	**3**	**1**	$-1/3$
$\bar{f}$	**N̄**	**3̄**	**1**	$1/3$
Ψ_u	**N**	**1**	**2**	$1/2$
$\bar{\Psi}_d$	**N̄**	**1**	**2**	$-1/2$
$X_m,\ Y_m$	**1**	**1**	**1**	0

Table 1. The matter content and charge assignment of a model of the composite NMSSM with dynamical SUSY breaking and its gauge mediation.

2.1 Composite NMSSM

Let us consider a supersymmetric $\mathrm{SU}(N)_H$ gauge theory with $(N+6)$ vector-like flavors. The five flavors, Ψ_u, $\bar{\Psi}_d$ and f, $\bar{f}$, are charged under the standard model gauge symmetries while the other $N_f = N + 1$ flavors are SM singlets. The matter content and charge assignment are summarized in table 1. The theory is in the conformal window, $3N/2 \leq N + 6 < 3N$. To maintain the perturbative gauge coupling unification, the rank of the gauge group is constrained as $N \lesssim 4$ with a vector-like pair of $\mathbf{5} + \bar{\mathbf{5}}$ messengers of gauge mediation at an intermediate scale around $100\,\mathrm{TeV}$–$1000\,\mathrm{TeV}$. In the case that the two pairs of the messengers exist, the perturbativity of the couplings holds up to $\sim 10^{15}\,\mathrm{GeV}$ with $N = 4$, evaluated by two-loop renormalization group equations. The constraint may be relaxed if there is a large positive anomalous dimension, giving negative contributions to the beta-functions of the SM gauge couplings [43].

We take the following superpotential in addition to the usual MSSM Yukawa couplings,

$$
\begin{aligned}
W = {} & \lambda_u H_u \bar{\Psi}_d Q_{N_f} + \lambda_d H_d \Psi_u \bar{Q}_{N_f} + \sum_I m_I Q_I \bar{Q}_I \\
& + \sum_{ij} \eta_{ij} X_m Q_i \bar{Q}_j + \sum_A \Big(\eta_c^A Y_m \Phi_c^A \bar{\Phi}_c^A + \eta_l^A Y_m \Phi_l^A \bar{\Phi}_l^A \\
& + M_c \Phi_c^A \bar{\Phi}_c^A + M_l \Phi_l^A \bar{\Phi}_l^A \Big) \\
& + m_\Psi \Psi_u \bar{\Psi}_d + m_f f \bar{f} + M_{XY} X_m Y_m + M_Y Y_m^2/2 \,,
\end{aligned}
\tag{2.1}
$$

where λ_u, λ_d, η_{ij} $(i,j = 1,\cdots,N_f - 1,\ I = 1,\cdots,N_f)$, η_c^A and η_l^A are dimensionless coupling constants and m_Ψ, m_f, M_{XY}, M_Y, m_I, M_c and M_l are mass parameters. The fields Φ_c^A, Φ_l^A, $\bar{\Phi}_c^A$, and $\bar{\Phi}_l^A$ act as messenger superfields after the SUSY is broken, and M_c and M_l should be smaller than the confinement scale Λ, which will be described later:

$$
M_c, M_l \ll \Lambda. \tag{2.2}
$$

The index A denotes the messenger number, and $A = 1, \ldots, N_5$. We assume the following hierarchies among the mass parameters,

$$m_\Psi \sim m_f > \Lambda \gg m_1 > m_2 > \cdots > m_{N_f-1} \gg m_{N_f}. \tag{2.3}$$

The last large hierarchy may be explained by imposing an approximate symmetry with a charge assignment, (charge of m_{Nf}) $\gg$ (charge of m_i). Also, couplings of the messenger fields to Q_{N_f} and $\bar{Q}_{N_f}$ are forbidden by this approximate symmetry.

The theory is first in the conformal window and the hidden gauge coupling is in a fixed point at sufficiently high energies. At the mass scale below $m_\Psi \sim m_f \sim M_{XY}$, Ψ_u, $\bar{\Psi}_d$, f, $\bar{f}$, X_m and Y_m, are integrated out. It is assumed that $M_Y \ll M_{XY}$. Solving the equations of motion of these fields at classical level, we obtain the following superpotential in the effective theory,

$$\begin{aligned}
W' = &-\frac{\lambda_u \lambda_d}{m_\Psi} Q_{N_f} \bar{Q}_{N_f} H_u H_d \\
&- \sum_{ij} \left(\frac{\eta_{ij}\eta_c}{M_{XY}} Q_i \bar{Q}_j \Phi_c \bar{\Phi}_c + \frac{\eta_{ij}\eta_l}{M_{XY}} Q_i \bar{Q}_j \Phi_l \bar{\Phi}_l \right) \\
&+ \sum_{ijkl} \frac{\eta_{ij}\eta_{kl} M_Y}{2 M_{XY}^2} Q_i \bar{Q}_j Q_k \bar{Q}_l \\
&+ \sum_I m_I Q_I \bar{Q}_I + M_c \Phi_c \bar{\Phi}_c + M_l \Phi_l \bar{\Phi}_l,
\end{aligned} \tag{2.4}$$

where the index A is omitted here and hereafter. Decoupling of the five flavors, the effective theory is a SQCD with $(N+1)$ flavors where the gauge coupling gets strong and finally the theory (s-)confines. The confinement scale Λ is near the mass scale $m_\Psi \sim m_f$ when the gauge coupling is sufficiently strong before the confinement. Below the confinement scale, the hidden quarks, Q_I, $\bar{Q}_J$, form meson chiral superfields and the first term of the above superpotential leads to the NMSSM cubic coupling, SH_uH_d ($S \sim Q_{N_f}\bar{Q}_{N_f}$), in a similar manner to the model of [34]. Here, $M_Y \ll M_{XY}$ is favored to suppress masses of meson superfields.

