

Dynamical GUT breaking and μ -term driven supersymmetry breaking

Ryuichiro Kitano

Stanford Linear Accelerator Center, Stanford University, Stanford, California 94309, USA

(Received 21 June 2006; published 6 December 2006)

Models for dynamical breaking of supersymmetric grand unified theories are presented. The doublet-triplet splitting problem is absent since the Higgs-doublet superfields can be identified with the massless mesons of the strong gauge group whereas there are no massless states corresponding to the colored-Higgs fields. Various strong gauge groups $SU(N_c)$, $Sp(N_c)$, and $SO(N_c)$ are examined. In a model with $SO(9)$ strong gauge group, adding a μ -term for the Higgs fields triggers to break supersymmetry in a metastable vacuum. The pattern of the supersymmetry breaking parameters is predicted to be of the gauge-mediation type with modifications in the Higgs sector.

DOI: [10.1103/PhysRevD.74.115002](https://doi.org/10.1103/PhysRevD.74.115002)

PACS numbers: 12.60.Jv, 12.10.Dm

I. INTRODUCTION

The Higgs boson, which has not been observed yet, is the most mysterious particle in the standard model although it plays an important role: electroweak symmetry breaking and the origin of the fermion masses. Successful electroweak symmetry breaking needs a negative mass squared for this particle and its size must be anomalously small compared to the cutoff scale of the theory. This situation motivated us to consider the supersymmetric standard model to protect the mass parameter from large quantum correction, but the Higgs particle is still left mysterious. The Higgs boson (or Higgsino) mass parameter, the μ -term, cannot be protected by gauge symmetry or supersymmetry (SUSY) although quantum corrections are successfully removed.

Embedding the supersymmetric standard model into grand unified theories (GUT), motivated by gauge coupling unification [1], makes the Higgs particle more mysterious. Even for the smallest group for grand unification, $SU(5)$ [2], the Higgs fields do not fit into a complete multiplet of the symmetry group, and therefore we need extra particles to appear at the GUT scale $M_G \sim 10^{16}$ GeV. One mysterious feature is that, it is not easy to realize this situation with nearly massless Higgs fields in models with the GUT symmetry breaking by the Higgs mechanism. The splitting of masses between the Higgs fields and extra particles must be done by a coupling to the fields whose vacuum expectation values (VEV) break the GUT symmetry group, but in simple models it gives masses to the Higgs fields of the order of M_G . This is the famous doublet-triplet splitting problem. Another mysterious feature of SUSY-GUT models is the absence or suppression of dimension-five proton-decay operators. For example, in the simplest $SU(5)$ model, the Higgs particles are embedded into the $\mathbf{5}$ and $\bar{\mathbf{5}}$ representations which contain two colored partner of the Higgs fields H_C and \bar{H}_C . If these colored-Higgs fields have masses by pairing to each other, i.e., $W \ni m_C H_C \bar{H}_C$, dimension-five operators suppressed by the scale m_C are generated by integrating out H_C and \bar{H}_C [3]. It has been studied that the colored-Higgs mass m_C has to be quite

large $m_C \gtrsim 10^{17}$ GeV, which is disfavored by the unification of the gauge couplings [4].

The above two problems, the doublet-triplet splitting and the proton decay, are actually related and there is a simple solution to these problems. If we are to avoid dimension-five proton-decay operators such as $QQQL$ in the superpotential, with Q and L being the quark and lepton doublets, an easy way is to impose a symmetry under which matter fields are charged, e.g., both Q and L have charge unity. In $SU(5)$ SUSY GUTs, this means that both of the Higgs fields in $\mathbf{5}$ and $\bar{\mathbf{5}}$ representations have charge -2 in order to have Yukawa interactions, and thus the mass term is forbidden. Now, to give mass terms to the colored-Higgs fields while keeping the symmetry unbroken we need to introduce another pair of $\mathbf{5}$ and $\bar{\mathbf{5}}$ field which have charge $+2$, but in this case, we have either zero or four Higgs-doublet fields at low energy which is unacceptable. Of course, adding another pair of $\mathbf{5}$ and $\bar{\mathbf{5}}$ field with charge -2 results in extra massless colored-Higgs fields. Therefore, in order to have only two Higgs doublets while forbidding the proton decay by continuous symmetry, we need to repeat the procedure of adding $\mathbf{5}$ and $\bar{\mathbf{5}}$ forever and end up with an infinite number of Higgs fields in $\mathbf{5}$ and $\bar{\mathbf{5}}$ representations [5].

While an infinite number of particles sounds unreasonable in field theory, it is quite possible to realize this situation in models with extra-dimensions. The infinite number of particles is identified with the Kaluza-Klein tower of the fields which are propagating into the bulk of the extra dimension. Indeed, simple GUT models have been constructed in higher dimensional space-time, where the boundary condition breaks the GUT symmetry and there is no doublet-triplet splitting [6] or the proton decay problem [7]. In this picture, the Higgs particles become less mysterious. They are just bulk fields.

On the other hand, there is another familiar mechanism of having an infinite tower of particles in field theory, that is actually happening in QCD. When an asymptotically free gauge theory becomes strong at low energy, the effective theory below that scale is described by gauge-singlet

particles such as mesons and baryons. These particles have also an infinite tower of excitation states. This fact naturally leads us to think of the possibility of realizing Higgs particles as composite fields in some strongly coupled gauge theory which breaks the GUT symmetry dynamically. This question is also interesting from the viewpoint of the anti-de Sitter (AdS)/conformal field theory (CFT) correspondence [8,9]. The extra-dimensional GUT models above may be interpreted as a dual picture of the strongly coupled theory. The explicit gauge symmetry breaking in the extra-dimensional picture may be justified by the presence of viable dynamical GUT breaking models.

Constructing GUT models associated with a strongly coupled gauge theory have been attempted by the group of Hotta, Izawa, and Yanagida [10–13]. (See also [14] for subsequent works.) Various gauge groups for the strong interaction, $SU(3) (\times U(1))$ [10,11], $SU(5)$ [13], and $SO(6)$ [12], were considered. It was found that the doublet-triplet splitting can be easily realized via the missing partner mechanism while preserving an (anomalous) $U(1)$ symmetry which forbids dimension-five proton decays. Along a similar line, a model with $Sp(2)$ gauge group has also been constructed recently in Ref. [15] where the model is quite simplified. (Our convention is such that $Sp(1) \simeq SU(2)$.) The model consists of six flavors of quarks of $Sp(2)$ and five of the flavors are identified with the $\mathbf{5}$ and $\bar{\mathbf{5}}$ representation fields of the $SU(5)$ GUT. The other flavor turns out to be a (constituent of) the Higgs doublets in low energy. In this model, the CFT nature of the $Sp(2)$ interaction above the GUT scale plays a crucial role. A similar approach in warped extra dimension can also be found in Ref. [16].

In this paper, we consider a generalization of the $Sp(2)$ model. We find that models with $Sp(N_c)$ with $N_c = 2$, and $SO(N_c)$ with $6 \leq N_c \leq 9$ work for dynamical GUT breaking while having massless doublet Higgs fields, and no viable $SU(N_c)$ group is found under the assumption on the particle content and superpotential. Of particular interest is the case with $SO(N_c)$ gauge group. There is no exotic particle left massless without adding superpotential terms to remove those particles.

With the success of the doublet-triplet splitting while forbidding the proton decay, the final missing piece for the Higgs mystery is the finite μ -term. In the gravity-mediated supersymmetry breaking scenario [17], it is possible to obtain a correct size of the μ -term in a simple way [18]. However, we find an alternative interesting possibility in the $SO(9)$ model. Instead of solving the μ -problem, if we add a small μ -term in the superpotential, supersymmetry breaks down at the intermediate scale $F \sim \mu M_G$. Although this is not the true vacuum, it is shown to be metastable [19]. It is amusing that the μ -term can drive supersymmetry breaking, which is the opposite direction to the usual thought. The smallness of the μ -term is “explained” by demanding a low supersymmetry breaking scale. The

Higgs fields can be responsible not only for electroweak symmetry breaking but also for GUT and supersymmetry breaking.

