Higgs boson from the meta-stable SUSY breaking sector

Yang Bai, Ji Ji Fan, and Zhenyu Han

1Department of Physics, Sloane Laboratory, Yale University, New Haven, CT 06520
2Department of Physics, University of California, Davis, CA 95616

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We construct a calculable model of electroweak symmetry breaking in which the Higgs doublet emerges from the meta-stable SUSY breaking sector as a pseudo Nambu-Goldstone boson. The Higgs boson mass is further protected by the little Higgs mechanism, and naturally suppressed by a two-loop factor from the SUSY breaking scale of 10 TeV. Gaugino and sfermion masses arise from standard gauge mediation, but the Higgsino obtains a tree-level mass at the SUSY breaking scale. At 1 TeV, aside from new gauge bosons and fermions similar to other little Higgs models and their superpartners, our model predicts additional electroweak triplets and doublets from the SUSY breaking sector.

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I. INTRODUCTION

Dynamical supersymmetry (SUSY) breaking has always been difficult due to the existence of nonzero Witten index of supersymmetric pure gauge theories [1]. It was believed before that gauge theories breaking SUSY dynamically must either be chiral or have massless matter [2]. Models built along this line turn out to be rather baroque. The situation is greatly improved by the recent discovery of long-lived metastable SUSY breaking vacua in $N = 1$ supersymmetric QCD theory by Intriligator, Seiberg and Shih (ISS) [3]. The ISS model opens up new possibilities to build realistic models. Using ISS’s idea, a number of new and simple models have been constructed to communicate SUSY breaking (directly) to the standard model [4, 5].

In an interesting attempt, the ISS model is embedded into the supersymmetric standard model by extending the Higgs sector [6]. But in that work, a strong coupling in the Higgs sector is needed to generate sufficiently large gluino masses. In this paper we will construct an electroweak symmetry breaking model with a weakly coupled Higgs sector embedded inside the meta-stable SUSY breaking sector.

In our model, the soft masses of gauginos and sfermions are generated by the standard gauge mediation. To obtain the soft masses above a few hundred GeV, the lowest allowed SUSY breaking scale is of order 10 TeV. We take two identical ISS sectors with $N_f = 4$ massive fundamental flavors and set the SUSY breaking and the global $U(4)^2$ breaking scale $\mu$ to be of order 10 TeV. We then gauge the diagonal subgroup $SU(3)_W \times U(1)_X$ of the unbroken global symmetry $U(3)^2$. By adding Yukawa couplings with extra vector-like fermions, both of the $U(3)$ symmetries are spontaneously broken at the scale $\mu \sim O(\mu/4\pi) \sim O(1 \text{ TeV})$. At the same time the gauge symmetry is broken to the electroweak symmetry. Among the resulting Goldstone boson fields, one doublet is eaten by the heavy gauge bosons, while the other doublet is a pseudo Nambu-Goldstone boson (PNGB) and identified as the Higgs doublet (this Higgs doublet is a light linear combination of the two doublets $H_u$ and $H_d$ in the minimal supersymmetric standard model). The one-loop effective potential of the Higgs doublet also breaks electroweak symmetry at $v \sim O(f/4\pi) \sim O(\mu/(4\pi)^2) \sim O(100 \text{ GeV})$ and gives a light Higgs boson with a mass from 100 GeV to 200 GeV. In short, all the mass scales arise from one single scale $\mu$ in our model. Since the scale $\mu$ can be generated dynamically [7], the hierarchy between the electroweak scale and the Planck scale may also be explained.

Below the SUSY breaking scale, the “collective symmetry breaking mechanism” protects the Higgs boson mass as in the little Higgs models [8]. Then in our model, like in the super-little Higgs models [9, 10], the Higgs boson mass is doubly protected by both SUSY and the little Higgs mechanism, and the fine-tuning problem has been alleviated compared to the minimal supersymmetric standard model. However, unlike the super-little Higgs models where the soft masses are introduced by hand, all the soft masses in our model are calculable. Our model contains new gauge bosons and fermions at 1 TeV that are also present in other little Higgs models. Furthermore, our model predicts additional electroweak triplets and doublets at 1 TeV from the pseudomoduli in the meta-stable SUSY breaking sector.

The paper is organized as follows. In Sec. II we will briefly review the ISS model, its symmetries and its field content. In Sec. III we present our model and describe how to break $SU(3)_W \times U(1)_X$ to $SU(2)_W \times U(1)_Y$. In Sec. IV we explain how to break the electroweak symmetry. We address the vacuum alignment and order one quartic Higgs coupling issues. We also explicitly minimize the full effective scalar potential to support our step-by-step analysis. In Sec. V we discuss the soft masses of gauginos and sfermions generated via gauge interactions. In Sec. VI we provide a sample of the mass spectrum of our model. We conclude in Sec. VII.
II. META-STABLE SUSY BREAKING

The ISS model is a deformed $N = 1$ supersymmetric $SU(N_c)$ QCD, with $N_f$ massive fundamental flavors. $N_f$ is taken to be in the free magnetic range, $N_c + 1 \leq N_f < \frac{3}{2} N_c$ for a controllable IR description of the theory. For concreteness, we will concentrate on the simplest model $N_f = N_c + 1$ where the magnetic gauge group is trivial.