The low-energy effective theory of a SQCD with $(N+1)$ flavors has the following dynamically generated superpotential as well as the superpotential eq. (2.4),

$$W_{\text{dyn}} = \frac{1}{\Lambda^{2N_f-3}} \left(\sum_{IJ} B_I M_{IJ} \bar{B}_J - \det M_{IJ} \right), \tag{2.5}$$

where $M_{IJ} = Q_I \bar{Q}_J$ are meson chiral superfields and $B_I = \epsilon_{II_1 \cdots I_N} Q_{I_1} \cdots Q_{I_N}/N!$, $\bar{B}_I = \epsilon_{II_1 \cdots I_N} \bar{Q}_{I_1} \cdots \bar{Q}_{I_N}/N!$ are (anti-)baryons. We can rewrite these composite chiral superfields in terms of canonically normalized fields,

$$S_{IJ} \sim \frac{M_{IJ}}{\Lambda}, \quad b_I \sim \frac{B_I}{\Lambda^{N_f-2}}, \quad \bar{b}_I \sim \frac{\bar{B}_I}{\Lambda^{N_f-2}}, \tag{2.6}$$

where the Kähler potential cannot be controlled and we have just put the dynamical scale Λ so that S_{IJ}, b_I and $\bar{b}_I$ have mass dimension one correctly. The total superpotential of

the theory is then given by

$$
\begin{aligned}
W_{\text{eff}} = {} & \lambda_S S H_u H_d + M_c \Phi_c \bar{\Phi}_c + M_l \Phi_l \bar{\Phi}_l \\
& + \sum_{ij} (\lambda_{ij}^c S_{ij} \Phi_c \bar{\Phi}_c + \lambda_{ij}^l S_{ij} \Phi_l \bar{\Phi}_l) \\
& + \sum_{IJ} \eta \, b_I S_{IJ} \bar{b}_J - \sum_I \eta' m_I \Lambda S_{II} - \eta'' \frac{\det S_{IJ}}{\Lambda^{N_f - 3}},
\end{aligned}
\tag{2.7}
$$

where $S \equiv S_{N_f N_f}$ and η, η', η'' in front of the terms of the last line are $\mathcal{O}(1)$ numerical factors. Mass terms, $S_{ij} S_{kl}$, will be discussed later. The other coupling constants are estimated as

$$
\lambda_S \sim \frac{\lambda_u \lambda_d \Lambda}{m_\Psi}, \quad \lambda_{ij}^c \sim \frac{\eta_{ij} \eta_c \Lambda}{M_{XY}}, \quad \lambda_{ij}^l \sim \frac{\eta_{ij} \eta_l \Lambda}{M_{XY}}.
\tag{2.8}
$$

For $N_f = N + 1 > 3$, the last term $\sim \det S_{IJ}$ is irrelevant and not important around $S_{IJ} = 0$. We then ignore this term in the discussion of SUSY breaking below.

2.2 Dynamical SUSY breaking

We now show that dynamical SUSY breaking occurs at a meta-stable vacuum around the origin of the meson field space in the present model. There is a local minimum at

$$
b = \bar{b} = \begin{pmatrix} \sqrt{m_1 \Lambda} \\ 0 \\ \vdots \\ 0 \end{pmatrix}, \qquad S_{IJ} = 0,
\tag{2.9}
$$

and the origin of the other fields in the theory. We have not written the $\mathcal{O}(1)$ numerical factors η, η' explicitly. The F-terms of the mesons S_{IJ} at this minimum are given by nonzero values,

$$
F_{S_{IJ}} = -\left(\frac{\partial W}{\partial S_{IJ}} \right)^* = m_I \delta_{IJ} \Lambda \neq 0 \qquad \text{for } I, J \neq 1,
\tag{2.10}
$$

which means SUSY is broken dynamically [38]. There is a supersymmetric vacuum far away from the origin of the meson field space, but the lifetime of the meta-stable vacuum is sufficiently long if $m_1 \ll \Lambda$ is satisfied. With these F-terms, the gravitino mass is given by

$$
m_{3/2}^2 \simeq \sum_{I \neq 1} \frac{|m_I \Lambda|^2}{3 M_P^2},
\tag{2.11}
$$

imposing the condition for the vanishing cosmological constant. Here, $M_P (\simeq 2.43 \times 10^{18} \, \text{GeV})$ is the reduced Planck mass.

Let us analyze the mass spectrum on the meta-stable vacuum. Here, we ignore the small mass parameter m_{N_f} just for simplicity. We expand the baryons around the vacuum,

$$
b = \begin{pmatrix} \sqrt{m_1 \Lambda} + \delta \chi_1 \\ \delta \chi_2 \\ \vdots \\ \delta \chi_{N_f} \end{pmatrix}, \quad \bar{b} = \begin{pmatrix} \sqrt{m_1 \Lambda} + \delta \bar{\chi}_1 \\ \delta \bar{\chi}_2 \\ \vdots \\ \delta \bar{\chi}_{N_f} \end{pmatrix}.
\tag{2.12}
$$

The supermultiplets of S_{11}, S_{I1} and S_{1I} ($I \neq 1$) form mass terms of order $\sqrt{m_1 \Lambda}$ with $(\delta\chi_1 + \delta\bar{\chi}_1)$, $\delta\chi_I$, $\delta\bar{\chi}_I$ respectively. The supermultiplet of $(\delta\chi_1 - \delta\bar{\chi}_1)$ is the massless Nambu-Goldstone (NG) multiplet of the spontaneously broken baryon number symmetry. When we gauge this symmetry, we just obtain the massive Abelian gauge multiplet by the super Higgs mechanism. Alternatively, we can explicitly break the symmetry and make the NG mode massive. The meson scalars S_{IJ} ($I, J \neq 1$) are flat directions at classical level but all of them except for the NMSSM singlet scalar S get nonzero masses by the 1-loop Coleman-Weinberg potential.

