Phenomenology of superstring-inspired SO(10) grand-unified-theory model with intermediate scale

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We examine in detail a superstring-inspired SO(10) grand-unified-theory model in which the mixing of Weinberg-Salam Higgs fields and large Majorana masses for right-handed neutrinos are generated by the higher-dimensional terms in the superpotential at the intermediate-scale symmetry breaking. Fixing the Higgs-field mixing strength μ as 1 TeV and assuming the usual soft terms of broken $N=1$ supergravity, we study the radiative symmetry breakings at the intermediate scale M_I and the scale M_W. Conditions and constraints for the symmetry breakings severely restrict the allowed range of the supergravity parameters and give bounds on the top-quark mass m_t. We have obtained $m_t \leq 141$ GeV if the identity $B=A-1$ is assumed. The low-energy particle spectra are analyzed. All squarks and sleptons are predicted to have masses in the TeV region. Some of the gaugino-related supersymmetric particles can be light, and lower bounds for these sparticle masses are given. One neutral Higgs boson is light and the other Higgs bosons are all very massive (> a few TeV).

I. INTRODUCTION

Along with the boom in superstring theories, considerable effort has been made in recent years to construct phenomenologically promising low-energy effective theories based on the ten-dimensional $E_8 \times E_{8'}$ heterotic superstring theory.[1] In order to make contact with low-energy physics in four dimensions, the six surplus dimensions of the superstring must be compactified. Up to the present, several compactification schemes have been proposed. The most popular scheme is Calabi-Yau compactification with identification of the gauge and spin connections, which leads to four-dimensional $N=1$ supersymmetric grand-unified-theory (GUT) models with gauge group $E_6 \times E_{8'}$. Such models have been studied extensively by many authors.[2-7]

Meanwhile, Witten[8] and Hull[9] have proposed a different compactification scheme which is based on manifolds with alternative embeddings of the spin connection in the gauge group. This scheme is called (2,0) compactification since the resulting two-dimensional σ model on the string world sheet has (2,0) supersymmetry, and it gives rise to SO(10) or SU(5) supersymmetric GUT models. Such models have been investigated in Refs. 10–15. The questions are then posed whether or not the latter compactification scheme is theoretically consistent and brings about phenomenologically acceptable models.[16]

A major theoretical problem is that the (2,0) compactification may be unstable by the presence of world-sheet instantons.[17] However, it was pointed out recently[18] that under certain conditions the destabilization of (2,0) compactifications by instantons does not occur. Under the present circumstances in which we do not have sufficient knowledge of the (2,0) compactification, it would seem premature to abandon superstring-inspired (SI) SO(10) and SU(5) GUT models.

Phenomenological difficulties of SI SO(10) and SI SU(5) models lie in the Higgs sector. The mixing between the Weinberg-Salam supersymmetric Higgs doublets H_d and H_u is essential to any supergravity model with realistic electroweak breaking.[19,20] A term $\mu H_d H_u$ in the effective superpotential generates a nontrivial minimum for the Higgs-scalar potential at nonzero vacuum expectation values (VEV's) of $\langle H_d \rangle$ and $\langle H_u \rangle$, and the spontaneous symmetry breaking of $SU(2)_W \times U(1)_Y$ is brought about. Also, $H_d H_u$ mixing is necessary to avoid the appearance of an unacceptable axion. Because phenomenological superstring models do not contain bilinear terms in the superpotential, μ is thought to be provided by an $SU(3)_C \times SU(2)_W \times U(1)_Y$-singlet field N which has a required Yukawa coupling NH_dH_u and acquires the VEV $\langle N \rangle$:

$$W \ni \lambda N H_d H_u, \quad \mu = \lambda \langle N \rangle . \tag{1.1}$$

In the ordinary SI E_6 models, chiral superfields of the **27** representation contain an $SU(3)_C \times SU(2)_W \times U(1)_Y$-singlet field N which plays the role in Eq. (1.1). On the other hand, in the SI SO(10) and SI SU(5) models, the chiral fields which couple to H_d and H_u should be SO(10) and SU(5) singlets, respectively. Then we are forced to include the singlet fields N in the SO(10) and SU(5) models. Indeed the SI SO(10) and SI SU(5) models of Refs. 12–15 accommodate these singlets. The introduction of the N fields, however, raises a problem. Generically, the singlet fields N have a Yukawa coupling λN^3 as well, which gives unacceptable domain walls.[16] In addition, more additional terms including the N fields are allowed in the superpotential and we need fine-tuning of these coupling parameters to make the models phenomenologically viable.

Recently one of the authors proposed[21] an alternative possibility of the H_dH_u mixing in the SI SO(10) GUT models without introducing any singlet field N. It was reported there that when a SI SO(10) GUT model has a symmetry breaking at the intermediate scale, the mixing of two Higgs doublets H_d and H_u could arouse through the nonrenormalizable terms in the superpotential which may be also responsible for the solution of the light-neutrino mass problem.

In fact, a possibility that the H_dH_u mixing could be generated by the nonrenormalizable terms has already been discussed by Kim and Nilles.[22] With a specific model they presented an example that solves the strong CP problem and the H_dH_u-mixing problem simultaneously. They showed that μ is comparable to a gravitino mass $m_{3/2}$. On the other hand, in Ref. 21, μ was linked with a large Majorana mass m_{ν_R} for the right-handed neutrino[23] and it was expected that $\mu \sim m_{\nu_R} \sim 1$ TeV.

In this paper, we will elaborate the model presented in Ref. 21 a little further and with this new version of the SI SO(10) GUT model we will explore a scenario in which a symmetry breaking occurs at the intermediate scale M_I, and the H_dH_u mixing and large Majorana masses for right-handed neutrinos turn up from the nonrenormalizable terms. We assume the usual general soft terms of broken $N=1$ supergravity models[20] and that symmetry breakings are generated by the radiative corrections. Then, writing down the evolution equations for the relevant couplings and parameters in the model in one-loop approximations, we will study in detail the constraints on the parameters, the allowed range for the top-quark mass, and the low-energy particle spectra.

The plan of this paper is as follows. In Sec. II we review SI SO(10) GUT models. In the analysis we make full use of the Cartan-Weyl weight method.[5,7,24] Our scenario for the H_dH_u mixing and a specific SI SO(10) GUT model to work with in this paper are presented in Sec. III. Since our specific model has two U(1) gauge symmetries, we also discuss U(1) mixings. In Sec. IV we deal with the evolution of the coupling constants and parameters which appear in the model and study the conditions imposed by the symmetry breakings both at the intermediate scale M_I and at the Weinberg-Salam scale M_W. In Sec. V we first discuss about the allowed regions for the supergravity parameters which are consistent with the conditions and constraints for the intermediate-scale symmetry breaking (ISSB) at M_I. Then examining the conditions for the $SU(2)_W \times U(1)_Y$ breaking at M_W, we give the low-energy particle mass spectra predicted by this model. In particular, the upper bound for the top-quark mass and the lower bounds for gluino and two lightest neutralino masses are given. Also the model predicts one light neutral Higgs boson whose mass is expressed by a simple formula. Section VI is devoted to conclusions and discussions.

II. SI SO(10) GUT MODELS

In this section we review SI SO(10) GUT models by means of the Cartan-Weyl weight method and clarify the notations for the following analyses. The gauge fields in the SI SO(10) models are the vector superfields which belong to the 45-dimensional representation of SO(10). The matter fields, on the other hand, are expected to be the chiral superfields in the following representations of SO(10):

$$N_f \mathbf{16} + \delta(\mathbf{16} + \mathbf{16}^*) + \epsilon \mathbf{10} + \gamma \mathbf{1} \ , \tag{2.1}$$

where N_f, δ, ϵ, and γ are non-negative integers related to topological invariants.[8] N_f flavors of the usual quarks and leptons lie in the $N_f\mathbf{16}$.

There are two possibilities expected for the symmetry-breaking patterns, which are

$$E_8 \underset{M_C}{\rightarrow} G \underset{M_I}{\rightarrow} G' \underset{M_W}{\rightarrow} SU(3)_C \times U(1)_{em} \ , \tag{2.2a}$$

$$E_8 \underset{M_C}{\rightarrow} G \underset{M_W}{\rightarrow} SU(3)_C \times U(1)_{em} \ . \tag{2.2b}$$

The first symmetry-breaking occurs in both (2.2a) and (2.2b) at the compactification scale M_C through the Hosotani mechanism,[2,25] and G is a subgroup of SO(10). The symmetry group G further breaks down into G' at an intermediate scale M_I (intermediate scale scenario) in (2.2a) or remains as an exact symmetry up to the Weinberg-Salam scale M_W (grand desert scenario) in (2.2b). The final symmetry breaking takes place at M_W in both cases.

The first stage of the symmetry breaking at M_C is due to the compactification of the ten-dimensional $E_8 \times E_8'$ superstring theory to $M_4 \times K$, where K is the Calabi-Yau manifold. We assume the existence of the new superstring compactification which gives rise to SO(10) GUT models. When K is multiply connected and is of the form $K_0/\tilde{G}_d$ where K_0 is a simply connected Calabi-Yau manifold and $\tilde{G}_d$ is a discrete group, then the Hosotani mechanism acts in the nontrivial direction of the Wilson loop U and the $E_8 \times E_8'$ gauge symmetry breaks down to $G \times E_8'$. The nontrivial U gives rise to the discrete subgroup G_d of SO(10), which is homomorphic to $\tilde{G}_d$. If U is given, the subgroup G and the massless superfields after the first stage breaking are determined.

It is very convenient in the following analysis to use the Cartan-Weyl weight method.[5,7,24] The Wilson loop U can be expressed as

$$U = \exp\left[i \left[\sum_i z_i H_i + \sum_j x_j E_j \right] \right] \ , \tag{2.3}$$

where the H_i and E_j are generators of SO(10) corresponding to zero roots and nonzero roots, respectively, and the real parameters $Z = \{z_i\}$ and $N = \{x_j\}$ show the symmetry-breaking directions. The SO(10) subgroup G in Eqs. (2.2a) and (2.2b) must contain in itself the standard gauge group G_{st}, which is

$$G_{st} = SU(3)_C \times SU(2)_W \times U(1)_Y \ . \tag{2.4}$$

TABLE I. The gauge field decomposition of the 45-dimensional representation of SO(10) under $G_{st}=SU(3)_C\times SU(2)_W\times U(1)_Y$.

Roots ξ	G_{st}	Number of fields	(Z,ξ)
(01000)	$(8,1,0)$	8	0
(0−1011)	$(1,3,0)$	3	0
(00000)	$(1,1,0)$	2	0
(0−12−1−1)	$(1,1,2)$	2	$4\alpha-2\beta$
(−10100)	$(3^*,2,\frac{5}{3})$	12	$-\beta$
(1−1100)	$(3,2,\frac{1}{3})$	12	$4\alpha-\beta$
(1001−1)	$(3^*,1,-\frac{4}{3})$	6	4α

The symmetry-breaking direction at the first stage, therefore, should preserve G_{st}. But Table I shows that any nonzero root direction breaks at least one of those $SU(3)_C$, $SU(2)_W$, and $U(1)_Y$ symmetries. Hence we conclude that only the zero-root breaking, i.e., $Z=\{z_i\}\neq 0$ and $N=\{x_j\}=0$, can occur at the first stage and, therefore, the rank of the subgroup G remains five, the same as the rank of SO(10). Also, $\tilde{G}_d$ is necessarily an Abelian discrete group.

The breaking direction Z can be determined in the following way: A SO(10) gauge boson with root vector ξ gains a mass which is proportional to (Z,ξ). Since the gauge group G includes G_{st} in itself, Z must satisfy $(Z,\xi)=0$ for the root vector ξ in G_{st}. Thus we find, in the dual basis,

$$Z=[2\alpha,0,2\alpha-\beta,\alpha,-\alpha]_d\ , \tag{2.5}$$

where the two parameters α,β are treated up to the modulus of the discrete group G_d.

Phenomenological constraints will further restrict the breaking direction Z. To forbid the fast proton decay, those gauge bosons with the representations $(3^*,2,\frac{5}{3})$ and $(3,2,\frac{1}{3})$ under G_{st} should be massive at the first stage. Then we find from Table I the conditions

$$-\beta\neq 0,\quad 4\alpha-\beta\neq 0\quad (\mathrm{mod}\,G_d)\ . \tag{2.6}$$

There are only four possible groups for G which satisfy the conditions (2.5) and (2.6). By counting the number of the massless gauge bosons from Table I, they are easily found as

$$\text{(i) } G=SU(4)_C\times SU(2)_W\times SU(2)_R\ , \tag{2.7a}$$

$$\text{(ii) } G=SU(4)_C\times SU(2)_W\times U(1)\ , \tag{2.7b}$$

$$\text{(iii) } G=SU(3)_C\times SU(2)_W\times SU(2)_R\times U(1)\ , \tag{2.7c}$$

$$\text{(iv) } G=G_{\min}=SU(3)_C\times SU(2)_W\times U(1)\times U(1)\ . \tag{2.7d}$$

Now let us consider matter fields which remain massless after the action of the Hosotani mechanism. The $N_f\,\mathbf{16}$ chiral superfields remain massless, because they are protected by the index theorem. On the other hand, among the superfields of the $\epsilon\mathbf{10}$ and $\delta(\mathbf{16}+\mathbf{16}^*)$ multiplets, only those components which are invariant under the action of $\tilde{G}_d+G_d$ remain massless. Under the action $\tilde{G}_d$ all components of an irreducible representation of SO(10) are multiplied by the same phase factor η(Ref. 11), and the state which belongs to the weight ρ in SO(10) acquires the phase (Z,ρ) under the action G_d (see Tables II and III).

III. A SPECIFIC MODEL WITH INTERMEDIATE SCALE SYMMETRY BREAKING

A. Scenario

We will pursue a possibility of H_dH_u mixing in the SI SO(10) models without introducing any singlet field N. The Higgs doublets H_d and H_u must belong to $\epsilon\mathbf{10}$, because the $\mathbf{10}\times\mathbf{16}\times\mathbf{16}$ couplings are allowed in the superpotential but the $(\mathbf{16})^3$ couplings are not. Hence, for the mixing between H_d and H_u, the superpotential should comprise a term that has two $\mathbf{10}$'s. But, as long as we

TABLE II. The decomposition of the $\mathbf{10}$ chiral superfields under $G_{st}=SU(3)_C\times SU(2)_W\times U(1)_Y$ and field assignments. Both $U(1)_R$ and $U(1)_C$ charges and $U(1)_A$ and $U(1)_B$ charges are presented for each field. The weights are given in terms of the Dynkin label.

Field assignments	G_{st}	$U(1)_R$ and $U(1)_C$ charges	$U(1)_A$ and $U(1)_B$ charges	Weight ρ	(Z,ρ)
g	$(3,1,-\frac{2}{3})$	$\left[0,-\frac{\sqrt{10}}{6}\right]$	$\left[\frac{1}{3},-\frac{\sqrt{6}}{6}\right]$	(10000)	2α
$\bar{g}$	$(3^*,1,\frac{2}{3})$	$\left[0,\frac{\sqrt{10}}{6}\right]$	$\left[-\frac{1}{3},\frac{\sqrt{6}}{6}\right]$	(−11000)	-2α
H_u	$(1,2,1)$	$\left[\frac{\sqrt{15}}{6},0\right]$	$\left[\frac{1}{2},\frac{\sqrt{6}}{6}\right]$	(0−1100)	$2\alpha-\beta$
H_d	$(1,2,-1)$	$\left[-\frac{\sqrt{15}}{6},0\right]$	$\left[-\frac{1}{2},-\frac{\sqrt{6}}{6}\right]$	(00−111)	$-2\alpha+\beta$

TABLE III. The decomposition of the **16** chiral superfields in **16**+**16*** under $G_{\rm st}={\rm SU}(3)_C\times{\rm SU}(2)_W\times{\rm U}(1)_Y$ and field assignments. Both ${\rm U}(1)_R$ and ${\rm U}(1)_C$ charges and ${\rm U}(1)_A$ and ${\rm U}(1)_B$ charges are presented for each field. The weights are given in terms of the Dynkin label.

