Phenomenological Aspects of a Direct-transmission Model of Dynamical Supersymmetry Breaking with the Gravitino Mass $m_{3/2} < 1$ keV

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Abstract

We analyze a direct-transmission model of dynamical SUSY breaking previously proposed. In the model the gravitino mass is naturally smaller than 1 keV, which is required from the standard cosmology. We find that there are many distinguishable features from in other models: for example the so-called GUT relation among the gaugino masses does not hold even if we consider the GUT models. Furthermore, the gauginos are always lighter than the sfermions since the gaugino masses have extra suppression factors. We also discuss a collider signature "$\gamma\gamma + \text{missing energy}"$ in the present model.
1 Introduction

Low-energy supersymmetry (SUSY) is very attractive since it stabilizes a hierarchy between the weak scale and a higher scale of new physics (say the Planck scale). It is also strongly supported by the observed unification of the standard model (SM) gauge couplings.

Recently, low-energy dynamical SUSY breaking (DSB) with gauge mediation has attracted attention since it may provide a natural explanation of large hierarchy between the weak scale and the Planck scale as well as a natural solution to the SUSY flavor changing neutral current (FCNC) and the CP problems [1, 4].

Several mechanisms for DSB have been discovered [3, 4, 5, 6] and their applications to realistic models have been proposed in the literatures [7, 8, 9]. Most of the models which have been proposed, however, need relatively large DSB scales $\Lambda \gtrsim 10^7$ GeV to get sufficiently large SUSY breaking masses in the minimal supersymmetric standard model (MSSM) sector. As a result, the gravitino mass becomes $m_{3/2} > 10$ keV. On the other hand, the gravitino should be lighter than 1 keV so that it does not overclose the universe [12]. To escape this bound, one may consider a late-time entropy production which, however, leads to a complicated cosmology [13, 14]. Furthermore, as recently pointed out in Ref. [15], a constraint on the cosmic X($\gamma$)-ray backgrounds from a dilaton decay requires that the gravitino should be lighter than 100 keV. To get rid of this bound, we have to construct a string theory without the dilaton. Such a string theory, however, has not been known. Therefore, it is very important to construct a model with $m_{3/2} < 1$ keV and to investigate phenomenological consequences of such a model.

In Ref. [16], a DSB model which transmits SUSY breaking effects to the MSSM sector directly has been proposed. In this model, we can lower the SUSY breaking scale and hence the gravitino mass can be set $m_{3/2} < 1$ keV to avoid an introduction of complicated non-standard cosmology keeping the advantage of gauge mediation. Moreover, such a light gravitino may suggest a distinct signature in the existing collider experiments, that is, the next-to-lightest superparticle (NLSP), mostly bino, can decay into the gravitino, ...

\footnote{In the model in Ref. [8] it seems possible to make the gravitino mass smaller than 1 keV due to strong gauge interaction. If a naive dimensional analysis of strongly coupled SUSY theory [14] is applicable, however, the gravitino mass is likely larger than 1 keV [11].}
which is the lightest superparticle (LSP), within detectors producing a observable \( \gamma\gamma + \) missing energy” signal as discussed in many literatures \[17, 18\]. Even if the signature \( \gamma\gamma + \) missing energy” is not observed within the detectors, the slow decay of the NLSP may be detectable in the near future experiments such as LHC as pointed out in Ref. \[19\]. Furthermore, in the model in Ref. \[16\], the mass spectrum of superparticles in the MSSM sector is quite different from in the ordinary gauge mediation models and in the gravity mediation models based on supergravity. In particular, the grand unified theory (GUT) relation among the gaugino masses does not hold even if we consider the GUT unification. Since the present model has many different features from the other ordinary models \[7, 9, 8\], it may be distinguishable.

The purpose of this paper is to investigate the mass spectrum in the model in Ref. \[16\] imposing experimental constraints and to show the existence of the phenomenologically viable parameter regions in which the gravitino mass is smaller than 1 keV. This paper is organized as follows. In Section 2, we briefly review the model in Ref. \[16\]. In Section 3, we consider the low-energy mass spectrum of the gauginos and sfermions in the MSSM sector, and argue their typical features. The radiative electroweak symmetry breaking is also discussed. In Section 4, we analyze the GUT models with and without the Yukawa unification. In section 5, we discuss other interesting features in the present model. Section 6 is devoted to our conclusions.

2 A direct-transmission model of SUSY breaking

In this section, we review a model which has been proposed in Ref. \[16\]. Let us first consider a dynamics for scale generation. To generate the scale, we adopt a SUSY SU(2) gauge theory with four doublet chiral superfields \( Q_i \), where \( i \) is a flavor index \( (i = 1, \cdots, 4) \). Without a superpotential, this theory has a flavor SU(4)\( _F \) symmetry. This SU(4)\( _F \) symmetry is explicitly broken down to a global SP(4)\( _F \) by a superpotential in this model. We add gauge singlets \( Y^a \) \( (a = 1, \cdots, 5) \) which constitute a five-dimensional representation of SP(4)\( _F \) to obtain a tree-level superpotential

\[
W_Y = \lambda_Y Y^a (QQ)_a, \quad (1)
\]
where \((QQ)_a\) denote a five-dimensional representation of \(\text{SP}(4)_F\) given by a suitable combination of gauge invariants \(Q_i Q_j\).

A low energy effective superpotential with \(W_Y\) in Eq. (1), which describes the dynamics of the \(\text{SU}(2)\) gauge interaction, may be given by

\[
W_{\text{eff}} = S(V^2 + V_a^2 - \Lambda^4) + \lambda_Y Y^a V_a
\]

in terms of low-energy degrees of freedom

\[
V \sim (QQ), \quad V_a \sim (QQ)_a,
\]

where \(S\) is an additional chiral superfield, \(\Lambda\) a dynamically generated scale, and a gauge invariant \((QQ)\) denotes a singlet of \(\text{SP}(4)_F\) defined by

\[
(QQ) = \frac{1}{2}(Q_1 Q_2 + Q_3 Q_4).
\]

The effective superpotential Eq. (2) implies that the singlet \(V \sim (QQ)\) condenses as

\[
\langle V \rangle = \Lambda^2,
\]

and SUSY is kept unbroken in this unique vacuum. Since the vacuum preserves the flavor \(\text{SP}(4)_F\) symmetry, we have no massless Nambu-Goldstone boson. The absence of flat direction at this stage is crucial for causing dynamical SUSY breaking as seen below.

Next we consider the dynamical SUSY breaking [5]. Let us first consider a model with one singlet chiral superfield \(Z\) for SUSY breaking which couples directly to \((QQ)\). Then, the superpotential is given by

\[
W = W_Y + \lambda Z(QQ).
\]

For a relatively large value of the coupling \(\lambda_Y\), we again obtain the condensation Eq. (5) with low-energy effective superpotential approximately given by

\[
W_{\text{eff}} \simeq \lambda \Lambda^2 Z.
\]

Then, \(F_Z \simeq \lambda \Lambda^2 \neq 0\) and SUSY is broken.

On the other hand, the effective Kähler potential is expected to take a form

\[
K = |Z|^2 - \frac{\eta}{4\Lambda^2} |\lambda Z|^4 + \cdots,
\]

3
where we expect that $\eta$ is a real constant of order one. Then the effective potential for the scalar $Z$ (with the same notation as the superfield) is given by

$$V_{\text{eff}} \simeq |\lambda|^2 \Lambda^4 (1 + \frac{\eta \Lambda}{|\lambda|^4} |Z|^2). \quad (9)$$

If $\eta > 0$, this implies $\langle Z \rangle = 0$. Otherwise we expect $|\langle Z \rangle| \sim \Lambda$, since the effective potential is lifted in the large $|Z|$ ($> \Lambda$) region \[5, 8, 21\].

In the following analyses, we assume the latter case $|\langle Z \rangle| \sim \Lambda$, which results in the breakdown of the $R$ symmetry\[3\] The spontaneous breakdown of the $R$ symmetry produces a Nambu-Goldstone $R$-axion. This $R$-axion is, however, cosmologically harmless, since it acquires a mass from the $R$-breaking constant term in the superpotential which is necessary to set the cosmological constant to zero \[22\].

To transmit the SUSY breaking effects to the MSSM sector, we introduce vector-like messenger quark multiplets $d, \bar{d}$ and lepton multiplets $l, \bar{l}$. We assume that the $d$ and $\bar{d}$ transform as the right-handed down quark and its antiparticle, respectively, under the SM gauge group. The $l$ and $\bar{l}$ are assumed to transform as the left-handed lepton doublet and its antiparticle, respectively.\[3\] We introduce coupling of the messenger quarks and leptons to the singlet $Z$ in order to directly transfer the SUSY breaking to the messenger sector. Then the effective superpotential is

$$W_{\text{eff}} \simeq Z(\lambda \Lambda^2 + k_d d \bar{d} + k_l l \bar{l}). \quad (10)$$

In this case, however, the condensations $\langle dd \rangle \neq 0$ and $\langle l \bar{l} \rangle \neq 0$ occur, and hence SUSY is not broken ($F_Z = 0$). To avoid such undesired condensations, we further introduce a pair of messenger quark and lepton multiplets $(d', l') + (\bar{d}', \bar{l}')$, and mass parameters $m_d, m_{\bar{d}}, m_l, m_{\bar{l}}$ as follows:

$$W_{\text{eff}} = Z(\lambda \Lambda^2 + k_d d \bar{d} + k_l l \bar{l})$$

$$+ m_d d d' + m_{\bar{d}} d' \bar{d} + m_l l l' + m_{\bar{l}} l' \bar{l}. \quad (11)$$

\[2\] Even if we take $|\langle Z \rangle| \sim 4\pi \Lambda$ the main conclusion in the present paper does not change.

\[3\] When $\langle Z \rangle = 0$, appropriate $R$-breaking mass terms such as $m d \bar{d} + m' l \bar{l}$ is necessary to give masses to the MSSM gauginos because the $R$ symmetry keeps the gauginos massless.

\[4\] One may consider that the messenger quark and lepton multiplets are embedded into SU(5) GUT multiplets $5$ and $5^*$ to preserve the unification of the SM gauge coupling constants.
Dynamical generation of these mass parameters $m_d, m_{\bar{d}}, m_l, \text{and } m_{\bar{l}}$ has been discussed in Ref. [16]. We will briefly review it below.

Owing to the mass parameters in Eq. (11), we can obtain a SUSY breaking vacuum with vacuum expectation values of the messenger squarks and sleptons vanishing,

$$\langle d \rangle = \langle \bar{d} \rangle = \langle l \rangle = \langle \bar{l} \rangle = \langle d' \rangle = \langle \bar{d}' \rangle = \langle l' \rangle = \langle \bar{l}' \rangle = 0, \quad \langle F_Z \rangle \simeq \lambda \Lambda^2.$$  

(12)

The condition for this desired vacuum to be the true one is given by examining the scalar potential,

$$V = |\lambda \Lambda^2 + k_d d \bar{d} + k_l l \bar{l}|^2 + |m_d d|^2 + |m_d d'|^2 + |m_l l|^2 + |m_l l'|^2 + |k_d Z d + m_d d'|^2 + |k_l Z l + m_l l'|^2,$$

(13)

as follows:

$$|m_d m_{\bar{d}}|^2 > |k_d \langle F_Z \rangle|^2,$$

$$|m_l m_{\bar{l}}|^2 > |k_l \langle F_Z \rangle|^2.$$  

(14)

Then, the soft SUSY breaking masses of the messenger squarks and sleptons are directly generated by $\langle F_Z \rangle \simeq \lambda \Lambda^2$ through the couplings $Z(k_d d \bar{d} + k_l l \bar{l})$. We will see later that such a direct-transmission of the SUSY breaking effects to the messenger sector makes it possible to realize the light gravitino $m_{3/2} < 1 \text{ keV}$, which is required from the standard cosmology.

