Phenomenological aspects of a direct-transmission model of dynamical supersymmetry breaking with the gravitino mass $m_{3/2}<1$ keV

Yasunori Nomura and Kazuhiro Tobe*
Department of Physics, University of Tokyo, Bunkyo-ku, Hongou, Tokyo 113, Japan
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We analyze a direct-transmission model of dynamical supersymmetry breaking previously proposed. In the model the gravitino mass is naturally smaller than 1 keV, which is required from standard cosmology. We find that there are many distinguishable features with 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 ''$\gamma\gamma+$ missing energy'' in the present model. [S0556-2821(98)07915-6]

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I. 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 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 the 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 [1,2].

Several mechanisms for DSB have been discovered [3–6] and their applications to realistic models have been proposed in the literature [7–9]. Most of the models that 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.[1] 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 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-ray (γ-ray) background from 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 as $m_{3/2}<1$ keV to avoid the introduction of a complicated nonstandard 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 b-ino, can decay into the gravitino, which is the lightest superparticle (LSP), within detectors producing an observable ''$\gamma\gamma+$ missing energy'' signal as discussed in many papers [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 experiments in the near future such as at the CERN 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 that in ordinary gauge mediation models and in 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 GUT unification. Since the present model has many different features from 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 Sec. II, we briefly review the model in Ref. [16]. In Sec. III, we consider the low-energy mass spectrum of the gauginos and sfermions in the MSSM sector, and argue their typical features. Radiative electroweak symmetry breaking is also discussed. In Sec. IV, we analyze GUT models with and without Yukawa unification. In Sec. V, we discuss other interesting features in the present model. Section VI is devoted to our conclusions.

II. DIRECT-TRANSMISSION MODEL OF SUSY BREAKING

In this section, we review a model which has been proposed in Ref. [16]. Let us first consider the dynamics for scale generation. To generate the scale, we adopt a SUSY

*Address after September 1: Department of Physics, The Ohio State University, Columbus, OH 43210.

[1]In the model in Ref. [8] it seems possible to make the gravitino mass smaller than 1 keV due to a strong gauge interaction. If a naive dimensional analysis of strongly coupled SUSY theory [10] is applicable, however, the gravitino mass is likely larger than 1 keV [11].

SU(2) gauge theory with four doublet chiral superfields Q_i, where i is the flavor index ($i=1,\ldots,4$). Without a superpotential, this theory has 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 ($a=1,\ldots,5$) which constitute a five-dimensional representation of $SP(4)_F$ to obtain a tree-level superpotential

$$W_Y=\lambda_Y Y^a (QQ)_a, \tag{1}$$

where $(QQ)_a$ denote a five-dimensional representation of $SP(4)_F$ given by a suitable combination of gauge invariants Q_iQ_j.

A low-energy effective superpotential with W_Y in Eq. (1), which describes the dynamics of the SU(2) gauge interaction, may be given by [20]

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

in terms of low-energy degrees of freedom

$$V\sim(QQ),\quad V_a\sim(QQ)_a, \tag{3}$$

where S is an additional chiral superfield, Λ a dynamically generated scale, and a gauge-invariant (QQ) denotes a singlet of $SP(4)_F$ defined by

$$(QQ)=\frac{1}{2}(Q_1Q_2+Q_3Q_4). \tag{4}$$

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

$$\langle V\rangle=\Lambda^2, \tag{5}$$

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

Next we consider 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

$$W=W_Y+\lambda Z(QQ). \tag{6}$$

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

$$W_{eff}\simeq\lambda\Lambda^2 Z. \tag{7}$$

Then, $F_Z\simeq\lambda\Lambda^2\neq0$ and SUSY is broken.

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

$$K=|Z|^2-\frac{\eta}{4\Lambda^2}|\lambda Z|^4+\cdots, \tag{8}$$

where we expect that η is a real constant of order 1. Then the effective potential for the scalar Z (with the same notation as the superfield) is given by

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

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

In the following analyses, we assume the latter case $|\lambda\langle Z\rangle|\sim\Lambda$, which results in the breakdown of the R symmetry.[2] 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].[3]

To transmit the SUSY breaking effects to the MSSM sector, we introduce vectorlike 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.[4] We introduce the 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

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

In this case, however, the condensations $\langle d\bar{d}\rangle\neq0$ and $\langle l\bar{l}\rangle\neq0$ 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, and $m_{\bar{l}}$ as follows:

$$\begin{aligned}W_{eff}=&Z(\lambda\Lambda^2+k_d d\bar{d}+k_l l\bar{l})\\&+m_d d\bar{d}'+m_{\bar{d}}d'\bar{d}+m_l l\bar{l}'+m_{\bar{l}}l'\bar{l}.\end{aligned} \tag{11}$$

The dynamical generation of these mass parameters m_d, $m_{\bar{d}}$, m_l, 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

[2]Even if we take $|\lambda\langle Z\rangle|\sim4\pi\Lambda$, the main conclusion in the present paper does not change.

[3]When $\langle Z\rangle=0$, appropriate R-breaking mass terms such as $md\bar{d}+m'l\bar{l}$ are 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.

a SUSY-breaking vacuum with vacuum expectation values of the messenger squarks and sleptons vanishing:

$$\langle d\rangle=\langle\overline{d}\rangle=\langle l\rangle=\langle\overline{l}\rangle=\langle d'\rangle=\langle\overline{d'}\rangle=\langle l'\rangle=\langle\overline{l'}\rangle=0,$$
$$\langle F_Z\rangle\simeq\lambda\Lambda^2. \tag{12}$$

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

$$V=|\lambda\Lambda^2+k_d d\overline{d}+k_l l\overline{l}|^2+|m_{\overline{d}}\overline{d}|^2+|m_d d|^2+|m_{\overline{l}}\overline{l}|^2+|m_l l|^2+|k_d Z\overline{d}+m_d\overline{d'}|^2+|k_d Z d+m_{\overline{d}} d'|^2+|k_l Z\overline{l}+m_l\overline{l'}|^2+|k_l Z l+m_{\overline{l}} l'|^2 \tag{13}$$

as follows:

$$|m_d m_{\overline{d}}|^2>|k_d\langle F_Z\rangle|^2,$$
$$|m_l m_{\overline{l}}|^2>|k_l\langle F_Z\rangle|^2. \tag{14}$$

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 $Z(k_d d\overline{d}+k_l l\overline{l})$. We will see later that such a direct transmission of SUSY-breaking effects to the messenger sector makes it possible to realize the light gravitino $m_{3/2}<1$ keV, which is required from standard cosmology.

Now we discuss the dynamical generation of the mass parameters m_d, $m_{\overline{d}}$, m_l, and $m_{\overline{l}}$. 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 a 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:[5]

$$W=W_Y+Z_1(\lambda_1(QQ)+k_{d1}d\overline{d}+k_{l1}l\overline{l}-f_1X^2)+Z_2(\lambda_2(QQ)+k_{d2}d\overline{d}+k_{l2}l\overline{l}-f_2X^2)+Z_3(\lambda_3(QQ)+k_{d3}d\overline{d}+k_{l3}l\overline{l}-f_3X^2)+X(f_d d\overline{d'}+f_{\overline{d}}d'\overline{d}+f_l l\overline{l'}+f_{\overline{l}}l'\overline{l}). \tag{15}$$

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,

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

	$Q,\psi,\overline{\psi},X$	$Z_i,\psi',\overline{\psi'}$
$U(1)_R$	0	2
$U(1)_\chi$	1	-2

Eq. (15), is a general one allowed by the global $U(1)_R\times U(1)_\chi$.[6] The charges for chiral superfields are given in Table I.

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 condition

$$|f_\psi f_{\overline{\psi}}(\lambda_1f_1+\lambda_3f_3)|^2>|k_{\psi2}\lambda_2(f_1^2+f_3^2)+k_{\psi3}f_1(\lambda_3f_1-\lambda_1f_3)|^2 \tag{16}$$

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

$$\langle X\rangle=\sqrt{\frac{\lambda_1f_1+\lambda_3f_3}{f_1^2+f_3^2}}\Lambda, \tag{17}$$

and the 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

$$F_{Z_1}=\frac{\lambda_1f_3-\lambda_3f_1}{f_1^2+f_3^2}f_3\Lambda^2,\quad F_{Z_2}=\lambda_2\Lambda^2,$$
$$F_{Z_3}=\frac{\lambda_3f_1-\lambda_1f_3}{f_1^2+f_3^2}f_1\Lambda^2, \tag{18}$$

and thus SUSY is broken.

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

$$m_\psi=f_\psi\langle X\rangle, \tag{19}$$

$$m_{\overline{\psi}}=f_{\overline{\psi}}\langle X\rangle \tag{20}$$

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, λ_i, f_ψ, $f_{\overline{\psi}}$ are $O(1)$ because they are generated by the same dynamics with the scale Λ.

[5]We can construct a model in which only two singlet chiral supermultiplets are needed to give a 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 nonrenormalizable operators as discussed later.

[6]This global symmetry may forbid mixings between the messenger quarks and the down-type quarks in the SM sector. This avoids naturally the FCNC problem [23]. Then there exists the lightest stable particle in the messenger sector [24].

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 [7,25,26]. We will discuss the obtained mass spectrum in the MSSM sector in the next section.

III. MASS SPECTRUM OF THE SUPERPARTICLES IN THE MSSM SECTOR

A. 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 ψ, $\overline{\psi}$, ψ', and $\overline{\psi}'$ for $\psi = d,l$ is represented as

$$W = \sum_{\psi = d,l} (\overline{\psi}, \overline{\psi}') M^{(\psi)} \begin{pmatrix} \psi \\ \psi' \end{pmatrix}, \tag{21}$$

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

$$M^{(\psi)} = \begin{pmatrix} m^{(\psi)} & m_{\overline{\psi}} \\ m_{\psi} & 0 \end{pmatrix}. \tag{22}$$

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

$$m^{(\psi)} = k_{\psi 2}\langle Z_2\rangle + k_{\psi 3}\langle Z_3\rangle, \tag{23}$$

$$m_{\psi} = f_{\psi}\langle X\rangle, \tag{24}$$

$$m_{\overline{\psi}} = f_{\overline{\psi}}\langle X\rangle. \tag{25}$$

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

$$\mathrm{diag}(M_1^{(\psi)}, M_2^{(\psi)}) = U^{(\psi)} M^{(\psi)} V^{(\psi)\dagger}. \tag{26}$$

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

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

where

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

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

$$\mathcal{L}_s = -\sum_{\psi = d,l} (\tilde{\psi}^*, \tilde{\psi}'^*, \tilde{\overline{\psi}}, \tilde{\overline{\psi}}') \tilde{M}^{2(\psi)} \begin{pmatrix} \tilde{\psi} \\ \tilde{\psi}' \\ \tilde{\overline{\psi}}^* \\ \tilde{\overline{\psi}}'^* \end{pmatrix}, \tag{29}$$

where the mass matrix $\tilde{M}^{2(\psi)}$ is given by

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

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

$$\mathrm{diag}(m_1^{(\psi)2}, m_2^{(\psi)2}, m_3^{(\psi)2}, m_4^{(\psi)2}) = T^{(\psi)} \tilde{M}^{2(\psi)} T^{(\psi)\dagger}. \tag{31}$$

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.