In terms of the superfields with normalized kinetic terms

$$W = y(\tilde{\psi}\Sigma\psi - \mu^2 \text{Tr} \Sigma),$$

(1)

where $y$ is dimensionless parameter of order one; $\mu$ is related to a holomorphic scale $\Lambda$ of the microscopic theory, but much smaller than $\Lambda$. $\Sigma, \psi, \tilde{\psi}$ are identified as mesons, baryons and antibaryons of the electric theory. In addition, there is an instanton generated operator $det \Sigma/\Lambda^{N_f-3}$. For $N_f > 3$, this term is irrelevant and will be neglected in discussions of the physics around the origin of $\Sigma$. Thus, below we will set $N_f = 4$. The reason for this choice will be discussed later.

The F-terms of the meson field $\Sigma$ are

$$F_{\Sigma ij} = y(\tilde{\psi}_i \psi_j - \mu^2 \delta_{ij}).$$

(2)

SUSY is broken since $\tilde{\psi}_i \psi_j$ has rank one while $\delta_{ij}$ has rank $N_f = 4$. Up to global transformations, the vacua are

$$\langle \Sigma \rangle = 0, \quad \langle \psi \rangle = \langle \tilde{\psi} \rangle = \begin{pmatrix} 0 \\ \mu \end{pmatrix}. \quad (3)$$

In these vacua, the global symmetry is broken to $SU(3) \times U(1) \times U(1)_R$.

To see what the light fields are, we expand around Eq. (3) using the following parametrization

$$\psi = \begin{pmatrix} H \\ \mu + \sigma \end{pmatrix}, \quad \tilde{\psi}^T = \begin{pmatrix} \tilde{H}^T \\ \mu + \tilde{\sigma} \end{pmatrix}, \quad \Sigma = \begin{pmatrix} \Phi & N \\ \tilde{N} & Y \end{pmatrix},$$

(4)

where the component fields transform under the unbroken global symmetry as

$$\begin{array}{c|cccc|cc|cc|c}
 & H & \tilde{H} & \sigma & \tilde{\sigma} & \Phi & N & \tilde{N} & Y \\
\hline
SU(3) & 3 & 3 & 1 & 1 & Ad+1 & 3 & 3 & 1 \\
U(1) & 1/3 & -1/3 & 0 & 0 & 0 & 1/3 & -1/3 & 0 \\
U(1)_R & 0 & 0 & 0 & 0 & 2 & 2 & 2 & 2 \\
\end{array}$$

and $\text{Tr} \Phi$ has a non-zero F-term

$$F_{\text{Tr} \Phi} = y \mu^2. \quad (5)$$

Another two fields, corresponding to the Cartan generators in the adjoint representation, have nonzero F-terms as well.

To identify the Goldstone fields, we use the following non-linear parametrization of scalar fields

$$\psi = e^{\frac{\sigma}{\sqrt{2}}} e^{\frac{\mu}{\sqrt{2}}} \Pi_L e^{\frac{\mu}{\sqrt{2}}} \Pi_K \left( 0, \mu + \frac{\sigma}{\sqrt{2}} \right)^T, \quad \tilde{\psi} = \left( 0, \mu + \frac{\sigma}{\sqrt{2}} \right) e^{\frac{\mu}{\sqrt{2}}} \Pi_K e^{-\frac{\mu}{\sqrt{2}}} \Pi_K, \quad (6)$$

with

$$\Pi_L = \begin{pmatrix} 0 & L \\ L^T & 0 \end{pmatrix}, \quad \Pi_K = \begin{pmatrix} 0 & K \\ K^T & 0 \end{pmatrix}.$$ 

To the leading order, $L$ and $K$ are related to $H, \tilde{H}$ as

$$L = \frac{1}{i\sqrt{2}} \left( H + \tilde{H} \right), \quad K = \frac{1}{i\sqrt{2}} \left( H - \tilde{H} \right). \quad (7)$$

The scalar mass spectrum is

$$m_{\sigma+}^2 = m_{\tilde{K}}^2 = m_{\tilde{\sigma}}^2 = 2g^2 \mu^2,$$

$$m_N^2 = m_{\tilde{N}}^2 = y^2 \mu^2,$$

$$m_{\Phi}^2 = m_{\tilde{\Phi}}^2 = m_L^2 = 0,$$

(8)

where $K$ is the heavy combination of the triplets $H, \tilde{H}$, while the massless combination $L$ together with the singlet $Im(\sigma_-)$ are the seven Nambu-Goldstone bosons (NGB)’s in the coset space $SU(4) \times U(1)^{4}/SU(3) \times U(1)$. $Re(\sigma_-), \Phi$ are pseudo moduli and obtain masses of order $g^2 \mu^2/(16\pi^2)$ through the one-loop Coleman-Weinberg potential

$$V_{\text{CW}}^{(1)} = \frac{1}{64\pi^2} \text{Str} M^4 \log \frac{M^2}{\Lambda^2}. \quad (9)$$

In the UV regime, the instanton term in the superpotential cannot be neglected and is crucial to generate supersymmetric vacua $\langle \Sigma \rangle \sim (\mu^2 \Lambda)^{\frac{3}{4}}$. The distance in field space between the local vacua and the global minima is controlled by a small parameter $\mu/\Lambda$, and thus the metastable vacuum can have a lifetime much longer than the age of the Universe.

III. THE MODEL

The ISS model provides a way to break SUSY and can be used to construct a gauge mediation model, where the soft masses of gauginos and sfermions are roughly of the same scale, $O(g_f^2 F/(16\pi^2 M))$. Here, $g_i (i = 1, 2, 3)$ are the gauge coupling constants in the standard model. $F$ is the value of an F-term indicating SUSY breaking and $M$ is the supersymmetric mass of the messenger. As the superpartner masses are of order 100 GeV or more, we need to have $F/M \gtrsim 10$ TeV. To avoid tachyonic
directions of the messenger fields, 
\( F < M^2 \) and the lowest allowed scale of \( F \) is of order \((10 \text{ TeV})^2\).

