

Dynamical supersymmetry breaking without messenger gauge interactions

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(Received 3 June 1996)

We investigate low-energy models of supersymmetry (SUSY) breaking by means of vectorlike gauge theories for dynamical SUSY breaking. It is not necessary to introduce messenger gauge interactions utilized so far to mediate the SUSY breaking to the standard-model sector, which reduces complications in the model building. We also consider various other ways of SUSY-breaking transmission. [S0556-2821(97)01201-0]

PACS number(s): 11.30.Pb, 11.30.Qc, 12.60.Jv

I. INTRODUCTION

Supersymmetry (SUSY) is a plausible candidate as an origin of the hierarchy between the weak scale and higher scales, say the grand unified theory (GUT) and/or the gravitational scales. It must be broken at low energy so that the observed particles do not have their superpartners with the same masses. We consider dynamical SUSY breaking (DSB) as the mechanism of SUSY breaking at low energy since DSB gives a complete solution to the hierarchy problem—it can induce a very tiny scale compared to the gravitational scale without any fine tunings of parameters in the theory. In particular, low-energy models of DSB (so-called visible sector models) are attractive also in that they can give a natural solution to the flavor-changing neutral current (FCNC) problem through the mass degeneracy of the squarks and sleptons [1].

There are two methods to mediate the low-energy SUSY breaking in the DSB sector to the SUSY standard-model sector. One is to identify a global symmetry in the DSB sector as gauge groups in the standard model.¹ This approach, however, has difficulty that the $SU(3)_C$ gauge coupling constant tends to blow up at too low an energy scale. The other is to introduce a messenger sector which couples to both the sectors.

The phenomenological models with low-energy DSB proposed so far [1] share the common structures: DSB is realized in chiral gauge theories and SUSY-breaking effects are mediated to the SUSY standard-model sector by an additional messenger gauge interaction which is introduced only for that purpose. The additional gauge interaction is needed since the DSB and the standard-model sectors are both chiral.

In this paper we investigate a model with a vectorlike gauge theory as the SUSY-breaking sector. We adopt vectorlike models of DSB proposed recently in Ref. [2]. We see that it is not necessary to introduce the messenger gauge interaction, whose coupling threatens to blow up. Note that this may largely reduce complication in the model building.

¹We note that this is easily realized in the $SP(N)$ vectorlike gauge theories of DSB proposed in Ref. [2], since the anomaly-free condition for the standard-model gauge interactions is manifestly satisfied without introducing additional fields.

We also consider various other ways of SUSY-breaking transmission.

II. THE SUSY-BREAKING SECTOR

Let us consider a SUSY $SU(2)$ gauge theory with four doublet chiral superfields Q_i and six singlet ones $Z^{ij} = -Z^{ji}$. Here, i and j denote the flavor indices ($i, j = 1, \dots, 4$). Without a superpotential, this model has a flavor $SU(4)_F$ symmetry.

The tree-level superpotential of the model [2] is given by

$$W_{\text{tree}} = \lambda_{ij}^{kl} Z^{ij} Q_k Q_l, \quad (1)$$

where λ_{ij}^{kl} denote generic coupling constants with $\lambda_{ij}^{kl} = -\lambda_{ji}^{kl} = -\lambda_{ij}^{lk}$. SUSY remains unbroken perturbatively in this model.

The exact effective superpotential of the model, which takes into account the full nonperturbative effects, may be written in terms of gauge-invariant low-energy degrees of freedom

$$V_{ij} = -V_{ji} \sim Q_i Q_j \quad (2)$$

as

$$W_{\text{eff}} = X(\text{Pf} V_{ij} - \Lambda^4) + \lambda_{ij}^{kl} Z^{ij} V_{kl}, \quad (3)$$

where X is an additional chiral superfield, $\text{Pf} V_{ij}$ denotes the Pfaffian of the antisymmetric matrix V_{ij} , and Λ is a dynamical scale of the $SU(2)$ gauge interaction. This is none other than a superpotential of O'Raifeartaigh-type. Namely, this effective superpotential yields conditions for SUSY vacua

$$\text{Pf} V_{ij} = \Lambda^4, \quad \lambda_{ij}^{kl} V_{kl} = 0, \quad (4)$$

which are incompatible as long as $\Lambda \neq 0$. Therefore, we conclude that SUSY is dynamically broken in this model [2].

