Soft supersymmetry breaking masses in a unified model with doublet-triplet splitting

Daijiro Suematsu*
Institute for Theoretical Physics, Kanazawa University, Kanazawa 920-1192, Japan

(Received 17 December 2002; published 25 April 2003)

We study soft supersymmetry breaking parameters in a supersymmetric unified model which potentially solves the doublet-triplet splitting problem. In the model the doublet-triplet splitting is solved by the discrete symmetry which is allowed to be introduced due to the direct product structure of the gauge group. The messenger fields for the gauge mediated supersymmetry breaking are naturally embedded in the model. The discrete symmetry required by the doublet-triplet splitting makes the gaugino masses nonuniversal and also induces a different mass spectrum for the scalar masses from the ordinary minimal gauge mediation model. Independent physical CP phases can remain in the gaugino sector even after the R transformation.

DOI: 10.1103/PhysRevD.67.075020 PACS number(s): 12.60.Jv, 12.10.Dm

I. INTRODUCTION

Supersymmetry is now considered to be the most promising candidate for the solution of the gauge hierarchy problem. Although we still have no direct evidence of supersymmetry, the unification shown by the gauge couplings in the minimal supersymmetric standard model (MSSM) may be considered as its indirect signal. When we consider grand unified models such as SU(5), SO(10), etc., based on this gauge coupling unification, we are often annoyed by the doublet-triplet splitting problem [1]. The reason is that, as stressed in [2], we cannot easily introduce a suitable symmetry to resolve the doublet-triplet degeneracy in a consistent way with the unified gauge structure. Recently, it has been pointed out that the doublet-triplet splitting problem can be solved by extending the gauge structure such as a deconstruction model [2] or introducing extra dimensions [3]. Since the doublet-triplet splitting problem is almost general in the grand unified models including the superstring model, it seems to be interesting to find models which can solve this problem and also to investigate the phenomenological features in such models.

In this article we propose a supersymmetry (SUSY) breaking scenario which is naturally introduced in a unified model which can solve the doublet-triplet splitting problem. The model is constructed by extending the deconstruction model given in [2]. The similar structure to the minimal gauge mediated supersymmetry breaking model [4–7] is automatically built in the model. In our model, gaugino masses seem to be generally nonuniversal and then have nonuniversal phases due to the gauge structure which is required to realize the doublet-triplet splitting. The soft scalar masses can also have a different spectrum from the ordinary one keeping the flavor blindness. We study the general feature of the supersymmetry breaking parameters in addition to the structure of the CP phases in this model.

This paper is organized as follows. In Sec. II we define our model and explain how the doublet-triplet splitting can be realized. In Sec. III we derive the soft supersymmetry breaking parameters based on such a feature of the model and give some comments on their phenomenological aspects. Section IV is devoted to the summary.

II. A SUSY MODEL WITH THE DOUBLET-TRIPLET SPLITTING

We consider a model with a direct product gauge structure such as $\mathcal{G}=\mathrm{SU}(5)'\times\mathrm{SU}(5)''$ and a global discrete symmetry F which commutes with this gauge symmetry [2]. Under this gauge structure we introduce bifundamental chiral superfields $\Phi_1(\bar{\mathbf{5}},\mathbf{5})$ and $\Phi_2(\mathbf{5},\bar{\mathbf{5}})$, an adjoint Higgs chiral superfield $\Sigma(\mathbf{1},\mathbf{24})$, three sets of chiral superfields $\Psi_{10}(\mathbf{10},\mathbf{1})+\Psi_{\bar{5}}(\bar{\mathbf{5}},\mathbf{1})$ which correspond to three generations of quarks and leptons, a set of chiral superfield $H(\mathbf{5},\mathbf{1})+\tilde{H}(\mathbf{1},\bar{\mathbf{5}})$ which contains Higgs doublets, and also a set of chiral superfield $\bar{\chi}(\bar{\mathbf{5}},\mathbf{1})+\chi(\mathbf{1},\mathbf{5})$ in order to cancel the gauge anomaly induced by the above contents. We additionally introduce several singlet chiral superfields S_α. These are summarized in Table I. In order to induce the symmetry breaking at the high energy scale we introduce a superpotential such as

$$W_1=M_\phi\mathrm{Tr}(\Phi_1\Phi_2)+\frac{1}{2}M_\sigma\mathrm{Tr}(\Sigma^2)+\lambda\,\mathrm{Tr}\left(\Phi_1\Sigma\Phi_2+\frac{1}{3}\Sigma^3\right). \tag{1}$$

The scalar potential based on this W_1 can be easily obtained as

$$V=\mathrm{Tr}|M_\phi\phi_1+\lambda\phi_1\sigma+y|^2+\mathrm{Tr}|M_\phi\phi_2+\lambda\sigma\phi_2+x|^2+\mathrm{Tr}|M_\sigma\sigma+\lambda\phi_1\phi_2+\sigma^2+z|^2, \tag{2}$$

where $\phi_{1,2}$ and σ are the scalar components of $\Phi_{1,2}$ and Σ, respectively. They are traceless and x, y, and z are the Lagrange multipliers for these traceless conditions.

