# Neutron Electric Dipole Moment under Non-Universal Soft SUSY Breaking Terms

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# Neutron Electric Dipole Moment under Non-Universal Soft SUSY Breaking Terms

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The electric dipole moment of the neutron (EDMN) is re-examined in the framework of non-universal soft supersymmetry breaking. We review some features of the relation between the EDMN and non-universal soft supersymmetry breaking terms. It is shown that the constraints on the soft scalar masses and soft $CP$ phases have a rather strong dependence on the non-universality of soft breaking terms. We also show that the soft $CP$ phase $\phi_B$ which has no natural suppression mechanism may not contribute significantly to the EDMN in a certain parameter region where the radiative symmetry breaking occurs successfully. If this is the case, $\phi_B$ may not need to be so small.

## § 1. Introduction

The minimal supersymmetric standard model (MSSM) is now considered as the most promising extension of the standard model (SM).[1] Although the origin of the supersymmetry breaking is still unknown, we can make various predictions by using a suitable parametrization of its breaking. This parametrization is known as the soft supersymmetry breaking terms. Phenomenological features of the MSSM are determined by these soft supersymmetry breaking parameters, which play a phenomenologically similar role to the vacuum expectation value of the Higgs field in the SM. Usually these soft supersymmetry breaking parameters are assumed as universal. Although universal soft breaking terms can be derived from a special type of supergravity theory, this assumption is based on the predictability of the model and also on phenomenological constraints. As is wellknown, the MSSM gives a new contribution to FCNC processes, such as $K^0$-$\bar{K}^0$ mixing, at the one-loop level.[2)~4)] To suppress such processes sufficiently, a suitable degeneracy is required among squarks as far as squark masses are assumed to be less than $O(1)$ TeV.[4)] Various models have been proposed to guarantee the degeneracy of squark masses.[5)~9)]

On the other hand, as recently stressed, soft supersymmetry breaking parameters are non-universal in the effective theories derived from the superstring theories and also general supergravity theories.[10)~12)] And various interesting features which are not seen in usual study are found under those soft breaking parameters.[13),27)] If we take such a situation seriously, it seems to be a very interesting problem to consider how this kind of non-universality does not conflict with the phenomenological constraints such as $K^0$-$\bar{K}^0$ mixing. One attractive proposal suggests that squark masses can be degenerate at low energy due to the heavy gaugino loop effect even if they are non-universal at $M_{\text{pl}}$.[6),7)] As another possibility, we can consider the non-universal squark masses which preserve their non-degeneracy in the low energy region but do

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not conflict with the FCNC constraints. This type of soft supersymmetry breaking parameter can be realized in rather constrained supergravity models. It is shown, however, that such non-universality shows very interesting effects, especially, in the gauge coupling unification.[13]

The electric dipole moment of the neutron (EDMN) in the MSSM also results in a severe constraint on the soft supersymmetry breaking parameters. Unlike the FCNC, however, it requires no degeneracy among squark masses, but constrains the absolute values of soft breaking parameters. The present experimental bound of the absolute value of the EDMN is $1.1\times10^{-25}$e cm.[14] As estimated in many pioneering works,[15] the contribution to the EDMN containing the gluino and squarks in the internal lines gives the following constraint,

$$\gamma\left|\frac{m_g A_U}{\tilde{m}^2_{q,q^c}}\left(\frac{100\,\mathrm{GeV}}{\mathrm{Max}(m_g,\,\tilde{m}_{q,q^c})}\right)^2\frac{m_{\mathrm{up}}}{10\,\mathrm{MeV}}\right|\lesssim 10^{-3}\,, \tag{1}$$

where $A_U$ is a soft supersymmetry breaking parameter corresponding to the up-quark Yukawa coupling. This bound is usually considered to be satisfied if either $CP$ phase $\gamma$ in the soft breaking parameters are unnaturally small as $O(10^{-3})$ or the gluino mass $m_g$ and averaged squark masses $\tilde{m}_{q,q^c}$ are heavier than $O(1)$TeV.[15),16] Although more detailed studies in various parameter settings give similar results,[17] these conditions seem to be inconsistent with the naturalness if we consider it seriously. This suggests that the EDMN may be a very important phenomenon to study the evidence of the supersymmetry and the origin of the supersymmetry breaking. As known from Eq. (1), the EDMN crucially depends on the structure of the soft supersymmetry breaking parameters. In fact, Eq. (1) shows that the non-universal scale-up of the soft breaking parameters can suppress the EDMN more effectively than their universal scale-up. Thus its detailed re-examination under non-universal soft supersymmetry breaking seems to be worthy for the study of the supersymmetric model building. It may also affect the electroweak baryogenesis scenario, since it is closely related to the largeness of soft $CP$ phases.[18]

In the present paper, we investigate the EDMN in the MSSM with non-universal soft supersymmetry breaking terms.*) We do not use the averaged values for the squark masses, unlike the usual study. The estimation based on the averaged squark masses is good for the case that squark masses are almost degenerate. However, if their degeneracy is not sufficient, the value of the EDMN can change by one order. And also, particularly, we carefully study the relation between the EDMN and the soft $CP$ phases. In § 2 we review the general feature of soft supersymmetry breaking terms induced by the $F$-term breaking of the dilaton and moduli fields. In the MSSM it is wellknown that the EDMN has a non-zero value due to the loop effect of the gluinos, the charginos and the neutralinos. In § 3 we concentrate on only the gluino contribution to the EDMN. Taking into account the non-universality of soft supersymmetry breaking terms, we give some formulae for the EDMN from which we can read off the effects of the non-universality. Features of the constraints on squark

*) The MSSM often includes an assumption of the universality on the soft supersymmetry breaking parameters. However, in this paper we will use this terminology to refer to the model without such an assumption.

masses and soft $CP$ phases $\phi_{Af_i}$ and $\phi_B$ are discussed on the basis of their numerical analysis. In § 4 a suppression mechanism of the contribution to the EDMN coming from the $CP$ phases in the soft supersymmetry breaking terms is discussed. We propose a new possibility to suppress the $\phi_B$ contribution to the EDMN even in the case that $\phi_B$ is not small. To confirm that such a mechanism works successfully, the chargino contribution should be taken into account. We will also comment on the chargino contribution. Section 5 will be devoted to a summary.

## § 2. Soft SUSY breaking parameters

In this section we briefly review the general formulae and new $CP$ phases of the soft supersymmetry breaking terms in the MSSM. The superpotential of the MSSM is written as

$$W^{\mathrm{eff}}=h_{IJ}^{U}Q^{I}H_2U^{J}+h_{IJ}^{D}Q^{I}H_1D^{J}+h_{IJ}^{E}L^{I}H_1E^{J}+\mu H_1H_2\,, \tag{2}$$

where $I$ and $J$ are the generation indices. The soft supersymmetry breaking terms are

$$\mathcal{L}_{\mathrm{soft}}=-\sum_i \tilde{m}_i^2|z_i|^2-(A_{IJ}^{U}h_{IJ}^{U}Q^{I}H_2U^{J}+A_{IJ}^{D}h_{IJ}^{D}Q^{I}H_1D^{J}+A_{IJ}^{E}h_{IJ}^{E}L^{I}H_1E^{J}$$
$$+B\mu H_1H_2+\sum_a\frac{1}{2}M_a\bar{\lambda}_a\lambda_a+\mathrm{h.c.})\,, \tag{3}$$

where the first term represents the mass terms of all the scalar components in the MSSM. In the last term, $\lambda_a$ are the gaugino fields for the gauge group specified by $a(a=3,2,1)$. The remaining terms are trilinear and bilinear scalar couplings. Although we use the same notation for the superfields and component fields here, they should not be confused.