There is an interesting possibility for the nature of the small μ -term. The small μ -term added by hand can come from a negative cosmological constant term in the supergravity action. By the Giudice-Masiero mechanism [18], the (supersymmetric) cosmological constant term induces a μ -term in the presence of a particular Kähler potential term. This μ -term, in turn, drives supersymmetry breaking which gives positive contribution to the vacuum energy and cancels the net cosmological constant.

We start the discussion of dynamical GUT breaking in Sec. II; there the general setup is defined. We analyze a successful model, the $SO(9)$ model, in Sec. III. In Sec. IV, the mechanism of μ -term driven supersymmetry breaking is presented and we discuss the mediation of the supersymmetry breaking to our sector. The generation of the μ -term through the cosmological constant is discussed in Sec. V.

II. DYNAMICAL GUT BREAKING

In Ref. [15], a simple model for the dynamical GUT breaking was constructed based on an $Sp(2)$ gauge theory. We study a generalization of the model with various gauge groups: $SU(N_c)$, $Sp(N_c)$, and $SO(N_c)$. Models with $Sp(N_c)$ and $SO(N_c)$ with a certain range of N_c are found to be viable but no viable $SU(N_c)$ group is found. Although they are not successful models, we start with the discussion of the $SU(N_c)$ models in which we can see the essential features of this class of models.

A. $SU(N_c)$ models

The model consists of six flavors, and five of which carry $SU(5)_{\text{GUT}}$ quantum numbers as listed in Table I. The quarks and leptons in the standard model are not charged under $SU(N_c)$ and are unified usually as $\mathbf{10}$ and $\bar{\mathbf{5}}$ of $SU(5)_{\text{GUT}}$. We introduce a superpotential for Q and \bar{Q} :

$$W = m \text{Tr}(Q\bar{Q}) - \frac{1}{M} \text{Tr}[(Q\bar{Q})(Q\bar{Q})] + \dots, \quad (1)$$

where $(Q\bar{Q})$ is the $SU(N_c)$ singlet 5×5 matrix, and “ \dots ” represents other higher dimensional operators such as $(\text{Tr}(Q\bar{Q}))^2$ and those are not important for the discussion. It is essential for the masslessness of the Higgs fields that the superfields T and \bar{T} do not have a superpotential at tree

TABLE I. The particle content of the $SU(N_c)$ model.

	$SU(N_c)$	$SU(5)_{\text{GUT}}$
Q	N_c	$\mathbf{5}$
\bar{Q}	\bar{N}_c	$\bar{\mathbf{5}}$
T	N_c	$\mathbf{1}$
\bar{T}	\bar{N}_c	$\mathbf{1}$

level since the Higgs fields will be identified with the meson fields $H \sim Q\bar{T}$ and $\bar{H} \sim \bar{Q}T$.

Before the analysis at the quantum level, it is helpful for the understanding of the model to discuss what happens at the classical level. The classical analysis is valid for $\Lambda \ll M_G$ with Λ being the dynamical scale of $SU(N_c)$. In this case, the picture becomes similar to models with product group unification [20]. At the classical level, there are vacua with $\text{rank}(Q\bar{Q}) = 0$ to $\min[5, N_c]$ which satisfy the conditions of $F_Q = F_{\bar{Q}} = 0$. We are interested in the vacuum with $\text{rank}(Q\bar{Q}) = 2$:

$$(Q\bar{Q}) = \begin{pmatrix} 0 & & & \\ & 0 & & \\ & & 0 & \\ & & & v^2 \\ & & & & v^2 \end{pmatrix}, \quad (2)$$

where $v^2 = mM/2$. At the vacuum, the $SU(N_c) \times SU(5)_{\text{GUT}}$ gauge symmetry is broken down to $SU(N_c - 2) \times SU(3)_C \times SU(2)_L \times U(1)_Y$ for $N_c \geq 3$ and the electroweak $SU(2)_L \times U(1)_Y$ is the diagonal subgroup of those in $SU(N_c)$ and $SU(5)_{\text{GUT}}$. Note that the vanishing components are not a consequence of the fine-tuning. The corresponding components in Q and \bar{Q} are charged under the unbroken gauge symmetry and that ensures the absence of linear terms in the potential, i.e., stable (or flat directions). Since the low energy $SU(2)_L \times U(1)_Y$ partly comes from $SU(N_c)$, two of the components in T and \bar{T} transform in exactly the same way as the Higgs fields in low energy whereas there is no colored component in T or \bar{T} . Therefore the double-triplet splitting problem and proton decay mediated by the colored Higgs are absent. The rest of the components in T and \bar{T} are fundamental and anti-fundamental representations of $SU(N_c - 2)$ and charged under $U(1)_Y$. All the components in Q and \bar{Q} are either eaten by gauge fields of the broken symmetry or obtain masses from the superpotential. The fate of the exotic particles in T and \bar{T} depends on the dynamics of the unbroken $SU(N_c - 2)$ group below the dynamical scale. Of course, with $\Lambda \ll v$, this is not a ‘‘unified’’ model. The three gauge coupling constants do not meet at the GUT scale since the embeddings of $SU(3)_C$ and $SU(2)_L \times U(1)_Y$ are different. The real unification picture arises when $\Lambda \geq v$, where quantum effects are important.

At the quantum level, the low energy physics is not very different, but some of the vacua are lifted. In particular, it is interesting to note that dynamical symmetry breaking has to happen once we take into account quantum effects [12]. With the above superpotential, the low energy theory of the $\text{rank}(Q\bar{Q}) = 0$ vacuum is $SU(N_c)$ with one flavor T and \bar{T} , which does not have the ground state [21]. The stability of the classical vacua of $\text{rank}(Q\bar{Q}) = 2$ depends on N_c . We show below that there is no N_c which is viable for low energy phenomenology.

For $N_c = 2$, the theory is not asymptotically free and the classical analysis is valid in low energy. However, in this case, the vacuum in Eq. (2) breaks the gauge symmetry into $SU(3) \times SU(2)$ which is not acceptable.

For $N_c = 3$, the quantum effect is easier to analyze in the dual gauge theory [22]. It is again an $SU(3)$ gauge theory but with a superpotential:

$$W = m \text{Tr} M_{Q\bar{Q}} - \frac{1}{M} \text{Tr}(M_{Q\bar{Q}} M_{Q\bar{Q}}) + \dots + \frac{1}{\Lambda} \bar{q} M_{Q\bar{Q}} q + \frac{1}{\Lambda} H \bar{q} t + \frac{1}{\Lambda} \bar{H} q \bar{t} + \frac{1}{\Lambda} S t \bar{t}, \quad (3)$$

where mesons are identified with the quark bilinears in the original (electric) theory: $M_{Q\bar{Q}} \sim Q\bar{Q}$, $H \sim Q\bar{T}$, $\bar{H} \sim \bar{Q}T$, and $S \sim T\bar{T}$. These mesons are singlets under $SU(3)$ and transforms as $\mathbf{1} + \mathbf{24}$, $\mathbf{5}$, $\bar{\mathbf{5}}$, and $\mathbf{1}$ under $SU(5)_{\text{GUT}}$, respectively. Dual quarks q , t and antiquarks \bar{q} , \bar{t} transform as q : $(\mathbf{3}, \mathbf{5})$, t : $(\mathbf{3}, \mathbf{1})$, \bar{q} : $(\mathbf{3}, \bar{\mathbf{5}})$, and \bar{t} : $(\mathbf{3}, \mathbf{1})$ under $SU(3) \times SU(5)_{\text{GUT}}$. The parameter $\hat{\Lambda}$ has dimension one since mesons have dimension two. Amazingly, this is almost identical to the model proposed in Ref. [11]. In the vacuum of our interest:

$$M_{Q\bar{Q}} = \begin{pmatrix} 0 & & & \\ & 0 & & \\ & & 0 & \\ & & & v^2 \\ & & & & v^2 \end{pmatrix}, \quad (4)$$