Field assignments	$G_{\rm st}$	${\rm U}(1)_R$ and ${\rm U}(1)_C$ charges	${\rm U}(1)_A$ and ${\rm U}(1)_B$ charges	Weight ρ	(Z,ρ)
	$(3,2,\frac{1}{3})$	$\left[0,\frac{\sqrt{10}}{12}\right]$	$\left[-\frac{1}{6},\frac{\sqrt{6}}{12}\right]$	(00001)	$-\alpha$
	$(3^*,1,-\frac{4}{3})$	$\left[-\frac{\sqrt{15}}{6},-\frac{\sqrt{10}}{12}\right]$	$\left[-\frac{1}{3},-\frac{\sqrt{6}}{4}\right]$	(01−110)	$-\alpha+\beta$
D_R	$(3^*,1,\frac{2}{3})$	$\left[\frac{\sqrt{15}}{6},-\frac{\sqrt{10}}{12}\right]$	$\left[\frac{2}{3},\frac{\sqrt{6}}{12}\right]$	(0010−1)	$3\alpha-\beta$
$L'=(N_L,E_L)$	$(1,2,-1)$	$\left[0,-\frac{\sqrt{10}}{4}\right]$	$\left[\frac{1}{2},-\frac{\sqrt{6}}{4}\right]$	(1−1010)	3α
	$(1,1,2)$	$\left[\frac{\sqrt{15}}{6},\frac{\sqrt{10}}{4}\right]$	$\left[0,\frac{5\sqrt{6}}{12}\right]$	(−101−10)	$-\alpha-\beta$
N_R	$(1,1,0)$	$\left[-\frac{\sqrt{15}}{6},\frac{\sqrt{10}}{4}\right]$	$\left[-1,\frac{\sqrt{6}}{12}\right]$	(−11−101)	$-5\alpha+\beta$

consider trilinear couplings in the superpotential and do not introduce any SO(10)-singlet field, there is no possibility for the existence of such a term. Then let us consider the higher-dimensional terms in the superpotential. These terms have already been introduced into model buildings to solve the neutrino mass problem:[6,26] A large Majorana mass is given to a right-handed neutrino ν_R at the ISSB through a nonrenormalizable term of the form $\mathbf{16}\times\mathbf{16}\times\mathbf{16}^*\times\mathbf{16}^*$ in the superpotential. Similarly, if we allow the existence of nonrenormalizable terms, we can find a solution to the H_dH_u-mixing problem in the SI SO(10) models. The mixing of the two Higgs superfields could arouse at the ISSB through a nonrenormalizable term of the form $\mathbf{10}\times\mathbf{10}\times\mathbf{16}\times\mathbf{16}^*$.

Now our scenario for the H_dH_u mixing is explained.[21]

(1) The symmetry group G is broken spontaneously at an intermediate scale M_I [thus we adopt the intermediate scale scenario of Eq. (2.2a)]. The residual gauge group G' is just the one for the standard model, i.e., $G'=G_{\rm st}$. Incidentally, the mass of the extra neutral vector boson Z' which exists in the usual SO(10) GUT models is expected to be very massive due to the ISSB, and the well-known Z-Z' mixing problem is avoided.

(2) The ISSB is caused by N_R and its conjugate $\overline{N_R}$ in $\delta(\mathbf{16}+\mathbf{16}^*)$ having the equal vacuum expectation values (VEV's)

$$\langle N_R\rangle=\langle\overline{N_R}\rangle\ . \tag{3.1}$$

Here N_R denotes a state in the **16** part of **16**+**16*** with the same quantum numbers of the right-handed sneutrino $\tilde{\nu}_R$. The two VEV's should be equal, since otherwise they would give nonvanishing contributions to the D^2 terms in the scalar potential and would break supersymmetry earlier than desired.[4]

(3) We will assume that the nonrenormalizable higher-dimensional terms appear in the superpotential somehow by the yet-unexplained mechanism.

(4) At the ISSB a right-handed neutrino ν_R gains a large Majorana mass through a term of the form

$$O(1/M_C)\mathbf{16}\times\mathbf{16}\times\mathbf{16}^*\times\mathbf{16}^*\to(\langle\overline{N_R}\rangle^2/M_C)\nu_R\times\nu_R\ . \tag{3.2}$$

Thus ν_R obtains a mass $m_{\nu_R}\sim\langle\overline{N_R}\rangle^2/M_C$. If the Dirac-neutrino mass is of the same order as the up-quark mass, say, 5 MeV (Ref. 27), then through the "seesaw" mechanism[23] a large Majorana mass $m_{\nu_R}\sim10^3$ GeV gives a neutrino mass $m_\nu\sim10$ eV which is within the present experimental limit.[28]

(4) When N_R and $\overline{N_R}$ get VEV's at the ISSB, a nonrenormalizable term of the form

$$O(1/M_C)\mathbf{10}\times\mathbf{10}\times\mathbf{16}\times\mathbf{16}^* \tag{3.3}$$

yields H_dH_u mixing of the form

$$(\langle N_R\rangle\langle\overline{N_R}\rangle/M_C)H_dH_u\ , \tag{3.4}$$

which can be written as μH_dH_u with $\mu=\langle N_R\rangle\langle\overline{N_R}\rangle/M_C$. The coupling strength μ is expected to be of the same order as m_{ν_R}, say, 10^3 GeV, which is considerably large compared to the value expected in the usual $N=1$ softly broken supergravity models.[29–33]

B. A specific SI SO(10) GUT model

We will now pick up a specific SI SO(10) model and demonstrate how plausibly the above scenario comes about. Among the four possible groups for G in Eqs. (2.7a)–(2.7d), the first two [(i) and (ii)] are discarded by the following reasons: In both cases the

$SU(3)_C \times SU(2)_W \times U(1)_Y$ singlet N_R from $\delta(\mathbf{16}+\mathbf{16}^*)$, which is responsible for the ISSB and, therefore, should remain massless below M_C, lies in the same **4** of $SU(4)_C$ as a colored triplet. If N_R is massless, this extra colored triplet is also massless and could bring about the fast proton decay. Besides, the models built with the gauge group (i) or (ii) give rather large one-loop predictions for $\sin^2\theta_W$. Studying next the remaining two groups for G closely, we finally choose the last one, i.e., case (iv), which gives a more preferable prediction for $\sin^2\theta_W$ than case (iii). The breaking direction Z which yields $G=G_{\min}$ is in the form of Eq. (2.5) with the conditions

$$4\alpha\neq 0,\quad -\beta\neq 0,\quad 4\alpha-2\beta\neq 0,\quad 4\alpha-\beta\neq 0 \pmod{G_d}\ , \tag{3.5a}$$

which are easily derived from Table I and the fact that the gauge bosons outside of $G_{\min}$ must be massive. The appropriate discrete subgroup G_d is thus given by

$$G_d = Z_m \times Z_n \quad (m\neq 4,\ 4\alpha\neq 2\beta,\ 4\alpha\neq\beta)\ . \tag{3.5b}$$

Once the specific group $G_{\min}$ has been chosen, our next task is to find the matter fields which remain massless modes after the Hosotani breaking at M_C. As stated in Sec. II, the N_f chiral superfields remain massless. Furthermore, we need at least a pair of Higgs doublets H_d and H_u from $\epsilon\mathbf{10}$ and a pair of superfields N_R and $\overline{N_R}$ from $\delta(\mathbf{16}+\mathbf{16}^*)$ to be massless mode below M_C.

Under the action $\tilde{G}_d+G_d$, a component with weight ρ in a **10** multiplet obtains the phase $\eta_{10}+(Z,\rho)$, where η_{10} is the phase acquired by the **10** under $\tilde{G}_d$. From the values of (Z,ρ) in Table II and the conditions in Eq. (3.5a) for $G_{\min}$, we find that, once one component in a **10** multiplet is set to be a massless mode, all other components in the same **10** become massive modes. Hence the required Higgs doublets H_d and H_u cannot belong to the same **10** multiplet and they come from different **10**'s (Ref. 13). Also the color-triplet fields g and $\bar{g}$ from the **10** to which the light Higgs bosons belong turn out to be massive modes and, therefore, the fast proton decays through these g-quark exchange are automatically avoided.

In order to determine the other massless matter fields, we must consider the symmetry-breaking mechanism. It is an interesting and popular idea in the low-energy supergravity models that the symmetry breakings are triggered by the radiative corrections.[20] We follow the same lines and assume that one has the usual general soft terms of $N=1$ supergravity models,[20] i.e., the universal mass m for all scalars and M for all gauginos and trilinear scalar couplings (parametrized by the parameter A) at the compactification scale M_C. Below M_C, these scalar masses, gaugino masses, and various coupling parameters evolve in accordance with the renormalization-group (RG) equations.

At the intermediate scale M_I, the symmetry breaking will be generated by N_R and $\overline{N_R}$ having large VEV's. The relevant parts of the tree-level scalar potential are of the form

$$V=m_{N_R}^2 N_R^2 + m_{\overline{N_R}}^2 \overline{N_R}^{\,2} + c(N_R^2-\overline{N_R}^{\,2})^2 + (|f|^2/M_C^2)(N_R^2\overline{N_R}^{\,4}+N_R^4\overline{N_R}^{\,2})\ , \tag{3.6}$$

where the third term is the D^2-term contribution, and the constant c is positive and depends on the two U(1) charges of N_R and $\overline{N_R}$ which will be given later [see Eq. (3.22)]. The last term results from a nonrenormalizable term of the form $(f/M_C)N_R^2\overline{N_R}^{\,2}$ in the superpotential,[4] and this term becomes relevant when N_R and $\overline{N_R}$ have large VEV's. The higher-dimensional terms are also possible, but they are expected to give much smaller contribution to the potential than the last term and are neglected.

Differentiating V by N_R^2 and $\overline{N_R}^{\,2}$, we find the condition for V having an extremum at nonzero VEV's of $\langle N_R\rangle=\langle\overline{N_R}\rangle$ [see Eq. (3.1)] as follows:

$$m_{N_R}^2 = m_{\overline{N_R}}^2 = -3(|f|^2/M_C^2)\langle N_R\rangle^4\ . \tag{3.7a}$$

In order for this extremum to be a minimum, we further need

$$(|f|^2/M_C^2)\langle N_R\rangle^2 < 2c\ . \tag{3.7b}$$

In our scenario $|f|\langle N_R\rangle^2/M_C$ is presumed to be of the same order as the Majorana-neutrino mass $m_{\nu_R}\sim 10^3$ GeV. Hence, if the constants f and c are of the order one, the condition Eq. (3.7b) is well satisfied, and the right-hand side (RHS) of Eq. (3.7a) is negative and expected to be around $-10^6\ \text{GeV}^2$. Thus, in order to generate the ISSB, the both scalar masses $m_{N_R}^2$ and $m_{\overline{N_R}}^2$, which are assumed to be positive and equal to the other scalar masses at the compactification scale M_C, must be driven to negative values at the intermediate energies. Accordingly, we need to provide appropriate superfields which have Yukawa couplings to N_R and $\overline{N_R}$, and those fields should be massless modes at the first stage.

As for the massless superfields in $\delta(\mathbf{16}+\mathbf{16}^*)$, we repeat the same analysis as we did for the $\epsilon\mathbf{10}$. Keeping in mind this time that the states in **16** acquire an equal phase under $\tilde{G}_d$, we find from Table III the massless modes in **16** (which should include N_R) and the conditions. There are three possible cases:

N_R,

$$8\alpha\neq\beta,\ 8\alpha\neq 2\beta\ \pmod{G_d}+\text{conditions in Eq. (3.5a)}\ ; \tag{3.8a}$$

N_R, D_R,

$$8\alpha=2\beta\ \pmod{G_d}+\text{conditions in Eq. (3.5a)}\ ; \tag{3.8b}$$

N_R, E_L, N_L,

$$8\alpha=\beta\ \pmod{G_d}+\text{conditions in Eq. (3.5a)}\ , \tag{3.8c}$$

where we use the notation (N_L,E_L) and D_R to denote the $SU(2)_W$-doublet and color-triplet states in the **16** part of $\mathbf{16}+\mathbf{16}^*$ with quantum numbers of (ν_L,e_L) and d_R, respectively. We will then examine the ISSB for the above three cases in Eqs. (3.8a)–(3.8c).

In the cases (a) and (b), the only possible G-invariant superpotential which contains N_R are of the form

$$W^{(a),(b)}=\sum_i \lambda_i L_i N_R H_u \ , \tag{3.9}$$

where L_i represents the $SU(2)_W$ lepton doublet (ν_{L_i}, e_{L_i}) from the ith family, and λ_i is the Yukawa coupling. If $\langle N_R \rangle$ develops, lepton doublets will in general acquire intermediate scale masses. So we need $\lambda_i=0$. In other words, N_R should not couple to usual leptons. Then, the running mass $m^2_{N_R}$ cannot be driven to negative values. Besides, there exists no G-invariant superpotential for $\overline{N_R}$, and, thus, $m^2_{\overline{N_R}}$ remains positive. The cases (a) and (b), therefore, do not work.

On the other hand, in the case (c), the superpotential can have the following terms containing N_R and $\overline{N_R}$:

$$W^{(c)}=\sum_j \kappa_j L' N_R H_{uj} + \sum_j \overline{\kappa_j}\,\overline{L'}\,\overline{N_R} H_{dj} \ , \tag{3.10}$$

where L' and $\overline{L'}$ are the $SU(2)_W$ lepton doublets from **16**+**16***, κ_j and $\overline{\kappa_j}$ are the Yukawa coupling constants, and we have introduced more Higgs doublets H_{uj} and H_{dj} $(j=1,2,\dots)$. Indeed the Yukawa couplings in Eq. (3.10) can drive the running masses of $m^2_{N_R}$ and $m^2_{\overline{N_R}}$ to negative values and bring about the ISSB. The case (c) works for our scenario. However, a plausible model should have at least one pair of light Higgs doublets H_u and H_d which break the electroweak symmetry at the scale M_W. Those Higgs doublets cannot couple to N_R and $\overline{N_R}$, because they would gain intermediate scale masses once $\langle N_R \rangle$ and $\langle \overline{N_R} \rangle$ develop. We will introduce only one such pair into our specific model. We assign $j=1$ for them and assume

$$\kappa_j=\overline{\kappa_j}=0 \quad \text{for } j=1 \ . \tag{3.11}$$

In addition the model should have other Higgs doublets H_{uj} and H_{dj} $(j \ge 2)$ which do couple to N_R and $\overline{N_R}$. So we need $\epsilon \ge 4$.

Now we present a specific SI SO(10) model with which the possibility of the scenario discussed in Sec. III A will be pursued in the following. We introduce one pair of mirror multiplets **16**+**16*** $(\delta=1)$, and four **10** $(\epsilon=4)$. This is due to the fact that we will obtain the more preferable value for $\sin^2\theta_W$ when the model has fewer matter fields aside from N_f**16**. The first symmetry breaking occurs at M_C by the Hosotani mechanism and the breaking direction Z satisfies the conditions in Eq. (3.8c). Below M_C the model has the gauge group

$$G=G_{\min}=SU(3)_C\times SU(2)_W \times U(1)^2 \ , \tag{3.12}$$

the massless matter modes

$$\begin{aligned} &N_f\mathbf{16} \ , \\ &N_R,\ L',\ \overline{N_R},\ \text{and } \overline{L}' \ \text{ from } \mathbf{16}+\mathbf{16}^* \ , \\ &H_{u1}, H_{d1}, H_{u2}, \text{and } H_{d2} \ \text{ from four } \mathbf{10}\text{'s} \ , \end{aligned} \tag{3.13}$$

and the superpotential

$$W=h_t H_{u1} Q t_R + \kappa L' N_R H_{u2} + \overline{\kappa}\,\overline{L'}\,\overline{N_R} H_{d2} + \text{possible nonrenormalizable terms} \ , \tag{3.14}$$

where Q denotes the superfield for the left-handed quark doublet (t_L, b_L). We only take into account the Yukawa coupling for the top quark with the coupling constant h_t, and the Yukawa coupling for other quarks and leptons are neglected. Also we omit the index 2 for the Yukawa coupling constants κ_2 and $\overline{\kappa}_2$.

The spontaneous symmetry breaking takes place at the intermediate scale M_I. Scalar components of N_R and $\overline{N_R}$ superfields acquire large VEV's and superfields L', $\overline{L}'$, H_{u2}, and H_{d2} become very massive. Hence the model has, below M_I, the gauge group

$$G'=SU(3)_C\times SU(2)_W\times U(1)_Y \tag{3.15}$$

and the massless matter modes

$$N_f\mathbf{16}, \quad H_{u1}, H_{d1} \ . \tag{3.16}$$

C. U(1) mixings

It is well known that in SI GUT models with two or more U(1) gauge groups, mixing between U(1)'s may occur under renormalization.[34] Since our specific model has two U(1)'s, we must choose the appropriate U(1) generators so that no mixing arises among them. These generators are uniquely determined by the contents of massless chiral superfields. The analysis of evolution equations for U(1) gauge couplings should be done for those which correspond to properly chosen U(1) generators.

Let us first take $U(1)_R$ and $U(1)_C$ for our two U(1)'s, where $U(1)_R$ is the U(1) subgroup of $SU(2)_R$ and $U(1)_C$ is the one which appears in the maximal decomposition of $SU(4)_C$, i.e., $SU(4)_C \supset SU(3)_C \times U(1)_C$. The $U(1)_R$ and $U(1)_C$ charges Q^i_R and Q^i_C of massless chiral superfields are normalized as follows:

$$\sum_{\text{all } i\in \mathbf{16}} (Q^i_R)^2 = \sum_{\text{all } i\in \mathbf{16}} (Q^i_C)^2 = \tfrac{10}{3} \tag{3.17a}$$

and

$$\sum_{\text{all } i\in \mathbf{10}} (Q^i_R)^2 = \sum_{\text{all } i\in \mathbf{10}} (Q^i_C)^2 = \tfrac{5}{3} \ . \tag{3.17b}$$

We list in Tables II and III the $U(1)_R$ and $U(1)_C$ charges of the fields in 10- and 16-dimensional representations. Next we define the quantities

$$p=\sum_i (Q^i_R)^2, \quad q=\sum_i (Q^i_C)^2, \quad r=\sum_i Q^i_R Q^i_C \ , \tag{3.18}$$

with i running through all massless chiral superfields. When $r\neq 0$, two U(1)'s mix under renormalization. Indeed if *all* chiral superfields in *a complete representation* take part in the summation for r, their contributions to r sum up to zero. But in the model concerned here only a few components in four **10**'s and **16**+**16*** remain massless below M_C and these "survivors" will give nonvanishing $r=-5\sqrt{6}/12$.