Various works based on superstring theories and also general supergravity theories suggest that these soft breaking parameters are generally non-universal.[10)~12)] In general, low energy effective supergravity theories are characterized in terms of the Kähler potential $K$, the superpotential $W$ and the gauge kinetic function $f_a$. Each of these is a function of ordinary massless chiral matter superfields $\Psi^I$ and gauge singlet fields $\Phi^i$ called moduli, whose potential is perturbatively flat as long as the supersymmetry is unbroken. Usually it is assumed that nonperturbative phenomena such as gaugino condensation occur in a hidden sector. After integrating out the fields relevant to these phenomena, the Kähler potential and the superpotential are expanded in the low energy observable matter fields $\Psi^I$ as

$$K=\kappa^{-2}\hat{K}(\Phi,\bar{\Phi})+Z(\Phi,\bar{\Phi})_{I\bar{J}}\Psi^I\bar{\Psi}^{\bar{J}}+\left(\frac{1}{2}Y(\Phi,\bar{\Phi})_{IJ}\Psi^I\Psi^J+\mathrm{h.c.}\right)+\cdots\,, \tag{4}$$

$$W=\hat{W}(\Phi)+\frac{1}{2}\tilde{\mu}(\Phi)_{IJ}\Psi^I\Psi^J+\frac{1}{3}\tilde{h}(\Phi)_{IJK}\Psi^I\Psi^J\Psi^K+\cdots\,, \tag{5}$$

where $\kappa^2=8\pi/M_{\mathrm{pl}}^2$. The ellipses stand for terms of higher order in $\Psi^I$. In Eq. (5), $\hat{W}(\Phi)$ and $\tilde{\mu}(\Phi)$ are induced by nonperturbative effects in the hidden sector. Using these functions, the scalar potential $V$ can be written as,[19)]

$$V=\kappa^{-2}e^{G}[G_\alpha(G^{-1})^{\alpha\bar{\beta}}G_{\bar{\beta}}-3\kappa^{-2}]+(D\text{-term}), \tag{6}$$

where $G=K+\kappa^{-2}\log\kappa^6|W|^2$, and the indices $\alpha$ and $\beta$ denote $\Psi^I$ as well as $\Phi^i$. The gravitino mass $m_{3/2}$ which characterizes the scale of supersymmetry breaking is expressed as

$$m_{3/2}=\kappa^2 e^{\hat{K}/2}|\hat{W}|. \tag{7}$$

In order to obtain the soft supersymmetry breaking terms in the low energy effective theory from Eq. (6), we take the flat limit $M_{\rm Pl}\to\infty$ with $m_{3/2}$ fixed. Through this procedure, we get the superpotential (2) and the soft supersymmetry breaking terms (3). In the effective superpotential (2), Yukawa couplings are rescaled as $h_{IJK}=e^{\hat{K}/2}\tilde{h}_{IJK}$, and the $\mu$ term is effectively expressed as

$$\mu=e^{\hat{K}/2}\tilde{\mu}+m_{3/2}Y-F^{\bar{J}}\partial_{\bar{J}}Y, \tag{8}$$

where $\mu$ should be understood as $\mu_{H_1H_2}$. If $|F^i|=O(m_{3/2})$, we can find that the appropriate $\mu$ scale of order $m_{3/2}$ can be remarkably induced through the second and third terms. This is originated from the $Y(\Phi,\bar{\Phi})$ in the Kähler potential.[20] However, it should also be noted that the scale of $\mu$ crucially depends on its origin. A case such as $|\mu|/m_{3/2}\ll 1$ can also occur if $Y(\Phi,\bar{\Phi})\ll 1$ and $|\tilde{\mu}|\ll m_{3/2}$. This case is interesting when considering the suppression of the effect of the soft $CP$ phases on the EDMN as seen later.

Each soft breaking term is expressed by using $K$ and $W$ as follows,[11),*)]

$$\tilde{m}^2_{I\bar{J}}=m^2_{3/2}Z_{I\bar{J}}-F^i\bar{F}^{\bar{J}}[\partial_i\partial_{\bar{J}}Z_{I\bar{J}}-(\partial_{\bar{J}}Z_{N\bar{J}})Z^{N\bar{L}}(\partial_i Z_{I\bar{L}})]+\kappa^2 V_0 Z_{I\bar{J}}, \tag{9}$$

$$A_{IJ}=F^i\left[\left(\partial_i+\frac{1}{2}\hat{K}_i\right)h_{IJN}-Z^{\bar{M}L}\partial_i Z_{\bar{M}(I}h_{JN)L}\right]/h_{IJN}, \tag{10}$$

$$\begin{aligned}B=&F^i\left[\left(\partial_i+\frac{1}{2}\hat{K}_i\right)\mu-Z^{\bar{M}L}\partial_i Z_{\bar{M}(I}\mu_{J)L}\right]/\mu-m_{3/2}\\&+\left[F^i\left(\partial_i+\frac{1}{2}\hat{K}_i\right)F^{\bar{J}}-2m_{3/2}F^{\bar{J}}\right]\partial_{\bar{J}}Y/\mu,\end{aligned} \tag{11}$$

where $F^i$ is an $F$-term of $\Phi^i$, and $\partial_i$ denotes $\partial/\partial\Phi^i$. The suffix $N$ in Eq. (10) represents the Higgs fields $H_1$ and $H_2$. $V_0$ in Eq. (9) is the cosmological constant expressed as $V_0=\kappa^{-2}(F^i\bar{F}^{\bar{J}}\partial_i\partial_{\bar{J}}\hat{K}-3m^2_{3/2})$. Requiring the cosmological constant to be zero or sufficiently small, we get $|F^i|=O(m_{3/2})$. From this we find that the soft breaking terms $\tilde{m}_{I\bar{J}}$, $A_{IJ}$ and $B$ are generally non-universal but characterized by the gravitino mass $m_{3/2}$. The structure of non-universality depends on the form of the Kähler potential. The gaugino masses $M_a$ are derived through the following formula,[19]

$$M_a=\frac{1}{2}({\rm Re}f_a)^{-1}F^j\partial_j f_a, \tag{12}$$

where the subscript $a$ represents a gauge group. This shows that $M_a$ is also char-

*) In these formulae we do not assume that the cosmological constant vanishes. It should be noted that these soft breaking parameters are not canonically normalized because the kinetic term of $\Psi_I$ is expressed as $Z_{I\bar{J}}\partial^\mu\Psi^I\partial_\mu\bar{\Psi}^{\bar{J}}$.

acterized by $m_{3/2}$.