Now we discuss the dynamical generation of the mass parameters $m_d, m_{\bar{d}}, m_l, \text{and } m_{\bar{l}}$. To generate these mass parameters dynamically, we introduce a singlet $X$ whose vacuum expectation value plays the role of mass parameters. Furthermore, in order to give the vacuum expectation value to $X$ keeping the SUSY breaking, three singlet chiral supermultiplets $Z_i (i = 1, \ldots, 3)$ which couple to $(QQ)$ are introduced as follows:

$$W = W_Y + Z_1 (\lambda_1 (QQ) + k_{d1} d \bar{d} + k_{l1} l \bar{l} - f_1 X^2) + Z_2 (\lambda_2 (QQ) + k_{d2} d \bar{d} + k_{l2} l \bar{l} - f_2 X^2)$$

$$+ Z_3 (\lambda_3 (QQ) + k_{d3} d \bar{d} + k_{l3} l \bar{l} - f_3 X^2) + X (f_{d1} d \bar{d} + f_{d2} d' \bar{d}' + f_{l1} l \bar{l} + f_{l2} l' \bar{l}')$$

(15)

5 We can construct a model in which only two singlet chiral supermultiplets are needed to give the vacuum expectation value to $X$ keeping the SUSY breaking if the GUT unification of the Yukawa couplings holds at the GUT scale. However, we consider three singlets because the Yukawa coupling unification is easily broken by non-renormalizable operators as discussed later.
Here, we should stress that the superpotential Eq.(15) is natural, since it has a global symmetry $U(1)_R \times U(1)_\chi$, where $U(1)_R$ is an $R$ symmetry. That is, the superpotential Eq.(15) is a general one allowed by the global $U(1)_R \times U(1)_\chi$. The charges for chiral superfields are given in Table 1.

Without loss of generality, we may set $k_{d1} = k_{l1} = f_2 = 0$ by an appropriate redefinition of $Z_1, Z_2,$ and $Z_3$. Under the following condition:

$$|f_\psi f_{\bar{\psi}}(\lambda_1 f_1 + \lambda_3 f_3)|^2 > |k_{\psi 2} \lambda_2 (f_1^2 + f_3^2) + k_{\bar{\psi} 3} f_1 (\lambda_3 f_1 - \lambda_1 f_3)|^2$$

for $\psi = d, l$, the superpotential yields a vacuum with

$$\langle X \rangle = \sqrt{\frac{\lambda_1 f_1 + \lambda_3 f_3}{f_1^2 + f_3^2}} \Lambda,$$

and vacuum expectation values of the messenger squarks and sleptons vanish. The condition Eq.(16) corresponds to the vacuum stability condition Eq.(14). In this vacuum, the $F$-components of $Z_i$ are given by

$$F_{Z_1} = \frac{\lambda_1 f_3 - \lambda_3 f_1}{f_1^2 + f_3^2} f_3 \Lambda^2, \quad F_{Z_2} = \lambda_2 \Lambda^2, \quad F_{Z_3} = \frac{\lambda_3 f_1 - \lambda_1 f_3}{f_1^2 + f_3^2} f_1 \Lambda^2,$$

and thus SUSY is broken.

Since the scalar component of $X$ superfield has the vacuum expectation value $\langle X \rangle$, the mass parameters $m_d, m_{\bar{d}}, m_l,$ and $m_{\bar{l}}$ are dynamically generated as

$$m_\psi = f_\psi \langle X \rangle,$$

$$m_{\bar{\psi}} = f_{\bar{\psi}} \langle X \rangle$$

---

Table 1: $U(1)_R \times U(1)_\chi$ charges for chiral superfields. Here, $\psi = d, l$ and $i = 1, 2, 3$.

| $U(1)_R$ | $U(1)_\chi$ |
|----------|-------------|
| 0        | 2           |
| 1        | -2          |
for $\psi = d, l$. Therefore, this model is reduced to the model described in Eq.(11) effectively. In a practical analysis we use the reduced model with Eq.(11). We should note that all of the mass parameters are at the same order of the SUSY breaking scale $\sqrt{F_{Z_i}}$ if the Yukawa couplings $f_i$, $\lambda_i$, $f_{\psi}$, $f_{\bar{\psi}}$ are $O(1)$ because they are generated by the same dynamics with the scale $\Lambda$.

The messenger sfermions receive the SUSY breaking mass squared as $k_{\psi_2} \langle F_{Z_2} \rangle + k_{\psi_3} \langle F_{Z_3} \rangle$. Therefore, the gauginos and sfermions in the MSSM sector acquire their masses through loop diagrams of the messenger multiplets \cite{7, 25, 26}. We will discuss the obtained mass spectrum in the MSSM sector in the next section.

3 Mass Spectrum of the superparticles in the MSSM sector

3.1 Mass Spectrum

In this section, we derive the low-energy mass spectrum of the gauginos, squarks, and sleptons in the MSSM sector. To calculate the masses for the gauginos and sfermions, we first consider the mass eigenstates of the messenger fermions and sfermions. To begin with, the superpotential for the mass terms of the messenger fields $\psi$, $\bar{\psi}$, $\psi'$, and $\bar{\psi}'$ for $\psi = d, l$ is represented as

$$W = \sum_{\psi = d, l} (\bar{\psi}, \bar{\psi}') M^{(\psi)} \begin{pmatrix} \psi \\ \psi' \end{pmatrix},$$

(21)

where the mass matrix $M^{(\psi)}$ is given by

$$M^{(\psi)} = \begin{pmatrix} m^{(\psi)} & m_{\psi} \\ m_{\bar{\psi}} & 0 \end{pmatrix}.$$  

(22)

In the present model, the above mass parameters are given by

$$m^{(\psi)} = k_{\psi_2} \langle Z_2 \rangle + k_{\psi_3} \langle Z_3 \rangle,$$

(23)

$$m_{\psi} = f_{\psi} \langle X \rangle,$$

(24)

$$m_{\bar{\psi}} = f_{\bar{\psi}} \langle X \rangle.$$  

(25)
Then, the messenger quark and lepton masses are given by diagonalizing the mass matrix $M^{(\psi)}$ as follows:

$$\text{diag} \left( M_1^{(\psi)}, M_2^{(\psi)} \right) = U^{(\psi)} M^{(\psi)} V^{(\psi)}\dagger. \quad (26)$$

On the other hand, the messenger squarks and sleptons receive the soft SUSY breaking mass terms,

$$\mathcal{L}_{soft} = \sum_{\psi=d,l} F^{(\psi)} \tilde{\psi} \tilde{\psi}, \quad (27)$$

where

$$F^{(\psi)} = k_{\psi 2} \langle F_{Z_2} \rangle + k_{\psi 3} \langle F_{Z_3} \rangle. \quad (28)$$

Then, the mass terms of the messenger squarks and sleptons are written as

$$\mathcal{L}_s = - \sum_{\psi=d,l} (\tilde{\psi}^*, \tilde{\psi}', \tilde{\psi}, \tilde{\psi}') \tilde{M}_2^{(\psi)} \begin{pmatrix} \tilde{\psi} \\ \tilde{\psi}' \\ \tilde{\psi}^* \\ \tilde{\psi}'^* \end{pmatrix}, \quad (29)$$

where the mass matrix $\tilde{M}_2^{(\psi)}$ is given by

$$\tilde{M}_2^{(\psi)} = \begin{pmatrix} |m^{(\psi)}|^2 + |m_{\psi}|^2 & m^{(\psi)*} m_{\tilde{\psi}} & F^{(\psi)*} & 0 \\ m^{(\psi)} m_{\tilde{\psi}}^* & |m_{\tilde{\psi}}|^2 & 0 & 0 \\ F^{(\psi)} & 0 & |m^{(\psi)}|^2 + |m_{\psi}|^2 & m^{(\psi)} m_{\psi}^* \\ 0 & 0 & m^{(\psi)*} m_{\psi} & |m_{\psi}|^2 \end{pmatrix} \quad (30)$$

for $\psi = d, l$. This can be diagonalized by the unitary matrix $T^{(\psi)}$ as

$$\text{diag} \left( m_1^{(\psi)^2}, m_2^{(\psi)^2}, m_3^{(\psi)^2}, m_4^{(\psi)^2} \right) = T^{(\psi)} \tilde{M}_2^{(\psi)} T^{(\psi)}\dagger. \quad (31)$$

When we take the mass eigenstates of the messenger sector, the interactions for the messenger fields are described in terms of the mixing matrices $U^{(\psi)}$, $V^{(\psi)}$, and $T^{(\psi)}$. (See Appendix A.) Then we calculate the masses for the superparticles in the MSSM sector.

The MSSM gauginos acquire masses through the one-loop diagrams of the messenger quark and lepton multiplets shown in Fig. [1]. Their masses are given by

$$m_{\tilde{g}_3} = \frac{\alpha_3}{2\pi} \mathcal{F}^{(d)},$$

$$m_{\tilde{g}_2} = \frac{\alpha_2}{2\pi} \mathcal{F}^{(l)},$$

$$m_{\tilde{g}_1} = \frac{\alpha_1}{2\pi} \left\{ \frac{2}{5} \mathcal{F}^{(d)} + \frac{3}{5} \mathcal{F}^{(l)} \right\}, \quad (32)$$
where we have adopted the SU(5) GUT normalization of the U(1)_Y gauge coupling ($\alpha_1 \equiv \frac{5}{3}\alpha_Y$), and $m_{\tilde{g}_i}$ ($i = 1, \cdots, 3$) denote the bino, wino, and gluino masses, respectively. Here $\mathcal{F}^{(d)}$ represents contributions from the messenger quark multiplets $d, d', \bar{d}$, and $d''$, and $\mathcal{F}^{(l)}$ contributions from the messenger lepton multiplets $l, l', \bar{l}$, and $l''$. The functions $\mathcal{F}^{(d)} (\psi = d, l)$ are given by

$$\mathcal{F}^{(d)} = \sum_{\alpha=1}^{2} \sum_{\beta=1}^{4} M^{(d)}_{\alpha}(U^{(d)}_{\alpha1}T^{(d)}_{3\beta} + U^{(d)}_{\alpha2}T^{(d)}_{4\beta})(T^{(d)}_{\beta1}V^{(d)}_{1\alpha} + T^{(d)}_{\beta2}V^{(d)}_{2\alpha}) \times \frac{m^{(d)}_{\beta}^2}{m^{(d)}_{\beta}^2 - M^{(d)}_{\alpha}^2} \ln \frac{m^{(d)}_{\beta}^2}{M^{(d)}_{\alpha}^2}. \quad (33)$$

where $M^{(d)}_{\alpha}$ and $m^{(d)}_{\beta}$ denote messenger fermion masses and messenger sfermion masses respectively. Under the condition of Eq. (14) the function $\mathcal{F}^{(d)}$ can be expanded in terms of $F^{(d)}/(m_{\psi}m_{\psi})$ as

$$\mathcal{F}^{(d)} = \frac{1}{2} \left| \frac{F^{(d)}}{m_{\psi}m_{\psi}} \right|^2 \mathcal{A}^{(d)}(\left| V^{(d)}_{11}/V^{(d)}_{12} \right|^2, \left| U^{(d)}_{11}/U^{(d)}_{12} \right|^2), \quad (34)$$

where $\mathcal{A}^{(d)}(a, b)$ is

$$\mathcal{A}^{(d)}(a, b) = \frac{(ab)^{\frac{1}{2}}}{6(1 - ab)^{\frac{1}{2}}(1 + ab)^{\frac{1}{2}} \left( 2(a + b)(-6ab - 8a^2b^2 + a^4b^4 + 12a^2b^2 \ln(ab)) - 6ab + 6a^2b^2 + 6a^3b^3 + a^4b^4 + a^5b^5 - 36a^2b^2(1 + ab) \ln(ab) \right)}. \quad (35)$$