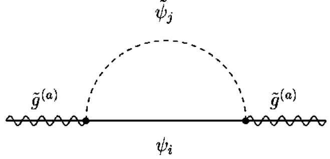

FIG. 1. Diagram contributing to the gaugino masses.

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

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

$$m_{\tilde{g}_2} = \frac{\alpha_2}{2\pi}\mathcal{F}^{(l)}, \tag{32}$$

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

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 $m_{\tilde{g}_i}$ ($i=1,\ldots,3$) denote the b-ino, W-ino, and gluino masses, respectively. Here $\mathcal{F}^{(d)}$ represents contributions from the messenger quark multiplets d, d', $\overline{d}$, and $\overline{d'}$ and $\mathcal{F}^{(l)}$ contributions from the messenger lepton multiplets l, l', $\overline{l}$, and $\overline{l'}$. The functions $\mathcal{F}^{(\psi)}$ ($\psi=d,l$) are given by

$$\mathcal{F}^{(\psi)}=\sum_{\alpha=1}^{2}\sum_{\beta=1}^{4} M_{\alpha}^{(\psi)}(U_{\alpha 1}^{(\psi)}T_{3\beta}^{(\psi)\dagger}+U_{\alpha 2}^{(\psi)}T_{4\beta}^{(\psi)\dagger}) \times(T_{\beta 1}^{(\psi)}V_{1\alpha}^{(\psi)\dagger}+T_{\beta 2}^{(\psi)}V_{2\alpha}^{(\psi)\dagger}) \times\frac{m_{\beta}^{(\psi)2}}{m_{\beta}^{(\psi)2}-M_{\alpha}^{(\psi)2}}\ln\frac{m_{\beta}^{(\psi)2}}{M_{\alpha}^{(\psi)2}}, \tag{33}$$

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

$$\mathcal{F}^{(\psi)}=\frac{1}{2}\left|\frac{F^{(\psi)}}{m_{\psi}m_{\overline{\psi}}}\right|^2\frac{F^{(\psi)}}{\sqrt{m_{\psi}m_{\overline{\psi}}}}\mathcal{A}^{(\psi)}(|V_{11}^{(\psi)}/V_{12}^{(\psi)}|^2,|U_{11}^{(\psi)}/U_{12}^{(\psi)}|^2), \tag{34}$$

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

$$\mathcal{A}^{(\psi)}(a,b)=\frac{(ab)^{1/4}}{6(1-ab)^4(1+a)^{3/2}(1+b)^{3/2}}\{2(a+b) \times[-1+8ab-8a^3b^3+a^4b^4+12a^2b^2\ln(ab)] -1-ab-64a^2b^2+64a^3b^3+a^4b^4 +a^5b^5-36a^2b^2(1+ab)\ln(ab)\}. \tag{35}$$

This function $\mathcal{A}^{(\psi)}(a,b)$ has the maximal value 0.1 at $a \simeq 3$ and $b\simeq 3$. We should note that the leading term of order $F^{(\psi)}/\sqrt{m_{\psi}m_{\overline{\psi}}}$ vanishes because $(M^{(\psi)})_{11}^{-1}=0$ [16].[7]

In ordinary gauge mediation models, contributions from the messenger quark multiplets are equal to those from 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, $\mathcal{F}^{(d)}\neq\mathcal{F}^{(l)}$ because the leading term of order $F^{(\psi)}/\sqrt{m_{\psi}m_{\overline{\psi}}}$ 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 is satisfied:

[7]This leading term cancellation of gaugino masses occurs generically, when we stabilize the SUSY-breaking vacuum by mass terms as in Eq. (11).

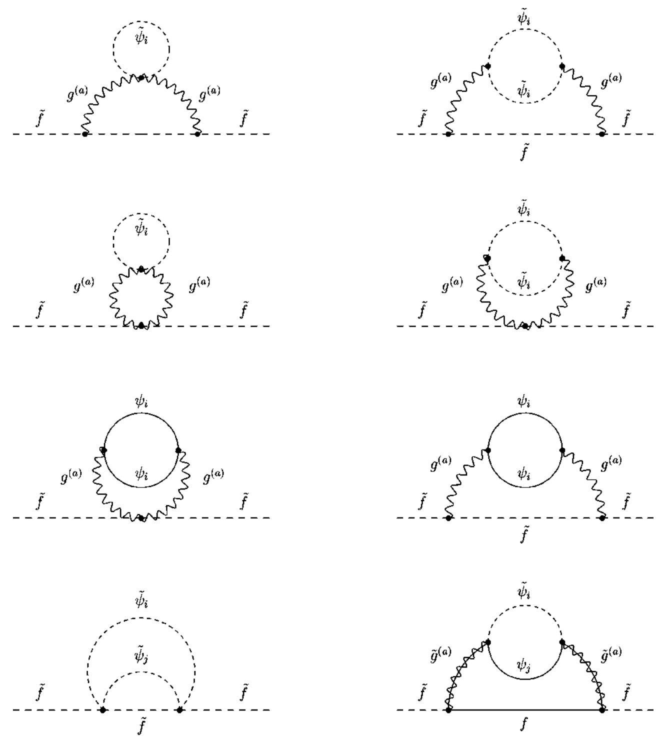

FIG. 2. Diagrams contributing to the squark, slepton, and Higgs boson masses.

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

This is a distinctive prediction in the present model.

The soft SUSY-breaking masses for squarks, sleptons, and Higgs bosons $\tilde{f}$ in the MSSM sector are generated by two-loop diagrams shown in Fig. 2 [7,25,26]. We obtain them as

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

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

Since global SUSY is spontaneously broken, there is a

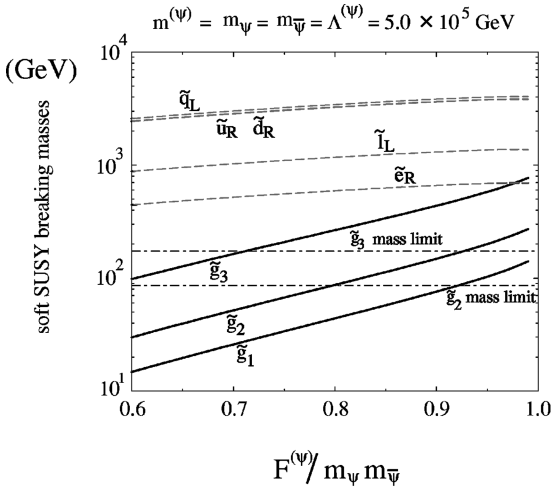

FIG. 3. The mass spectrum of the gauginos and sfermions as a function of the parameter $|F^{(\psi)}/m_\psi m_{\bar\psi}|$. Here we have assumed the mass parameters in the messenger sector as $m^{(\psi)}=m_\psi=m_{\bar\psi}\equiv\Lambda^{(\psi)}$ for $\psi=d,l$. We have set the scales $\Lambda^{(d)}=\Lambda^{(l)}=5.0\times10^5$ GeV. Solid lines represent the gluino $\tilde g_3$, W-ino $\tilde g_2$, and b-ino $\tilde g_1$ masses; the dashed lines denote the doublet squark $\tilde q_L$, right-handed up squark $\tilde u_R$, right-handed down squark $\tilde d_R$, 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 W-ino masses (dash-dotted lines).

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

$$m_{3/2}=\frac{F_{Z_1}+F_{Z_2}+F_{Z_3}}{\sqrt{3}M_*} =\left(\frac{\sqrt{F_{Z_1}+F_{Z_2}+F_{Z_3}}}{2\times10^6\ \text{GeV}}\right)^2 \ \text{keV}. \tag{38}$$

Here, we have imposed the vanishing cosmological constant and $M_*=2.4\times10^{18}$ GeV is the reduced Planck mass. If the SUSY-breaking scale is smaller than 2×10^6 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

$$m_{\tilde q}\sim\frac{\alpha_3}{4\pi}\sqrt{F_Z}. \tag{39}$$

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 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:

$$m^{(\psi)}=m_\psi=m_{\bar\psi}\equiv\Lambda^{(\psi)} \tag{40}$$

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

$$\Lambda^{(d)}=\Lambda^{(l)}=5.0\times10^5\ \text{GeV}, \tag{41}$$

which corresponds to the gravitino mass $m_{3/2}\sim0.1$ 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 $\mathcal{A}^{(\psi)}$ in Eq. (35) of the mixing angles of the messenger mass matrices. Thus, the gauginos are always lighter than sfermions in the present model.[8]

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 ($m_{\tilde g_3}>173$ GeV) [28] and the W-ino mass ($m_{\tilde g_2}>86$ GeV) [29] in Fig. 3.[9] The parameters $|F^{(d)}/m_d m_{\bar d}|$ and $|F^{(l)}/m_l m_{\bar l}|$ are independent from each other in general. Thus, they are constrained by the experimental bounds on the gluino and W-ino masses independently. From Fig. 3, we see that the parameters $|F^{(d)}/m_d m_{\bar d}|$ and $|F^{(l)}/m_l m_{\bar l}|$ are constrained as

[8]The mass spectrum that the squarks are relatively heavy compared with the W-ino is desirable from the viewpoint of the proton decay in the GUT case. The experimental constraint from the proton decay is about $m_{\tilde q}^2/m_{\tilde g_2}\gtrsim10^4$ GeV in the relevant region [27].

