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N. Kitazawa – Tokyo Metropolitan University

N. Okada – Tokyo Metropolitan University

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Supersymmetric composite model with dynamical supersymmetry breaking

Noriaki Kitazawa* and Nobuchika Okada†

Department of Physics, Tokyo Metropolitan University, Hachioji-shi, Tokyo 192-03, Japan

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We present a supersymmetric composite model with dynamical supersymmetry breaking. The model is based on the gauge group $SU(2)_S \times SU(2)_H \times SU(3)_c \times SU(2)_L \times U(1)_Y$. Supersymmetry is dynamically broken by the nonperturbative effect of the $SU(2)_S$ ‘‘supercolor’’ interaction. The large top quark Yukawa coupling is naturally generated by the $SU(2)_H$ ‘‘hypercolor’’ interaction as recently proposed by Nelson and Strassler. The supersymmetry breaking is mediated to the standard model sector by a new mechanism. The electroweak symmetry breaking is caused by the radiative correction due to the large top quark Yukawa coupling with supersymmetry breaking. This is the ‘‘radiative breaking scenario,’’ which originates from the dynamics of the supercolor and hypercolor gauge interactions. [S0556-2821(97)02417-X]

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I. INTRODUCTION

In most of the supersymmetric models, if supersymmetry is broken, the electroweak symmetry breaking is caused by the radiative correction due to the large top quark Yukawa coupling. This mechanism is known as the ‘‘radiative breaking scenario’’ [1]. However, to understand the scenario completely, we should investigate two subjects. One is a mechanism of supersymmetry breaking, and the other is a natural explanation of the large top quark Yukawa coupling.

The method proposed by Seiberg and co-workers [2] is remarkable, when we consider the dynamical supersymmetry breaking in ($N=1$) supersymmetric gauge theories. By the method, we can evaluate the nonperturbative effect of a strong gauge interaction and exactly determine the dynamically generated superpotential. There are many models [2–5] in which the supersymmetry is broken by the nonperturbative effect of the strong gauge interaction.

On the other hand, the supersymmetric composite model recently proposed by Nelson and Strassler [6] is remarkable to understand the large top quark Yukawa coupling. They introduce the $SU(2)$ gauge interaction with six doublet superfields, the ‘‘preon’’ superfields. There is $SU(6) \times U(1)_R$ global symmetry. The $SU(5)$ subgroup is gauged and identified with $SU(5) \supset G_{SM}$, where $G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$ is the gauge group of the standard model. Then, the five doublet superfields transform as **5** of the $SU(5)$, and the last one is singlet. Below the dynamical scale of the $SU(2)$ gauge interaction, the preon superfields are confined into the superfields of **5** and **10** of $SU(5)$, which are identified with the up-type Higgs boson and the ten-dimensional superfields in the conventional supersymmetric $SU(5)$ GUT framework. The top quark Yukawa coupling is dynamically generated by the nonperturbative effect of the $SU(2)$ gauge interaction, and it can be understood as the exchange of the preon superfields in a microscopical point of view. Therefore, the top quark Yukawa coupling is naturally expected as order 1.

In this paper, we present a supersymmetric composite

model with dynamical supersymmetry breaking. The model is based on the gauge group $SU(2)_S \times SU(2)_H \times G_{SM}$. We assume that dynamical scales of these gauge interactions are ordered as $\Lambda_S \gg \Lambda_H \gg E_{EW}$, where Λ_S and Λ_H are the dynamical scale of the $SU(2)_S$ and $SU(2)_H$ gauge interactions, respectively, and E_{EW} is the scale of the electroweak symmetry breaking. It is useful to separate the model into three building blocks by these energy scales, namely, the dynamical supersymmetry-breaking sector, the preon sector proposed by Nelson and Strassler, and the standard model sector.

The dynamical supersymmetry-breaking sector is based on the $SU(2)_S$ ‘‘supercolor’’ gauge group with four doublet and eight singlet superfields. There is $Sp(2) \times U(1)_R$ global symmetry. The $SU(2)$ subgroup is gauged, which is identified with $SU(2)_H$. Therefore, there are some $SU(2)_H$ doublet superfields in this sector. The supersymmetry is dynamically broken at the scale Λ_S by the nonperturbative effect of the supercolor interaction. In addition, the $U(1)_R$ symmetry is broken by the effective Kähler potential generated by the Yukawa coupling in the superpotential. Both the scalar and fermion fields in the $SU(2)_H$ doublet superfields get their masses due to the breaking of the supersymmetry and $U(1)_R$ symmetry.

The preon sector is based on the $SU(2)_H$ ‘‘hypercolor’’ gauge group with six doublet preon superfields as proposed by Nelson and Strassler. The large top quark Yukawa coupling is naturally generated by the nonperturbative effect of the hypercolor gauge interaction. Since we focus on the relation among the dynamical supersymmetry breaking, the large top quark Yukawa coupling, and the electroweak symmetry breaking, we concentrate only on the third generation. Therefore, our model is a ‘‘toy model,’’ but it is possible to include the last two generations as the elementary particles.

The supersymmetry breaking is mediated to the preon sector and the standard model sector through the $SU(2)_H \times G_{SM}$ gauge interaction. The hypercolor gaugino and the scalar preons get the soft breaking masses through radiative corrections by the $SU(2)_H$ doublet fields in the dynamical supersymmetry-breaking sector. Here, the doublet fields play a role of the ‘‘messenger fields’’ [3]. Scalar fields composed of scalar preons in the standard model get masses,

*Electronic address: kitazawa@phys.metro-u.ac.jp

†Electronic address: n-okada@phys.metro-u.ac.jp

because of the masses of the scalar preons. Once the hypercolor gaugino and scalar preons become massive, the standard model gauginos get masses through the radiative correction, since the preon superfields have both charges of the $SU(2)_H$ and G_{SM} . Here, the preon and scalar preon fields, which are the $SU(2)_H$ doublet fields in the preon sector, also play a role of the ‘‘messenger fields.’’ This is a new mechanism to give the standard model gauginos the soft breaking masses. At low energy where the preon superfields are confined, the mechanism can be understood as the mixing between the gauginos and the composite fermions in an adjoint representation.

The electroweak symmetry is broken by the effect of the large top quark Yukawa coupling as same as the radiative breaking scenario. It must be noticed that the breaking is triggered by the dynamics of the $SU(2)_S \times SU(2)_H$ gauge interaction, since both the supersymmetry breaking and large Yukawa coupling owe their origin to the dynamics.

In Sec. II, we introduce the preon sector which is recently proposed by Nelson and Strassler. In Sec. III, the dynamical supersymmetry-breaking sector is constructed. In Sec. IV, we discuss the mediation of the supersymmetry breaking to low energy. In Sec. V, the standard model sector is discussed with an emphasis on the spectrum of the mediated soft breaking masses. In Sec. VI, we show that the electroweak symmetry breaking really occurs in our model by including the effect of the large top quark Yukawa coupling. In Sec. VII, the proton decay caused by the colored Higgs boson in our model is discussed, and we show that there is no contradiction with the experiment. In Sec. VIII, we briefly discuss the R -axion problem in our model and give an example of the solution. In Sec. IX, we summarize our model and comment on the extension of the model to include the first and second generations.

