Vanishing squark and slepton masses in a class of supergravity models

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We study a class of supergravity models where an observable sector is separated from a hidden sector in *the superspace density which defines the supergravity Lagrangian.* It is shown that soft supersymmetry-breaking mass terms for squarks and sleptons vanish at the tree level if the cosmological constant vanishes. Since their masses are induced by radiative corrections mostly due to gauge interactions, the sfermions with the same gauge-quantum numbers are highly degenerate in mass, which is required for the suppression of the flavor-changing neutral currents. By a numerical calculation, we find that the $SU(2)_L \times U(1)_Y$ gauge symmetry can be broken by quantum corrections for suitable values of the free parameters. Some phenomenological features of this model are investigated. In particular, this model predicts right-handed sleptons below 150 GeV.

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I. INTRODUCTION

The striking agreement [1] on the electroweak mixing angle $\sin^2\hat{\theta}_W(m_Z)$ between the prediction of supersymmetric (SUSY) grand unified theories (GUT's) [2] and the experimental result has caused renewed interest in the low-energy SUSY standard model. However, there are two potential problems in the low-energy SUSY model; one is the presence of lower-dimensional operators ($d=4,5$) for baryon-number violation [3], the other is the unsuppressed flavor-changing neutral currents (FCNC's) [4].

Observed suppression of FCNC's suggests an accurate mass degeneracy of squarks and sleptons in different families [4,5]. However, in the general framework of $N=1$ supergravity [6], no symmetry guarantees the mass degeneracy. Indeed if gauge singlets in a hidden sector [7] which are responsible for the SUSY breaking couple to the squarks and sleptons in the Kähler potential, the masses of the squarks and sleptons are not generally degenerate [8,9]. In order to avoid the disastrous mass splitting, one often considers supergravity models where the hidden and observable sectors are separated not only in the superpotential but also in the Kähler potential. Since gravitational interactions do not distinguish the flavors of the matter fields, the soft SUSY-breaking masses are all degenerate at the tree level. The so-called minimal supergravity model [10] belongs to this category and it has been investigated closely.

In this paper we will consider an alternative class of models in which the decoupling between the observable sector and the hidden sector is realized in the *superspace density which defines the supergravity Lagrangian,* instead of the decoupling in the Kähler potential. This class of supergravity models includes an $SU(n,1)/SU(n)\times U(1)$ no-scale model [11] as a special case as explained later. We investigate the structure of soft SUSY-breaking terms [12] in a scalar potential. It turns out that *chirality-conserving* soft SUSY-breaking mass terms $m_i^2\varphi_i^*\varphi_i$ for scalar components φ_i of chiral supermultiplets vanish if and only if the cosmological constant vanishes. At the same time, soft SUSY-breaking trilinear couplings also vanish in this type of supergravity.

We apply this class of supergravity models to the minimal SUSY standard $SU(3)_C\times SU(2)_L\times U(1)_Y$ model. We show that the electroweak gauge symmetry can be spontaneously broken by radiative corrections [13] for suitable values of free parameters. In this radiative breaking scenario, we find indeed a parameter region consistent with the present experimental limits. We will also discuss some characteristics of low-energy predictions of our models. We note that the squarks and sleptons in the first and the second generations acquire their masses mostly through gauge-multiplet loops. Since the gauge interactions are family blind, the particles with the same $SU(3)_C\times SU(2)_L\times U(1)_Y$ quantum numbers are highly degenerate in mass and therefore we have no FCNC problem.

The plan of this paper is the following. In the following section we develop theoretical aspects of the class of supergravity models we propose. In Sec. III we investigate some phenomenological features of this model in the radiative electroweak-breaking scenario. Section IV is devoted to conclusions.

II. VANISHING COSMOLOGICAL CONSTANT AND MASSLESS SQUARKS AND SLEPTONS

We start with a brief review of the scalar potential of supergravity. The relevant terms of the supergravity Lagrangian are given by [14]

$$\int d^4x\, d^4\theta\, E \left[\Phi(S,\bar{S}) + \mathrm{Re}\left(\frac{1}{R} W(S) \right) \right] , \tag{1}$$

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where S is a chiral supermultiplet, E is a superspace determinant, and R is a chiral scalar curvature superfield. Φ is a real function of S and $\bar{S}$ and the superpotential W is a holomorphic function of S. After eliminating auxiliary fields, we obtain the standard form of the scalar potential [15]:

$$V=e^G\hat{V}\ ,\quad \hat{V}=G_A(G^{-1})^A_B G^B-3\ , \tag{2}$$

where we have defined the total Kähler potential as

$$G(x,x^*)=-3\ln\Phi(x,x^*)+\ln|W(x)|^2\ , \tag{3}$$

and $G_A\equiv\partial G/\partial x^A$, etc., and x is the first component of the superfield S. Unless mentioned explicitly, we set $M\equiv m_{\rm Planck}/\sqrt{8\pi}=1$, where $m_{\rm Planck}$ is the Planck mass.

Let us assume that we have two sectors which decouple from each other in the superpotential: one is a hidden sector $\{z^\alpha\}$ and the other is an observable sector $\{y^i\}$. The hidden sector fields are assumed to be gauge singlets and responsible for the SUSY breaking. Throughout this paper we consider the F-term breaking of the SUSY and do not consider the Fayet-Iliopoulos D term. The hidden sector fields are characterized by

$$\langle G_\alpha\rangle=\left\langle\frac{\partial G}{\partial z^\alpha}\right\rangle\lesssim 1\ , \tag{4}$$

where $\langle A\rangle$ denotes a vacuum expectation value of a quantity A. The observable sector includes ordinary quark, lepton, and Higgs supermultiplets. It may include singlets which are usually superheavy [16]. In this observable sector, we assume [17]

$$\langle G_i\rangle=\left\langle\frac{\partial G}{\partial y^i}\right\rangle=0\quad\text{or at least}\ \ll 1\ . \tag{5}$$

When a scalar in the hidden sector couples to squarks and sleptons in the Kähler potential $K\equiv-3\ln\Phi$, the mass splitting of the sfermions occurs in general. A way to avoid this notorious nondegeneracy is therefore to completely separate the hidden sector from the observable sector in the Kähler potential. A typical example for this type of decouplings is minimal supergravity [10], in which the following Kähler potential is taken:

$$K=z^*_\alpha z^\alpha+y^*_i y^i\ , \tag{6}$$

leading to the canonical kinetic terms. In this model, all masses of the squarks and sleptons are equal to a gravitino mass when the cosmological constant vanishes. Low-energy properties of minimal supergravity have been widely investigated in the literature.

In this paper we propose an alternative way of separating the hidden sector from the observable sector; i.e., consider the following separation in *the supergravity Lagrangian* (1):

$$\Phi=I(z,z^*)+J(y,y^*)\ , \tag{7}$$

where I and J are real functions of z and y, respectively. This separation seems quite natural, since the hidden and observable sectors couple with each other only gravitationally at the superspace density level [18]. Notice that the Kähler potential of an $SU(n,1)/SU(n)\times U(1)$ noscale model [11,19] is a special example of our ansatz (7), where I and J are given by

$$I=z+z^*\ ,\quad J=-y^*_i y^i\ . \tag{8}$$

Concerning the superpotential W we take

$$W=h(z)+g(y)\ , \tag{9}$$

as mentioned earlier. Therefore a totality of our assumptions is a separation of the hidden and observable sectors at the superspace density level (1) and Eqs. (4), (5), and (12) given below. Under our ansatz (7) and (9), we calculate a low-energy scalar potential consisting of light particles in the observable sector and study the structure of the soft SUSY-breaking terms in it.