$SU(5)_{\text{GUT}}$ is broken down to the standard model gauge group, and two of the dual quarks in q and \bar{q} obtain masses. After integrating out the massive quarks, the theory becomes an $SU(3)$ gauge theory with four flavors, which is a confining theory with a superpotential [21]:

$$W = m \text{Tr} M_{Q\bar{Q}} - \frac{1}{M} \text{Tr}(M_{Q\bar{Q}} M_{Q\bar{Q}}) + \dots + \frac{1}{\Lambda} M_{Q\bar{Q}}^{(3 \times 3)} M_{q\bar{q}}^{(3 \times 3)} + \frac{1}{\Lambda} H_C \bar{H}'_C + \frac{1}{\Lambda} \bar{H}_C H'_C + \frac{1}{\Lambda} S S' - \frac{1}{v^2 \hat{\Lambda}} H_D \bar{H}_D S' + \dots + \frac{\det M^{(4 \times 4)}}{\hat{\Lambda}^5} + \frac{1}{\hat{\Lambda}^5} B_{(4)} M^{(4 \times 4)} \bar{B}_{(4)}. \quad (5)$$

The superscript (3×3) represents the 3×3 meson matrix made of colored parts of $Q\bar{Q}$ or $q\bar{q}$. The fields H' , \bar{H}' , and S' are mesons made of dual quarks; $H' \sim q\bar{t}$, $\bar{H}' \sim \bar{q}t$, and $S' \sim t\bar{t}$. The subscripts C and D represent the $SU(3)_C$ colored and the $SU(2)_L$ doublet part of the corresponding meson fields, respectively. $M^{(4 \times 4)}$ is the matrix:

$$M^{(4 \times 4)} = \begin{pmatrix} M_{q\bar{q}}^{(3 \times 3)} & H'_C \\ \bar{H}'_C & S' \end{pmatrix}, \quad (6)$$

and the baryons are

$$B_{(4)} = \begin{pmatrix} B_C \\ B^- \end{pmatrix}, \quad \bar{B}_{(4)} = \begin{pmatrix} \bar{B}_C \\ B^+ \end{pmatrix}, \quad (7)$$

with the identification of $B_C \sim qqt$, $B^- \sim qq\bar{q}$, $\bar{B}_C \sim \bar{q}\bar{q}\bar{t}$, and $B^+ \sim \bar{q}\bar{q}\bar{q}$. Finally $\hat{\Lambda}$ is the dynamical scale of the four-flavor SU(3) theory.

Note here that the doublet Higgses H_D and \bar{H}_D do not obtain a mass term as long as $S' = 0$ whereas the colored Higgses H_C and \bar{H}_C already have mass terms accompanied with the dual mesons H'_C and \bar{H}'_C . Indeed $S' = 0$ is ensured by the condition of $F_S = 0$. The missing partner for the doublets, say H'_D and \bar{H}'_D , have masses since the dual quarks are massive. This is the realization of the doublet-triplet splitting with an infinite number of particles. If we define the T -number with the charge assignment $Q: 0, \bar{Q}: 0, T: +1$, and $\bar{T}: +1$, the Higgs fields H and \bar{H} have the same T -number, $+1$. In this case, an infinite tower of the Higgs fields should be necessary as in Fig. 1 (left). We need an extra Higgs field with charge -1 to make the colored Higgs massive, but it introduces an additional unwanted massless doublet. Repeating this procedure forever is the only possibility of realizing doublet-triplet splitting in this situation. These infinite particles are realized here by the hadron tower of the Higgs fields. The mismatching of the level, i.e., no zero mode only for doublet part of H' and \bar{H}' , happened because of the mass of the constituent quarks q_D and \bar{q}_D without violating T -number. A schematic picture of the hadron spectrum is shown in Fig. 1 (right).

In fact, in this SU(3) model, the situation is a little bit different from the story in the introduction. If the T -number violating term is absent in the superpotential, the dangerous dimension-five proton-decay operators are forbidden. However, since the T -number is anomalous with respect to the SU(3) gauge theory, the nonperturbatively generated superpotential in Eq. (5) violates the T -number. As we see later, $M_{q\bar{q}}^{(3 \times 3)}$ acquires a VEV and it gives T -number violating masses to the colored-Higgs fields. Consequently, dimension-five proton-decay opera-

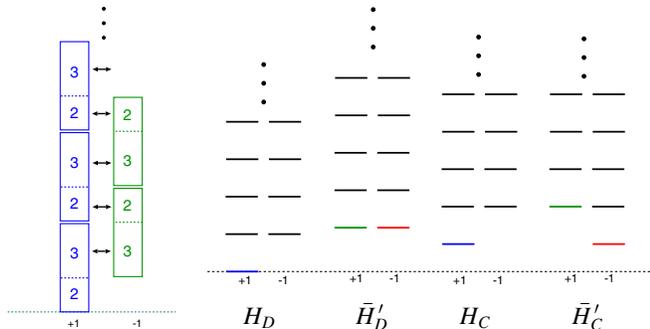


FIG. 1 (color online). A schematic view of the hadron spectrum of the model. Only the doublet part H_D remains massless whereas the colored Higgs H_C have masses by pairing up with \bar{H}'_C .

tors are generated by the colored-Higgs-exchange diagrams as usual.

The stability of the vacuum can be checked by solving the $F = 0$ conditions for all the fields. We can find a solution with Eq. (4) and

$$M_{q\bar{q}}^{(3 \times 3)} = \begin{pmatrix} -m\hat{\Lambda} & & \\ & -m\hat{\Lambda} & \\ & & -m\hat{\Lambda} \end{pmatrix}, \quad (8)$$

$$\frac{S}{\hat{\Lambda}} + \frac{B^+ B^-}{\hat{\Lambda}^5} = -\frac{m^3 \hat{\Lambda}^3}{\hat{\Lambda}^5}. \quad (9)$$

The second equation indicates that the vacuum is not uniquely determined and there is a flat direction. Correspondingly, there are massless particles B^+ and B^- which are charged under $U(1)_Y$. This is the same situation as in the classical analysis, in which the remaining gauge symmetry is just the standard model gauge group but a pair of $U(1)_Y$ charged particles originate from T and \bar{T} are left massless in addition to the Higgs doublets. In order to avoid the massless charged baryons, we need to add a mass term for T and \bar{T} , but that also makes the Higgs doublets heavy. Therefore, the $N_c = 3$ case is not acceptable.

For $N_c = 4$, there is a hope that the exotic states confine and form standard model singlet states so that it is phenomenologically viable. However, unfortunately, it is not the case. As in the case of $N_c = 3$, we can analyze the model by taking a dual gauge group and integrate out the heavy flavors. For $N_c = 4$, the dual theory becomes SU(2) with four flavors. As expected, doublet-triplet splitting happens in the same way as above. By taking the dual again and going back to the electric theory, we find another SU(2) theory with four flavors and superpotential interactions. The solution of the $F = 0$ equations can be found with the same vacuum in Eqs. (4) and (8), which gives mass terms for all the quarks except for T and \bar{T} . The low energy theory in this vacuum is then an SU(2) gauge theory with one flavor which has no ground state. Although we could avoid the charged exotic state, the vacuum is lifted at the quantum level.

For $N_c \geq 5$, the models are confining theories and acquire a nonperturbatively generated superpotential. By the effect of the superpotential, there is no ground state corresponding to the vacuum with Eq. (4).

In summary, there is no phenomenologically viable model with the particle content in Table I and superpotential in Eq. (1), although the doublet-triplet splitting problem is solved in a simple way. This result motivates us to consider the case with a different type of groups such as $Sp(N_c)$ and $SO(N_c)$.