We try to find a nontrivial and physically interesting solution of the minimum of this scalar potential. The conditions for it can be written in such a way as

$$\phi_2=\frac{x}{y}\phi_1, \tag{3}$$

*Email address: suematsu@hep.s.kanazawa-u.ac.jp

TABLE I. Discrete charge assignment for the chiral superfields.

	$\mathcal{F}(\mathcal{G}$ representation)	F	F'	
			$\mathbf{3}\in\mathbf{5}$ or $\overline{\mathbf{3}}\in\overline{\mathbf{5}}$	$\mathbf{2}\in\mathbf{5}$ or $\overline{\mathbf{2}}\in\overline{\mathbf{5}}$
Quarks/leptons	$\Psi_{10}^{j}(\mathbf{10},\mathbf{1})$	α	α	α
$(j=1\sim3)$	$\Psi_{5}^{j}(\overline{\mathbf{5}},\mathbf{1})$	β	β	β
Higgs fields	$H(\mathbf{5},\mathbf{1})$	γ	γ	γ
	$\widetilde{H}(\mathbf{1},\overline{\mathbf{5}})$	ξ	$\xi+2a$	$\xi-3a$
Messenger fields	$\overline{\chi}(\overline{\mathbf{5}},\mathbf{1})$	δ	δ	δ
	$\chi(\mathbf{1},\mathbf{5})$	ζ	$\zeta+2b$	$\zeta-3b$
Bifundamental field	$\Phi_1(\overline{\mathbf{5}},\mathbf{5})$	η	$\eta+2c$	$\eta-3c$
	$\Phi_2(\mathbf{5},\overline{\mathbf{5}})$	σ	$\sigma+2d$	$\sigma-3d$
Adjoint Higgs field	$\Sigma(\mathbf{1},\mathbf{24})$	0	0	
Singlets	$S_1(\mathbf{1},\mathbf{1})$	e	e	
	$S_2(\mathbf{1},\mathbf{1})$	f	f	

$$M_\phi\phi_1+\lambda\,\phi_1\sigma+y=0, \tag{4}$$

$$M_\sigma\sigma+\lambda\left(\sigma^2+\frac{x}{y}\phi_1^2\right)+z=0, \tag{5}$$

where the Lagrange multipliers y and z are determined as

$$y=-\frac{\lambda}{5}\mathrm{Tr}(\phi_1\sigma),\quad z=-\frac{\lambda}{5}\mathrm{Tr}\left(\sigma^2-\frac{5x}{\lambda\,\mathrm{Tr}(\phi_1\sigma)}\phi_1^2\right), \tag{6}$$

where x remains as a free parameter. We restrict ourselves to consider a special direction in the field space such that $\phi_1=\kappa\sigma$ and we also assume that $M_\sigma=M_\phi(1+x\kappa^2/y)$ is satisfied. Then, along this direction Eqs. (4) and (5) become consistent with each other and they are reduced to an interesting equation such as

$$M_\phi\sigma+\lambda\,\sigma^2-\frac{\lambda}{5}\mathrm{Tr}(\sigma^2)=0. \tag{7}$$

This equation has the same form as the potential minimum condition for the adjoint Higgs scalar in the ordinary supersymmetric SU(5). It is well known that there are three supersymmetric degenerate independent solutions in this equation. The most interesting one can be written as

$$\sigma=\widetilde{M}\,\mathrm{diag}(2,\ 2,\ 2,\ -3,\ -3), \tag{8}$$

where $\widetilde{M}$ is defined as $\widetilde{M}=M_\phi/\lambda$. In the present discussion we adopt this solution. Using this σ, other fields can be determined as

$$\phi_1=\kappa\sigma,\quad \phi_2=\frac{1}{\kappa}\left(\frac{M_\sigma}{M_\phi}-1\right)\sigma. \tag{9}$$

We have an unfixed parameter κ in this solution. However, if we assume that this model is obtained as a result of a suitable deconstruction, κ can be determined as discussed below. There is no D-term contribution to V from these vacuum expectation values (VEVs) in Eqs. (8) and (9) and then the supersymmetry is conserved at this stage.

It is convenient to use the deconstruction method in order to see what kind of discrete symmetry remains unbroken when these VEVs are induced [2]. We consider the theory space represented by the moose diagram which is composed of the n sites Q_i placed on the vertices of an n-polygon and one site on its center P of this polygon. We assign SU(5)$'$ on the site P and SU(5)$''$ on each site Q_i and also put a bifundamental chiral superfield Φ_i on each link from P to Q_i. On each link from Q_i to Q_{i+1} we put the adjoint Higgs chiral superfield Σ of SU(5)$''$. For the later discussion, we may consider the unitary link variables $U_i\equiv\exp(i\phi_i/\widetilde{M})$ and $\mathcal{W}\equiv\exp(-i\sigma/\widetilde{M})$. Here we introduce an equivalence relation only for the boundary points of the polygon by the $2\pi/n$ rotation and we identify this Z_n symmetry with the above mentioned discrete symmetry F. The equivalence relation defined by F makes Σ independent of i or, equivalently, invariant under F. This makes us consider the reduced theory space composed of only three sites P, Q_1, and Q_2, in which the field contents become equivalent to the one given in Table I. If we use $\mathcal{W}$ introduced above, this equivalence relation requires that $\mathcal{W}^n=1$ is satisfied. Thus we can write $\mathcal{W}$ as