II. PREON SECTOR

In this section we construct the preon sector which is based on the model recently proposed by Nelson and Strassler [6]. Only the third generation is considered as mentioned above. We introduce the $SU(2)_H$ hypercolor gauge group with the six doublet preon superfields. The maximal nonanomalous global symmetry of this sector is $SU(6) \times U(1)_R$, and the subgroup $SU(5) \supset G_{SM}$ is gauged, where $G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$ is nothing but the gauge group of the standard model. The preon superfields P and N transform under $SU(2)_H \times SU(5)$ as

	$SU(2)_H$	$SU(5)$
P	2	5
N	2	1

It is convenient to decompose the preon superfield P as $P = (dh)$, where d and h transform as $(\mathbf{3}, \mathbf{1}, -1/3)$ and $(\mathbf{1}, \mathbf{2}, 1/2)$, respectively, under the standard model gauge group G_{SM} .

In addition, we introduce two elementary superfields Φ_1 and Φ_2 , both of which are in $\mathbf{5}^*$ representation of the $SU(5)$. The field $\Phi_1 = (\bar{b}\bar{\ell})$ represents the superfields of the antibottom quark and the lepton doublet $\bar{\ell} = (\bar{\tau} - \bar{\nu}_\tau)$, and Φ_2 is

corresponds to the down-type Higgs superfield $\Phi_2 = (\bar{D}\bar{H})$ with $\bar{D} = (\mathbf{3}^*, \mathbf{1}, 1/3)$ and $\bar{H} = (\mathbf{1}, \mathbf{2}, -1/2)$.

At high energy where the hypercolor interaction is weak, the superpotential is generally given by (including dimension-4 operators)

$$W = \eta \bar{H}[hN] + \eta_D \bar{D}[dN] + \frac{\kappa_b}{M_b} \bar{H}[dh]\bar{b} + \frac{\kappa_\ell}{M_\ell} \bar{H}[hh]\bar{\ell}, \quad (1)$$

where square brackets denote the contraction of the $SU(2)_H$ index by the ϵ tensor, and M_b and M_ℓ are parameters with dimension 1. Although the third and fourth terms in Eq. (1) are needed to generate the bottom and τ Yukawa couplings, we neglect these terms in the following, and consider only the top quark Yukawa coupling.

Below the scale Λ_H where the $SU(2)_H$ interaction becomes strong, this theory is described by the massless $SU(2)_H$ singlet effective superfields M_{ij} , where i and j are flavor indices. The effective superfield is the 15-component antisymmetric tensor given by

$$M = \epsilon^{\alpha\beta} \begin{pmatrix} d \\ h \\ N \end{pmatrix}_\alpha (dhN)_\beta = \begin{pmatrix} [dd] & [dh] & [dN] \\ [hd] & [hh] & [hN] \\ [Nd] & [Nh] & 0 \end{pmatrix}, \quad (2)$$

where α and β are the $SU(2)_H$ indices. This tensor is decomposed as $\mathbf{5} + \mathbf{10}$ of $SU(5)$:

$$\mathbf{5}: \begin{pmatrix} [dN] \\ [hN] \end{pmatrix}, \quad \mathbf{10}: \begin{pmatrix} [dd] & [dh] \\ [hd] & [hh] \end{pmatrix}. \quad (3)$$

Here, the superfield in $\mathbf{5}$ representation can be identified with the up-type Higgs superfield ($[dN][hN] \sim \Lambda_H(DH)$), where $D = (\mathbf{3}, \mathbf{1}, -1/3)$ and $H = (\mathbf{1}, \mathbf{2}, 1/2)$. The components of the superfield in $\mathbf{10}$ representation are also identified as $[dd] \sim \Lambda_H \bar{t}$, $[dh] \sim \Lambda_H (tb) = \Lambda_H q$, and $[hh] \sim \Lambda_H \bar{\tau}$, respectively.

The dynamically generated superpotential is exactly given by

$$W_{\text{dyn}} = -\frac{1}{\Lambda_H^3} \text{Pf}(M) = -\frac{1}{\Lambda_H^3} \frac{1}{2^3 3!} \epsilon^{ijklmn} M_{ij} M_{kl} M_{mn}. \quad (4)$$

From Eqs. (1) and (4), the total effective superpotential is given by

$$W_{\text{eff}} = \alpha H q \bar{t} + \beta q q D + \gamma \bar{D} \bar{\tau} + \mu \bar{H} \bar{H} + \mu_D \bar{D} D. \quad (5)$$

Note that $\alpha \sim \beta \sim \gamma \sim 1$ is naturally expected, since it can be understood that these Yukawa couplings are constructed by the exchange of the preon superfields. The first term in Eq. (5) is nothing but the top quark Yukawa coupling, which is naturally large. The fourth and fifth terms are the μ terms of the $SU(2)_L$ doublet and the color triplet Higgs superfields, respectively, where $\mu \sim \Lambda_H \eta$ and $\mu_D \sim \Lambda_H \eta_D$ from Eq. (1).

Note that this composite model is the only one which satisfies the following three requirements. The first is that the model is based on the gauge group $SU(N)$ with vectorlike N_f matter fields in the fundamental representation. The sec-

ond is that there is only the ‘‘mesonic’’ effective superfields but not the ‘‘baryonic’’ one. The third is that the Yukawa coupling has to be dynamically generated. These can be satisfied only if $N_f = N + 1$ and $N = 2$.

III. DYNAMICAL SUPERSYMMETRY BREAKING SECTOR

In order for the supersymmetric composite model discussed in the previous section to be realistic, the supersymmetry should be broken. We construct the dynamical supersymmetry-breaking sector in this section.

It is recently pointed out that the theory of the $SU(2)$ gauge interaction with four doublet and six singlet superfields breaks supersymmetry by the nonperturbative effect of the gauge interaction [5]. The dynamical supersymmetry-breaking sector in our model has the same structure as the theory.

This sector is based on the $SU(2)_S$ supercolor gauge group with four doublet and eight singlet superfields. The maximal nonanomalous global symmetry of this sector is $Sp(2) \times U(1)_R$. The $SU(2)$ subgroup of $Sp(2)$ is gauged and identified with the hypercolor gauge group $SU(2)_H$. Note that we treat the hypercolor interaction as weak in the discussion of the dynamical symmetry breaking, because of the assumption $\Lambda_S \gg \Lambda_H$.