In fact, there is another interesting way of realizing massless doublet Higgs fields in this class of models. If we impose a global SU(6) symmetry in the superpotential where $Q \equiv (Q, T)$ and $\bar{Q} \equiv (\bar{Q}, \bar{T})$ transform as **6** and $\bar{\mathbf{6}}$

and if the global symmetry is broken down to $SU(4) \times SU(2) \times U(1)$, a pair of doublet Higgs fields is ensured to be massless since these are pseudo-Goldstone particles [16,23,24]. With the similar superpotential:

$$W = m \text{Tr}(Q \bar{Q}) - \frac{1}{M} \text{Tr}[(Q \bar{Q})(Q \bar{Q})] + \dots, \quad (10)$$

the mechanism should work and the unwanted exotic particles can be massive by the superpotential terms if such a vacuum exists. Although it is an interesting possibility, we do not pursue this direction further in this paper partly because it is incompatible with the later discussion of supersymmetry breaking.

B. $Sp(N_c)$ and $SO(N_c)$ models

We can indeed find viable models for $Sp(N_c)$ and $SO(N_c)$. We show the result of the analysis for these cases. The particle content of the $Sp(N_c)$ models is listed in Table II where the field T_2 is necessary to avoid the Witten anomaly [25]. We assumed the same superpotential for Q and \bar{Q} as in Eq. (1) with the matrix $(Q\bar{Q})$ being the $Sp(N_c)$ singlet combination with 5×5 flavor indices this time. The analysis can go through in the similar fashion to the $SU(N_c)$ case, and we find that only the $N_c = 2$ case has a stable minimum with Eq. (4). The massless modes of the $Sp(2)$ model consist of four Higgs doublets, $H_D \sim Q_D T_1$, $\bar{H}_D \sim \bar{Q}_D T_1$, $H_{D2} \sim Q_D T_2$, and $\bar{H}_{D2} \sim \bar{Q}_D T_2$. Two of them H_{D2} and \bar{H}_{D2} can be made massive by adding a superpotential term $W \ni (QT_2)(\bar{Q}T_2)$ without giving a mass for H_D and \bar{H}_D . This is the model found in Ref. [15]. Similar to the $SU(N_c)$ case, $Sp(N_c)$ models with $N_c \geq 3$ do not have a vacuum with $\text{rank}(M_{Q\bar{Q}}) = 2$ due to the nonperturbatively generated superpotential [26].

The $SO(N_c)$ model can also be constructed, and turns out to be the most interesting case. A detailed analysis will be presented in the next section. The particle content is given in Table III where we have to introduce only one T field in contrast to the case of $SU(N_c)$ or $Sp(N_c)$. Again, the form of the superpotential is the same as that in Eq. (1). With the same analysis, we find that there are stable vacua for $4 \leq N_c \leq 9$ with massless Higgs-doublet fields and there are no unwanted massless fields at low energy. The theory is asymptotically free for $N_c \geq 6$. Additional singlet fields under the standard model gauge group appear for $N_c \geq 5$.

TABLE II. The particle content of the $Sp(N_c)$ model.

	$Sp(N_c)$	$SU(5)_{\text{GUT}}$
Q	$2N_c$	5
\bar{Q}	$2N_c$	$\bar{\mathbf{5}}$
T_1	$2N_c$	1
T_2	$2N_c$	1

TABLE III. The particle content of the $SO(N_c)$ model.

	$SO(N_c)$	$SU(5)_{\text{GUT}}$
Q	N_c	5
\bar{Q}	N_c	$\bar{\mathbf{5}}$
T	N_c	1

III. $SO(9)$ MODEL

We discuss in more detail the most interesting model among those in the previous section: the $SO(9) \times SU(5)_{\text{GUT}}$ model. The phenomenological aspects of the model such as gauge coupling unification, proton decay, and Yukawa interactions for matter fields will be addressed. Many of these features are shared with the model of Ref. [15].

The particle content and the tree-level superpotential are defined in the previous section in Table III and in Eq. (1). The model is an $SO(9)$ gauge theory with 11 flavors which is in the conformal window [27]. This fact becomes important for the discussion of the phenomenological issues. We take a picture in which $\Lambda \gg \nu$ where Λ is the scale where the $SO(9)$ gauge theory flows into the fixed point. Since confinement does not happen until the fields decouple, the actual confinement scale coincides with the GUT scale ν which is set by the parameters m and M . Therefore we have an energy region with a CFT between $M_G \sim \nu$ and Λ .

The analysis of the vacuum can be done along the same line in the $SU(N_c)$ case. We first take a dual picture of the theory which is an $SO(6)$ gauge theory with 11 flavors with a superpotential [27]:

$$\begin{aligned}
 W = & m \text{Tr} M_{Q\bar{Q}} - \frac{1}{M} \text{Tr}(M_{Q\bar{Q}} M_{Q\bar{Q}}) + \dots + \frac{1}{\Lambda} \bar{q} M_{Q\bar{Q}} q \\
 & + \frac{1}{\Lambda} \bar{q} M_{QQ} \bar{q} + \frac{1}{\Lambda} q M_{\bar{Q}\bar{Q}} q + \frac{1}{\Lambda} H \bar{q} t + \frac{1}{\Lambda} \bar{H} q t \\
 & + \frac{1}{\Lambda} S t t, \quad (11)
 \end{aligned}$$

where $M_{Q\bar{Q}}$, M_{QQ} , and $M_{\bar{Q}\bar{Q}}$ are mesons made of Q and \bar{Q} which are singlet under $SO(6)$ and **1** + **24**, **15**, and $\bar{\mathbf{15}}$ under $SU(5)_{\text{GUT}}$. The mesons involving T are $H \sim QT$, $\bar{H} \sim \bar{Q}T$, and $S \sim TT$, and these are **5**, $\bar{\mathbf{5}}$, and **1** under $SU(5)_{\text{GUT}}$. The dual quarks q , \bar{q} , and t transform as q : (**6**, **5**), \bar{q} : (**6**, $\bar{\mathbf{5}}$), and t : (**6**, **1**) under $SO(6) \times SU(5)_{\text{GUT}}$. Again $\hat{\Lambda}$ is introduced such that dimensionality of the superpotential is correct. It is interesting that this dual picture is similar to the $SO(6) \times SO(10)_{\text{GUT}}$ model proposed in Ref. [12]. By the VEV of $M_{Q\bar{Q}}$ in Eq. (4), four flavors (doublet part of q and \bar{q}) obtain masses and the low energy theory becomes an $SO(6)$ theory with seven flavors (colored part of q and \bar{q} and t). This is still an interacting theory.

When we take the dual again, the theory now comes back to the original electric theory but the gauge group is reduced to $SO(5)$. The superpotential is

$$\begin{aligned}
 W = & m \text{Tr} M_{Q\bar{Q}} - \frac{1}{M} \text{Tr}(M_{Q\bar{Q}} M_{Q\bar{Q}}) + \dots \\
 & + \frac{1}{\tilde{\Lambda}} M_{Q\bar{Q}}^{(3\times 3)} M_{q\bar{q}}^{(3\times 3)} + \dots - \frac{1}{v^2 \tilde{\Lambda}} H_D \bar{H}_D S' \\
 & + \frac{1}{\tilde{\Lambda}} H_C \bar{H}'_C + \frac{1}{\tilde{\Lambda}} \bar{H}_C H'_C + \frac{1}{\tilde{\Lambda}} S S' - \frac{1}{\tilde{\Lambda}} \tilde{Q}_C M_{q\bar{q}}^{(3\times 3)} \tilde{Q}_C \\
 & + \dots - \frac{1}{\tilde{\Lambda}} \tilde{Q}_C H'_C \tilde{T} - \frac{1}{\tilde{\Lambda}} \tilde{Q}_C \bar{H}'_C \tilde{T} - \frac{1}{\tilde{\Lambda}} S' \tilde{T} \tilde{T}, \quad (12)
 \end{aligned}$$

where we wrote down only terms relevant for the discussion. The quarks \tilde{Q} , \tilde{Q} , and \tilde{T} will be identified with original quarks upon integrating out the massive fields. The doublet-triplet splitting happens in the same way as the $SU(N_c)$ examples. The massless doublet is obtained with $S' = 0$ which is ensured by the $F_S = 0$ condition, and the triplets get masses by pairing with dual mesons. In contrast to the case of the $SU(3)$ model in the previous section, there is no nonperturbatively generated superpotential. Therefore, the symmetry of the superpotential at the classical level, the T -number ($T: +1$), is respected. This is exactly the situation discussed in the introduction. This fact becomes important in the discussion of the proton decay.