Since the Higgs boson mass is likely to be \( O(100 \text{ GeV}) \), it has to be at least two-loop factor below the SUSY breaking scale in the ISS model. Below the SUSY breaking scale, we introduce the collective symmetry breaking mechanism as in little Higgs models to protect the mass of the Higgs boson. In the simplest little Higgs model \([11]\), the Higgs boson mass is lighter than the \( SU(3)'s \) breaking scale \( f \approx 1 \text{ TeV} \) by a factor of \( 4 \pi \). For our purpose, we need to achieve the \( SU(3)'s \) breaking scale \( f \) one loop factor lower than the SUSY breaking scale \( \sqrt{F} \approx 10 \text{ TeV}. \)

To achieve this purpose and to have enough light degrees of freedom containing the Higgs doublet as a PNGB, we adopt two ISS sectors with \( U(4)^2 \) global symmetry. For simplicity, we choose two identical ISS sectors by imposing a \( Z_2 \) symmetry between them. In each ISS sector, the global symmetry \( U(4) \) and SUSY are broken at the same scale, \( \sqrt{F} \approx \mu \), where the coupling \( y \) in the ISS sectors is chosen to be of order one. Then there are two massless triplets \( L1 \) and \( L2 \), which are the NGB’s in each sector. Hence, the effective field theory below \( \mu \) resembles the simplest little Higgs model with the unbroken global symmetry \( U(3)^2 \). After gauging the diagonal \( SU(3)_W \times U(1)_X \) subgroup, \( L1 \) and \( L2 \) become PNGB’s. We will later introduce a superpotential to spontaneously break the approximate global \( U(3)^2 \) symmetry to \( U(2)^2 \).

The Higgs doublet is a PNGB and contained in the light triplets as
\[
L_1^T = f \begin{pmatrix} i \frac{h}{|h|} \sin \frac{|h|}{2f} & \cos \frac{|h|}{2f} \end{pmatrix},
L_2^T = f \begin{pmatrix} -i \frac{h}{|h|} \sin \frac{|h|}{2f} & \cos \frac{|h|}{2f} \end{pmatrix},
\]
where \( U(3)_1 \) and \( U(3)_2 \) are broken at the same scale \( f \); \( \sqrt{2} \) is chosen to have properly normalized kinetic term of the Higgs doublet.

A. Higgs sector and D-term

For the two identical ISS sectors, the superpotential is
\[
W = y (\text{Tr} \tilde{\psi}_1 \Sigma_1 \psi_1 - \mu^2 \text{Tr} \Sigma_1) + y (\text{Tr} \tilde{\psi}_2 \Sigma_2 \psi_2 - \mu^2 \text{Tr} \Sigma_2) .
\]

The symmetry breaking and the vacua in both sectors are the same as described in section \([11]\).

The parametrization around the ISS vacua are then
\[
\psi_i = \begin{pmatrix} H_i \\ \mu + \sigma_i \end{pmatrix}, \tilde{\psi}_i^T = \begin{pmatrix} H_i^T \\ \mu + \tilde{\sigma}_i \end{pmatrix}, \Sigma_i = \begin{pmatrix} \Phi_i & N_i \\ \tilde{N}_i & Y_i \end{pmatrix}
\]
where \( i = 1, 2 \). The unbroken global symmetry is \([SU(3)_1 \times U(1)_1] \times [SU(3)_2 \times U(1)_2]\). The diagonal subgroup of this global symmetry is gauged and denoted as \( SU(3)_W \times U(1)_X \) with the electroweak gauge group as a subgroup. Under the gauge symmetries, the field content in the Higgs sector is described in Table\([11]\).

| \(SU(3)_W\) | \(H_i\) | \(\tilde{H}_i\) | \(\sigma_i\) | \(\tilde{\sigma}_i\) | \(\Phi_i\) | \(N_i\) | \(\tilde{N}_i\) | \(Y_i\) |
|---|---|---|---|---|---|---|---|---|
| \(SU(3)_c\) | 1 | 1 | 1 | 1 | 1 | 1 | 1 | 1 |
| \(U(1)_X\) | 1/3 | -1/3 | 0 | 0 | 1/3 | -1/3 | 0 |

TABLE I: Field content in the Higgs sector. \(SU(3)_c\) is the QCD group.