The scalar of S remains massless even at 1-loop level. However, there is an important 2-loop correction for this scalar field, which gives a runaway potential for the NMSSM singlet scalar. Near the origin, the generated negative soft mass-squared is given by [44]

$$V_{2-\text{loop}} \approx -1.26\, \eta^6 \frac{m_2 \Lambda}{(16\pi^2)^2}\, S^\dagger S, \tag{2.13}$$

which makes it possible to explain the correct EWSB with the non-zero $\langle S \rangle$, and then the Higgsino mass term is generated by the superpotential interaction $\lambda_S S H_u H_d$ as in the usual NMSSM. Here, $m_2 \gg m_3$ is assumed. The potential (2.13) indicates that the size of the soft mass depends on the SUSY breaking scale $\sqrt{m_2 \Lambda}$ while that tachyonic soft mass has to be smaller than the electroweak scale to preserve naturalness. Then, if the coupling η is $\mathcal{O}(1)$, $\sqrt{m_2 \Lambda}$ is not bigger than $\sim 100\,\text{TeV}$ in the present model. In this case, the cosmologically safe gravitino with a mass around $10\,\text{eV}$ is predicted.

We need to stabilize the runaway direction of the NMSSM scalar from the potential (2.13). In addition, some of the fermion components in S_{IJ} are massless and need SUSY mass terms. For these reasons, we consider higher dimensional operators,

$$\Delta W = \frac{\lambda_{IJKL}}{M_0} Q_I \bar{Q}_J Q_K \bar{Q}_L = \tilde{m}_{IJKL} S_{IJ} S_{KL}, \tag{2.14}$$

where M_0 is some ultra-violet (UV) mass scale and

$$\tilde{m}_{IJKL} \sim \frac{\lambda_{IJKL} \Lambda^2}{M_0}. \tag{2.15}$$

Here, among the higher dimensional operators, $W \ni Q_i \bar{Q}_j Q_k \bar{Q}_l$ arises from eq. (2.1). The higher dimensional operators in eq. (2.14) can give nonzero masses for all the meson scalars and fermions except for the massless goldstino of SUSY breaking. The mass terms

$$\tilde{m}_{N_f N_f pp}, \ \tilde{m}_{ppN_f N_f} \ (p =, 2\ldots, N_f - 1) \tag{2.16}$$

should be suppressed, since otherwise the EWSB scale is destabilized by $\Delta V \sim F_{pp}(\tilde{m}_{ppN_f N_f} + \tilde{m}_{N_f N_f pp})S + \text{h.c.}$ We can suppress these dangerous mass terms by imposing a symmetry which forbids $(Q_{N_f} \bar{Q}_{N_f})(Q_p \bar{Q}_q)$, where $q = 2, \ldots, N_f - 1$.

In addition, the linear terms $\Delta V \sim F_{pp}(\tilde{m}_{ppqr} + \tilde{m}_{qrpp})S_{qr} + \text{h.c.}$ shift the vacuum expectation values of the pseudo-moduli mesons from the origin, $\Delta S_{qr} \sim 16\pi^2 \Lambda^2/M_0 \sim 16\pi^2 \tilde{m}$ as discussed in [39, 45]. This must be smaller than $4\pi\sqrt{m_2 \Lambda}$ so that the analysis of SUSY breaking above is valid and also smaller than the messenger masses $M_c \sim M_l$ for the messengers not to be tachyonic (see below). Consequently, $\tilde{m}_{ppqr}$ and $\tilde{m}_{qrpp}$ can not be significantly larger than 0.1–1 TeV.

	Q_{N_f}	$\bar{Q}_{N_f}$	Q_p	$\bar{Q}_q$	X_1	X_2	X_3	X_m	Y_m	$(H_u\bar{\Psi}_d)$	$(H_d\Psi_u)$	$(\Phi_{l,c}\bar{\Phi}_{l,c})$
U(1)	$-4/5$	$-1/5$	$-1/5$	0	1	$4/5$	$2/5$	$1/5$	$-1/15$	$4/5$	$1/5$	$1/15$
	M_1	M_2	M_3	M_{XY}	m_{N_f}	m_p	M_Y	$M_{l,c}$				
U(1)	-2	$-8/5$	$-4/5$	$-2/15$	1	$1/5$	$2/15$	$-1/15$				

Table 2. The example of the U(1) charge assignment. The mass parameters m_{N_f}, m_p, M_{XY}, M_Y and $M_a (a = 1, 2, 3)$ are considered to be spurions of U(1) symmetry breaking. (Also, m_Ψ (m_f) has a non-vanishing charge.).

The higher dimensional operators (2.14) to satisfy the constraints can be UV completed in the following renormalizable superpotential for instance,

$$W = X_1 Q_{N_f}\bar{Q}_{N_f} + \sum_p (X_2 Q_{N_f}\bar{Q}_p + X_3 Q_p\bar{Q}_{N_f})$$
$$+ \sum_{p,q} X_m(Q_p\bar{Q}_q) + M_{XY}X_m Y_m + \frac{1}{2}M_Y Y_m^2$$
$$+ \frac{1}{2}(M_1 X_1^2 + M_2 X_2^2 + M_3 X_3^2), \tag{2.17}$$

where $p, q = 2, \ldots, N_f-1$, and X_a ($a = 1, 2, 3$) are chiral superfields which are singlet under the standard model and $SU(N)_H$ gauge symmetries. The mass parameters are larger than the confinement scale:

$$M_1, M_2, M_3, M_{XY}, M_Y > \Lambda. \tag{2.18}$$

The above superpotential is explained by an approximate U(1) symmetry, with a charge assignment summarized in table 2. A Yukawa coupling is implicitly multiplied by each of cubic terms. Linear terms of X_1, X_m and Y_m have been removed by shifts of the fields. The charge of $(H_u H_d)$ is chosen to be non-zero such that the bare mass term, so-called μ-term, is prohibited. The MSSM matter fields, which are not shown here, are also charged under this approximate symmetry.