By solving the $F = 0$ conditions, we can find a solution in Eqs. (4) and (8) that gives a mass term to \tilde{Q}_C and \tilde{Q}_C . After integrating out the heavy fields, the theory ends up with an $SO(5)$ theory with one flavor \tilde{T} without a superpotential, and we have massless Higgs-doublet fields. This result is the same as the classical analysis in the Higgs phase.

The decoupling of the fields immediately make the $SO(5)$ interaction strong and the quark \tilde{T} confines. This $SO(5)$ one flavor theory has been known to have two branches [27]. In one branch, a superpotential $W = (\tilde{\Lambda}^8/\tilde{S})^{1/2}$ is generated with $\tilde{\Lambda}$ being the dynamical scale of one flavor $SO(5)$ theory and $\tilde{S} \sim \tilde{T} \tilde{T}$. This branch is unacceptable because there is no ground state. In the other branch, no superpotential is generated and there is no singularity at the origin of the meson \tilde{S} even though the gauge symmetry is enhanced there at the classical level. Therefore there is a stable vacuum in this branch. The low energy spectrum is just a pair of doublet Higgs fields with a massless meson \tilde{S} and the superpotential is $W = 0$.

Now we start to discuss the phenomenological issues. First, we need to check whether gauge coupling unification is maintained in this model. The first order answer to this question is yes. There are no exotic massless fields in the spectrum, and the running of the gauge coupling constants are the same as that of the minimal supersymmetric standard model. However, the threshold correction is also important for precise unification, which is not obvious.

The question depends on the spectrum of the heavy fields and that cannot be estimated without the knowledge of the Kähler potential. However, the qualitative discussion is still possible by assuming that the Kähler potential is not very different from the classical one. In this case, the mass spectrum can be estimated by explicitly calculating the mass of the components in Q and \bar{Q} at the classical level. There are three classes of fields: fields eaten by the $SO(9)/SO(5)$ gauge fields, ones eaten by the $SU(5)_{\text{GUT}}/(SU(3)_C \times SU(2)_L \times U(1)_Y)$ gauge fields, and others that obtain masses by the superpotential. The first two classes of fields have masses of order $v = \sqrt{mM}/2$ and the last ones have $O(m)$. Therefore, in order not to destroy gauge coupling unification due to the mass splitting between those two, the mass parameter m is required to be around the GUT scale $M_G \sim v$ which means the scale M should also be of the order of the GUT scale.

This sounds unreasonable since we expect that the higher dimensional operators are suppressed by the Planck scale M_{Pl} that is 2 orders of magnitude larger than the GUT scale. Also, with GUT scale suppressed operators, we are not allowed to discuss the high energy theory above the GUT scale, which we are doing. However, in this model, the mechanism of suppressing M is already built in. As we discussed before, this $SO(9)$ theory is in the conformal window, and we expect an energy range of CFT above the GUT scale. In this case, the meson fields have a large negative anomalous dimension which enhances the couplings in low energy. In other words, the interaction never gets strong at high energy. The anomalous dimension of the meson fields are calculated by using the relation between the nonanomalous R -charge and the dimension of the operator as follows [22]:

$$\gamma(Q\bar{Q}) = D(Q\bar{Q}) - 2 = \frac{3}{2}R(Q\bar{Q}) - 2 = -\frac{10}{11}. \quad (13)$$

With this anomalous dimension, the coefficient $1/M$ enhances almost quadratically with the scale towards low energy. Therefore it is natural to have a large enhancement. If we assume that the original operator is suppressed by the Planck scale M_{Pl} , the factor of 100 enhancement is easily realized by a small CFT range such as from 10^{16} GeV to 10^{17} GeV. In the same way, the coefficient m enhances almost linearly with the energy scale towards low energy. Therefore, the original scale of the model m evaluated at the Planck scale was smaller than the GUT scale by a factor of 10 or so. This is an interesting scale for the right-handed neutrino masses in the seesaw model [28].

The gauge coupling of $SU(5)_{\text{GUT}}$ above the GUT scale can be in the perturbative region all the way up to the Planck scale even accounting for the large anomalous dimension of Q and \bar{Q} [29].

The Yukawa coupling constants between matter and the Higgs fields originate from higher dimensional operators since the Higgs fields are identified with the meson fields. The gauge invariant terms:

$$W_{\text{Yukawa}} = \frac{f_u}{M_Y} (\mathbf{10})(\mathbf{10})(QT) + \frac{f_d}{M_Y} (\mathbf{10})(\bar{\mathbf{5}})(\bar{Q}T), \quad (14)$$

become the Yukawa interactions at low energy. The matter fields are represented by $(\mathbf{10})$ and $(\bar{\mathbf{5}})$. The low energy Yukawa coupling is roughly $y_u \sim f_u M_G / M_Y$ and $y_d \sim f_d M_G / M_Y$, where M_G is the GUT scale. The Yukawa coupling constant for the top quark is necessary to be $O(1)$, which again requires the scale M_Y to be the GUT scale. This is not a problem for the same reason as before. These operators are almost marginal operators and thus the coefficient is enhanced linearly in low energy by the large anomalous dimensions.

After integrating out the massive fields with the Yukawa interactions in Eq. (14), the final low energy effective superpotential is:

$$W = W_{\text{MSSM}} + \frac{y_u y_d}{m} \frac{\tilde{S}}{M_G^2} \times (\hat{Q} \hat{Q} \hat{Q} \hat{L} + \hat{U} \hat{U} \hat{D} \hat{E} + \hat{Q} \hat{Q} \hat{U} \hat{D} + \hat{U} \hat{E} \hat{Q} \hat{L}), \quad (15)$$

with \hat{Q} , \hat{U} , \hat{D} being the quark superfields and \hat{L} , \hat{E} are the lepton superfields. In addition to the usual Yukawa interactions W_{MSSM} , the baryon-number-violating terms (the first two terms in the parenthesis) appeared. However, as long as \tilde{S} is stabilized near the origin, these terms do not cause rapid proton decay. The particle content of the low energy effective theory is just that of the minimal supersymmetric standard model and a singlet field \tilde{S} which only couples to the higher dimensional operators. The value of \tilde{S} depends on the shape of the Kähler potential and how supersymmetry is broken. We will discuss these in the next section.

We briefly mention the case with $N_c = 6, 7$, and 8 , where the vacuum with Eq. (4) exists and the doublet-triplet splitting happens for these cases. All of the models are in the conformal window. For $N_c = 6$, the low energy effective theory has an extra U(1) gauge symmetry with two charged (but the standard model singlet) fields t^+ and t^- . Since this theory is at the edge of the conformal window (barely asymptotically free), we do not expect the large enhancement of the Yukawa coupling constants. For $N_c = 7$, the low energy theory again has an extra U(1) factor. There are three standard model singlet fields S , t^+ , and t^- , where t^+ and t^- are monopoles, with superpotential $W = f(S) S t^+ t^-$. The unknown function $f(S)$ is non-vanishing at the origin. For $N_c = 8$, there is no extra gauge symmetry and there are two branches as in the $N_c = 9$ case. In one branch, a stable vacuum exists with a superpotential term $W = f(x) S t t$ where S and t are the standard model singlet fields. Again, $f(x)$ ($x = S^2 t^2$) is an unknown function but is nonzero at the origin.