[9]The gluino mass is constrained as $m_{\tilde g_3}>173$ GeV for large squark masses by the Tevatron experiment [28]. The lightest chargino mass is constrained as $m_{\tilde\chi_1^+}>86$ GeV for large sneutrino mass and large μ parameter by the LEP experiment at $\sqrt{s}=172$ GeV [29]. In our model, the squark and slepton masses are much larger than the gaugino masses. The μ parameter also tends to be large as we will discuss in Sec. III B. Therefore, we use the gluino mass bound $m_{\tilde g_3}>173$ GeV and the W-ino mass bound $m_{\tilde g_2}>86$ GeV here.

$$0.71 < \left| \frac{F^{(d)}}{m_d m_{\bar{d}}} \right| < 1,$$

$$0.79 < \left| \frac{F^{(l)}}{m_l m_{\bar{l}}} \right| < 1. \qquad (42)$$

Here, the upper bounds come from the vacuum stability condition, Eq. (14). As we discussed in the previous section, the masses m_ψ 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 $|F^{(d)}/m_d m_{\bar{d}}|$ and $|F^{(l)}/m_l m_{\bar{l}}|$ are independent of each other. For example, when $|F^{(d)}/m_d m_{\bar{d}}| = 0.75$ and $|F^{(l)}/m_l m_{\bar{l}}| = 0.95$, the gluino and W-ino have almost the same mass: $m_{\tilde{g}_3} \simeq m_{\tilde{g}_2} \simeq 210$ GeV as seen in Fig. 3. It is remarkable that even in this case the gaugino mass relation, Eq. (36), holds and hence the b-ino 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 μ 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 μ parameter tends to be large (~ 1 TeV). Therefore the lightest neutralino and chargino are almost gauginos. In particular, the lightest neutralino is mostly b-ino because the b-ino mass is smaller than the W-ino mass in the relevant parameter space.

So far we have assumed the mass parameter relation, Eq. (40). Next we discuss the $m^{(\psi)}$ 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 $\psi = d,l$. 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 W-ino masses, the parameters $m^{(\psi)}/\sqrt{m_\psi m_{\bar{\psi}}}$ for $\psi = d,l$ are constrained as

$$0.21 < \frac{m^{(d)}}{\sqrt{m_d m_{\bar{d}}}} < 4.8,$$

$$0.36 < \frac{m^{(l)}}{\sqrt{m_l m_{\bar{l}}}} < 3.0. \qquad (43)$$

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 $k_{\psi i}\langle Z_i \rangle$ be almost at the same order as the strong SU(2) dynamical scale Λ.

The constraints in Eqs. (42), (43) will be weakened if the scales $\Lambda^{(\psi)}$ become larger. When we fix the parameter $|F^{(d)}/m_d m_{\bar{d}}|$ ($|F^{(l)}/m_l m_{\bar{l}}|$) and increase the scale $\Lambda^{(d)}$ ($\Lambda^{(l)}$), then the gluino mass (the W-ino mass) becomes larger proportionally to the scale $\Lambda^{(d)}$ ($\Lambda^{(l)}$) and the function $\mathcal{G}^{(d)2}$ ($\mathcal{G}^{(l)2}$) 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 a value

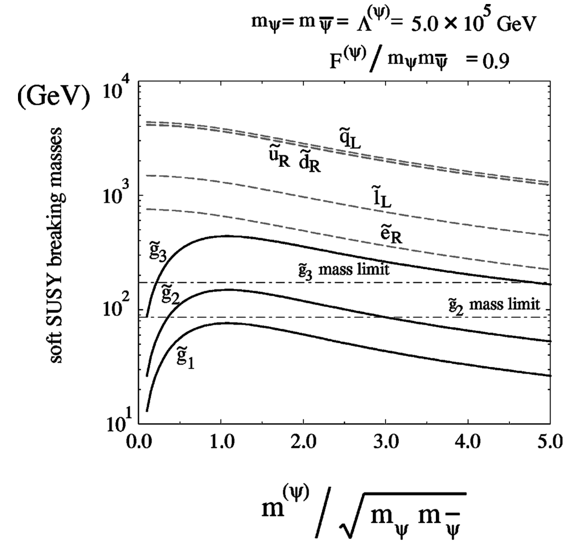

FIG. 4. The mass spectrum of the gauginos and sfermions in the MSSM sector as a function of the parameter $m^{(\psi)}/\sqrt{m_\psi m_{\bar{\psi}}}$ for $\psi = d,l$. Here we have assumed the mass parameters in the messenger sector as $m_\psi = m_{\bar{\psi}} \equiv \Lambda^{(\psi)}$ for $\psi = d,l$. We have set the scales $\Lambda^{(d)} = \Lambda^{(l)} = 5.0 \times 10^5$ GeV. Solid lines represent the gluino $\tilde{g}_3$, W-ino $\tilde{g}_2$, and b-ino $\tilde{g}_1$ masses; the dashed lines denote the doublet squark $\tilde{q}_L$, right-handed up squark $\tilde{u}_R$, right-handed down squark $\tilde{d}_R$, 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 W-ino masses.

$\Lambda^{(\psi)}$ 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 $\Lambda^{(d)} = \Lambda^{(l)} \leqslant 2 \times 10^6$ GeV. When we set $\Lambda^{(d)} = \Lambda^{(l)} = 2 \times 10^6$ GeV, the constraints on the parameters $|F^{(d)}/m_d m_{\bar{d}}|$ and $|F^{(l)}/m_l m_{\bar{l}}|$ get weaker than those in Eq. (42) as

$$0.47 < \left| \frac{F^{(d)}}{m_d m_{\bar{d}}} \right| < 1,$$

$$0.54 < \left| \frac{F^{(l)}}{m_l m_{\bar{l}}} \right| < 1. \qquad (44)$$

The constraints on the parameters $|m^{(d)}/\sqrt{m_d m_{\bar{d}}}|$ and $|m^{(l)}/\sqrt{m_l m_{\bar{l}}}|$ are also much weaker:

$$0.05 < \frac{m^{(d)}}{\sqrt{m_d m_{\bar{d}}}} < 19,$$

$$0.08 < \frac{m^{(l)}}{\sqrt{m_l m_{\bar{l}}}} < 12, \qquad (45)$$

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

B. Radiative electroweak symmetry breaking

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

The soft SUSY-breaking masses for the squarks, sleptons, and Higgs bosons 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:

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

where m_{H_1} and $m_{\tilde{t}}$ are the soft SUSY-breaking masses for another Higgs doublet and the top squark, respectively. Λ_{mess} denotes the messenger scale, and f_t represents the top quark Yukawa coupling constant. In the present model, the top squark is heavier than the SU(2) doublet Higgs bosons 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 $\Lambda_{mess} \sim 10^5$ GeV and the electroweak scale $m_{weak} \sim 10^2$ GeV is not so long. Therefore, electroweak symmetry is broken radiatively.

Requiring that the tree-level potential have 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

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

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

When we fix $\tan\beta$, the μ parameter is determined from Eq. (47) to reproduce the correct value for the Z boson mass m_Z. Once the parameter μ 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 radiative electroweak symmetry breaking. In our numerical analysis, we fix $\tan\beta$, and we use the one-loop effective potential [31] to determine the μ and B parameters. We include all of the third-generation Yukawa couplings. Here, we assume that some unknown dynamics generates the μ 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 Λ, $A_{\tilde{f}}(\Lambda) \simeq 0$, 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 μ parameters which satisfy the condition of 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)]. The solid line, dashed line, and long-dashed line correspond to the W-ino mass $m_{\tilde{g}_2} = 86$ GeV (the experimental lower bound), 150 GeV, and 200 GeV, respectively. As one sees from Fig. 5, the μ 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 a suppression like gauginos. Since the squarks, especially the top squark, become much heavier, the mass squared of H_2 becomes too negative as one can see in Eq. (46). Then one needs a large μ parameter to reproduce the correct value of the Z boson mass. [See Eq. (47).] For example, in the case with $\tan\beta = 3$ and $|F^{(\psi)}/m_\psi m_{\bar{\psi}}| = 0.6$, we need a large μ parameter as $|\mu| \gtrsim 3100$ GeV. In this way, we may need a fine-tuning of the μ parameter when the parameter $|F^{(\psi)}/m_\psi m_{\bar{\psi}}|$ 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 μ parameter gets smaller because the mass squared of H_2 becomes less negative due to the relatively smaller top quark Yukawa coupling.

We finally comment on the μ-term generation [7,33,34]. In the present model, the μ term may be generated in the same way as the generation of the messenger mass parameters: if the superfield X couples to $H_1 H_2$, the SUSY-invariant mass μ for Higgs bosons 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 λ_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 λ_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

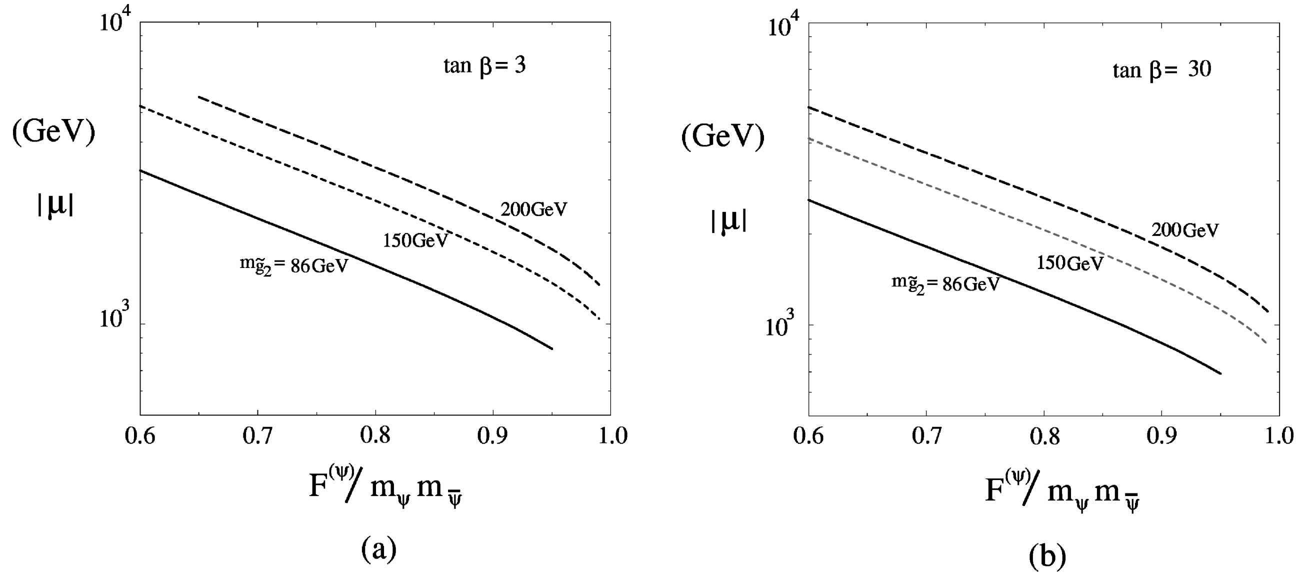

FIG. 5. The μ parameters which satisfy the condition of the radiative electroweak symmetry breaking as a function of the parameter $F^{(\psi)}/m_\psi m_{\bar{\psi}}$. Here we take (a) $\tan\beta=3$ and (b) 30. Solid line, dashed line, and long-dashed line correspond to the W-ino mass $m_{\tilde{g}_2}=86$ GeV (the experimental lower bound), 150 GeV, and 200 GeV, respectively.

small.[10] Hence the scale μ may originate from the same dynamics as SUSY breaking. In this case, a large $\tan\beta\sim 50$ is required because of the small B parameter to break electroweak symmetry radiatively.