The particle contents are as follows:

	$SU(2)_S$	$SU(2)_H$	$U(1)_R$
Q	2	2	0
\tilde{Q}_1	2	1	0
\tilde{Q}_2	2	1	0
Z	1	1	2
Z'	1	1	2
Z_1	1	2	2
Z_2	1	2	2
X	1	1	0
Y	1	1	2

The tree-level superpotential is introduced as

$$W_{\text{tree}} = \lambda \text{tr}[\hat{Z}\hat{V}] + 4\lambda_V ZV + \lambda_Z ZX^2 + \lambda_Y YX^2, \quad (6)$$

where \hat{Z} and \hat{V} are the antisymmetric 4×4 matrices defined by

$$\hat{Z} = \begin{bmatrix} 0 & Z+Z' & Z_1^1 & Z_2^1 \\ -(Z+Z') & 0 & Z_1^2 & Z_2^2 \\ -Z_1^1 & -Z_1^2 & 0 & Z-Z' \\ -Z_2^1 & -Z_2^2 & -(Z-Z') & 0 \end{bmatrix} \quad (7)$$

and

$$\hat{V} = \epsilon^{\alpha\beta} \begin{pmatrix} Q \\ \tilde{Q}_1 \\ \tilde{Q}_2 \end{pmatrix}_\alpha (Q\tilde{Q}_1\tilde{Q}_2)_\beta = \begin{bmatrix} 0 & V+V' & V_{11} & V_{21} \\ -(V+V') & 0 & V_{12} & V_{22} \\ -V_{11} & -V_{12} & 0 & V-V' \\ -V_{21} & -V_{22} & -(V-V') & 0 \end{bmatrix}, \quad (8)$$

respectively. Here, $a=1,2$ of Z_i^a and V_{ia} ($i=1,2$) is the $SU(2)_H$ index, and α and β are the $SU(2)_S$ indices. The field \hat{Z} is in the reducible representation $\mathbf{1} + \mathbf{5}$ of $Sp(2)$: Z is singlet, and Z' and Z_i are in $\mathbf{5}$. This is the same for the field \hat{V} . Note that Eq. (6) is the general superpotential which possesses $Sp(2) \times U(1)_R$ global symmetry.

Below the scale Λ_S where the supercolor interaction becomes strong, this sector is described by the effective superfield \hat{V} . The dynamically generated superpotential is exactly given by

$$W_{\text{dyn}} = A(\text{Pf}\hat{V} - \Lambda_S^4) = A(V^2 - V'^2 - [V_1 V_2] - \Lambda_S^4), \quad (9)$$

where A is the Lagrange multiplier superfield, and the square brackets denote the contraction of the $SU(2)_H$ indices with the ϵ tensor. The total effective superpotential is given by $W_{\text{eff}} = W_{\text{tree}} + W_{\text{dyn}}$.

In the analysis of the vacuum, Eq. (9) is regarded as a constraint, $\text{Pf}\hat{V} - \Lambda_S^4 = 0$. By this constraint one superfield should be eliminated from W_{eff} . We eliminate the $Sp(2)$ singlet superfield V from W_{eff} by the constraint, since it is expected that the vacuum is realized with the largest global symmetry. Then, the effective superpotential is rewritten by

$$W_{\text{eff}} = -4(\lambda - \lambda_V)Z\sqrt{\Lambda_S^4 + V'^2 + [V_1 V_2]} - \lambda\{4Z'V' + 2[Z_1 V_1] + 2[Z_2 V_2]\} + \lambda_Z ZX^2 + \lambda_Y YX^2, \quad (10)$$

where the ambiguity of the sign in front of the square root is absorbed by the $U(1)_R$ phase rotation. By factoring out Λ_S to normalize the dimension of the effective superfields, and expanding the square root, we obtain

$$W_{\text{eff}} \simeq -4(\lambda - \lambda_V)Z\left\{\Lambda_S^2 + \frac{1}{2}V'^2 + \frac{1}{2}[V_1 V_2]\right\} - \lambda\Lambda_S\{4Z'V' + 2[Z_1 V_1] + 2[Z_2 V_2]\} + \lambda_Z ZX^2 + \lambda_Y YX^2. \quad (11)$$

This effective superpotential is one of the type of O’Raifeartaigh models [7], and five supersymmetric vacuum conditions $\partial W_{\text{eff}}/\partial Z = 0$, $\partial W_{\text{eff}}/\partial Z' = 0$, $\partial W_{\text{eff}}/\partial Z_i = 0$ ($i=1,2$) and $\partial W_{\text{eff}}/\partial Y = 0$ cannot be satisfied simultaneously. As a result, the supersymmetry is broken.

Note that the effective potential has the ‘‘pseudoflat direction’’ as O’Raifeartaigh models have. For example, the vacuum energy remains minimum along the direction of ar-

bitrary vacuum expectation value $\langle Z \rangle$,¹ when the other fields have no vacuum expectation values. However, this flatness will disappear, if the quantum correction to the Kähler potential is considered. The effective Kähler potential at low energy cannot be reliably evaluated, because of the strong supercolor interaction. Therefore, at the high energy where the supercolor is weak and negligible, we evaluate it by considering only the Yukawa couplings in the tree-level superpotential. Note that the superfield Z couples to the massless superfield X in Eq. (6).

In the nonsupersymmetric massless $\lambda\phi^4$ theory, it is well known that an infrared divergence appears in the loop correction of the massless field. Although the infrared divergence disappears with summing up the one-loop diagrams, the potential with the one-loop correction becomes singular at the point $\langle\phi\rangle=0$, as a result. The problem of the infrared divergence or the singularity at the origin is solved by the fact that the nonzero vacuum expectation value $\langle\phi\rangle\neq 0$ emerges in the effective potential. This mechanism is pointed out by Coleman and Weinberg [8].

When the effective Kähler potential of the superfield Z is evaluated by considering the loop correction of the massless superfield X , we encounter the same situation with the infrared divergence, since the superfield is treated as the boson field in the supergraph method [9]. Unfortunately, the one-loop effective Kähler potential is not reliable (see the Appendix). However, we can expect that the effective Kähler potential of Z is nontrivial, and the minimum of the effective potential is realized at the point $\langle Z \rangle \neq 0$,² according to the mechanism of Coleman and Weinberg. This discussion is also applied to the effective Kähler potential of the superfield Y , and $\langle Y \rangle \neq 0$ is also expected.

In the following, we assume that Kähler potentials of all the fields except for Z and Y are naive, and their scalar components have no vacuum expectation values. By this assumption, the vacuum energy is given by

$$E_{\text{vac}} = 16(\lambda - \lambda_V)^2 \Lambda_S^4 \left/ \left[\frac{\partial}{\partial(Z^\dagger Z)} \left((Z^\dagger Z) \frac{\partial K_Z}{\partial(Z^\dagger Z)} \right) \right]_{Z=\langle Z \rangle} \right., \quad (12)$$

where K_Z is the nontrivial Kähler potential of the superfield Z . The supersymmetry is broken by the $SU(2)_S$ supercolor gauge dynamics, unless the denominator becomes singular.