IV. μ -TERM DRIVEN SUPERSYMMETRY BREAKING

In the previous section, we have seen that the SO(9) model is quite successful in obtaining massless Higgs fields in the low energy spectrum. However, to be phenomenologically completely viable, the Higgs fields have to have a mass term of $O(100 \text{ GeV})$, the μ -term, otherwise Higgsinos become massless which is excluded by experiment. As we see below, it is possible to obtain a μ -term by giving a small mass term for T , but it causes a dramatic effect in the dynamical system: dynamical supersymmetry breaking [30,31].

The μ -term can be obtained by adding a superpotential term:

$$W \ni \hat{\mu} T T, \quad (16)$$

which becomes $\hat{\mu} S$ in Eq. (12). With this term, the $F_S = 0$ condition leads to

$$S' = -\hat{\mu} \hat{\Lambda}, \quad (17)$$

and it induces a mass term for the Higgs doublets:

$$W \ni \frac{\hat{\mu}}{v^2} H_D \bar{H}_D. \quad (18)$$

In terms of the canonically normalized fields \hat{H}_D and $\hat{\bar{H}}_D$, this is nothing but the μ -term, $\mu \hat{H}_D \hat{\bar{H}}_D$ with $\mu \sim \hat{\mu}$. It is obvious that the mesons made of T become massive once we introduce the mass term for the quark T . This is also easy to understand in the classical level analysis. Since the Higgs fields are simply the components in T , the term in Eq. (16) is the μ -term.

The term in Eq. (16) also gives a potential term for \tilde{S} . With Eq. (17) and the superpotential in Eq. (12), \tilde{T} obtains a term in low energy:

$$\hat{\mu} \tilde{T} \tilde{T}, \quad (19)$$

and after confinement of the SO(5) gauge theory, it becomes a linear superpotential for the meson $\tilde{S} \sim \tilde{T} \tilde{T}$:

$$W \ni \hat{\mu} \tilde{S}. \quad (20)$$

If we ignore the higher dimensional operators in Eq. (15), there is no solution for $F_{\tilde{S}} = 0$.

Interestingly, this does not mean that the vacuum is destabilized or quarks and leptons must condense by the presence of the small μ -term. First, we start the discussion by ignoring the higher dimensional operators. In this case, it was shown by Intriligator, Seiberg, and Shih [19] that the vacuum is metastable and supersymmetry is broken there. The argument is pretty easy. Whether the vacuum is stable or not depends on the shape of the Kähler potential for \tilde{S} , but since we know that for large \tilde{S} , where the classical analysis is valid, the potential grows by the mass term. Therefore, there must be a local minimum somewhere. The true supersymmetric vacua exists in the other branch where

\tilde{S} is stabilized far away from the origin $M_G(M_G/\hat{\mu})^{2/3}$ (this is meaningless because it is much larger than the Planck scale) and also at different vacua from that in Eq. (4). Since the true vacua exist far away or the energy difference between the true vacuum and the metastable one is much smaller than the height of the potential barrier of $O(M_G^4)$, we expect that the lifetime of this metastable vacuum is long enough for us [32].

Once we include the higher dimensional operators, another supersymmetric vacuum appears where quarks and leptons acquire VEVs. However, if we assume that the scalar components of quarks and leptons obtain positive supersymmetry breaking mass-squared terms, the vacuum is again metastable and its lifetime is very long since the peak of the potential barrier is located far from the origin $Q \sim L \sim (\mu M_G)^{1/2}$ compared to the height of the potential $V^{1/4} \sim O(\mu^3 M_G)^{1/4}$ [33]. Therefore, we conclude that there is a supersymmetry breaking metastable vacuum.

The size of supersymmetry breaking is $F_{\hat{S}} \sim O(\mu M_G)$ with a canonically normalized field $\hat{S} \sim M_G \tilde{S}$. For the Higgs fields, terms in the Kähler potential such as $\hat{S}^\dagger \hat{S} H^\dagger H / \tilde{\Lambda}^2$ are expected to be generated by the non-perturbative effect and supersymmetry breaking can be mediated directly (in the sense of gravity mediation). In this case, the soft scalar masses for the Higgs fields are obtained with a similar size to the μ -term. If the term $\hat{S} H^\dagger H / \tilde{\Lambda} + \text{H.c.}$ is generated, which should be possible since there is no unbroken symmetry to protect the term, the trilinear A and the bilinear B -term is also nonvanishing and of the same order with the μ -term. If the matter fields in the third generation couple to the Higgs fields strongly, the soft scalar masses for those fields can also be obtained directly.

Gauge mediation [34] also happens if \hat{S} is stabilized away from the origin, where the colored-Higgs fields play a role of the messenger field.¹ By integrating out \tilde{Q}_C \tilde{Q}'_C in Eq. (12), we obtain mass terms for the colored-Higgs fields:

$$W \simeq \hat{S} \hat{H}'_C \hat{H}'_C + M_C \hat{H}_C \hat{H}'_C + M_C \hat{H}_C \hat{H}'_C, \quad (21)$$

where M_C is the colored-Higgs mass of order M_G and meson fields are canonically normalized. Unfortunately, with this structure of superpotential, the leading order contribution to the gaugino masses of $O(F_{\hat{S}}/\langle \hat{S} \rangle)$ cancels out [36],² and moreover there is no contribution to the $SU(2)_L$ gauginos.

In order to obtain gaugino masses, there must be a gauge-mediation effect since the supersymmetry breaking scale is too low ($O(\mu M_G)$) for gravity mediation. A simple example for gauge mediation is to assume an interaction term:

¹A similar structure of the model can be found in Ref. [35].

²I thank Y. Nomura for discussion on this point.

$$W_{\text{messenger}} = \frac{1}{M_X} T^2 \Phi \bar{\Phi}, \quad (22)$$

where Φ and $\bar{\Phi}$ transform under the $SU(5)_{\text{GUT}}$ such as $\mathbf{5}$ and $\bar{\mathbf{5}}$ and the singlet under $SO(9)$. By the enhancement of the $1/M_X$ suppressed term due to the large anomalous dimension of T^2 , this term effectively becomes

$$W_{\text{messenger}} \rightarrow \lambda \hat{S} \Phi \bar{\Phi} \quad (23)$$

with $\lambda \sim O(0.1 - 1)$ even if M_X is $O(M_{\text{Pl}})$. With this superpotential, another supersymmetric true vacuum with $\Phi = \bar{\Phi} \neq 0$ appears. However, assuming that the \hat{S} field is stabilized away from the origin, which is reasonable since we expect the presence of a linear term in the Kähler metric, the vacuum with $\Phi = \bar{\Phi} = 0$ is metastable.³ At the metastable vacuum the gaugino masses are generated to be

$$m_{1/2} = \frac{\alpha}{4\pi} \frac{F_{\hat{S}}}{\langle \hat{S} \rangle}. \quad (24)$$

In order for the gaugino masses to be similar to μ in size, a somewhat small value of $\langle \hat{S} \rangle \sim 10^{14-15}$ GeV is necessary. That is consistent with the suppression of the coefficient of proton-decay operators in Eq. (15).