IV. GRAND UNIFIED MODEL

A. GUT model with Yukawa unification

The messenger quarks and leptons (d,l), (d',l'), $(\overline{d},\overline{l})$, and $(\overline{d'},\overline{l'})$ are embedded in **5** and **5*** representations of the SU(5) group. Therefore, if we extend our model to the GUT model (without nonrenormalizable interactions), certain Yukawa coupling constants in Eq. (15) are unified as

$$k_{di}=k_{li} \quad (i=1\text{–}3), \tag{49}$$

$$f_d=f_l, \tag{50}$$

$$f_{\bar{d}}=f_{\bar{l}}, \tag{51}$$

at the GUT scale $M_{\text{GUT}}\simeq 2\times 10^{16}$ GeV. Under 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 $f_{\bar{d}}$ and $f_{\bar{l}}$ 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 $g_3(\mu)\geqslant g_2(\mu)$ at $\mu\leqslant M_{\text{GUT}}$. We consider the RGEs for the ratios of the Yukawa couplings k_{di}/k_{li}, f_d/f_l, and $f_{\bar{d}}/f_{\bar{l}}$:

$$16\pi^2\mu\frac{d}{d\mu}\left(\frac{k_{di}}{k_{li}}\right)=\frac{k_{di}}{k_{li}}\left\{-\frac{16}{3}g_3^2+3g_2^2+\frac{1}{3}g_1^2+2\sum_{j=1}^{3}(k_{dj}^2-k_{lj}^2)+f_d^2-f_l^2+f_{\bar{d}}^2-f_{\bar{l}}^2\right\}$$
$$+\sum_{j=1}^{3}\frac{k_{dj}}{k_{li}}\left(1-\frac{k_{di}}{k_{li}}\frac{k_{lj}}{k_{dj}}\right)(\lambda_i\lambda_j+3k_{di}k_{dj}+2k_{li}k_{lj}+2f_if_j), \tag{52}$$

[10]In this case, the SUSY CP problem is also solved. When we consider 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 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 Higgs bosons since there are no A terms and no B term at the tree level.

$$16\pi^2\mu\frac{d}{d\mu}\left(\frac{f_d}{f_l}\right)=\frac{f_d}{f_l}\left\{-\frac{16}{3}g_3^2+3g_2^2+\frac{1}{3}g_1^2+\sum_{i=1}^{3}(k_{di}^2-k_{li}^2)+2(f_d^2-f_l^2)\right\},\tag{53}$$

$$16\pi^2\mu\frac{d}{d\mu}\left(\frac{f_{\bar{d}}}{f_{\bar{l}}}\right)=\frac{f_{\bar{d}}}{f_{\bar{l}}}\left\{-\frac{16}{3}g_3^2+3g_2^2+\frac{1}{3}g_1^2+\sum_{i=1}^{3}(k_{di}^2-k_{li}^2)+2(f_{\bar{d}}^2-f_{\bar{l}}^2)\right\}.\tag{54}$$

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 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 scale. Then we can always set $k_{d1}=k_{l1}=k_{d2}=k_{l2}=0$ and $f_2=0$ at any scale by a 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 $f_{\bar{d}}/f_{\bar{l}}$ equal those of the mass parameters for the messenger sector $m^{(d)}/m^{(l)}$, m_d/m_l, and $m_{\bar{d}}/m_{\bar{l}}$, respectively.

To numerically analyze the relations between mass parameters in the messenger sector, we fix some parameters[11] as $\lambda_Y=0.3$, $(\lambda_1,\lambda_2,\lambda_3)=(0.2,0.1,0.16)$, $(f_1,f_2,f_3)=(0.5,0,0.6)$, $(k_{d1},k_{d2})=(k_{l1},k_{l2})=(0,0)$ at the GUT scale, $\Lambda=2\times10^6$ GeV, which is the dynamical scale of the strong SU(2) gauge interaction, and $\langle Z_3\rangle=1\times10^6$ GeV, which is the vacuum expectation value of the singlet field Z_3. The strong SU(2) gauge coupling g is taken as 2π 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 $f_{\bar{d}}=f_{\bar{l}}=f_{\bar{\psi}}$ at the GUT scale. We find the ratios of the mass parameters given by

$$\frac{m^{(d)}}{m^{(l)}}\simeq1.4\ \ (1.4),\tag{55}$$

$$\frac{m_d}{m_l}=\frac{m_{\bar{d}}}{m_{\bar{l}}}\simeq1.4\ \ (1.4),\tag{56}$$

$$0.5\ \ (0.4)<\frac{m^{(l)}}{\sqrt{m_l m_{\bar{l}}}}\simeq\frac{m^{(d)}}{\sqrt{m_d m_{\bar{d}}}}<1.2\ \ (0.9),\tag{57}$$

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 $F^{(l)}/m_l m_{\bar{l}}>0.7$. The ratio of $m^{(d)}/m^{(l)}$ is almost the same as the ratio of $m_d/m_l=m_{\bar{d}}/m_{\bar{l}}$ because the running effects for the ratios of the Yukawa couplings k_d/k_l, f_d/f_l, and $f_{\bar{d}}/f_{\bar{l}}$ 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 $F^{(l)}/m_l m_{\bar{l}}$ as follows:

$$\frac{F^{(l)}}{m_l m_{\bar{l}}}\simeq1.4\ \frac{F^{(d)}}{m_d m_{\bar{d}}},\tag{58}$$

since

$$\left(\frac{F^{(l)}}{m_l m_{\bar{l}}}\right)\Bigg/\left(\frac{F^{(d)}}{m_d m_{\bar{d}}}\right)=\frac{k_{l3}}{k_{d3}}\frac{m_d}{m_l}\frac{m_{\bar{d}}}{m_{\bar{l}}}=\frac{m^{(l)}}{m^{(d)}}\frac{m_d}{m_l}\frac{m_{\bar{d}}}{m_{\bar{l}}}.\tag{59}$$

Because of the Yukawa unification, $F^{(l)}/m_l m_{\bar{l}}$ becomes larger than $F^{(d)}/m_d m_{\bar{d}}$ 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 $F^{(l)}/m_l m_{\bar{l}}$. Here we have taken the same parameter set as the 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 $m_{3/2}\simeq0.12$ 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 Yukawa unification, Eqs. (49)–(51), the parameter $F^{(d)}/m_d m_{\bar{d}}$ is smaller than $F^{(l)}/m_l m_{\bar{l}}$ as shown in Eq. (58). Therefore the gluino mass receives a larger suppression than the W-ino mass, and hence the gluino tends to be relatively light. The experimental lower bound on the gluino mass constrains the parameter $F^{(l)}/m_l m_{\bar{l}}$ to be $F^{(l)}/m_l m_{\bar{l}}>0.89$. This leads to a constraint on the Yukawa coupling k at the GUT scale as $k>0.65$.[12] 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 W-ino.

[11]The Yukawa couplings $k_i, f_\psi, f_{\bar{\psi}}$ do not significantly affect the runnings of λ_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 $k_i, f_\psi, f_{\bar{\psi}}$.

[12]When we regard the masses for the gauginos and sfermions as a function of $F^{(l)}/m_l m_{\bar{l}}$, 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 $F^{(l)}/m_l m_{\bar{l}}$ and $F^{(d)}/m_d m_{\bar{d}}$ is almost independent of $f_\psi=f_{\bar{\psi}}$.

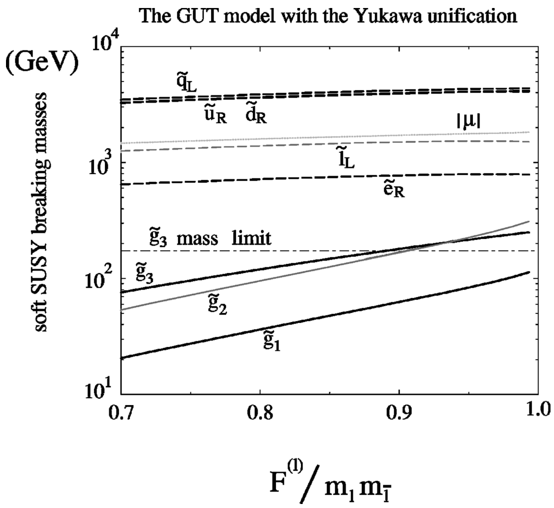

FIG. 6. The mass spectrum of the gauginos and sfermions in the MSSM sector as a function of $F^{(l)}/m_l m_{\bar{l}}$. Here we have taken $\lambda_Y=0.3$, $(\lambda_1,\lambda_2,\lambda_3)=(0.2,0.1,0.16)$, $(f_1,f_2,f_3)=(0.5,0,0.6)$, $(k_{d1},k_{d2})=(k_{l1},k_{l2})=(0,0)$ at the GUT scale. The dynamical scale Λ of the strong SU(2) gauge interaction and the vacuum expectation value $\langle Z_3\rangle$ have been taken as $\Lambda=2.0\times10^6$ GeV and $\langle Z_3\rangle=1.0\times10^6$ GeV, respectively. We have set $f_d=f_l=f_{\bar{d}}=f_{\bar{l}}=0.2$ 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 $m_{3/2}\simeq0.12$ keV. Solid lines represent the gaugino masses: the b-ino $\tilde{g}_1$, W-ino $\tilde{g}_2$, and gluino $\tilde{g}_3$ masses. Dashed lines represent the sfermion masses: the left-handed squark $\tilde{q}_L$, right-handed up squark $\tilde{u}_R$, right-handed down squark $\tilde{d}_R$, doublet slepton $\tilde{l}_L$, and right-handed selectron $\tilde{e}_R$ masses. The μ 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).