The D-term potential of \(U(1)_X\) preserves the global \(SU(3)_1 \times SU(3)_2\) symmetry, so it will not give a potential to the Higgs doublet and we neglect it here. To study the Higgs doublet potential, we ignore fields \(N_i\), \(\tilde{N}_i\) and \(\Phi_i\) and only consider the following part of the \(SU(3)_W\) D-term potential
\[
V_D = \frac{g_2^2}{2} \sum_a (\sum_i H_i^a H_i - \tilde{H}_i^a \tilde{H}_i)^2 ,
\]
\[
= \frac{g_2^2}{2} \left( (L_1^a L_1)(K_1^a K_2) + (K_1^a L_2)(K_2^a L_1) \right) + \text{h.c.} \right) + \cdots
\]
where in the second line we expand the D-term potential in terms of \(L_1, K_1\), which are defined in section \([11]\). \(g\) is the \(SU(3)_W\) gauge coupling; \(t_a\) are the generators of \(SU(3)_W\) and \(\text{Tr} [t^a, \theta] = \frac{1}{2} \theta \lambda^{ab}\); the dots represent other quartic terms of \(L_1\) and \(L_2\). Substituting \(K_i = 0\), the D-term does not provide a self-interaction potential of \(L_i\) and the Higgs doublet at tree level (the Higgs doublet is embedded in \(L_1\) as in Eq. \([10]\)). Actually, this result comes from the same vacuum expectation values (VEV’s) of \(\psi_i\) and \(\tilde{\psi}_i\) in the ISS model. After integrating out the heavy modes, the one-loop effective potential of the two light triplets from the gauge interaction is of the form
\[
V = O(\frac{g^2}{16\pi^2}) (\mu^2 |L_1|^2 + \mu^2 |L_2|^2 + |L_1|^4 + |L_2|^4
+ |L_1|^2 |L_2|^2 + |L_1^2 L_2|^2) + \cdots
\]
where order one numbers are neglected in the parenthesis. Substituting the parameterization of \(L_1\) and \(L_2\) in Eq. \([10]\) into this potential, the Higgs doublet mass is \(O(g/4\pi f)\) and it is at most of order 100 GeV.

B. Breaking \(SU(3)_W \times U(1)_X\) to \(SU(2)_w \times U(1)_Y\)

To spontaneously break the global \(SU(3)\) symmetries, we employ the trick of adding vector-like fermions that has been widely used in the little Higgs model building. The superpotential we propose is
\[
\sum_{i=1}^2 [y_1 Q_i \psi_i \rho_i + y_2 \text{Tr} \Phi_i T_i \tilde{T}_i + y_3 (\mu T_i X_i^c + \mu T_i^c X_i)
+ y_4 \mu T_i T_i^c],
\]
\[(15)\]
where the superfields $Q_1 \equiv (Q, T_1)$ and $Q_2 \equiv (Q, T_2)$ with $Q \equiv (t, b, p)$. Here $t$ and $b$ are top and bottom left-handed quarks. The charge assignments of those new fields are listed in Table II. We only list fields in the top sector, which provides the dominate contribution to the Higgs doublet mass. The full list of fields without gauge anomalies can be found in [5].

In order not to change the “rank condition” of the ISS model and not to generate VEV’s for the SU(3)$_c$ charged fields, the following relation among Yukawa couplings for charged fields, the following relation among Yukawa couplings for

\begin{equation}
L_i \quad T_i \quad T_i^c \quad \vec{p}_i \quad X_i \quad X_i^c
\end{equation}

\begin{tabular}{|c|c|c|c|c|c|c|}
\hline
SU(3)$_W$ & SU(3)$_C$ & U(1)$_X$ & 1/3 & 1/3 & -1/3 & -2/3 & 1/3 & -1/3 \\
\hline
3 & 3 & 3 & 3 & 3 & 3 & 3 & 3 & 3 \\
\hline
\end{tabular}

\begin{table}
\caption{Field content in the top quark sector.}
\end{table}

The potential of those two triplets is

\begin{equation}
V_{\text{triplets}} = V_{\text{CW}}(L_1, L_2) + \frac{y_1^2 |L_1|^4}{4} + \frac{y_2^2 |L_2|^4}{4}.
\end{equation}

The fermion and scalar mass matrices are too complicated to be diagonalized analytically, and hence we present a numerical study here. Choosing $y = y_2 = y_4 = 1$, $y_1 = y_3 = 2$ (satisfying the condition in Eq. (16)) and $y_5 = 0.5$, we plot the triplet potential along the $L_1$ direction in Fig. 1. Fig. 1 shows that the triplet $|L_1|$ develops a nonzero VEV at 0.09\mu, which parametrically is

\begin{equation}
\langle |L_1| \rangle = f = O(\mu/4\pi),
\end{equation}

and breaks $U(3)_1$ to $U(2)_1$. A similar result can be obtained for $L_2$ to break $U(3)_2$ to $U(2)_2$.

A few comments are in order about the superpotential in Eq. (15) and Fig. 1:

- The scale $O(10^{-6})$ on the vertical axis of Fig. 1 comes from the product of the one-loop suppression factor and small values of $\langle L_1 \rangle$ near the minimum.
- The potential of $L_i$ is periodic as $L_i$ are PNBG's. It is found that the minimum in Fig. 1 is the global one, so we only consider the parameter space near the origin.
- If the operator $L_i^{\dagger}L_2$ were generated in our effective potential, naive dimensional analysis suggests it would give an $O(f^2)$ contribution to the Higgs doublet mass. However, our superpotential preserves the $U(1) \times U(1)_2$ symmetry which forbids the $L_1^{\dagger}L_2$ term.
- The gauge couplings are smaller than the Yukawa couplings in Eq. (15), so we neglect the contributions from the gauge interactions to the potential of $L_1$ in Fig. 1.
- Logarithmic divergence for the one loop potential of the triplets from Eq. (15) is absent. This is due to the double protection of SUSY and the collective symmetry breaking mechanism on the masses of PNBG's or triplets here. Without SUSY, more than one Yukawa coupling is needed to generate a potential for PNBG's, and then there is not any quadratic divergent potential for the PNBG's at one loop. SUSY protects the mass of PNBG's furthermore and leaves us one finite potential for those triplets at one loop.