When new fields X_i are integrated out, we obtain the higher dimensional operators (2.14) with forbidden $\lambda_{pqN_fN_f}$ and $\lambda_{N_fN_fpq}$, correctly. (Therefore, $\tilde{m}_{N_fN_fpp}$ and $\tilde{m}_{ppN_fN_f}$ vanish.) Other mass terms which mix the NMSSM singlet S to other mesons e.g. $m_{N_fN_fpN_f}$ are prohibited. Since the symmetry forbids the messengers-Q_{N_f} ($\bar{Q}_{N_f}$) couplings such as $\Phi_l\bar{\Psi}_d Q_{N_f}$ in the superpotential eq. (2.1), couplings of S to the messenger pairs, $S\Phi_c\bar{\Phi}_c$ and $S\Phi_l\bar{\Phi}_l$, vanish: the EWSB scale is not destabilized.

2.3 Gauge mediation

We now consider gauge mediation of SUSY breaking to ordinary superparticles. Integrating out the messenger fields, the gaugino mass and scalar masses are generated as

$$m_{\rm soft} \sim \frac{g^2}{16\pi^2}\frac{\bar{m}\Lambda}{M}, \tag{2.19}$$

where g denotes the standard model gauge coupling and we have defined $\bar{m} \equiv \sum_{i\neq 1}\lambda_{ii}m_i$. Here, $M \sim M_c \sim M_l$ is the messenger mass scale. To obtain the soft masses at the

electroweak scale, we set

$$\Lambda_{\text{mess}} \equiv \frac{\bar{m}\Lambda}{M} \sim 100\,\text{TeV}. \tag{2.20}$$

The condition that the messenger fields are not tachyonic gives $\bar{m}\Lambda < M^2$. The messengers also give the 1-loop Coleman-Weinberg potential to the meson scalars S_{ij}, which also leads to the shifts of the pseudo-moduli [39]. Since there is a vacuum with lower energy where the messenger scalars are condensed, a tachyonic direction appears around $S_{ij} \sim M$ (the existence of this vacuum is essential for non-vanishing leading order gaugino masses. See e.g. [45–47]). Then, the stabilized point of the pseudo-moduli has to be smaller than the messenger mass scale M, which gives a constraint, $M > \lambda^2\sqrt{m\Lambda}$. Note that the transition between the SUSY breaking local minimum to the minimum with $S_{ij} \sim M$ provides a similar but more stringent constraint, $M \gtrsim 3\lambda\sqrt{m\Lambda}$ [48],[4] which will be discussed in the next section.

3 Phenomenology

We now turn to discuss the phenomenology of the model. As mentioned in the previous section, gaugino, squark and slepton masses are the same as those in usual gauge mediation with a messenger mass M and B-term of $\bar{m}\Lambda = M\Lambda_{\text{mess}}$. On the other hand, values of soft breaking parameters in the extended Higgs sector are different from those of usual gauge mediation scenarios.

The low-energy effective superpotential of the extended Higgs sector can be written as

$$W \supset \xi_F S + \frac{1}{2}\mu' S^2 + \lambda_S S H_u H_d, \tag{3.1}$$

where $\xi_F = -m_{N_f}\Lambda$ in eq. (2.7) and $\mu' = 2\tilde{m}_{N_f N_f N_f N_f}$ in eq. (2.14), thus, we treat these and λ_S as free parameters here. Note that the large λ_S of ~ 1 is quite natural as shown in appendix B.

The coupling of S^3 term will be suppressed because it is provided from a higher dimensional operator, $(Q_{N_f}\bar{Q}_f)^3$, and hence, it has been neglected. The corresponding SUSY breaking terms are defined as

$$V \supset \left(\xi_S S + \frac{1}{2}m_S'^2 S^2 + A_\lambda \lambda_S S H_u H_d + \text{h.c.}\right) \\ + m_S^2|S|^2 + m_{H_u}^2|H_u|^2 + m_{H_d}^2|H_d|^2, \tag{3.2}$$

where the m_S^2 receives the negative two-loop contribution $\sim \eta^6 m_2\Lambda/(16\pi^2)^2 \sim \eta^6(m_{3/2}M_P)/(16\pi^2)^2$ written in eq. (2.13). It turns out that this negative m_S^2 is important for the successful EWSB. The other soft mass parameters are

$$A_\lambda \approx 0\,\text{GeV}, \quad \xi_S \approx 0\,\text{GeV}, \quad m_S' \approx 0\,\text{GeV}, \tag{3.3}$$

at the messenger scale because the singlet S does not couple messenger fields directly. Hereafter we assume all parameters are real for simplicity.

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[4]If the messenger superfields are not thermalized, the constraint becomes weaker as $M \gtrsim 1.5\lambda\sqrt{m\Lambda}$, provided that the SUSY breaking local minimum is selected in the early universe.