Here we should comment on the possibility of the hierarchically different soft breaking parameters. The values of these soft breaking terms are at most $O(m_{3/2})$ as shown in Eqs. (9)~(12). However, it should be noted that this does not imply the non-existence of the large difference among soft breaking parameters. In fact, it is shown that in some string models, soft breaking terms with different orders of magnitude can be realized.[12),21)] This result arises from the question of whether these soft breaking terms are yielded by tree level effects or loop effects. The soft breaking terms presented in Eqs. (9)~(12) are the values at $M_{\rm Pl}$. The non-universality at high energy scale does not necessarily imply non-universality in the low energy region when the quantum corrections are taken into account. For example, the differences among the soft scalar masses can be diluted by renormalization effects due to the heavy gauginos, as stressed in Ref. 6). Although such a situation is preferred to suppress the FCNC sufficiently, in the following study we consider a kind of non-universality such that the difference among scalar masses is not diluted away in the weak scale region. As shown in Ref. 13) the MSSM with this kind of soft scalar masses has the interesting feature of the gauge coupling unification.

As stressed in the Introduction, non-universal soft supersymmetry breaking is restricted by the low energy phenomena. In particular, it is well known that the non-universality among the scalar masses in the low energy region is strictly constrained by the FCNC phenomena.[4)] For example, the real part of the $K^0$-$\bar{K}^0$ mixing leads to the constraints

$$\begin{aligned}
&\mathrm{Re}\left[\frac{\mathcal{A}^2}{\mathrm{Max}(m_g,\tilde{m}_L,\tilde{m}_R)^2}\left(\frac{\delta\tilde{m}_L{}^2}{\tilde{m}_L{}^2}\right)^2\right]\lesssim 5\times10^{-9}\,\mathrm{GeV}^{-2}\,,\\
&\mathrm{Re}\left[\frac{\mathcal{A}^2}{\mathrm{Max}(m_g,\tilde{m}_L,\tilde{m}_R)^2}\left(\frac{\delta\tilde{m}_L{}^2}{\tilde{m}_L{}^2}\right)\left(\frac{\delta\tilde{m}_R{}^2}{\tilde{m}_R{}^2}\right)\right]\lesssim 10^{-10}\,\mathrm{GeV}^{-2}\,,
\end{aligned}\tag{13}$$

where $m_g$, $\tilde{m}_L$ and $\tilde{m}_R$ are gaugino mass, and averaged left-handed and right-handed squark masses, respectively. The squared mass difference between the left-handed (right-handed) squarks is represented as $\delta\tilde{m}^2_{L(R)}$. $\mathcal{A}^2$ represents the product of the factors coming from the mixing angles of quarks and squarks. For our purpose, we should confine ourselves to the case where the low energy non-universality of the soft scalar masses remains without contradicting the FCNC constraints. As such a simple example, we consider the case where the soft scalar masses belonging to the same flavor are equal but the left-handed squark mass is different from the right-handed up and down squark masses,

$$\tilde{m}^2_{QL}\neq\tilde{m}^2_{UR},\ \tilde{m}^2_{DR}\,.\tag{14}$$

We can easily satisfy the conditions (13) in this case.*) To realize this kind of soft scalar mass, the form of Kähler potential is restricted. As seen from Eq. (9), $Z(\Phi,\bar{\Phi})_{I\bar{J}}$ should be proportional to $\delta_{I\bar{J}}$ in $Q_L{}^I$, $U_R{}^I$ and $D_R{}^I$ sectors, respectively. This

*) This type of soft scalar mass may induce the large contribution to nuclear parity violation.[22)] However, in this paper squark mass and gluino mass are considered to be more than 100 GeV. In this range, nuclear parity violation will not yield a strict constraint.

constraint on the Kähler potential is weaker than that in the universal case and is often satisfied in the superstring models, as suggested in Ref. 12).

Next we review the $CP$ phases in the soft supersymmetry breaking parameters. The soft breaking parameters $A^U_{IJ}$, $A^D_{IJ}$, $B$ and $M_a$ are generally complex and become the new origins of the $CP$ violation which do not exist in the SM. It should be noted that soft scalar masses are real when $Z(\Phi, \bar{\Phi})_{I\bar{J}}$ is diagonal. All of the phases of these parameters are known not to be physically independent. We can extract the physically independent phases from them in the usual way.[16] We take the VEVs of the Higgs fields $H_1$ and $H_2$ to be real by an appropriate redefinition of $H_1$ and $H_2$ so as to make $B\mu$ real. If we note that the complex phases in the gaugino masses are common in Eq. (12), we can make the gaugino mass real by the use of the $R$-transformation and summarize the new $CP$ phases associated with the soft breaking terms in the following form:

$$\phi_{Af_I} = \arg(A^f_{IJ} M^*)\,, \quad \phi_B = \arg(BM^*)\,, \tag{15}$$

where $M^*$ is a complex conjugate of gluino mass $M$. Generally, $A^f_{IJ}$ in Eq. (10) has a different phase structure because of its non-universality and introduces the new $CP$ phase $\phi_{Af_I}$ for each Yukawa coupling. These new $CP$ phases can cause new contributions to the EDMN. In the following study we also assume, for simplicity, that

$$A^U_{IJ} = A_U\,, \quad A^D_{IJ} = A_D\,. \tag{16}$$

This assumption imposes a certain condition on the moduli dependence of Yukawa couplings $h_{IJN}$. The constant Yukawa couplings satisfy the above assumption providing the proportionality of $Z(\Phi, \bar{\Phi})_{I\bar{J}}$ to $\delta_{I\bar{J}}$. If we take $A_U = A_D \equiv A$, the independent $CP$ phases are reduced to $\phi_A$ and $\phi_B$ as in the case of usual universal soft breaking.

## § 3. Electric dipole moment of a neutron

We now proceed to express the formula of the EDM of quarks explicitly keeping the above type of non-universality of the soft scalar masses.*) In order to see the effects of such a non-universality to the EDM of quarks, in this section we only consider the gluino contribution whose Feynmann diagram is shown in Fig. 1. To calculate this diagram, we need an explicit form of a squark mass matrix $M_f{}^2$. For the $f$-type squark ($f = U, D$) it is explicitly written as

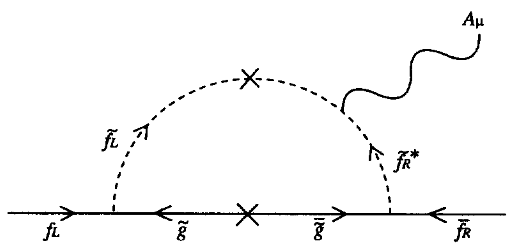


Fig. 1. A Feynman diagram of the gluino contribution to the EDMN.

$$\begin{pmatrix} |m_f|^2 + \tilde{m}^2_{fL} + m_Z{}^2\cos 2\beta(T_f{}^3 - Q_f \sin^2\theta_W) & m_f(A_f + R_f\mu^*) \\ m^*_f(A^*_f + R_f\mu) & |m_f|^2 + \tilde{m}^2_{fR} + m_Z{}^2\cos 2\beta Q_f \sin^2\theta_W \end{pmatrix}, \tag{17}$$

*) In Ref. 17) this calculation has been done. However, the squark masses are replaced by their averaged value at the final stage there.

where $m_f$, $\tilde{m}_{fL}$ and $\tilde{m}_{fR}$ are masses of the $f$-quark, the corresponding left-handed squark and the right-handed squark, respectively. $T_f^3$ is the third component of the weak isospin of the left-handed quark $f$ and $Q_f$ is the electric charge of the quark $f$. $R_f$ is defined by using $\tan\beta \equiv \langle H_2 \rangle / \langle H_1 \rangle$ as

$$R_f = \begin{cases} \cot\beta & (\text{for} \quad f=U), \\ \tan\beta & (\text{for} \quad f=D). \end{cases} \tag{18}$$

Although $M_f^2$ is a $6\times 6$ matrix, here we extract the part corresponding to the first generation to estimate the EDMN. This treatment can be justified because the generation mixing off-diagonal components of $M_f^2$ should be suppressed from the FCNC constraints.