$$\mathcal{W}=\mathrm{diag}(e^{2i\rho},\ e^{2i\rho},\ e^{2i\rho},\ e^{-3i\rho},\ e^{-3i\rho}), \tag{10}$$

where $e^{i\rho}$ is the nth root of unity. If we assume that our model is obtained as a result of the above discussed deconstruction, the condition $U_i\mathcal{W}U_{i+1}^{-1}=1$ should be satisfied for $i=1$, which means that the holonomy around each two-dimensional plaquette is equal to 1.[1] This requirement is interpreted in our vacuum defined by Eqs. (8) and (9) as an additional condition $-\phi_1+\sigma+\phi_2=0$, which can be transformed into a condition for the κ in $\kappa^2-\kappa+1-M_\sigma/M_\phi=0$.

[1]This corresponds to the energy minimum condition from the viewpoint of the lattice gauge [2].

Now we consider the transformation property of this vacuum under the gauge transformation such as

$$U_i' = \omega' U_i \, (\omega'')^{-1}, \quad \mathcal{W}' = \omega'' \, \mathcal{W} \, (\omega'')^{-1}, \tag{11}$$

where ω' and ω'' are the group elements of SU(5)$'$ and SU(5)$''$, respectively. The invariance of U_i and $\mathcal{W}$ shows that the group elements ω of the unbroken gauge group satisfy the condition: $\omega = \omega' = \omega''$ and $[\omega, \mathcal{W}] = 0$. Since we take the VEVs of Higgs scalar fields in such a way as Eqs. (8) and (9), the unbroken gauge group is $\mathcal{H} = \mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ which is a subgroup of the diagonal sum SU(5) of $\mathcal{G}$. Next we consider a discrete symmetry F' as a diagonal subgroup of $F \times G_{\mathrm{U}(1)''}$ where $G_{\mathrm{U}(1)''}$ is a discrete subgroup of a hypercharge in SU(5)$''$. If we write the group elements of F and $G_{\mathrm{U}(1)''}$ as f and ω_D, the transformation of U_i due to F' can be written as

$$U_i' = (fU_i)\,\omega_D^{-1} = U_{i+1}\,\omega_D^{-1}. \tag{12}$$

If we take ω_D as $\mathcal{W}$, we find that U_i is invariant under this transformation due to the relation $U_i \mathcal{W} U_{i+1}^{-1} = 1$ and F' remains unbroken. The invariance of $\mathcal{W}$ is also clear. Thus we can conclude that in this model the symmetry $\mathcal{G} \times F$ breaks down into $\mathcal{H} \times F'$ by considering the vacuum defined by Eqs. (8) and (9). Since the definition of F' contains the discrete subgroup of U(1)$''$ in SU(5)$''$ as its component, every field which has a nontrivial transformation property with respect to SU(5)$''$ can have different charges. We assign the charges of F' for every field as shown in Table I.

In order to solve the doublet-triplet splitting problem, only the color triplet Higgs chiral superfields H_3 and $\tilde{H}_3$ except for the ordinary Higgs chiral superfields H_2 and $\tilde{H}_2$ should become massive when the above discussed symmetry breaking occurs. We should also require the conditions on F' to satisfy various phenomenological constraints in a consistent way with this realization. We impose the following conditions.

(i) Each term in the superpotential W_1 should exist and this requirement imposes the conditions

$$\eta + \sigma + 2(c+d) = 0, \quad \eta + \sigma - 3(c+d) = 0. \tag{13}$$

(ii) The gauge invariant bare mass terms of the fields such as $\Psi_{\bar{5}} H$, $H\bar{\chi}$, $\tilde{H}\chi$, and $S_\alpha S_\beta$ should be forbidden. These conditions are summarized as

$$\beta + \gamma \neq 0, \quad \gamma + \delta \neq 0, \quad \xi + \zeta + 2(a+b) \neq 0,$$

$$\xi + \zeta - 3(a+b) \neq 0, \quad 2e \neq 0, \quad 2f \neq 0, \quad e + f \neq 0. \tag{14}$$

(iii) To realize the doublet-triplet splitting, Yukawa coupling $\Phi_1 H_2 \tilde{H}_2$ should be forbidden although $\Phi_1 H_3 \tilde{H}_3$ is allowed. This gives the conditions such as

$$\gamma + \xi - 3a + \eta - 3c \neq 0, \quad \gamma + \xi + 2a + \eta + 2c = 0. \tag{15}$$