We also show the μ parameter in Fig. 6. As we discussed in Sec. III B, the μ parameter becomes much larger than the gaugino masses.

B. GUT model without Yukawa unification

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

$$W=\sum_{i=1}^{3} Z_i\left(k'_{\psi_i}\psi\bar{\psi}+k''_{\psi_i}\frac{\Sigma}{M_*}\psi\bar{\psi}+\cdots\right)+X\left(f'_\psi\psi\bar{\psi}'+f'_{\bar{\psi}}\psi'\bar{\psi}+f''_\psi\frac{\Sigma}{M_*}\psi\bar{\psi}'+f''_{\bar{\psi}}\frac{\Sigma}{M_*}\psi'\bar{\psi}\right), \tag{60}$$

where the fields ψ and ψ' [$\bar{\psi}$ and $\bar{\psi}'$] are **5** [**5***] dimensional representation fields of the SU(5) group, which contain the messenger quark and lepton multiplets (d,l) and (d',l') [$(\overline{d},\overline{l})$ and $(\overline{d',l'})$], respectively. The field Σ 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 $\langle\Sigma\rangle=V\mathrm{diag}(2,2,2,-3,-3)$. We should notice that the nonrenormalizable interactions in Eq. (60) are not forbidden by any symmetries provided that the Σ is a trivial representation of $U(1)_R\times U(1)_\chi$. Then, the Yukawa coupling constants k_ψ, f_ψ, and $f_{\bar{\psi}}$ receive corrections of order of $O(\langle\Sigma\rangle/M_*)$ as

$$k_{di}=k'_{\psi i}+2k''_{\psi i}\frac{V}{M_*}, \tag{61}$$

$$k_{li}=k'_{\psi i}-3k''_{\psi i}\frac{V}{M_*}, \tag{62}$$

$$f_d=f'_\psi+2f''_\psi\frac{V}{M_*}, \tag{63}$$

$$f_l=f'_\psi-3f''_\psi\frac{V}{M_*}, \tag{64}$$

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 nonrenormalizable terms, and the relation in Eq. (58) does not hold. Thus the mass relation among the gauginos in the case with Yukawa unification may easily change. Therefore, the existence of nonrenormalizable 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:

$$k_{d2}=k_2-\delta,\quad k_{l2}=k_2,$$
$$k_{d3}=k_3-\delta,\quad k_{l3}=k_3,$$
$$f_d=f-\delta,\quad f_l=f,$$
$$f_{\bar{d}}=\bar{f}-\delta,\quad f_{\bar{l}}=\bar{f}. \tag{65}$$

The parameter δ represents the difference between the Yukawa couplings of messenger quarks and leptons due to corrections from nonrenormalizable interactions. The case $\delta=0$ corresponds to the GUT model with Yukawa unification. We show the mass spectrum of the gauginos and sfermions as a function of the parameter δ 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 Λ 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 $\Lambda=1.0\times10^6$ GeV and $\langle Z_2\rangle=\langle Z_3\rangle=5.0\times10^5$ GeV, respectively. We have taken the param-

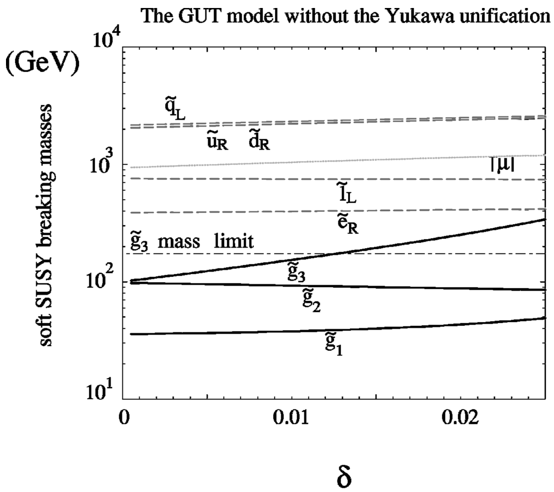

FIG. 7. The mass spectrum of the gauginos and sfermions as a function of the parameter δ. Here we have taken the Yukawa couplings at the GUT scale as follows: $\lambda_Y=0.3$, $(\lambda_1,\lambda_2,\lambda_3)=(0.2,0.1,0.16)$, $(f_1,f_2,f_3)=(0.5,0,0.6)$, $k_{d1}=k_{l1}=0$. We have set $k_2=0$, $k_3=0.8$, and $f=\overline{f}=0.2$ in Eq. (65). Solid lines represent the gaugino masses: the b-ino $\tilde{g}_1$, W-ino $\tilde{g}_2$, and gluino $\tilde{g}_3$ masses. Dashed lines represent the sfermion masses: the left-handed squark $\tilde{q}_L$, right-handed up squark $\tilde{u}_R$, right-handed down squark $\tilde{d}_R$, doublet slepton $\tilde{l}_L$, and right-handed selectron $\tilde{e}_R$ masses. The μ 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).

eters $k_2=0$, $k_3=0.8$, and $f=\overline{f}=0.2$ in Eq. (65). From Fig. 7, we see that the case with 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 nonrenormalizable terms may significantly change the mass spectrum. This is because the parameter $F^{(d)}/m_d m_{\overline{d}}$ gets larger than that in the case with Yukawa unification with $F^{(l)}/m_l m_{\overline{l}}$ 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 nonrenormalizable 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 a suppression factor.

V. OTHER INTERESTING FEATURES IN THE PRESENT MODEL

From the naturalness point of view as discussed in Sec. III B, 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, a lighter gravitino which satisfies the cosmological requirement $m_{3/2}<1$ keV. Such a light gravitino may bring us other interesting consequences in collider experiments. When the gravitino is so light that the cosmological requirement is satisfied, the gravitino becomes the LSP. Thus the NLSP, primarily the b-ino, decays into a gravitino emitting a photon. The Goldstino component of the gravitino has the following interaction with matter [35]:

$$\mathcal{L}=-\frac{1}{F}j^{\alpha\mu}\partial_\mu\tilde{G}_\alpha,$$

$$=\cos\,\theta\frac{m_{\tilde{g}_1}}{2\sqrt{2}F}\tilde{B}\overline{\sigma^\mu}\sigma^\nu\tilde{G}F_{\mu\nu}+\cdots. \quad (66)$$

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 a gravitino may occur within the detector.

From the interaction, Eq. (66), we find the decay rate of the b-ino as

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

Then the decay length L of the b-ino with energy E in the laboratory frame is given by

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

For example, when $\sqrt{F}=5\times10^5$ GeV which corresponds to the gravitino mass with $m_{3/2}=0.06$ keV and $m_{\tilde{g}_1}=100$ GeV, the decay length of the b-ino is 8 cm 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 $m_{3/2}<1$ 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'' are not known except for the present model. The phenomenological investigations have been done in many papers [17,18] only under the assumption of the existence of a very light gravitino.[13]

Recently, however, constraints on the SUSY models with a ''$\gamma\gamma+$ missing energy'' signal have been reported in Ref.

[13]The signature ''$\gamma\gamma+$ missing energy'' does not necessarily suggest a light gravitino. It has been known that a light axino in the framework of no-scale supergravity also brings us such a signal in collider experiments [36]. Therefore, other experiments, axion search and direct superparticle search, for example, are also necessary to distinguish the present model from other models.

[37]. The lower bounds on the lightest neutralino and the lightest chargino masses are obtained as

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

$$m_{\tilde{g}_2} \gtrsim 150 \ \text{GeV}, \tag{69}$$

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 ($m_{\tilde{f}} \sim 10$ TeV) or the parameter $F^{(\psi)}/m_\psi m_{\bar{\psi}}$ should be very close to 1 in the present model. If the squarks and sleptons are very heavy, $m_{\tilde{f}} \sim 10$ TeV, a fine-tuning may be needed to obtain the correct electroweak scale as discussed in Sec. III B. 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 b-ino mass as well as the SUSY-breaking scale. For example, when $\sqrt{F} = 5 \times 10^5$ GeV and $m_{\tilde{g}_1} = 40$ GeV, $L = 8$ m for $(E^2/m_{\tilde{g}_1}^2 - 1)^{1/2} = 1$ which is larger than the typical detector size. Therefore, the signal ''$\gamma\gamma+$ missing energy'' cannot be observed within the detector when the b-ino mass is much smaller than 100 GeV. In the case where the b-ino slowly decays outside the detector, the bounds in Eq. (69) are not applicable because they are derived under the assumption that the b-ino completely decays into a 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 existing experiments, pair production of light gauginos can be observed directly in 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 b-ino may be detectable in experiments in the near future such as LHC, even for the case of a long L.[14]

Finally we remark on the unification of the gauge coupling constants. To break 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 Λ as $\Lambda \simeq 10^{5-6}$ GeV in order that $m_{3/2} \sim 0.1$–1 keV [16].

[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 cannot be observable at LHC even with the same L.

VI. 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 is no SUSY FCNC, all mass scales are generated from the strong SU(2) dynamics (the μ 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 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 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 μ parameter, and we may need a fine-tuning. If we consider the GUT model with Yukawa unification, the gluino mass tends to be lighter than the W-ino mass. However, Yukawa unification is broken in the presence of nonrenormalizable 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 b-ino is relatively heavy ($m_{\tilde{g}_1} \gtrsim 100$ GeV). From the naturalness point of view, however, light gauginos are most likely in our model. Thus the collider signature will not be observed in present experiments because the b-ino decays outside the detector. Even in this case, b-ino decay may be detectable [19] in colliders of the near future because of the relatively lower SUSY-breaking scale.