The contribution to the EDM of a quark $f$ from the diagram in Fig. 1 is

$$d_f^g/e = \frac{\alpha_S}{3\pi}\sum_{i=1}^{2}\mathrm{Im}((S^f)_{2i}(S^f)^*_{1i})\frac{Q_f}{m_g}r_{fi}\int_0^1 dx\frac{x(1-x)}{1-x+r_{fi}-x(1-x)s_{fi}}, \tag{19}$$

where $r_{fi} = m_g{}^2/\tilde{m}^2_{fi}$, $s_{fi} = m_f{}^2/\tilde{m}^2_{fi}$, and $m_g$ is the gluino mass. $(S^f)_{ij}$ is the element of the unitary matrix $S^f$ which diagonalizes the squark mass matrix $M_f^2$ as $M^2_{\text{diag}} = (S^f)^\dagger M_f^2 (S^f)$. The eigenvalues of this matrix are represented as $\tilde{m}^2_{fi}$. They are explicitly written as

$$\tilde{m}^2_{f1} = \frac{m_g{}^2}{2}\left[X_f + \sqrt{Y_f{}^2 + Z_f{}^2}\right],$$

$$\tilde{m}^2_{f2} = \frac{m_g{}^2}{2}\left[X_f - \sqrt{Y_f{}^2 + Z_f{}^2}\right]. \tag{20}$$

Here $X_f$, $Y_f$ and $Z_f$ are defined as

$$X_f = \frac{\tilde{m}^2_{Q_L}}{m_g{}^2} + \frac{\tilde{m}^2_{f_R}}{m_g{}^2} + \frac{m_Z{}^2}{m_g{}^2}\cos 2\beta T_f{}^3,$$

$$Y_f = \frac{\tilde{m}^2_{Q_L}}{m_g{}^2} - \frac{\tilde{m}^2_{f_R}}{m_g{}^2} + \frac{m_Z{}^2}{m_g{}^2}\cos 2\beta (T_f{}^3 - 2Q_f \sin^2\theta_W),$$

$$Z_f = \frac{2}{m_g{}^2}|m_f(A_f + R_f \mu^*)|. \tag{21}$$

After evaluating the matrix elements $(S^f)_{ij}$, we obtain the final form of $d_f^g$ as

$$d_f^g/e\sin\gamma_f = \frac{\alpha_S}{3\pi}\frac{Q_f}{m_g}\left[\frac{Z_f{}^2}{Y_f{}^2 + Z_f{}^2}\right]^{1/2}[r_{f2}I(r_{f2}) - r_{f1}I(r_{f1})] \equiv F_f. \tag{22}$$

In the derivation of these formulae, we used the fact $m_f \ll m_g$ for $f=U$, $D$. The function $I(r)$ has the following form,[23]

$$I(r) = \frac{1}{2(1-r)^2}\left[1 + r + \frac{2r\ln r}{1-r}\right]. \tag{23}$$

In this expression, $s_{fi} = m_f{}^2/\tilde{m}^2_{fi}$ is neglected because the quark mass $m_f$ is sufficiently small in comparison to the soft breaking scalar masses $\tilde{m}_{fi}$.*)

In Eq. (22) the angle $\gamma_f$ can be written as**)

$$\tan\gamma_f = \frac{|A_f|\sin\phi_{A_f} + |R_f\mu^*|\sin\phi_B}{|A_f|\cos\phi_{A_f} + |R_f\mu^*|\cos\phi_B}. \tag{24}$$

As is easily seen from this, $\sin\gamma_f$ is of order 1 as long as both of $\tan\phi_{A_f}$ and $\tan\phi_B$ are $O(1)$. This is independent of the value of $|A_f|$ and $|R_f\mu^*|$. In the case that $\phi_{Af}$ and $\phi_B$ are sufficiently small, Eq. (24) is reduced to the usual form:

$$\gamma_f \sim \frac{|A_f|}{|A_f| + |R_f\mu^*|}\phi_{A_f} + \frac{|R_f\mu^*|}{|A_f| + |R_f\mu^*|}\phi_B. \tag{25}$$

This shows that $\gamma_f$ is approximately expressed as $\gamma_f = O(\phi_{A_f})$ or $O(\phi_B)$ depending on whether $|A_f| > |R_f\mu|$ or $|A_f| < |R_f\mu|$.

To reconstruct the EDMN from the dipole moments of the quarks, we can follow the conventional method and use the result of the nonrelativistic quark model

$$d_n = \frac{1}{3}(4d_D - d_U). \tag{26}$$

For our purpose to investigate the effects of the non-universality of soft breaking parameters on the EDMN, however, it will be sufficient to consider the behavior of $d_f^g$.

The parameters relevant to $F_f$ in Eq. (22) are $m_g$, $\tilde{m}_{Q_L}$, $\tilde{m}_{f_R}$, $A_f$, $\tan\beta$ and $\mu$. In these parameters, only $\mu$ cannot be directly related to $m_{3/2}$ in some type of model, as remarked previously.

At first we consider the treatment of the $\tan\beta$ dependence. The radiative symmetry breaking scenario due to the large top Yukawa coupling usually predicts $\tan\beta > 1$, but not a particularly large value.***) Except when $1 \lesssim \tan\beta \lesssim 2$, we can safely, use the approximation $\cos 2\beta \sim -1$. If we consider $\tilde{m}_{Q_L}$ and $\tilde{m}_{f_R}$ to be heavier than $m_Z$, the $\cos 2\beta$ dependence of $X_f$ and $Y_f$ in Eq. (21) is expected to be small, except for the special case with respect to $Y_f$ around the region $\tilde{m}_{Q_L} \sim \tilde{m}_{f_R}$. Then the $\tan\beta$ dependence of $F_f$ is considered to come from $Z_f$ mainly. From this we can summarize the $\tan\beta$ ($R_f$) dependence of $F_f$ as follows. If $|A_f| > |\mu R_f|$, $F_f$ is proportional to $A_f$, and there is no significant $\tan\beta$ dependence. On the other hand, if $|A_f| < |\mu R_f|$, $F_f$ depends on $R_f$ linearly. Taking into account these features, it seems to be convenient to use the combination $|A_f + R_f\mu^*|$ $(= Z_f m_g^2 / 2|m_f|)$. Using this parametrization, we can safely estimate $F_f$ for any value of $\tan\beta$. In the following numerical study, we will put $\tan\beta = 2$. We wish to stress again that this does not yield significant influence on $X_f$ and $Y_f$.