3.1 EWSB and Higgs mass

To show the viable parameter space and typical mass spectra, we discuss the feature of the extended Higgs sector. We define neutral scalar components of Higgs as $H_u^0 = v_u + (h_u + ia_u)/\sqrt{2}$, $H_d^0 = v_d + (h_d + ia_d)/\sqrt{2}$ and $S = v_S + (s_R + is_I)/\sqrt{2}$, respectively. These v_u, v_d, v_S denote vacuum expectation values and $\tan\beta \equiv v_u/v_d$ and $v \equiv \sqrt{v_u^2 + v_d^2} \simeq 174.1\,\text{GeV}$. The vacuum conditions can be obtained by

$$
\begin{aligned}
\frac{m_Z^2}{2} &\approx -\mu_{\text{eff}}^2 + \frac{m_{H_d}^2 - m_{H_u}^2 \tan^2\beta}{\tan^2\beta - 1}, \\
\sin 2\beta &\approx \frac{2(\mu_{\text{eff}}\mu' + \lambda_S \xi_F)}{m_{H_u}^2 + m_{H_d}^2 + 2\mu_{\text{eff}}^2 + \lambda_S^2 v^2}, \\
\mu_{\text{eff}} &\approx -\frac{\lambda_S \mu'}{2}\frac{2\xi_F - \lambda_S v^2 \sin 2\beta}{m_S^2 + \mu'^2 + \lambda_S^2 v^2},
\end{aligned}
\tag{3.4}
$$

at tree level. Here $m_Z^2 = (g_Y^2 + g_2^2)v^2/2$ and $\mu_{\text{eff}} \equiv \lambda_S v_S$. Note that without the negative m_S^2, μ_{eff} is predicted to be around $-\lambda_S \xi_F/\mu'$. As a result, the predicted value of $\tan\beta$ is huge: the successful EWSB does not occur, unless the Lagrangian with eq. (3.1) and (3.2) becomes a MSSM limit by $\lambda_S \to 0$ and $(\lambda_S \xi_F/\mu') = \text{fixed}$.[5]

Using the following base

$$
\begin{pmatrix} h_1' \\ h_2' \\ h_3' \end{pmatrix} = \begin{pmatrix} \cos\beta & -\sin\beta & 0 \\ \sin\beta & \cos\beta & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} h_d \\ h_u \\ s_R \end{pmatrix},
\tag{3.5}
$$

the CP-even Higgs mass squared matrix can be written as

$$
M_{\text{H}}^2 = \begin{pmatrix} M_{\text{H}11}^2 & M_{\text{H}12}^2 & M_{\text{H}13}^2 \\ & M_{\text{H}22}^2 & M_{\text{H}23}^2 \\ & & M_{\text{H}33}^2 \end{pmatrix},
\tag{3.6}
$$

where

$$
\begin{aligned}
M_{\text{H}11}^2 &\approx m_Z^2 (\cos 2\beta)^2 + \lambda_S^2 v^2 (\sin 2\beta)^2, \\
M_{\text{H}22}^2 &\approx 2(\mu_{\text{eff}}\mu' + \lambda_S \xi_F)/\sin 2\beta \\
&\quad + \left(m_Z^2 - \lambda_S^2 v^2\right)(\sin 2\beta)^2, \\
M_{\text{H}33}^2 &\approx \mu'(-\lambda_S \xi_F + \lambda_S^2 v^2 \sin\beta/2)/\mu_{\text{eff}}, \\
M_{\text{H}12}^2 &\approx (-m_Z^2 + \lambda_S^2 v^2)\sin 4\beta/2, \\
M_{\text{H}13}^2 &\approx \lambda_S v(2\mu_{\text{eff}} - \mu'\sin 2\beta), \\
M_{\text{H}23}^2 &\approx -\lambda_S v\mu' \cos 2\beta,
\end{aligned}
\tag{3.7}
$$

at tree level. As we can see, there are positive and negative contributions to the lightest Higgs boson mass: a positive contribution is the additional F-term in $M_{\text{H}11}^2$, $\lambda_S^2 v^2 (\sin 2\beta)^2$

---

[5]There are several gauge mediation models explaining the successful EWSB, where the effective μ-term and B_μ terms are described by the VEV and the NMSSM-like singlet [49–53].
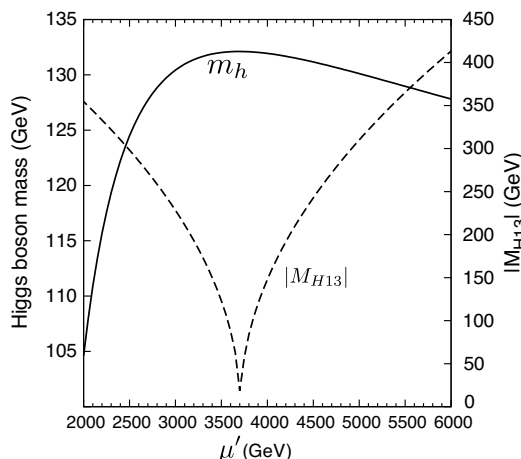
Figure 1. The lightest Higgs boson mass m_h and an off-diagonal component of the neutral CP-even Higgs mass matrix M_{H13}. We take $(\lambda_S, \tan\beta, \mu_{\text{eff}}/|\mu_{\text{eff}}|) = (1.0, 4.0, 1)$ and $(M, \Lambda_{\text{mess}}, N_5) = (300\,\text{TeV}, 180\,\text{TeV}, 1)$. Here, $\alpha_s(M_Z) = 0.1185$ and $m_t(\text{pole}) = 173.34\,\text{GeV}$.

and a negative contribution comes from $M_{\text{H}13}^2$ after diagonalizing the mass squared matrix. Therefore, small $\tan\beta$ and $\mu_{\text{eff}} \sim \mu' \sin 2\beta/2$ are favored to push the lightest Higgs boson mass up.