This function $\mathcal{A}^{(d)}(a, b)$ has the maximal value 0.1 at $a \simeq 3$ and $b \simeq 3$. We should note that the leading term of order $F^{(d)}/(m_{\psi}m_{\psi})$ vanishes because $(M^{(d)})_{11}^{-1} = 0$. [10] [11] [12] [13] [14] [15] [16]

\footnote{This leading term cancellation of gaugino masses occurs generically, when we stabilize the SUSY breaking vacuum by mass terms as in Eq. (14).}
In the ordinary gauge mediation models, contributions from the messenger quark multiplets are equal to those from the messenger lepton multiplets in the leading order of $F/m$, and hence the GUT relation among the gaugino masses, $m_{\tilde{g}_1}/\alpha_1 = m_{\tilde{g}_2}/\alpha_2 = m_{\tilde{g}_3}/\alpha_3$, holds. In the present model, however, $F(d) \neq F(l)$ because the leading term of order $F(\psi)/\sqrt{m_\psi m_{\bar{\psi}}}$ is canceled out. Therefore, even if the unification of the Yukawa couplings and mass parameters is assumed at the GUT scale, the GUT relation does not hold in general. From Eq.(32) the following relation among the gaugino masses are satisfied:

$$m_{\tilde{g}_1}/\alpha_1 = \frac{3}{5} m_{\tilde{g}_2}/\alpha_2 + \frac{2}{5} m_{\tilde{g}_3}/\alpha_3.$$  (36)

This is a distinctive prediction in the present model.

The soft SUSY breaking masses for squarks, sleptons, and Higgses $\tilde{f}$ in the MSSM sector are generated by two-loop diagrams shown in Fig.2 [7, 23, 24]. We obtain them as

$$m_f^2 = \frac{1}{2} \left[ C_3^{\tilde{f}} \left( \frac{\alpha_3}{4\pi} \right)^2 G^{(d)2} + C_2^{\tilde{f}} \left( \frac{\alpha_2}{4\pi} \right)^2 G^{(l)2} + \frac{3}{5} Y^2 \left( \frac{\alpha_1}{4\pi} \right)^2 \left( \frac{2}{5} G^{(d)2} + \frac{3}{5} G^{(l)2} \right) \right],$$  (37)

where $C_3^{\tilde{f}} = \frac{4}{3}$ and $C_2^{\tilde{f}} = \frac{4}{3}$ when $\tilde{f}$ is in the fundamental representation of SU(3)$_C$ and SU(2)$_L$, and $C_l^{\tilde{f}} = 0$ for the gauge singlets, and $Y$ denotes the U(1)$_Y$ hypercharge ($Y \equiv Q - T_3$). Here $G^{(d)2}$ and $G^{(l)2}$ represent the contributions from the messenger quark and lepton multiplets, respectively, and they are given in Appendix A in details. In contrast to the gaugino masses, the leading term of order $F(\psi)/\sqrt{m_\psi m_{\bar{\psi}}}$ is not canceled out for the sfermion masses. Therefore, the gaugino masses have an extra suppression factor $|F(\psi)/(m_\psi m_{\bar{\psi}})|^2$ compared with the sfermion masses. The lighter gauginos are a typical feature of this model.

Since the global SUSY is spontaneously broken, there is a Nambu-Goldstone fermion (goldstino) for the SUSY breaking. In the framework of local SUSY (supergravity), the goldstino becomes the longitudinal component of the gravitino. Then, the gravitino has a mass given by

$$m_{3/2} = \frac{F_{Z_1} + F_{Z_2} + F_{Z_3}}{\sqrt{3} M_*},$$

$$= \left( \frac{\sqrt{F_{Z_1} + F_{Z_2} + F_{Z_3}}}{2 \times 10^6 \text{ GeV}} \right)^2 \text{ keV}.$$  (38)
Figure 2: Diagrams contributing to the squark, slepton and Higgs masses
Here, we have imposed the vanishing cosmological constant and $M_* = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. If the SUSY breaking scale is smaller than $2 \times 10^6$ GeV, the gravitino mass is smaller than 1 keV and hence the cosmological requirement is satisfied. If all Yukawa couplings $k_i, \lambda_i, f_i, f_\psi, f_\bar{\psi}$ are $O(1)$ the squark mass $m_{\tilde{q}}$, for example, is roughly given by

$$m_{\tilde{q}} \sim \frac{\alpha_3}{4\pi} \sqrt{F_Z}.$$  \hspace{1cm} (39)

We expect $\sqrt{F_Z} \sim O(10^5)$ GeV to obtain $m_{\tilde{q}} \sim O(10^3)$ GeV. Thus $m_{3/2} < 1$ keV is a natural result in the present model.

The gaugino and sfermion masses in Eqs. (32, 37) are given only at the messenger scale. Therefore, we must reevaluate them at the electroweak scale by using the renormalization group equations (RGEs). Here we present numerical results of the mass spectrum of the gauginos and sfermions including the running effects to the electroweak scale from the messenger scale. To see the dependence of $|F(\psi)/m_\psi m_\bar{\psi}|$, we assume the following mass parameter relation in the messenger sector Eq.(22) for simplicity:

$$m(\psi) = m_\psi = m_{\bar{\psi}} \equiv \Lambda^{(\psi)}$$ \hspace{1cm} (40)

for $\psi = d, l$. The mass spectrum of the superparticles in the MSSM sector is shown as a function of $|F(\psi)/m_\psi m_\bar{\psi}|$ in Fig.3. Here we have set the scales $\Lambda^{(\psi)}$ for $\psi = d, l$ as

$$\Lambda^{(d)} = \Lambda^{(l)} = 5.0 \times 10^5 \text{ GeV},$$ \hspace{1cm} (41)

which corresponds to the gravitino mass $m_{3/2} \sim 0.1$ keV for $\lambda_1 = \lambda_2 = \lambda_3 = f_1 = f_3 = f_\psi = f_{\bar{\psi}} = 1$. As one can see from Fig.3, the gaugino masses rapidly decrease as the parameter $|F(\psi)/m_\psi m_\bar{\psi}|$ gets smaller because there is the suppression factor $|F(\psi)/m_\psi m_\bar{\psi}|^2$ in the gaugino masses in Eq.(34). On the other hand, the sfermion masses weakly depend on $|F(\psi)/m_\psi m_\bar{\psi}|$ since the leading term of order $F(\psi)/\sqrt{m_\psi m_\bar{\psi}}$ does not vanish in contrast to the gaugino masses. Moreover, even if the parameter $|F(\psi)/m_\psi m_\bar{\psi}|$ gets close to 1, the gluino remains lighter than the squarks. The reason is that there is a further suppression factor in a function $A^{(\psi)}$ in Eq.(35) of the mixing angles of the messenger mass matrices. Thus, the gauginos are always lighter than sfermions in the present model.\footnote{The mass spectrum that the squarks are relatively heavy compared with the wino is desirable from the viewpoint of the proton decay in the GUT case. The experimental constraint from the proton decay is about $m_{\tilde{q}}^2/m_{\tilde{g}_2} \gtrsim 10^4$ GeV in the relevant region.}
Figure 3: The mass spectrum of the gauginos and sfermions as a function of the parameter $|F^{(\psi)}/m_\psi m_\bar{\psi}|$. Here we have assumed the mass parameters in the messenger sector as $m^{(\psi)} = m_\psi = m_\bar{\psi} = \Lambda^{(\psi)}$ for $\psi = d, l$. We have set the scales $\Lambda^{(d)} = \Lambda^{(l)} = 5.0 \times 10^5$ GeV. Solid lines represent the gluino $\tilde{g}_3$, wino $\tilde{g}_2$, and bino $\tilde{g}_1$ masses, the dashed lines denote the doublet squark $\tilde{q}_L$, right-handed sup $\tilde{u}_R$, right-handed sdown $\tilde{d}_R$, doublet slepton $\tilde{l}_L$, and right-handed selectron $\tilde{e}_R$. The renormalization effects from the messenger scale to the electroweak scale have been taken into account. We also show the experimental lower bounds on the gluino and wino masses (dash-dotted lines).
Since the gaugino masses are much smaller than the sfermion masses, the constraints on this model come from the experimental bounds on the gaugino masses. Thus we also show the experimental lower bounds on the gluino mass \( m_{\tilde{g}} > 173 \text{ GeV} \) \[28\] and the wino mass \( m_{\tilde{g}_2} > 86 \text{ GeV} \) \[29\] in Fig. 3. The parameters \( |F^{(d)}/m_d m_{\tilde{d}}| \) and \( |F^{(l)}/m_l m_{\tilde{l}}| \) are independent from each other in general. Thus, they are constrained by the experimental bounds on the gluino and wino masses independently. From Fig. 3 we see the parameters \( |F^{(d)}/m_d m_{\tilde{d}}| \) and \( |F^{(l)}/m_l m_{\tilde{l}}| \) are constrained as

\[
\begin{align*}
0.71 &< \frac{|F^{(d)}|}{m_d m_{\tilde{d}}} < 1, \\
0.79 &< \frac{|F^{(l)}|}{m_l m_{\tilde{l}}} < 1. 
\end{align*}
\]

Here, the upper bounds come from the vacuum stability condition Eq. (14). As we discussed in the previous section, the masses \( m_\psi \) and \( m_{\bar{\psi}} \) originated from the same dynamics as SUSY breaking. Therefore it is natural that the parameters \( |F^{(\psi)}/m_\psi m_{\bar{\psi}}| \) for \( \psi = d, l \) are close to 1.

One can see that the so-called GUT relation among the gaugino masses does not hold since the parameters \( |F^{(d)}/m_d m_{\tilde{d}}| \) and \( |F^{(l)}/m_l m_{\tilde{l}}| \) are independent from each other. For example, when \( |F^{(d)}/m_d m_{\tilde{d}}| = 0.75 \) and \( |F^{(l)}/m_l m_{\tilde{l}}| = 0.95 \), the gluino and wino have almost the same mass: \( m_{\tilde{g}_3} \simeq m_{\tilde{g}_2} \simeq 210 \text{ GeV} \) as seen in Fig. 3. It is remarkable that even in this case the gaugino mass relation Eq. (36) holds and hence the bino mass is determined by the other gaugino masses. Since the sfermion masses weakly change with \( |F^{(\psi)}/m_\psi m_{\bar{\psi}}| \), they are almost the same as seen in Fig. 3.

In this analysis, we can not determine the \( \mu \)-parameter unless both the messenger quark and lepton mass parameters are fixed. However, since the scalar masses are much larger than the gaugino masses, the \( \mu \)-parameter tends to be large (\( \sim 1 \text{ TeV} \)). Therefore the lightest neutralino and chargino are almost gauginos. In particular, the lightest neutralino is mostly bino because the bino mass is smaller than the wino mass in the relevant

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9 The gluino mass is constrained as \( m_{\tilde{g}_3} > 173 \text{ GeV} \) for large squark masses by the Tevatron experiment \[28\]. The lightest chargino mass is constrained as \( m_{\tilde{\chi}_1^+} > 86 \text{ GeV} \) for large sneutrino mass and large \( \mu \)-parameter by the LEP experiment at \( \sqrt{s} = 172 \text{ GeV} \) \[29\]. In our model, the squark and slepton masses are much larger than the gaugino masses. The \( \mu \)-parameter also tends to be large as we will discuss in section 3.2. Therefore, we use the gluino mass bound \( m_{\tilde{g}_3} > 173 \text{ GeV} \) and the wino mass bound \( m_{\tilde{g}_2} > 86 \text{ GeV} \) here.
parameter space.