The difference between $F_D$ and $F_U$ comes from the differences in $\tilde{m}_{f_R}$, $A_f$ and $R_f$. In the case $\tilde{m}_{U_R} = \tilde{m}_{D_R}$ and $A_D = A_U$, we can expect $|F_D| > |F_U|$ because of the fact that $\tan\beta > 1$ as found in the RGE study. In this case we can approximate the EDMN as

*) This approximation yields an apparent singularity in $I(r)$ at $r=1$. We need a careful treatment on that point.

**) Here we assume that $\arg(m_f)$ is small enough that the problem of strong $CP$ is not encountered.

***) Here we do not consider the Yukawa unification and therefore the large $\tan\beta$ solution will not be needed.

$d_n^g/e \sim (4/3)F_D \sin\gamma_D$. In the case $\tilde{m}_{U_R} \neq \tilde{m}_{D_R}$ and/or $A_D \neq A_U$, the question of whether $F_D$ or $F_U$ mainly contributes to the EDMN completely depends on the relative size of $A_U$ and $A_D$ and also $\tilde{m}_{U_R}$ and $\tilde{m}_{D_R}$. As is easily seen from Eq. (22), $F_f$ is larger when $A_f$ becomes larger and/or $\tilde{m}_{f_R}$ becomes smaller. In any case we can give the numerical constraints on the soft breaking parameters by comparing the $d_f^g/e$ and the present experimental bound of $d_n^g/e$.

We plot the contour lines of $F_D(\equiv d_D^g/e\sin\gamma_D)$ in the $(\tilde{m}_{D_R}/m_g)$-$(\tilde{m}_{Q_L}/m_g)$ plane in Figs. 2~4. Each graph corresponds to the various values of $m_g$ and $|A_D+\mu^*\tan\beta|$ for $\tan\beta=2$. We can confirm the constraints on the soft breaking parameters from these

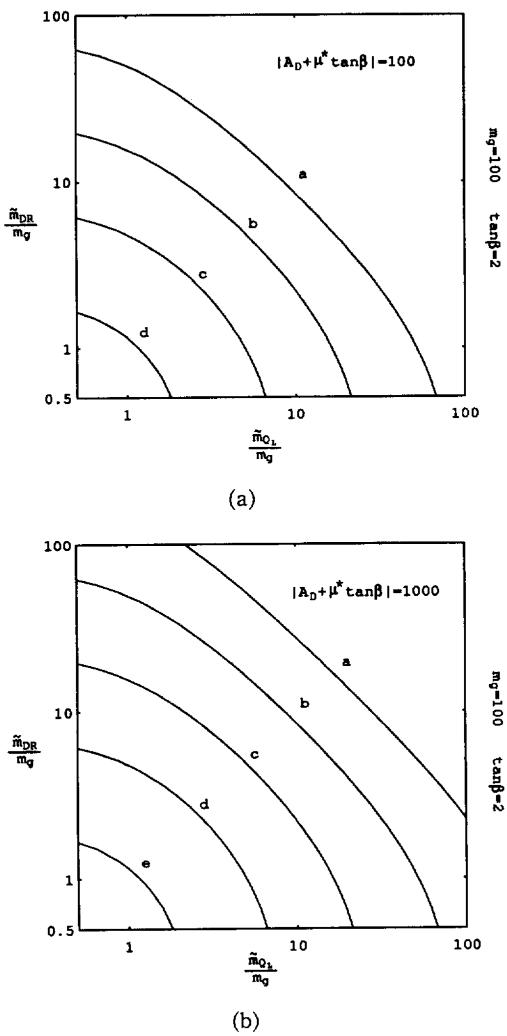


Fig. 2. The contours of $F_D \equiv d_D^g/e\sin\gamma_D$ in the $(\tilde{m}_{D_R}/m_g)$-$(\tilde{m}_{Q_L}/m_g)$ plane at $m_g=100$ GeV and $\tan\beta=2$. $|A_D+\mu^*\tan\beta|$ is chosen as 100 GeV and 1000 GeV in Figs. 2(a) and (b) respectively. Each contour corresponds to a) $10^{-26}$ cm, b) $10^{-25}$ cm, c) $10^{-24}$ cm, d) $10^{-23}$ cm and e) $10^{-22}$ cm.

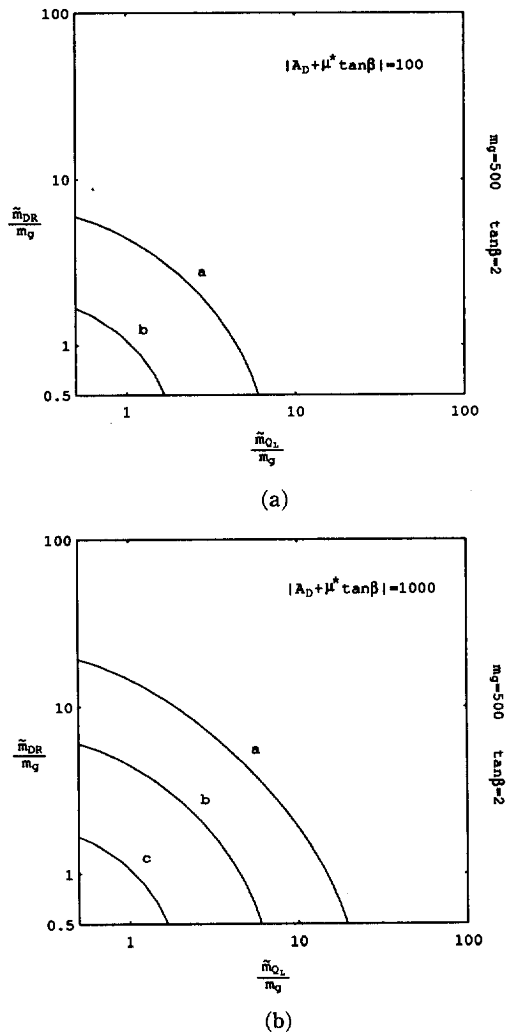


Fig. 3. The contours of $F_D \equiv d_D^g/e\sin\gamma_D$ in the $(\tilde{m}_{D_R}/m_g)$-$(\tilde{m}_{Q_L}/m_g)$ plane at $m_g=500$ GeV and $\tan\beta=2$. The value of $|A_D+\mu^*\tan\beta|$ is the same as in Fig. 2. Each contour represents the same value as in Fig. 2.

figures. For simplicity, we now consider the case $\tilde{m}_{U_R}=\tilde{m}_{D_R}=m_R$ and $A_D=A_U=A$. In this case we can use the approximation $d_n{}^g/e\sim(4/3)F_D\sin\gamma_D$. The constraints usually quoted in the universal soft breaking case are easily obtained by assuming that $\tilde{m}_{Q_L}\sim\tilde{m}_R\sim m_g\sim|A+R_f\mu^*|$. If we take these values as 100 GeV, we have $F_D=10^{-23}$ from Fig. 2, and then $\sin\gamma_D\lesssim 10^{-2}$ is needed to be consistent with the experimental bound. This means that at least $\phi_A$ or $\phi_B$ should be less than $O(10^{-2})$, depending on whether $|A_f|>|R_f\mu|$ or $|A_f|<|R_f\mu|$ as known from Eq. (25).