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APPENDIX 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 Sec. III A, gaugino-fermion-scalar interaction is

$$\mathcal{L}_s = \sqrt{2} i g \sum_{\alpha=1}^{2} \sum_{\beta=1}^{4} [(T_{\beta 1}^{(\psi)} V_{1\alpha}^{(\psi)\dagger} + T_{\beta 2}^{(\psi)} V_{2\alpha}^{(\psi)\dagger}) \tilde{\psi}_{i\beta}^{*} \mathrm{T}_{ij}^{(a)} \psi_{j\alpha} \tilde{g}^{(a)} - (V_{\alpha 1}^{(\psi)} T_{1\beta}^{(\psi)\dagger} + V_{\alpha 2}^{(\psi)} T_{2\beta}^{(\psi)\dagger}) \tilde{g}^{(a)*} \psi_{i\alpha}^{*} \mathrm{T}_{ij}^{(a)} \tilde{\psi}_{j\beta} - (U_{\alpha 1}^{(\psi)} T_{3\beta}^{(\psi)\dagger} + U_{\alpha 2}^{(\psi)} T_{4\beta}^{(\psi)\dagger}) \overline{\psi}_{i\alpha} \mathrm{T}_{ij}^{(a)} \tilde{\psi}_{j\beta} \tilde{g}^{(a)} + (T_{\beta 3}^{(\psi)} U_{1\alpha}^{(\psi)\dagger} + T_{\beta 4}^{(\psi)} U_{2\alpha}^{(\psi)\dagger}) \tilde{g}^{(a)*} \tilde{\psi}_{i\beta}^{*} \mathrm{T}_{ij}^{(a)} \overline{\psi}_{j\alpha}^{*}], \tag{A1}$$

where ψ_α and $\overline{\psi}_\alpha$ denote the messenger fermions in the mass eigenstates and $\tilde{\psi}_\alpha$ is the messenger sfermions in the mass eigenstates as follows:

$$\begin{pmatrix} \psi_1 \\ \psi_2 \end{pmatrix} = V^{(\psi)} \begin{pmatrix} \psi \\ \psi' \end{pmatrix}, \tag{A2}$$

$$\begin{pmatrix} \overline{\psi}_1 \\ \overline{\psi}_2 \end{pmatrix} = U^{(\psi)*} \begin{pmatrix} \overline{\psi} \\ \overline{\psi}' \end{pmatrix}, \tag{A3}$$

$$\begin{pmatrix} \tilde{\psi}_1 \\ \tilde{\psi}_2 \\ \tilde{\psi}_3 \\ \tilde{\psi}_4 \end{pmatrix} = T^{(\psi)} \begin{pmatrix} \tilde{\psi} \\ \tilde{\psi}' \\ \tilde{\overline{\psi}}^{*} \\ \tilde{\overline{\psi}}'^{*} \end{pmatrix}, \tag{A4}$$

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

$$m_{\tilde{g}_3} = \frac{\alpha_3}{2\pi} \mathcal{F}^{(d)}, \tag{A5}$$

$$m_{\tilde{g}_2} = \frac{\alpha_2}{2\pi} \mathcal{F}^{(l)}, \tag{A6}$$

$$m_{\tilde{g}_1} = \frac{\alpha_1}{2\pi} \left\{ \frac{2}{5} \mathcal{F}^{(d)} + \frac{3}{5} \mathcal{F}^{(l)} \right\}, \tag{A7}$$

where the masses $m_{\tilde{g}_i}$ $(i=1,\ldots,3)$ denote the b-ino, W-ino, and gluino masses, respectively, and we have adopted the SU(5) GUT normalization of the $\mathrm{U}(1)_Y$ gauge coupling $(\alpha_1 \equiv \frac{5}{3} \alpha_Y)$. The functions $\mathcal{F}^{(\psi)}$ are given by

$$\begin{aligned} \mathcal{F}^{(\psi)} &= 32\pi^2 i \sum_{\alpha=1}^{2} \sum_{\beta=1}^{4} (U_{\alpha 1}^{(\psi)} T_{3\beta}^{(\psi)\dagger} + U_{\alpha 2}^{(\psi)} T_{4\beta}^{(\psi)\dagger}) (T_{\beta 1}^{(\psi)} V_{1\alpha}^{(\psi)\dagger} + T_{\beta 2}^{(\psi)} V_{2\alpha}^{(\psi)\dagger}) \mathrm{T}_{ij}^{(a)} \mathrm{T}_{ji}^{(a)} M_\alpha^{(\psi)} \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(p-k)^2 + m_\beta^{(\psi)2}} \frac{1}{k^2 + M_\alpha^{(\psi)2}} \Bigg|_{p=0} \\ &= 16\pi^2 i \sum_{\alpha=1}^{2} \sum_{\beta=1}^{4} (U_{\alpha 1}^{(\psi)} T_{3\beta}^{(\psi)\dagger} + U_{\alpha 2}^{(\psi)} T_{4\beta}^{(\psi)\dagger}) (T_{\beta 1}^{(\psi)} V_{1\alpha}^{(\psi)\dagger} + T_{\beta 2}^{(\psi)} V_{2\alpha}^{(\psi)\dagger}) \\ &\quad \times M_\alpha^{(\psi)} \left\{ -\frac{i}{16\pi^2} \int_0^1 dx \ln[x m_\beta^{(\psi)2} + (1-x) M_\alpha^{(\psi)2} - p^2 x(1-x)] \right\} \Bigg|_{p=0} \\ &= \sum_{\alpha=1}^{2} \sum_{\beta=1}^{4} M_\alpha^{(\psi)} (U_{\alpha 1}^{(\psi)} T_{3\beta}^{(\psi)\dagger} + U_{\alpha 2}^{(\psi)} T_{4\beta}^{(\psi)\dagger}) (T_{\beta 1}^{(\psi)} V_{1\alpha}^{(\psi)\dagger} + T_{\beta 2}^{(\psi)} V_{2\alpha}^{(\psi)\dagger}) \int_0^1 dx \ln[x m_\beta^{(\psi)2} + (1-x) M_\alpha^{(\psi)2}] \\ &= \sum_{\alpha=1}^{2} \sum_{\beta=1}^{4} M_\alpha^{(\psi)} (U_{\alpha 1}^{(\psi)} T_{3\beta}^{(\psi)\dagger} + U_{\alpha 2}^{(\psi)} T_{4\beta}^{(\psi)\dagger}) (T_{\beta 1}^{(\psi)} V_{1\alpha}^{(\psi)\dagger} + T_{\beta 2}^{(\psi)} V_{2\alpha}^{(\psi)\dagger}) \frac{m_\beta^{(\psi)2}}{m_\beta^{(\psi)2} - M_\alpha^{(\psi)2}} \ln \frac{m_\beta^{(\psi)2}}{M_\alpha^{(\psi)2}}. \end{aligned} \tag{A8}$$

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

$$m_{\tilde{f}}^2 = \frac{1}{2} \left[C_3^{\tilde{f}} \left(\frac{\alpha_3}{4\pi} \right)^2 \mathcal{G}^{(d)2} + C_2^{\tilde{f}} \left(\frac{\alpha_2}{4\pi} \right)^2 \mathcal{G}^{(l)2} + \frac{3}{5} Y^2 \left(\frac{\alpha_1}{4\pi} \right)^2 \left(\frac{2}{5} \mathcal{G}^{(d)2} + \frac{3}{5} \mathcal{G}^{(l)2} \right) \right], \tag{A9}$$

where $C_3^{\tilde{f}} = \frac{4}{3}$ and $C_2^{\tilde{f}} = \frac{3}{4}$ when $\tilde{f}$ is in the fundamental representation of $\mathrm{SU}(3)_C$ and $\mathrm{SU}(2)_L$, and $C_i^{\tilde{f}} = 0$ for the gauge singlets, and Y denotes the $\mathrm{U}(1)_Y$ hypercharge $(Y \equiv Q - T_3)$. Here $\mathcal{G}^{(\psi)2}$ are given by [25]

$$\mathcal{G}^{(\psi)2}=(4\pi)^4\Bigg\{-\sum_{\alpha=1}^{4}\langle m_\alpha^{(\psi)}|m_\alpha^{(\psi)}|0\rangle-4\sum_{\alpha=1}^{4}\langle m_\alpha^{(\psi)}|m_\alpha^{(\psi)}|0,0\rangle-4\sum_{\alpha=1}^{2}\langle M_\alpha^{(\psi)}|M_\alpha^{(\psi)}|0\rangle+8\sum_{\alpha=1}^{2}\langle M_\alpha^{(\psi)}|M_\alpha^{(\psi)}|0,0\rangle$$
$$-\sum_{\alpha=1}^{4}\sum_{\beta=1}^{4}(T_{\alpha1}^{(\psi)}T_{1\beta}^{(\psi)\dagger}+T_{\alpha2}^{(\psi)}T_{2\beta}^{(\psi)\dagger}-T_{\alpha3}^{(\psi)}T_{3\beta}^{(\psi)\dagger}-T_{\alpha4}^{(\psi)}T_{4\beta}^{(\psi)\dagger})$$
$$\times(T_{\beta1}^{(\psi)}T_{1\alpha}^{(\psi)\dagger}+T_{\beta2}^{(\psi)}T_{2\alpha}^{(\psi)\dagger}-T_{\beta3}^{(\psi)}T_{3\alpha}^{(\psi)\dagger}-T_{\beta4}^{(\psi)}T_{4\alpha}^{(\psi)\dagger})\langle m_\alpha^{(\psi)}|m_\beta^{(\psi)}|0\rangle$$
$$+4\sum_{\alpha=1}^{4}\sum_{\beta=1}^{2}[(V_{\beta1}^{(\psi)}T_{1\alpha}^{(\psi)\dagger}+V_{\beta2}^{(\psi)}T_{2\alpha}^{(\psi)\dagger})(T_{\alpha1}^{(\psi)}V_{1\beta}^{(\psi)\dagger}+T_{\alpha2}^{(\psi)}V_{2\beta}^{(\psi)\dagger})+(U_{\beta1}^{(\psi)}T_{3\alpha}^{(\psi)\dagger}+U_{\beta2}^{(\psi)}T_{4\alpha}^{(\psi)\dagger})(T_{\alpha3}^{(\psi)}U_{1\beta}^{(\psi)\dagger}+T_{\alpha4}^{(\psi)}U_{2\beta}^{(\psi)\dagger})]$$
$$\times[\langle m_\alpha^{(\psi)}|M_\beta^{(\psi)}|0\rangle+(m_\alpha^{(\psi)2}-M_\beta^{(\psi)2})\langle m_\alpha^{(\psi)}|M_\beta^{(\psi)}|0,0\rangle]\Bigg\}, \tag{A10}$$

where

$$\langle m_1|m_2|m_3\rangle=\int\frac{d^4q}{(2\pi)^4}\frac{d^4k}{(2\pi)^4}\frac{1}{(q^2+m_1^2)(k^2+m_2^2)([k-q]^2+m_3^2)}, \tag{A11}$$

$$\langle m_1|m_2|m_3,m_3\rangle=\int\frac{d^4q}{(2\pi)^4}\frac{d^4k}{(2\pi)^4}\frac{1}{(q^2+m_1^2)(k^2+m_2^2)([k-q]^2+m_3^2)^2}. \tag{A12}$$