After $L_1$ and $L_2$ develop nonzero VEV's as $L_1^T = f (0, 0, 1)$ (the alignment issue will be discussed later), the global symmetry $U(3)^2$ spontaneously breaks to $U(2)^2$ and generates two doublets as NGB's. The gauge symmetry $SU(3)_W \times U(1)_X$ spontaneously breaks to $SU(2)_W \times U(1)_Y$. One of those two NGB doublets is
Therefore, the coefficient and we choose it to be $O$ with a positive assumption that the VEV’s of $L_1$ and $L_2$ are aligned and parameterized in Eq. (10). The D-term also gives an additional mass to one linear combination of $K_1$ and $K_2$ to complete the “super Higgs mechanism”.

IV. ELECTROWEAK SYMMETRY BREAKING

Other than providing negative mass terms to $L_1$ and $L_2$, the one-loop Coleman-Weinberg potential of $L_i$ calculated from Eq. (10) also contains an operator, $|L_1L_2^\dagger|^2$, with a positive $O(1/(16\pi^2))$ coefficient. For example, using the same values of Yukawa couplings as in the previous section, the coefficient is $+0.5/(16\pi^2)$. Substituting the parametrization of $L_1$ and $L_2$ in terms of the Higgs doublet in Eq. (10), we have

$$V_{CW}(L_1, L_2) = O(\frac{1}{16\pi^2})|L_1L_2^\dagger|^2 + \cdots,$$

$$= -O(\frac{f^2}{16\pi^2})hh^\dagger + \cdots.$$  (20)

Here $L_1L_2^\dagger = f^2 - hh^\dagger + \cdots$. The Higgs doublet mass is negative and the electroweak symmetry successively breaks to $U(1)_{em}$. However, this result is based on the assumption that the VEV’s of $L_1$ and $L_2$ are aligned (the VEV’s of both triplets have only nonzero values at the third entries). Therefore, we need to address the vacuum alignment issue first.

A. Vacuum alignment

Although the operator $|L_1L_2^\dagger|^2$ with a positive coefficient gives a negative mass to the Higgs doublet, it prefers anti-aligned VEV’s of $L_1$ and $L_2$. Thus, we need to introduce a new operator together with $|L_1L_2^\dagger|^2$ to achieve the alignment and a negative mass for the Higgs doublet at the same time. One possibility is to include an operator, $y'Z_3(H_1H_2 - \mu^2)$, which is the only operator explicitly breaking one linear combination of $U(1)_1 \times U(1)_2$. Therefore, the coefficient $y'$ can be as small as possible and we choose it to be $O(1/16\pi^2)$. Combining the $F$-term potential of $Z_3$ and $O(1/16\pi^2)|L_1L_2^\dagger|^2$, we have $SU(3)_1 \times SU(3)_2$ breaking parts of the potential as

$$V = -v_1^2(L_1L_2^\dagger + L_2L_1^\dagger)/2 + \frac{v_2^2}{f^2}|L_1L_2^\dagger|^2 + \cdots,$$  (21)

where we define $v_1 \equiv y'\mu = O(1/16\pi^2)\mu = O(1/4\pi)f$ and the positive $O(1/16\pi^2)$ coefficient in front of $|L_1L_2^\dagger|^2$ as $v_2^2/f^2$. For this potential, under the following condition

$$v_2^2 < v_1^2 < 2v_2^2,$$  (22)

the VEV’s of $L_1$ and $L_2$ are aligned and at the same time the Higgs doublet has a negative mass

$$-\frac{1}{2}(4v_2^2 - 2v_1^2)hh^\dagger.$$  (23)

B. Order one quartic Higgs coupling

To obtain the electroweak symmetry breaking scale at 246 GeV, we need to introduce an order one quartic Higgs coupling, which is difficult without generating a large Higgs mass term. This problem is present not only in a realistic super little Higgs model but also in the simplest little Higgs model. In our model, we use the “sliding singlet mechanism” \(^12\) to reach this goal by including the following renormalizable superpotential:

$$y_6[Z_4(S_1^2 - \bar{H}_1H_1) + Z_5(S_2^2 - \bar{H}_2H_2) + Z_6(S_3S_4 - \bar{H}_1H_1) + Z_7(S_3S_4 - \bar{H}_2H_2) + Z_8(S_1S_3 - \bar{H}_1H_2) + Z_9(S_2S_4 - \bar{H}_2H_1)].$$  (24)

By assigning appropriate charges to $Z_i$ and $S_i$, this superpotential also preserves the global $U(1)_1 \times U(1)_2$ symmetry. The $F$-term potential of the $Z_i$’s is

$$y_6^2(|S_1^2 - L_1L_2^\dagger/2|^2 + |S_2^2 - L_2L_1^\dagger/2|^2 + |S_3S_4 - L_2L_2^\dagger/2|^2 + |S_2S_4 - L_1L_2^\dagger/2|^2).$$  (25)

Minimizing this potential, one can see that the VEV’s of $S_i$ cancel the VEV’s of $L_i$. After substituting the parametrization of $L_i$ in Eq. (10), we have a potential of the Higgs doublet without any VEV’s. Hence, the quartic coupling of the Higgs doublet is of order one, and with a negative mass from Eq. (23), the Higgs doublet develops a VEV at $O(1/4\pi)f$.