IV. MEDIATION OF SUPERSYMMETRY BREAKING

In this section, we investigate the mediation of the supersymmetry breaking. The breaking is mediated to the preon sector and the standard model sector by the radiative corrections due to the $SU(2)_H$ and G_{SM} charged particles. The

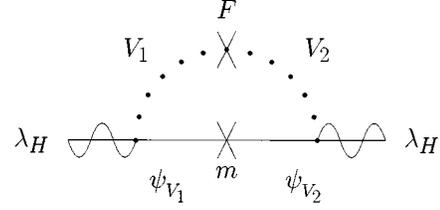


FIG. 1. The one-loop diagram for the soft breaking mass of the hypercolor gaugino. The solid and dotted internal lines denote the propagators of ψ_{V_i} and V_i ($i=1,2$), respectively. The crosses denote the insertions of F and m for the internal boson and fermion lines, respectively.

$SU(2)_H$ doublet fields play the role of the ‘‘messenger fields’’ [3].

Let us consider the mass spectrum of the $SU(2)_H$ doublet fields from the superpotential of Eq. (11). These are given by

$$\begin{aligned} \mathcal{L}_{\text{mass}} = & -(m^2 + m'^2)(V_1^\dagger V_1 + V_2^\dagger V_2) - F[V_1 V_2] + \text{H.c.} \\ & - m'^2(Z_1^\dagger Z_1 + Z_2^\dagger Z_2) - mm'(Z_1^\dagger V_2 + Z_2^\dagger V_1) + \text{H.c.} \\ & - m[\psi_{V_1} \psi_{V_2}] - m'([\psi_{Z_1} \psi_{V_1}] + [\psi_{Z_2} \psi_{V_2}]) + \text{H.c.}, \end{aligned} \quad (13)$$

where

$$\begin{aligned} F \equiv & 8(\lambda - \lambda_V)^2 \Lambda_S^2 \left/ \left[\frac{\partial}{\partial(Z^\dagger Z)} \left((Z^\dagger Z) \frac{\partial K_Z}{\partial(Z^\dagger Z)} \right) \right]_{Z=\langle Z \rangle} \right., \\ m \equiv & 2(\lambda - \lambda_V) \langle Z \rangle, \\ m' \equiv & 2\lambda \Lambda_S. \end{aligned} \quad (14)$$

We can regard F , m , and m' as three independent free parameters instead of the four independent free parameters λ , λ_V , λ_Z , and Λ_S in the following.³ Note that the $SU(2)_H$ doublet superfields Z_i and V_i ($i=1,2$) should be decoupled at the energy scale of Λ_H , in order for the mechanism by Nelson and Strassler really to work. This fact demands that we take $\Lambda_H^2 \ll F$, m^2 , and m'^2 . In the following discussion, we set $m' \ll m$ and $F \ll m^2$, and treat the $SU(2)_H$ coupling perturbatively.

The hypercolor gaugino gets its soft breaking mass through the one-loop diagram in Fig. 1. It is given by

$$m_{\lambda_H} \sim \frac{\alpha_H F}{8\pi m}, \quad (15)$$

where $\alpha_H = g_H^2/4\pi$ is the coupling constant of the hypercolor interaction. This is the first order correction of α_H .

¹We use the same notation for the superfield itself and the scalar component of the superfield.

²The possibility that the scalar component of Z has a large vacuum expectation value without the additional Yukawa coupling λ_Z is pointed out by Shirman [10]. If it is true, we can start a more simple $\text{Sp}(2) \times \text{U}(1)_R$ symmetric superpotential, $W_{\text{tree}} = \lambda \text{tr}[\hat{Z}\hat{V}] + 4\lambda_V ZV$, without introducing the superfields X and Y .

³Note that the Kähler potential K_Z and the vacuum expectation value $\langle Z \rangle$ are determined by the coupling λ_Z .

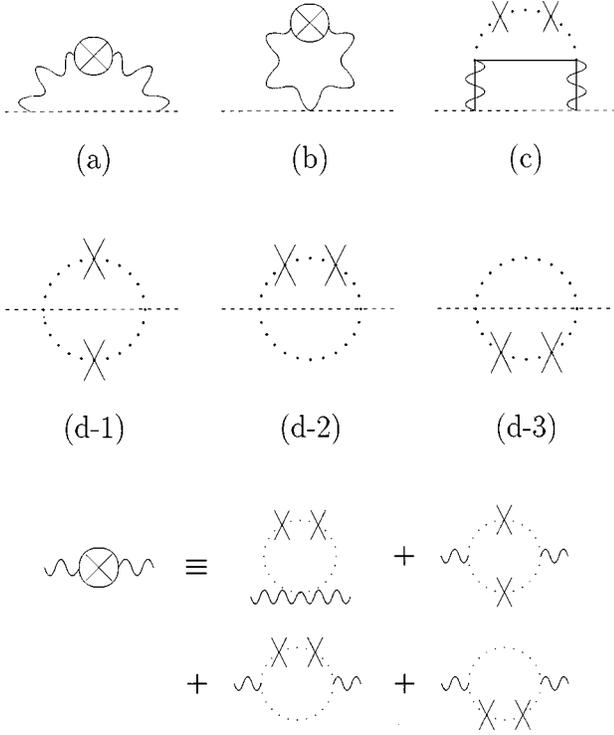


FIG. 2. The two-loop diagrams for the soft breaking masses of the scalar preons. The solid and dotted internal lines denote the propagators of ψ_{V_i} and V_i ($i=1,2$), respectively. The dashed internal line denotes the propagator of the scalar preon \tilde{P} or \tilde{N} . The wavy line denotes the hypercolor gauge boson propagator. The propagator of the hypercolor gaugino is denoted by the wavy line with the solid line. The crosses denote the insertions of F .

The scalar preons \tilde{P} and \tilde{N} also get their masses through the two-loop diagrams in Fig. 2. Since their $SU(2)_H$ charge is the same, they have the same masses given by

$$m_{\tilde{P}}^2 = m_{\tilde{N}}^2 \sim \frac{3}{2} \left(\frac{\alpha_H}{4\pi} \right)^2 \left(\frac{F}{m} \right)^2. \quad (16)$$

This is the second order correction of α_H . Note that the scalar preon masses are the same order of the hypercolor gaugino mass. Since the scalar preons have the masses, the composite scalars, the scalar quarks \tilde{q} and \tilde{t} , the scalar lepton $\tilde{\tau}$, and the up-type Higgs boson also have the same soft breaking masses. We assume that their masses of the composite scalars are twice of the scalar preon mass, for simplicity.

Furthermore, the gauginos in the standard model get their masses through the radiative correction including the preon superfield P , since it has not only $SU(2)_H$ charge but also G_{SM} charge. By the three-loop diagram⁴ in Fig. 3, the mass is given by

$$m_{\lambda_N} \sim \frac{3\alpha_N}{8\pi} \left(\frac{\alpha_H}{4\pi} \right)^2 \frac{F}{m} \ln \left(\frac{m}{\Lambda_H} \right), \quad (17)$$

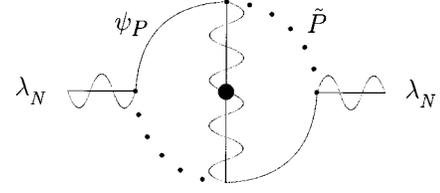


FIG. 3. The three-loop diagram for the soft breaking masses of the standard model gauginos. The solid and dotted internal lines denote the propagators of the preon ψ_P and the scalar preon \tilde{P} , respectively. The internal line of the hypercolor gaugino is denoted by the wavy line with the solid line. The dot on the propagator of the hypercolor gaugino denotes the one-loop diagram in Fig. 1.

where α_N ($N=1,2,3$) is the coupling constant of the $SU(N)$ interaction in the standard model.⁵ This three-loop correction is the second order for α_H , and the first for α_N .

