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FERMILAB-Pub-99/335-T

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November 1999

Submitted to *Physics Letters B*

Operated by Universities Research Association Inc. under Contract No. DE-AC02-76CH03000 with the United States Department of Energy

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FERMILAB-PUB-99/335-T
UICHEP-TH/99-8
hep-ph/9910465
September 1999

Additional two–loop contributions to electric dipole moments in supersymmetric theories

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ABSTRACT

We calculate the two–loop contributions to the electric dipole moments of the electron and the neutron mediated by charged Higgs in a generic supersymmetric theories. The new contributions are originated from the potential CP violation in the trilinear couplings of the charged Higgs bosons to the scalar-top or the scalar-bottom quarks. These couplings did not receive stringent constraints directly. We find observable effects for a sizeable portion of the parameter space related to the third generation scalar-quarks in the minimal supersymmetric standard model.

PACS numbers: 11.30.Er, 14.80.Er

It is widely believed theoretically that supersymmetry will play a significant role at certain fundamental scale that can be probed in the near future. As a result, it is one of the major goal of the new and the future colliders to look for signature of supersymmetry.

The strongest constraint on supersymmetry is from the flavor changing neutral current process such as $K - \bar{K}$ mixing and $\mu \to e\gamma$, or from the flavor neutral CP violating quantities such as the electric dipole moments (EDM's) of the electron and the neutron[1]. These quantities put strong limits on new parameters related to the first two generations of quarks and leptons in any generic supersymmetric theories resulting in the famous SUSY flavor problem and the SUSY CP problem. A solution will require theory to explain why these parameters are so small. One scenario out of these problems is to assume that the first two generations of squarks and sleptons are heavy, while keeping the third generation relatively light in order to produce interesting physical consequence such as low energy electroweak baryogenesis[2]. Such scenario is consistent with the sparticle spectrum resulting from renormalization group evolution starting from an high energy unified theory. It is therefore interesting to look for effective constraint on the new SUSY parameters related to the third generation using the low energy data or collider searches.

In a recent paper[3], it has been shown that by considering the two–loop contributions to the electric dipole moment of the electron and the neutron, it is possible to put non-trivial constraint on the CP violating parameters related to the third generation directly. In fact it was shown later[5] that such constraint actually helps rule out a scenario of the supersymmetric theory in which the CP violating ϵ parameter in the kaon system was proposed to be originated purely from a new source beyond the standard model.[6].

In Ref.[3], only the two–loop diagrams with the second loop mediated by a neutral Higgs boson have been considered. These diagrams are of Barr-Zee type[4]. However, there are scenarios in supersymmetric theories in which other scalar bosons, such as the charged Higgs boson, may be light enough to produce larger effect. In this paper we wish to extend the earlier analysis to include two–loop contributions with the second loop involving the exchange of a charged Higgs boson. It is possible to analyse the two sets of two–loop contributions separately because they are gauge invariant independently. The charged Higgs set is more technically sophisticated and has only been studied in non-supersymmetric theories in Ref.[7].

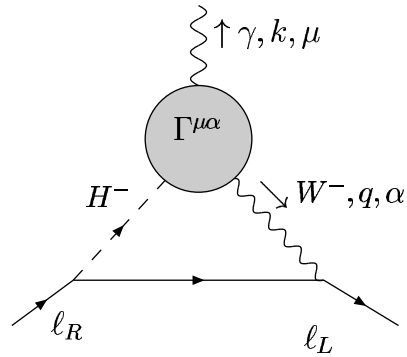

Fig. 1. Structure of the two–loop EDM diagram via Barr-Zee mechanism in supersymmetric theories with charged Higgs exchange in the second loop.

The two–loop diagrams are shown as in Fig. 1, in which both the stop and the sbottom are involved in the first shaded loop. In a generic gauge, there are also other diagrams involving the unphysical Higgs boson $G^{\pm}$. However it is known[8] that in R_ξ *nonlinear* Landau gauge (NLLG), such diagrams will not contribute. For this reason, our calculation is based on NLLG.

We shall first discuss the first loop that creates an $H\gamma W^*$ vertex and then integrate the second loop in the EDM calculation.