Let us impose a flavor $SP(4)_F$ symmetry on the above model to make our analysis simpler, where we adopt the notation $SP(4)_F \subset SU(4)_F$. Then, the effective superpotential can be written as

$$W_{\text{eff}} = X(V^2 + V_a V_a - \Lambda^4) + \lambda_Z Z V + \lambda_Z^a V_a, \quad (5)$$

where V and Z are singlets and V_a and Z^a are five-dimensional representations of $\text{SP}(4)_F$, respectively, in V_{ij} and Z^{ij} , which constitute six-dimensional representations of $\text{SU}(4)_F$. Here, $a=1, \dots, 5$ and λ_Z and λ denote coupling constants which are taken to be positive.

When the coupling λ_Z is small, the effective superpotential W_{eff} implies that we obtain the following vacuum expectation values:

$$\langle V \rangle \approx \Lambda^2, \quad \langle V_a \rangle \approx 0. \quad (6)$$

Then, the low-energy effective superpotential may be approximated by²

$$W_{\text{eff}} \approx \lambda_Z \Lambda^2 Z. \quad (7)$$

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

$$K = ZZ^* - \frac{\eta}{4\Lambda^2} \lambda_Z^4 (ZZ^*)^2 + \dots, \quad (8)$$

where η is a real constant of order one.

Then, the effective potential of the scalar Z (with the same notation as the superfield) is given by

$$V_{\text{eff}} \approx \lambda_Z^2 \Lambda^4 \left(1 + \frac{\eta}{\Lambda^2} \lambda_Z^4 ZZ^* + \dots \right). \quad (9)$$

When $\eta > 0$, this leads to $\langle Z \rangle = 0$. Otherwise, we suspect that $\langle Z \rangle$ is of order $\lambda_Z^{-1} \Lambda$ since the effective potential is expected to be lifted in the region far from the origin ($Z=0$), as is the case for the O’Raifeartaigh model [3]. In any event, the DSB scale is given by $F_Z \approx \lambda_Z \Lambda^2$, where F_Z denotes the F component of the superfield Z , and the mass m_Z of the scalar Z is seen to be of order $\sqrt{\eta} \lambda_Z^3 \Lambda$ except for the potential massless R axion.

We, henceforth, assume $\eta > 0$ for definiteness.

III. THE MESSENGER SECTOR

The messenger sector consists of chiral superfields $Y, d, \bar{d}, l, \bar{l}$, which are all singlets under the strong $\text{SU}(2)$ and the global $\text{SP}(4)_F$. Y is also a singlet under the standard-model gauge group. As for $d, \bar{d}$ and $l, \bar{l}$ we tentatively assume that they transform as the down quark, the antilepton doublet, and their antiparticles, respectively. The interactions between the DSB and the messenger sectors are described by a superpotential

$$W_{\text{int}} = \frac{\lambda_Y}{2} \epsilon_{\alpha\beta} (Q_1^\alpha Q_2^\beta + Q_3^\alpha Q_4^\beta) Y - \frac{f}{3} Y^3 + (k_1 \bar{d} \bar{d} + k_2 \bar{l} \bar{l}) Y, \quad (10)$$

where α and β denote $\text{SU}(2)$ gauge indices and the couplings λ_Y , f , k_1 , and k_2 are taken to be positive. In the following, we set $k \equiv k_1 = k_2$ for simplicity, which approximately holds in SUSY GUT’s.

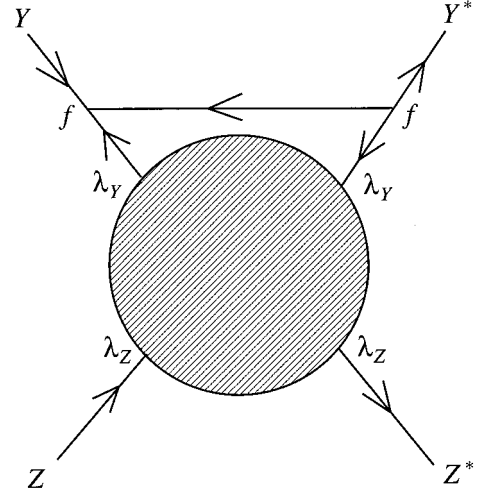


FIG. 1. Typical Feynman diagrams generating f -dependent corrections to the Kähler potential.