Note that the preon and scalar preon fields play the role of “messenger fields,” and this is a new mechanism to give soft breaking masses to the standard model gauginos. At low energy where the preon and scalar preon are confined, the mechanism can be understood as the mixing between gauginos and the composite fermions in adjoint representation. This is a special picture of this supersymmetric composite model.

We have shown that the soft breaking masses of the particles in the standard model are really generated. However, in the discussion of the electroweak symmetry breaking, the above perturbative estimation cannot be reliable, since the hypercolor interaction becomes strong at low energy. In the following, we use the masses given in Eqs. (16) and (17) with $\alpha_H/4\pi=1$ for order estimation. The condition $m_{\lambda_H}, m_{\tilde{P}} \ll \Lambda_H$ must be satisfied in order for the mechanism by Nelson and Strassler really to work at low energy.

V. STANDARD MODEL SECTOR

Let us consider the standard model sector at the scale where the preon superfields are confined by the $SU(2)_H$ interaction. The large top quark Yukawa coupling is generated as discussed in Sec. II. The scalar quarks (\tilde{q} and \tilde{t}), the scalar lepton ($\tilde{\tau}$), and the up-type Higgs boson have the same soft breaking masses

$$m_{\tilde{q}} = m_{\tilde{t}} = m_{\tilde{\tau}} = m_H \sim 2m_{\tilde{P}} \sim \sqrt{6} \left(\frac{F}{m} \right), \quad (18)$$

where we use Eq. (16) with $\alpha_H/4\pi=1$. The standard model gauginos have masses as given by Eq. (17) with $\alpha_H/4\pi=1$. Note that the ratio of the masses of two gauginos is determined by the ratio of two gauge couplings:

$$m_{\lambda_N}/m_{\lambda_M} = \alpha_N/\alpha_M, \quad (19)$$

⁵In this calculation, the infrared divergence occurs, which means that the mass at zero momentum cannot be defined. We have taken Λ_H as the cutoff parameter of the infrared divergence, since the preon superfield P is confined at the energy lower than Λ_H .

⁴A similar diagram was considered by Randall in Ref. [11].

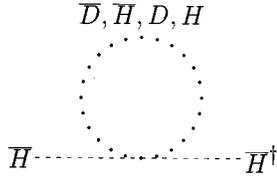


FIG. 4. The one-loop diagram for the soft breaking mass of the down-type Higgs boson. The dotted internal line denotes the propagators of \bar{D} , \bar{H} , D , and H . It is not needed to consider the corrections of the scalar quarks and leptons, because their contributions cancel out each other.

where $N, M = 1, 2, 3$. This relation is the same as obtained in the gauge-mediated model of Ref. [3].

Furthermore, other soft breaking masses in the standard model are also induced by the radiative correction through the standard model gauge interaction. By the one-loop diagram in Fig. 4, the mass of the down-type Higgs boson \bar{H} is given by

$$m_{\bar{H}}^2 \sim \frac{\alpha_1}{8\pi} \left[\mu_D^2 \ln \left(1 + \frac{m_q^2}{\mu_D^2} \right) + m_q^2 \ln \left(\frac{\mu_D^2 + m_q^2}{\mu^2 + m_q^2} \right) - \mu^2 \ln \left(1 + \frac{m_q^2}{\mu^2} \right) \right]. \quad (20)$$

The soft breaking term $B\bar{H}H + \text{H.c.}$ in the Higgs potential is also induced by the one-loop diagram in Fig. 5. We obtain

$$B \sim \frac{3\alpha_2}{2\pi} m_{\lambda_2} \mu \ln \left(\frac{\Lambda_H}{m_{\lambda_2}} \right), \quad (21)$$

where the ultraviolet divergence is cut off by the composite scale Λ_H , and we assume $\mu^2/m_{\lambda_2}^2 \ll 1$.

By considering all the estimated soft breaking masses, the Higgs potential is given by

$$V = \frac{1}{2}(\mu^2 + m_{\bar{H}}^2)\bar{h}^2 + \frac{1}{2}(\mu^2 + m_q^2)h^2 + B\bar{h}h + \frac{g_1^2 + g_2^2}{32}(\bar{h}^2 - h^2)^2, \quad (22)$$

where g_1 and g_2 are the gauge coupling constants, and \bar{h} and h are the real parts of the neutral component of \bar{H} and H , respectively. Since $m_{\bar{H}}^2 > 0$ and $B < m_q^2$ from Eqs. (17), (18), (20), and (21) with $\mu^2 < \mu_D^2$, the condition $B^2 < (\mu^2 + m_{\bar{H}}^2)(\mu^2 + m_q^2)$ is satisfied. Therefore, the electroweak symmetry cannot be broken at this level.

VI. ELECTROWEAK SYMMETRY BREAKING

We have seen that the electroweak symmetry cannot be broken in Eq. (22). However, if we consider the radiative corrections due to the large top quark Yukawa coupling, the electroweak symmetry breaking really emerges according to the ‘‘radiative breaking scenario’’ [1].

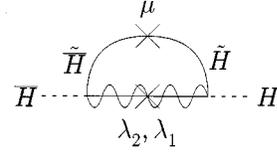


FIG. 5. The one-loop diagram to generate the soft breaking term $B\bar{H}H + \text{H.c.}$ in the Higgs potential. The solid internal lines denote the propagators of the Higgsinos \tilde{H} or $\tilde{\bar{H}}$. The wavy line with the solid line denotes the propagator of the $SU(2)_L$ gaugino or the $U(1)_Y$ gaugino. Since the contribution of the $U(1)_Y$ gaugino loop is far smaller than that of the $SU(2)_L$ gaugino, it can be neglected in Eq. (21).

Let us consider the one-loop radiative corrections induced by the top quark and scalar top quark loops. The one-loop effective potential is given by

$$V_{\text{one loop}} = \frac{3}{8\pi^2} \int_0^{\Lambda_H^2} dk^2 k^2 \left[\ln \left(1 + \frac{g_t^2 h^2/2}{k^2 + m_q^2} \right) - \ln \left(1 + \frac{g_t^2 \bar{h}^2/2}{k^2} \right) \right], \quad (23)$$

where we use Λ_H as a cutoff parameter because of the compositeness of the top quark, the scalar top quark, and the up-type Higgs doublet. In Eq. (23), the first and second terms in the square brackets are given by the loop corrections of the scalar top and top quarks, respectively. One can check $V_{\text{one loop}} \rightarrow 0$ in the supersymmetric limit $m_q^2 \rightarrow 0$. Now, we obtain the effective potential $V_{\text{eff}} = V + V_{\text{one loop}}$ from Eqs. (22) and (23).