The first loop of $H\gamma W^*$ Vertex

The relevant interaction Lagrangian in our study is given by,

$$\mathcal{L} = vH^{-*}\tilde{t}^*\mathcal{E}\tilde{b} + \frac{g_2 \tan\beta m_\ell}{\sqrt{2}M_W}\bar{\nu}_L\ell_R H^{-*} - \tfrac{g_2}{\sqrt{2}}W_\sigma^{-*}(\bar{\ell}_L\gamma^\sigma\nu_L + \tilde{t}^*\mathcal{K}i\overleftrightarrow{\partial^\sigma}\tilde{b}) + \text{H.c.} \tag{1}$$

plus mass terms. Here $\tilde{t}$ and $\tilde{b}$ carry indices $1, 2$ labeling the corresponding mass eigenstates which are superpositions of the bosonic chirality states,

$$\begin{pmatrix} \tilde{b}_L \\ \tilde{b}_R \end{pmatrix} = U_b \begin{pmatrix} \tilde{b}_1 \\ \tilde{b}_2 \end{pmatrix} , \quad \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix} = U_t \begin{pmatrix} \tilde{t}_1 \\ \tilde{t}_2 \end{pmatrix} . \tag{2}$$

These unitary transformation U_t and U_b diagonalize the mass matrices

$$\mathcal{M}^2_{\tilde{q}} = \begin{pmatrix} M^2_{\tilde{q}_L} + m_q^2 + \Delta_{q_L} M_Z^2 \cos 2\beta & -(\hat{m}_q\mu + m_q A_q^*) \\ -(\hat{m}_q\mu^* + m_q A_q) & M^2_{\tilde{q}_R} + m_q^2 + \Delta_{q_R} M_Z^2 \cos 2\beta \end{pmatrix} , \tag{3}$$

with $q = t, b$ with $\hat{m}_b = m_b \tan\beta$, $\hat{m}_t = m_t \cot\beta$, $\Delta_q = T_q^3 - e_q \sin^2\theta_W$ and the $SU(2)$ relation $M_{\tilde{t}_L} = M_{\tilde{b}_L}$. The complex phase can be factored out as

$$U_q = \begin{pmatrix} 1 & 0 \\ 0 & e^{i\delta_q} \end{pmatrix} \begin{pmatrix} c_q & s_q \\ -s_q & c_q \end{pmatrix} , \tag{4}$$

with $c_q \equiv \cos\theta_q$, $s_q \equiv \sin\theta_q$ for $q = b, t$, and

$$\hat{m}_q \mu^* + m_q A_q = |\hat{m}_q \mu^* + m_q A_q| e^{i\delta_q} \ . \tag{5}$$

The charged current coupling matrix $\mathcal{K}$ in this basis of the $\tilde{t}$–$\tilde{b}$ sector is *real*,

$$(\mathcal{K}) = U_t^\dagger \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix} U_b = \begin{pmatrix} c_t c_b & c_t s_b \\ s_t c_b & s_t s_b \end{pmatrix} \ . \tag{6}$$

The tri-linear scalar terms for charged Higgs boson can be written as

$$(\mathcal{E}) = \frac{\sqrt{2}}{v^2} U_t^\dagger \begin{pmatrix} -M_W^2 \sin 2\beta + m_b^2 \tan\beta + m_t^2 \cot\beta & m_b \mu - A_b^* \hat{m}_b \\ m_t \mu^* - A_t \hat{m}_t & 2 m_b m_t / \sin 2\beta \end{pmatrix} U_b \ . \tag{7}$$

The imaginary entries are CP violating,

$$\begin{aligned} \mathrm{Im}\,(\mathcal{E}) \ &= \ \frac{\sqrt{2}}{v^2} \frac{2 m_b m_t}{\sin 2\beta} \sin(\delta_b - \delta_t) \begin{pmatrix} s_t s_b & -s_t c_b \\ -c_t s_b & c_t c_b \end{pmatrix} \\ &\quad + \begin{pmatrix} -\, c_t s_b\, \lambda_b \,+\, s_t c_b\, \lambda_t & c_t c_b\, \lambda_b \,+\, s_t s_b\, \lambda_t \\ -\, s_t s_b\, \lambda_b \,-\, c_t c_b\, \lambda_t & s_t c_b\, \lambda_b \,-\, c_t s_b\, \lambda_t \end{pmatrix} \ . \end{aligned} \tag{8}$$

where

$$\lambda_q = \frac{\sqrt{2} \hat{m}_q}{v^2 \sin\beta \cos\beta} \mathrm{Im}\left(\mu e^{i\delta_q}\right) \ . \tag{9}$$