Before proceeding we calculate the effective superpotential which is relevant at low energies, postulating that the gravitino mass lies at the Fermi scale which is much below the Planck scale. In order to do this we consider the bilinear terms of chiral fermions in the supergravity Lagrangian [14]

$$e^{G/2}(G_{ij}+G_iG_j-G_A(G^{-1})^A_B G^B_{ij})\psi^i\psi^j\equiv U_{ij}\psi^i\psi^j\ , \tag{10}$$

where ψ^i is a chiral fermion in the observable sector. The indices A and B run over both the observable and hidden sectors. Define

$$\begin{aligned}\mu_{ij}&\equiv\langle U_{ij}\rangle=m_{3/2}(\langle G_{ij}\rangle-\langle G_\alpha\rangle\langle G^\alpha_{ij}\rangle)\ ,\\ f_{ijk}&\equiv\langle U_{ij,k}\rangle=m_{3/2}\langle G_{ijk}\rangle+O(m_{3/2}/M)\ ,\end{aligned} \tag{11}$$

where $m_{3/2}\equiv\langle e^{G/2}\rangle$ is the gravitino mass and $m_{3/2}\ll M=m_{\rm Planck}/\sqrt{8\pi}$. In deriving the above equations we have normalized the Kähler metric as

$$\begin{aligned}&\langle G^i_j\rangle=\delta^i_j\ ,\quad\langle G^\alpha_\beta\rangle=\delta^\alpha_\beta\ ,\\ &\langle G^i_\alpha\rangle=\langle G^\alpha_i\rangle=0\ ,\end{aligned} \tag{12}$$

and have used Eq. (5). Then at low energies, Eq. (10) reduces to

$$(\mu_{ij}+f_{ijk}y^k)\psi^i\psi^j\ , \tag{13}$$

and therefore we can see that the effective superpotential of the low-energy global SUSY theory is given by

$$\tilde{g}(Y)=\tfrac{1}{2}\mu_{ij}Y^iY^j+\tfrac{1}{6}f_{ijk}Y^iY^jY^k\ , \tag{14}$$

where Y^i is the superfield whose first component is y^i. Notice that the effective superpotential $\tilde{g}$ is generally different from the superpotential g presented in the supergravity Lagrangian [20]. One might think that the coupling constants in Eq. (14) vanish at the supersymmetric limit $m_{3/2}\to 0$. However it is not the case. In order to see this we should note that

$$\begin{aligned}&\langle W\rangle=O(m_{3/2})\ ,\\ &\langle G_\alpha\rangle\ ,\langle K_{ij}\rangle\ ,\langle K^\alpha_{ij}\rangle\ ,\langle K_{ijk}\rangle\lesssim 1\ ,\end{aligned} \tag{15}$$

and

$$\langle G_{ij}\rangle=\langle K_{ij}\rangle+\frac{\langle W_{ij}\rangle}{\langle W\rangle}=\langle K_{ij}\rangle+\frac{\langle g_{ij}\rangle}{\langle W\rangle}\,,$$
$$\langle G_{ij}^{\alpha}\rangle=\langle K_{ij}^{\alpha}\rangle\,, \tag{16}$$
$$\langle G_{ijk}\rangle=\langle K_{ijk}\rangle+\frac{\langle W_{ijk}\rangle}{\langle W\rangle}=\langle K_{ijk}\rangle+\frac{\langle g_{ijk}\rangle}{\langle W\rangle}\,.$$

Then it is easy to see that μ_{ij} and f_{ijk} in Eq. (11) survive even at the SUSY limit $m_{3/2}\to 0$ and in this limit the effective superpotential $\tilde{g}$ coincides with the original g up to an irrelevant overall normalization.

Let us now calculate the scalar potential for the light particles in the observable sector at the flat limit, i.e., taking $M=m_{\rm Planck}/\sqrt{8\pi}\to\infty$ with $m_{3/2}$ fixed. For this purpose we will evaluate the coupling constants of the renormalizable interactions in the scalar potential $\langle V_j^i\rangle\equiv\langle\partial^2V/\partial y_i^*\partial y^j\rangle$, $\langle V_{ij}\rangle\equiv\langle\partial^2V/\partial y^i\partial y^j\rangle$, etc.

We first calculate $\langle V_j^i\rangle$, which corresponds to *chirality-conserving* scalar mass terms. Differentiating Eq. (2) twice we obtain

$$\langle V_j^i\rangle=\langle V\rangle\langle G_j^i\rangle+\langle e^G\rangle\langle\hat{V}_j^i\rangle=\langle V\rangle\delta_j^i+m_{3/2}^2\langle\hat{V}_j^i\rangle\,. \tag{17}$$

A straightforward calculation shows

$$\langle\hat{V}_j^i\rangle=\delta_j^i+\langle G_{jk}\rangle\langle(G^{-1})_\alpha^{k,i}\rangle\langle G^\alpha\rangle+\langle G_\alpha\rangle\langle[(G^{-1})_\beta^\alpha]_j^i\rangle\langle G^\beta\rangle+\langle G_\alpha\rangle\langle(G^{-1})_j^{\alpha,i}\rangle+\langle G^{ik}\rangle\langle G_{kj}\rangle+\langle G_\alpha\rangle\langle(G^{-1})_{k,j}^\alpha\rangle\langle G^{ki}\rangle+\langle G_\alpha\rangle\langle G_j^{\alpha i}\rangle\,. \tag{18}$$

Using

$$\langle(G^{-1})_{A,C}^B\rangle=-\langle G_{AC}^B\rangle\,,\quad\langle[(G^{-1})_A^B]_C^D\rangle=\langle G_{AC}^E\rangle\langle G_E^{BD}\rangle+\langle G_A^{DE}\rangle\langle G_{EC}^B\rangle-\langle G_{AC}^{BD}\rangle\,, \tag{19}$$

the above equation reads

$$\langle\hat{V}_j^i\rangle=\delta_j^i+\langle G_\alpha\rangle\{\langle G_k^{\alpha i}\rangle\langle G_{\beta j}^k\rangle-\langle G_{\beta j}^{\alpha i}\rangle\}\langle G^\beta\rangle+(\langle G^{ik}\rangle-\langle G^\alpha\rangle\langle G_\alpha^{ik}\rangle)(\langle G_{kj}\rangle-\langle G_\beta\rangle\langle G_{kj}^\beta\rangle)\,. \tag{20}$$

Therefore, noticing Eq. (11), we find

$$\langle V_j^i\rangle=(\langle V\rangle+m_{3/2}^2)\delta_j^i+m_{3/2}^2\langle G_\alpha\rangle\{\langle G_k^{\alpha i}\rangle\langle G_{\beta j}^k\rangle-\langle G_{\beta j}^{\alpha i}\rangle\}\langle G^\beta\rangle+\mu^{ik}\mu_{kj}\,. \tag{21}$$

Here μ^{ik} is the complex conjugate of μ_{ik}. In our ansatz (7) one can see that

$$\langle G_k^{\alpha i}\rangle\langle G_{\beta j}^k\rangle-\langle G_{\beta j}^{\alpha i}\rangle=-\tfrac{1}{3}\delta_\beta^\alpha\delta_j^i\,. \tag{22}$$

Substituting this into Eq. (21) we obtain

$$\langle V_j^i\rangle=\tfrac{2}{3}\langle V\rangle\delta_j^i+\mu^{ik}\mu_{kj}\,. \tag{23}$$

The second term of the right-hand side (RHS) is a SUSY-invariant mass term which is zero for squarks and sleptons. The soft SUSY-breaking term is proportional to the cosmological constant $\langle V\rangle$ and therefore $\langle V\rangle=0$ results in the vanishing soft SUSY-breaking mass terms for the squarks and sleptons. From now on, we consider the case of the vanishing cosmological constant $\langle V\rangle=0$.