C. Complete triplet potential

Up to now, we have first studied the spontaneously breaking of $SU(3)_W \times U(1)_X$ to $SU(2)_w \times U(1)_Y$, and then studied the Higgs doublet potential to break the electroweak symmetry. On the other hand, we can also study the full $SU(3)_W \times U(1)_Y$ invariant potential of the two light triplets $L_i$, minimize their potential and derive the final VEV structure. Combining Eqs. (15), (21), (25) and choosing $y_5 = y_6 = 1/2$ (the reason to choose a smaller $y_6$ than the Yukawa couplings in Eq. (15) is the same as $y_5$), we have the following complete potential

$$V = \frac{(|L_1|^2 - f)^2}{16} + \frac{(|L_2|^2 - f)^2}{16} - \frac{v_1^2(L_1L_1^\dagger + L_2L_2^\dagger)}{2} + \frac{v_2^2}{f^2}|L_1L_2^\dagger|^2 + (|S_1^2 - L_1L_2^\dagger|^2 + |S_2^2 - L_2L_1^\dagger|^2 + |S_3S_4 - L_1L_2^\dagger|^2 + |S_2S_4 - L_2L_2^\dagger|/2)^2/4,$$  (26)

with $f = O(1/4\pi)\mu$ and $v_1, v_2 = O(1/4\pi)f$ determined by the Yukawa couplings in Eq. (15). The crucial “−” signs in the first two terms and the “+” sign before the fourth term are also derived from Eq. (15). Here, we neglect the contributions to the effective potential from the
gauge interaction, due to the smallness of the gauge coupling compared to the Yukawa couplings. For simplicity, we neglect the quartic operator $|L_1|^2|L_2|^2$ generated at one loop, which does not change our final result significantly.

Minimizing this potential, we derive

$$\langle |S_i| \rangle = \sqrt{\frac{f^2 + 8v^2 + 6v_1^2}{2 + 40v_1^2/f^2}},$$

$$\langle L_1 \rangle^T = f' \left( i \sin \left( \frac{h}{\sqrt{2}f} \right), 0, \cos \left( \frac{h}{\sqrt{2}f} \right) \right),$$

$$\langle L_2 \rangle^T = f' \left( -i \sin \left( \frac{h}{\sqrt{2}f} \right), 0, \cos \left( \frac{h}{\sqrt{2}f} \right) \right),$$

where

$$f' = \sqrt{\frac{f^2 + 12v_2^2 + 4v_1^2}{1 + 20v_1^2/f^2}},$$

$$\langle |h| \rangle = \sqrt{2f'} \arctan \sqrt{\frac{6v_2^2 - 3v_1^2}{f^2 + 6v_2^2 + 7v_1^2}}. \quad (28)$$

Approximately we have

$$f' \approx f,$$

$$\langle |h| \rangle \approx 12v_2^2 - 6v_1^2 \equiv v. \quad (29)$$

Eqs. (23), (29) show a relation between the Higgs doublet mass and its VEV as $m_h = v/\sqrt{3}$. Using $v = 246$ GeV, the Higgs boson mass is around 150 GeV. Actually the ratio of the Higgs boson mass over $v$ depends on order one Yukawa couplings in our model. Thus, the Higgs boson mass can vary from 100 GeV to 200 GeV by choosing $y_t$ from 1/3 to 2/3.

After electroweak symmetry breaking, the lightest quark in the top sector, which is identified as the top quark in the standard model, has top Yukawa coupling

$$y_t \approx \sqrt{\frac{y_1^2 + y_3^2}{2(y_1^2 + y_3^2)}}. \quad (30)$$

Choosing the same numerical values $y_1 = y_3 = 2$ as in Fig. 1, $y_t \approx 1$.

V. DIRECT GAUGE MEDIATION

In our model, the soft masses of gauginos and sfermions arise from gauge mediation. The messenger fields are embedded inside the SUSY breaking sector, and therefore our model belongs to a direct gauge mediation model.

The couplings in the superpotential in Eq. (15) explicitly break the continuous $U(1)_R$ symmetry to $R$ parity. We calculate the one-loop Coleman-Weinberg potential of the pseudo-moduli $\text{Tr} \Phi_i$, and find that they develop non-zero VEV’s as $\langle \text{Tr} \Phi_i \rangle = O(\mu)$ (similar results can be found in [3]). In our model, it is crucial to have non-zero VEV’s of $\text{Tr} \Phi_i$ to generate masses for $SU(3)_W$ and $U(1)_X$ gauginos, which include winos and bino. But the masses of gluinos can be generated independent of $\langle \text{Tr} \Phi_i \rangle$.

For the gluino masses, $T_i$, $T_i^c$, $X_i$, and $X_i^c$, are the “messengers” in our model. We only need to consider the last four operators in Eq. (15)

$$\begin{pmatrix} T_i & X_i \\ y_1' \mu & y_3' \mu & 0 \end{pmatrix},$$

where $y_i' \equiv y_i + y_2(\langle \text{Tr} \Phi_i \rangle)/\mu$. The SUSY breaking mass terms of the messenger fields are given as $y_2 F_{\text{Tr} \Phi_i}(T_i T_i^c + h.c.)$. From the standard one-loop Feynman diagram calculation, we have the formula for the masses of gluinos as

$$m_{\text{gluino}} = 2 \frac{g_3^2}{16\pi^2} \{ \cos \alpha_1 \cos \alpha_2 [B(m_1, M_1^2, M_2^2) + B(m_2, M_2^2, M_2^2)] + \sin \alpha_1 \sin \alpha_2 [B(m_1, M_1^2, M_2^2) + B(m_2, M_2^2, M_2^2)] \}, \quad (32)$$

where the factor of 2 comes from the presence of two ISS sectors and the function $B(a, b^2, c^2)$ is defined as

$$B(a, b^2, c^2) = a \left( b^2 \log \frac{b^2}{a^2} + c^2 \log \frac{c^2}{a^2} \right).$$