So far we have assumed the mass parameter relation Eq.(40). Next we discuss $m^{(\psi)}$ dependence. We assume that $m_\psi = m_\bar{\psi} = \Lambda^{(\psi)}$ for simplicity, and fix the parameter $|F^{(\psi)}/m_\psi m_\bar{\psi}| = 0.9$ and $\Lambda^{(\psi)} = 5 \times 10^5$ GeV for $\psi = d, l$. Then the parameter $m^{(\psi)}$ dependence on the mass spectrum of the superparticles in the MSSM sector is shown in Fig.4. From the experimental bounds on the gluino and wino masses, the parameters $m^{(\psi)}/\sqrt{m_\psi m_\bar{\psi}}$ for $\psi = d, l$ are constrained as

\begin{align}
0.21 & < \frac{m^{(d)}}{\sqrt{m_d m_d}} < 4.8, \\
0.36 & < \frac{m^{(l)}}{\sqrt{m_l m_l}} < 3.0.
\end{align}

(43)

Since the parameters $m^{(\psi)}$ come from the vacuum expectation values of the singlet fields $k_{\psi_i} Z_i$ (See Eq.(23)), the above constraints require that the scales $k_{\psi_i} \langle Z_i \rangle$ are almost at the same order as the strong SU(2) dynamical scale $\Lambda$.

The constraints in Eq.(42, 43) will be weaken if the scales $\Lambda^{(\psi)}$ becomes larger. When we fix the parameter $|F^{(d)}/m_d m_d| (|F^{(l)}/m_l m_l|)$ and increase the scale $\Lambda^{(d)} (\Lambda^{(l)})$, then the gluino mass (the wino mass) becomes larger proportionally to the scale $\Lambda^{(d)} (\Lambda^{(l)})$ and the function $G^{(d)2} (G^{(l)2})$ in the sfermion mass squared also gets larger proportionally to $\Lambda^{(d)2} (\Lambda^{(l)2})$. However, since the gravitino mass increases proportionally to $\Lambda^{(\psi)2}$, too large value $\Lambda^{(\psi)}$ conflicts with the cosmological bound on the gravitino mass $m_{3/2} < 1$ keV. If the Yukawa couplings are set as $\lambda_1 = \lambda_2 = \lambda_3 = f_1 = f_3 = f_\psi = f_{\bar{\psi}} = 1$ and we take the mass parameter relation Eq.(40), the cosmological bound $m_{3/2} < 1$ keV corresponds to $\Lambda^{(d)} = \Lambda^{(l)} \leq 2 \times 10^6$ GeV. When we set $\Lambda^{(d)} = \Lambda^{(l)} = 2 \times 10^6$ GeV, the constraints on the parameters $|F^{(d)}/m_d m_d|$ and $|F^{(l)}/m_l m_l|$ get weaker than those in Eq.(42) as

\begin{align}
0.47 & < \left| \frac{F^{(d)}}{m_d m_d} \right| < 1, \\
0.54 & < \left| \frac{F^{(l)}}{m_l m_l} \right| < 1.
\end{align}

(44)

The constraints on the parameters $m^{(d)}/\sqrt{m_d m_d}$ and $m^{(l)}/\sqrt{m_l m_l}$ are also much weaken,

\begin{align}
0.05 & < \frac{m^{(d)}}{\sqrt{m_d m_d}} < 19, \\
0.06 & < \frac{m^{(l)}}{\sqrt{m_l m_l}} < 19.
\end{align}
Figure 4: The mass spectrum of the gauginos and sfermions in the MSSM sector as a function of the parameter $m^{(\psi)} / \sqrt{m_{\psi} m_{\bar{\psi}}}$ for $\psi = d, l$. Here we have assumed the mass parameters in the messenger sector as $m_{\psi} = m_{\bar{\psi}} \equiv \Lambda^{(\psi)}$ for $\psi = d, l$. We have set the scales $\Lambda^{(d)} = \Lambda^{(l)} = 5.0 \times 10^5$ GeV. Solid lines represent the gluino $\tilde{g}_3$, wino $\tilde{g}_2$, and bino $\tilde{g}_1$ masses, the dashed lines denote the doublet squark $\tilde{q}_L$, right-handed sup $\tilde{u}_R$, right-handed sdown $\tilde{d}_R$, doublet slepton $\tilde{l}_L$, and right-handed selectron $\tilde{e}_R$. The renormalization effects from the messenger scale to the electroweak scale have been taken into account. We also show the experimental lower bounds on the gluino and wino masses.
\[ 0.08 < \frac{m_{(\ell)}}{\sqrt{m_{\ell} \tilde{m}_{\ell}}} < 12, \quad (45) \]

where we have taken the mass parameters \( m_\psi = m_{\tilde{\psi}} = \Lambda^{(\psi)} = 2 \times 10^6 \text{ GeV} \) and \( |F^{(\psi)}/m_\psi m_{\tilde{\psi}}| = 0.9 \). In the case of the maximal scale \( \Lambda^{(\psi)} \simeq 2 \times 10^6 \text{ GeV} \), however, the squarks become much heavier than the electroweak scale \( m_\tilde{q} \simeq 10 \text{ TeV} \). Therefore, we need a fine tuning of \( \mu \)-parameter in order to break the electroweak symmetry correctly as we will discuss in the next subsection.

### 3.2 Radiative electroweak symmetry breaking

In the framework of the low-energy gauge-mediated SUSY breaking scenario, the radiative electroweak symmetry breaking is also realized as discussed in Refs. [17, 30]. In the model we consider, the radiative electroweak symmetry breaking occurs as in the ordinary gauge mediation models.

The soft SUSY breaking masses for the squarks, sleptons, and Higgses are generated at the messenger scale as discussed in the previous subsection. When we include the running effects of the RGEs, the soft SUSY breaking masses receive significant corrections from the Yukawa interactions as well as the gauge interactions. The soft SUSY breaking masses, especially for \( H_2 \) which is the MSSM Higgs doublet and couples to the top quark, receive significant corrections from the large top Yukawa interaction. The approximate solution to the RGE of the soft SUSY breaking mass squared of \( H_2, m_{H_2}^2 \), is given by an iteration as follows:

\[
m_{H_2}^2 \simeq m_{H_1}^2 - \frac{f_t^2}{16\pi^2} 12m_{\tilde{t}}^2 \log \left( \frac{\Lambda_{\text{mess}}}{m_{\tilde{t}}} \right), \quad (46)
\]

where \( m_{H_1} \) and \( m_{\tilde{t}} \) are the soft SUSY breaking masses for another Higgs doublet and the stop, respectively. \( \Lambda_{\text{mess}} \) denotes the messenger scale, and \( f_t \) represents the top Yukawa coupling constant. In the present model, the stop is heavier than the SU(2) doublet Higgses because \( \alpha_3 > \alpha_2 \). Then it drives \( m_{H_2}^2 \) to a negative value even if the running distance between the messenger scale \( \Lambda_{\text{mess}} \sim 10^5 \text{ GeV} \) and the electroweak scale \( m_{\text{weak}} \sim 10^2 \text{ GeV} \) is not so long. Therefore, the electroweak symmetry is broken radiatively.
Requiring that the tree level potential has an extremum at vacuum expectation values for the two Higgs doublets as $\langle H_1 \rangle = v \cos \beta / \sqrt{2}$ and $\langle H_2 \rangle = v \sin \beta / \sqrt{2}$, one finds

$$\frac{m_Z^2}{2} = \frac{m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \mu^2, \quad (47)$$

$$\sin 2\beta = -\frac{B\mu}{m_{H_1}^2 + m_{H_2}^2 + 2\mu^2}. \quad (48)$$

When we fix $\tan \beta$, the $\mu$-parameter is determined from Eq.(47) to reproduce the correct value for the $Z$ boson mass $m_Z$. Once the parameter $\mu$ is fixed, the $B$-parameter (soft SUSY breaking mass for the Higgs doublet defined as $\mathcal{L} = B\mu H_1 H_2$) is also determined from Eq.(48).

Here we present numerical results of the radiative electroweak symmetry breaking. In our numerical analysis, we fix $\tan \beta$, and we use the one-loop effective potential \[31\] to determine the $\mu$- and $B$-parameters. We include all of the third generation Yukawa couplings. Here, we assume that some unknown dynamics generates the $\mu$- and $B$- terms \[32\]. As for $A$-parameters, which are trilinear couplings of scalars, we assume that they are very small at the messenger scale $\Lambda : A_f(\Lambda) \approx 0$, because they are generated only through higher loop diagrams in all known gauge mediation models. Under this initial condition at the messenger scale, we solve the RGEs for the $A$-parameters and calculate them at the electroweak scale to evaluate the one-loop effective potential.

We show the $\mu$-parameters which satisfy the condition of the radiative electroweak symmetry breaking as a function of the parameter $|F(\psi)/m_\psi m_\bar{\psi}|$ in Fig.5. Here we have assumed the mass parameter relation Eq.(40) and $\Lambda^{(d)} = \Lambda^{(l)}$ in the messenger sector and we have taken $\tan \beta = 3$ (Fig.5(a)), 30 (Fig.5(b)). Solid line, dashed line, and longdashed line correspond to the wino mass $m_{\tilde{g}_w} = 86$ GeV (the experimental lower bound), 150 GeV, and 200 GeV, respectively. As one sees from Fig.5, the $\mu$-parameter gets larger as the parameter $|F(\psi)/m_\psi m_\bar{\psi}|$ becomes smaller when we fix the gaugino masses. The reason is as follows: if the parameter $|F(\psi)/m_\psi m_\bar{\psi}|$ is too small, the gauginos are much lighter than their experimental mass bounds since the gaugino masses have the suppression factor $|F(\psi)/m_\psi m_\bar{\psi}|^2$. To exceed the experimental bound on the gaugino masses, the messenger scale is required to be much larger. Then, the soft SUSY breaking masses for the sfermions get larger since they do not have the suppression like gauginos.
Figure 5: The $\mu$-parameters which satisfy the condition of the radiative electroweak symmetry breaking as a function of the parameter $F^{(\psi)}/m_{\psi}m_{\bar{\psi}}$. Here we take (a) $\tan \beta = 3$ and (b) 30. Solid line, dashed line, and long-dashed line correspond to the wino mass $m_{\tilde{g}_2} = 86$ GeV (the experimental lower bound), 150 GeV, and 200 GeV, respectively.
Since the squarks, especially stop, become much heavier, the mass squared of $H_2$ becomes too negative as one can see in Eq. (46). Then one need a large $\mu$-parameter to reproduce the correct value of the $Z$ boson mass. (See Eq. (47).) For example, in the case with $\tan \beta = 3$ and $\left| F(\psi) / m_\psi m_{\bar{\psi}} \right| = 0.6$, we need a large $\mu$-parameter as $|\mu| \gtrsim 3100$ GeV. In this way, we may need a fine tuning of $\mu$-parameter when the parameter $\left| F(\psi) / m_\psi m_{\bar{\psi}} \right|$ is small (i.e. the SUSY breaking scale is large). Therefore, from such a naturalness point of view the light gravitino $m_{3/2} \sim 0.01 - 1$ keV is implied. On the other hand, as $\tan \beta$ becomes larger, the $\mu$-parameter gets smaller because the mass squared of $H_2$ becomes less negative due to the relatively smaller top Yukawa coupling.