Next consider $\langle V_{ij}\rangle$. In the minimal SUSY standard model, this term is allowed only for a mixing term of two Higgs scalars. A similar calculation gives

$$\langle V_{ij}\rangle=m_{3/2}^2\langle\hat{V}_{ij}\rangle=m_{3/2}^2(2\langle G_{ij}\rangle-\langle G_\alpha\rangle\langle G_{ij}^\alpha\rangle+\langle G^\alpha\rangle\langle G_{ij\alpha}\rangle-\langle G_{ik}\rangle\langle G_{j\alpha}^k\rangle\langle G^\alpha\rangle-\langle G_{jk}\rangle\langle G_{i\alpha}^k\rangle\langle G^\alpha\rangle+\langle G_\alpha\rangle\langle(G^{-1})_{\beta,ij}^\alpha\rangle\langle G^\beta\rangle)\,, \tag{24}$$

which are generally nonzero and expected to be $O(m_{3/2}^2)$.

A SUSY-breaking trilinear coupling $\langle V_{ijk}\rangle$ is computed as

$$\langle V_{ijk}\rangle=m_{3/2}^2\langle\hat{V}_{ijk}\rangle=m_{3/2}^2[3\langle G_{ijk}\rangle+\langle G_{ijk\alpha}\rangle\langle G^\alpha\rangle-(\langle G_{ijl}\rangle\langle G_{\alpha k}^l\rangle\langle G^\alpha\rangle+\text{permutations})]\,. \tag{25}$$

Note that $\langle G_{ijk}\rangle$, $\langle G_{ijk\alpha}\rangle=O(m_{3/2}^{-1})$. Our assumptions (7) and (9) lead to

$$\langle G_{ijk\alpha}\rangle=-\frac{\langle g_{ijk}\rangle}{\langle W\rangle}\frac{\langle h_\alpha\rangle}{\langle W\rangle}+O(1)=-\langle G_{ijk}\rangle\frac{\langle h_\alpha\rangle}{\langle W\rangle}+O(1)\,,$$
$$\langle G_{k\alpha}^l\rangle=-\frac{\langle I_\alpha\rangle}{\langle\phi\rangle}\delta_k^l\,. \tag{26}$$

Then

$$\langle V_{ijk}\rangle=m_{3/2}^2\langle G_{ijk}\rangle\{3-\langle G^\alpha\rangle\langle G_\alpha\rangle\}+O(m_{3/2}^2/M)=-\langle G_{ijk}\rangle\langle V\rangle+O(m_{3/2}^2/M)=O(m_{3/2}^2/M)\,, \tag{27}$$

where we have recovered M from the dimensional analysis. Therefore the trilinear coupling also vanishes at the flat limit $M \to \infty$ with the vanishing cosmological constant.

Other terms are calculated as

$$\langle V_{ij}^{k} \rangle = f_{ijl}\mu^{lk} , \tag{28}$$

$$\langle V_{ij}^{kl} \rangle = f_{ijm} f^{mkl} , \tag{29}$$

at the flat limit. μ^{ij} and f^{ijk} are the complex conjugates of μ_{ij} and f_{ijk}, respectively. One can easily check that the hard SUSY-breaking terms $\langle V_{ijkl} \rangle$ and $\langle V_{jkl}^{i} \rangle$ vanish at the flat limit.

To summarize, we obtain the scalar potential of light particles in the observable sector as

$$V = \sum_i |\tilde{g}_i|^2 + \tfrac{1}{2} D^a D^a + \tfrac{1}{2}\langle V_{ij} \rangle y^i y^j + \tfrac{1}{2}\langle V^{ij} \rangle y_i^* y_j^* , \tag{30}$$

where we have added the D terms D^a which we had not taken into account [21]. The first line of the RHS of the above equation is supersymmetric, while the terms in the second line are the soft SUSY-breaking terms.

Concerning gauge fermion masses, we will assume a universal gauge fermion mass (which should be justified for a GUT). The masses of squarks and sleptons are mainly induced by the gauge interactions and therefore they are degenerate in the different families, which will result in the sufficient suppression of the FCNC's.

Finally we should comment on a crucial difference between our model and the no-scale model. In the no-scale model [22], the scalar potential identically vanishes, since the constant superpotential in the hidden sector is assumed. The SUSY-breaking scale and consequently the gravitino mass and the other SUSY-breaking terms are not determined at the tree level but determined through loop corrections. On the other hand, in our model we assume a nontrivial superpotential in which the SUSY-breaking terms are determined at the tree level. Furthermore we introduce a SUSY-invariant Higgs-boson mass, which admits a relatively heavy top quark (see Sec. III).

III. PREDICTIONS AT THE FERMI SCALE

In this section we will investigate phenomenological features of our model. We begin with a brief review of the Higgs sector in the minimal SUSY standard $SU(3)_C \times SU(2)_L \times U(1)_Y$ model [23,24]. In this model, there are two chiral superfields H_1 and H_2, which transform, respectively, as $H_1 = (1,2,-\frac{1}{2})$ and $H_2 = (1,2,\frac{1}{2})$ under the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge group. The Higgs potential is given by

$$V = \frac{g^2}{8}(A_1^\dagger \tau_a A_1 + A_2^\dagger \tau_a A_2)^2 + \frac{g'^2}{8}(A_1^\dagger A_1 - A_2^\dagger A_2)^2 + m_1^2 A_1^\dagger A_1 + m_2^2 A_2^\dagger A_2 - m_3^2 (A_1 A_2 + A_1^* A_2^*) , \tag{31}$$

where $A_1(A_2)$ is the first component of H_1 (H_2), $g[g']$ is the gauge coupling constant of $SU(2)_L$ [$U(1)_Y$], and τ_a represents a Pauli matrix. m_1^2 and m_2^2 in Eq. (31) are written as

$$m_1^2 = m_s^2 + \Delta_1^2 , \quad m_2^2 = m_s^2 + \Delta_2^2 . \tag{32}$$

Here m_s is a SUSY-invariant mass term of Higgs fields and $\Delta_1^2(\Delta_2^2)$ is a soft SUSY-breaking mass term of A_1 (A_2).

The stability of the potential requires

$$m_1^2 + m_2^2 > 2|m_3^2| . \tag{33}$$

The $SU(2)_L \times U(1)_Y$ gauge symmetry breaks down to the $U(1)_{em}$ symmetry if the condition

$$m_3^4 > m_1^2 m_2^2 \tag{34}$$

is satisfied. We define an angle θ by a ratio of the vacuum expectation values of the two Higgs scalars

$$\tan\theta = \langle A_2 \rangle / \langle A_1 \rangle . \tag{35}$$

Then θ satisfies

$$\sin 2\theta = 2m_3^2/(m_1^2 + m_2^2) , \tag{36}$$

and the W-boson mass is given by

$$M_W^2 = \frac{g^2}{g^2 + g'^2}\left[\frac{m_2^2 - m_1^2}{\cos 2\theta} - m_1^2 - m_2^2\right] \tag{37}$$

at the tree level.