In order for the electroweak symmetry to be broken, two minimization conditions $\partial V_{\text{eff}}/\partial \bar{h} = 0$ and $\partial V_{\text{eff}}/\partial h = 0$ should be satisfied with nonzero values of $\langle \bar{h} \rangle$ and $\langle h \rangle$. Furthermore, these vacuum expectation values should satisfy the condition $\langle \bar{h} \rangle^2 + \langle h \rangle^2 = v^2$ to realize the correct scale of the electroweak symmetry breaking, where $v \approx 246$ GeV. By considering this condition, two minimization conditions are described by

$$\frac{\partial V_{\text{eff}}}{\partial \bar{h}} = 0 \Rightarrow \mu^2 + m_{\bar{H}}^2 + \frac{1}{2} M_Z^2 \cos 2\beta + B \tan \beta = 0 \quad (24)$$

and

$$\begin{aligned} \frac{\partial V_{\text{eff}}}{\partial h} = 0 \Rightarrow & \mu^2 + m_q^2 - \frac{1}{2} M_Z^2 \cos 2\beta + B \cot \beta \\ & + \frac{3}{4\pi^2} \frac{m_t^2}{v^2 \sin^2 \beta} \left[m_t^2 \ln \left(\frac{m_q^2 + m_t^2}{m_t^2} \right) \right. \\ & \left. - m_q^2 \ln \left(\frac{\Lambda_H^2}{m_q^2 + m_t^2} \right) \right] = 0, \end{aligned} \quad (25)$$

where we use the definitions $\langle \bar{h} \rangle = v \cos \beta$, $\langle h \rangle = v \sin \beta$, and $m_t^2 = g_t^2 \langle h \rangle^2/2$ for the top quark mass and M_Z^2

$=(g_1^2 + g_2^2)v^2/4$ for the Z boson mass. Note that $\mu < 0$ is required for $\tan\beta$ to be positive from Eqs. (21) and (24).

Now we show that there is a realistic solution by substituting the concrete values to the parameters in Eqs. (24) and (25). There are five independent unknown parameters in the two equations: μ , $m_{\bar{H}}$, $m_{\bar{q}}$, B , and $\tan\beta$. These parameters are described by the six parameters μ , μ_D , $m_{\bar{q}}$, Λ_H , $r \equiv m/\Lambda_H$, and $\tan\beta$ through Eqs. (17), (18), (20), and (21). We set the values of the four parameters as $\mu = -200$ GeV, $\mu_D = 1$ TeV, $m_{\bar{q}} = 20$ TeV, and $r = 1000$. The values of μ_D and $m_{\bar{q}}$ have to be large enough, since there is no evidence of the production of the exotic colored particle. The value of r has been set as $r \gg 1$ to keep the hierarchy which is assumed in the previous section. The remaining two parameters Λ_H and $\tan\beta$ are determined by solving Eqs. (24) and (25). By numerical calculation, we can find the solution

$$\tan\beta \sim 3.7, \quad \Lambda_H/m_{\bar{q}} \sim 9.3 \times 10^4, \quad (26)$$

with $\alpha_3 = 0.12$, $\alpha_2 = 0.033$, $\alpha_1 = 0.017$, $m_t = 180$ GeV, and $v = 246$ GeV. Note that the hierarchy $m_{\lambda_H}, m_{\bar{p}} \ll \Lambda_H$, which must be satisfied to work the mechanism proposed by Nelson and Strassler, is realized, because of $m_{\lambda_H} \sim m_{\bar{p}} \sim m_{\bar{q}} \ll \Lambda_H$ [see Eqs. (15), (16), and (18)]. As a result, the compositeness scale is very high.

The values of $\tan\beta$ and $\Lambda_H/m_{\bar{q}}$ are increasing functions of $|\mu|$ and μ_D , while decreasing functions of $m_{\bar{q}}$ and r . That parameter dependence of the solution can be seen analytically by the rough estimation of the solution. We regard Eq. (24) as the equation to determine the value of $\tan\beta$. From Eqs. (17), (18), and (21), we obtain

$$\begin{aligned} \tan\beta &\sim -(\mu^2 + m_{\bar{H}}^2)/B \\ &\sim -(\mu^2 + m_{\bar{H}}^2) \left/ \left[\frac{3\sqrt{6}\alpha_2^2}{32\pi^2} m_{\bar{q}} \mu \ln(r) \right] \right., \end{aligned} \quad (27)$$

where the term $1/2M_Z^2 \cos 2\beta$ is neglected with $\mu = -200$ GeV, and the logarithm in Eq. (21) is considered as order 1. Substituting Eq. (20) into (27), one can check the parameter dependence of $\tan\beta$ mentioned above. On the other hand, we obtain, from Eq. (25),

$$\frac{\Lambda_H}{m_{\bar{q}}} \sim \exp \left[\frac{2\pi^2 v^2}{3m_t^2} \frac{\tan^2 \beta}{1 + \tan^2 \beta} \right], \quad (28)$$

where the relation $m_{\bar{q}}^2 \gg \mu^2$, M_Z^2 , and B is used. The value of $\Lambda_H/m_{\bar{q}}$ has the same dependence as $\tan\beta$, since it is the increasing function of $\tan\beta$.

By using $m_{\bar{q}}$, the mass of the gauginos in the standard model sector is given by

$$m_{\lambda_N} \sim \frac{\sqrt{6}\alpha_N}{16\pi} m_{\bar{q}} \ln(r). \quad (29)$$

Then, we obtain the values of the masses of the gauginos for $m_{\bar{q}} = 20$ TeV as

$$m_{\lambda_3} \sim 810 \text{ GeV}, \quad m_{\lambda_2} \sim 220 \text{ GeV}, \quad m_{\lambda_1} \sim 110 \text{ GeV}. \quad (30)$$

These values are experimentally acceptable [12].

VII. PROTON DECAY PROBLEM

In our model there are colored Higgs superfields \bar{D} and D which may cause proton decay. The baryon-number-violating interaction at the tree level in the first and second generations is forbidden by imposing usual $U(1)_{B-L}$ symmetry, under which \bar{D} and D are singlets. On the other hand, there is another $U(1)_{B-L}$ symmetry in the third generation, under which \bar{D} and D have charge $2/3$ and $-2/3$, respectively. This symmetry allows the baryon-number-violating interactions as the second and third terms in Eq. (5). These interactions may cause too rapid proton decay through the flavor mixing between the third and other generations which reduces these two symmetries to the unique $U(1)_{B-L}$ symmetry.⁶

To discuss the problem of proton decay, we should know the quark and lepton mass matrices: in other words, we need a model which explains the origin of masses of all quarks and leptons. Although such a model is not yet constructed, we can discuss proton decay with some assumptions which can be reasonably expected in our model.