Actually, we can see that the lightest Higgs boson mass is maximized around the minimum of $|M_{\text{H}13}|$ in figure 1. In our numerical calculations, the Higgs boson mass and SUSY mass spectra are calculated using NMGMSB [54], which is contained in NMSSMTools 4.8.2. The NMSSM parameters are taken as $(\lambda_S, \tan\beta, \mu_{\text{eff}}/|\mu_{\text{eff}}|) = (1.0, 4.0, 1)$, while parameters in the messenger sector are $(\Lambda_{\text{mess}}, M, N_5) = (180\,\text{TeV}, 300\,\text{TeV}, 1)$, where N_5 is a number of the messenger pairs. It should be noted that considering above discussion on the Higgs boson mass and vacuum conditions in eq. (3.4), the values of μ_{eff}, μ', $\sqrt{\xi_F}$ and m_S would be the same order in the viable parameter space.

We also estimate the fine-tuning of the EWSB scale using the following fine-tuning measure [55, 56]:[6]

$$\Delta = \max\left\{\left|\frac{\partial \ln v}{\partial \ln |a|}\right|\right\}, \left(a \in \begin{array}{c} \text{fundamental mass} \\ \text{parameters} \end{array}\right) \tag{3.8}$$

where $a = \xi_F, \mu', \Lambda_{\text{mess}}, |m_S^2|$ in our case. ($d\ln|\xi_F|$, $d\ln|\Lambda_{\text{mess}}|$ and $d\ln|m_S^2|$ correspond to $d\ln|m_{N_f}|$, $d\ln|\bar{m}|$ and $d\ln|m_2|$, respectively.)

In figure 2, the fine-tuning measure Δ is shown on the gluino mass (pole mass)-M plane. Within the green (shaded) region, the lightest Higgs boson mass is in a range of 122–128 GeV. It is noticed that, thanks to the additional F-term contribution, the observed Higgs boson mass around 125 GeV is easily explained even with the 1.2 TeV gluino mass in the gauge mediation scenario, as pointed out in ref. [31]. The larger gluino mass can of course also be consistent with the observed Higgs boson mass by changing the λ_S value

---

[6]The definition of Δ here differs from the original one by a factor 2. With the definition of eq. (3.8), $\partial \ln v/\partial \ln |\mu| \simeq 2\mu^2/m_Z^2$ in the MSSM.
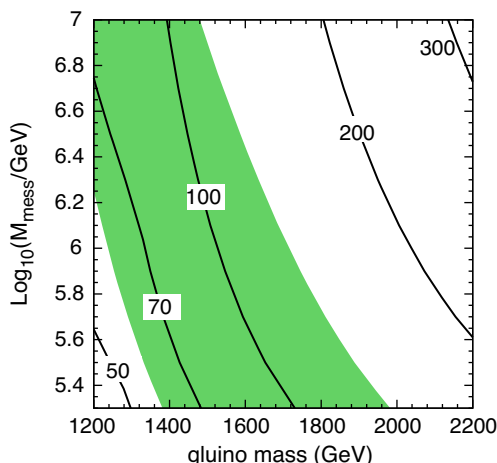
Figure 2. The fine-tuning measure Δ (black solid). We take $\lambda_S = 1.0$, $\tan\beta = 4$, $\mu_{\text{eff}} > 0$, $\mu' = 6\,\text{TeV}$ and $N_5 = 2$, while the other parameters are the same as in figure 1. In the green (shaded) region, the lightest Higgs boson mass is in a range of 122–128 GeV.

to slightly smaller. As a result, the fine-tuning of the EWSB scale is drastically improved, compared to the MSSM in gauge mediation; $\Delta \lesssim 100$ for $M \simeq 200\,\text{TeV}$.

3.2 Cosmology

Before closing this section, let us discuss cosmological aspects and implications to the collider signals. In our model, the gravitino mass is estimated as

$$m_{3/2} \simeq \frac{\bar{m}\Lambda}{\lambda} \frac{1}{\sqrt{3}M_P}$$
$$\approx 7\,\text{eV} \cdot \left(\frac{\Lambda_{\text{mess}}}{100\,\text{TeV}}\right)\left(\frac{M}{300\,\text{TeV}}\right)\frac{1}{\lambda}, \tag{3.9}$$

where λ is a typical value of λ_{ii}, provided $\lambda_{ii} \sim \lambda_{jj}(i \neq j)$.

To satisfy the warm dark matter constraint, the gravitino should be lighter than 16 eV [57], and in this range there is no constraint on the reheating temperature. Then, the messenger scale is bounded from above as

$$M < 6.7 \times 10^5\,\text{GeV} \cdot \lambda \left(\frac{100\text{TeV}}{\Lambda_{\text{mess}}}\right), \tag{3.10}$$

with

$$\lambda \gtrsim 0.15 \cdot \left(\frac{\Lambda_{\text{mess}}}{100\,\text{TeV}}\right)^2, \tag{3.11}$$

which is required not to conflict with $M^2 > \bar{m}\Lambda$. On the other hand, to avoid the unstable SUSY breaking minimum with a life-time shorter than the age of the universe, the messenger scale should be [48]

$$M \gtrsim 9 \times 10^5\text{GeV} \cdot \lambda^2 \left(\frac{\Lambda_{\text{mess}}}{100\,\text{TeV}}\right). \tag{3.12}$$

Combining the above three conditions, we get the upper-bound on Λ_{mess} as

$$\Lambda_{\text{mess}} \lesssim 150\,\text{TeV}. \tag{3.13}$$

For $N_5 = 1$, it is difficult to satisfy this upper-bound taking into account null results of latest LHC SUSY searches.[7] By demanding that the gluino mass, $m_{\tilde{g}}$, be larger than 1.4 (1.6) TeV, the lower-bound on Λ_{mess} is

$$\Lambda_{\text{mess}} \gtrsim 180\,(205)\,\text{TeV}/N_5, \tag{3.14}$$

for $\Lambda_{\text{mess}}/M = 1/2$. Therefore, $N_5 > 1$ is required. With the very light gravitino, the next-to-lightest SUSY particle (NLSP), which is likely to be the stau in our model, decays promptly (for a very light gravitino spectrum search, see [58]). In our case, right-handed sleptons are almost degenerated in mass, and the strong constraint comes from the SUSY searches in final states with multi-jets, multi-leptons ($\gtrsim 3$) and missing transverse momentum [59]. The chargino needs to be heavier than 850–900 GeV, resulting in the lower-bound on the gluino mass as $m_g > 2.4$–2.7 TeV with a GUT relation among gaugino masses. Here, $m_{\tilde{g}}$ is the gluino mass.