Let us now apply our supergravity model given in the previous section to the minimal SUSY standard model with three generations and the two Higgs doublets. We have $m_1^2 = m_2^2 = m_s^2$ at the renormalization point $\mu = M_X$ (we take as M_X the GUT scale which is of order 10^{16} GeV in this case) and the $SU(2)_L \times U(1)_Y$ gauge symmetry is not broken at the tree level because Eq. (34) cannot be satisfied while keeping Eq. (33). Therefore the gauge symmetry is expected to be broken by radiative corrections [13]. Here we will study this possibility.

The soft SUSY-breaking parameters and other coupling constants follow their renormalization-group equations (RGE's) [25] and their values at $\mu \sim M_W$ are obtained by solving the RGE's. In our numerical analysis we neglect all Yukawa couplings but for the top quark. As input parameters we use [1]

$$\alpha_{em}^{-1} = 127.9, \quad \sin^2\hat{\theta}_W(m_Z) = 0.2334 . \tag{38}$$

Let us summarize the boundary conditions on the parameters at the scale M_X, which we use when solving the RGE's (we follow the notation in Ref. [26]).

(i) The gauge fermion masses:

$$M_{1X} = M_{2X} = M_{3X} \equiv M_{gX} . \tag{39}$$

(ii) The supersymmetric Higgs-boson mass:

$$m_s(M_X) = m_{sX} . \tag{40}$$

(iii) The soft SUSY-breaking masses for scalars:

$$\Delta_{1X}^2 = \Delta_{2X}^2 = m^2(\tilde{q}_r)_X = m^2(\tilde{u}_r)_X = m^2(\tilde{d}_r)_X = m^2(\tilde{l}_r)_X = m^2(\tilde{e}_r)_X = 0 , \tag{41}$$

where r ($r=1,2,3$) denotes the generation.

(iv) The soft SUSY-breaking trilinear top-squark–top-squark–Higgs-boson coupling:

$$m_{10X}=0 \ . \tag{42}$$

(v) The Higgs-boson mixing term:

$$m_{sX}\Delta_{3X}=m_{3X}^2 \ . \tag{43}$$

Here $\Delta_{3X}=\gamma m_{3/2}$ where the coefficient γ depends on the detail of the hidden sector.

(vi) The Yukawa coupling for the top quark:

$$h_t(M_X)=h_X \ . \tag{44}$$

As was shown in the previous section, all chirality-conserving SUSY-breaking mass parameters and the trilinear scalar coupling vanish at $\mu \sim M_X$.

On the above parameters we must impose the stability condition of the Higgs potential (33) at the GUT scale. In our model, it reads

$$m_{sX}^2 > |m_{3X}^2| \ . \tag{45}$$

We have four parameters M_{gX}, m_{sX}, m_{3X}^2, and h_X. However, one of the combinations should be fixed so that the correct Fermi scale is produced. In practice, it is convenient to determine the Yukawa coupling from Eq. (37). Therefore we have three free parameters M_{gX}, m_{sX}, and m_{3X}^2. Notice again that the SUSY is softly broken by only two mass parameters M_{gX} and m_{3X}^2.

It is instructive to mention some of the properties of solutions of the RGE's. The gauge fermion masses $M_i (i=1,2,3)$ at $\mu=M_W$ are easily obtained as

$$\begin{aligned} M_1 &= \tfrac{5}{3}M_{gX}g'^2/g_X^2 \ , \\ M_2 &= M_{gX}g^2/g_X^2 \ , \\ M_3 &= M_{gX}g_c^2/g_X^2 \ , \end{aligned} \tag{46}$$

where g_X is the unified gauge coupling constant at $\mu=M_X$. The soft SUSY-breaking mass parameters for the sleptons are given by

$$\begin{aligned} m^2(\tilde{l}_r) &= M_{gX}^2[3(1-g^4/g_X^4)/2 \\ &\qquad +(1-25g'^4/9g_X^4)/22] \ , \\ m^2(\tilde{e}_r) &= M_{gX}^2 2(1-25g'^4/9g_X^4)/11 \quad (r=1,2,3) \end{aligned} \tag{47}$$

and those for the squarks are

$$\begin{aligned} m^2(\tilde{q}_r) &= M_{gX}^2[-8(1-g_c^4/g_X^4)/9+3(1-g^4/g_X^4)/2+(1-25g'^4/9g_X^4)/198] \ , \\ m^2(\tilde{u}_r) &= M_{gX}^2[-8(1-g_c^4/g_X^4)/9+8(1-25g'^4/9g_X^4)/99] \ , \\ m^2(\tilde{d}_r) &= M_{gX}^2[-8(1-g_c^4/g_X^4)/9+2(1-25g'^4/9g_X^4)/99] \quad (r=1,2) \ . \end{aligned} \tag{48}$$

The masses for the top squarks are smaller than those for the squarks in the other generations because of the top-quark-loop contributions. Numerically we obtain the relations

$$\begin{aligned} &M_2/M_1=2.01, \quad M_3/M_1=7.19 \ , \\ &m(\tilde{l}_r)/M_1=1.78 \ , \quad m(\tilde{e}_r)/M_1=0.95 \ , \\ &m(\tilde{q}_r)/M_1=6.60 \ , \\ &m(\tilde{u}_r)/M_1=6.40 \ , \quad m(\tilde{d}_r)/M_1=6.38 \ . \end{aligned} \tag{49}$$

The physical masses $\overline{m}$ for the sleptons and squarks are then given by

$$\begin{aligned} \overline{m}^2(\tilde{l}_r)_1 &= m^2(\tilde{l}_r)+M_W^2\cos2\theta(\tan^2\theta_W+1)/2 \ , \\ \overline{m}^2(\tilde{l}_r)_2 &= m^2(\tilde{l}_r)+M_W^2\cos2\theta(\tan^2\theta_W-1)/2 \ , \\ \overline{m}^2(\tilde{e}_r) &= m^2(\tilde{e}_r)-M_W^2\cos2\theta\tan^2\theta_W \ , \\ \overline{m}^2(\tilde{q}_r)_1 &= m^2(\tilde{q}_r)-M_W^2\cos2\theta(\tan^2\theta_W/6-\tfrac{1}{2}) \ , \\ \overline{m}^2(\tilde{q}_r)_2 &= m^2(\tilde{q}_r)-M_W^2\cos2\theta(\tan^2\theta_W/6+\tfrac{1}{2}) \ , \\ \overline{m}^2(\tilde{u}_r) &= m^2(\tilde{u}_r)+2M_W^2\cos2\theta\tan^2\theta_W/3 \ , \\ \overline{m}^2(\tilde{d}_r) &= m^2(\tilde{d}_r)-M_W^2\cos2\theta\tan^2\theta_W/3 \ . \end{aligned} \tag{50}$$

To the solutions of the RGE's, we impose the following cuts.