CP violation is a result of the mismatch in the phase between the $\mathcal{E}$ vertex of the charged Higgs boson and the $\mathcal{K}$ vertex of the W boson. Our convention happens to give rise to real $\mathcal{K}$ and complex $\mathcal{E}$. Note that while there are two stop and two sbottom eigenstates, we actually need only one of them each to give contribution to the EDM of leptons or quarks. Therefore, one can in principle consider the numerically simpler limit in which one of the squark of each flavor is much lighter than the other one. Also, even if the squark mass matrices are accidentally real, the CP violation can still be generated through the weak-basis-charged-Higgs-coupling matrix in Eq.(7). In the first term of Eq. (8), the CP violation is due to the relative phase between stop and sbottom mass matrices, while in the second term CP violation is due to the relative phase between the weak-basis-charged-Higgs-coupling matrix and squark mass matrices. Note, however, that if $\lambda_t = \lambda_b = 0$, the phase of μ may not be zero, however $\sin(\delta_b - \delta_t)$ will be zero and there will be no CP violation in the charged Higgs interaction in our approximation (of ignoring the squark flavor mixing).

We work out the 1-loop amplitude for the process $H(k+q) \to \gamma(k,\mu) + W^*(q,\alpha)$. Besides the 3-prong irreducible diagram, Fig. 2(a),

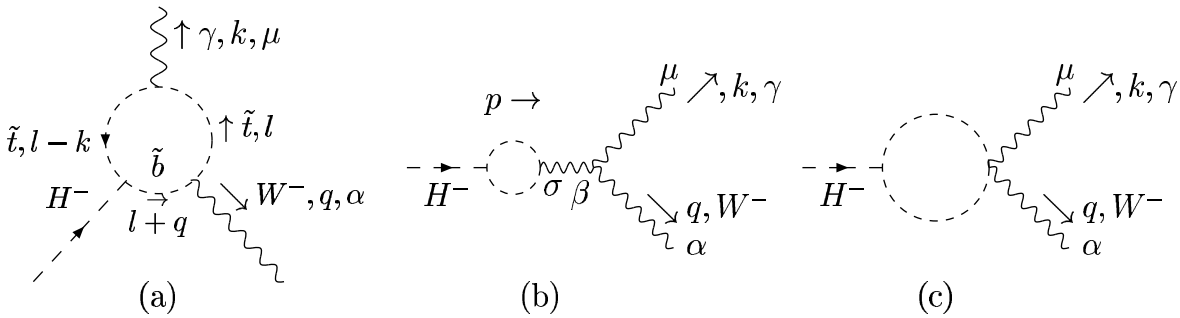

Fig. 2 One loop $H^- \to W^-\gamma$ Feynman diagrams in the *non-linear* Landau gauge.

there is a one-particle reducible H^-–W^- bubble diagram, Fig. 2(b), with the external photon k attached to the outgoing leg W^-. This amplitude is not vanishing even at the limit of the gauge parameter $\xi \to 0$ in our calculation based on the R_ξ *non-linear* Landau gauge (NLLG), where the unphysical Goldstone boson coupling in $GW^-\gamma$ is absent. One notices that the W propagator has the form

$$\mathcal{P}^{\sigma\beta}(p) = -\frac{g^{\sigma\beta} + \frac{p^\sigma p^\beta(\xi-1)}{p^2-\xi M_W^2}}{p^2 - M_W^2} \xrightarrow{\xi\to 0} -\frac{g^{\sigma\beta} - p^\sigma p^\beta/p^2}{p^2 - M_W^2} - \frac{\xi p^\sigma p^\beta}{(p^2)^2} \ . \tag{10}$$