The appropriate U(1)'s, which we denote as $U(1)_A$ and

$U(1)_B$, are obtained by rotating the old ones:

$$U(1)_A = U(1)_R \cos\phi - U(1)_C \sin\phi \;,$$
$$U(1)_B = U(1)_R \sin\phi + U(1)_C \cos\phi \;. \tag{3.19}$$

The angle ϕ is determined by the condition that the products of two new U(1) charges Q_A^i and Q_B^i of massless chiral superfields add up to zero, i.e., $\sum_i Q_A^i Q_B^i = 0$. This condition can be written as

$$\tan 2\phi = -\frac{2r}{p-q} \tag{3.20}$$

and we obtain $\tan 2\phi = 2\sqrt{6}$. If ϕ is chosen to be $0 \le \phi \le \pi/2$, we find

$$\cos\phi = \sqrt{3/5}, \quad \sin\phi = \sqrt{2/5} \;. \tag{3.21}$$

Now it is an easy task to calculate the new $U(1)_A$ and $U(1)_B$ charges for the components in **10** and **16** representations and they are listed also in Tables II and III.

Referring to the $U(1)_A$ and $U(1)_B$ charges of the fields N_R and $\overline{N_R}$ in Table III, we find that the positive constant c in the third term of Eq. (3.6) is given by

$$c = \tfrac{1}{2}(g_{1A}^2 + \tfrac{1}{24} g_{1B}^2) \;, \tag{3.22}$$

where g_{1A} and g_{1B} are the gauge coupling constants for $U(1)_A$ and $U(1)_B$, respectively.

IV. EVOLUTION OF COUPLINGS AND PARAMETERS

Various coupling constants and parameters which appear in the model change with energy according to the RG equations.[19,29,31] For completeness the relevant one-loop RG equations for these parameters from the scale M_C to M_I and from M_I to M_W are shown in Appendixes A and B, respectively. Solving these coupled differential equations, we also give in the Appendixes the expressions for all the relevant renormalized couplings and parameters.

At the compactification scale M_C we take the usual boundary conditions,[20] i.e., the unification of all gauge coupling constants, the universal mass M for all gauginos and m for all scalars, and the universal A for all soft-breaking trilinear couplings:

$$g_3^2 = g_2^2 = \tfrac{5}{3} g_{1A}^2 = \tfrac{5}{3} g_{1B}^2 \;, \tag{4.1a}$$

$$M_3 = M_2 = M_{1A} = M_{1B} = M \;, \tag{4.1b}$$

$$m_i^2 = m \quad (i = \text{all scalars}) \;, \tag{4.1c}$$

$$A_t = A_\kappa = A_{\bar\kappa} = A \;, \tag{4.1d}$$

where g_3, g_2, g_{1A}, and g_{1B} are gauge coupling constants, and M_3, M_2, M_{1A}, and M_{1B} are gaugino masses for $SU(3)_C$, $SU(2)_W$, $U(1)_A$, and $U(1)_B$, respectively. The appearance of the factor $\frac{5}{3}$ in Eq. (4.1a) is due to our normalization convention of two U(1) charges in Eqs. (3.17a) and (3.17b). The parameters A_t, A_κ, and $A_{\bar\kappa}$ are the trilinear scalar couplings which correspond to the Yukawa couplings for the t quark, N_R, and $\overline{N_R}$, respectively. The Yukawa coupling constants h_t, κ, and $\bar\kappa$ at the scale M_C cannot be fixed and remain as free parameters. However, as discussed in Sec. III, N_R and $\overline{N_R}$ should have equal VEV's of $\langle N_R \rangle = \langle \overline{N_R} \rangle$ at the intermediate scale M_I. So we assume, at the scale M_C,

$$\kappa = \bar\kappa \;. \tag{4.1e}$$

With this assumption the scalar masses $m_{N_R}^2$ and $m_{\overline{N_R}}^2$ remain equal during the evolution from the scale M_C to M_I, and the equal VEV's of $\langle N_R \rangle = \langle \overline{N}_R \rangle$ are developed at M_I.

A. The Weinberg angle $\sin^2\theta_W$

The values of M_C, M_I and the Weinberg angle $\sin^2\theta_W$ are closely interconnected. At the intermediate scale M_I, the gauge symmetry $U(1)_A \times U(1)_B$ is spontaneously broken down to $U(1)_Y$. In the model concerned, the weak hypercharge Y is related to the $U(1)_A$ and $U(1)_B$ charges X_A and X_B as

$$Y = \tfrac{1}{5} X_A + \frac{2\sqrt{6}}{5} X_B \tag{4.2}$$

and, hence, the $U(1)_Y$ gauge coupling constant g_Y at the scale M_I is given by

$$g_Y^{-2} = \tfrac{1}{25} g_{1A}^{-2} + \tfrac{24}{25} g_{1A}^{-2} \quad \text{at} = E = M_I \;. \tag{4.3}$$

Now using Eq. (4.3) and the evolution equations for the gauge coupling constants above and below M_I which are listed in Eqs. (A40) and (B27) together with the β-function coefficients b_i ($i = 2, 1A$, and $1B$) and b_j' ($j = 2$ and Y) listed in Eqs. (A4) and (B2), we find that the Weinberg angle $\sin^2\theta_W$ at the scale M_W is expressed as

$$\sin^2\theta_W(M_W) = \frac{3}{8} - \frac{\alpha_{em}(M_W)}{4\pi}\left[6\ln\frac{M_C}{M_W} + \ln\frac{M_I}{M_W}\right] , \tag{4.4}$$

where α_{em} ($= e^2/4\pi$) is the electromagnetic coupling and it is estimated that $\alpha_{em} = \frac{1}{128}$ at $E = M_W \approx 80$ GeV. Also the $SU(3)_C$ gauge coupling constant $\alpha_3 = g_3^2/4\pi$ at the scale M_W is expressed as

$$\alpha_3(M_W)^{-1} = \tfrac{3}{8}\alpha_{em}(M_W)^{-1} - \frac{1}{4\pi}\left[18\ln\frac{M_C}{M_W} - 3\ln\frac{M_I}{M_W}\right] . \tag{4.5}$$

It is interesting to note here that the predictions of both $\sin^2\theta_W(M_W)$ and $\alpha_3(M_W)$ in one-loop approximations do not depend on the number of families N_f. This is one of the features of SO(10) GUT models.

In Table IV we show the predicted values of $\sin^2\theta_W(M_W)$ and $\alpha_3(M_W)^{-1}$ for several choices of M_C and M_I when M_W is fixed at 80 GeV and $\alpha_{em}(M_W) = \frac{1}{128}$ is used. The numerical values of $\sin^2\theta_W(M_W)$ tend to increase as M_I takes smaller values while the values of $\alpha_3(M_W)^{-1}$ show the opposite tendency. The recent analysis of existing data on the weak neutral-current events and the W and Z masses gives[35] $\sin^2\theta_W(M_W) = 0.230 \pm 0.0048$. For the color interaction, it has been estimated[36] that $\alpha_3(M_W) = 0.12^{+0.01}_{-0.02}$ corre-

TABLE IV. Predicted values of $\sin^2\theta_W$ and α_3^{-1} at M_W (=80 GeV) for several choices of M_C and M_I.

M_C (GeV)	M_I (GeV)	$\sin^2\theta_W(M_W)$	$\alpha_3(M_W)^{-1}$
	10^{16}	0.2252	5.96
10^{17}	10^{15}	0.2266	5.41
	10^{14}	0.2280	4.86
	10^{15}	0.2352	8.70
10^{16}	10^{14}	0.2366	8.16
	10^{13}	0.2380	7.61
	10^{14}	0.2452	11.45
10^{15}	10^{13}	0.2466	10.90
	10^{12}	0.2481	10.35

sponding to the QCD scale $\Lambda_{\rm QCD}=150^{+150}_{-100}$ MeV.

Taking into account the fact that the evolution equations for the gauge coupling constants used here are those which have been obtained by one-loop approximations to the β functions and by replacing all threshold effects of heavy particles with θ functions, it may be reasonable to regard the model as phenomenologically acceptable if its predicted values of $\sin^2\theta_W(M_W)$ and $\alpha_3(M_W)^{-1}$ fall into the following regions: $0.222<\sin^2\theta_W(M_W)<0.238$ and $7.5<\alpha_3(M_W)^{-1}<10$. These criteria restrict the allowed domain for the scale M_C and M_I quite severely, which is seen from Table IV. The compactification scale M_C must be around 10^{16} GeV. With M_C fixed at 10^{16} GeV, the preferable value of $\sin^2\theta_W(M_W)$ is obtained if M_I takes the value closer to M_C. However, we must keep in mind that we are here considering the model with an ISSB which is generated by radiative corrections. If two scales M_C and M_I are too close, then we need large Yukawa couplings κ and $\bar{\kappa}$ in order to drive the scalar masses $M^2_{N_R}$ and $m^2_{\overline{N_R}}$ quickly to negative values. And they may exceed the limits for the justification of perturbative (one-loop) approximations. After these considerations, we choose

$$M_C=1.5\times10^{16}\ \text{GeV}\ ,\quad M_I=1.0\times10^{14}\ \text{GeV}\ ,\quad M_W=80\ \text{GeV}\ , \tag{4.6}$$

for the numerical analysis in the next section. These inputs give $\sin^2\theta_W(M_W)=0.235$ and $\alpha_3(M_W)^{-1}=7.6$

B. Constraints

There are some restrictions which should be obeyed by the parameters in the model. To ensure that the scalar potential does not have unwanted minima which break the $SU(3)_C\times U(1)_{\rm em}$ symmetry, we impose

$$A_\kappa^2 m^2\le 3(m^2_{N_R}+m^2_{L'}+m^2_{H_{u2}})\ , \tag{4.7a}$$

$$A_{\bar\kappa}^2 m^2\le 3(m^2_{\overline{N_R}}+m^2_{\overline{L'}}+m^2_{H_{u2}})\ , \tag{4.7b}$$

$$A_t^2 m^2\le 3(m^2_Q+m^2_{u_R}+m^2_{H_{u1}})\ . \tag{4.7c}$$

Restrictions (4.7a) and (4.7b) should be observed from the scale M_C to M_I, while (4.7c) must be assured from M_C all the way down to M_W. Especially, at the compactification scale M_C the above restrictions reduce to

$$|A|\le 3\ . \tag{4.8}$$

C. Symmetry breaking at the scale M_I

The ISSB is generated by radiative corrections. Scalar masses $m^2_{N_R}$ and $m^2_{\overline{N_R}}$ are driven to negative values and the scalar potential will have a nontrivial minimum with nonzero VEV's $\langle N_R\rangle=\langle\overline{N_R}\rangle$, which breaks $U(1)_A\times U(1)_B$ symmetry down to $U(1)_Y$.

The expression of the running mass $m^2_{N_R}$ is given in Appendix A. Its behavior with energy depends on the input parameters M, m, A, and $Y_\kappa(M_C)=\kappa(M_C)^2/(4\pi)^2$. The conditions which must be satisfied for the symmetry breaking at M_I have been presented in Eqs. (3.7a) and (3.7b). The condition (3.7b) is believed to be well satisfied, and the RHS of (3.7a) is expected to be around -10^6 GeV^2. For definiteness we take, in the following numerical analysis,

$$m^2_{N_R}(M_I)=m^2_{\overline{N_R}}(M_I)=-10^6\ \text{GeV}^2\ , \tag{4.9}$$

for the condition of the ISSB. Equation (4.9) then enables us to solve $Y_\kappa(M_C)$ as a function of the parameters M, m, and A. In fact, in the next section we will assign some sets of numerical values to M, m, and A and calculate the necessary $Y_\kappa(M_C)$ for the symmetry breaking at M_I. However, we should keep in mind that the acquired $Y_\kappa(M_C)$ must be an appropriate one in the sense that it satisfies the constraint (4.7a) and at the same time is small enough to be justified for perturbative approximations.

The gauge symmetry $U(1)_A\times U(1)_B$ breaks into $U(1)_Y$ at the intermediate scale M_I. The $U(1)_Y$ gauge coupling constant g_Y at M_I is expressed in terms of $U(1)_A$ and $U(1)_B$ couplings g_{1A} and g_{1B} as given in Eq. (4.3). Now let us consider the mass M_Y which the gaugino λ_Y, a superpartner of the $U(1)_Y$ gauge boson, acquires at M_I.

At the scale M_I the scalar fields N_R and $\overline{N_R}$ get the equal VEV's $\langle N_R\rangle=\langle\overline{N_R}\rangle$, and $U(1)_A$ and $U(1)_B$ vector superfields V_A and V_B obtain masses. First we study the situation in which the supersymmetry is an exact symmetry. Then the 2×2 mass matrix for V_A and V_B fields is expressed as

$$\widehat{M}^2=(\langle N_R\rangle^2+\langle\overline{N_R}\rangle^2)\begin{bmatrix} g_{1A}^2 & -\frac{\sqrt6}{12}g_{1A}g_{1B} \\ -\frac{\sqrt6}{12}g_{1A}g_{1B} & \frac{1}{24}g_{1B}^2\end{bmatrix}\ , \tag{4.10}$$

where we have used the fact that $U(1)_A$ and $U(1)_B$ charges of N_R and $\overline{N_R}$ fields are $(-1,\sqrt6/12)$ and $(1,-\sqrt6/12)$, respectively. Two eigenvalues of $\widehat{M}^2$ are $\widehat{M}_Y{}^2=0$ and $\widehat{M}_X{}^2=(\langle N_R\rangle^2+\langle\overline{N_R}\rangle^2)(g^2_{1A}+\frac{1}{24}g^2_{1B})$, and corresponding eigenstates V_Y and V_X are given by

$$V_Y = \cos\xi V_A + \sin\xi V_B ,$$
$$V_X = -\sin\xi V_A + \cos\xi V_B , \tag{4.11}$$

with

$$\tan\xi = 2\sqrt{6} g_{1A}/g_{1B} . \tag{4.12}$$

Thus we find the gaugino λ_Y is massless and the fermion component λ_X of the superfield V_X has a mass $\hat{M}_X{}^2$ provided that the supersymmetry is an exact symmetry.

Now recall that at the compactification scale M_C the supersymmetry has been softly broken and all gauginos have been given an equal mass M. These gaugino masses evolve similarly as the gauge coupling constants [see Eq. (A6)]. The fermion components of superfields V_A and V_B, therefore, keep the nonzero masses $M_{1A}(M_I)$ and $M_{1B}(M_I)$, respectively, at the intermediate scale M_I. Taking these remnant masses into consideration, the gauginos λ_Y and λ_X are no longer eigenstates of the gaugino mass matrix. However, in the limit of $M_{1A}(M_I)$, $M_{1B}(M_I) \ll \hat{M}_X$, in other words, $M_{1A}(M_I), M_{1B}(M_I) \ll \langle N_R \rangle$ which is the situation we are considering now, λ_Y and λ_X can be regarded in good approximation as eigenstates of the gaugino mass matrix, and λ_Y has a mass

$$M_Y(M_I) = M_{1A}(M_I)\cos^2\xi + M_{1B}(M_I)\sin^2\xi$$
$$= \frac{M_{1A}(M_I) g_{1B}^2(M_I) + 24 M_{1B}(M_I) g_{1A}^2(M_I)}{24 g_{1A}^2(M_I) + g_{1B}^2(M_I)} , \tag{4.13}$$

at the scale M_I. Equation (4.13) will then serve as the initial condition for the RG equation of the gaugino mass M_Y from the scale M_I to M_W.

Superfields L', $\overline{L'}$, H_{u2}, and H_{d2}, which have the Yukawa couplings to N_R and $\overline{N_R}$, obtain large masses of the same order as $\langle N_R \rangle$ at the ISSB. Hence, massless modes below the scale M_I are N_f **16**, H_{u1}, and H_{d1}. Since the relevant Higgs fields below M_I are only those of the Weinberg-Salam Higgs doublets H_{u1} and H_{d1}, we omit writing the index 1 from now on.

At the ISSB, the right-handed neutrinos ν_R gain large Majorana masses m_{ν_R} and there appears a mixing between the Weinberg-Salam Higgs doublets H_d and H_u. Thus, the $SU(3)_C \times SU(2)_W \times U(1)_Y$-symmetric superpotential below M_I has the form

$$W = h_t H_u Q t_R + (\text{other usual Yukawa coupling terms})$$
$$+ m_{\nu_R} \nu_R \nu_R + \mu H_d H_u , \tag{4.14}$$

where family indices for the ν_R term are omitted. The coupling strength μ for the $H_d H_u$ mixing is expected to be of the same order as m_{ν_R} at the scale M_I. For definiteness in the following numerical analysis we fix

$$\mu(M_I) = 10^3 \text{ GeV} . \tag{4.15}$$

The superfields ν_R, H_u, and H_d now become massive modes because of the new bilinear terms in Eq. (4.14). However, these fields can still be safely treated in the RG equations as massless modes for most of the energy range from M_I to M_W.