We finally comment on the $\mu$-term generation \[7, 33, 34\]. In the present model, the $\mu$-term may be generated in the same way as the generation of the messenger mass parameters: if the superfield $X$ couples to $H_1H_2$, the SUSY invariant mass $\mu$ for Higgses $H_1$ and $H_2$ is generated. Since $\langle X \rangle \simeq 10^{5-6}$ GeV, we need a small coupling constant $\lambda_h \simeq 10^{-3}$, where $\lambda_h$ is defined by $W = \lambda_h X H_1 H_2$, to have the desired value $\mu \simeq (10^2 - 10^3)$ GeV. The small $\lambda_h$ is natural in the sense of \'t Hooft. We note that no large $B$-term ($B\mu H_1 H_2$) is induced since the $F$-component of $X$ is very small.\[10\] Hence the scale $\mu$ may originate from the same dynamics as the SUSY breaking. In this case, the large $\tan \beta \sim 50$ is required because of the small $B$-parameter to break electroweak symmetry radiatively.

4 The Grand Unified Model

4.1 The GUT model with the Yukawa unification

The messenger quarks and leptons $(d, l)$, $(d', l')$, $(\bar{d}, \bar{l})$, and $(\bar{d}', \bar{l}')$ are embedded in $\mathbf{5}$ and $\mathbf{5}^*$ representations of the SU(5) group. Therefore, if we extend our model to the GUT model (without non-renormalizable interactions), certain Yukawa coupling constants in

\[10\] In this case, the SUSY CP problem is also solved. When we consider the GUT models, the phases of the gaugino masses can be eliminated by a common rotation of the gauginos since $k_{d2}/k_{d3} \simeq k_{l2}/k_{l3}$ holds even at the low-energy scales. Then the rotation of the gauginos gives rise to a phase in the Yukawa-type gauge couplings of the gauginos. Such a phase can be eliminated by a rotation of the sfermions and Higgses since there are no $A$-terms and no $B$-term at the tree level.

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Eq.(14) are unified as
\[
k_{di} = k_{li}, \quad (i = 1 - 3), \quad (49)
\]
\[
f_d = f_l, \quad (50)
\]
\[
f_{\bar{d}} = f_{\bar{l}}, \quad (51)
\]
at the GUT scale \(M_{\text{GUT}} \approx 2 \times 10^{16} \) GeV. Under the Yukawa unification Eqs.(49−51), we can obtain the relations between the Yukawa couplings \(k_{di}\) and \(k_{li}\), \(f_d\) and \(f_l\), and \(f_{\bar{d}}\) at the messenger scale using the RGEs for these Yukawa coupling constants. We list the RGEs for the coupling constants in this model in Appendix B. In general, the gauge interactions increase the Yukawa coupling constants at the lower scale. Since the messenger quarks have the SU(3)\(_C\) gauge interaction but the messenger leptons do not, the Yukawa coupling constants \(k_{di}, f_d\), and \(f_{\bar{d}}\) related to the messenger quarks tend to be larger than the Yukawa couplings \(k_{li}, f_l\), and \(f_{\bar{l}}\) related to the messenger leptons because \(g_3(\mu) \geq g_2(\mu)\) at \(\mu \leq M_{\text{GUT}}\). We consider the RGEs for the ratios of the Yukawa couplings \(k_{di}/k_{li}\), \(f_d/f_l\), and \(f_{\bar{d}}/f_{\bar{l}}\):
\[
16\pi^2 \mu \frac{d}{d\mu} \left( \frac{k_{di}}{k_{li}} \right) = \frac{k_{di}}{k_{li}} \left\{ -\frac{16}{3} g_3^2 + 3 g_2^2 + \frac{1}{3} g_1^2 \\
+ 2 \sum_{j=1}^{3} (k_{dj}^2 - k_{lj}^2) + f_d^2 - f_l^2 + \sum_{j=1}^{3} (\lambda_i \lambda_j + 3 k_{di} k_{dj} + 2 k_{li} k_{lj} + 2 f_i f_j) \right\}, \quad (52)
\]
\[
16\pi^2 \mu \frac{d}{d\mu} \left( \frac{f_d}{f_l} \right) = \frac{f_d}{f_l} \left\{ -\frac{16}{3} g_3^2 + 3 g_2^2 + \frac{1}{3} g_1^2 + \sum_{i=1}^{3} (k_{di}^2 - k_{li}^2) + 2(f_d^2 - f_l^2) \right\}, \quad (53)
\]
\[
16\pi^2 \mu \frac{d}{d\mu} \left( \frac{f_{\bar{d}}}{f_{\bar{l}}} \right) = \frac{f_{\bar{d}}}{f_{\bar{l}}} \left\{ -\frac{16}{3} g_3^2 + 3 g_2^2 + \frac{1}{3} g_1^2 + \sum_{i=1}^{3} (k_{di}^2 - k_{li}^2) + 2(f_{\bar{d}}^2 - f_{\bar{l}}^2) \right\}. \quad (54)
\]
One can see that under the condition Eqs.(49−51), the ratios \(k_{di}/k_{li}\), \(f_d/f_l\), and \(f_{\bar{d}}/f_{\bar{l}}\) become larger at the lower scale because the SU(3)\(_C\) gauge interaction dominates over in Eqs.(52−54) as long as the perturbative description of the Yukawa couplings is valid. When we impose the unification condition Eqs.(49−51), the right-hand side in Eq.(52) is independent of the index \(i\). Therefore the ratio \(k_{di}/k_{li}\) does not depend on the index \(i\) at any scales. Then we can always set \(k_{d1} = k_{l1} = k_{d2} = k_{l2} = 0\) and \(f_2 = 0\) at any scales.
by an unitary transformation of the singlet fields $Z_i$. Thus we work in this basis below. Then the ratios of the Yukawa couplings $k_{d3}/k_{l3}$, $f_d/f_l$, and $f_{d}/f_{l}$ equal to those of the mass parameters for the messenger sector $m^{(d)}/m^{(l)}$, $m_d/m_l$, and $m_{d}/m_{l}$, respectively.

To numerically analyze the relations between mass parameters in the messenger sector, we fix some parameters as $\lambda_Y = 0.3$, $(\lambda_1, \lambda_2, \lambda_3) = (0.2, 0.1, 0.16)$, $(f_1, f_2, f_3) = (0.5, 0.0, 0.6)$, $(k_{d1}, k_{d2}) = (k_{l1}, k_{l2}) = (0, 0)$ at the GUT scale, $\Lambda = 2 \times 10^6$ GeV which is the dynamical scale of the strong SU(2) gauge interaction, and $\langle Z_3 \rangle = 1 \times 10^6$ GeV which is the vacuum expectation value of the singlet field $Z_3$. The strong SU(2) gauge coupling $g$ is taken as $2\pi$ at the messenger scale (i.e. $\alpha = g^2/4\pi = \pi$). We solve the RGEs numerically varying $k_{d3} = k_{l3} \equiv k$, $f_d = f_l = f_\psi$, and $f_{d} = f_{l} = f_\bar{\psi}$ at the GUT scale. We find the ratios of the mass parameters given by

$$\frac{m^{(d)}}{m^{(l)}} \simeq 1.4 \ (1.4),$$

$$\frac{m_d}{m_l} = \frac{m_{d}}{m_{l}} \simeq 1.4 \ (1.4),$$

$$0.5 \ (0.4) < \frac{m^{(l)}}{\sqrt{m_l m_l}} \simeq \frac{m^{(d)}}{\sqrt{m_d m_d}} < 1.2 \ (0.9),$$

for $0.42 < k < 1.0$ and $f_\psi = f_{\bar{\psi}} = 0.20$ (for $0.32 < k < 0.54$ and $f_\psi = f_{\bar{\psi}} = 0.18$). Here we have taken $f_\psi = f_{\bar{\psi}}$ for simplicity. These parameters correspond to the parameter region $F^{(l)}/m_l m_l > 0.7$. The ratio of $m^{(d)}/m^{(l)}$ is almost the same as the ratio of $m_d/m_l = m_{d}/m_{l}$ because the running effects for the ratios of the Yukawa couplings $k_d/k_l$, $f_d/f_l$, and $f_d/f_l$ dominantly come from the SU(3)$_C$ gauge interaction as one can see from Eqs.(52–54). The parameter $F^{(d)}/m_d m_{\bar{d}}$ is also related to $F^{(l)}/m_l m_l$ as follow:

$$\frac{F^{(l)}}{m_l m_l} \simeq 1.4 \ \frac{F^{(d)}}{m_d m_{\bar{d}}},$$

since

$$\left( \frac{F^{(l)}}{m_l m_l} \right) / \left( \frac{F^{(d)}}{m_d m_{\bar{d}}} \right) = \frac{k_{l3}}{k_{d3}} \frac{m_{d}}{m_l} \frac{m_{\bar{d}}}{m_{\bar{d}}} = \frac{m^{(l)}}{m^{(d)}} \frac{m_{d}}{m_{l}} \frac{m_{d}}{m_{d}}.$$

The Yukawa couplings $k_i, f_\psi, f_{\bar{\psi}}$ do not significantly affect the runnings of $\lambda_i, f_i$ as long as the perturbative description of the Yukawa couplings is valid. Thus the vacuum expectation values $\langle X \rangle$ and $\langle F_{Z_i} \rangle$ hardly depend on $k_i, f_\psi, f_{\bar{\psi}}$. 

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Because of the Yukawa unification, \( F(l)/m_l m_l \) becomes larger than \( F(d)/m_d m_d \) at the messenger scale. This yields an interesting consequence on the mass spectrum of the gauginos as we will see below.

We are now at the position to show the mass spectrum of the gauginos and sfermions in the MSSM sector. In Fig.6, the mass spectrum is shown as a function of the parameter \( F(l)/m_l m_l \). Here we have taken the same parameter set as above mentioned one and we have set \( f_\psi = f_\bar{\psi} = 0.2 \) and varied \( k \) with \( 0.4 < k < 1 \). Note that this parameter set corresponds to the gravitino mass \( m_{3/2} \approx 0.12 \text{ keV} \). As we have seen in the previous section, the gaugino masses strongly depend on the parameters \( F(\psi)/m_\psi m_\bar{\psi} \) for \( \psi = d, l \) since they have an extra suppression factor \( (F(\psi)/m_\psi m_\bar{\psi})^2 \). When we impose the Yukawa unification Eqs.(49–51), the parameter \( F(d)/m_d m_d \) is smaller than \( F(l)/m_l m_l \) as shown in Eq.(58). Therefore the gluino mass receives a larger suppression than the wino mass, and hence the gluino tends to be relatively light. The experimental lower bound on the gluino mass constrains the parameter \( F(l)/m_l m_l \) to be \( F(l)/m_l m_l > 0.89 \). This leads to a constraint on the Yukawa coupling \( k \) at the GUT scale as \( k > 0.65 \). Therefore, the GUT relation among the gaugino masses does not hold even though we consider the GUT model. We note that the gluino can be lighter than the wino.

We also show the \( \mu \)-parameter in Fig.6. As we discussed in section 3.2, the \( \mu \)-parameter becomes much larger than the gaugino masses.