In view of Eqs. (6) and (7), the full effective superpotential of the DSB and messenger sectors is obtained as

$$W_{\text{eff}} \approx \lambda_Z \Lambda^2 Z + \lambda_Y \Lambda^2 Y - \frac{f}{3} Y^3 + k(\bar{d} \bar{d} + \bar{l} \bar{l}) Y, \quad (11)$$

where we have used

$$V \sim \frac{1}{2} \epsilon_{\alpha\beta} (Q_1^\alpha Q_2^\beta + Q_3^\alpha Q_4^\beta). \quad (12)$$

The Kähler potential of Z and Y is expected to have the form

$$K = ZZ^* + YY^* - \frac{\eta}{4\Lambda^2} |\lambda_Z Z + \lambda_Y Y|^4 - \frac{\delta}{4\Lambda^2} \frac{\lambda_Y^2}{16\pi^2} |\lambda_Z Z + \lambda_Y Y|^2 \{ 4f^2 YY^* + k^2 (\bar{d} \bar{d}^* + \bar{d} \bar{d}^* + \bar{l} \bar{l}^* + \bar{l} \bar{l}^*) \} + \dots, \quad (13)$$

where δ is a real constant of order one. We notice that nonanomalous R and discrete symmetries with the coupling f as an external field³ may be utilized to see that there are no trilinear terms such as ZYY^* , Zdd^* , $Z\bar{d}\bar{d}^*$, Zll^* , and $Zl\bar{l}^*$ in the above Kähler potential. As we will see below, we obtain vanishing F component of Y , $\langle F_Y \rangle = 0$, in a limit $f \rightarrow 0$. In that case, the SUSY breaking is not transmitted to the standard-model sector by means of the singlet Y . Typical Feynman diagrams generating f -dependent corrections to the Kähler potential in Eq. (13) are shown in Fig. 1.

We obtain an effective potential from Eqs. (11) and (13):

²The fields Z^a and V_a form massive multiplets with mass of order $\lambda \Lambda$ and they are integrated out.

³The R charges of Q, Z, Y, f read $0, 2, 2, -4$, respectively. The discrete symmetry amounts to a transformation $Q \rightarrow iQ$, $Z \rightarrow -Z$, $Y \rightarrow -Y$, $f \rightarrow -f$.

$$\begin{aligned}
V_{\text{eff}} \approx & \lambda_Z^2 \Lambda^4 + |\lambda_Y \Lambda^2 - f Y^2 + k(\overline{d}d + \overline{l}l)|^2 + |k\overline{d}Y|^2 + |k\overline{l}Y|^2 \\
& + |k\overline{l}Y|^2 + |k\overline{l}Y|^2 + \eta \lambda_Z^6 \Lambda^2 Z Z^* \\
& + \eta \lambda_Z^5 \lambda_Y \Lambda^2 (Z Y^* + Z^* Y) + \eta \lambda_Z^4 \lambda_Y^2 \Lambda^2 Y Y^* \\
& + \delta \frac{f^2}{64\pi^2} \lambda_Z^4 \Lambda^2 \{4\lambda_Y^2 Y Y^* + k^2(d\overline{d}^* + \overline{d}d^* + l\overline{l}^* \\
& + \overline{l}l^*)\}, \tag{14}
\end{aligned}$$

where we have taken into account only the leading corrections in the coupling λ_Y .

We restrict ourselves to the vacua with $\langle d \rangle = \langle \overline{d} \rangle = \langle l \rangle = \langle \overline{l} \rangle = 0$ in the following consideration since, otherwise, the standard-model gauge group would be broken by their vacuum expectation values. Notice that the above vacuum is realized in a certain range of the coupling constants λ_Y and k . Then, the potential equation (14) yields a vacuum given by

$$\langle F_Y \rangle = \langle \lambda_Y \Lambda^2 - f Y^2 \rangle = \delta \frac{f^2}{16\pi^2} \frac{\lambda_Z^4 \lambda_Y^2}{2f} \Lambda^2, \quad \langle Y^2 \rangle = \frac{\lambda_Y}{f} \Lambda^2. \tag{15}$$

The messenger quarks and leptons have masses $m_{d,l}$ of order Λ in this vacuum:

$$m_{d,l} \approx k \langle Y \rangle = k \sqrt{\frac{\lambda_Y}{f}} \Lambda. \tag{16}$$

Since $\langle F_Y \rangle$ and $\langle Y \rangle$ are nonvanishing, the SUSY breaking is mediated to the standard-model sector as seen in the next section.