A. Experimental cuts

We require that the obtained mass spectrum is consistent with the following experimental limits: top-quark mass: $m_t > 89$ (GeV) [27]; gluino mass: $m_{\tilde{g}} > 150$ (GeV) [28]; squark masses: $m_{\tilde{q}} > 170$ (GeV) [28,29]; chargino masses: $m_{\chi^{\pm}} > 45$ (GeV) [30]; selectron mass: $m_{\tilde{e}} > 43$ (GeV) [31]; contribution of sneutrino to the Z width [32],

$$\frac{\Gamma(Z\to\tilde{\nu}\tilde{\nu})}{\Gamma(Z\to\nu\nu)} < 0.18 \ ; \tag{51}$$

contribution of neutralinos to the Z width [32],

$$\frac{\Gamma(Z\to\tilde{N}\tilde{N}')}{\Gamma(Z\to\nu\nu)} < 0.18 \ , \tag{52}$$

where $\tilde{N}\tilde{N}'$ represents all pairs of neutralinos relevant to the Z decays.

Note that we do not impose any condition on the neutral-Higgs-boson masses, because it has recently been recognized that the mass of the lightest Higgs boson may receive large radiative corrections due to top and top-squark loops [33] and becomes easily above the experimental lower bound.

B. An astrophysical cut

Since the lightest superparticle (LSP) is stable in the minimal SUSY standard model, we require that it should be $SU(3)_C \times U(1)_{em}$ neutral.

C. A theoretical cut

When solving the RGE's, we neglect the Yukawa coupling h_b for the bottom quark. The ratio of the Yukawa couplings and that of the quark masses are related as

$$\frac{h_b}{h_t}=\frac{m_b}{m_t}\tan\theta \; . \tag{53}$$

Thus when $\tan\theta$ is large, our approximation of neglecting the bottom-quark Yukawa coupling is invalid. Therefore we impose that $(h_b/h_t)^2<0.1$ at $\mu=M_W$ as a theoretical cut.

Here we should note that a similar analysis was done in Ref. [34], where only the gaugino mass is assumed as a SUSY-breaking term, which corresponds to the $m_{3X}=0$ case in our models. They used experimental constraints on the masses of superparticles, etc., and searched the allowed region of the parameter space. However, recent accelerator experiments give stricter constraints than those imposed in [34]. Therefore the lower bounds on the masses of the superparticles, etc., predicted by our analysis are more severe than those given by [34].

We have searched solutions for the radiative breaking scenario in the parameter range of $|M_{gX}|$, $|m_{sX}|$, $|m_{3X}|\leq 10$ TeV. In our convention we take $M_{gX}>0, m_{3X}^2>0$. On the other hand, m_{sX} can be posi-

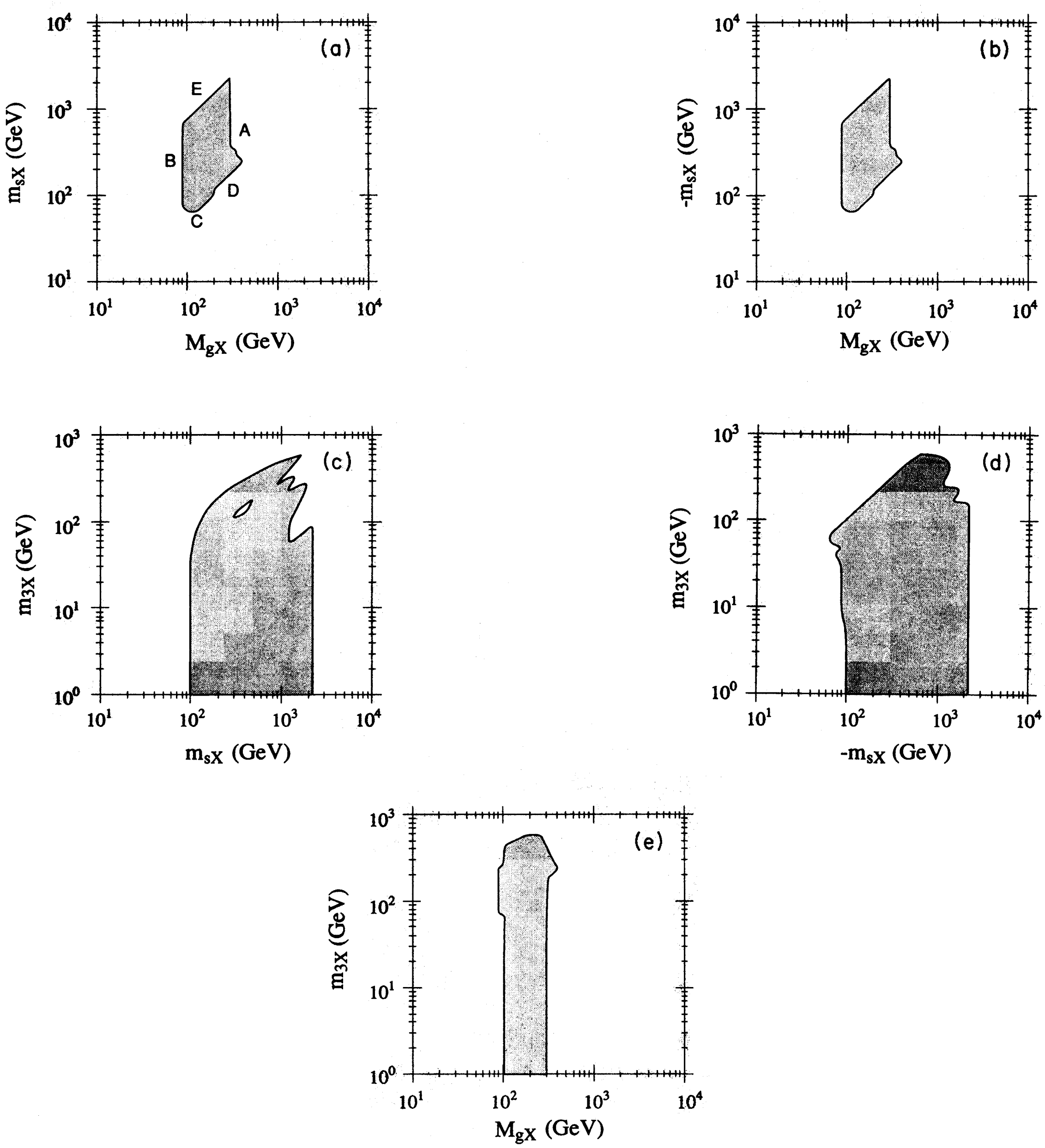

FIG. 1. Allowed regions for the parameters at M_{gX}. In our convention we take $M_{gX}, M_{3X}>0$. (a) and (b) show the (M_{gX}, m_{sX}) plane for positive and negative m_{sX}, respectively. (c) and (d) show the (m_{sX}, m_{3X}) plane for positive and negative m_{sX}. (e) is the (M_{gX}, m_{3X}) plane.