The first term gives vanishing result when it is convoluted with the bubble amplitude of the form Bp_σ. The second term will give finite result even when $\xi \to 0$ as there are singular terms from the tri-gauge-boson coupling in the Non-Linear Landau gauge,

$$\xi(-B/p^2)[(p+k+q/\xi)^\alpha p^\mu + (-k+q-p/\xi)\cdot p g^{\mu\alpha} + (-q-p)^\mu p^\alpha] \longrightarrow B g^{\mu\alpha} \ , \tag{11}$$

where we have extracted the leading constant term in the limit $\xi \to 0$. We need also to include a sea-gull diagram, Fig. 2(c), with the external photon k sticked to the $\tilde{t}$–$\tilde{b}$–W vertex. The photon (k, μ) is attached to the external electric field for the EDM measurement. We only keep the lowest order of k in our calculation.

$$i\Gamma^{\mu\alpha} = -i\sum_{c,d} \frac{\mathcal{E}^*_{c,d} v g_2 \mathcal{K}_{c,d} e}{8\pi^2\sqrt{2}} \int_0^1 \frac{3[e_t(1-y)+e_b y]y(1-y)dy}{(1-y)m_{\tilde{t}_c}^2 + y m_{\tilde{b}_d}^2 - q^2 y(1-y)} (k^\alpha q^\mu - k\cdot q g^{\mu\alpha}) \ . \tag{12}$$

Note that the amplitude in NLLG satisfies the Ward identity $k_\mu \Gamma^{\mu\alpha} = 0$ even when $q^2 \neq M_W^2$.

The second loop and EDM

The final 2-loop EDM amplitude (Fig. 1) of the lepton becomes

$$\begin{aligned} i\mathcal{M}_{H^+} &= i\sum_{c,d} e \frac{\mathrm{Im}(\mathcal{E}^*_{cd})\mathcal{K}_{cd} m_\ell \tan\beta g_2^3 v i}{512\pi^4\sqrt{2}M_W} \int \frac{3(e_t(1-y)+e_b y)y(1-y)dy}{(1-y)m_{\tilde{t}_c}^2 + y m_{\tilde{b}_d}^2 + Q^2 y(1-y)} \\ &\quad \times \frac{dQ^2}{Q^2+M_{H^-}^2} \frac{Q^2}{Q^2+M_W^2} \frac{i}{2} \sigma^{\mu\nu} k_\nu \gamma_5 \ . \end{aligned} \tag{13}$$

Here we integrate Q^2 over $(0,\infty)$, and y over $(0,1)$. In our EDM convention, $\mathcal{M} = -\sigma^{\mu\nu}k_\nu\gamma_5 d_\ell$, we have the contribution to EDM from the charged Higgs sector,

$$\left(\frac{d_\ell}{e}\right)_{H^\pm} = \sum_{c,d}\frac{\mathrm{Im}(\mathcal{E}^*_{c,d})\mathcal{K}_{c,d}m_\ell\tan\beta g_2^2}{512\pi^4\sqrt{2}}\int\frac{3[e_t(1-y)+e_b y]\frac{y(1-y)Q^2dydQ^2}{(Q^2+M^2_{H^-})(Q^2+M^2_W)}}{(1-y)m^2_{\tilde{t}_c}+ym^2_{\tilde{b}_d}+Q^2y(1-y)}$$

$$= -3\sum_{c,d}\frac{\mathrm{Im}(\mathcal{E}^*_{c,d})\mathcal{K}_{c,d}\alpha_{\mathrm{em}}\tan\beta}{256\sqrt{2}\pi^3\sin^2\theta_W}\frac{m_\ell}{M^2_{H^-}}\left[e_t\mathcal{F}\left(\frac{M^2_W}{M^2_{H^-}},\frac{M^2_{\tilde{t}_c}}{M^2_{H^-}},\frac{M^2_{\tilde{b}_d}}{M^2_{H^-}}\right)\right.$$
$$\left.+\,e_b\mathcal{F}\left(\frac{M^2_W}{M^2_{H^-}},\frac{M^2_{\tilde{b}_d}}{M^2_{H^-}},\frac{M^2_{\tilde{t}_c}}{M^2_{H^-}}\right)\right]\ , \quad (14)$$

$$\mathcal{F}(w,\tau,\beta) = \frac{1}{1-w}\left[\mathcal{F}_0(\tau,\beta)-\mathcal{F}_0\left(\frac{\tau}{w},\frac{\beta}{w}\right)\right]\ , \quad (15)$$

$$\mathcal{F}_0(\tau,\beta) = \int_0^1\frac{2y(1-y)^2\ln\frac{y(1-y)}{(1-y)\tau+y\beta}}{\tau(1-y)+\beta y-y(1-y)}dy\ . \quad (16)$$