On the other hand, in the case of non-universal soft breaking, the constraints seem to be rather weakened by its various combined effects. As an example, we consider Fig. 2 where $m_g=|A+R_f\mu^*|=100$ GeV. As long as $\tilde{m}_{Q_L}>1000$ GeV, $d_n{}^g/e\sim 10^{-25}$ can be realized even if $\sin\gamma_D\sim O(1)$ and $\tilde{m}_R=100$ GeV. When non-universality is suitably combined, $F_D$ can be as small as $O(10^{-25})$ in a more natural soft breaking parameter region. This result can also be found in previous works.[15] However, we would like to stress that the suitable combination of non-universality may open an interesting possibility for the suppression of the EDMN. As seen from Eq. (25), for example, if the $|A_f|$ are not the same order as $|R_f\mu^*|$, either $\phi_{A_f}$ or $\phi_B$ will mainly contribute to the EDMN. This may give a new way for arriving at a natural solution of the soft $CP$ phases, as seen in the next section.

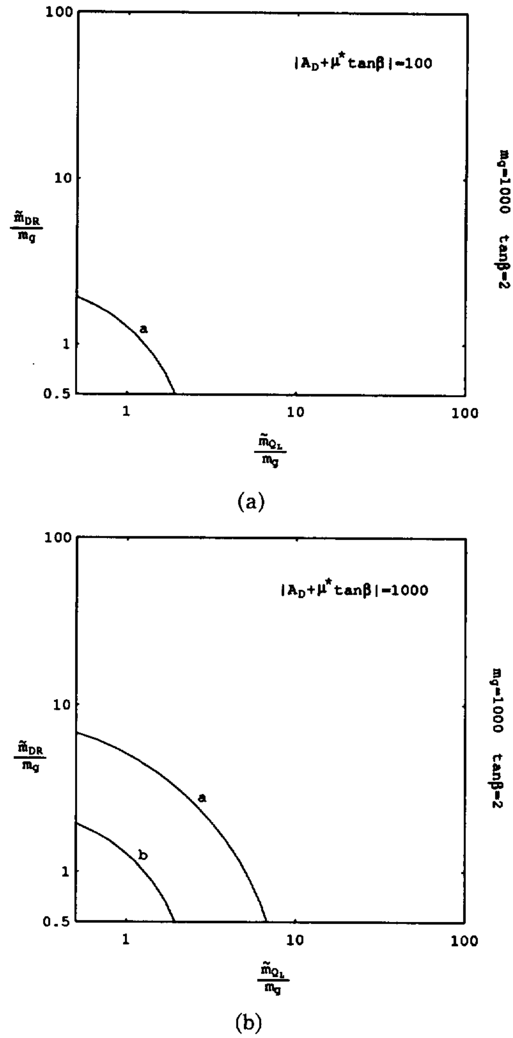


Fig. 4. The contours of $F_D\equiv d_D{}^g/e\sin\gamma_D$ in the $(\tilde{m}_{D_R}/m_g)$-$(\tilde{m}_{Q_L}/m_g)$ plane at $m_g=1000$ GeV and $\tan\beta=2$. The value of $|A_D+\mu^*\tan\beta|$ is the same as in Fig. 2. Each contour represents the same value as in Fig. 2.

Before closing this section, we will summarize the features of the dependence of the EDMN on $\tilde{m}_{Q_L}$, $\tilde{m}_{f_R}$, $m_g$ and $|A_f+R_f\mu^*|$.

(1) Generally large values of $\tilde{m}_{Q_L}$ or $m_{f_R}$ are required to suppress the EDM of the $f$-quark and then the EDMN. However, it should be remarkable that if only one of them is sufficiently heavy, the EDMN can be largely suppressed and satisfy the experimental bound. This is expected from consideration for the squark mass eigenvalues. This feature seems to be favorable for gauge coupling unification in

superstring models. Such non-universal soft scalar masses can shift the gauge coupling unification scale upward from that of the MSSM as shown in Ref. 13).

(2) As $|A_f+R_f\mu^*|$ increases, $d_f^g/e\sin\gamma_f$ increases proportionally. This feature is due to a fact that $|A_f+R_f\mu^*|$ characterizes the left-right mixing of the squark mass matrix. Depending on whether the main contribution comes from $A_f$ or $R_f\mu^*$, the constraints on the soft $CP$ phases $\phi_A$ and $\phi_B$ can be different. Also we should note that a significant cancellation between $A_f$ and $R_f\mu^*$ can occur depending on the relative phase between $A_f$ and $R_f\mu^*$. In such a case, the EDMN can become an extremely small value even if soft breaking parameters are relatively small.

(3) The value of $d_f^g/e\sin\gamma_f$ decreases almost inversely proportionally to the increase of the gluino mass. This feature seems to be irrelevant to the squark masses. It is worth noting that this cannot be read off from Eq. (1) in the range $m_g<\tilde{m}_{q,q^c}$.

(4) Following the usual RGE study, the soft masses satisfying $m_g\gg\tilde{m}_{Q_L}$, $\tilde{m}_{f_R}$ seem to be difficult to realize in the low energy region. The large gluino mass will make the squark masses the same order as the gluino mass through the renormalization effect. This suggests that the soft scalar masses including the RGE effects favor a region satisfying $\tilde{m}_{f_R}/m_g$ and $\tilde{m}_{Q_L}/m_g\gtrsim 1$. It is notable that for a certain value of $m_g$, the EDMN in such a region has a smaller value compared to the one in the region which the RGE effects disfavor.

## § 4. A new suppression mechanism of the EDMN

In this section we study the necessity of the natural suppression of the soft $CP$ phases. As was shown previously, $\sin\gamma_f$ should be small enough not to exceed the experimental bound of the EDMN if all soft supersymmetry breaking parameters are $O(100)$ GeV. This is usualy considered to be equivalent to the condition that both of the soft $CP$ phases $\phi_{A_f}$ and $\phi_B$ are less than $10^{-2}\sim 10^{-3}$, depending on the detailed values of soft breaking parameters. From the viewpoint of the naturalness, such small phases seem to be unexpected in the general soft supersymmetry breaking schemes. We shall propose a natural explanation for this problem for certain types of models.

Recently it was suggested that the phase $\phi_{A_f}$ can be small enough not to contradict the EDMN bound in the models derived from the superstring theories associated with the supersymmetry breaking due to the $F$-terms of a dilaton and moduli. In Ref. 12) it was shown that the dilaton dominated supersymmetry breaking suppresses the phase $\phi_{A_f}$ sufficiently. This is because the phase structures of $A_f$ and $M_a$ in Eqs. (9) and (12) have a certain similarity. As is wellknown, this kind of model induces universal soft breaking terms such as $\tilde{m}_i=-A/\sqrt{3}=M_a/\sqrt{3}=m_{3/2}$. On the other hand, Choi pointed out in Ref. 24) that the various complex phases contributing to $\phi_{A_f}$ are tuned to a value less than $O(10^{-3})$ by the dynamical mechanism based on the Peccei-Quinn symmetry on the dilaton and moduli. In this case any kind of universality of soft breaking terms is not required. This situation is very different from the suppression mechanism based on the dilaton dominance. In both cases, unfortunately, there is no such general suppression mechanisms for the phase $\phi_B$.