It might be possible and would be great if all the gaugino masses are obtained by (really) direct gauge mediation [38] without having the messenger particles above by extending the gauge group and/or matter content. But in any case, the pattern of the supersymmetry breaking parameters in this scenario is essentially of gauge-mediation type except for the Higgs sector, since the Higgs fields can feel the supersymmetry breaking directly. The soft scalar masses $m_{H_u}^2$, $m_{H_d}^2$, μ , B , and A -terms can be taken as free parameters at the GUT scale. (Probably $m_{H_u}^2 \simeq m_{H_d}^2$ because of the parity symmetry $Q \leftrightarrow \bar{Q}$ which is only broken by the Yukawa interactions.) The soft masses for the third generation fields may also be modified depending on the size of the Yukawa coupling constants. This prediction should be testable at future colliders. The source of flavor and CP violation in this model is only in the Yukawa coupling constants, which is the desired situation taking into account the stringent constraints on the supersymmetry breaking parameters. The gravity-mediation effect gives an $O(1\%)$ correction to the parameters. This is interesting for the detection of the flavor and CP violating processes.⁴

³Since the linear term in the Kähler metric has nonvanishing R -charge, it is suppressed when the explicit R -symmetry breaking by the superpotential is small, i.e., $m \ll \tilde{\Lambda}$. Even in this case, the supergravity effect shifts the vacuum to $\langle \hat{S} \rangle \sim \tilde{\Lambda}^2 / M_{\text{Pl}} \sim 10^{14}$ GeV, which is numerically consistent with the phenomenological requirements. See [37] for detailed discussion.

⁴One should take into account the possible conformal sequestering effect [39] to the gravity-mediated contribution.

The lightest supersymmetric particle is the gravitino. The gravitino mass is estimated to be

$$m_{3/2} = \frac{F_{\hat{S}}}{\sqrt{3}M_{\text{Pl}}} \sim \mu \left(\frac{M_G}{M_{\text{Pl}}} \right) \sim O(1) \text{ GeV}. \quad (25)$$

This mass range is interesting for cosmology [40] and also for collider experiments [41].

For other choices of N_c , the situation is different. For $N_c = 6$, the addition of the μ -term just gives a mass term for t^+ and t^- which does not cause supersymmetry breaking. For $N_c = 7$ and 8, the term μS appears but the supersymmetry unbroken vacuum exists where t^\pm (or t for $N_c = 8$) fields acquire nonvanishing VEV. However, it is possible that the S is stabilized far from the origin where t^\pm or t is heavy, and supersymmetry is broken there.

V. COSMOLOGICAL CONSTANT DRIVEN SUPERSYMMETRY BREAKING

Because of the large anomalous dimension, the size of the μ -parameter is originally smaller than $O(100 \text{ GeV})$ by a factor of 10 or so. This is about the size of the gravitino mass. Therefore it is possible that the origin of the μ -term can actually be the cosmological constant by the Giudice-Masiero mechanism [18]. Assuming a presence of a T^2 ($\sim \tilde{\Lambda} \hat{S}$) term in the Kähler potential, this effectively becomes the $\hat{\mu}$ -term in the superpotential. It is equivalent to study a model with

$$K = \hat{S}^\dagger \hat{S} + a \tilde{\Lambda} (\hat{S} + \text{H.c.}) - \frac{(\hat{S}^\dagger \hat{S})^2}{\tilde{\Lambda}^2} + \dots \quad (26)$$

and $W = c$,

where c is a constant term which represents the negative cosmological constant term, $V_{\text{AdS}} \simeq -3|W|^2$, of order $c^2 \sim m_{3/2}^2$ in the unit of $M_{\text{Pl}} = 1$. This term is always necessary to cancel the positive vacuum energy from the supersymmetry breaking, $V_F \simeq |F|^2$. The parameters have a hierarchical structure: $a \gg 1$, $\tilde{\Lambda} \ll 1$, and $c \ll 1$, where a represents the enhancement of the coupling through the large anomalous dimension.

By the Kähler transformation, $K \rightarrow K - x - x^\dagger$ and $W \rightarrow W e^x$ with x being a chiral superfield, this is identical to the system:

$$K = \hat{S}^\dagger \hat{S} - \frac{(\hat{S}^\dagger \hat{S})^2}{\tilde{\Lambda}^2} + \dots \quad \text{and} \quad W = c e^{a \tilde{\Lambda} \hat{S}}. \quad (27)$$

By expanding the superpotential, we obtain a $\hat{\mu}$ -term of order $ac \simeq am_{3/2}$. Within the range $|\hat{S}| \lesssim \tilde{\Lambda}$, where the effective theory makes sense, the minimum of the potential exists near the origin $\hat{S} \simeq \tilde{\Lambda}/(4a)$, and the cosmological constant can be cancelled when $a \simeq \sqrt{3}/\tilde{\Lambda} \sim O(100)$. Supersymmetry is broken at the minimum with $F_{\hat{S}} \simeq ac \tilde{\Lambda} \simeq \mu \tilde{\Lambda}$. The value of \hat{S} at the minimum is $O(10^{14} \text{ GeV})$ which is consistent with the required value

for gauge mediation in Eq. (24) and also the suppression of the coefficient of the dimension-five proton-decay operators in Eq. (15).

In the conventional scenario of supersymmetry breaking, some mass scales, such as a dynamical scale, determine the size of $|F|^2$ and the net cosmological constant is cancelled by an independent negative contribution from the c -term. However, in this scenario, the c -term drives supersymmetry breaking, and therefore these are related. In particular, it is interesting that supersymmetry is recovered in the flat limit ($c \rightarrow 0$), resulting in the supersymmetric flat space rather than the supersymmetry broken de Sitter space.

VI. DISCUSSION AND CONCLUSIONS

From the consideration of the mystery of the Higgs particle, we arrived at a rather unified picture. At every stage of the phase transitions, GUT breaking, supersymmetry breaking, and electroweak symmetry breaking, the Higgs field may be playing a crucial role. We have succeeded to construct a realistic GUT model with dynamical symmetry breaking, and found that, in the SO(9) model, the inclusion of the μ -term for the Higgs fields triggers supersymmetry breaking in a metastable vacua by the same dynamics.

We discuss possible generalizations of the model here. Although we discussed the GUT breaking and supersymmetry breaking in a unified picture, we can separately discuss the following mechanisms:

- (i) Doublet-triplet splitting through dynamical GUT breaking,
- (ii) μ -term driven supersymmetry breaking.

Dynamical GUT breaking without supersymmetry breaking is possible. The μ -term can be obtained from a separate SUSY breaking sector by, e.g., the Giudice-Masiero mechanism. The SO(10)_{GUT} extension of the model should be straightforward and is interesting for the discussion of the neutrino masses. Considering different types of particle content and assumption on the superpotential is also worth investigating.

It is possible to obtain massless colored-Higgs fields instead of the doublet fields. In this case, the doublet-triplet splitting can be done by introducing a pair of elementary Higgs fields with the coupling to the meson operators [10–13]. This is possible in the rank($M_{Q\bar{Q}}$) = 3 vacuum with SO(N_c) with $6 \leq N_c \leq 11$.

μ -term driven supersymmetry breaking can be discussed without GUT. For example, in the SO(9) model we can gauge only the SU(2)_L × U(1)_Y subgroup of SU(5)_{GUT}. Then the dynamical scale can be lowered (or even raised) as long as the gauge couplings of the standard model gauge group maintains the perturbativity. Adding a μ -term breaks supersymmetry in the same way but the scale can be much lower (or higher). Gauge mediation through the colored-Higgs fields might be able to be gen-

eralized as a realistic direct gauge-mediation model. We leave those questions for future studies. In any case, the prediction to the low energy spectrum is a modification of the gauge-mediation type in the Higgs sector. In the SO(9) model we presented, there is a cosmological problem associated with the modulus \hat{S} [42]. The mass of \hat{S} is of the same order of μ , i.e., $O(100 \text{ GeV})$ independent of the dynamical scale. A realistic cosmological scenario needs to be considered.

We see that two mechanisms nontrivially fit into a picture: dynamical GUT and supersymmetry breaking. Although it is not likely that we can directly probe the

GUT theory by experiments, the spectrum of the supersymmetric particles in low energy gives us a hint for high energy theories. In this model, the direct connection between the Higgs fields and the supersymmetry breaking sector provides a characteristic feature in the low energy spectrum.