In the meantime, if the gravitino mass is in a range

$$m_{3/2} \gtrsim \mathcal{O}(10)\,\text{keV}, \tag{3.15}$$

the gravitino can be a cold dark matter [60, 61] with an appropriate reheating temperature or the late-time entropy production [62, 63]. The lower bound on the gravitino mass in turn leads to the lower bound on the messenger scale:

$$M > 4.2 \times 10^8\text{GeV} \cdot \lambda \left(\frac{m_{3/2}}{10\,\text{keV}}\right)\left(\frac{100\,\text{TeV}}{\Lambda_{\text{mess}}}\right), \tag{3.16}$$

or equivalently

$$(\bar{m}\Lambda/\lambda)^{1/2} \gtrsim 2.1 \times 10^8\,\text{GeV}\left(\frac{m_{3/2}}{10\,\text{keV}}\right)^{1/2}. \tag{3.17}$$

In this case, the stability bound of the SUSY breaking minimum, eq. (3.12), is no longer important. Note that the messenger scale can be still low as $\sim 10^6$ GeV for $\lambda \sim 0.01$, and $\eta \sim 0.2$ (see eq. (2.7)).

Since the typical decay length of the NLSP is $\mathcal{O}(10)$ m, the stau NLSP is strongly constrained as $m_{\tilde{\tau}_1} \gtrsim 450$–500 GeV [64, 65], where $m_{\tilde{\tau}_1}$ is the (lighter) stau mass. Therefore, the bino NLSP predicted with $N_5 = 1$ may be favored if the gravitino is the cold dark matter. In this case, the LHC signatures as well as limits of SUSY particle masses are similar to those in gravity mediation, and the strongest constraint comes from the SUSY searches with multi-jets and missing transverse momentum. So far, it is expected that $m_{\tilde{g}} \gtrsim 1.4$–1.6 TeV is required [66], depending on the squark mass.

Finally, we show the typical mass spectra and fine-tuning measure Δ in table 3. At the point **I** with $N_5 = 1$, the bino-like neutralino is the NLSP and $m_{3/2} \gtrsim \mathcal{O}(10)$ keV which

---

[7]It may be possible to avoid this constraint if the reheating temperature is sufficiently low or the gravitino abundance is diluted by the late-time entropy production, allowing the larger gravitino mass than 16 eV.

Parameters	Point **I**	Point **II**
Λ_{mess} (TeV)	185	130
M (TeV)	1000	200
N_5	1	3
$\tan\beta$	4.5	5.0
λ_S	0.9	0.92
μ_{eff} (GeV)	783	893
μ' (GeV)	6000	4000
ξ_F (GeV)^2	-4.5×10^6	-3.3×10^6
m_S^2 (GeV)^2	-1.95×10^6	-1.41×10^6
Δ	165	235
Particles	Mass (GeV)	Mass (GeV)
$\tilde{g}$	1430	2890
$\tilde{q}$	1840–1940	2570–2680
$\tilde{t}_{2,1}$	1860, 1660	2610, 2410
$\tilde{\chi}^\pm_{2,1}$	807, 486	1130, 886
$\tilde{\chi}^0_5$	5160	3610
$\tilde{\chi}^0_4$	807	1130
$\tilde{\chi}^0_3$	796	910
$\tilde{\chi}^0_2$	486	886
$\tilde{\chi}^0_1$	253	580
$\tilde{e}_{L,R}(\tilde{\mu}_{L,R})$	652, 326.6	810, 392.9
$\tilde{\tau}_{2,1}$	652, 325.9	810, 392.3
$H^\pm$	851	1137
$A_{2,1}$	4940, 847	3380, 1133
h_1	124.7	125.1

Table 3. Mass spectrum in sample points. Values of λ_S and $\tan\beta$ show the values at the SUSY scale and values of ξ_F, μ' and M_S^2 show the values at the messenger scale in this table. Here, A_1 (A_2) denotes the lighter (heavier) CP-odd Higgs.

enables us to explain the observed dark matter by the gravitino. As discussed above, the gravitino mass with $m_{3/2} < 16\,\mathrm{eV}$ is difficult to be achieved in this case. On the other hand, at the points **II**, the stau is the NLSP; therefore, $m_{3/2} < 16\,\mathrm{eV}$ is required to avoid the strong constraint on the stable stau without increasing the SUSY mass scale. In such a very light gravitino region, there is no constraint on the reheating temperature and the observed dark matter relic would be explained by another particle (e.g. QCD axion).

4   Conclusion and discussion

We have proposed a new scheme for the NMSSM in gauge mediation, where in general the successful EWSB does not occur due to the absence of the soft SUSY breaking mass

parameter for S. In our framework, S is a composite meson in the hidden QCD, which is responsible for the dynamical SUSY breaking. The required soft SUSY breaking mass for S naturally arises after SUSY is broken dynamically, and the electroweak symmetry is successfully broken.

Although the UV Lagrangian eq. (2.1) contains two singlets, it is possible to construct a model without singlets as shown in appendix A, where the required particle content becomes less. In this case, there are additional contributions to the soft masses for the singlet and Higgs doublets, and all trilinear couplings. The stop, sbottom and stau masses are also modified. Therefore, the quite different phenomenology is expected to appear.