The $\mu$ term generally has various origins as shown in Eq. (8). Therefore the structure of $\phi_B$ completely depends on its origin as seen from Eq. (11). It seems very difficult to suppress $\phi_B$ naturally in a way independent of its origin.*) However, the existence of natural suppression mechanisms of $\phi_{A_f}$ can open the new possibility of the sufficient suppression of the EDMN. Instead of finding the suppression mechanism of $\phi_B$, it seems more promising to investigate this new possibility that the EDMN may be sufficiently suppressed even if the phase $\phi_B$ is not small.

For this purpose we will consider the case $|A_f|\gg|R_f\mu|$. We assume that the smallness of $\phi_{A_f}$ is guaranteed by the above mentioned Choi's mechanism because we need the non-universality among $\tilde{m}_i$, $A_f$, $B$ and $M_a$. In such a case the value of $\gamma_f$ can be estimated as

$$\gamma_f \sim \phi_{A_f} + \frac{|R_f\mu^*|}{|A_f|}\sin\phi_B\,. \tag{27}$$

The contribution from $\phi_B$ can be suppressed by a prefactor $|R_f\mu^*|/|A_f|$ even if $\phi_B$ is $O(1)$. The main issue of this scenario is the consistency between the radiative symmetry breaking and the smallness of $|R_f\mu^*|/|A_f|$. Using Eqs. (22) and (27), we can estimate the contribution to the EDMN from $\phi_B$ as

$$d_n{}^g/e\sin\phi_B = \frac{1}{3}\frac{|\mu|}{|A|}(-F_U\cot\beta + 4F_D\tan\beta) \sim \frac{4}{3}\frac{|\mu|}{|A|}F_D\tan\beta\,, \tag{28}$$

where $A_f$ is assumed as $A_U = A_D \equiv A$, for simplicity again. However, this assumption is not essential for our discussion. As discussed in the previous section, we can find $|F_D| \gtrsim |F_U|$ from numerical analysis, and also $\tan\beta > 1$ is generally expected in the radiative symmetry breaking scenario due to the large top Yukawa coupling. Thus the second similarity in Eq. (28) is deduced. The approximate value of $F_D$ can be read off from Figs. 2∼4 for each value of soft breaking parameters. As far as $\tan\beta \sim O(1)$ and the masses of all superpartners are $\sim 100$ GeV, the necessary condition to satisfy the experimental bound of the EDMN is estimated as $|\mu|/|A| < 10^{-2}$. In fact, we can also see this from numerical evaluation.

As suggested by the previous argument on the soft breaking terms, $A$ is expected to be $O(m_{3/2})$, where the magnitude of $m_{3/2}$ is dependent on the supersymmetry breaking mechanism. On the other hand, the scale of $\mu$ depends on its origin as we stressed previously. Then $|\mu|/|A| < 10^{-2} \sim 10^{-3}$ may be naturally realized in a certain class of models. However, the small $|\mu|$ may yield the light chargino and neutralino which conflict with the present experimental bounds. The mass matrices for charginos and neutralinos can be written as

$$\begin{pmatrix} |\mu|e^{-i\phi_B} & \sqrt{2}m_W\sin\beta \\ \sqrt{2}m_W\cos\beta & M_2 \end{pmatrix}. \tag{29}$$

The mass eigenvalues of a squared hermitian mass matrix of charginos are represent-

*) The only known mechanism to suppress $\phi_B$ naturally is to replace the $\mu$-term by a Yukawa coupling of a singlet field and Higgs fields. In this model $\phi_B$ is reduced to the $\phi_A$ type phase and then $\phi_B$ will also be small automatically.[12)]

ed by the following formula,

$$m^2_{\omega_2 \atop \omega_1}=\frac{1}{2}[|\mu|^2+M_2{}^2+2m_W{}^2$$
$$\pm\sqrt{(|\mu|^2-M_2{}^2)^2+4m_W{}^2(|\mu|^2+M_2{}^2+2M_2|\mu|\sin2\beta\cos\phi_B)+4m_W{}^4\cos^22\beta}\,]\,. \tag{30}$$

If we consider the region where both of $|\mu|$ and $M_2$ are sufficiently small, and also $\tan\beta$ is not significantly larger than one, we easily find that the chargino mass does not conflict with the present experimental bound 45 GeV. In order to confirm the existence of the window in the small $|\mu|$ and $M_2$ region which is not excluded from the mass bounds of the charginos, we draw the contours of the smallest eigenvalue in the $(|\mu|, M_2)$ plane. From Fig. 5 we can see that there is a window which satisfies our demand. The constraint from the neutralinos was recently presented by L3.[25] Their analysis shows that there still remains a small window in a region of the $(|\mu|, M_2)$ plane where we are focusing as long as $\tan\beta\sim1$. Based on these considerations, we concentrate our attention on the parameter region satisfying

$$|B|, |A|, \tilde{m}_i \gg |\mu|, M_2 \tag{31}$$

in the low energy region. In this region we practice the RGE study to examine the radiative $SU(2)\times U(1)$ breaking and estimate the top quark mass. As is well known, reducing $|\mu|$ increases the value of $\tan\beta$ significantly. This effect may cancel the smallness of $|\mu|$ and make our scenario less attractive. To avoid such a situation, we take $|B(M_{\rm pl})|$ somewhat larger than other soft parameters. This is because $\tan\beta$ does not depend on $\mu$ directly but depends on $B\mu$. At the tree level analysis, $\beta$ is expressed as[26)]

$$\sin2\beta=\frac{2B\mu}{\tilde{m}^2_{H_1}+\tilde{m}^2_{H_2}+2|\mu|^2}\,. \tag{32}$$

From this, one can find that the value of $m_{H_1}$ influences $\tan\beta$ in the same way as $B$.*) The non-universality of soft scalar masses may be applied not only to squarks and sleptons but also to the Higgs sector.**) By choosing $\tilde{m}_{H_1}(M_{\rm pl})$ smaller than $\tilde{m}_{H_2}(M_{\rm pl})$, we can reduce $\tan\beta$ further. The non-universality between $\tilde{m}_U$ and $\tilde{m}_D$ also affects the running of Higgs masses through RGEs and one can expect effects similar to those mentioned above. However, such effects are indirect and negligible unless

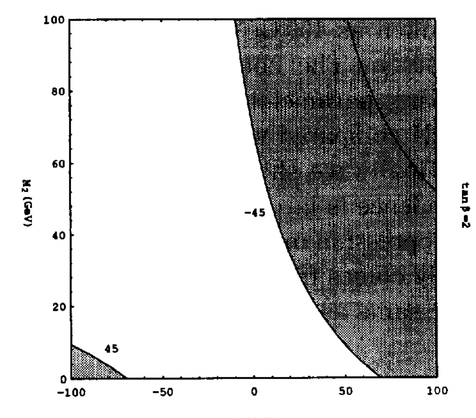


Fig. 5. The contours of mass eigenvalues of the lighest chargino in the $(|\mu|, M_2)$ plane. The shaded region is excluded experimentally.