Because of the $H_d H_u$ mixing, the Higgs scalar masses acquire the additional contribution μ^2 below the scale M_I. Then the squared masses μ_u^2 and μ_d^2 of Higgs scalars H_u and H_d, respectively, are given by

$$\mu_u^2 = m_{Hu}^2 + \mu^2 , \tag{4.16a}$$

$$\mu_d^2 = m_{Hd}^2 + \mu^2 . \tag{4.16b}$$

The RG equations followed by these mass parameters μ_u^2 and μ_d^2 and their solutions are presented in Appendix B. The initial conditions for μ_u^2 and μ_d^2 at the scale M_I are

$$\mu_u^2(M_I) = m_{Hu}^2(M_I) + \mu^2(M_I) , \tag{4.17a}$$

$$\mu_d^2(M_I) = m_{Hd}^2(M_I) + \mu^2(M_I) . \tag{4.17b}$$

The relevant Higgs-scalar potential along the neutral direction is now given by

$$V_H = \tfrac{1}{8}(g_2^2 + g_Y^2)(|H_u|^2 - |H_d|^2)^2 + \mu_u^2 |H_u|^2 + \mu_d^2 |H_d|^2$$
$$- \mu_3^2 (H_d H_u + \text{H.c.}) , \tag{4.18}$$

where the last term is again due to the $H_d H_u$ mixing in Eq. (4.14). The potential V_H is of the usual form in $N=1$ supergravity models with a mixed Higgs-boson mass term.[19,29–33] The evolution equation for μ_3^2 and its solution are also presented in Appendix B. The initial condition for μ_3^2 at the scale M_I is

$$\mu_3^2(M_I) = -B m \mu(M_I) , \tag{4.19}$$

where B is the bilinear scalar coupling parameter. It is well known that in the case of the "minimal" $N=1$ supergravity model, there exists a relation[37]

$$B = A - 1 . \tag{4.20}$$

But in the case of a general Kähler manifold, Eq. (4.20) will not hold in general. We will carry out numerical analyses for both the "minimal" and "nonminimal" cases in the next section. Finally, the value m_{ν_R} is hardly altered upon renormalization and remains as large as that given at M_I, which is expected to be around 10^3 GeV.

D. Symmetry breaking at the scale M_W

The $SU(2)_W \times U(1)_Y$ symmetry breaks down spontaneously to $U(1)_{em}$ at the scale M_W. The breaking is induced by the same mechanism as the one used in softly broken $N=1$ supergravity models. The Higgs potential V_H in Eq. (4.18) is minimized for[19]

$$\sin 2\theta = \frac{2\mu_3^2}{\mu_u^2 + \mu_d^2} , \tag{4.21}$$

$$v^2 = v_u^2 + v_d^2 = 2\frac{\mu_d^2 - \mu_u^2 - (\mu_d^2 + \mu_u^2)\cos 2\theta}{(g_2^2 + g_Y^2)\cos 2\theta} , \tag{4.22}$$

where $v_u = \langle H_u^0 \rangle$, and $v_d = \langle H_d^0 \rangle$, and the angle θ is defined as $\cot\theta = v_u/v_d$. Since the W-gauge-boson mass is given by $M_W^2 = \frac{1}{2} g_2^2 v^2$, the condition (4.22) is converted to

$$\cos 2\theta = \frac{\mu_d^2 - \mu_u^2}{\mu_d^2 + \mu_u^2 + M_Z^2} , \tag{4.23}$$

with M_Z being the Z-boson mass. It is noted that the allowed range of $\cos 2\theta$ is $0 < \cos 2\theta < 1$, since RG equations for μ_d^2 and μ_u^2 claim that $[\mu_d^2(M_W) - \mu_u^2(M_W)]$ is always positive.

Further we must impose two more conditions:[19]

$$(\mu_3)^4 > \mu_u^2 \mu_d^2 , \tag{4.24}$$

$$2|\mu_3|^2 < \mu_u^2 + \mu_d^2 . \tag{4.25}$$

The former assures nonexistence of a lower symmetric minimum at the scale M_W, and the latter is required in order that the potential V_H in Eq. (4.18) may be bounded below. The condition (4.25) must be satisfied all over the energy range from M_I to M_W. Meanwhile all the parameters and couplings in Eqs. (4.21)–(4.24) should be evaluated at the scale M_W.

E. The low-energy particle spectra

We write down some formulas for the particle masses which appear in the model. Most of them have already been given in the literature, but they are presented here for the sake of discussions in the next section. It is understood that the following mass and coupling parameters are the ones evaluated at the scale M_W, unless they are specified.

Let us start with fermion masses in the model. Fermions consists of quarks, leptons, gauginos, and Higgsinos. The top-quark mass is given by

$$m_t = h_t v_u . \tag{4.26}$$

There exist an upper bound for m_t (Refs. 19, 29, and 31):

$$m_t \le h_t v \le 4\pi \left[\frac{E_t(M_I)\tilde{E}_t(M_W)}{12[F_t(M_I) + E_t(M_I)\tilde{F}_t(M_W)]} \right]^{1/2} \times \left[\frac{2\sin^2\theta_W M_W^2}{e^2(M_W)} \right]^{1/2} . \tag{4.27}$$

where $E_t(M_I)$, $F_t(M_I)$, $\tilde{E}_t(M_W)$, and $\tilde{F}_t(M_W)$ are defined in Appendixes A and B. It is noted that the RHS of (4.27) is independent of the initial value $h_t(M_C)$, M, and m. The model considered here gives the upper bound 209 GeV for m_t. Meanwhile, we cannot tell much about the masses of the other quarks and leptons, since we have neglected their Yukawa couplings. Neutrinos obtain tiny masses through the well-known "seesaw" mechanism.[23]

All gluinos ($\tilde{g}$) have the same Majorana mass M_3. On the other hand, the W-inos and the B-inos will mix with Higgsinos. The charged W-inos and Higgsinos combine to form a couple of Dirac fermions (charginos $\tilde{\chi}_1^{\pm}$ and $\tilde{\chi}_2^{\pm}$) with masses[38]

$$m^2_{\tilde{\chi}_1^{\pm},\tilde{\chi}_2^{\pm}} = \tfrac{1}{2}\{M_2^2 + \mu^2 + 2M_W^2 \pm [(M_2^2 - \mu^2)^2 + 4M_W^4 \cos^2 2\theta + 4M_W^2(M_2^2 + \mu^2 + 2M_2\mu \sin 2\theta)]^{1/2}\} . \tag{4.28}$$

In the limit of M_2, $\mu \gg M_W$, we have $m_{\tilde{\chi}_1^{\pm},\tilde{\chi}_2^{\pm}} \simeq M_2$ and μ. When M_2, $M_W \ll \mu$, the chargino masses are given approximately by

$$m_{\tilde{\chi}_1^{\pm}} \approx M_2 - \left[\frac{M_W^2}{\mu} \right] \sin 2\theta , \tag{4.29a}$$

$$m_{\tilde{\chi}_2^{\pm}} \approx \mu , \tag{4.29b}$$

where we have denoted a lighter chargino as $\tilde{\chi}_1^{\pm}$ and a heavier one as $\tilde{\chi}_2^{\pm}$. Thus we may have one light chargino. Concerning the neutral gaugino-Higgsino sector, the neutral W-ino $\tilde{W}_3$, the B-ino $\tilde{B}$, and two Higgsinos $\tilde{H}_d^0$ and $\tilde{H}_u^0$ will combine to form four neutralinos ($\tilde{\chi}_1^0$, $\tilde{\chi}_2^0$, $\tilde{\chi}_3^0$, and $\tilde{\chi}_4^0$). The relevant 4×4 mass matrix is given as follows:[38]

$$\begin{array}{c|cccc|} & \tilde{W}_3 & \tilde{B} & \tilde{H}_d^0 & \tilde{H}_u^0 \\ \hline \tilde{W}_3 & M_2 & 0 & M_W\sin\theta & -M_W\cos\theta \\ \tilde{B} & 0 & M_Y & -M_W\tan\theta_W\sin\theta & M_W\tan\theta_W\cos\theta \\ \tilde{H}_d^0 & M_W\sin\theta & -M_W\tan\theta_W\sin\theta & 0 & -\mu \\ \tilde{H}_u^0 & -M_W\cos\theta & M_W\tan\theta_W\cos\theta & -\mu & 0 \end{array} . \tag{4.30}$$

In the case of M_2, M_Y, $\mu \gg M_W$, four eigenvalues of the above mass matrix are given approximately by M_2, M_Y, and $\pm\mu$. Then all the neutralinos get large masses. Meanwhile, when $M_2, M_Y, M_W \ll \mu$, we obtain the eigenvalues

$$M_2 - \frac{M_W^2}{\mu}\sin 2\theta, \quad M_Y - \frac{M_W^2}{\mu}\sin 2\theta \tan^2\theta_W, \quad \pm\mu , \tag{4.31}$$

and hence we may have two light and two heavy neutralinos.

Let us turn to the spectra of scalar particles. There have been four neutral Higgs bosons in the model. One is swallowed up by the Z gauge boson. The masses of the other three neutral Higgs bosons are given by the well-known formula[19]

$$m_{H_c}^2 = \mu_d^2 + \mu_u^2 , \tag{4.32a}$$

$$m^2_{H_a,H_b}=\tfrac{1}{2}[m^2_{H_c}+M_Z^2 \pm\sqrt{(m^2_{H_c}+M_Z^2)^2-4m^2_{H_c}M_Z^2\cos^2 2\theta}] \ . \tag{4.32b}$$

Since we expect that $m^2_{H_c}$ is very heavy, we will have one very light neutral Higgs boson (H_b). In fact, in the limit of $m^2_{H_c} \gg M_Z^2$, we obtain

$$m_{H_b}=M_Z\cos 2\theta \ . \tag{4.33}$$

Concerning the four real fields contained in the charged Higgs bosons H_u^+ and H_d^-, two are swallowed by the $W^\pm$ gauge bosons, and the other two form a complex charged scalar with a mass[19]

$$m^2_{H^\pm}=M_W^2+\mu_d^2+\mu_u^2 \ . \tag{4.34}$$

For the masses of all squarks and sleptons except for t squarks and sneutrinos we obtain much the same formulas as those given in Appendix C of Ref. 31. The only difference exists in the terms of the gaugino contributions to the scalar masses, since our model has the ISSB. If the initial value of m or M (or both) is too large compared to M_Z, these squark and slepton masses are of the same order of m or M, and are very massive. In the case of the t squarks, $\tilde{t}_L$ and $\tilde{t}_R$ form a 2×2 mass matrix since the $\tilde{t}_L$-$\tilde{t}_R$ mixing terms cannot be neglected.[20,29] The eigenvalues of this matrix are

$$m^2_{\tilde{t}_{h,l}}=m_t^2 +\tfrac{1}{2}((m_Q^2+m^2_{t_R}-\tfrac{1}{2}\cos 2\theta M_Z^2) \pm\{[m_Q^2-m^2_{t_R}+\cos 2\theta M_Z^2(-\tfrac{1}{2}+\tfrac{4}{3}\sin^2\theta_W)]^2 +4m_t^2(A_t m+\mu\tan\theta)^2\}^{1/2}) \ , \tag{4.35}$$

where the expressions for m_Q^2 and $m^2_{t_R}$ are given in Appendix B.

Finally we discuss about the masses of sneutrinos. Since we expect the relevant Yukawa coupling for neutrinos to be extremely small, the $\tilde{\nu}_L$-$\tilde{\nu}_R$ mixings are negligible. Then $\tilde{\nu}_L$ and $\tilde{\nu}_R$ are considered to be the physical eigenstates. The $\tilde{\nu}_L$ mass has the similar expression as the other sleptons. As for the $\tilde{\nu}_R$ mass, new contributions come from the term $m_{\nu_R}\nu_R\nu_R$ in the superpotential W in Eq. (4.14) and we obtain

$$m^2_{\tilde{\nu}_R}=m^2-M^2[2f_{1A}(M_I)+\tfrac{1}{12}f_{1B}(M_I)]+4m^2_{\nu_R} +mBm_{\nu_R} \tag{4.36}$$

where the expressions for $f_{1A}(M_I)$ and $f_{1B}(M_I)$ are given in Appendix A. Because of the large Majorana mass m_{ν_R}, we expect that $m_{\tilde{\nu}_R}$ should be no lighter than 1 TeV.

V. NUMERICAL ANALYSIS

A. Parameters A, m, and R

First, we have investigated numerically the range of the parameters M, m, and A which are consistent with the conditions for the ISSB. The condition (4.9) provides us with a relation among the parameters M, m, A, and $Y_\kappa(M_C)$. This enables us to calculate the necessary $Y_\kappa(M_C)$ for the symmetry breaking at M_I when a set of numerical values are assigned to M, m, and A. The acquired $Y_\kappa(M_C)$ must be real and positive, and also be small enough to be justified for the perturbative approximations. Furthermore, it must be compatible with the constraint (4.7a) along with a given set of M, m, and A. In the following numerical analysis we have put a criterion

$$Y_\kappa(M_C)<0.02 \ , \tag{5.1}$$

for the justification of the perturbative calculations.

These constraints and condition for the ISSB severely restrict the allowed domain of parameters M, m, and A. Figures 1(a)–1(c) show the possible range of the parameters A and m for several fixed values of $R=M/m$ ($R=0,2,5$). As $|A|$ gets smaller, larger values of $Y_\kappa(M_C)$ are necessary to satisfy the condition (4.9) and they may conflict with the criterion (5.1). Hence the allowed values for A fall into two separate regions. We have found that one should have at least

$$1.5<|A|\leq 3 \ , \tag{5.2}$$

for the consistent ISSB. The plots also show that the possible values for m are bounded from below. This is because with smaller values of m, the condition (4.9) requires larger values of $Y_\kappa(M_C)$ and the constraint (4.7a) or the bound (5.1) or both are violated. The close numerical analysis demonstrates that one should have

$$m>1.1\times 10^3 \text{ GeV} \ . \tag{5.3}$$

In Figs. 1(a)–1(c) we have represented the plots only for several non-negative values of R. The plots for negative R are easily obtained from those with positive $|R|$ by exchanging values of A for $-A$. This is due to the fact that the condition (4.9) and, thus, the value of $Y_\kappa(M_C)$ solved from (4.9) are invariant under the simultaneous replacement $A\to -A$ and $R\to -R$. This symmetry may be manifestly exhibited when we plot the allowed domain for the ISSB in the A-R plane with fixed values of m. Figures 2(a)–2(c) show the range of possible values for A and R with several fixed values of m ($m=1.5\times 10^3$, 2×10^3, and 10^4 GeV). It is apparent from these plots that the allowed domain for A and R are symmetric under 180-degree rotation around the origin.

When m is close to the bound value 1.1×10^3 GeV, the constraint (4.7a) may not be satisfied if R is small. For example, with $m=1.5\times 10^3$ GeV, the range of $-1.5<R<1.5$ is excluded [see Fig. 2(a)]. On the other hand, when $|R|$ gets too large, the criterion (5.1) may not be satisfied. In fact, we have found

$$|R|<8.8 \ , \tag{5.4}$$

in order to have the satisfactory symmetry breaking at the intermediate scale. Thus, the model considered in this paper puts strong restrictions, i.e., Eqs. (5.2)–(5.4), on the allowed values for the parameters A, m, and R. This is one of the features of this model.

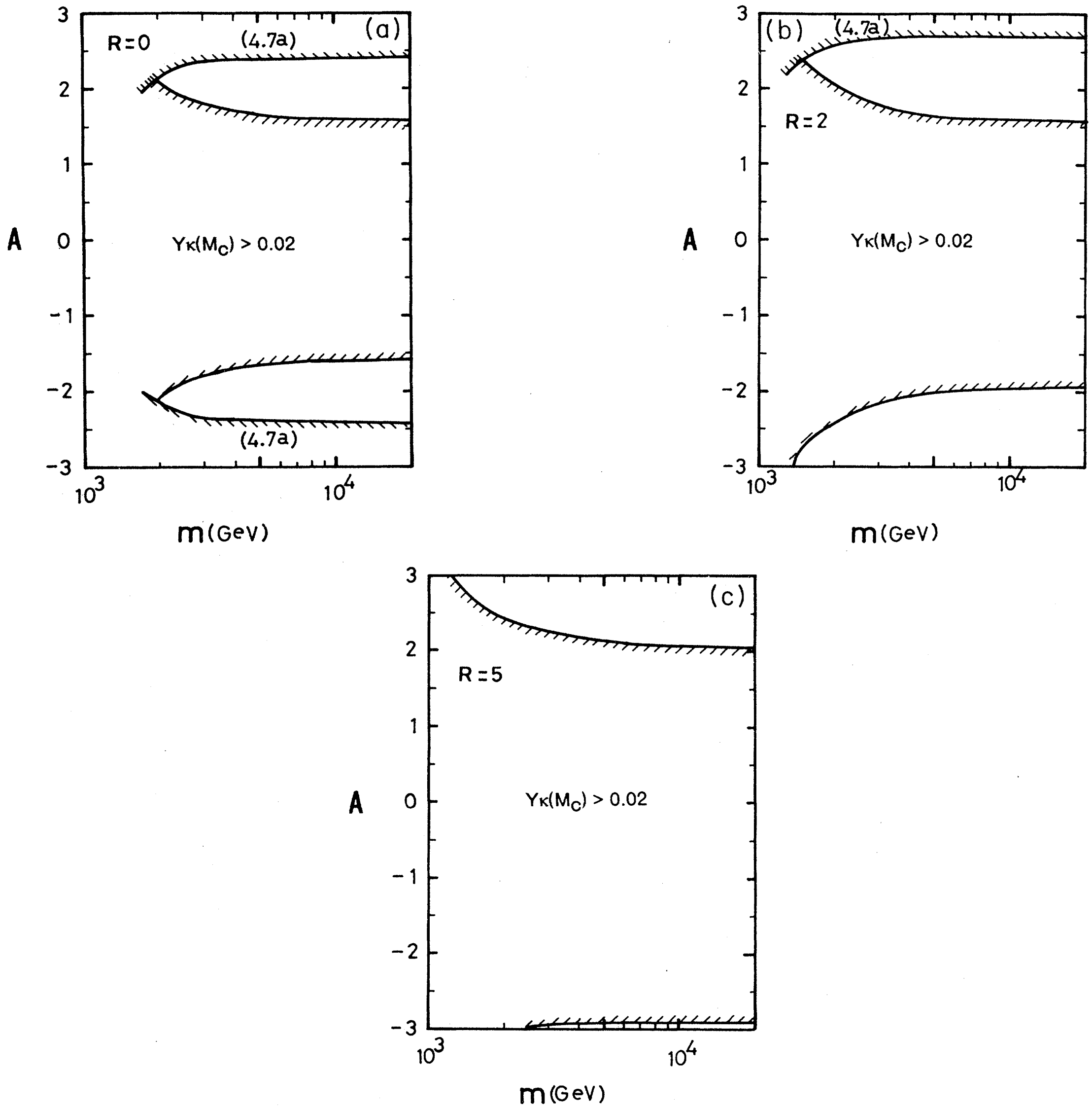

FIG. 1. Values of A and m consistent with the radiative symmetry breaking at the intermediate scale for fixed R: (a) $R=0$, (b) $R=2$, and (c) $R=5$. The hatched areas are forbidden and (4.7a) in the plots indicates regions where the constraint (4.7a) is not satisfied.