(iv) Yukawa couplings of quarks and leptons, that is, $\Psi_{10}\Psi_{10}H_2$ and $\Psi_{10}\Psi_{\bar{5}}\tilde{H}_2\Phi_1$ should exist. This requires

$$2\alpha + \gamma = 0, \quad \alpha + \beta + \xi - 3a + \eta - 3c = 0. \tag{16}$$

(v) The fields χ and $\bar{\chi}$ should be massless at the $\mathcal{G}$ breaking scale and play the role of the messenger fields of the supersymmetry breaking which is assumed to occur in the S_α sector. These require both the absence of $\Phi_2 \chi \bar{\chi}$ and the existence of the coupling $\Phi_2 S_\alpha \chi \bar{\chi}$. These conditions can be written as

$$\delta + \zeta + 2b + \sigma + 2d \neq 0, \quad \delta + \zeta - 3b + \sigma - 3d \neq 0,$$

$$\delta + \zeta + 2b + \sigma + 2d + e = 0, \quad \delta + \zeta - 3b + \sigma - 3d + f = 0. \tag{17}$$

(vi) The neutrino should be massive and the proton should be stable. This means that $\Phi_{\bar{5}}^2 H_2^2$ should exist and $\Psi_{10}\Psi_{\bar{5}}^2$ and $\Psi_{10}^3 \Psi_{\bar{5}}$ should be forbidden [2]. These require

$$2(\beta + \gamma) = 0, \quad \alpha + 2\beta \neq 0, \quad 3\alpha + \beta \neq 0. \tag{18}$$

Every equation should be understood up to modulus n when we take $F' = Z_n$.

We can easily find an example of the consistent solution for these constraints. For example, if we take $F' = Z_{20}$, such an example can be given as

$$\alpha = \delta = \eta = b = -e = 1, \quad \sigma = \xi = \zeta = -a = 3,$$

$$\gamma = -c = -2, \quad d = -f = 6, \quad \beta = -8, \tag{19}$$

where these charges should be understood up to the modulus 20. We have not taken account of the anomaly of F'. Although this anomaly cancellation seems to require the introduction of new fields and impose the additional constraints on the charges, it does not affect the result of the present phenomenological study of the model. So we do not discuss this problem further here. It should be noted that the existence of the different singlet fields $S_{1,2}$ are generally required in order to make χ and $\bar{\chi}$ play a role of messengers of the supersymmetry breaking. In fact, the F' charges of χ and $\bar{\chi}$ satisfy

$$e - f = -5(b+d) \neq 0 \pmod{n} \tag{20}$$

which is derived from Eq. (17). This feature is caused by the direct product structure of the gauge group which is motivated to realize the doublet-triplet splitting.

We can now consider the physics at the scale after the symmetry breaking due to the VEVs in Eqs. (8) and (9). The massless degrees of freedom are composed of the contents of the MSSM and the fields (q, l) and $(\bar{q}, \bar{\ell})$ which come from $\chi(\mathbf{1},\mathbf{5})$ and $\bar{\chi}(\bar{\mathbf{5}},\mathbf{1})$. We can expect the successful gauge coupling unification for these field contents in the similar way to the MSSM. Under the discrete symmetry F', the superpotential for these fields can be written as

$$W_2=h_1\Psi_{10}\Psi_{10}H_2+h_2\Psi_{10}\Psi_{\bar{5}}\tilde{H}_2+\lambda_1 S_1 q\bar{q}+\lambda_2 S_2 \ell\bar{\ell}. \tag{21}$$

The last three terms effectively appear through the nonrenormalizable terms as a result of the symmetry breaking due to $\langle\phi_1\rangle$ and $\langle\phi_2\rangle$. This feature makes the second term favorable to explain the hierarchy between the masses of the top and bottom quarks [2]. The messenger fields q, $\bar{q}$ and ℓ, $\bar{\ell}$ couple with the different singlet fields $S_{1,2}$. If both their scalar components and F components get the VEVs, they can play the role of messenger fields for the supersymmetry breaking in the observable sector as in the minimal gauge mediation model [4,5]. This is discussed in the next section. Although it seems to be difficult to produce a weak scale μ term within the field contents given in Table I, it may be expected to be generated associated with the supersymmetry breaking by extending the model. We may consider various ways for generating the μ term such as the mechanism of Giudice-Masiero [8] or the model based on the VEV of the singlet field like the next MSSM [9,5]. However, we do not discuss its origin here and treat it only as an effective parameter in the following discussion.