The above integral is derived based on the relation,

$$\int_0^\infty\frac{xdx}{(x+a)(x+b)(x+c)} = \frac{1}{a-b}\left[\frac{\ln(a/c)}{1-c/a}-\frac{\ln(b/c)}{1-c/b}\right]\ .$$

The value of $\mathcal{F}_0$ is negative and approaches zero as the arguments τ,β approach infinity. Therefore $\mathcal{F}$ is also nagetive. Note that $\sum_{c,d}\mathrm{Im}(\mathcal{E}^*_{c,d})\mathcal{K}_{c,d} = 0$ as expected because if $M_{\tilde{q}_1} = M_{\tilde{q}_2}(q=t,b)$ the mass matrix $\mathcal{M}^2_{\tilde{q}}$ is proportional to the diagonal matrix and the CP violating phase δ_q disappears. In that case, even though the charged Higgs coupling matrix in Eq.(7) may appear to be complex, but the off diagonal phases can easily be absorbed into $\tilde{t}_R$ and $\tilde{b}_R$. Note also that if $M_{\tilde{b}_1} = M_{\tilde{b}_2}$, the CP violation will depends on λ_t only and the contribution from the first term in Eq.(8) will disappear as expected.

Comparison with Pseudoscalar Contribution

The neutral A^0 couples to stop and sbottom in the form

$$\mathcal{L}_{A^0} = vA^0\left[\tilde{t}^\dagger(\xi_t)\tilde{t}+\tilde{b}^\dagger(\xi_b)\tilde{b}\right]+\frac{g_2m_\ell}{2M_W}\tan\beta A^0\bar{\ell}(i\gamma_5)\ell\ . \quad (17)$$

The real part of the matrix (ξ) is CP violating,

$$\mathrm{Re}\,(\xi_q) = \frac{\lambda_q}{\sqrt{2}}\sin 2\theta_q\begin{pmatrix}1 & -\cot 2\theta_q\\ -\cot 2\theta_q & -1\end{pmatrix}\ . \quad (18)$$

We obtain

$$\left(\frac{d_\ell}{e}\right)_{A^0\gamma} = -\frac{3e_\ell\alpha_{\mathrm{em}}}{32\pi^3}\frac{m_\ell}{M^2_{A^0}}\tan\beta\sum_{q=t,b}\frac{\lambda_q}{\sqrt{2}}\sin 2\theta_q e_q^2\left[F\left(\frac{M^2_{\tilde{q}_1}}{M^2_{A^0}}\right)-F\left(\frac{M^2_{\tilde{q}_2}}{M^2_{A^0}}\right)\right]\ , \quad (19)$$

with

$$F(\tau) = \mathcal{F}(0,\tau,\tau) = \mathcal{F}_0(\tau,\tau) = \int_0^1 \frac{y(1-y)}{\tau - y(1-y)} \ln \frac{y(1-y)}{\tau} dy \ . \tag{20}$$

Note that, as expected, if both $M_{\tilde{q}_1} = M_{\tilde{q}_2}$ for $(q = t, b)$ CP violation disappear in our approximation. This explains the cancelation in Eq.(19).

Similar contribution from the Z exchange diagram can be derived based on the interaction,

$$\mathcal{L}_Z = -\frac{g_2}{\cos\theta_W} Z \left(\bar{f}\gamma^\sigma (T_f^3 - e_f \sin^2\theta_W) f + \tilde{t}^* \mathcal{N}_t i \overleftrightarrow{\partial^\sigma} \tilde{t} + \tilde{b}^* \mathcal{N}_b i \overleftrightarrow{\partial^\sigma} \tilde{b} \right) \ . \tag{21}$$

$$\mathcal{N}^q = \begin{pmatrix} T_{q_L}^3 \cos^2\theta_q - e_q \sin^2\theta_W & \frac{1}{2} T_{q_L}^3 \sin 2\theta_q \\ \frac{1}{2} T_{q_L}^3 \sin 2\theta_q & T_{q_L}^3 \sin^2\theta_q - e_q \sin^2\theta_W \end{pmatrix} \ . \tag{22}$$