### 4.2 The GUT model without the Yukawa unification

So far we have considered the GUT model with the Yukawa unification. In this section, we consider the GUT model with non-renormalizable interactions. The superpotential which contributes to the Yukawa couplings \( k_\psi, f_\psi, \) and \( f_\bar{\psi} (\psi = d, l) \) at low energy scales is given by

\[
W = \sum_{i=1}^{3} Z_i \left( k'_{\psi_i} \psi_i \bar{\psi} + k''_{\psi_i} \frac{\Sigma}{M_e} \psi_i \bar{\psi} + \cdots \right) + X \left( f'_{\psi} \psi_i \bar{\psi} + f''_{\psi} \frac{\Sigma}{M_e} \psi_i \bar{\psi} + f'''_{\psi} \frac{\Sigma}{M_e} \psi_i \bar{\psi} + \cdots \right),
\]

\[\text{(60)}\]

12When we regard the masses for the gauginos and sfermions as a function of \( F(l)/m_l m_l \), we obtain almost the same result as in Fig.6 even if we take a different value of the Yukawa coupling \( f_\psi = f_\bar{\psi} \). The reason is that the ratio between \( F(l)/m_l m_l \) and \( F(d)/m_d m_d \) is almost independent of \( f_\psi = f_\bar{\psi} \).
Figure 6: The mass spectrum of the gauginos and sfermions in the MSSM sector as a function of $F^{(l)}/m_{1/2}$. Here we have taken $\lambda_Y = 0.3$, $(\lambda_1, \lambda_2, \lambda_3) = (0.2, 0.1, 0.16)$, $(f_1, f_2, f_3) = (0.5, 0, 0.6)$, $(k_{d1}, k_{d2}) = (k_{l1}, k_{l2}) = (0, 0)$ at the GUT scale. The dynamical scale $\Lambda$ of the strong SU(2) gauge interaction and the vacuum expectation value $\langle Z_3 \rangle$ have been taken as $\Lambda = 2.0 \times 10^6$ GeV and $\langle Z_3 \rangle = 1.0 \times 10^6$ GeV, respectively. We have set $f_d = f_l = f_{\tilde{d}} = f_{\tilde{d}} = 0.2$ at the GUT scale for simplicity and varied $k_{l3} = k_{d3} \equiv k$ with $0.4 < k < 1.0$. Note that this parameter set corresponds to the gravitino mass $m_{3/2} \simeq 0.12$ keV. Solid lines represent the gaugino masses: the bino $\tilde{g}_1$, wino $\tilde{g}_2$, and gluino $\tilde{g}_3$ masses, and dashed lines represent the sfermion masses: the left-handed squark $\tilde{g}_{L}$, right-handed squark $\tilde{u}_R$, right-handed sbottom $\tilde{d}_R$, doublet slepton $\tilde{l}_L$, and right-handed slepton $\tilde{e}_R$ masses. The $\mu$-parameter is also shown (dotted line). The renormalization effects from the messenger scale to the electroweak scale have been taken into account. We also show the experimental lower bound on the gluino mass (dash-dotted line).
where the fields $\psi$ and $\psi'$ ($\bar{\psi}$ and $\bar{\psi}'$) are $5$ ($5^*$) dimensional representation fields of SU(5) group, which contain the messenger quark and lepton multiplets ($d, l$) and ($d', l'$) ($\bar{d}, \bar{l}$) and ($\bar{d}', \bar{l}'$), respectively. The field $\Sigma$ is a $24$ dimensional representation which breaks SU(5) down to SU(3)$_C \times$SU(2)$_L \times$U(1)$_Y$ with expectation value $\langle \Sigma \rangle = V \text{diag}(2, 2, 2, -3, -3)$. We should notice that the non-renormalizable interactions in Eq.(60) are not forbidden by any symmetries provided that the $\Sigma$ is a trivial representation of U(1)$_R \times$U(1)$_\chi$. Then, the Yukawa coupling constants $k_\psi$, $f_\psi$, and $f_{\bar{\psi}}$ receive corrections of order of $O(\langle \Sigma \rangle/M_*)$ as

\begin{align*}
k_{di} &= k'_{\psi d} + 2k''_{\psi d} \frac{V}{M_*}, \\
k_{li} &= k'_{\psi l} - 3k''_{\psi l} \frac{V}{M_*}, \\
f_d &= f'_{\psi d} + 2f''_{\psi d} \frac{V}{M_*}, \\
f_l &= f'_{\psi l} - 3f''_{\psi l} \frac{V}{M_*},
\end{align*}

at the GUT scale, and hence the Yukawa coupling unification of Eqs.(49–51) is broken in general.

In such a case, the relations between the messenger quark and lepton mass parameters depend on the corrections from the non-renormalizable terms, and the relation in Eq.(58) does not hold. Thus the mass relation among the gauginos in the case with the Yukawa unification may easily change. Therefore, the existence of non-renormalizable terms weakens our prediction on the mass spectrum of the superparticles. However, it is important to study how the mass spectrum in the previous subsection changes.

To demonstrate changes of the mass spectrum, we consider a simple example. Here we take the following initial condition of the Yukawa couplings at the GUT scale,

\begin{align*}
k_{d2} &= k_2 - \delta, \\
k_{l2} &= k_2, \\
k_{d3} &= k_3 - \delta, \\
k_{l3} &= k_3, \\
f_d &= f - \delta, \\
f_l &= f, \\
f_{\bar{d}} &= \bar{f} - \delta, \\
f_{\bar{l}} &= \bar{f}.
\end{align*}

The parameter $\delta$ represents the difference between the Yukawa couplings of messenger quarks and leptons due to corrections from non-renormalizable interactions. The case
$\delta = 0$ corresponds to the GUT model with the Yukawa unification. We show the mass spectrum of the gauginos and sfermions as a function of the parameter $\delta$ in Fig.7. Here the other Yukawa couplings have been set as $\lambda_Y = 0.3$, $(\lambda_1, \lambda_2, \lambda_3) = (0.2, 0.1, 0.16)$, $(f_1, f_2, f_3) = (0.5, 0, 0.6)$, $k_{d1} = k_{l1} = 0$ at the GUT scale and the strong SU(2) gauge coupling $g$ has been taken as $g = 2\pi$ at the messenger scale. The dynamical scale $\Lambda$ of the strong SU(2) gauge interaction and the vacuum expectation values $\langle Z_2 \rangle$ and $\langle Z_3 \rangle$ have been taken as $\Lambda = 1.0 \times 10^6$ GeV and $\langle Z_2 \rangle = \langle Z_3 \rangle = 5.0 \times 10^5$ GeV, respectively. We have taken the parameters $k_2 = 0$, $k_3 = 0.8$, and $f = \bar{f} = 0.2$ in Eq.(65). From Fig.6, we see that the case with the Yukawa unification (i.e. $\delta = 0$) is excluded by the experimental lower bound on the gluino mass. However, the small corrections ($\delta \sim 0.013$) from the non-renormalizable terms may significantly change the mass spectrum. This is because the parameter $F^{(d)}/m_d m_{\tilde{d}}$ gets larger than that in the case with the Yukawa unification while $F^{(l)}/m_l m_{\tilde{l}}$ being kept unchanged. Since the suppression factor in the gluino mass gets closer to 1, the gluino becomes heavier to escape the experimental lower bound.

It is relatively difficult to predict the mass spectrum in the presence of the non-renormalizable terms. However, we should stress that this model still has many distinguishable features as discussed in the previous sections: the GUT relation among the gaugino masses does not hold, and the gauginos tend to be lighter than the sfermions since the gaugino masses have the suppression factor.

5 Other interesting features in the present model

From the naturalness point of view as discussed in Section 3.2, much large sfermion masses compared with the weak scale are not preferable. Then, it is natural to have a relatively lower SUSY breaking scale, that is, the lighter gravitino which satisfies the cosmological requirement $m_{3/2} < 1$ keV. Such a light gravitino may bring us other interesting consequences in the collider experiments. When the gravitino is so light that the cosmological requirement is satisfied, the gravitino becomes the LSP. Thus the NLSP, primarily the bino, decays into the gravitino emitting a photon. The goldstino component of the gravitino has the following interaction with the matters [35],

$$\mathcal{L} = -\frac{1}{F} j^{\alpha\mu} \partial_{\mu} \tilde{G}_{\alpha},$$
Figure 7: The mass spectrum of the gauginos and sfermions as a function of the parameter \( \delta \). Here we have taken the Yukawa couplings at the GUT scale as follows: \( \lambda_Y = 0.3 \), \( (\lambda_1, \lambda_2, \lambda_3) = (0.2, 0.1, 0.16) \), \( (f_1, f_2, f_3) = (0.5, 0, 0.6) \), \( k_{d_1} = k_{l_1} = 0 \). We have set \( k_2 = 0 \), \( k_3 = 0.8 \), and \( f = \bar{f} = 0.2 \) in Eq.(65). Solid lines represent the gaugino masses: the bino \( \tilde{g}_1 \), wino \( \tilde{g}_2 \), and gluino \( \tilde{g}_3 \) masses, and dashed lines represent the sfermion masses: the left-handed squark \( \tilde{q}_L \), right-handed sup \( \tilde{u}_R \), right-handed sdown \( \tilde{d}_R \), doublet slepton \( \tilde{l}_L \), and right-handed selectron \( \tilde{e}_R \) masses. The \( \mu \)-parameter is also shown (dotted line). The renormalization effects from the messenger scale to the electroweak scale have been taken into account. We also show the experimental lower bound on the gluino mass (dash-dotted line).
\[
\text{cos} \theta \frac{m_{\tilde{g}_1}}{2\sqrt{2}F} \tilde{B}\bar{\sigma}^\nu \sigma^\mu \tilde{G} F_{\mu\nu} + \cdots.
\] (66)

An important point is that the couplings of the goldstino are suppressed by only the SUSY breaking scale, not the Planck scale. Therefore, the couplings of the goldstino get larger as the SUSY breaking scale becomes smaller. If the SUSY breaking scale is sufficiently small, the decay of the superparticles into the gravitino may occur within the detector.

From the interaction Eq.(66), we find the decay rate of the bino as

\[
\Gamma(\tilde{g}_1 \rightarrow \tilde{G} \gamma) = \frac{m_{\tilde{g}_1}^5 \cos^2 \theta_W}{16\pi F^2}.
\] (67)

Then the decay length \( L \) of the bino with energy \( E \) in the laboratory frame is given by

\[
L = 1.3 \left( \frac{100 \text{ GeV}}{m_{\tilde{g}_1}} \right)^5 \left( \frac{\sqrt{F}}{10^6 \text{ GeV}} \right)^4 \left( \frac{E^2}{m_{\tilde{g}_1}^2} - 1 \right)^{\frac{1}{2}} \text{ m}.
\] (68)

For example, when \( \sqrt{F} = 5 \times 10^5 \text{ GeV} \) which corresponds to the gravitino mass with \( m_{3/2} = 0.06 \text{ keV} \) and \( m_{\tilde{g}_1} = 100 \text{ GeV} \), the decay length of the bino is 8cm for \( (E^2/m_{\tilde{g}_1}^2 - 1)^{1/2} = 1 \). Since it is possible for the gravitino mass to be much lighter than \( m_{3/2} < 1 \text{ keV} \), we may find signature “\( \gamma \gamma + \) missing energy” in the collider experiments. We should notice that any realistic models with a sizable signature “\( \gamma \gamma + \) missing energy” have not been known except for the present model. The phenomenological investigations have been done in many literatures \[17, 18\] only under the assumption of the existence of very light gravitino.\[13\]

Recently, however, the constraints on the SUSY models with “\( \gamma \gamma + \) missing energy” signal have been reported in Ref. \[37\]. The lower bounds on the lightest neutralino and the lightest chargino masses are obtained as

\[
m_{\tilde{g}_1} \gtrsim 75 \text{ GeV},
\]

\[
m_{\tilde{g}_2} \gtrsim 150 \text{ GeV},
\] (69)

where the GUT relation among the gaugino masses is assumed. To satisfy the mass bounds Eq.(69), the squarks and sleptons should be very heavy \( (m_{\tilde{f}} \sim 10 \text{ TeV}) \) or the

\[\text{The signature “\( \gamma \gamma + \) missing energy” does not necessarily suggest the light gravitino. It has been known that the light axino in the framework of the no-scale supergravity also brings us such a signal in the collider experiments [34]. Therefore, other experiments, axion search and direct superparticle search for example, are also necessary to distinguish the present model from other models.}\]
parameter $F(\psi)/m_\psi m_\bar{\psi}$ should be very close to 1 in the present model. If the squarks and sleptons are so heavy, $m_f \sim 10$ TeV, a fine tuning may be needed to obtain the correct electroweak scale as discussed in section 3.2. On the other hand, in the case where the parameter $F(\psi)/m_\psi m_\bar{\psi}$ is very close to 1, it already means a fine tuning. From the naturalness point of view, we may exclude the parameter regions where the “$\gamma\gamma+$ missing energy” signal is detectable in the present experiments.