III. SOFT SUSY BREAKING PARAMETERS

In this section we study the soft supersymmetry breaking parameters in the present model. The gauge anomaly cancellation for $\mathcal{G}$ requires us to introduce a set of vectorlike fields χ and $\bar{\chi}$ as mentioned above. Using these fields as the messenger fields, fortunately, we can apply the well-known minimal gauge mediation supersymmetry breaking scenario [4–7] to our model. If the singlet chiral superfields S_1 and S_2 couple with the hidden sector fields which break down the supersymmetry, $q,\bar{q}$ and $\ell,\bar{\ell}$ play the role of messenger fields as in the ordinary scenario. The only difference from the ordinary minimal gauge mediation scenario is that q, $\bar{q}$ and ℓ, $\bar{\ell}$ couple with different singlet chiral superfields S_1 and S_2 in the superpotential W_2 because of the discrete symmetry F'. If we assume that both S_α and F_{S_α} get the VEVs, the gaugino masses and the soft scalar masses are generated through one-loop and two-loop diagrams, respectively. However, the mass formulas are modified from the usual ones since the messenger fields q, $\bar{q}$ and ℓ, $\bar{\ell}$ couple with the different singlets.

The massless vector supermultiplet of SU(5) is written as

$$V=V'\cos\theta+V''\sin\theta, \tag{22}$$

where V' and V'' are the vector supermultiplets of SU(5)$'$ and SU(5)$''$. A mixing angle θ and a new gauge coupling constant g of SU(5) are determined as

$$\frac{1}{g^2}=\frac{1}{g'^2}+\frac{1}{g''^2},\qquad \tan\theta=\frac{g'}{g''}, \tag{23}$$

where g' and g'' are the gauge coupling constants of SU(5)$'$ and SU(5)$''$. The same relations are satisfied for each factor group of $\mathcal{H}$ at the symmetry breaking scale $\tilde{M}$. The gauge coupling constants of $\mathcal{H}$ follow the unification relation $g=g_3=g_2=\sqrt{5/3}g_1$ at $\tilde{M}$. The information on the direct product gauge structure at the high energy region is included in the mixing angle θ. The gauginos become massive due to the mixing between the gauginos of SU(5)$'$ and SU(5)$''$ through the one-loop diagrams shown in Fig. 1. These mass mixings can be estimated as

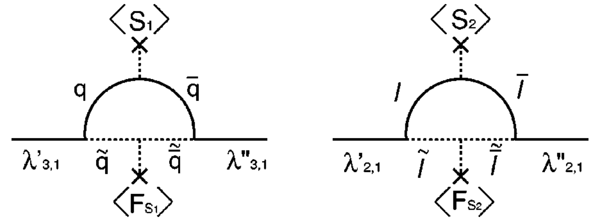

FIG. 1. One-loop diagrams contributing to gaugino masses.

$$M_{\lambda_3'\lambda_3''}=\frac{g_3'g_3''}{16\pi^2}\Lambda_1,\qquad M_{\lambda_2'\lambda_2''}=\frac{g_2'g_2''}{16\pi^2}\Lambda_2,$$

$$M_{\lambda_1'\lambda_1''}=\frac{g_1'g_1''}{16\pi^2}\left(\frac{2}{3}\Lambda_1+\Lambda_2\right), \tag{24}$$

where $\Lambda_\alpha=\langle F_{S_\alpha}\rangle/\langle S_\alpha\rangle$. These can be transformed into the masses M_r of the gauginos λ_r of the gauge group $\mathcal{H}$ by taking account of Eqs. (22) and (23). They can be written in the form as

$$M_3=\frac{\alpha_3}{4\pi}\Lambda_1,\quad M_2=\frac{\alpha_2}{4\pi}\Lambda_2,\quad M_1=\frac{\alpha_1}{4\pi}\left(\frac{2}{3}\Lambda_1+\Lambda_2\right), \tag{25}$$

where $\alpha_r=g_r^2/4\pi$.

These give the same sum rule among gaugino masses as the one of the usual minimal gauge mediation scenario at the supersymmetry breaking scale such as

$$\frac{M_1}{\alpha_1}=\frac{2}{3}\frac{M_3}{\alpha_3}+\frac{M_2}{\alpha_2}. \tag{26}$$

However, depending on the ratio of the scale Λ_1/Λ_2, each mass ratio can be different from the ordinary ones, that is,

$$\frac{M_2}{M_3}=\frac{\alpha_2}{\alpha_3}\frac{\Lambda_2}{\Lambda_1},\quad \frac{M_1}{M_3}=\frac{\alpha_1}{\alpha_3}\left(\frac{2}{3}+\frac{\Lambda_2}{\Lambda_1}\right). \tag{27}$$

These formulas show that M_3 can be much smaller than $M_{1,2}$ in the case of $\Lambda_2>\Lambda_1$. If we take account of the evolution effect by the renormalization group, their values at the weak scale M_W, for example, can be obtained as

$$M_r(M_W)=M_r(\Lambda)\frac{\alpha_r(M_W)}{\alpha_r(\Lambda)}, \tag{28}$$

where Λ is a scale at which the supersymmetry breaking is introduced. Since Λ_α is generally independent, the phases contained in the gaugino masses are nonuniversal even in the case of $|\Lambda_1|=|\Lambda_2|$. We cannot remove them completely by using the R transformation unlike the universal gaugino mass case. In fact, if we define the phases as $\Lambda_\alpha \equiv |\Lambda_\alpha| e^{i\varphi_\alpha}$ and make M_2 real by the R transformation, the phases of M_3 and M_1 are written as

$$\arg(M_3)=\varphi_1-\varphi_2,$$

$$\arg(M_1)=\arctan\left(\frac{2|\Lambda_1|\sin(\varphi_1-\varphi_2)}{3|\Lambda_2|+2|\Lambda_1|\cos(\varphi_1-\varphi_2)}\right). \tag{29}$$