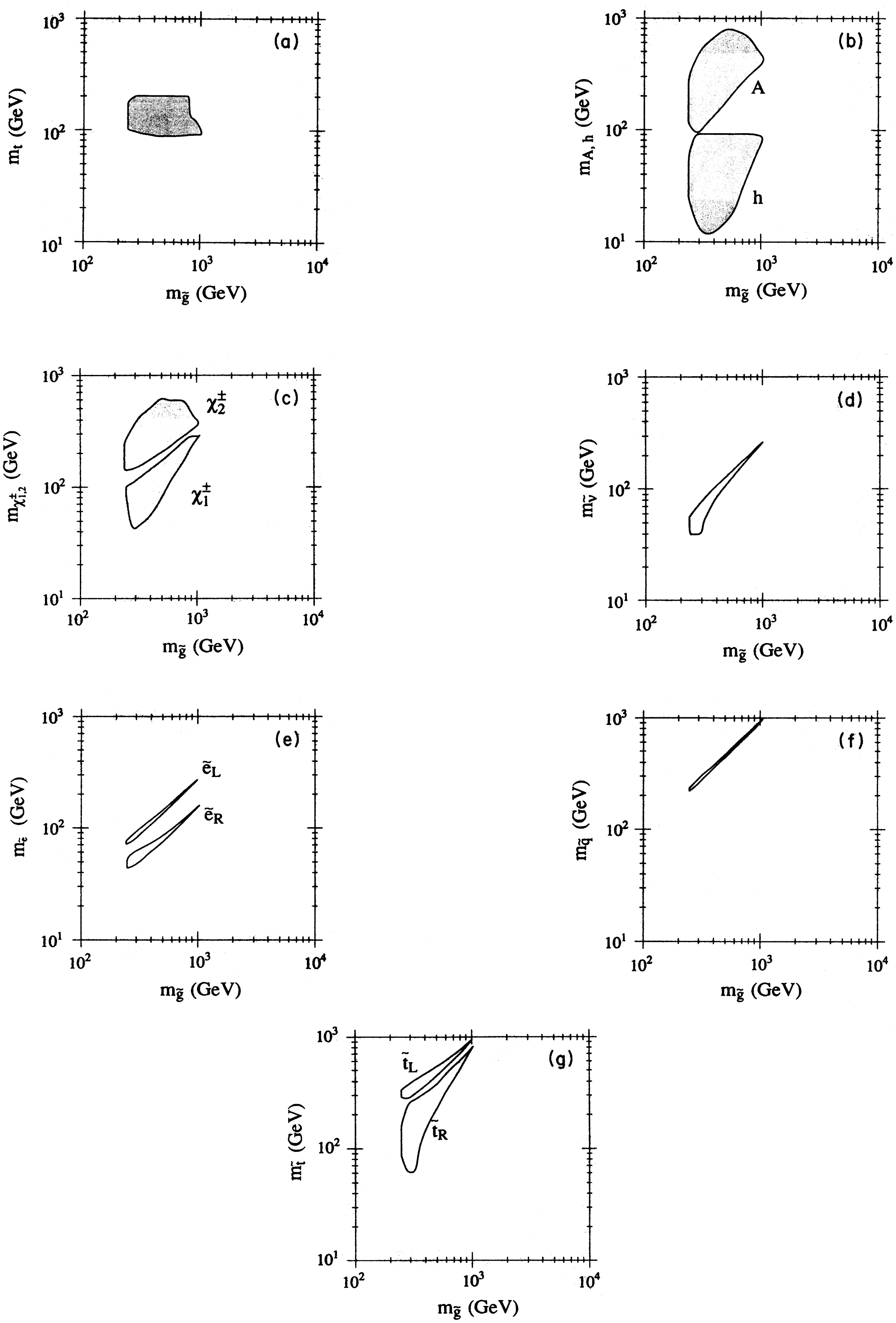

FIG. 2. Low-energy predictions of our model. Masses of various particles are plotted versus the gluino mass. (a) indicates the top-quark mass, (b) the mass of the lightest Higgs boson h and that of the pseudoscalar Higgs boson A, (c) two chargino masses, (d) the sneutrino mass, (e) the mass of the right-handed sleptons $\tilde{e}_R$, and that of the left-handed sleptons $\tilde{e}_L$, (f) the masses of the squarks in the first and the second generations, and (g) the top-squark masses.

tive or negative. We find that there exist solutions which survive the constraints **A**–**C**. In Figs. 1(a)–1(e), we show the allowed region of the parameters at M_X. We observe that

$$80 \lesssim |M_{gX}| \lesssim 400 \text{ GeV} ,$$
$$60 \lesssim |m_{sX}| \lesssim 2000 \text{ GeV} , \qquad (54)$$
$$0 \lesssim |m_{3X}| \lesssim 500 \text{ GeV} .$$

An important feature is that the allowed region of the universal gauge fermion mass at $\mu = M_X$ is tightly restricted from above, which results in relatively light gaugino and slepton masses.

It is worthwhile to observe what determines the boundaries in the figures. Let us consider, for example, Fig. 1(a). We can see from Eq. (49) that a gaugino mass parameter M_1 is slightly larger than the right-handed scalar electron mass parameter $m_{\tilde{e}_R}$. The lightest neutralino consists of a mixture of the neutralinos in our case and it becomes lighter than $\tilde{e}_R$ due to the mixing effect. However when the gaugino mass is large, the mixing effect reduces and hence the lightest neutralino becomes heavier than $\tilde{e}_R$. Indeed this is the case in the region A of Fig. 1(a). This means that the charged particle $\tilde{e}_R$ becomes the LSP and therefore this region is astrophysically excluded. The region B is excluded because the sneutrino is too light to satisfy $\Gamma_{\tilde{\nu}\tilde{\nu}}/\Gamma_{\nu\nu} < 0.18$. In the region C, the chargino becomes the LSP. The region D invades the lower limit of the top-quark mass 89 GeV. Finally in the region E, we cannot find a solution of the radiative $SU(2)_L \times U(1)_Y$ breaking. Remark that the experimental limits on the masses of the gluino and the squarks have not been used to restrict the parameter region. Namely, in the allowed region, the experimental constraints on their masses are automatically satisfied. Similar observations can be done for other figures but will not be stated here.

Figure 2 shows low-energy predictions of this model. The main characteristics of our model are the following.

(a) The gluino mass is predicted as 250 GeV $\lesssim m_{\tilde{g}} \lesssim 1$ TeV. The lower and the upper bounds of it are given as explained above. The upper bound of the gluino mass $m_{\tilde{g}} < 400$ GeV was also derived by Ellis and Zwirner [34]. This upper bound comes from the imposition that the neutralino LSP solves the dark-matter problem; i.e., the contribution to the cosmic density from the LSP is equal to the critical one. However, this imposition also predicts a light top quark with $m_t < 90$ GeV, which is very unlikely at present [27].

(b) Our model has little predictive power on the top-quark mass [Fig. 2(a)] [35]. The upper bound merely reflects that the Yukawa coupling blows up before the GUT scale if the top-quark mass violates the bound.

(c) In Fig. 2(b), the lightest Higgs-boson mass m_h is 10 GeV $\lesssim m_h \lesssim 90$ GeV at the tree level. As mentioned in the constraint **A**, we have not imposed an experimental lower bound on the mass, because radiative corrections easily raise the mass above the experimental limit [33]. The mass of the pseudoscalar Higgs boson m_A lies in the wide range of 90–800 GeV. The mass of the other neutral scalar boson is found to be almost degenerate with m_A. The charged-Higgs-boson mass $m_{H^\pm}$ is given as $m_{H^\pm}^2 = m_A^2 + m_{W^\pm}^2$.

(d) There are two charginos $\chi_1^\pm$ and $\chi_2^\pm$. The mass of the lighter one is below 300 GeV and that of the heavier one is below 600 GeV [Fig. 2(c)].

(e) Figures 2(d) and 2(e) indicate that the masses of the sleptons almost linearly depend on the gluino mass. This is because the soft SUSY-breaking slepton mass parameters are given as a function of the universal gaugino mass M_{gX} as can be seen in Eqs. (47) and (49). We find that they are relatively light and, in particular, the right-handed selection which is the lightest slepton in our model weighs less than 150 GeV.