Here, $m_i$ are fermion masses and $M_i$ are scalar masses. Corresponding expressions are

$$m_{1,2} = \frac{y_4 + \sqrt{y_4^2 + 4y_3^2}}{2} \mu,$$

$$M_{1,3}^2 = \frac{-F + m_1^2 + m_2^2}{2} + \frac{\sqrt{F^2 + (m_1^2 - m_2^2)^2 + 2F(m_1^2 - m_2^2) \cos 2\theta}}{2},$$

$$M_{2,4}^2 = \frac{-F + m_1^2 + m_2^2}{2} + \frac{\sqrt{F^2 + (m_1^2 - m_2^2)^2 + 2F(m_1^2 - m_2^2) \cos 2\theta}}{2},$$

with

$$\theta = \frac{1}{2} \arctan \frac{2y_3}{y_4},$$

$$\alpha_{1,2} = \frac{1}{2} \arctan \frac{\pm F \sin 2\theta}{m_2^2 - m_1^2 \pm F \cos 2\theta},$$

$$F \equiv y_2 F_{\text{Tr} \Phi_i} = y_2 y_\mu^2.$$
which agrees with $^{13}$. In terms of the fermion masses, the leading term of gluino masses is

$$\frac{g_3^2}{8\pi^2} F^3 \{[m_i^2 - m_j^2]m_{12} - 4m_1m_2 + m_3^2 + m_4^2]m_6^2 + m_7^2 + m_8^2; 8m_5^2m_7^2 + m_6^2 + m_2^2 + m_4^2 + m_1^2 + m_2^2 + 12m_1^2m_2^2(m_1^2 + m_2^2) \log (m_1^2/m_2^2)\}/(6m_7^2m_2^2(m_1 - m_2)^2(m_1 + m_2)^4)\right\}. \quad (33)$$

- When $y_3 = 0$, this is the simplest case of gauge mediation models with only one messenger. In this case, $\alpha_1 = \alpha_2 = 0$, $m_1 = 3y_4\mu$, $m_2 = 0$, $M_1 = M_2 = M_4 = 0$. After algebraic manipulations, the masses of the gluino are $m_{\text{gluino}} = 2g_3^2/16\pi^2 \times F/m_1 \times g(F/m_1)^2$, with $g(x) = [(1 + x) \log (1 + x) + (1 - x) \log (1 - x)]/x^2$. This agrees with the result in the literature $^{14}$.

- When $y_4 = 0$, the continuous $U(1)_R$ symmetry is unbroken. We have $m_{12} = \pm y_3\mu$, $\theta = -\pi/4$ and $\alpha_1 = \alpha_2 = \pi/4$. Considering that $B(a, b^2, c^2) = B(-a, b^2, c^2)$, we have $m_{\text{gluino}} = 0$.

In our model, all the Yukawa couplings are generally order one numbers, so $F$ and $\mu^2$ are at the same scale. With $\mu = O(10 \text{ TeV})$ and $\sqrt{F} = O(10 \text{ TeV})$, we have

$$m_{\text{gluino}} = \frac{2g_3^2}{16\pi^2} O\left(\frac{F^3}{\mu^5}\right) = O(200 \text{ GeV}). \quad (34)$$

For the $SU(3)_W$ gauginos, the mixed messengers are $H_i, \tilde{H}_i, N_i, \tilde{N}_i$. Considering the non-zero VEV $\langle \text{Tr} \Phi \rangle$ and applying the same formula in Eq. $^{12}$ by substituting $y_4 \rightarrow y/\langle \text{Tr} \Phi \rangle/\mu$, $y_3 \rightarrow y$ and $F \rightarrow y\mu^2$, we have the masses of those $SU(3)_W$ gauginos to be also of $O(200 \text{ GeV})$. Similar results are derived for the $U(1)_X$ gaugino with two sets of mixed messengers: one is $T_i, T'_i, X_i$ and $X'_i$; the other one is $H_i, \tilde{H}_i, N_i, \tilde{N}_i$. The winos and bino are parts of $SU(3)_W \times U(1)_X$ gauginos, and hence have masses of $O(200 \text{ GeV})$.

Finally, the masses of squarks and sleptons are generated through the traditional two-loop diagrams and are also of $O(200 \text{ GeV})$.

At the $U(3)^2$ breaking scale $f = O(1 \text{ TeV})$, there are the pseudo-moduli $\Phi_i$ and their fermionic partners containing electroweak triplets and doublets, and massive vector superfields $W's$ and $Z'$ corresponding to $SU(3)_W \times U(1)_X$. Colored superfields $p_i, p'_i$ also have masses of this scale.

The gauginos and sfermions in the minimal supersymmetric standard model obtain masses from direct gauge mediation as described in the previous section and depend on order a few hundred GeV. Finally, we comment on the masses of the standard model singlets in our model. Most singlets in the ISS sectors, $Z_i$ and $S_i$ have masses at or above $1 \text{ TeV}$. The singlet PNGB in the coset space $U(1)_1 \times U(1)_2/\text{det} \times U(1)_Y$ obtains a mass of $O(100 \text{ GeV})$ from the operator $y'Z_3(H_1H_2 - \mu^2)$ added to achieve the vacuum alignment.