IV. THE STANDARD-MODEL SECTOR

The standard-model gauginos obtain their masses from radiative corrections through loops of the messenger quarks and leptons [1]:

$$\begin{aligned}
m_{\tilde{g}_i} & \approx \frac{g_i^2}{16\pi^2} k^2 \frac{\langle Y \rangle \langle F_Y \rangle}{m_{d,l}^2} \approx \frac{\alpha_i}{4\pi} \frac{\langle F_Y \rangle}{\langle Y \rangle} \\
& \approx \frac{\alpha_i}{4\pi} \delta \frac{f^2}{16\pi^2} \frac{\lambda_Z^4 \lambda_Y^{3/2}}{2f^{1/2}} \Lambda, \tag{17}
\end{aligned}$$

where $\alpha_i = g_i^2/4\pi$ denote the standard-model gauge couplings.

The masses of the squarks and sleptons are given by [1]

$$\tilde{m}^2 \approx 2 \left(C_3 m_{\text{gluino}}^2 + C_2 m_{W\text{-ino}}^2 + \frac{3}{5} C_1 m_{B\text{-ino}}^2 \right), \tag{18}$$

where $C_3 = 0$ or $4/3$ and $C_2 = 0$ or $3/4$, whose nonvanishing values are the quadratic Casimir invariants for the fundamental representations of $SU(3)_C$ and $SU(2)_L$, respectively, and C_1 denotes the corresponding hypercharges squared. Namely, the squarks and sleptons have the masses of the same order as the gaugino masses. Notice here that the squarks and sleptons with the same quantum numbers of the

standard-model gauge group degenerate in mass, which leads to the suppression of FCNC as noted in the Introduction.

We now proceed to consider phenomenological constraints on the parameters in the potential equation (14).

First, the gluino mass should be $(10^2 - 10^3)$ GeV to maintain the weak scale of order 10^2 GeV. This is because growth of the gluino mass increases the top squark mass [see Eq. (18)] and the large top squark mass raises the Higgs boson masses which determine the weak scale. Thus, we demand

$$m_{\text{gluino}} \approx \frac{\alpha_3}{4\pi} \delta \frac{f^2}{16\pi^2} \frac{\lambda_Z^4 \lambda_Y^{3/2}}{2f^{1/2}} \Lambda \approx (10^2 - 10^3) \text{ GeV}. \tag{19}$$

Second, the gravitino mass $m_{3/2}$ should be less than about 1 GeV to keep the mass degeneracy of the squarks and sleptons explaining the required suppression of FCNC [4]:⁴

$$m_{3/2} \approx \frac{\langle F_Z \rangle}{\sqrt{3}M} \approx \frac{\lambda_Z \Lambda^2}{\sqrt{3}M} \leq 1 \text{ GeV}, \tag{20}$$

where M is the gravitational scale, that is, $M \approx 2 \times 10^{18}$ GeV.

In terms of the parameters, $\lambda_Z \approx f \approx 1$, for instance, we get

$$\lambda_Y^{3/2} \Lambda \approx (10^6 - 10^7) \text{ GeV} \tag{21}$$

from Eq. (19). By means of Eq. (20), this gives

$$\lambda_Y \geq 10^{-2}, \quad \Lambda \leq 10^9 \text{ GeV}. \tag{22}$$

For $\lambda_Y \leq 1$ we get

$$10^{-2} \leq \lambda_Y \leq 1, \quad 10^6 \text{ GeV} \leq \Lambda \leq 10^9 \text{ GeV}. \tag{23}$$

The gravitino mass is given by

$$m_{3/2} \approx 1 \text{ keV} - 1 \text{ GeV}. \tag{24}$$

In the case of $m_{3/2} = (1 - 100)$ keV, the reheating temperature T_R after inflation should be very low as $T_R \leq 10^2$ GeV to avoid the overclosure of the Universe [5]. Thus, we must invoke a late-time baryogenesis such as in the Affleck-Dine mechanism [6].