*) In most cases, $|\tilde{m}^2_{H_2}+|\mu|^2|$ is about $O(10^{-1})(\tilde{m}^2_{H_1}+|\mu|^2)$ so that we need not take into account its effect.

**) From the viewpoint of radiative $SU(2)\times U(1)$ breaking, it is interesting to vary initial values of Higgs masses from other scalar ones.[27)] we assume extremely large squark mass hierarchy, which often causes of color $SU(3)$ breaking. Combining these effects, we can find a suitable parameter region on the basis of RGEs study. For such an illustration, we list a typical relevant set of soft supersymmetry breaking parameters at $m_Z$,

$$|A|\sim 530\,\text{GeV}\,,\quad |\mu|/|A|\sim 1.3\times 10^{-2}\,,\quad \tan\beta\sim 2.6\,,$$

$$\tilde{m}_Q\sim 370\,\text{GeV}\,,\quad \tilde{m}_U\sim 99\,\text{GeV}\,,\quad \tilde{m}_D\sim 510\,\text{GeV}\,,$$

$$m_g\sim 110\,\text{GeV}\,,\quad M_2\sim 41\,\text{GeV}\,.$$

Here we use a handy method to evaluate the contribution from the 1-loop effective potential.[28] For the initial values giving these results, the top quark mass becomes 162 GeV and the masses of the lightest chargino and neutralino are almost on the present experimental bound. The small top quark mass relates to the smallness of $\tan\beta$. Such a top mass seems to be too small when compared with the results of recent experiments.[29]

To avoid such a situation, one can consider a relatively large top Yukawa coupling in order to raise the top quark mass. This prescription may cause the additional problem that the right-handed stop mass becomes too small to satisfy the experimental bounds. This is because of the renormalization effects due to a large Yukawa coupling.

However, we want to stress that this difficulty can again be solved by non-universality. It allows us to set $\tilde{m}_U>\tilde{m}_D$ at $M_{\rm st}$ so as to make $\tilde{m}_U$ large enough at $M_W$.

Using Figs. 2~4, we can estimate $F_D$ as $\sim 7\times 10^{-25}$. Then the constraint on $\phi_B$ seems to disappear for these parameters. This is due to the combined effect of rather large squark mass values and the suppression factor $|\mu|/|A|$. Even if squark masses are taken as 100 GeV, the present value of $|\mu|/|A|$ can reduce the constraint to $\phi_B \sim O(10^{-1})$. The combined effects of the non-universality of soft supersymmetry breaking parameters can weaken the constraints on the soft $CP$ phase from the EDMN compared with those usually considered.

Finally we should comment on the chargino contribution. In this case the $\phi_A$ dependence is largely suppressed due to the small Yukawa couplings even if there is no suppression mechanism of $\phi_A$ and then $\phi_A\sim O(1)$. Thus the chargino contribution mainly comes from the $\phi_B$ dependent part of the $d$-quark EDM. The reason for this is the same as that which is used to derive Eq. (28). This contribution is estimated as[17]

$$\begin{aligned} d_d^c/e\sin\phi_B &= \frac{\alpha_{em}}{4\pi\sin^2\theta_W}\frac{M_2|\mu|\tan\beta}{m_{\omega_2}^2-m_{\omega_1}^2}\frac{m_f}{\tilde{m}_i^2}C \\ &\sim 5.1\times 10^{-25}\left(\frac{1\,\text{TeV}}{\tilde{m}_i}\right)^2\left(\frac{m_f}{10\,\text{MeV}}\right)\frac{M_2|\mu|\tan\beta}{m_{\omega_2}^2-m_{\omega_1}^2}C\ \text{cm}\,, \end{aligned} \tag{33}$$

where $C$ is a numerical factor depending on the value of $t\equiv(m_{\omega_i}^2/m_i^2)$. It is of order 1, assuming a value $1.6<|C|<3.5$ for $10^{-1}>t>10^{-2}$. This chargino contribution is expected to be sufficiently suppressed within the experimental bound for $\phi_B\sim 10^{-1}$ even if $\tilde{m}_i \sim O(100)$ GeV. This is due to the appearance of the additional suppression factor

$$\frac{M_2|\mu|\tan\beta}{m_{\omega_2}^2 - m_{\omega_1}^2}$$

$$= \frac{M_2|\mu|\tan\beta}{\sqrt{(|\mu|^2 - M_2^2)^2 + 4m_W^2(|\mu|^2 + M_2^2 + 2M_2|\mu|\sin 2\beta\cos\phi_B) + 4m_W^4\cos^2 2\beta}}$$

$$\lesssim 6.2\times 10^{-2}\,, \tag{34}$$

for the above parameters.*) Moreover, for the parameters presented above as an example, $\tilde{m}_i$ is larger than 100 GeV by some factor and then $\phi_B \sim O(1)$ will be allowed. Thus the present parameter region can suppress both the gluino and chargino contributions successfully. The large $\phi_B$ models can be consistent with the present bound of the EDMN. In the models with large $\phi_B$, it seems to be interesting to calculate the baryon number asymmetry produced by the anomalous electroweak processes.[18]

Following the above mechanism, the present experimental bound of the EDMN may be reconciled with the MSSM without introducing unnatural assumptions on the *CP* phases of the soft breaking terms. We need only consider a certain kind of non-universal soft supersymmetry breaking. The feature of models with these soft breaking parameters is a relatively light top mass in comparison with the center value of the CDF result and also a light chargino and neutralino close to the experimental bound. Although the favorable parameter region of the present models does not seem to be wide, it may be interesting enough to study it in more detail.

## § 5. Summary

We re-examined the EDMN under the non-universal soft supersymmetry breaking parameters. We discussed the features of relations between the EDMN and soft supersymmetry breaking parameters. They are consistent with the ones already studied in the universal soft breaking framework. The effects of non-universality of these parameters can weaken somewhat the constraints from the EDMN. We also showed that the soft *CP* phase $\phi_B$ whose natural suppression mechanism is not presently known does not yield a large contribution to the EDMN in the parameter space where the radiative symmetry breaking occurs successfully. This may be an interesting non-universal parameter region of the MSSM. In such a parameter region we may not need to require $\phi_B$ to be small. This may be very convenient for the electroweak baryogenesis scenario.

The FCNC severely constrains the soft masses of the squarks and requires degeneracy among the masses of squarks with the same charge at the $m_Z$ scale. On the other hand, the study of the EDMN may provide us with some other knowledge for the squark masses as is shown in this paper. Although these subjects have been extensively studied, if we combine these from various viewpoints, we may get a new useful insight for the overall structure of the soft squark masses. A more precise

*) In Ref. 17) the largeness of the chargino contribution is stressed in the case that $\mu$ and $M_2$ are rather large. Here we are interested in the parameter region where $\mu$ and $M_2$ are small.

study of the soft breaking parameters on the basis of the EDMN seems to be very important still now. Also the improvement of the experimental bound of the EDMN is strongly desired.

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