(f) Because of the QCD corrections, the squarks are generally heavier than the sleptons. Their masses (except for the top-squark mass) also linearly depend on the gluino mass and 200 GeV $\lesssim m_{\tilde{q}} \lesssim 1$ TeV [Fig. 2(f)]. The top-squark masses are widely distributed [Fig. 2(g)] since the left-right mixing term gives an extra contribution.

(g) The mass of the LSP is predicted as 2 GeV $\lesssim m_{\text{LSP}} \lesssim 150$ GeV. The lighter LSP survives the constraint from the Z decay because it is dominated by the photino. In many regions of the parameter space, the LSP and the right-handed scalar electron are almost degenerate in mass. We find that the LSP heavier than 40 GeV are occupied by a b-ino. For the b-ino LSP, we calculated the cosmic density Ω and found that $0.01 h_0^{-2} \lesssim \Omega \lesssim 0.2 h_0^{-2}$, where h_0 is the present-day Hubble constant in the unit of 100 km^{-1}s^{-1}Mpc. Photino-dominant LSP gives a slightly large cosmic density. With the astrophysical observations $0.5 \lesssim h_0 \lesssim 1$, $\Omega \lesssim 1$ is automatically satisfied. In Ref. [34], it has been pointed out

TABLE I. Predictions of the masses of particles for the right-handed selectron mass $m_{\tilde{e}_R} = 80$ GeV. As explained in the text we should not take m_h at a face value, since it receives a large radiative correction [33].

m_t	90–190 GeV	$m_{\tilde{\nu}}$	110–140 GeV
m_h	20–90 GeV	$m_{\tilde{e}_L}$	130–15 GeV
m_A	170–720 GeV	$m_{\tilde{q}}$	410–490 GeV
$m_{H^\pm}$	190–720 GeV	$m_{\tilde{t}_L}$	450–550 GeV
$m_{\tilde{g}}$	470–540 GeV	$m_{\tilde{t}_R}$	270–340 GeV
$m_{\chi^\pm}$	70–170 GeV	m_{LSP}	60–80 GeV

that if the top-quark mass is above 90 GeV the density Ω is always smaller than the critical density $\Omega_c = 1$ and therefore our result is consistent with theirs.

The masses of the squarks in the first and second generations are highly degenerate. On the other hand, the masses of the top squarks are different from those of the squarks in the other generations. We checked, however, that because the relevant elements of the Kobayashi-Maskawa matrix are very small (for a review, see [36]), the contribution of the top squark to the K-$\bar{K}$ mixing through the W-ino exchange is small enough and is consistent with the experimental value.

The right-handed selectron whose mass is roughly below 90 GeV can be searched at the CERN e^+e^- collider LEP II. In Table I we show the predictions of our model for $m_{\tilde{e}_R} = 80$ GeV. We can see for a given $m_{\tilde{e}_R}$, other masses of the superpartners can be well determined in our model.

IV. CONCLUSIONS

In this paper we have considered a supergravity model in which the observable sector and the hidden sector are separated in *the superspace density*. This separation guarantees the mass degeneracy of the squarks and sleptons, as required for the suppression of FCNC's. This model has a special form of the soft SUSY-breaking terms. That is, at the tree level, mass terms of the squarks and sleptons and the cosmological constant simultaneously vanish. At the same time, soft SUSY-breaking trilinear couplings become zero.

We have performed a phenomenological analysis on this model. Namely, we have combined this supergravity with the minimal SUSY standard model and have found that the $SU(2)_L \times U(1)_Y$ gauge symmetry can be radiatively broken for suitable values of the free parameters. An interesting feature of our models is that LSP (which is mostly a neutralino) is nearly degenerate with the right-handed scalar electron in the wide region of the parameter space. When the gaugino mass M_{gX} at M_X is large, the neutralino becomes heavier than the right-handed selectron, which is forbidden by an astrophysical reason since the charged particle becomes the LSP in this case. Thus our models have an upper bound of the universal gaugino mass $M_{gX} \lesssim 400$ GeV and consequently an upper bound of the masses of the gluino, squarks, and sleptons. In particular the right-handed scalar electron which is the lightest slepton in our case weighs less than 150 GeV. The selectron with $m_{\tilde{e}_R} \lesssim 90$ GeV can be explored by the LEP II, although the whole range can be covered by e^+e^- collider experiments in the next generation. On the masses of the squarks and the gluino, our models predict 200 GeV $\lesssim m_{\tilde{q},\tilde{g}} \lesssim 1$ TeV. The masses up to, say, 300 GeV or so will be covered by the Collider Detector at Fermilab (CDF). Heavier squarks and gluino will be searched by new hadronic accelerators.

There is an alternative possibility [37] of realizing massless scalar bosons at the fundamental scale. That is, in a nonlinear σ model coupled to supergravity [38], the Nambu-Goldstone bosons remain massless, if the total Kähler potential is invariant under relevant transformations [39]. In this case, soft SUSY-breaking trilinear couplings do not necessarily vanish. However, we may expect that the Nambu-Goldstone hypothesis will share main features such as the presence of light sleptons in the present model.

We would like to mention that the Higgs potential may receive large effects associated with the top-squark mass threshold [40,33] when the Yukawa coupling h_t is considerably large. Then the tree-level relation (37) must be modified by taking account of the threshold effects. This will shift the estimate of the top-quark Yukawa coupling and consequently the top-quark mass in our numerical calculation. However, the change of the top-quark mass does not affect much the bound of the gaugino mass. Since the gaugino mass is the most responsible for yielding the main results in our models, we expect that the threshold effects will not change our present results drastically except for the lightest Higgs-boson mass. These threshold effects, however, deserve further study and will be discussed elsewhere.

Finally we should stress that the presence of the light sleptons below 150 GeV is common in a large class of models including the no-scale model [11], the Nambu-Goldstone hypothesis, and the minimal SUSY-breaking model with gaugino mass [24,34]. If the sleptons are not found below 150 GeV in future experiments, we believe that all these models as well as ours will be confronted with a serious difficulty.

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[18] We do not know, however, what mechanism will ensure this separation between the two sectors.

[19] There is a different type of no-scale model in which only the SU(1,1) symmetry is imposed. There the Kähler potential is assumed as $K=-3\ \ln(z+z^*)+y_i^*y^i$ which leads to nonvanishing SUSY-breaking masses for the scalar bosons in the observable sector. Therefore the phenomenology is completely different from ours.

[20] The μ_{ij} includes a SUSY-invariant mass term of two Higgs doublets in the minimal SUSY standard model. It is interesting to observe that even if the mass term vanishes in the superpotential g in the supergravity Lagrangian (1), it will emerge in the low-energy superpotential $\tilde{g}$ through the Kähler potential. Remarkably the magnitude of the SUSY-invariant mass term becomes of order $m_{3/2}$ without any fine-tuning. See G. F. Giudice and A. Masiero, Phys. Lett. B **206**, 480 (1988).