$$\begin{aligned} \left(\frac{d_\ell}{e}\right)_{A^0 Z} &= -\frac{3\alpha_{\rm em} \tan\beta (\frac{1}{2} T_{\ell_L}^3 - e_\ell \sin^2\theta_W)}{32\pi^3 \sin^2\theta_W \cos^2\theta_W} \frac{m_\ell}{M_{A^0}^2} \\ &\times \sum_{q=t,b} \sum_{a,b} \mathcal{N}_{a,b}^q e_q \ {\rm Re}(\xi_q)_{a,b} \mathcal{F}\left(\frac{M_Z^2}{M_{A^0}^2}, \frac{M_{\tilde{q}_a}^2}{M_{A^0}^2}, \frac{M_{\tilde{q}_b}^2}{M_{A^0}^2} \right) \ . \end{aligned} \tag{23}$$

Although we explicitly use the lepton ℓ for our EDM calculation, the formalism can be easily extended to the quark. The color EDM of quark can also be obtained from Eq. (19) by changing color factors. After including the renormalization group evolution factor, we use the simple nonrelativistic quark model to estimate the neutron EDM[3].

In Fig. 3 we show contributions to the EDM of the electron and the neutron. For EDMs of the electron and the quarks, contributions from the charged Higgs boson (plus W boson) and from the pseudoscalar neutral Higgs boson (plus either photon and Z boson) are all shown for comparison. For the EDM of the neutron, there is additional contribution due to chromoelectric dipole moments of quarks from two–loop diagrams with the pseudoscalar neutral Higgs boson plus gluons. For simplicity, we assume the μ parameter is real and $M_{\tilde{q}_L}^2 = M_{\tilde{q}_R}^2 = 0.6$ TeV, $A_t = A_b = i(1$ TeV) as in Ref.[3]. In Figs. 3a and 3b, we display the $\tan\beta$ dependence of each contribution to the EDM for the electron and the neutron respectively for both $M_A = 150$ GeV and 300 GeV ($\mu = 1$ TeV). In Figs. 3c and 3d we display the μ dependence for the same choice of parameters with $\tan\beta = 20$. In the minimal supersymmetric Standard Model (MSSM), the masses of the charged Higgs boson and the pseudoscalar boson are related by

$$M_{H^\pm}^2 = M_{A^0}^2 + M_W^2 + \epsilon \ ,$$

where $\epsilon = 3g_2^2 m_t^4 \ln(1 + (\tilde{m}^2/m_t^2))/(8\pi^2 M_W^2)$ is the the quantum correction[9]. We have used the tree level result in plotting Figs. 3a–d. However, if quantum corrections are included, $M_{H^\pm}$ and M_A become less closely related. Therefore in Figs. 3e and 3f we display

the dependence on $M_{H^\pm}$ and the dependence on M_A for their respective contributions. Overall, the EDMs are not very sensitive functions of $M_{H^\pm}$ and M_A. For the EDM of the neutron, the chromoelectric dipole moment contribution (due to pseudoscalar boson exchanges) still dominates over other sources. For the EDM of the electron, the contribution of the charged Higgs boson exchange is in general small than that of the pseudoscalar exchange by about an order of magnitude. The A^0Z exchange diagrams, which have not been included in Ref.[3] contribute even smaller.

Conclusion

We have calculated the charged Higgs related two–loop contributions to the electric dipole moments of electron and neutron. For numerical simplicity, we ignore the generational mixing between squarks, however, it should be quite easy to incorporate if necessary. We find that the charged Higgs contribution are generally smaller than the neutral Higgs contribution calculated earlier[3] in MSSM without the squark flavor mixing which we assume. However, one can imagine that in theories beyond MSSM, the two contributions may involve independent parameters that should be constrained separately. It is straightforward to generalize our analytic result to accommodate theories beyond MSSM or to include squark flavor mixing.

While we were preparing this manuscript, we became aware of a preprint[10] which aims to calculate the same contribution. Our analytic results in Eq. (8) and Eq. (14) differ from those in Ref.[10]. In particular our amplitude in Eq. (12) has the gauge invariant form, while the corresponding one in Ref.[10] is not. This difference does not affect the order of magnitude of the numerical results very much. However, the result for the charged Higgs contribution to EDM of electron in Ref.[10] starts decreasing for large $\tan\beta(> 35)$, but not in our gauge invaraiant result as shown in Fig. 3a.