B. Top-quark mass

We have next investigated numerically the $SU(2)_W \times U(1)_Y$ breaking at the scale M_W. Since we have found above that m should be larger than 1 TeV, it is expected that the scalar partners of the relevant superfields may possess masses around or more than 1 TeV at low energies. Also gauginos could be very massive. Then we must stop the evolution of these mass parameters before they reach the scale M_W. Moreover, we should turn off the contribution of these particles to the RG equations for the gauge coupling constants and the parameters Y_t and A_t when energy scale has lowered to their masses. In the actual calculation, however, we have evaluated all the relevant parameters (including masses and gauge couplings) down at the M_W scale. This procedure simplifies computations considerably and does not have in general much influence on the numerical results except for gluino and squark masses. For example, if we stop the evolution of gluino mass M_3 at 1 TeV, we obtain a smaller value than the one evaluated at the scale M_W by 15%. Also when the parameter M is comparable to m, the evaluation of squark masses at 1 TeV gives smaller values than those at M_W by as far as 30%. However, in a case in which $|R| \ll 1$, the underestimation of squark masses reduces to less than 10%. As for the masses of other particles, such as, sleptons, Higgs scalars, charginos, and neutralinos, the differences between the values evaluated at 1 TeV and at M_W come within 10%. We may already have this order of uncertainties anyway by neglecting two-loop corrections and replacing all threshold effects of heavy particles with θ functions in the RG equations.

As we stated in Sec. IV C, unless we assume $B = A - 1$, the constant B in Eq. (4.19) is a free parameter. In this case Eq. (4.21), which is one of the conditions for the $SU(2)_W \times U(1)_Y$ breaking at the scale M_W, does not serve as a restriction upon the parameters $Y_t(M_C)$, A, R, and m. Indeed Eq. (4.21) can always be satisfied by adjusting the value of the B parameter. On the other hand, another condition for the $SU(2)_W \times U(1)_Y$ breaking, i.e., Eq. (4.23), will provide us with a relation among parameters. When we combine this condition with the restrictions

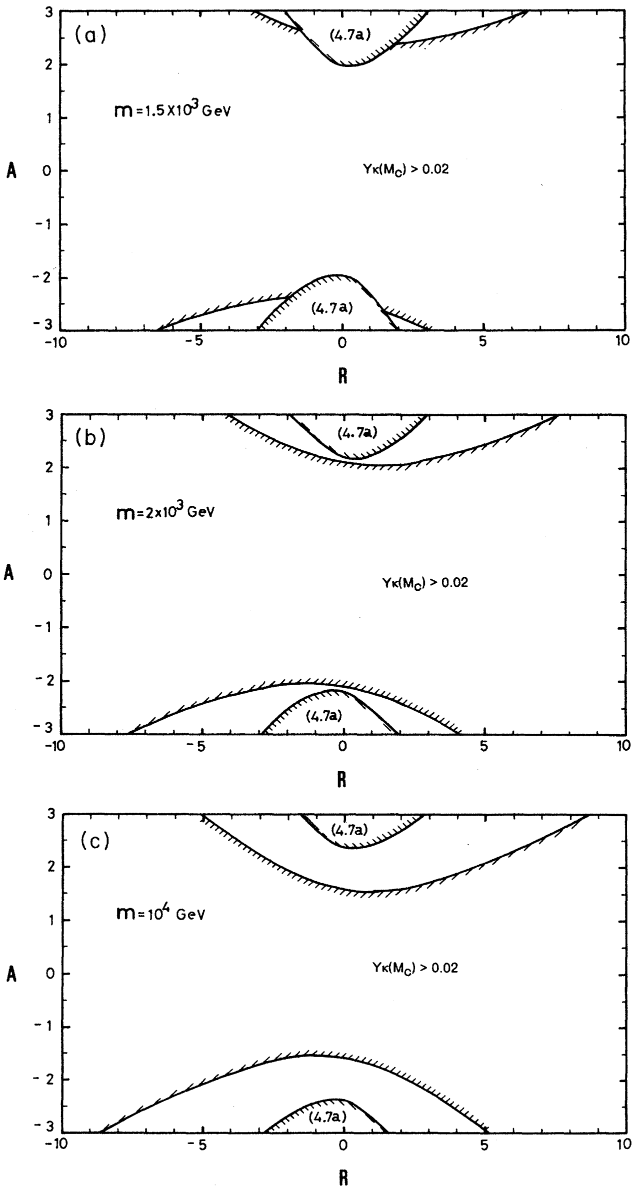

FIG. 2. Values of A and R consistent with the radiative symmetry breaking at the intermediate scale for fixed m: (a) $m = 1.5\times10^3$ GeV, (b) $m = 2\times10^3$ GeV, and (c) $m = 10^4$ GeV. The hatched areas are forbidden and (4.7a) in the plots indicates regions where the constraint (4.7a) is not satisfied.

that the parameters A and m should fall into the allowed region for the ISSB [i.e., Eqs. (5.2) and (5.3)], we can find an extent of possible values of top-quark mass $m_t(M_W)$ for given parameters R and $\cos2\theta$.

We represent in Figs. 3(a)–3(c) typical plots showing the allowed range of $m_t(M_W)$ vs $\cos2\theta$ for several fixed values of R ($R = -1,0,1$). The plots show that the $SU(2)_W \times U(1)_Y$ breaking requires lighter top-quark mass for smaller values of $\cos2\theta$. Actually, the symmetry can be broken with an arbitrary small $m_t(M_W)$ by having a smaller and smaller $\cos2\theta$. As $|R|$ decreases with $\cos2\theta$ fixed, larger top-quark mass is required. These are well-known features of the usual $N=1$ supergravity GUT models in which the $SU(2)_W \times U(1)_Y$ symmetry is broken radiatively.[31]

The upper bound for $m_t(M_W)$ is obtained when $\cos2\theta = 1$ and $R = 0$. We have found

$$m_t(M_W) < 148 \text{ GeV} , \tag{5.5}$$

which is the result of the "nonminimal" case (i.e., without the assumption of $B = A - 1$). It should be remarked, however, that when $\cos2\theta \to 1$ (i.e., $v_d/v_u \to 0$), we cannot neglect anymore the Yukawa coupling constant h_b for the bottom quark since in this limit the finiteness of the bottom-quark mass requires $h_b \to \infty$. Inclusion of a non-negligible h_b facilitates the $SU(2)_W \times U(1)_Y$ breaking and we would have obtained a smaller upper bound for $m_t(M_W)$ than that of Eq. (5.5).

There are some other points of interest in Figs. 3(a)–3(c). The allowed domain for $m_t(M_W)$ is made up of two regions, one corresponding to $A > 0$ and the other $A < 0$. When R is positive, the $SU(2)_W \times U(1)_Y$ breaking claims a heavier top-quark mass in the case $A < 0$ than in the case $A > 0$. The situation changes conversely, when R is negative. With large $|R|$, two allowed regions for $m_t(M_W)$ are separated. But they get closer as $|R|$ decreases, and finally they coincide when $R = 0$. These results relate to the fact that Eq. (4.23), i.e., one of the conditions for the $SU(2)_W \times U(1)_Y$ breaking, is invariant under the simultaneous replacement $A \to -A$ and $R \to -R$.

Now we examine the case in which the same relationship $B = A - 1$ holds as in the case of the "minimal" $N=1$ supergravity model. This time, Eq. (4.21) serves as another relation among the parameters. Hence, when a set of suitable numerical values are assigned to R, A, and $\cos2\theta$, one can obtain from Eqs. (4.21) and (4.23) the appropriate $m_t(M_W)$ and m for the $SU(2)_W \times U(1)_Y$ breaking. We represent in Figs. 3(a)–3(c) typical plots for the "minimal" case (shaded areas) showing the allowed range of $m_t(M_W)$ vs $\cos2\theta$ for several fixed values of R ($R = -1,0,1$).

It is noteworthy from these plots that the allowed regions for $m_t(M_W)$ and $\cos2\theta$ in the case of positive A are quite different from those for negative A. When $A > 0$ (actually, $1.5 < A \leq 3$), $\sin2\theta$ takes a negative value and its absolute value should stay small for the $SU(2)_W \times U(1)_Y$ breaking. Thus the possible values of $\cos2\theta$ are restricted to the range

$$\cos2\theta > 0.6 \quad \text{for } B = A - 1 \text{ and } 1.5 < A \leq 3 . \tag{5.6}$$

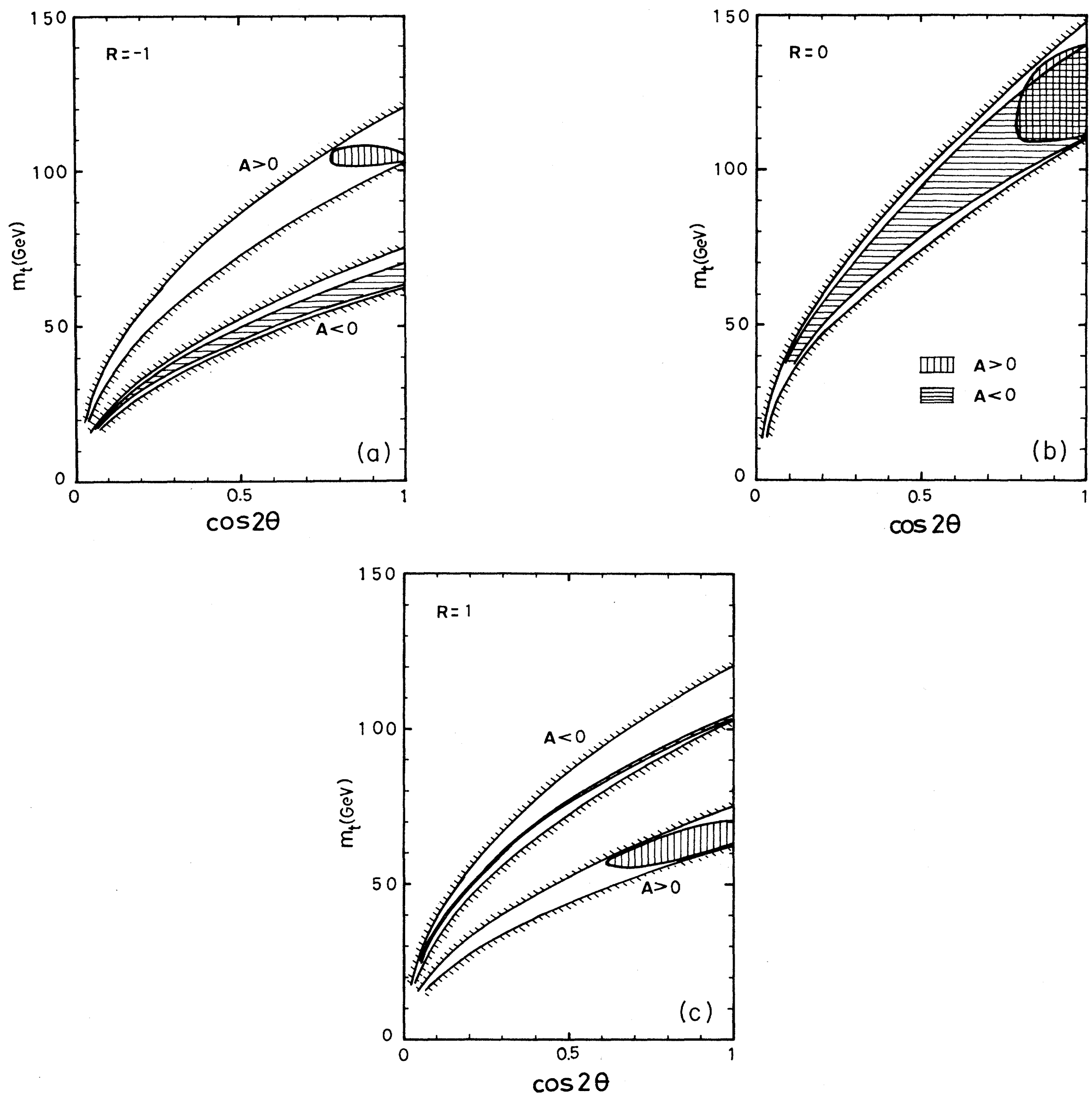

FIG. 3. Allowed values of $m_t(M_W)$ vs cos2θ with and without the assumption $B=A-1$. Results are shown for (a) $R=-1$, (b) $R=0$, and (c) $R=1$. The hatched areas (oblique lines) are forbidden. The shaded areas (vertical and horizontal lines) are allowed regions in the case of $B=A-1$.

In consequence, the allowed values of $m_t(M_W)$ are bounded from both below and above. Examining other cases of R, we have found, for the "minimal" case with positive A,

$$47\ \text{GeV} < m_t(M_W) < 141\ \text{GeV} \quad \text{for } B=A-1 \text{ and } 1.5 < A \le 3 . \tag{5.7}$$

The top-quark mass $m_t(M_W)$ takes the largest value at $R=0$ and the lowest around $R=2$.

On the other hand, for negative A $(-3 \le A < -1.5)$, cos2θ takes any positive value and thus the $SU(2)_W \times U(1)_Y$ symmetry can be broken with an arbitrarily small $m_t(M_W)$. The allowed region for $m_t(M_W)$ is given by

$$m_t(M_W) < 141\ \text{GeV} \quad \text{for } B=A-1 \text{ and } -3 \le A < -1.5 , \tag{5.8}$$

and the upper bound is obtained when $R=0$. In Figs. 4(a)–4(c) we plot the allowed range of $m_t(M_W)$ vs m with several fixed values of R $(R=-1,0,1)$ for the "minimal" case. These plots show how large the parameter m should be. Note that the bound in Eq. (5.3) is respected. Again the allowed domain in the m_t-m plane for positive A are quite distinct from those for negative A.

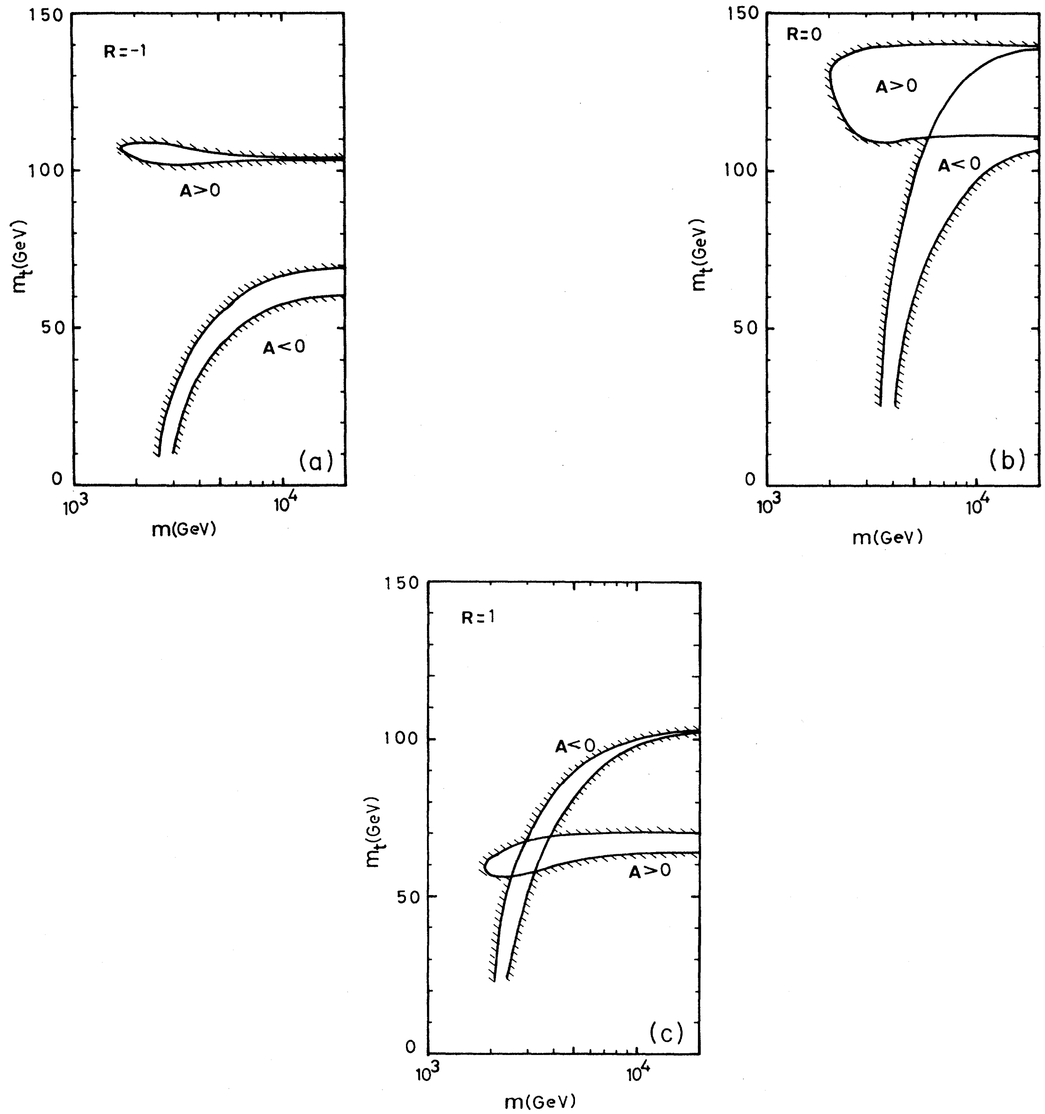

FIG. 4. Allowed values of $m_t(M_W)$ vs m in the assumption of $B = A - 1$. Results are shown for (a) $R = -1$, (b) $R=0$, and (c) $R=1$. The hatched areas are forbidden.