This becomes

$$\mathcal{G}^{(\psi)2}=\sum_{\alpha=1}^{4}m_\alpha^{(\psi)2}(4\ln m_\alpha^{(\psi)2}+\ln^2 m_\alpha^{(\psi)2})+\sum_{\alpha=1}^{2}M_\alpha^{(\psi)2}(-8\ln M_\alpha^{(\psi)2}+4\ln^2 M_\alpha^{(\psi)2})$$
$$+\sum_{\alpha=1}^{4}\sum_{\beta=1}^{4}(T_{\alpha1}^{(\psi)}T_{1\beta}^{(\psi)\dagger}+T_{\alpha2}^{(\psi)}T_{2\beta}^{(\psi)\dagger}-T_{\alpha3}^{(\psi)}T_{3\beta}^{(\psi)\dagger}-T_{\alpha4}^{(\psi)}T_{4\beta}^{(\psi)\dagger})(T_{\beta1}^{(\psi)}T_{1\alpha}^{(\psi)\dagger}+T_{\beta2}^{(\psi)}T_{2\alpha}^{(\psi)\dagger}-T_{\beta3}^{(\psi)}T_{3\alpha}^{(\psi)\dagger}-T_{\beta4}^{(\psi)}T_{4\alpha}^{(\psi)\dagger})$$
$$\times m_\alpha^{(\psi)2}\left[-\ln^2 m_\beta^{(\psi)2}+2\ln m_\alpha^{(\psi)2}\ln m_\beta^{(\psi)2}-2\,\mathrm{Li}_2\left(1-\frac{m_\alpha^{(\psi)2}}{m_\beta^{(\psi)2}}\right)\right]$$
$$+2\sum_{\alpha=1}^{4}\sum_{\beta=1}^{2}[(V_{\beta1}^{(\psi)}T_{1\alpha}^{(\psi)\dagger}+V_{\beta2}^{(\psi)}T_{2\alpha}^{(\psi)\dagger})(T_{\alpha1}^{(\psi)}V_{1\beta}^{(\psi)\dagger}+T_{\alpha2}^{(\psi)}V_{2\beta}^{(\psi)\dagger})+(U_{\beta1}^{(\psi)}T_{3\alpha}^{(\psi)\dagger}+U_{\beta2}^{(\psi)}T_{4\alpha}^{(\psi)\dagger})(T_{\alpha3}^{(\psi)}U_{1\beta}^{(\psi)\dagger}+T_{\alpha4}^{(\psi)}U_{2\beta}^{(\psi)\dagger})]$$
$$\times\Bigg\{m_\alpha^{(\psi)2}\left[\ln^2 M_\beta^{(\psi)2}-\ln m_\alpha^{(\psi)2}\ln M_\beta^{(\psi)2}+\mathrm{Li}_2\left(1-\frac{m_\alpha^{(\psi)2}}{M_\beta^{(\psi)2}}\right)-\mathrm{Li}_2\left(1-\frac{M_\beta^{(\psi)2}}{m_\alpha^{(\psi)2}}\right)\right]$$
$$+M_\beta^{(\psi)2}\left[\ln^2 m_\alpha^{(\psi)2}-\ln m_\alpha^{(\psi)2}\ln M_\beta^{(\psi)2}+\mathrm{Li}_2\left(1-\frac{M_\beta^{(\psi)2}}{m_\alpha^{(\psi)2}}\right)+\mathrm{Li}_2\left(1-\frac{m_\alpha^{(\psi)2}}{M_\beta^{(\psi)2}}\right)\right]\Bigg\}. \tag{A13}$$

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

APPENDIX B: RENORMALIZATION GROUP EQUATIONS FOR THE MODEL

In this appendix, we list the RGEs for the model. The superpotential at the scale above Λ is

$$
\begin{aligned}
W=&\frac{\lambda}{\sqrt{2}}\epsilon_{\alpha\beta}\left\{Q_1^\alpha Q_3^\beta Y_1+Q_1^\alpha Q_4^\beta Y_2+Q_2^\alpha Q_3^\beta Y_3+Q_2^\alpha Q_4^\beta Y_4+\frac{1}{\sqrt{2}}(Q_1^\alpha Q_2^\beta-Q_3^\alpha Q_4^\beta)Y_5\right\}\\
&+\frac{\lambda_1}{\sqrt{2}}\epsilon_{\alpha\beta}(Q_1^\alpha Q_2^\beta+Q_3^\alpha Q_4^\beta)Z_1+k_{d1}d\overline{d}Z_1+k_{l1}l\overline{l}Z_1-f_1X^2Z_1+\frac{\lambda_2}{\sqrt{2}}\epsilon_{\alpha\beta}(Q_1^\alpha Q_2^\beta+Q_3^\alpha Q_4^\beta)Z_2+k_{d2}d\overline{d}Z_2+k_{l2}l\overline{l}Z_2-f_2X^2Z_2\\
&+\frac{\lambda_3}{\sqrt{2}}\epsilon_{\alpha\beta}(Q_1^\alpha Q_2^\beta+Q_3^\alpha Q_4^\beta)Z_3+k_{d3}d\overline{d}Z_3+k_{l3}l\overline{l}Z_3-f_3X^2Z_3+f_d d\overline{d'}X+f_{\overline{d}}d'\overline{d}X+f_l l\overline{l'}X+f_{\overline{l}}l'\overline{l}X.
\end{aligned}
\tag{B1}
$$

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

$$\mu\frac{d\lambda}{d\mu}=\frac{1}{16\pi^2}\frac{1}{2}\lambda(\lambda_1^2+\lambda_2^2+\lambda_3^2+7\lambda^2-6g^2),$$

$$\mu\frac{d\lambda_1}{d\mu}=\frac{1}{16\pi^2}\left[\frac{1}{2}\lambda_1(\lambda_1^2+\lambda_2^2+\lambda_3^2+5\lambda^2-6g^2)+\lambda_i(\lambda_i\lambda_1+3k_{di}k_{d1}+2k_{li}k_{l1}+2f_if_1)\right],$$

$$\mu\frac{dk_{d1}}{d\mu}=\frac{1}{16\pi^2}\left[k_{d1}\left(2k_{d1}^2+2k_{d2}^2+2k_{d3}^2+f_d^2+f_{\overline{d}}^2-\frac{16}{3}g_3^2-\frac{4}{15}g_1^2\right)+k_{di}(\lambda_i\lambda_1+3k_{di}k_{d1}+2k_{li}k_{l1}+2f_if_1)\right],$$

$$\mu\frac{dk_{l1}}{d\mu}=\frac{1}{16\pi^2}\left[k_{l1}\left(2k_{l1}^2+2k_{l2}^2+2k_{l3}^2+f_l^2+f_{\overline{l}}^2-3g_2^2-\frac{3}{5}g_1^2\right)+k_{li}(\lambda_i\lambda_1+3k_{di}k_{d1}+2k_{li}k_{l1}+2f_if_1)\right],$$

$$\mu\frac{df_1}{d\mu}=\frac{1}{16\pi^2}[2f_1(4f_1^2+4f_2^2+4f_3^2+3f_d^2+3f_{\overline{d}}^2+2f_l^2+2f_{\overline{l}}^2)+f_i(\lambda_i\lambda_1+3k_{di}k_{d1}+2k_{li}k_{l1}+2f_if_1)],$$

$$\mu\frac{d\lambda_2}{d\mu}=\frac{1}{16\pi^2}\left[\frac{1}{2}\lambda_2(\lambda_1^2+\lambda_2^2+\lambda_3^2+5\lambda^2-6g^2)+\lambda_i(\lambda_i\lambda_2+3k_{di}k_{d2}+2k_{li}k_{l2}+2f_if_2)\right],$$

$$\mu\frac{dk_{d2}}{d\mu}=\frac{1}{16\pi^2}\left[k_{d2}\left(2k_{d1}^2+2k_{d2}^2+2k_{d3}^2+f_d^2+f_{\overline{d}}^2-\frac{16}{3}g_3^2-\frac{4}{15}g_1^2\right)+k_{di}(\lambda_i\lambda_2+3k_{di}k_{d2}+2k_{li}k_{l2}+2f_if_2)\right],$$

$$\mu\frac{dk_{l2}}{d\mu}=\frac{1}{16\pi^2}\left[k_{l2}\left(2k_{l1}^2+2k_{l2}^2+2k_{l3}^2+f_l^2+f_{\overline{l}}^2-3g_2^2-\frac{3}{5}g_1^2\right)+k_{li}(\lambda_i\lambda_2+3k_{di}k_{d2}+2k_{li}k_{l2}+2f_if_2)\right],$$

$$\mu\frac{df_2}{d\mu}=\frac{1}{16\pi^2}[2f_2(4f_1^2+4f_2^2+4f_3^2+3f_d^2+3f_{\overline{d}}^2+2f_l^2+2f_{\overline{l}}^2)+f_i(\lambda_i\lambda_2+3k_{di}k_{d2}+2k_{li}k_{l2}+2f_if_2)],$$

$$\mu\frac{d\lambda_3}{d\mu}=\frac{1}{16\pi^2}\left[\frac{1}{2}\lambda_3(\lambda_1^2+\lambda_2^2+\lambda_3^2+5\lambda^2-6g^2)+\lambda_i(\lambda_i\lambda_3+3k_{di}k_{d3}+2k_{li}k_{l3}+2f_if_3)\right],$$

$$\mu\frac{dk_{d3}}{d\mu}=\frac{1}{16\pi^2}\left[k_{d3}\left(2k_{d1}^2+2k_{d2}^2+2k_{d3}^2+f_d^2+f_{\overline{d}}^2-\frac{16}{3}g_3^2-\frac{4}{15}g_1^2\right)+k_{di}(\lambda_i\lambda_3+3k_{di}k_{d3}+2k_{li}k_{l3}+2f_if_3)\right],$$

$$\mu\frac{dk_{l3}}{d\mu}=\frac{1}{16\pi^2}\left[k_{l3}\left(2k_{l1}^2+2k_{l2}^2+2k_{l3}^2+f_l^2+f_{\overline{l}}^2-3g_2^2-\frac{3}{5}g_1^2\right)+k_{li}(\lambda_i\lambda_3+3k_{di}k_{d3}+2k_{li}k_{l3}+2f_if_3)\right],$$