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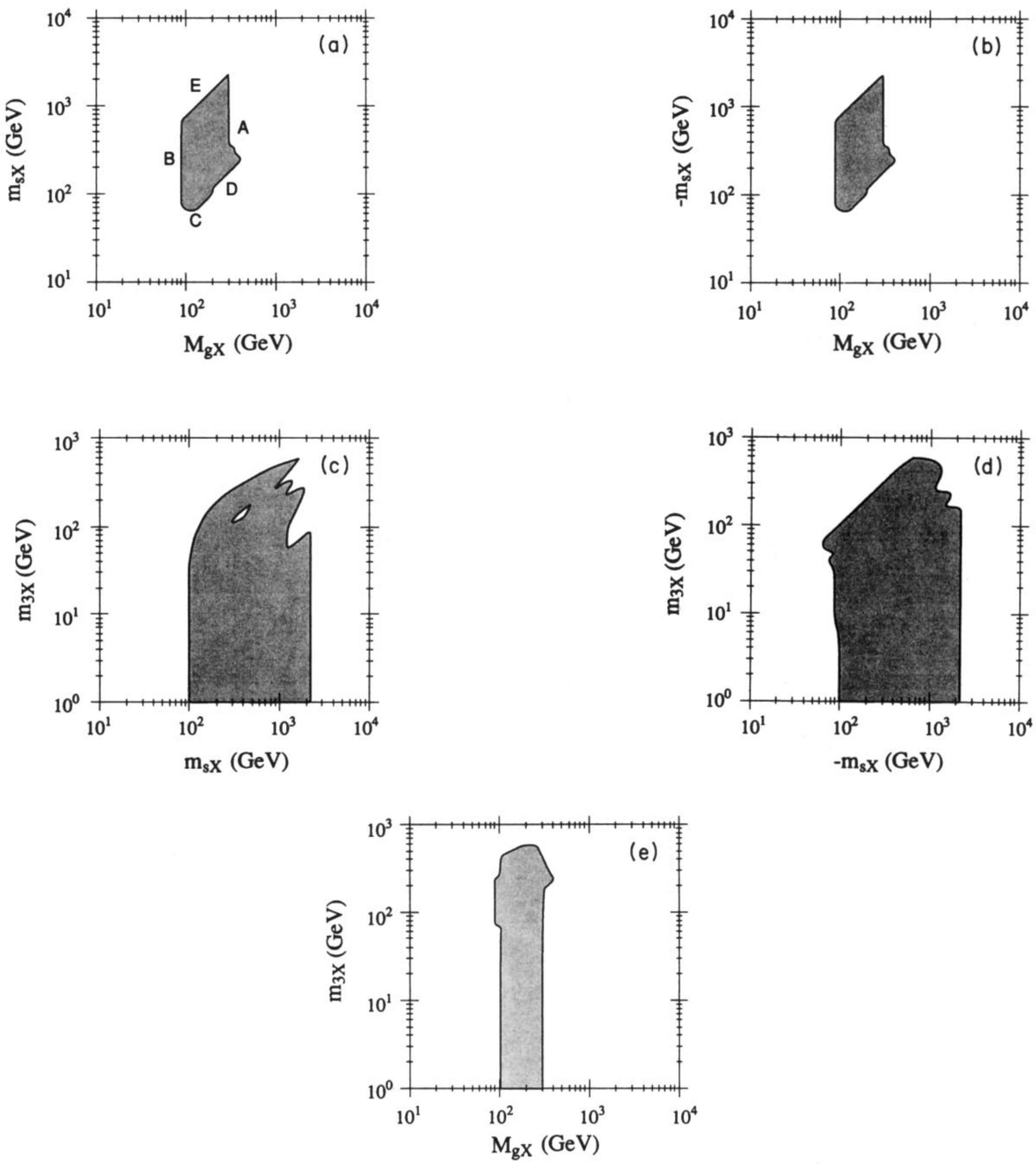

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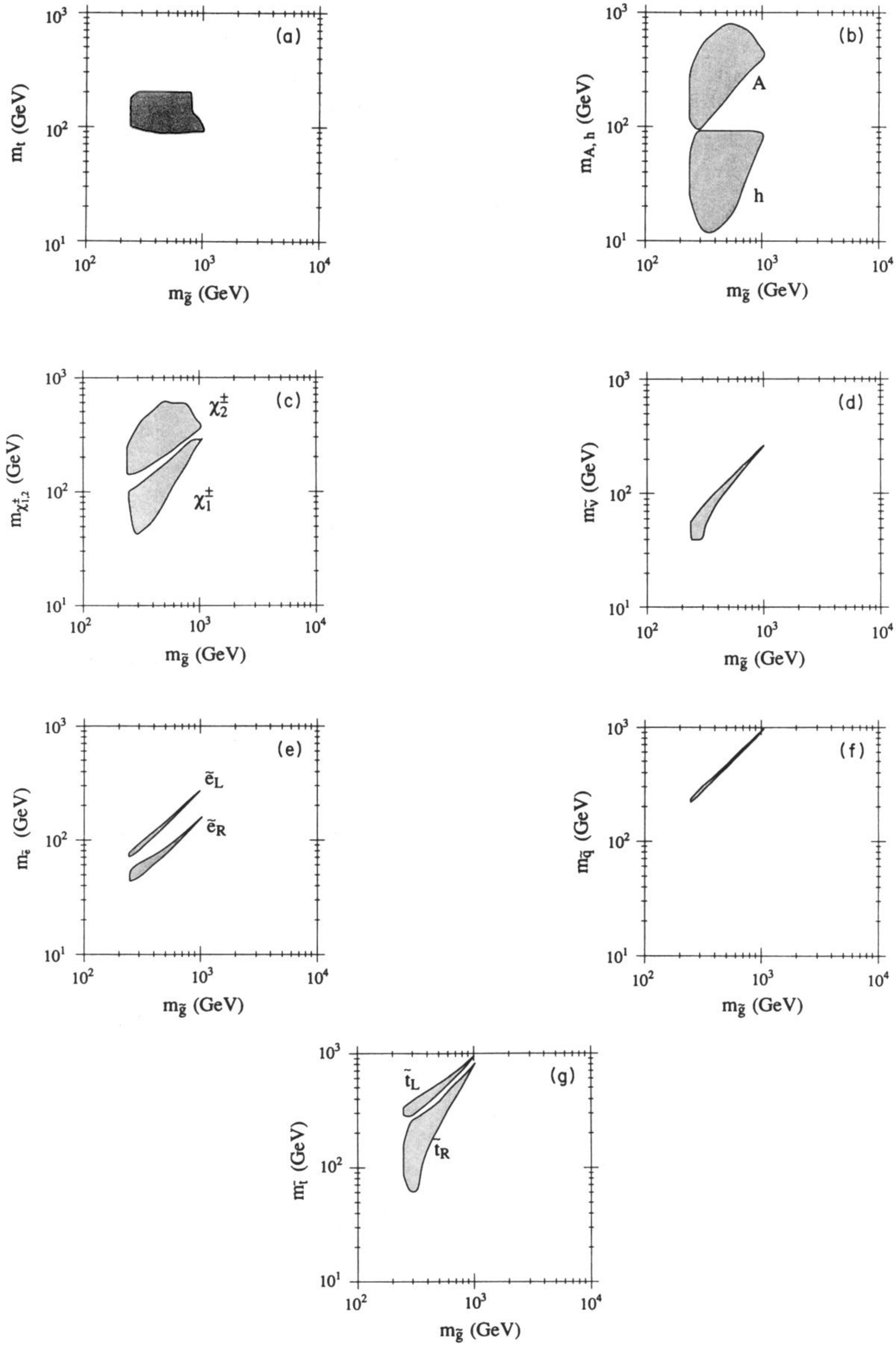

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