V. CONCLUSION

In this final section, we deal with a few remaining aspects of our model.

First, we consider the Polonyi problem on the singlet Z . The mass of Z is given by

$$m_Z \approx \sqrt{\eta} \lambda_Z^3 \Lambda \approx (10^6 - 10^9) \text{ GeV}. \tag{25}$$

This shows that the singlet causes no cosmological problem [7].

Second, we argue a naturalness problem in the messenger sector. The relevant superpotential is given by

⁴We assume that there exists an additional sector to set the cosmological constant vanishing. Then, the squarks and sleptons in the SUSY standard model acquire soft SUSY-breaking masses of order $m_{3/2}$ due to supergravity effects.

$$W = \lambda_Z(QQ)Z + \lambda_Y(QQ)Y - \frac{f}{3}Y^3 + (k_1\overline{d}d + k_2\overline{l}l)Y, \quad (26)$$

where (QQ) denotes the right-hand side in Eq. (12). We can consistently impose $U(1)_R$ symmetry⁵ with the charges of $Q_i, Y, Z, d, \overline{d}, l, \overline{l}$ as $2/3$. However, this R symmetry allows additional terms such as $f_Z Z^3$. In order for our SUSY-breaking vacuum to be (meta)stable, the coupling f_Z should be extremely small.⁶

Third, we comment on the μ problem. Let us consider an interaction $hYHH$ as a direct source for the μ term μHH in the superpotential:

$$\mu = h\langle Y \rangle. \quad (27)$$

Then, we also obtain a soft mass term $\overline{B}HH$ in the potential with

$$B = h\langle F_Y \rangle. \quad (28)$$

Together with Eq. (17), these yield

$$\mu m_{\text{gluino}} \simeq \frac{\alpha_3}{4\pi} B, \quad (29)$$

which implies that the value B is too large. Thus, we should give up the direct generation of the μ term. When we obtain

⁵This R symmetry is anomalous and there is no light R axion.

⁶This may be natural in the sense of 't Hooft under a symmetry which imposes $\lambda_Y = f_Z = 0$. The fact that $\lambda_Y \gg |f_Z|$ may originate from the charge difference between the two terms $\lambda_Y(QQ)Y$ and $f_Z Z^3$.

an appropriate μ term by one way or other [1], the $SU(2)_L \times U(1)_Y$ breaking may be induced by radiative corrections.

We note that even the messenger singlet Y may be unnecessary if an appropriate mass terms for the messenger quarks and leptons $d, \overline{d}, l, \overline{l}$ are generated such as the μ term. Then, the SUSY breaking may be mediated by introducing a messenger gauge interaction for the messenger quarks and leptons as well as Q_i .

We note that many extensions of the above models are possible by introducing various singlets to realize DSB. In fact, generic O'Raifeartaigh modes may be converted into dynamical ones by replacing the SUSY-breaking scales with the condensed operators (QQ) .

Finally, we comment on hidden sector models of DSB with vectorlike gauge theories. Since SUSY breaking is communicated by gravity in hidden sector models, we need to introduce no messenger sector (no fields such as $Y, d, \overline{d}, l, \overline{l}$), which enables us to respect naturalness thoroughly. Sizeable gaugino masses stem from terms of the form $(Z/M)W_\alpha W^\alpha$, which is compatible with $U(1)_R$ -charge assignments of Z and (QQ) as zero and two, respectively. Then, we get an effective superpotential

$$W_{\text{eff}} = \lambda_Z \Lambda^2 Z \left[1 + O\left(\frac{Z}{M}\right) \right]. \quad (30)$$

We expect a SUSY-breaking local minimum at $|Z/M| \ll 1$ for the effective potential corresponding to Eq. (30) in supergravity [8] in view of Sec. II. When $\Lambda \simeq 10^{11}$ GeV and $\lambda_Z \simeq 1$, the gravitino mass turns out to be $m_{3/2} \simeq 10^3$ GeV [see Eq. (20)], which characterizes the SUSY-breaking scale in the standard-model sector. Note that $m_Z \simeq 10^{11}$ GeV in this case so that the singlet Z causes no cosmological problem. However, in contrast with the visible sector models, the FCNC problem is not automatically resolved in the hidden sector models.

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