In order to generate the up quark mass and the mixing between the up and top quarks, the following nonrenormalizable interactions must be introduced in the superpotential:

$$\begin{aligned} W &= \frac{\kappa_{uu}}{M_{uu}} [hN] q_1 \bar{u} + \frac{\kappa_{ut}}{M_{ut}^2} [hN] q_1 [dd] + \frac{\kappa_{tu}}{M_{tu}^2} [hN] [dh] \bar{u} \\ &\sim \left(\frac{\Lambda_H}{M_{uu}} \right) H q_1 \bar{u} + \left(\frac{\Lambda_H}{M_{ut}} \right)^2 H q_1 \bar{t} + \left(\frac{\Lambda_H}{M_{tu}} \right)^2 H q_3 \bar{u}, \end{aligned} \quad (31)$$

where q_i denotes the quark doublet in the i th generation, M_{uu} , M_{ut} , and M_{tu} are parameters with dimension 1, and we take $\kappa_{uu} \sim \kappa_{ut} \sim \kappa_{tu} \sim 1$ in the second line. We would be able to assume that $M_{uu} \sim M_{ut} \sim M_{tu}$, and the mixing angle between the up and top quarks is given by

$$\theta_{ut} \sim \left(\frac{m_u}{m_t} \right)^2. \quad (32)$$

The same argument is applied to the second generation, and the mixing angle between the charm and top quarks is given by $(m_c/m_t)^2$. Note that the flavor mixings are highly suppressed in the up-type quark sector.

On the other hand, for the down-type quarks and charged leptons, we cannot obtain such relations, since \bar{H} is the elementary particle and no higher dimensional terms are needed to introduce the Yukawa interactions. However, we can assume that the mixing angle between the down and bottom quarks is of the order of the (1,3) or (3,1) element of

⁶It is possible to identify $\bar{\tau}$ in Eq. (5) with a singlet heavy particle instead of the right-handed τ , as done in the original paper by Nelson and Strassler [6]. In this case, there is no problem on the proton decay.

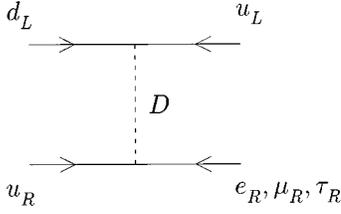


FIG. 6. The four-fermion interaction caused by the colored Higgs boson exchange. D denotes the colored Higgs boson.

the Kobayashi-Maskawa matrix ($\theta_{13}^{\text{KM}} \sim 10^{-3}$ [12]), since the mixing angles between the up-type quarks are very small as assumed above.

Let us discuss the proton lifetime which is caused by the four-fermion interaction mediated by the colored Higgs boson and Higgsino. The lifetime is roughly described as

$$\tau_p \sim \frac{G_D^{-2}}{m_p^5}, \quad (33)$$

where G_D is the coupling constant of the four-fermion interaction, and m_p is the proton mass. The experimental bound is obtained by

$$G_D < 2 \times 10^{-32} \text{ GeV}^{-2} \quad (34)$$

from $\tau_p > 10^{32}$ yr [12].

At the tree level, there are decay modes $p \rightarrow \pi^0 e^+ (\mu^+)$ through lepton mixing and $p \rightarrow \pi^0 (\tau^+)^*$ (the asterisk means off shell) caused by the four-fermion interaction through the colored Higgs boson exchange depicted in Fig. 6. For the mode $p \rightarrow \pi^0 e^+ (\mu^+)$, G_D is roughly estimated as

$$G_D \sim \frac{2}{m_D^2} \theta_{ut}^2 \theta_{13}^{\text{KM}} \theta_{e\tau(\mu\tau)}, \quad (35)$$

where $\theta_{e\tau(\mu\tau)}$ is the mixing angle between the right-handed electron (muon) and τ , and m_D is the colored Higgs boson mass [note that $m_D^2 = m_{\tilde{q}}^2 + \mu_D^2$ from Eqs. (5) and (18)]. By substituting the numerical values used in Sec. VI and $m_u = 5$ MeV, we obtain $G_D \sim 3 \times 10^{-30} \theta_{e\tau(\mu\tau)}$. Therefore, there is no contradiction with the experiments, if $\theta_{e\tau(\mu\tau)} < 10^{-2}$. In the case of the $p \rightarrow \pi^0 (\tau^+)^*$ mode, $\theta_{e\tau(\mu\tau)}$ is replaced by the factor of the off-shell τ ‘‘decay,’’ $G_F m_p^3 / m_\tau \sim 10^{-5}$. We obtain the sufficiently small value $G_D \sim 3 \times 10^{-35}$.

At the one-loop level, the dominant decay modes are $p \rightarrow \pi^0 e^+ (\mu^+)$ and $p \rightarrow \pi^0 (\tau^+)^*$, which are caused by the four-fermion interaction through the colored Higgsino exchange depicted in Fig. 7. For the mode $p \rightarrow \pi^0 e^+ (\mu^+)$, we obtain

$$G_D \sim \frac{\alpha_3}{\pi} \frac{1}{m_{\tilde{q}}^2} \theta_{ut}^2 \theta_{13}^{\text{KM}} \theta_{e\tau(\mu\tau)} \sim 6 \times 10^{-32} \theta_{e\tau(\mu\tau)}. \quad (36)$$

Here, note that the dependence on μ_D and m_{λ_3} is absent because of $\mu_D^2, m_{\lambda_3}^2 \ll m_{\tilde{q}}^2$. This value of G_D is allowed, if

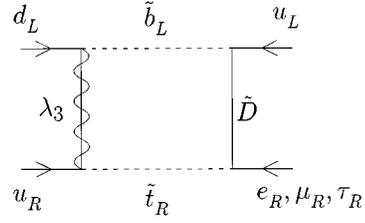


FIG. 7. The four-fermion interactions caused by the colored Higgsino exchange. \tilde{D} denotes the colored Higgsino. The diagram by replacement $d_L \leftrightarrow u_L$ and $\tilde{b}_L \rightarrow \tilde{t}_L$ is also possible.

$\theta_{e\tau(\mu\tau)} < 10^{-2}$. The mode $p \rightarrow \pi^0 (\tau^+)^*$ is also substantially suppressed, since the angle $\theta_{e\tau(\mu\tau)}$ is replaced by the factor of the off-shell τ ‘‘decay.’’

From the following argument, one can expect that proton decay is sufficiently suppressed in any process. Since the baryon-number-violating interactions exist only in the third generation, the baryon-number-violating interactions in the first and second generations can be induced only through the flavor mixing between the third and other generations. Therefore, in any proton decay process, G_D is always suppressed by the flavor mixing angles. The angles in the up-type quark sector can be extremely small as reasonably expected above. Moreover, the angles for the flavor mixing of the scalar particles are very small, because of the decoupling effect due to the heavy scalars in the third generation. The value of G_D becomes very small by these small mixing angles.