The decay length Eq.(68) strongly depends on the bino mass as well as the SUSY breaking scale. For example, when $\sqrt{F} = 5 \times 10^5$ GeV and $m_{\tilde{g}} = 40$ GeV, $L = 8$ m for $(E^2/m_{\tilde{g}}^2 - 1)^{1/2} = 1$ which is larger than the typical detector size. Therefore, the signal “$\gamma\gamma+$ missing energy” can not be observed within the detector when the bino mass is much smaller than 100 GeV. In the case where the bino slowly decays outside the detector, the bounds in Eq.(69) are not applicable because they are derived under the assumption that the bino completely decays into the photon and gravitino within the detector. Since our model suggests the light gauginos as discussed in the previous sections, it is likely the case. Although the “$\gamma\gamma+$ missing energy” signature is not expected in the existing experiments, the pair production of the light gauginos can be observed directly in the future collider experiments. Furthermore, if we require that $m_{3/2} \lesssim 1$ keV or $m_\tilde{f} \lesssim 10$ TeV, the SUSY breaking scale becomes as $\sqrt{F} \lesssim 2 \times 10^6$ GeV. Then, according to Ref. [19], the slow decay of the bino may be detectable in the near future experiments such as LHC, even for the case of long $L_{\tilde{f}}$.

Finally we remark on the unification of the gauge coupling constants. To break the SUSY dynamically, we assume the strong SU(2) gauge interaction with the four fundamental representation fields. It is remarkable that all gauge coupling constants, not only $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge couplings but also strong SU(2) gauge coupling, meet at the scale $\sim 10^{16}$ GeV when we set the strong SU(2) dynamical scale $\Lambda$ as $\Lambda \simeq 10^{5-6}$ GeV in order that $m_{3/2} \sim 0.1 - 1$ keV [16].

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14 In the experiment proposed in Ref. [19], the “$\gamma\gamma+$ missing energy” signature can be detectable at LHC as long as $\sqrt{F} \lesssim 10^7$ GeV. Thus such a signal will be observed in our model. In other models which have the large $\sqrt{F} > 10^7$ GeV, however, this signal can not be observable at LHC even with the same $L$.
6 Conclusion

We have performed a detailed analysis of a direct-transmission model of the dynamical SUSY breaking previously proposed in Ref. [16]. This model possesses many remarkable points: there are no SUSY FCNC, all mass scales are generated from the strong SU(2) dynamics (\(\mu\)-term may also originate from the same dynamics), and the unification of all gauge coupling constants of SU(3)\(_C\), SU(2)\(_L\), U(1)\(_Y\), and the strong SU(2) gauge groups may be realized. Furthermore it is quite natural for the gravitino mass to be smaller than 1 keV as required from the standard cosmology. We notice that this cosmological requirement is not satisfied by any other models which have been proposed.

In the present model, there are many distinguishable low-energy features from in other models: the so-called GUT relation among the gaugino masses does not hold even if we consider the GUT models. Furthermore, the gauginos become lighter than sfermions because the gaugino masses have extra suppression factors. When the suppression of the gaugino masses is large, squarks tend to be heavy in order to satisfy the experimental lower bounds on the gaugino masses. In this case, the radiative electroweak symmetry breaking requires a large \(\mu\)-parameter, and we may need a fine tuning. If we consider the GUT model with the Yukawa unification, the gluino mass tends to be lighter than the wino mass. However, the Yukawa unification is broken in the presence of non-renormalizable interactions. Then the prediction of the mass spectrum of the superparticles is weakened.

Moreover, the light gravitino brings us a fascinating signature, that is, \(\gamma\gamma+\) missing energy” in the collider experiments when the bino is relatively heavy (\(m_{\tilde{g}_1}\gtrsim 100\) GeV). From the naturalness point of view, however, the light gauginos are most likely in our model. Thus the collider signature will not be observed in the present experiments because the bino decays outside the detector. Even in this case, the bino decay may be detectable [19] in the near future colliders because of the relatively lower SUSY breaking scale.

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A Calculation of the gaugino and sfermion masses

In this Appendix, we show the detailed calculation of the gaugino and sfermion masses in the MSSM sector.

When we take the mass eigenbasis in the messenger fermions and sfermions as discussed in section 3.1, gaugino-fermion-scalar interaction is

\[ \mathcal{L}_s = \sqrt{2} \lambda \sum_{\alpha=1}^{3} \sum_{\beta=1}^{4} \left[ (T^{(\psi)}_{\beta \gamma} V^{\dagger}_{\beta \alpha} + T^{(\psi)}_{\gamma \alpha} V^{\dagger}_{\gamma \beta}) \tilde{\psi}^{(a)}_{\beta} \tilde{\psi}^{(a)}_{\gamma} \tilde{g}^{(a)} \right. \\
- \left( V^{(\psi)}_{\alpha \beta} T^{(\psi)}_{\beta \gamma} \right) \tilde{\psi}^{(a)}_{\beta} \tilde{\psi}^{(a)}_{\gamma} \tilde{g}^{(a)} \right] \]

where \( \psi_{\alpha} \) and \( \tilde{\psi}_{\alpha} \) denote the messenger fermions in the mass eigenstates and \( \tilde{\psi}_{\alpha} \) is the messenger sfermions in the mass eigenstates as follows:

\[
\begin{pmatrix}
\psi_1 \\
\psi_2
\end{pmatrix} = V^{(\psi)} \begin{pmatrix}
\psi \\
\psi'
\end{pmatrix}, \quad \quad \quad (71)
\]

\[
\begin{pmatrix}
\tilde{\psi}_1 \\
\tilde{\psi}_2
\end{pmatrix} = U^{(\psi)} \begin{pmatrix}
\tilde{\psi} \\
\tilde{\psi}'
\end{pmatrix}, \quad \quad \quad (72)
\]

\[
\begin{pmatrix}
\tilde{\psi}_1 \\
\tilde{\psi}_2 \\
\tilde{\psi}_3 \\
\tilde{\psi}_4
\end{pmatrix} = T^{(\psi)} \begin{pmatrix}
\tilde{\psi} \\
\tilde{\psi}' \\
\tilde{\psi}^{*} \\
\tilde{\psi}'^{*}
\end{pmatrix}, \quad \quad \quad (73)
\]

and \( \tilde{g}^{(a)} \) is the gaugino and \( T^{(a)}_{ij} \) represents the generator of the gauge group. Then the MSSM gaugino masses coming from Fig.1 becomes

\[
m_{\tilde{g}_3} = \frac{\alpha_3}{2\pi} F^{(d)}, \quad \quad \quad (74)
\]

\[
m_{\tilde{g}_2} = \frac{\alpha_2}{2\pi} F^{(l)}, \quad \quad \quad (75)
\]

\[
m_{\tilde{g}_1} = \frac{\alpha_1}{2\pi} \left\{ \frac{2}{5} F^{(d)} + \frac{3}{5} F^{(l)} \right\}, \quad \quad \quad (76)
\]

where the masses \( m_{\tilde{g}_i} \) \((i = 1, \cdots, 3)\) denote the bino, wino, and gluino masses, respectively, and we have adopted the SU(5) GUT normalization of the U(1)$_{Y}$ gauge coupling \((\alpha_1 \equiv \cdots \equiv \alpha_3)\).
\[ F^{(\psi)} = 32\pi^2 i \sum_{a=1}^{4} \sum_{\beta=1}^{4} (U^{(\psi)\alpha}_{a1} T^{(\psi)\beta}_{3\beta} + U^{(\psi)\alpha}_{a2} T^{(\psi)\beta}_{4\beta})(T^{(\psi)\alpha}_{\beta1} V^{(\psi)\alpha}_{1\alpha} + T^{(\psi)\alpha}_{\beta2} V^{(\psi)\alpha}_{2\alpha}) \]

\[ \times T^{(a)}_{ij} T^{(a)}_{ji} M^{(\psi)}_{\alpha} \int \frac{d^4k}{(2\pi)^4} \frac{1}{(p-k)^2 + m^{(\psi)\beta}_\beta} k^2 + M^{(\psi)\alpha}_\alpha |p=0 = 16\pi^2 i \sum_{a=1}^{4} \sum_{\beta=1}^{4} (U^{(\psi)\alpha}_{a1} T^{(\psi)\beta}_{3\beta} + U^{(\psi)\alpha}_{a2} T^{(\psi)\beta}_{4\beta})(T^{(\psi)\alpha}_{\beta1} V^{(\psi)\alpha}_{1\alpha} + T^{(\psi)\alpha}_{\beta2} V^{(\psi)\alpha}_{2\alpha}) \]

\[ \times M^{(\psi)}_\alpha \{ -\frac{i}{16\pi^2} \int_0^1 dx \ln[x M^{(\psi)\beta}_\beta + (1-x) M^{(\psi)\alpha}_\alpha] \} |_{p=0} \]

\[ = \sum_{a=1}^{4} \sum_{\beta=1}^{4} M^{(\psi)}_\alpha (U^{(\psi)\alpha}_{a1} T^{(\psi)\beta}_{3\beta} + U^{(\psi)\alpha}_{a2} T^{(\psi)\beta}_{4\beta})(T^{(\psi)\alpha}_{\beta1} V^{(\psi)\alpha}_{1\alpha} + T^{(\psi)\alpha}_{\beta2} V^{(\psi)\alpha}_{2\alpha}) \]

\[ \times \frac{m^{(\psi)\beta}_\beta}{m^{(\psi)\alpha}_\alpha - M^{(\psi)\alpha}_\alpha} \ln \frac{m^{(\psi)\beta}_\beta}{M^{(\psi)\alpha}_\alpha}. \] (77)