ACKNOWLEDGMENTS

I thank Roni Harnik for reading the manuscript. This work was supported by the U.S. Department of Energy under Contract No. DE-AC02-76SF00515.

-
- [1] S. Dimopoulos and H. Georgi, Nucl. Phys. **B193**, 150 (1981); S. Dimopoulos, S. Raby, and F. Wilczek, Phys. Rev. D **24**, 1681 (1981). N. Sakai, Z. Phys. C **11**, 153 (1981).
- [2] H. Georgi and S.L. Glashow, Phys. Rev. Lett. **32**, 438 (1974).
- [3] S. Weinberg, Phys. Rev. D **26**, 287 (1982); N. Sakai and T. Yanagida, Nucl. Phys. **B197**, 533 (1982).
- [4] T. Goto and T. Nihei, Phys. Rev. D **59**, 115009 (1999); H. Murayama and A. Pierce, Phys. Rev. D **65**, 055009 (2002).
- [5] T. Watari, in Proceedings of the KEK Theory Meeting on Collider Physics, Tsukuba, Japan, 2002 (unpublished).
- [6] Y. Kawamura, Prog. Theor. Phys. **105**, 999 (2001).
- [7] G. Altarelli and F. Feruglio, Phys. Lett. B **511**, 257 (2001); L. J. Hall and Y. Nomura, Phys. Rev. D **64**, 055003 (2001).
- [8] J.M. Maldacena, Adv. Theor. Math. Phys. **2**, 231 (1998); Int. J. Theor. Phys. **38**, 1113 (1999); S.S. Gubser, I.R. Klebanov, and A.M. Polyakov, Phys. Lett. B **428**, 105 (1998); E. Witten, Adv. Theor. Math. Phys. **2**, 253 (1998).
- [9] N. Arkani-Hamed, M. Porrati, and L. Randall, J. High Energy Phys. 08 (2001) 017; R. Rattazzi and A. Zaffaroni, J. High Energy Phys. 04 (2001) 021.
- [10] T. Yanagida, Phys. Lett. B **344**, 211 (1995); T. Hotta, K. I. Izawa, and T. Yanagida, Phys. Rev. D **53**, 3913 (1996); J. Hisano and T. Yanagida, Mod. Phys. Lett. A **10**, 3097 (1995); K. I. Izawa and T. Yanagida, Prog. Theor. Phys. **97**, 913 (1997).
- [11] T. Hotta, K. I. Izawa, and T. Yanagida, Phys. Lett. B **409**, 245 (1997); Prog. Theor. Phys. **95**, 949 (1996).
- [12] T. Hotta, K. I. Izawa, and T. Yanagida, Phys. Rev. D **54**, 6970 (1996).
- [13] K. I. Izawa and T. Yanagida, Prog. Theor. Phys. **99**, 423 (1998).
- [14] H. C. Cheng, Phys. Lett. B **410**, 45 (1997); M. Graesser, Phys. Rev. D **59**, 035007 (1999).
- [15] R. Kitano and G. D. Kribs, J. High Energy Phys. 03 (2005) 033.
- [16] Y. Nomura, D. Poland, and B. Tweedie, hep-ph/0605014.
- [17] A. H. Chamseddine, R. Arnowitt, and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); R. Barbieri, S. Ferrara, and C. A. Savoy, Phys. Lett. B **119**, 343 (1982); L. J. Hall, J. D. Lykken, and S. Weinberg, Phys. Rev. D **27**, 2359 (1983).
- [18] G.F. Giudice and A. Masiero, Phys. Lett. B **206**, 480 (1988); J. A. Casas and C. Munoz, Phys. Lett. B **306**, 288 (1993).
- [19] K. Intriligator, N. Seiberg, and D. Shih, J. High Energy Phys. 04 (2006) 021.
- [20] R. Barbieri, G.R. Dvali, and A. Strumia, Phys. Lett. B **333**, 79 (1994).
- [21] N. Seiberg, Phys. Rev. D **49**, 6857 (1994); K. A. Intriligator, R. G. Leigh, and N. Seiberg, Phys. Rev. D **50**, 1092 (1994).
- [22] N. Seiberg, Nucl. Phys. **B435**, 129 (1995).
- [23] K. Inoue, A. Kakuto, and H. Takano, Prog. Theor. Phys. **75**, 664 (1986).
- [24] H. C. Cheng, Phys. Rev. D **60**, 075015 (1999).
- [25] E. Witten, Phys. Lett. B **117**, 324 (1982).
- [26] K. A. Intriligator and P. Pouliot, Phys. Lett. B **353**, 471 (1995).
- [27] K. A. Intriligator and N. Seiberg, Nucl. Phys. **B444**, 125 (1995).
- [28] T. Yanagida, in Proceedings of the Workshop on the Unified Theory and the Baryon Number in the Universe, Tsukuba, Japan, 1979, edited by O. Sawada and A. Sugamoto; M. Gell-Mann, P. Ramond, and R. Slansky, CERN, Print-80-0576.
- [29] V. A. Novikov, M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Nucl. Phys. **B229**, 381 (1983); M. A. Shifman and A. I. Vainshtein, Nucl. Phys. **B277**, 456 (1986); Zh. Eksp. Teor. Fiz. **91**, 723 (1986) [Sov. Phys. JETP **64**, 428 (1986)].
- [30] E. Witten, Nucl. Phys. **B188**, 513 (1981).
- [31] I. Affleck, M. Dine, and N. Seiberg, Phys. Rev. Lett. **51**, 1026 (1983); Nucl. Phys. **B241**, 493 (1984); **B256**, 557 (1985).
- [32] S. R. Coleman, Phys. Rev. D **15**, 2929 (1977); **16**, 1248(E) (1977); C. G. Callan and S. R. Coleman, Phys. Rev. D **16**, 1762 (1977).
- [33] M. J. Duncan and L. G. Jensen, Phys. Lett. B **291**, 109 (1992).
- [34] M. Dine, A. E. Nelson, and Y. Shirman, Phys. Rev. D **51**,

- 1362 (1995); M. Dine, A.E. Nelson, Y. Nir, and Y. Shirman, Phys. Rev. D **53**, 2658 (1996).
- [35] K. I. Izawa and T. Yanagida, Prog. Theor. Phys. **114**, 433 (2005).
- [36] K. I. Izawa, Y. Nomura, K. Tobe, and T. Yanagida, Phys. Rev. D **56**, 2886 (1997).
- [37] R. Kitano, Phys. Lett. B **641**, 203 (2006).
- [38] E. Poppitz and S. P. Trivedi, Phys. Rev. D **55**, 5508 (1997); L. Randall, Nucl. Phys. **B495**, 37 (1997); N. Arkani-Hamed, J. March-Russell, and H. Murayama, Nucl. Phys. **B509**, 3 (1998); H. Murayama, Phys. Rev. Lett. **79**, 18 (1997).
- [39] M. A. Luty and R. Sundrum, Phys. Rev. D **65**, 066004 (2002).
- [40] J. L. Feng, A. Rajaraman, and F. Takayama, Phys. Rev. Lett. **91**, 011302 (2003); J. L. Feng, S. f. Su, and F. Takayama, Phys. Rev. D **70**, 063514 (2004); J. L. Feng, S. Su, and F. Takayama, Phys. Rev. D **70**, 075019 (2004).
- [41] W. Buchmuller, K. Hamaguchi, M. Ratz, and T. Yanagida, Phys. Lett. B **588**, 90 (2004); K. Hamaguchi, Y. Kuno, T. Nakaya, and M. M. Nojiri, Phys. Rev. D **70**, 115007 (2004); J. L. Feng and B. T. Smith, Phys. Rev. D **71**, 015004 (2005); **71**, 019904(E) (2005).
- [42] T. Banks, D. B. Kaplan, and A. E. Nelson, Phys. Rev. D **49**, 779 (1994).