As a concrete model of the dynamical SUSY breaking, we have utilized the ISS model. However, it is possible to consider other dynamical SUSY breaking models. For instance, if the IYIT model [67, 68] is adopted, a model similar to the Dirac NMSSM [69] appears as a low-energy effective theory [70], which also increases the Higgs boson mass in a similar but different way. It needs to be checked whether the correct EWSB is explained.

In our model, the gravitino can be either very light as $\sim 10\,\mathrm{eV}$ or $\mathcal{O}(10)\,\mathrm{keV}$. In the former case, there is no upper bound for the reheating temperature, but the gravitino can not explain the observed dark matter abundance. One needs another candidate for a dark matter. For instance, the stable baryon in the hidden QCD may explain the observed abundance of the dark matter. In the latter case, the gravitino is cold enough and a dark matter candidate. Although the gravitino tends to be over-produced in general, it is possible to fit for the standard cosmology if the late-time entropy production exists [62, 63].

A   UV model without singlets

An UV model without singlets is shown. The particle contents are less than eq. (2.1). The superpotential is

$$
\begin{aligned}
W = {}& \lambda_u H_u \bar{\Psi}_d Q_{N_f} + \lambda_d H_d \Psi_u \bar{Q}_{N_f} \\
& + \sum_A \Big(\eta_i^c \Phi_c^A \bar{f} Q_i + \bar{\eta}_i^c \bar{\Phi}_c^A f \bar{Q}_i + \eta_i^l \Phi_l^A \bar{\Psi}_d Q_i + \bar{\eta}_i^l \bar{\Phi}_l^A \Psi_u \bar{Q}_i \\
& \qquad + M_c \Phi_c^A \bar{\Phi}_c^A + M_l \Phi_l^A \bar{\Phi}_l^A \Big) \\
& + m_\Psi \Psi_u \bar{\Psi}_d + m_f f \bar{f} + m_I Q_I \bar{Q}_I .
\end{aligned}
\tag{A.1}
$$

After integrating out f, $\bar{f}$, Ψ_u and $\bar{\Psi}_d$, the number of the flavor becomes $N_f = N + 1$ and the theory confines. Then, the low-energy effective Lagrangian is

$$
\begin{aligned}
W_{\text{eff}} = {} & \lambda_S S H_u H_d + M_c \Phi_c \bar{\Phi}_c + M_l \Phi_l \bar{\Phi}_l \\
& + \lambda^c_{ij} S_{ij} \Phi_c \bar{\Phi}_c + \lambda^l_{ij} S_{ij} \Phi_l \bar{\Phi}_l \\
& + \tilde{\lambda}^u_i S_{N_f i} H_u \bar{\Phi}_l + \tilde{\lambda}^d_i S_{i N_f} H_d \Phi_l \\
& + \eta\, b_I S_{IJ} \bar{b}_J - \eta' m_I \Lambda S_{II} - \eta'' \frac{\det S_{IJ}}{\Lambda^{N_f - 3}}.
\end{aligned}
\tag{A.2}
$$

In this model, there are messenger-Higgs couplings, $S_{N_f i} H_u \bar{\Phi}_l$ and $S_{i N_f} H_d \Phi_l$. These couplings generate two-loop negative contributions to m_S^2 other than the contribution shown in eq. (2.13). The soft masses of H_u and H_d as well as those of the stop, sbottom and stau are also modified by two-loop effects. The trilinear couplings A_λ, A_t, A_b and A_τ are generated at the one-loop level, where A_t, A_b and A_τ are those of the stop, sbottom, and stau, respectively. Size of one-loop corrections to $m_{H_u}^2$ and $m_{H_d}^2$ are parametrically similar to those of two-loop contributions. These one-loop corrections to $m_{H_{u,d}}^2$ vanish for $\Lambda_{\text{mess}}/M \to 0$.

B Couplings at the fixed point

In this appendix, we estimate the couplings relevant to λ_S. Provided that $\lambda_u H_u \bar{\Psi}_d Q_{N_f}$, $\lambda_d H_d \Psi_u \bar{Q}_{N_f}$ and $\eta_{ij} X_m Q_i \bar{Q}_j$ in eq. (2.1) are at the fixed point, using the a-maximization technique [71], the anomalous dimensions of the fields are

$$
\begin{aligned}
\gamma_{H_u} = \gamma_{H_d} &= 0.148, \qquad \gamma_{\Psi_u} = \gamma_{\bar{\Psi}_d} = -0.088, \\
\gamma_f = \gamma_{\bar{f}} &= -0.114, \quad \gamma_{Q_{N_f}} = \gamma_{\bar{Q}_{N_f}} = -0.060, \\
\gamma_{Q_i} = \gamma_{\bar{Q}_i} &= -0.106, \quad \gamma_{X_m} = 0.211,
\end{aligned}
\tag{B.1}
$$

for $N = 4$. On the other hand, the one-loop calculations give

$$
\begin{aligned}
\gamma_{H_u} &= \frac{1}{16\pi^2}(N\lambda_u^2), \quad \gamma_{H_d} = \frac{1}{16\pi^2}(N\lambda_d^2), \\
\gamma_{X_m} &= \frac{1}{16\pi^2}(N_f - 1)N\eta_Q^2,
\end{aligned}
\tag{B.2}
$$

where $\eta_{ij} = \eta_Q$ is taken for simplicity. Then, the fixed point values of λ_u, λ_d and η_Q are estimated as $\lambda_u = \lambda_d \approx 2.42$, $\eta_Q \approx 1.44$. After f and $\bar{f}$ are integrated out, the theory becomes stronger but still in the conformal window. The couplings become larger as $\lambda_u = \lambda_d \approx 5.11$, $\eta_Q \approx 2.71$ and we obtain a sizable λ_S coupling.

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