In summary, we list a sample of the mass spectrum of the particles charged under the standard model gauge group in Table $^{[III]}$.

| Mass (GeV) | Description |
|-----------|-------------|
| $\sim 20 \text{ TeV}$ | heavy colored superfields $T, T^c, X, X^c$ |
| $\sim 1 \text{ TeV}$ | heavy singlets $K_i$ |
| $250 \sim 500 \text{ GeV}$ | triplets and doublets in $\Phi_i$ |
| $\sim 100 \text{ GeV}$ | vector superfields: $W's$ and $Z'$ |
| $\sim 100 \text{ GeV}$ | colored superfields $p, p'^c$ |
| $\sim 100 \text{ GeV}$ | singlets $N_i, \tilde{N}_i$ |
| $\sim 100 \text{ GeV}$ | the Higgs boson |

TABLE III: A sample of the mass spectrum of the particles in our model. This is only the order of magnitude estimate. The Higgs boson mass can vary from $100 \text{ GeV}$ to $200 \text{ GeV}$ depending on the order one Yukawa couplings in our model.

The scales in Table $^{[III]}$ are only an order-of-magnitude estimate. The detailed spectrum depends on order one parameters in the model, and subject to constraints from electroweak precision measurements. To estimate the constraints, we compare our model with the simplest little Higgs model. Previous analyses on the simplest little Higgs model $^{15}$ yield lower bounds for the masses of $W's$, and then a lower bound for $\sqrt{F_1^2 + F_2^2} \sim f$ about $3 \sim 6 \text{ TeV}$, depending on different choices of fermion charge assignments. Here, $f_1$ and $f_2$ are the magnitudes of the two $U(3)$ breaking VEV’s (we have assumed $f_1 = f_2 = f$ for the sake of simplicity throughout the paper, but this is not essential for our model). As discussed below Eq. $^{20}$, the Higgs mass remains variable when the scale of $f$ is fixed. Therefore, we can achieve $f \gtrsim 3 \text{ TeV}$ without introducing significant fine tuning. At and below the TeV scale, in addition to particles present in
the simplest little Higgs model, there are extra scalars and superparticles from the SUSY breaking sector. The TeV scale scalars contribute to electroweak observables only at loop levels unless the $SU(2)_W$ triplets develop VEV’s to break the custodial symmetry \(16\), which is not the case in our model. The superparticles contribute to the electroweak observables only at loop levels. The full analyses of their effects again depend on order one parameters and are beyond the interests of this paper.

A lot of particles in our model are at or below the TeV scale and thus can be produced at the upcoming LHC. The spectrum bears a resemblance of the super-little Higgs models, namely, there exist heavy gauge bosons and fermions predicted by the little Higgs mechanism, and superpartners of the SM fields predicted by SUSY. Indeed, our model can be viewed as a UV completion of the super-simplest little Higgs model. However, there are also important differences between our model and previous super-little Higgs models. First, as discussed above, the Higgsino is necessarily absent from TeV scale spectra. Second, our model predicts extra scalars from the Higgsino sector and its effect is only mediated to the standard model Higgs doublet through messenger fields. In this paper, we have explored another approach which is to identify the Higgs sector plays the role of a messenger field, which distinguishes our approach from traditional gauge mediation models.

The Higgs doublet, which is a PNGB inside \(L_1\) and \(L_2\), has a negative mass with another loop factor below the triplet masses and triggers the electroweak symmetry breaking. With order one quartic potential generated through the sliding singlet mechanism, the correct Higgs boson mass can vary from 100 GeV to 200 GeV depending on order one Yukawa couplings in our model.

VII. CONCLUSIONS

The SUSY breaking sector is usually treated as a “hidden” sector and its effect is only mediated to the standard model through messenger fields. In this paper, we have explored another approach which is to identify the meta-stable SUSY breaking sector as an extended Higgs sector. In our model, the standard model Higgs doublet emerges from the meta-stable SUSY breaking sector as a PNGB, and the electroweak scale is naturally two-loop factor below the SUSY breaking scale.

We have employed two identical meta-stable SUSY breaking sectors with \(U(4)^2\) global symmetry. At order 10 TeV, SUSY is broken and the global symmetry is broken to \(U(3)^2\). Most fields in the SUSY breaking sectors have tree-level masses of order 10 TeV. Among those massless Nambu-Goldstone boson fields, there are two triplets \(L_1\) and \(L_2\), in which the standard model Higgs boson sits.

The electroweak gauge group, \(SU(2)_w \times U(1)_Y\), is extended to a \(SU(3)_W \times U(1)_X\) gauge group embedded in the unbroken global symmetry. The gauge interaction explicitly breaks the global \(U(3)^2\) symmetry and makes the triplets \(L_1\) and \(L_2\) PNGB’s. Introducing additional vector-like quarks in the top sector, we have studied the one-loop effective potential of \(L_1\) and \(L_2\) and found that \(L_1\) and \(L_2\) obtain negative masses. With order one quartic potential, \(U(3)^2\) spontaneously breaks to \(U(2)^2\) at the scale of 1 TeV. At the same time the gauge symmetry \(SU(3)_W \times U(1)_X\) is broken to \(SU(2)_w \times U(1)_Y\) and a light doublet PNGB is identified as the Higgs doublet.

At order 1 TeV, aside from the massive vector superfields \(W\)’s and \(Z\)’s and additional colored superfields, there are pseudo-moduli fields including electroweak triplets and doublets from the dynamical SUSY breaking sectors. The soft masses of gauginos and sfermions in the minimal supersymmetric standard model arise at one loop and two loops, respectively, in a similar way as in gauge mediation, and are of order a few hundred GeV. It is interesting that the Higgs sector plays the role of a messenger field, which distinguishes our approach from traditional gauge mediation models.

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