We next consider the squark and slepton masses. These masses arise from Fig. 2 and are given by

\[ m_j^2 = \frac{1}{2} \left[ C^{(d)}_3 \left( \frac{\alpha_3}{4\pi} \right)^2 G^{(d)2} + C^{(t)}_3 \left( \frac{\alpha_3}{4\pi} \right)^2 G^{(t)2} + \frac{3}{5} Y^2 \left( \frac{\alpha_1}{4\pi} \right)^2 \left( \frac{2}{5} G^{(d)2} + \frac{3}{5} G^{(t)2} \right) \right], \] (78)

where \( C^{(d)}_3 = \frac{4}{3} \) and \( C^{(t)}_3 = \frac{3}{4} \) when \( f \) is in the fundamental representation of SU(3)_C and SU(2)_L, and \( C^{(t)}_t = 0 \) for the gauge singlets, and \( Y \) denotes the U(1)_Y hypercharge \( (Y \equiv Q - T_3) \). Here \( G^{(\psi)2} \) are given by [23]

\[ G^{(\psi)2} = (4\pi)^4 \{ -\sum_{a=1}^{4} < m^{(\psi)}_\alpha | m^{(\psi)}_\alpha | 0 > -4 \sum_{a=1}^{4} < m^{(\psi)}_\alpha | m^{(\psi)}_\alpha | 0, 0 > \]

\[ -4 \sum_{a=1}^{2} < M^{(\psi)}_\alpha | M^{(\psi)}_\alpha | 0 > +8 \sum_{a=1}^{2} < M^{(\psi)}_\alpha | M^{(\psi)}_\alpha | 0, 0 > \]

\[ - \sum_{a=1}^{4} \sum_{\beta=1}^{4} (T^{(\psi)\alpha}_{a1} T^{(\psi)\beta}_{1\beta} + T^{(\psi)\alpha}_{a2} T^{(\psi)\beta}_{2\beta} - T^{(\psi)\alpha}_{a3} T^{(\psi)\beta}_{3\beta} - T^{(\psi)\alpha}_{a4} T^{(\psi)\beta}_{4\beta}) \]

\[ \times (T^{(\psi)\alpha}_{\beta1} T^{(\psi)\alpha}_{1\alpha} + T^{(\psi)\alpha}_{\beta2} T^{(\psi)\alpha}_{2\alpha} - T^{(\psi)\alpha}_{\beta3} T^{(\psi)\alpha}_{3\alpha} - T^{(\psi)\alpha}_{\beta4} T^{(\psi)\alpha}_{4\alpha}) < m^{(\psi)}_\alpha | m^{(\psi)}_\beta | 0 > \]
In this Appendix, we list the RGEs for the model.

Here $\text{Li}_2(x) = -\int_0^x (dt/t) \ln(1-xt)$ is a dilog function.

B Renormalization group equations for the model

In this Appendix, we list the RGEs for the model.
The superpotential at the scale above $\Lambda$ is

$$W = \frac{\lambda}{\sqrt{2}} \epsilon_{\alpha\beta} \{ Q_1^\alpha Q_3^\beta Y_1 + Q_1^\alpha Q_1^\beta Y_2 + Q_3^\alpha Q_3^\beta Y_3 + Q^\alpha_3 Q^\beta_3 Y_4 + \frac{1}{\sqrt{2}} (Q_1^\gamma Q_2^\beta - Q_3^\gamma Q_3^\beta) Y_5 \}
+ \frac{\lambda_1}{\sqrt{2}} \epsilon_{\alpha\beta} (Q_1^\alpha Q_2^\beta + Q_3^\alpha Q_4^\beta) Z_1 + k_{d1} \bar{d} d Z_1 + k_{l1} \bar{l} l Z_1 - f_1 X^2 Z_1
+ \frac{\lambda_2}{\sqrt{2}} \epsilon_{\alpha\beta} (Q_1^\alpha Q_2^\beta + Q_3^\alpha Q_4^\beta) Z_2 + k_{d2} \bar{d} d Z_2 + k_{l2} \bar{l} l Z_2 - f_2 X^2 Z_2
+ \frac{\lambda_3}{\sqrt{2}} \epsilon_{\alpha\beta} (Q_1^\alpha Q_2^\beta + Q_3^\alpha Q_4^\beta) Z_3 + k_{d3} \bar{d} d Z_3 + k_{l3} \bar{l} l Z_3 - f_3 X^2 Z_3
+ f_d \bar{d} d X + f_d \bar{d} d X + f_l \bar{l} l X + f_l \bar{l} l X. \quad (83)$$

Then, one-loop RGEs for the Yukawa coupling constants are given by

$$\mu \frac{d \lambda}{d \mu} = \frac{1}{16 \pi^2} \frac{1}{2} \lambda (\lambda_1^2 + \lambda_2^2 + \lambda_3^2 + 7 \lambda^2 - 6 g^2),$$
$$\mu \frac{d \lambda_1}{d \mu} = \frac{1}{16 \pi^2} \frac{1}{2} \lambda_1 (\lambda_1^2 + \lambda_2^2 + \lambda_3^2 + 5 \lambda^2 - 6 g^2)
+ \lambda_i (\lambda_i \lambda_1 + 3k_i \bar{k} k d_1 + 2k_i k l_1 + 2f_i f_1),$$
$$\mu \frac{d k_{d1}}{d \mu} = \frac{1}{16 \pi^2} [k_{d1} (2k_{d1}^2 + 2k_{d2}^2 + 2k_{d3}^2 + f_{d}^2 + f_{d}^2 - 16 \frac{3}{9} g_3^2 - \frac{4}{15} g_1^2)]
+ k_{d1} (\lambda_i \lambda_1 + 3k_i \bar{k} k d_1 + 2k_i k l_1 + 2f_i f_1),$$
$$\mu \frac{d k_{l1}}{d \mu} = \frac{1}{16 \pi^2} [k_{l1} (2k_{l1}^2 + 2k_{l2}^2 + 2k_{l3}^2 + f_{l}^2 + f_{l}^2 - 3 g_2^2 - \frac{3}{5} g_1^2)]
+ k_{l1} (\lambda_i \lambda_1 + 3k_i \bar{k} k d_1 + 2k_i k l_1 + 2f_i f_1),$$
$$\mu \frac{d f_1}{d \mu} = \frac{1}{16 \pi^2} [2f_1 (4f_{1}^2 + 4f_{2}^2 + 4f_{3}^2 + 3f_{d}^2 + 3f_{d}^2 + 2f_{l}^2 + 2f_{l}^2)]
+ f_i (\lambda_i \lambda_1 + 3k_i \bar{k} k d_1 + 2k_i k l_1 + 2f_i f_1),$$
$$\mu \frac{d \lambda_2}{d \mu} = \frac{1}{16 \pi^2} \frac{1}{2} \lambda_2 (\lambda_1^2 + \lambda_2^2 + \lambda_3^2 + 5 \lambda^2 - 6 g^2)
+ \lambda_i (\lambda_i \lambda_2 + 3k_i \bar{k} k d_2 + 2k_i k l_2 + 2f_i f_2),$$
$$\mu \frac{d k_{d2}}{d \mu} = \frac{1}{16 \pi^2} [k_{d2} (2k_{d1}^2 + 2k_{d2}^2 + 2k_{d3}^2 + f_{d}^2 + f_{d}^2 - 16 \frac{3}{9} g_3^2 - \frac{4}{15} g_1^2)]
+ k_{d2} (\lambda_i \lambda_2 + 3k_i \bar{k} k d_2 + 2k_i k l_2 + 2f_i f_2),$$
$$\mu \frac{d k_{l2}}{d \mu} = \frac{1}{16 \pi^2} [k_{l2} (2k_{l1}^2 + 2k_{l2}^2 + 2k_{l3}^2 + f_{l}^2 + f_{l}^2 - 3 g_2^2 - \frac{3}{5} g_1^2)]
+ k_{l2} (\lambda_i \lambda_2 + 3k_i \bar{k} k d_2 + 2k_i k l_2 + 2f_i f_2),$$
$$\mu \frac{d f_2}{d \mu} = \frac{1}{16 \pi^2} [2f_2 (4f_{1}^2 + 4f_{2}^2 + 4f_{3}^2 + 3f_{d}^2 + 3f_{d}^2 + 2f_{l}^2 + 2f_{l}^2)]$$
\( \mu \frac{d\lambda_3}{d\mu} = \frac{1}{16\pi^2} [\frac{1}{2} \lambda_3 (\lambda_1^2 + \lambda_2^2 + \lambda_3^2 + 5\lambda^2 - 6g^2) + \lambda_i (\lambda_i \lambda_3 + 3k_i k_d d_3 + 2k_i k_i l_3 + 2f_i f_3)], \)

\( \mu \frac{dk_{l3}}{d\mu} = \frac{1}{16\pi^2} [k_{l3} (2k_{f1}^2 + 2k_{f2}^2 + 2k_{f3}^2 + f_{l1}^2 + f_{l2}^2 - \frac{16}{3} g_3^2 - \frac{4}{15} g_1^2) + k_i (\lambda_i \lambda_3 + 3k_i k_d d_3 + 2k_i k_i l_3 + 2f_i f_3)], \)

\( \mu \frac{df_3}{d\mu} = \frac{1}{16\pi^2} [2f_3 (4f_{l1}^2 + 4f_{l2}^2 + 4f_{l3}^2 + 3f_{d1}^2 + 3f_{d2}^2 + 2f_{l1}^2 + 2f_{l2}^2) + f_i (\lambda_i \lambda_3 + 3k_i k_d d_3 + 2k_i k_i l_3 + 2f_i f_3)], \)

\( \mu \frac{df_d}{d\mu} = \frac{1}{16\pi^2} f_d (k_{d1}^2 + k_{d2}^2 + k_{d3}^2 + 4f_{l1}^2 + 4f_{l2}^2 + 4f_{l3}^2 + 3f_{d1}^2 + 3f_{d2}^2 + 2f_{l1}^2 + 2f_{l2}^2 - \frac{16}{3} g_3^2 - \frac{4}{15} g_1^2), \)

\( \mu \frac{df_{l1}}{d\mu} = \frac{1}{16\pi^2} f_i (k_{l1}^2 + k_{l2}^2 + k_{l3}^2 + 4f_{l1}^2 + 4f_{l2}^2 + 4f_{l3}^2 + 3f_{d1}^2 + 3f_{d2}^2 + 2f_{l1}^2 + 2f_{l2}^2 - \frac{4}{15} g_1^2), \)

\( \mu \frac{df_{l2}}{d\mu} = \frac{1}{16\pi^2} f_i (k_{l1}^2 + k_{l2}^2 + k_{l3}^2 + 4f_{l1}^2 + 4f_{l2}^2 + 4f_{l3}^2 + 3f_{d1}^2 + 3f_{d2}^2 + 2f_{l1}^2 + 2f_{l2}^2 - \frac{3}{5} g_2^2 - \frac{3}{5} g_1^2), \)

\( \mu \frac{df_{l3}}{d\mu} = \frac{1}{16\pi^2} f_i (k_{l1}^2 + k_{l2}^2 + k_{l3}^2 + 4f_{l1}^2 + 4f_{l2}^2 + 4f_{l3}^2 + 3f_{d1}^2 + 3f_{d2}^2 + 2f_{l1}^2 + 2f_{l2}^2 - \frac{3}{5} g_2^2 - \frac{3}{5} g_1^2), \)

where \( g, g_3, g_2, g_1 \) denote strong SU(2), SU(3)\(_C\), SU(2)\(_L\), U(1)\(_Y\) gauge coupling constants respectively and the index \( i \) is summed over 1, 2, 3.

Further, one-loop RGEs for the gauge coupling constants are given by

\[ \mu \frac{dg}{d\mu} = -\frac{4}{16\pi^2} g^3, \]
\[ \mu \frac{dg_3}{d\mu} = -\frac{1}{16\pi^2} g_3^3, \]
\[ \mu \frac{dg_2}{d\mu} = \frac{3}{16\pi^2} g_2^3, \]
\[ \mu \frac{dg_1}{d\mu} = \frac{1}{16\pi^2} \frac{43}{5} g_1^3. \]
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