C. Low-energy particle spectra

We now proceed to discuss in some detail the low-energy spectra predicted by this model. We also refer to the existing experimental results on new particle searches and study whether further constraints could be placed upon the parameters in the model. As already explained in Sec. V A, the parameters m, R, and A are severely restricted by the conditions and constraints for the ISSB. Especially, m should be no lighter than 1.1 TeV. Thus all the squarks and sleptons are expected to be as massive as or heavier than 1 TeV.

Let us first examine the "minimal" case in which the assumption $B = A - 1$ enables us to make more specific predictions on the low-energy particle spectra than in general situations. There are three free parameters and we have chosen R, A, and $\cos2\theta$. Having assigned several sets of numerical values to these parameters, we obtained for each set the low-energy mass spectra as well as appropriate m and $m_t(M_W)$ for the $SU(2)_W \times U(1)_Y$ breaking. The results are shown in Table V.

Spectra (a) and (b) represent cases when $R \ll 1$. In particular, (a) typifies a situation where the parameter M $(=Rm)$ is small (≤ 50 GeV). R is set to be one in the rest of the examples which feature situations where M is the same order of magnitude with m, i.e., in the TeV region. Spectra (d)–(f) represent cases in which A is negative, and each has a different $\cos2\theta$.

The mass spectra for gaugino-related SUSY particles such as charginos ($\tilde{\chi}_1^{\pm}$ and $\tilde{\chi}_2^{\pm}$), neutralinos ($\tilde{\chi}_1^0$, $\tilde{\chi}_2^0$, $\tilde{\chi}_3^0$, and $\tilde{\chi}_4^0$), and gluinos ($\tilde{g}$) depend heavily on the parameter M. The gaugino masses evaluated at the scale M_W are related to M through the following numerical formulas:

TABLE V. Six examples of low-energy spectra for the "minimal" case with the assumption of $B=A-1$.

	a	b	c	d	e	f
R	0.008	0.05	1	1	1	1
A	2.1	2.1	2.1	-2.3	-2.3	-2.3
$\cos 2\theta$	0.9	0.9	0.9	0.9	0.5	0.15
m (TeV)	3.42	3.61	4.78	10.4	3.98	2.46
m_t (GeV)	122	117	63.6	99.5	76.3	44.0
$\tilde{g}$ (TeV)	8.40×10^{-2}	5.55×10^{-1}	14.7	31.9	12.2	7.56
$\tilde{\chi}_1^{\pm}$ (TeV)	2.35×10^{-2}	1.41×10^{-1}	1.28	1.23	1.25	1.27
$\tilde{\chi}_2^{\pm}$ (TeV)	1.18	1.20	3.70	8.05	3.08	1.91
$\tilde{\chi}_1^{0}$ (TeV)	1.17×10^{-2}	7.25×10^{-2}	1.28	1.23	1.24	9.66×10^{-1}
$\tilde{\chi}_2^{0}$ (TeV)	2.33×10^{-2}	1.41×10^{-1}	1.28	1.23	1.25	1.29
$\tilde{\chi}_3^{0}$ (TeV)	1.18	1.19	1.89	4.12	1.58	1.29
$\tilde{\chi}_4^{0}$ (TeV)	1.18	1.19	3.70	8.05	3.08	1.91
H_a (TeV)	3.71	3.91	6.09	13.0	5.75	4.29
H_b (GeV)	83.3	83.3	83.3	83.3	46.3	13.9
H_c (TeV)	3.71	3.91	6.09	13.0	5.75	4.29
$H^{\pm}$ (TeV)	3.71	3.91	6.09	13.0	5.75	4.29
$\tilde{t}_h$ (TeV)	2.77	2.97	13.9	30.3	11.6	7.30
$\tilde{t}_l$ (TeV)	1.84	2.02	13.2	28.7	11.1	7.08

$$M_3(M_W)=3.07M \ , \qquad (5.9a)$$

$$M_2(M_W)=0.773M \ , \qquad (5.9b)$$

$$M_Y(M_W)=0.396M \ . \qquad (5.9c)$$

In the case $|M_2|, M_W \ll \mu$ [spectra (a) and (b)], one chargino ($\tilde{\chi}_1^{\pm}$) is light and its mass is approximately given by the formula Eq. (4.29a). For the heavier one ($\tilde{\chi}_2^{\pm}$) we have $m_{\tilde{\chi}_2^{\pm}} \simeq \mu$ (≈ 1 TeV). Experimentally a lower bound $m_{\tilde{\chi}_1^{\pm}} > 22.5$ GeV for the chargino mass has been reported.[39] When one combines this experimental result with Eq. (4.29a), one gets a new constraint on the parameter M. Remembering that $-0.8 < \sin 2\theta < 0$ for $A>0$ [see Eq. (5.6)] and $\sin 2\theta > 0$ for $A<0$, as well as $\mu \approx 1$ TeV, we obtain a lower bound $|M_2| > 17$ GeV. Using Eq. (5.9b), we find a limit

$$|M| > 22 \text{ GeV} \ . \qquad (5.10)$$

Spectra (a) serves as an example in which $\tilde{\chi}_1^{\pm}$ has a mass very close to the experimental limit value. In the situations when $|M_2|, \mu \gg M_W$ [spectra (c)–(f)], chargino masses are given by $m_{\tilde{\chi}_1^{\pm}} \simeq \mu$ and $m_{\tilde{\chi}_2^{\pm}} \simeq |M_2|$, and hence both lie in TeV region. It is interesting to note that we always have one chargino with a mass around 1 TeV.

As for the neutralinos, we have two light and two heavy ones when $|M_2|, |M_Y|, M_W \ll \mu$ [spectra (a) and (b)]. In this case we find from Eq. (4.31) that the masses for two light neutralinos (which we denote as $\tilde{\chi}_1^0$, and $\tilde{\chi}_2^0$ with $\tilde{\chi}_1^0 < \tilde{\chi}_2^0$) are approximately given by

$$m_{\tilde{\chi}_1^0} \simeq |M_Y| - \frac{M_W^2}{\mu} \sin 2\theta \tan^2\theta_W \ , \qquad (5.11a)$$

$$m_{\tilde{\chi}_2^0} \simeq |M_2| - \frac{M_W^2}{\mu} \sin 2\theta \ , \qquad (5.11b)$$

and two heavy ones both have mass around μ. Then using Eqs. (5.9b) (5.9c), and (5.10) we obtain lower bounds

$$m_{\tilde{\chi}_1^0} > 6.8 \text{ GeV} \ , \qquad (5.12a)$$

$$m_{\tilde{\chi}_2^0} > 11 \text{ GeV} \ , \qquad (5.12b)$$

for the two light neutralinos. When $\tilde{\chi}_1^0$ is light, it is mostly made up of the photino ($\tilde{\gamma}$) state. There is a report on the photino mass limit from the single-photon production experiment at SLAC (Ref. 40), where the lower bound for $m_{\tilde{\gamma}}$ was obtained under the assumption of pure $\tilde{\gamma}$ eigenstate and equal mass for the right- and left-handed selectron ($\tilde{e}$). But it has been concluded there that there is no limit on $m_{\tilde{\gamma}}$ if selectron mass is larger than 58 GeV. The model studied in this paper indeed predicts very massive (> 1 TeV) selectrons.

In the case $|M_2|, |M_Y|, \mu \gg M_W$, four neutralino masses are approximately given by $|M_2|, |M_Y|, \mu$ and μ, and they are all massive [spectra (c)–(f)]. As is easily seen from Table V, $\tilde{\chi}_1^0$ is not only the lightest neutralino but also the lightest SUSY particle in this model.

The gluino mass $M_{\tilde{g}} = |M_3|$ is given by Eq. (5.9a) and the following lower bound is obtained from Eq. (5.10):

$$M_{\tilde{g}} > 68 \text{ GeV} \ . \qquad (5.13)$$

Experimentally, UA1 group[41] ruled out the range 4–53 GeV for the gluino mass (at 90% C.L. under the assumption that all squarks are heavier than gluinos) from the SUSY particle searches in the $p\bar{p}$ collider at CERN.

There are five Higgs particles in the low-energy spectra: three neutral (H_a, H_b, H_c) and two charged ($H^{\pm}$). Table V shows that four of them (H_a, H_c, and $H^{\pm}$) have almost the same mass and are very massive ($\geq$ a few TeV). On the other hand, H_b has a lighter mass than the

TABLE VI. Summary of the SI SO(10) GUT model studied in this paper: the allowed range of the supergravity parameters and the model predictions on the low-energy particle spectra.

	"Minimal" case with $B=A-1$	General case
m	$m>1.1$ TeV	$m>1.1$ TeV
A	$1.5<\|A\|\leq 3$	$1.5<\|A\|\leq 3$
R	$\|R\|<8.8$	$\|R\|<8.8$
$\cos 2\theta$	$0.6<\cos 2\theta<1$ (for $A>0$) $\cos 2\theta<1$ (for $A<0$)	$\cos 2\theta<1$
t quark	47 GeV $<m_t<$ 141 GeV (for $A>0$) $m_t<141$ GeV (for $A<0$)	$m_t<148$ GeV
$\tilde{g}$	$m_{\tilde{g}}>68$ GeV[a]	$m_{\tilde{g}}>63$ GeV[a]
$\tilde{\chi}_1^0$	$m_{\tilde{\chi}_1^0}>6.8$ GeV[a]	$m_{\tilde{\chi}_1^0}>6.1$ GeV[a]
$\tilde{\chi}_2^0$	$m_{\tilde{\chi}_2^0}>11$ GeV[a]	$m_{\tilde{\chi}_2^0}>9.3$ GeV[a]
H_b	$M_Z\cos 2\theta$	$M_Z\cos 2\theta$
All $m_{\tilde{q}}$ and $m_{\tilde{l}}$ $\tilde{\chi}_3^0,\tilde{\chi}_4^0$	>1 TeV	>1 TeV
$H_a,H_c,H^{\pm}$	$>$a few TeV	$>$a few TeV

[a]In deriving the lower bounds, the experimental result $\tilde{\chi}_1^{\pm}>22.5$ GeV is used.

Z boson. It is true that the usual supersymmetric GUT models predict a neutral Higgs boson whose mass is less than M_Z. But in this model we have a very massive H_c and, hence, the H_b mass is effectively given by Eq. (4.33), which is easily ascertained in Table V. If we respect the UA1 experimental result on the top-quark mass bound $m_t>44$ GeV, then referring to Fig. 3 for the allowed domains of m_t and the parameter $\cos 2\theta$ we obtain $\cos 2\theta>0.1$. Combining this constraint with Eq. (4.33), we get a bound

$$m_{H_b}>9.1 \text{ GeV}\ , \tag{5.14}$$

for the lightest neutral Higgs boson H_b, which is still above the experimentally acquired lower bounds.[28]

We cannot say much about the squark and slepton masses in this model. They are all very massive (>1 TeV). We have listed in Table V some examples of predicted mass spectra for the top-squarks $\tilde{t}_h$ and $\tilde{t}_l$. When $|R|\geq 1$ all squark masses including $m_{\tilde{t}_h}$ and $m_{\tilde{t}_l}$ are roughly given by a single formula $m_{\tilde{q}}\approx m(1+7.5R^2)^{1/2}$ [see spectra (c)–(f)], and they are heavier than sleptons. In the case $|R|\simeq 0$, $\tilde{t}_h$ and $\tilde{t}_l$ have smaller masses than other squarks and sleptons due to a non-negligible Yukawa coupling.

We now discuss the low-energy particle spectra in the general situation, i.e., when the assumption $B=A-1$ is not made. Since B is a free parameter in this case, even when one fixes the parameters R, A, and $\cos 2\theta$ one cannot solve m (and hence M) uniquely from the $SU(2)_W\times U(1)_Y$-breaking conditions. But m should observe the bound (5.3) and be very massive. We, therefore, expect to have spectra which are similar to those in the case $B=A-1$. Squarks and sleptons are all very massive (heavier than 1 TeV).

As for the gaugino-related SUSY particles, it could happen again that $\tilde{\chi}_1^{\pm},\tilde{\chi}_1^0,\tilde{\chi}_2^0$ and $\tilde{g}$ are light. The lower bounds for the masses of the last three SUSY particles are obtained in the similar way as before. Since there is no constraint on $\sin 2\theta$ in the general case, we get from Eq. (4.29a) and the experimental result $m_{\tilde{\chi}_1^{\pm}}>22.5$ GeV, a bound $|M_2|>15.9$ GeV. This gives through Eqs. (5.11a), (5.11b), and (5.9a)–(5.9c) the following lower bounds for $\tilde{\chi}_1^0,\tilde{\chi}_2^0$ and $\tilde{g}$: $m_{\tilde{\chi}_1^0}>6.1$ GeV, $m_{\tilde{\chi}_2^0}>9.3$ GeV, and $m_{\tilde{g}}>63$ GeV. Other gaugino-related SUSY particles are very heavy with masses in TeV region.

In the general case, we have again one light neutral Higgs boson H_b whose mass is given by Eq. (4.33) and four heavy Higgs bosons $H^{\pm}$, H_a, and H_c whose masses are almost the same and lie in the TeV region. Before finishing up this section, we summarize in Table VI the results on the numerical analysis of the SI SO(10) GUT model studied in this paper.

VI. CONCLUSIONS AND DISCUSSION

We have presented in this paper a superstring-inspired SO(10) GUT model in which the mixing of the Weinberg-Salam supersymmetric Higgs doublets H_d and H_u and large Majorana masses for right-handed neutrinos are generated by the higher-dimensional terms in the superpotential at the intermediate scale symmetry breaking. We fix the mixing strength parameter μ for H_d and H_u as $\mu=10^3$ GeV, expecting that it would be of the same order as Majorana-neutrino masses. Introducing into the model the usual general soft terms of $N=1$ supergravity, we have examined a scenario that symmetry breakings are induced by radiative corrections both at the intermediate scale M_I and at the Weinberg-Salam scale M_W.

Although the model turns out to have five parameters

m, M, A, B, and h_t [or $m_t(M_W)$], the conditions and constraints for the symmetry breakings severely restrict allowed regions of these parameters. We find that in order to generate the symmetry breaking satisfactorily at the intermediate scale one should have at least $1.5 < |A| \leq 3$, $m > 1.1$ TeV, and $|M/m| < 8.8$. Together with these restrictions for the parameter A, m, and R $(=M/m)$, the $SU(2)_W \times U(1)_Y$-breaking conditions give bounds on the top-quark mass. We have obtained $m_t < 148$ GeV. If we further assume the relation $B = A - 1$ as in the case of the "minimal" $N=1$ supergravity, we get $m_t < 141$ GeV. Actually, m_t is bounded also from below when A takes a positive value (i.e., $1.5 < A \leq 3$), and we find 47 GeV $< m_t < 141$ GeV.

We also examined the low-energy particle spectra consistent with the above restrictions. We studied two cases: the "minimal" case with $B = A - 1$ and the "nonminimal" case with B as a free parameter. In both cases we get similar spectra. All squarks and sleptons gain quite heavy masses which lie in TeV region. We may have some light gaugino-related SUSY particles, depending on the initial value M for gaugino masses. If M is small enough, then gluinos, two charginos, and two neutralinos turn out to be light. When we use the experimental result on the chargino mass $m_{\tilde{\chi}_1^\pm} > 22.5$ GeV, we obtain the lower bounds for the masses of two neutralinos and gluinos which are listed in Table VI. All the Higgs scalars are quite heavy with masses in TeV region except for the neutral Higgs boson H_b whose mass is given by the formula $m_{H_b} = M_Z \cos 2\theta$ and thus lighter than M_Z. In fact the usual standard SUSY GUT models always predict one light neutral Higgs boson H_b (Ref. 19) and its mass squared is in general written in terms of M_Z, m_{H_c}, and $\cos 2\theta$ such as in Eq. (4.32b). However, in this model m_{H_c} is quite heavy (> a few TeV) and hence the above simple expression is derived for m_{H_b}.