VIII. R-AXION PROBLEM

In Sec. III, we assume that $\langle Z \rangle \neq 0$ and $\langle Y \rangle \neq 0$. This means the spontaneous breaking of the global $U(1)_R$ symmetry, and the Nambu-Goldstone boson called the R axion appears. Since the quark superfields have nonzero R charge, they couple to the R axion. This coupling may cause an astrophysical or a cosmological problem, such as an overcooling of the red giants or a breakdown of the successful predictions from big bang nucleosynthesis. We consider the possibility to avoid the problem.

Let us introduce a $\lambda_X X^3$ term into the tree-level superpotential of Eq. (6). Note that the new term does not disturb the supersymmetry breaking. The R charge of the superfields is changed as follows:

$$\begin{array}{ll} U(1)_R & \\ \hat{V} & 4/3 \\ \hat{Z} & 2/3 \\ X & 2/3 \\ Y & 2/3 \end{array}$$

Note that the superfields Q and \tilde{Q}_i must have R charge, because of the nonzero R charge of \hat{V} . Therefore, the new $U(1)_R$ symmetry is explicitly broken by the $SU(2)_S$ supercolor anomaly, and the light R axion disappears. However, the generality of the superpotential is lost by this strategy.

For example, we cannot introduce a $\lambda' Y^3$ term, which is allowed by the symmetry, into the superpotential, since the term recovers the supersymmetry.

IX. SUMMARY AND COMMENTS

We have constructed a supersymmetric composite model with dynamical supersymmetry breaking. The model is based on the gauge group $SU(2)_S \times SU(2)_H \times G_{SM}$. There are three building blocks, namely, the dynamical supersymmetry-breaking sector, the preon sector recently proposed by Nelson and Strassler, and the standard model sector.

In the dynamical supersymmetry-breaking sector, the nonzero vacuum expectation value of the F component of the gauge singlet superfield Z arises by the nonperturbative effect of the $SU(2)_S$ supercolor interaction. Therefore, supersymmetry is dynamically broken. In addition, the scalar component of Z is assumed to have a nonzero vacuum expectation value by the effective Kähler potential including the radiative corrections due to the massless gauge singlet superfield X . The $U(1)_R$ symmetry is spontaneously broken by this vacuum expectation value.

In the preon sector, the preon superfields are confined into the up-type Higgs boson in **5** and the superfield in **10** of $SU(5) \supset G_{SM}$ by the nonperturbative effect of the $SU(2)_H$ hypercolor interaction. The large top quark Yukawa coupling is dynamically generated, and it is understood as the exchange of the preons in the top quark and Higgs doublet.

The supersymmetry breaking is mediated to the preon sector and the standard model sector by the radiative corrections. The hypercolor gaugino gets the soft breaking mass through the one-loop diagram in Fig. 1, and the scalar preons also get the masses from the two-loop diagrams in Fig. 2. The soft breaking masses of the composite scalars, the scalar quarks and lepton, and the up-type Higgs boson originate from the mass of the scalar preons. The gauginos in the standard model get the soft breaking masses through the three-loop diagram in Fig. 3. It is crucial that the preon superfield P have both charges of the hypercolor and the standard model gauge groups, and that it play a role of the ‘‘messenger field.’’ This is a new mechanism to generate the soft breaking masses of the standard model gauginos. Below the scale where the hypercolor interaction becomes strong, this mechanism can be understood as the mixing between the gauginos and the massive composite fermions in adjoint representation.

This model predicts the spectrum of the soft breaking masses in the standard model sector. All composite scalar fields have the same soft mass, and the masses of the gauginos satisfy the relation of Eq. (19). All possible soft breaking masses in the Higgs potential are generated by the radiative correction due to the standard model gauge interactions. Although the Higgs potential with all the soft breaking masses cannot cause the electroweak symmetry breaking, the breaking is realized by the effect of the large top quark Yukawa coupling. This is the radiative breaking scenario, which originates from the dynamics of the supercolor and the hypercolor interactions. We have shown that electroweak symmetry breaking can occur with experimentally acceptable values of the soft breaking masses in the standard model.

In our model, there is the colored Higgs boson which is strongly coupled to quarks and leptons as Eq. (5). It seems that the interaction causes too rapid proton decay through flavor mixing. However, in according to Eq. (31), it can be reasonably assumed that the mixing between the third and other generations is highly suppressed. As a result, there is no dangerous proton decay.

There is a potential R -axion problem in the model, since the nonanomalous $U(1)_R$ global symmetry is spontaneously broken by the nonzero vacuum expectation value of the scalar component of Z . For example, this problem can be avoided by introducing the new term $\lambda_X X^3$ into the superpotential of Eq. (6), because this term requires us to change the charge assignment of $U(1)_R$. Since the new $U(1)_R$ symmetry is explicitly broken by the supercolor anomaly, no light R axion emerges. However, the generality of the superpotential is lost.

Finally, we comment on the flavor-changing neutral current between the first and second generations. The scalar partners in these generations can get the soft breaking masses by radiative corrections as in Fig. 4 and/or Fig. 5. Since the scalar partners with the same gauge charge are degenerate, the flavor-changing neutral current between these two generations is suppressed. This feature is the same as the gauge-mediated model [3].

APPENDIX: THE EFFECTIVE KÄHLER POTENTIAL

We discuss the effective Kähler potential of the superfield Z generated by the Yukawa coupling λ_Z in the superpotential of Eq. (6). We calculate the infinite series of the one-loop graphs, and obtain the effective Kähler potential as

$$K_{\text{eff}} = Z^\dagger Z \left[1 - \frac{\lambda_Z^2}{8\pi^2} \ln \left(\frac{\lambda_Z^2 Z^\dagger Z}{\mu^2} \right) \right], \quad (\text{A1})$$

where μ is the scale introduced by the wave function renormalization of the superfield Z . Then, the vacuum energy is given by

$$E_{\text{vac}} = \left| \frac{\partial W}{\partial Z} \right|^2 \bigg/ \left[1 - \frac{\lambda_Z^2}{8\pi^2} \left\{ 2 + \ln \left(\frac{\lambda_Z^2 Z^\dagger Z}{\mu^2} \right) \right\} \right]_{Z=\langle Z \rangle}. \quad (\text{A2})$$

One may expect that the vacuum is realized at $\langle Z \rangle = 0$, since the denominator diverges and E_{vac} vanishes. However, the one-loop approximation cannot be valid near the origin, since the logarithm becomes too large in Eq. (A1). Although the complete calculation seems to be impossible like in the nonsupersymmetric massless $\lambda\phi^4$ theory discussed by Coleman and Weinberg, we expect that $\langle Z \rangle \neq 0$ is realized. The appearance of the nonzero vacuum expectation value seems to be a common mechanism to solve the problem of the infrared divergence connected with the singularity of the effective potential at the origin. Indeed, Coleman and Weinberg have shown that such mechanism really works in the nonsupersymmetric massless $\lambda\phi^4$ theory with $U(1)$ gauge interaction in their original paper.

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