The scalar masses are induced as the values of $O(|\Lambda_\alpha|^2)$ through the two-loop diagrams as in the ordinary case. Again, only difference comes from the fact that the model has the direct product gauge structure at the high energy region and the messengers $(q,\ \ell)$ and $(\bar{q},\ \bar{\ell})$ are the representations of the different factor groups. This brings the mixing factor between the vector superfields V_r and V_r', V_r'' as in the gaugino mass case. Taking account of this, their formulas can be written as

$$\tilde{m}_f^2=2\left[C_3\left(\frac{\alpha_3}{4\pi}\right)^2+\frac{2}{3}\left(\frac{Y}{2}\right)^2\left(\frac{\alpha_1}{4\pi}\right)^2\right]|\Lambda_1|^2+2\left[C_2\left(\frac{\alpha_2}{4\pi}\right)^2+\left(\frac{Y}{2}\right)^2\left(\frac{\alpha_1}{4\pi}\right)^2\right]|\Lambda_2|^2, \tag{30}$$

where $C_3=4/3$ and 0 for the SU(3) triplet and singlet fields, and $C_2=3/4$ and 0 for the SU(2) doublet and singlet fields, respectively. The hypercharge Y is expressed as $Y=2(Q-T_3)$. These formulas can give rather different mass spectrum for the scalar fields depending on the values of Λ_1/Λ_2. In fact, if we assume $\Lambda_1<\Lambda_2$, for example, the mass difference between the color singlet fields and the colored fields tends to be smaller in comparison with the one in the ordinary scenario. Let us take $\Lambda_1=40$ TeV and $\Lambda_2=100$ TeV to show a typical spectrum of the superpartners at the supersymmetry breaking scale. Then we can have the following spectrum as

$$\begin{aligned}
&M_3=273\ \text{GeV},\quad M_2=279\ \text{GeV},\quad M_1=111\ \text{GeV},\\
&\tilde{m}_Q=562\ \text{GeV},\\
&\tilde{m}_U=455\ \text{GeV},\quad \tilde{m}_D=449\ \text{GeV},\quad \tilde{m}_L=347\ \text{GeV},\\
&\tilde{m}_E=130\ \text{GeV},\\
&m_1=m_2=347\ \text{GeV},
\end{aligned} \tag{31}$$

where m_1 and m_2 are masses of the Higgs scalars that couple with the down and up sectors of quarks and leptons, respectively. These masses are somewhat affected by the renormalization group running effect, although the running region is not so large. For example, the modifications due to this effect can be solved analytically for the masses of sleptons and H_1, for which Yukawa coupling effects can be neglected, as [7]

$$\begin{aligned}
\tilde{m}^2_{\ell_L}(M_W)=&\tilde{m}^2_{\ell_L}(\Lambda)-\tfrac{3}{2}|M_2(\Lambda)|^2\left(\frac{\alpha_2^2(M_W)}{\alpha_2^2(\Lambda)}-1\right)\\
&-\tfrac{1}{22}|M_1(\Lambda)|^2\left(\frac{\alpha_1^2(M_W)}{\alpha_1^2(\Lambda)}-1\right),\\
\tilde{m}^2_{\ell_R}(M_W)=&\tilde{m}^2_{\ell_R}(\Lambda)-\tfrac{2}{11}|M_1(\Lambda)|^2\left(\frac{\alpha_1^2(M_W)}{\alpha_1^2(\Lambda)}-1\right),
\end{aligned} \tag{32}$$

where we do not write the D-term contribution explicitly. The mass m_1^2 of the Higgs scalar has the same formula as $\tilde{m}^2_{\ell_L}$.

As in the minimal gauge mediation model discussed in [7], the soft supersymmetry breaking A_f and B parameters can also be expected to be induced through the radiative correction such as

$$A_f\simeq A_f(\Lambda)+M_2(\Lambda)(-1.85+0.34|h_t|^2),$$

$$\frac{B}{\mu}\simeq\frac{B}{\mu}(\Lambda)-\tfrac{1}{2}A_t(\Lambda)+M_2(\Lambda)(-0.12+0.17|h_t|^2), \tag{33}$$

where we should omit a term of h_t in the expression of A_f except for the top sector $(f=t)$. In the case of $A_f(\Lambda)=B(\Lambda)=0$, which are expected in many gauge mediation scenarios, A_f and B are proportional to M_2 and then the CP phases in the soft supersymmetry breaking parameters are completely rotated away as far as gaugino masses are universal [7]. However, in the present model this situation is broken even in the case of $A_f(\Lambda)=B(\Lambda)=0$, since the phases in the gaugino masses are not universal. Although the generation of B should be considered on the basis of the various mechanisms like the μ term [9] also in the present model, it is completely model dependent and we do not discuss it further.