$$\mu\frac{df_3}{d\mu}=\frac{1}{16\pi^2}[2f_3(4f_1^2+4f_2^2+4f_3^2+3f_d^2+3f_{\bar{d}}^2+2f_l^2+2f_{\bar{l}}^2)+f_i(\lambda_i\lambda_3+3k_{di}k_{d3}+2k_{li}k_{l3}+2f_if_3)],$$

$$\mu\frac{df_d}{d\mu}=\frac{1}{16\pi^2}f_d\left(k_{d1}^2+k_{d2}^2+k_{d3}^2+4f_1^2+4f_2^2+4f_3^2+5f_d^2+3f_{\bar{d}}^2+2f_l^2+2f_{\bar{l}}^2-\frac{16}{3}g_3^2-\frac{4}{15}g_1^2\right),$$

$$\mu\frac{df_{\bar{d}}}{d\mu}=\frac{1}{16\pi^2}f_{\bar{d}}\left(k_{d1}^2+k_{d2}^2+k_{d3}^2+4f_1^2+4f_2^2+4f_3^2+3f_d^2+5f_{\bar{d}}^2+2f_l^2+2f_{\bar{l}}^2-\frac{16}{3}g_3^2-\frac{4}{15}g_1^2\right),$$

$$\mu\frac{df_l}{d\mu}=\frac{1}{16\pi^2}f_l\left(k_{l1}^2+k_{l2}^2+k_{l3}^2+4f_1^2+4f_2^2+4f_3^2+3f_d^2+3f_{\bar{d}}^2+4f_l^2+2f_{\bar{l}}^2-3g_2^2-\frac{3}{5}g_1^2\right),$$

$$\mu\frac{df_{\bar{l}}}{d\mu}=\frac{1}{16\pi^2}f_{\bar{l}}\left(k_{l1}^2+k_{l2}^2+k_{l3}^2+4f_1^2+4f_2^2+4f_3^2+3f_d^2+3f_{\bar{d}}^2+2f_l^2+4f_{\bar{l}}^2-3g_2^2-\frac{3}{5}g_1^2\right), \tag{B2}$$

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

$$\mu\frac{dg}{d\mu}=-\frac{4}{16\pi^2}g^3,\quad \mu\frac{dg_3}{d\mu}=-\frac{1}{16\pi^2}g_3^3,\quad \mu\frac{dg_2}{d\mu}=\frac{3}{16\pi^2}g_2^3,\quad \mu\frac{dg_1}{d\mu}=\frac{1}{16\pi^2}\frac{43}{5}g_1^3. \tag{B3}$$

[1] M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189**, 575 (1981); S. Dimopoulos and S. Raby, *ibid.* **B192**, 353 (1981); **B219**, 479 (1983); M. Dine and W. Fischler, Phys. Lett. **110B**, 227 (1982); Nucl. Phys. **B204**, 346 (1982); M. Dine and M. Srednicki, *ibid.* **B202**, 238 (1982); L. Alvarez-Gaumé, M. Claudson, and M. Wise, *ibid.* **B207**, 96 (1982); C. Nappi and B. Ovrut, Phys. Lett. **113B**, 175 (1982).

[2] For a review, M. Dine, hep-ph/9612389.

[3] I. Affleck, M. Dine, and N. Seiberg, Nucl. Phys. **B256**, 557 (1986).

[4] K. Intriligator, N. Seiberg, and S. H. Shenker, Phys. Lett. B **342**, 152 (1995).

[5] Izawa K.-I. and T. Yanagida, Prog. Theor. Phys. **95**, 829 (1996); K. Intriligator and S. Thomas, Nucl. Phys. **B473**, 121 (1996).

[6] See also C. Csaki, L. Randall, and W. Skiba, Nucl. Phys. **B479**, 65 (1996); R. G. Leigh, L. Randall, and R. Rattazzi, *ibid.* **B501**, 375 (1997).

[7] M. Dine and A. E. Nelson, Phys. Rev. D **48**, 1277 (1993); M. Dine, A. E. Nelson, and Y. Shirman, *ibid.* **51**, 1362 (1995); M. Dine, A. E. Nelson, Y. Nir, and Y. Shirman, *ibid.* **53**, 2658 (1996).

[8] T. Hotta, Izawa K.-I., and T. Yanagida, Phys. Rev. D **55**, 415 (1997).

[9] E. Poppitz and S. Trivedi, Phys. Rev. D **55**, 5508 (1997); L. Randall, Nucl. Phys. **B495**, 37 (1997); N. Haba, N. Maru, and T. Matsuoka, *ibid.* **B497**, 31 (1997); Phys. Rev. D **56**, 4207 (1997); N. Arkani-Hamed, J. March-Russell, and H. Murayama, Nucl. Phys. **B509**, 3 (1998); Izawa K.-I., Prog. Theor. Phys. **98**, 443 (1997); H. Murayama, Phys. Rev. Lett. **79**, 18 (1997).

[10] M. A. Luty, Phys. Rev. D **57**, 1531 (1998); A. G. Cohen, D. B. Kaplan, and A. E. Nelson, Phys. Lett. B **412**, 301 (1997).

[11] Y. Nomura, K. Tobe, and T. Yanagida (unpublished).

[12] H. Pagels and J. R. Primack, Phys. Rev. Lett. **48**, 223 (1982).

[13] A. Gouvêa, T. Moroi, and H. Murayama, Phys. Rev. D **56**, 1281 (1997).

[14] T. Moroi, H. Murayama, and M. Yamaguchi, Phys. Lett. B **303**, 289 (1993).

[15] J. Hashiba, M. Kawasaki, and T. Yanagida, Phys. Rev. Lett. **79**, 4525 (1997).

[16] Izawa K.-I., Y. Nomura, K. Tobe, and T. Yanagida, Phys. Rev. D **56**, 2886 (1997).

[17] K. S. Babu, C. Kolda, and F. Wilczek, Phys. Rev. Lett. **77**, 3070 (1996); S. Dimopoulos, S. Thomas, and J. D. Wells, Nucl. Phys. **B488**, 39 (1997); J. A. Bagger, K. T. Matchev, D. M. Pierce, and R. Zhang, Phys. Rev. D **55**, 3188 (1997); D. A. Dicus, B. Dutta, and S. Nandi, Phys. Rev. Lett. **78**, 3055 (1997).

[18] S. Dimopoulos, M. Dine, S. Raby, and S. Thomas, Phys. Rev. Lett. **76**, 3494 (1996); S. Ambrosanio, G. L. Kane, G. D. Kribs, and S. P. Martin, *ibid.* **76**, 3498 (1996); Phys. Rev. D **54**, 5395 (1996); S. Dimopoulos, S. Thomas, and J. D. Wells, *ibid.* **54**, 3283 (1996); J. Lopez, D. Nanopoulos, and A. Zichichi, Phys. Rev. Lett. **77**, 5168 (1996); Phys. Rev. D **55**, 5813 (1997); H. Baer, M. Brhlik, C. Chen, and X. Tata, *ibid.* **55**, 4463 (1997); A. Ghosal, A. Kundu, and B. Mukhopadhyaya, *ibid.* **56**, 504 (1997); S. Ambrosanio, G. D. Kribs, and S. P.

Martin, *ibid.* **56**, 1761 (1997); D. A. Dicus, B. Dutta, and S. Nandi, *ibid.* **56**, 5748 (1997); A. Datta, A. Datta, A. Kundu, B. Mukhopadhyaya, and S. Roy, Phys. Lett. B **416**, 117 (1998); C. H. Chen and J. F. Gunion, *ibid.* **420**, 77 (1998); J. Kim, J. L. Lopez, D. V. Nanopoulos, R. Rangarajan, and A. Zichichi, Phys. Rev. D **57**, 554 (1998).

[19] K. Maki and S. Orito, Phys. Rev. D **57**, 554 (1998).

[20] For a review, K. Intriligator and N. Seiberg, Nucl. Phys. B (Proc. Suppl.) **45BC**, 1 (1996).

[21] See also Y. Shirman, Phys. Lett. B **389**, 287 (1996); N. Arkani-Hamed and H. Murayama, Phys. Rev. D **57**, 6638 (1998).

[22] J. Bagger, E. Poppitz, and L. Randall, Nucl. Phys. **B426**, 3 (1994).

[23] M. Dine, Y. Nir, and Y. Shirman, Phys. Rev. D **55**, 1501 (1997).

[24] S. Dimopoulos, G. F. Giudice, and A. Pomarol, Phys. Lett. B **389**, 37 (1996).

[25] S. P. Martin, Phys. Rev. D **55**, 3177 (1997).

[26] S. Dimopoulos and G. F. Giudice, Phys. Lett. B **393**, 72 (1997).

[27] J. Hisano, H. Murayama, and T. Yanagida, Nucl. Phys. **B402**, 46 (1993); J. Hisano, T. Moroi, K. Tobe, and T. Yanagida, Mod. Phys. Lett. A **10**, 2267 (1995).

[28] CDF Collaboration, L. Montanet *et al.*, Phys. Rev. D **56**, 1357 (1997).

[29] OPAL Collaboration, K. Ackerstaff *et al.*, Eur. Phys. J. C **2**, 213 (1998); L3 Collaboration, M. Acciarri *et al.*, Report No. CERN-PPE-97-130; ALEPH Collaboration, R. Barate *et al.*, Eur. Phys. J. C **2**, 417 (1998).

[30] R. Rattazzi and U. Sarid, Nucl. Phys. **B501**, 297 (1997).

[31] S. Coleman and E. Weinberg, Phys. Rev. D **7**, 1888 (1973).

[32] S. Dimopoulos, G. Dvali, and R. Rattazzi, Phys. Lett. B **413**, 336 (1997).

[33] G. Dvali, G. F. Giudice, and A. Pomarol, Nucl. Phys. **B478**, 31 (1996).

[34] T. Yanagida, Phys. Lett. B **400**, 109 (1997).

[35] P. Fayet, Phys. Lett. **84B**, 416 (1979); Phys. Rep. **105**, 21 (1984).

[36] J. Hisano, K. Tobe, and T. Yanagida, Phys. Rev. D **55**, 411 (1997).

[37] J. Ellis, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B **394**, 354 (1997); D0 Collaboration, B. Abbott *et al.*, Phys. Rev. Lett. **80**, 442 (1998).