Most of the SUSY particles and Higgs scalars in this model are predicted to be too massive to be detected in the accelerator experiments before the advent of the Superconducting Super Collider (SSC). However, it is highly expected that the top quark and one light neutral Higgs boson will be soon observed. The discovery of these particles and the knowledge of their masses will further constrain the allowed ranges of the parameters, and may possibly clarify the viability of this model.

There are other points of interest in this SI SO(10) GUT model. Except for right-handed neutrinos, extra (or "exotic") fields do not appear in the low-energy spectra. In particular, the color-triplet fields g and $\bar{g}$ from the **10** multiplets, to which the light Higgs doublets belong, become all massive modes at the compactification. Hence the fast proton decays do not occur through g-quark exchange. Also the extra neutral vector boson Z' acquires a large mass due to the ISSB, and thus the well-known Z-Z' mixing problem is avoided.

It may be fair to comment also on the less attractive features of this model. First, we have introduced four Higgs fields H_{u1}, H_{d1}, H_{u2}, and H_{d2}. H_{u1} and H_{d1} are the usual Weinberg-Salam Higgs doublets and we have assumed that they do not couple to N_R and $\overline{N_R}$, so that H_{u1} and H_{d1} remain light after the ISSB. Meanwhile, H_{u2} and H_{d2} couple to N_R and $\overline{N_R}$, respectively, and we have assumed equal strength for these Yukawa coupling constants, i.e., $\kappa = \bar{\kappa}$ at the compactification scale M_C, so that N_R and $\overline{N_R}$ may have equal VEV's $\langle N_R \rangle = \langle \overline{N_R} \rangle$ at the intermediate scale. These two assumptions seem to be rather artificial and need to get some theoretical support from, for example, certain discrete symmetries.

Second, we have chosen $M_C = 1.5 \times 10^{16}$ GeV and $M_I = 1.0 \times 10^{14}$ GeV in performing the numerical analysis. Such a choice has been made since with these values for M_C and M_I, the predicted value for the Weinberg angle $\sin^2\theta_W(M_W)$ and the $SU(3)_C$ gauge coupling constant $\alpha_3(M_W)$ both fall into the phenomenologically acceptable range. However, the value of M_I seems too close to that of M_C. In other words we need to have a rather early ISSB after the compactification. Here it should be remarked that the above figures for M_C and M_I have been obtained by analyzing the evolution equations for the gauge coupling constants in one-loop approximations. It might be possible that inclusion of two-loop effects to the evolution equations would improve the situation; i.e., the interval between M_C and M_I would be widened. We have not studied this possibility yet.

In this paper we have correlated the $H_d H_u$ mixing with the light-neutrino mass problem. Meanwhile, Kim and Nilles[22] discussed the $H_d H_u$ mixing in conjunction with the strong CP problem. Also the incorporation of the light-neutrino masses and the resolution of the strong CP problem was proposed by Kang and Shin.[42] Then, it would be a very interesting possibility that the strong CP problem of quarks, the mixing problem of the Higgs fields H_d and H_u and the light-neutrino mass problem may be all intimately related to one another.

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APPENDIX A: EVOLUTION OF PARAMETERS ABOVE M_I

1. RG equations for parameters *from* M_C *to* M_I

We collect here the RG equations followed by the parameters in the SI SO(10) GUT model presented in Sec. III from the compactification scale M_C to the intermediate scale M_I. First, we define the quantities

$$t = \ln(M_C/E), \quad \tilde{\alpha}_\rho = g_\rho^2/(4\pi)^2 \quad (\rho = 3,\ 2,\ 1A,\ \text{and}\ 1B)\ , \tag{A1}$$

$$Y_t = h_t^2/(4\pi)^2, \quad Y_\kappa = \kappa^2/(4\pi)^2, \quad Y_{\bar{\kappa}} = \bar{\kappa}^2/(4\pi)^2\ . \tag{A2}$$

The gauge coupling constants are subject to the RG equations

$$\frac{d\tilde{\alpha}_\rho}{dt} = 2b_\rho \tilde{\alpha}_\rho^2 \quad (\rho = 3,\ 2,\ 1A,\ \text{and}\ 1B)\ , \tag{A3}$$

where

$$b_3=9-2N_f,\quad b_2=3-2N_f\ ,$$
$$b_{1A}=-5-10N_f/3,\quad b_{1B}=-35/12-10N_f/3 \tag{A4}$$

for N_f families. At the scale M_C (i.e., $t=0$) they obey

$$\tilde{\alpha}_3(M_C)=\tilde{\alpha}_2(M_C)=\tfrac{5}{3}\tilde{\alpha}_{1A}(M_C)=\tfrac{5}{3}\tilde{\alpha}_{1B}(M_C)\equiv\tilde{\alpha}_0\ . \tag{A5}$$

The gaugino masses evolve similarly as the gauge coupling constants:

$$\frac{M_\rho(t)}{\tilde{\alpha}_\rho(t)}=\text{const}\quad(\rho=3,\ 2,\ 1A,\ \text{and}\ 1B)\ , \tag{A6}$$

with the initial conditions at M_C,

$$M_3(M_C)=M_2(M_C)=M_{1A}(M_C)=M_{1B}(M_C)\equiv M\ . \tag{A7}$$

For the Yukawa couplings Y_t and Y_κ, we have

$$\frac{dY_t}{dt}=-12Y_t^2+Y_t(\tfrac{32}{3}\tilde{\alpha}_3+6\tilde{\alpha}_2+\tfrac{14}{9}\tilde{\alpha}_{1A}+\tfrac{7}{3}\tilde{\alpha}_{1B})\ , \tag{A8}$$

$$\frac{dY_\kappa}{dt}=-8Y_\kappa^2+Y_\kappa(6\tilde{\alpha}_2+6\tilde{\alpha}_{1A}+\tfrac{7}{3}\tilde{\alpha}_{1B})\ , \tag{A9}$$

and the RG equation for $Y_{\bar{\kappa}}$ is obtained from (A9) by replacing Y_κ with $Y_{\bar{\kappa}}$.

For the soft-breaking trilinear scalar couplings, we have

$$\frac{dA_t}{dt}=-12Y_tA_t+\frac{1}{m}(\tfrac{32}{3}\tilde{\alpha}_3M_3+6\tilde{\alpha}_2M_2+\tfrac{14}{9}\tilde{\alpha}_{1A}M_{1A}+\tfrac{7}{3}\tilde{\alpha}_{1B}M_{1B})\ , \tag{A10}$$

$$\frac{dA_\kappa}{dt}=-8Y_\kappa A_\kappa+\frac{1}{m}(6\tilde{\alpha}_2M_2+6\tilde{\alpha}_{1A}M_{1A}+\tfrac{7}{3}\tilde{\alpha}_{1B}M_{1B})\ , \tag{A11}$$

and the one for $A_{\bar{\kappa}}$ is obtained from (A11) by replacing A_κ and Y_κ with $A_{\bar{\kappa}}$ and $Y_{\bar{\kappa}}$, respectively. We assume the usual initial conditions at M_C:

$$A_t(M_C)=A_\kappa(M_C)=A_{\bar{\kappa}}(M_C)\equiv A\ . \tag{A12}$$

For the scalar masses, we have

$$\frac{dm_Q^2}{dt}=-2Y_t(m_Q^2+m_{t_R}^2+m_{m_{H_{u1}}}^2+m^2A_t^2)+\tfrac{32}{3}\tilde{\alpha}_3M_3^2+6\tilde{\alpha}_2M_2^2+\tfrac{2}{9}\tilde{\alpha}_{1A}M_{1A}^2+\tfrac{1}{3}\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A13}$$

$$\frac{dm_{t_R}^2}{dt}=-4Y_t(m_Q^2+m_{t_R}^2+m_{m_{H_{u1}}}^2+m^2A_t^2)+\tfrac{32}{3}\tilde{\alpha}_3M_3^2+\tfrac{8}{9}\tilde{\alpha}_{1A}M_{1A}^2+3\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A14}$$

$$\frac{dm_{H_{u1}}^2}{dt}=-6Y_t(m_Q^2+m_{t_R}^2+m_{m_{H_{u1}}}^2+m^2A_t^2)+6\tilde{\alpha}_2M_2^2+2\tilde{\alpha}_{1A}M_{1A}^2+\tfrac{4}{3}\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A15}$$

$$\frac{dm_{H_{d1}}^2}{dt}=6\tilde{\alpha}_2M_2^2+2\tilde{\alpha}_{1A}M_{1A}^2+\tfrac{4}{3}\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A16}$$

$$\frac{dm_{N_R}^2}{dt}=-4Y_\kappa(m_{N_R}^2+m_{L'}^2+m_{m_{H_{u2}}}^2+m^2A_\kappa^2)+8\tilde{\alpha}_{1A}M_{1A}^2+\tfrac{1}{3}\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A17}$$

$$\frac{dm_{L'}^2}{dt}=-2Y_\kappa(m_{N_R}^2+m_{L'}^2+m_{m_{H_{u2}}}^2+m^2A_\kappa^2)+6\tilde{\alpha}_2M_2^2+2\tilde{\alpha}_{1A}M_{1A}^2+3\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A18}$$

$$\frac{dm_{H_{u2}}^2}{dt}=-2Y_\kappa(m_{N_R}^2+m_{L'}^2+m_{M_{H_{u2}}}^2+m^2A_\kappa^2)+6\tilde{\alpha}_2M_2^2+2\tilde{\alpha}_{1A}M_{1A}^2+\tfrac{4}{3}\tilde{\alpha}_{1B}M_{1B}^2\ , \tag{A19}$$

The RG equations for $m_{\overline{N_R}}^2$, $m_{\overline{L'}}^2$ and $m_{H_{d2}}^2$ are obtained from (A17), (A18), and (A19), respectively, by replacement of $m_{N_R}^2$, $m_{L'}^2$, $m_{H_{u2}}^2$, Y_κ, and A_κ with $m_{\overline{N_R}}^2$, $m_{\overline{L'}}^2$, $m_{H_{d2}}^2$, $Y_{\bar{\kappa}}$, and $A_{\bar{\kappa}}$. The initial conditions for scalar masses at M_C are

$$m_i^2(M_C)=m^2\quad i=\text{all scalars}\ . \tag{A20}$$

2. Formulas for the renormalized parameters *above* M_I

Here we collect some formulas required for the expressions of the renormalized parameters above the scale M_I. We use the similar notations which have been given in Ref. 30. They are

$$\beta_\rho=2\tilde{\alpha}_\rho(M_C)b_\rho\ , \tag{A21}$$

$$e_\rho(t)=\frac{1}{1-\beta_\rho t}\ , \tag{A22}$$

$$f_\rho(t)=\frac{1}{b_p}[1-e_\rho(t)^2]\ , \tag{A23}$$

with $\rho=3$, 2, $1A$, and $1B$;

$$E_t(t)=e_3(t)^{16/3b_3}e_2(t)^{3/b_2}e_{1A}(t)^{7/9b_{1A}}e_{1B}(t)^{7/6b_{1B}}\ , \tag{A24}$$

$$E_\kappa(t)=e_2(t)^{3/b_2}e_{1A}(t)^{3/b_{1A}}e_{1B}(t)^{7/6b_{1B}}\ , \tag{A25}$$

$$F_t(t)=\int_0^t E_t(t')dt'\ , \tag{A26}$$

$$F_\kappa(t)=\int_0^t E_\kappa(t')dt'\ , \tag{A27}$$

$$D_t(t)=1+12Y_t(0)F_t(t)\ , \tag{A28}$$

$$D_\kappa(t)=1+8Y_\kappa(0)F_\kappa(t)\ , \tag{A29}$$

$$H_2^t(t)=t\tilde{\alpha}_0\frac{M}{m}[\tfrac{32}{3}e_3(t)+6e_2(t)+\tfrac{14}{9}(\tfrac{3}{5})e_{1A}(t)+\tfrac{7}{3}(\tfrac{3}{5})e_{1B}(t)]\ , \tag{A30}$$

$$H_2^\kappa(t)=t\tilde{\alpha}_0\frac{M}{m}[6e_2(t)+6(\tfrac{3}{5})e_{1A}(t)+\tfrac{7}{3}(\tfrac{3}{5})e_{1B}(t)]\ , \tag{A31}$$

$$H_3^t(t)=\int_0^t H_2^t(t')E_t(t')dt'=\frac{M}{m}[tE_t(t)-F_t(t)]\ , \tag{A32}$$

$$H_3^\kappa(t)=\int_0^t H_2^\kappa(t')E_\kappa(t')dt' =\frac{M}{m}[tE_\kappa(t)-F_\kappa(t)]\ , \tag{A33}$$

$$H_6^t(t)=\int_0^t (H_2^t(t'))^2E_t(t')dt'\ , \tag{A34}$$

$$H_6^\kappa(t)=\int_0^t (H_2^\kappa(t'))^2E_\kappa(t')dt'\ , \tag{A35}$$

$$G_1^t(t)=-M^2[\tfrac{16}{3}f_3(t)+3f_2(t)+\tfrac{7}{9}f_{1A}(t)+\tfrac{7}{6}f_{1B}(t)]\ , \tag{A36}$$

$$G_1^\kappa(t)=-M^2[3f_2(t)+3f_{1A}(t)+\tfrac{7}{6}f_{1B}(t)]\ , \tag{A37}$$

$$G_2^t(t)=\int_0^t G_1^t(t')E_t(t')dt'\ , \tag{A38}$$

$$G_2^\kappa(t)=\int_0^t G_1^\kappa(t')E_\kappa(t')dt'\ . \tag{A39}$$

3. The renormalized parameters *above* M_I

We collect here the expressions for the renormalized parameters after solving the RG equations given in A1. They are expressed in terms of the formulas defined in A2 as follows:

$$\tilde\alpha_\rho(t)=\tilde\alpha_\rho(M_C)e_\rho(t),\quad M_\rho(t)=M_\rho(M_C)e_\rho(t) \quad (\rho=3,\ 2,\ 1A,\ \text{and}\ 1B)\ , \tag{A40}$$

$$Y_t(t)=Y_t(M_C)E_t(t)/D_t(t)\ ,\quad Y_\kappa(t)=Y_\kappa(M_C)E_\kappa(t)/D_\kappa(t)\ , \tag{A41}$$

$$A_t(t)=\frac{A}{D_t(t)}+H_2^t(t)-\frac{12Y_t(M_C)H_3^t(t)}{D_t(t)}\ , \tag{A42}$$

$$A_\kappa(t)=\frac{A}{D_\kappa(t)}+H_2^\kappa(t)-\frac{8Y_\kappa(M_C)H_3^\kappa(t)}{D_\kappa(t)}\ , \tag{A43}$$

$$m_Q^2(t)=\tfrac{1}{6}m_1^2(t)+\tfrac{1}{2}m_2^2(t)-\tfrac{1}{3}m_3^2(t)\ , \tag{A44}$$

$$m_{t_R}^2(t)=\tfrac{1}{3}m_1^2(t)-m_2^2(t)+\tfrac{1}{3}m_3^2(t)\ , \tag{A45}$$

$$m_{H_{u1}}^2(t)=\tfrac{1}{2}m_1^2(t)+\tfrac{1}{2}m_2^2(t)\ , \tag{A46}$$

$$m_{H_{d1}}^2(t)=m^2-M^2[\tfrac{3}{2}f_2(t)+\tfrac{1}{2}f_{1A}(t)+\tfrac{1}{3}f_{1B}(t)]\ , \tag{A47}$$

$$m_{N_R}^2(t)=\tfrac{1}{2}m_4^2(t)-\tfrac{1}{2}m_6^2(t)\ , \tag{A48}$$

$$m_{L'}^2(t)=\tfrac{1}{4}m_4^2(t)+\tfrac{1}{2}m_5^2(t)+\tfrac{1}{4}m_6^2(t)\ , \tag{A49}$$

$$m_{H_{u2}}^2(t)=\tfrac{1}{4}m_4^2(t)-\tfrac{1}{2}m_5^2(t)+\tfrac{1}{4}m_6^2(t)\ , \tag{A50}$$

where

$$m_1^2(t)=G_1^t(t)+\frac{1}{D_t(t)}\{3m^2-12Y_t(M_C)[m^2H_6^t(t)+G_2^t(t)]\} -\frac{12Y_t(M_C)m^2}{[D_t(t)]^2}[A^2F_t(t)+2AH_3^t(t) -12Y_t(M_C)H_3^t(t)^2]\ , \tag{A51}$$

$$m_2^2(t)=-m^2+M^2[\tfrac{16}{3}f_3(t)-\tfrac{2}{9}f_{1A}(t)+\tfrac{1}{2}f_{1B}(t)]\ , \tag{A52}$$

$$m_3^2(t)=-3m^2+M^2[\tfrac{40}{3}f_3(t)+3f_2(t)-\tfrac{5}{9}f_{1A}(t) +\tfrac{5}{12}f_{1B}(t)]\ , \tag{A53}$$

$$m_4^2(t)=G_1^\kappa(t)+\frac{1}{D_\kappa(t)}\{3m^2-8Y_\kappa(M_C)[m^2H_6^\kappa(t)+G_2^\kappa(t)]\} -\frac{8Y_\kappa(M_C)m^2}{[D_\kappa(t)]^2}[A^2F_\kappa(t)+2AH_3^\kappa(t) -8Y_\kappa(M_C)H_3^\kappa(t)^2]\ , \tag{A54}$$

$$m_5^2(t)=-\tfrac{5}{12}M^2f_{1B}(t)\ , \tag{A55}$$

$$m_6^2(t)=-m^2+M^2[3f_2(t)-f_{1A}(t)+f_{1B}(t)]\ , \tag{A56}$$

APPENDIX B: EVOLUTION OF PARAMETERS BELOW M_I

1. RG equations for parameters *from* M_I *to* M_W

We collect here the RG equations which are obeyed by the relevant parameters from the intermediate scale M_I to the Weinberg-Salam scale M_W. For the gauge coupling constants, we have

$$\frac{d\tilde\alpha_\rho}{dt'}=2b'_\rho\tilde\alpha_\rho^2\quad(\rho=3,2,Y)\ , \tag{B1}$$

where

$$b'_3=9-2N_f,\quad b'_2=5-2N_f,\quad b'_Y=-1-\tfrac{10}{3}N_f\ , \tag{B2}$$

for N_f families and $t'=\ln(M_I/E)$. The initial condition for $\tilde\alpha_Y$ at M_I (i.e., $t'=0$) is

$$\tilde\alpha_Y(M_I)^{-1}=\tfrac{1}{25}\tilde\alpha_{1A}(M_I)^{-1}+\tfrac{24}{25}\tilde\alpha_{1B}(M_I)^{-1}\ . \tag{B3}$$