Finally we comment on some phenomenological aspects on these soft breaking parameters. At present it seems to be difficult to relate the supersymmetry breaking parameters to the observed values. Only exception might be found in the electroweak symmetry breaking. As is well known, the minimum condition of the tree level scalar potential in the MSSM can be written as

$$m_Z^2=-2\mu^2+2\frac{m_1^2-m_2^2\tan^2\beta}{\tan^2\beta-1}. \tag{34}$$

Supersymmetry breaking parameters in the right-hand side can be estimated by using the one-loop renormalization group equations (RGEs). Through the semianalytic calculation [10], their weak scale values can be expressed using various soft parameters at the supersymmetry breaking scale Λ whose examples are shown in Eq. (32). If we take $\Lambda=100$ TeV, $m_{\rm top}=170$ GeV, and $\tan\beta=5$ and use the nu-

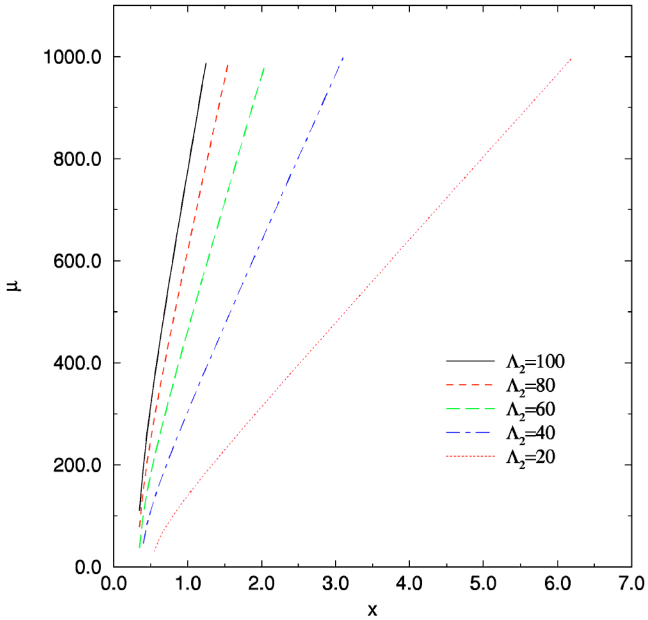

FIG. 2. Values of μ required to realize a correct vacuum for the various SUSY breaking scales $\Lambda_{1,2}$ TeV.

merical coefficient obtained through the RGEs in this case [11], Eq. (34) is expressed as[2]

$$m_Z^2 = -1.8\mu^2 - 0.2M_2^2 + 0.4M_3^2 + 0.2A_t^2 + 0.4\tilde{m}_{Q_3}^2 + 0.4\tilde{m}_{U_3}^2 - 1.7m_2^2 - 0.2A_tM_3 + \cdots, \quad (35)$$

where the ellipses represent the subdominant contributions. Assuming $\Lambda_{1,2}$ to be real and substituting soft parameters given by Eqs. (25) and (30), we obtain

$$m_Z^2 = (115x^2 + 6.1x - 13.2)(10^{-3}\Lambda_2)^2 - 1.8\mu^2, \quad (36)$$

where $x = \Lambda_1/\Lambda_2$ and we use Eq. (33) with $A_t(\Lambda) = 0$. In Fig. 2 we plot the values of μ satisfying this relation for various values of $\Lambda_{1,2}$. This shows that μ can take a reasonable value as far as we can set up $\Lambda_{1,2}$ appropriately. As a general feature we find that the large Λ_2 tends to require the large value of μ. The sensitivity of μ against x seems to be almost independent of the value of x in the $x \gtrsim 0.5$ region. In the $x \lesssim 0.5$ region, we can obtain a small value of μ such as $\mu \sim 100$ GeV as the consistent solution. However, it is necessary to tune carefully the value of x to be 0.35–0.5 depending on Λ_2 to obtain the smaller value of μ. This required tuning is finer for the larger Λ_2 value. Anyway, this feature looks different from the ordinary minimal gauge mediation model in which $x = 1$ is satisfied and then $\mu \gtrsim 300$ GeV is required for $\Lambda_2 \geqslant 40$ TeV as seen from Fig. 2.[3]

There is another interesting feature in this case. In the usual minimal gauge mediation scenario the lightest superparticle in the whole spectrum except for the gravitino is the right-handed slepton as far as we do not take account of the radiative effect. However, in the present scenario the photino can be the lightest one in this situation even at the tree level as we can see it in the example given in Eq. (31). This feature might be relevant to the event such as $e^+e^-\gamma\gamma +$ missing E_T [6,7]. Our model might be discriminated from other gauge mediation models by using this aspect.