As before, the gaugino masses behave similarly as the gauge coupling constants and we have

$$\frac{M_\rho(t')}{\tilde\alpha_\rho(t')}=\text{const}\quad(\rho=3,2,Y) \tag{B4}$$

and the initial condition for M_Y at the scale M_I is given by

$$M_Y(M_I)=\frac{M_{1A}(M_I)\tilde\alpha_{1B}(M_I)+24M_{1B}(M_I)\tilde\alpha_{1A}(M_I)}{24\tilde\alpha_{1A}(M_I)+\tilde\alpha_{1B}(M_I)}\ . \tag{B5}$$

For the Yukawa coupling Y_t and the soft-breaking trilinear scalar coupling A_t, we have

$$\frac{dY_t}{dt'}=-12Y_t^2+Y_t(\tfrac{32}{3}\tilde{\alpha}_3+6\tilde{\alpha}_2+\tfrac{26}{9}\tilde{\alpha}_Y)\ , \tag{B6}$$

$$\frac{dA_t}{dt'}=-12Y_t A_t+\frac{1}{m}(\tfrac{32}{3}\tilde{\alpha}_3M_3+6\tilde{\alpha}_2M_2+\tfrac{26}{9}\tilde{\alpha}_YM_Y)\ . \tag{B7}$$

For the scalar masses, we have

$$\frac{dm_Q^2}{dt'}=-2Y_t(m_Q^2+m_{t_R}^2+\mu_u^2-\mu^2+m^2A_t^2)+\tfrac{32}{3}\tilde{\alpha}_3M_3^2+6\tilde{\alpha}_2M_2^2+\tfrac{2}{9}\tilde{\alpha}_YM_Y^2\ , \tag{B8}$$

$$\frac{dm_{t_R}^2}{dt'}=-4Y_t(m_Q^2+m_{t_R}^2+\mu_u^2-\mu^2+m^2A_t^2)+\tfrac{32}{3}\tilde{\alpha}_3M_3^2+\tfrac{32}{9}\tilde{\alpha}_YM_Y^2\ , \tag{B9}$$

$$\frac{d\mu_u^2}{dt'}=-6Y_t(m_Q^2+m_{t_R}^2+\mu_u^2+m^2A_t^2)+6\tilde{\alpha}_2M_2^2+2\tilde{\alpha}_YM_Y^2+2\mu^2(3\tilde{\alpha}_2+\tilde{\alpha}_Y)\ , \tag{B10}$$

$$\frac{d\mu_d^2}{dt'}=6\tilde{\alpha}_2M_2^2+2\tilde{\alpha}_YM_Y^2+2\mu^2(3\tilde{\alpha}_2+\tilde{\alpha}_Y-3Y_t)\ . \tag{B11}$$

The initial conditions for μ_u^2 and μ_d^2 at the scale M_I are given by

$$\mu_u^2(M_I)=m_{H_{u1}}^2(M_I)+\mu^2(M_I)\ , \tag{B12}$$

$$\mu_d^2(M_I)=m_{H_{d1}}^2(M_I)+\mu^2(M_I)\ . \tag{B13}$$

The mass parameter μ^2 obeys the RG equation

$$\frac{d\mu^2}{dt'}=2\mu^2(3\tilde{\alpha}_2+\tilde{\alpha}_Y-3Y_t) \tag{B14}$$

and for its boundary value at the scale M_I we have chosen $\mu^2(M_I)=10^6$ GeV^2.

2. Formulas for the renormalized parameters *below* M_I

Here we collect some formulas required for the expressions of the renormalized parameters below the scale M_I. They are

$$\tilde{\beta}_\rho=2\tilde{\alpha}_\rho(M_I)b'_\rho\ , \tag{B15}$$

$$\tilde{e}_\rho(t')=1/(1-\tilde{\beta}_\rho t')\ , \tag{B16}$$

$$\tilde{f}_\rho(t')=\frac{1}{b'_\rho}[1-\tilde{e}_\rho(t')^2]\ , \tag{B17}$$

with $\rho=3$, 2, and Y,

$$\tilde{E}_t(t')=\tilde{e}_3(t')^{16/3b'_3}\tilde{e}_2(t')^{3/b'_2}\tilde{e}_Y(t')^{13/9b'_Y}\ , \tag{B18}$$

$$\tilde{F}_t(t')=\int_0^{t'}\tilde{E}_t(t'')dt''\ , \tag{B19}$$

$$\tilde{D}_t(t')=1+12Y_t(M_I)\tilde{F}_t(t')\ , \tag{B20}$$

$$\tilde{H}_2^t(t')=\frac{t'}{m}[\tfrac{32}{3}M_3(M_I)\tilde{\alpha}_3(t')+6M_2(M_I)\tilde{\alpha}_2(t')+\tfrac{26}{9}M_Y(M_I)\tilde{\alpha}_Y(t')]\ , \tag{B21}$$

$$\tilde{H}_3^t(t')=\int_0^{t'}\tilde{H}_2^t(t'')\tilde{E}_t(t'')dt''\ , \tag{B22}$$

$$\tilde{H}_6^t(t')=\int_0^{t'}[\tilde{H}_2^t(t'')]^2\tilde{E}_t(t'')dt''\ , \tag{B23}$$

$$\tilde{G}_1^t(t')=-\tfrac{16}{3}M_3^2(M_I)\tilde{f}_3(t')-3M_2^2(M_I)\tilde{f}_2(t')-\tfrac{13}{9}M_Y^2(M_I)\tilde{f}_Y(t')\ , \tag{B24}$$

$$\tilde{G}_2^t(t')=\int_0^{t'}\tilde{G}_1^t(t'')\tilde{E}_t(t'')dt''\ , \tag{B25}$$

$$q(t')=\tilde{e}_2(t')^{3/2b'_2}\tilde{e}_Y(t')^{1/2b'_Y}/[\tilde{D}_t(t')]^{1/4}\ . \tag{B26}$$

3. The renormalized parameters *below* M_I

We collect here the expressions for the renormalized parameters after solving the RG equations given in B1. They are expressed in terms of the formulas defined in B2 as

$$\tilde{\alpha}_\rho(t')=\tilde{\alpha}_\rho(M_I)\tilde{e}_\rho(t'),\quad M_\rho(t')=M_\rho(M_I)\tilde{e}_\rho(t')\quad (\rho=3,\ 2,\ \text{and}\ Y)\ , \tag{B27}$$

$$Y_t(t')=Y_t(M_I)\tilde{E}_t(t')/\tilde{D}_t(t')\ , \tag{B28}$$

$$A_t(t')=\frac{A_t(M_I)}{\tilde{D}_t(t')}+\tilde{H}_2^t(t')-\frac{12Y_t(M_I)\tilde{H}_3^t(t')}{\tilde{D}_t(t')}\ , \tag{B29}$$

$$m_Q^2(t')=\tfrac{1}{6}m_7^2(t')-\tfrac{1}{4}m_8^2(t')-\tfrac{5}{12}m_9^2(t')\ , \tag{B30}$$

$$m_{t_R}^2(t')=\tfrac{1}{3}m_7^2(t')-\tfrac{1}{2}m_8^2(t')+\tfrac{1}{6}m_9^2(t')\ , \tag{B31}$$

$$\mu_u^2(t')=\tfrac{1}{2}m_7^2(t')+\tfrac{3}{4}m_8^2(t')+\tfrac{1}{4}m_9^2(t')+\mu^2(t')\ , \tag{B32}$$

$$\mu_d^2(t')=m_{H_{d1}}^2(M_I)-\tfrac{3}{2}M_2^2(M_I)\tilde{f}_2(t')-\tfrac{1}{2}M_Y^2(M_I)\tilde{f}_Y(t')+\mu^2(t')\ , \tag{B33}$$

where

$$m_7^2(t')=\frac{1}{\tilde{D}_t(t')}\{m_Q^2(M_I)+m_{t_R}^2(M_I)+m_{H_{u1}}^2(M_I)-12Y_t(M_I)[m^2\tilde{H}_6^t(t')+\tilde{G}_2^t(t')]\}-\frac{12Y_t(M_I)m^2}{\tilde{D}_t(t')^2}[A_t^2(M_I)\tilde{F}(t')+2A_t(M_I)\tilde{H}_3^t(t')-12Y_t(M_I)\tilde{H}_3^t(t')^2]+\tilde{G}_1^t(t') , \tag{B34}$$

$$m_8^2(t')=-m_{t_R}^2(M_I)+\tfrac{2}{3}m_{H_{u1}}^2(M_1)+\tfrac{8}{3}M_3^2(M_I)\tilde{f}_3(t')-M_2^2(M_I)\tilde{f}_2(t')+\tfrac{5}{9}M_Y^2(M_I)\tilde{f}_Y(t') , \tag{B35}$$

$$m_9^2(t')=-2m_Q^2(M_I)+m_{t_R}^2(M_I)+\tfrac{8}{3}M_3^2(M_I)\tilde{f}_3(t')+3M_2^2(M_I)\tilde{f}_2(t')-\tfrac{7}{9}M_Y^2(M_I)\tilde{f}_Y(t') , \tag{B36}$$

$$\mu^2(t')=\mu^2(M_I)q^2(t') , \tag{B37}$$

$$\mu_3^2(t')=q(t')\left\{\mu_3^2(M_I)-2\mu(M_I)t'[3M_2(M_I)\tilde{\alpha}_2(t')+M_Y(M_I)\tilde{\alpha}_Y(t')]+6\mu(M_I)mY_t(M_I)\times\frac{A_t(M_I)\tilde{F}_t(t')+\tilde{H}_3^t(t')}{\tilde{D}_t(t')}\right\} , \tag{B38}$$

with

$$\mu_3^2(M_I)=-Bm\mu(M_I) . \tag{B39}$$

[1]D. J. Gross, J. A. Harvey, E. Martinec, and R. Rohm, Phys. Rev. Lett. **54**, 502 (1985); Nucl. Phys. **B256**, 253 (1985); **B267**, 75 (1987).

[2]P. Candelas, G. T. Horowitz, A. Strominger, and E. Witten, Nucl. Phys. **B258**, 46 (1985).

[3]E. Witten, Nucl. Phys. **B258**, 75 (1985); S. Cecotti, J. P. Derendinger, S. Ferrara, L. Girardello, and M. Roncadelli, Phys. Lett. **156B**, 318 (1985); M. Mangano, Z. Phys. C **28**, 613 (1985); J. Ellis, K. Enqvist, D. V. Nanopoulos, and F. Zwirner, Mod. Phys. Lett. A **1**, 57 (1986); B. Greene, K. H. Kirklin, P. J. Miron, and G. G. Ross, Nucl. Phys. **B278**, 667 (1986); **B292**, 606 (1987); Phys. Lett. B **180**, 69 (1986).

[4]M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi, and N. Seiberg, Nucl. Phys. **B259**, 519 (1985).

[5]J. D. Breit, B. A. Ovrut, and G. C. Sègre, Phys. Lett. **158B**, 33 (1985).

[6]J. P. Derendinger, L. E. Ibáñez, and H. P. Nilles, Phys. Lett. **155B**, 65 (1985); Nucl. Phys. **B267**, 365 (1986).

[7]T. Matsuoka and D. Suematsu, Nucl. Phys. **B274**, 106 (1986); Prog. Theor. Phys. **76**, 886 (1986).

[8]E. Witten, Nucl. Phys. **B268**, 79 (1986).

[9]C. M. Hull, Phys. Lett. B **178**, 357 (1986).

[10]D. Bailin, A. Love, and S. Thomas, Phys. Lett. B **178**, 15 (1986); B. R. Greene, K. H. Kirklin, and P. J. Miron, Nucl. Phys. **B274**, 574 (1986).

[11]R. Holman and D. B. Reiss, Phys. Lett. B **176**, 74 (1986).

[12]D. Bailin, A. Love, and S. Thomas, Phys. Lett. B **176**, 81 (1986).

[13]D. Bailin and A. Love, Phys. Lett. B **181**, 273 (1986).

[14]M. Drees and M. Glück, Phys. Lett. B **180**, 66 (1986).

[15]J. P. Ader and J. C. Wallet, Phys. Lett. B **191**, 375 (1987).

[16]J. Ellis, K. Enqvist, D. V. Nanopoulos, K. Olive, M. Quirós, and F. Zwirner, Phys. Lett. B **176**, 403 (1986), and references therein.

[17]M. Dine, N. Seiberg, X. G. Wen, and E. Witten, Nucl. Phys. **B278**, 769 (1986).

[18]J. Distler, Phys. Lett. B **188**, 431 (1987).

[19]K. Inoue, A. Kakuto, H. Komatsu, and S. Takeshita, Prog. Theor. Phys. **67**, 1889 (1982); **68**, 927 (1982); **71**, 413 (1984).

[20]H. P. Nilles, Phys. Rep. **110**, 1 (1984).

[21]K. Sasaki, Phys. Lett. B **199**, 395 (1987).

[22]J. E. Kim and H. P. Nilles, Phys. Lett. **138B**, 150 (1984).

[23]M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by P. van Nieuwenhuizen and D. Z. Freedman (North-Holland, Amsterdam, 1979); Y. Yanagida, Prog. Theor. Phys. **59**, 66 (1978).

[24]R. Slansky, Phys. Rep. **79**, 1 (1981).

[25]Y. Hosotani, Phys. Lett. **126B**, 309 (1983); **129B**, 193 (1984).

[26]S. Nandi and U. Sarkar, Phys. Rev. Lett. **56**, 564 (1986).

[27]J. Gasser and H. Leutwyler, Phys. Rep. **79**, 77 (1981).

[28]Particle Data group, G. P. Yost *et al.*, Phys. Lett. B **204**, 138 (1988).

[29]L. Alvarez-Gaumé, J. Polchinski, and M. D. Wise, Nucl. Phys. **B211**, 495 (1983).

[30]J. Ellis, D. V. Nanopulos, and K. Tamvakis, Phys. Lett. **121B**, 123 (1983).

[31]L. E. Ibáñez and C. Lopez, Phys. Lett. **126B**, 54 (1983); Nucl. Phys. **B233**, 511 (1984).

[32]L. E. Ibáñez, C. Lopez, and C. Munoz, Nucl. Phys. **B256**, 218 (1985).

[33]C. Kounnas, A. B. Lahanas, D. V. Nanopulos, and M. Quiros, Nucl. Phys. **B236**, 438 (1984).

[34]B. Holdom, Phys. Lett. **166B**, 196 (1986); T. Matsuoka and D. Suematsu, Prog. Theor. Phys. **76**, 901 (1986).

[35]U. Amaldi, A. Böhm, L. S. Durkin, P. Langacker, A. K. Mann, W. J. Marciano, A. Sirlin, and H. H. William, Phys. Rev. D **36**, 1385 (1987).

[36]D. W. Duke and R. G. Roberts, Phys. Rep. **120**, 275 (1985).

[37]R. Barbieri, S. Ferrara, and C. A. Savoy, Phys. Lett. **119B**, 343 (1982); H. P. Nilles, M. Srednicki, and D. Wyler, *ibid.* **120B**, 346 (1983); L. Hall, J. Lykken, and S. Weinberg, Phys. Rev. D **27**, 2359 (1983).

[38]See, for example, H. E. Haber, and G. L. Kane, Phys. Rep. **117**, 75 (1985), and references therein.

[39]JADE Collaboration, W. Bartel *et al.*, Z. Phys. C **31**, 359 (1986).

[40]C. Hearty *et al.*, Phys. Rev. Lett. **58**, 1711 (1987).

[41]UA1 Collaboration, C. Albajar *et al.*, Phys. Lett. B **198**, 261 (1987).

[42]K. Kang and M. Shin, Mod. Phys. Lett. A **1**, 585 (1986).