The gaugino mass universality seems to be a rather general result in various supersymmetry breaking scenarios. However, the present model naturally induces nonuniversal gaugino masses as a result of intrinsic nature of the model. We generally have physical CP phases in the gaugino sector. This may be dangerous since it can give a large contribution to the electric dipole moment of a neutron and an electron. However, they could be within the experimental bound even if the CP phases are $O(1)$. It is expected that there can be an effective cancellation between the chargino and neutralino contributions to them [12]. We can check this in the present model and the result will be presented elsewhere [13]. In the case that there is no contradiction with the electric dipole moment, these large CP phases may be important when we consider the electroweak baryogenesis [14].

IV. SUMMARY

We investigated the soft supersymmetry breaking masses in the supersymmetric unified model which can solve the doublet-triplet splitting problem. The model is constructed through the deconstruction by extending the gauge structure into the direct product group $SU(5)' \times SU(5)''$. The low energy spectrum is the one of the MSSM with the additional chiral superfields which can play a role of messengers in the gauge mediated supersymmetry breaking. The gauge anomaly cancellation requires to introduce these chiral superfields. The discrete symmetry can be introduced to realize the doublet-triplet splitting because of the the direct product gauge structure. It forces the color triplet and color singlet messengers to couple with the different singlet chiral superfields whose scalar and auxiliary components are assumed to get the VEVs due to the hidden sector dynamics. This can make the different structure of the soft supersymmetry breaking masses from the ones of the ordinary minimal gauge mediation scenario. One of the interesting feature is that the mass difference between the colored fields and the color singlet fields can be smaller in comparison with the ordinary gauge mediation scenario. Another interesting point is that the gaugino masses become nonuniversal generally and the nonuniversal phases are introduced in the gaugino masses.

[2]In this expression we do not take account of the difference between Λ_1 and Λ_2 for the estimation of the numerical coefficients. However, we can expect that there is no substantial difference even if we take account of it.

[3]The small Λ_2 makes M_2 too small and it will be excluded from the fact that neutralinos and charginos have not been found at the CERN e^+e^- collider LEP.

The CP phases can remain in the gaugino sector as the physical phases after the R transformation. This feature may discriminate this model from others since it is rather difficult to construct the well-motivated model with the nonuniversal gaugino masses. Further phenomenological study of the model seems to be worthy since it is constructed on the basis of the reasonable motivation to solve the doublet-triplet splitting problem in the grand unified model.

ACKNOWLEDGMENTS

This work is supported in part by a Grant-in-Aid for Scientific Research (C) from Japan Society for Promotion of Science (Grant No. 14540251) and also by a Grant-in-Aid for Scientific Research on Priority Areas (A) from The Ministry of Education, Science, Sports and Culture (Grant No. 14039205).

[1] For a review, see for example, L. Randall and C. Csáki, in *Particles, Strings, and Cosmology*, edited by J. Bagger (World Scientific, Singapore, 1996).

[2] E. Witten, hep-ph/0201018.

[3] Y. Kawamura, Prog. Theor. Phys. **105**, 999 (2001); G. Altarelli and F. Feruglio, Phys. Lett. B **511**, 257 (2001).

[4] M. Dine and A.E. Nelson, Phys. Rev. D **51**, 1362 (1995); M. Dine, A.E. Nelson, Y. Nir, and Y. Shirman, *ibid.* **53**, 2658 (1996).

[5] M. Dine, Y. Nir, and Y. Shirman, Phys. Rev. D **55**, 1501 (1997).

[6] S. Dimopoulos, M. Dine, S. Raby, and S. Thomas, Phys. Rev. Lett. **76**, 3494 (1996).

[7] K.S. Babu, C. Kolda and F. Wilczek, Phys. Rev. Lett. **77**, 3070 (1996).

[8] G.F. Giudice and A. Masiero, Phys. Lett. B **206**, 480 (1988).

[9] M. Leurer, Y. Nir, and N. Seiberg, Nucl. Phys. **B420**, 468 (1994); T. Gherghetta, G. Jungman, and E. Poppitz, hep-ph/9511317. G. Dvali, G.F. Giudice, and A. Pomarol, Nucl. Phys. **B478**, 31 (1996).

[10] M. Carena, P. Chankowski, O. Olechowski, S. Pokorski, and C.E.M. Wagner, Nucl. Phys. **B491**, 103 (1997); G.L. Kane and S.F. King, Phys. Lett. B **451**, 113 (1991).

[11] G.L. Kane, J. Lykken, B.D. Nelson, and L.-T. Wang, Phys. Lett. B **551**, 146 (2003).

[12] M. Brhlik, G.J. Good, and G.L. Kane, Phys. Rev. D **59**, 115004 (1999); M. Brhlik, L. Everett, G.L. Kane, and J. Lykken, Phys. Rev. Lett. **83**, 2124 (1999); Phys. Rev. D **62**, 035005 (2000); R. Arnowitt, B. Dutta, and Y. Santoso, *ibid.* **64**, 113010 (2001).

[13] D. Suematsu and H. Tsuchida (in preparation).

[14] M. Brhlik, G.J. Good, and G.L. Kane, Phys. Rev. D **63**, 035002 (2001).