Acknowledgment

We thank A. Pilaftsis for a discussion. DC and WFC are supported partially by a grant from National Science Council of Republic of China and WYK partially by a grant from US Department of Energy. DC and WYK wish to thank KIAS (Seoul) and Seoul National University for hospitality during the SSS-99 workshop when this work was initiated.

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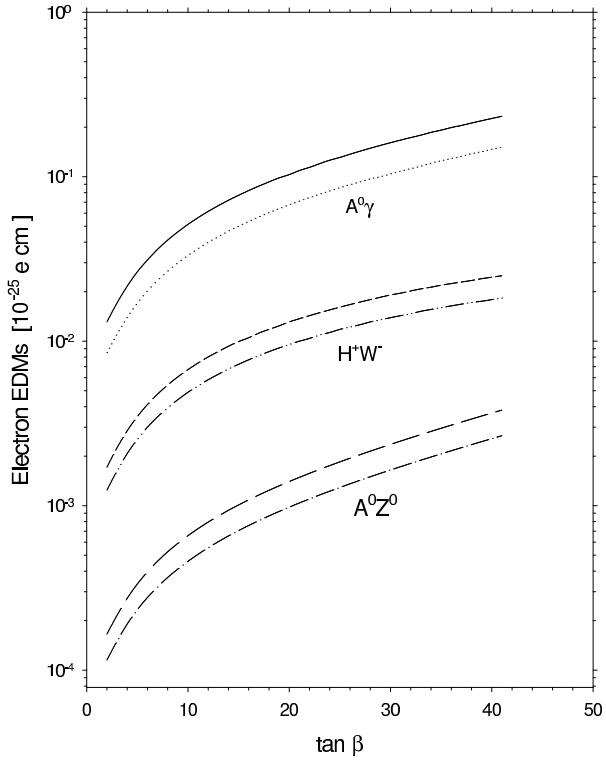

Fig. 3a

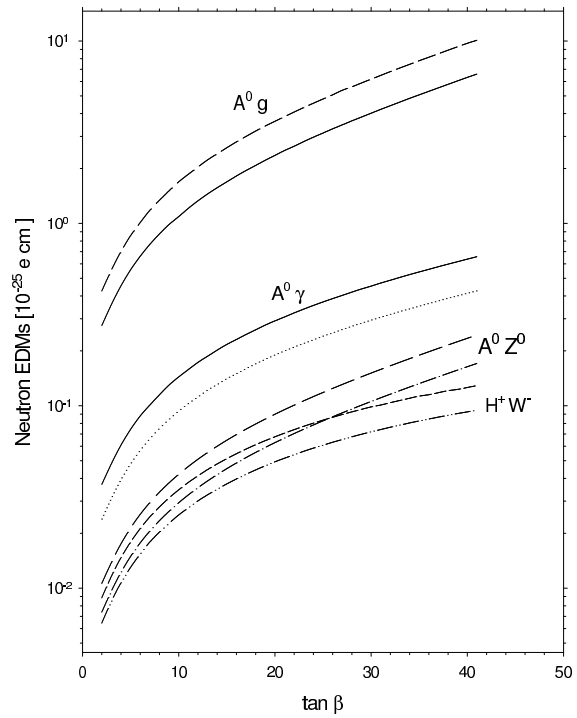

Fig. 3b

Fig. 3a Numerical estimates of various contributions to the EDM of the electron as a function of $\tan\beta$, for the case $M^2_{\tilde{q}_L} = M^2_{\tilde{q}_R} = 0.6$ TeV, $M_A = 150$ GeV and 300 GeV, real $\mu = 1$ TeV and imaginary $A_t = A_b = i(1\ \text{TeV})$.

Fig. 3b Numerical estimates of various contributions to the EDM of the neutron as a function of $\tan\beta$, for the case $M^2_{\tilde{q}_L} = M^2_{\tilde{q}_R} = 0.6$ TeV, $M_A = 150$ GeV and 300 GeV, real $\mu = 1$ TeV and imaginary $A_t = A_b = i(1\ \text{TeV})$.

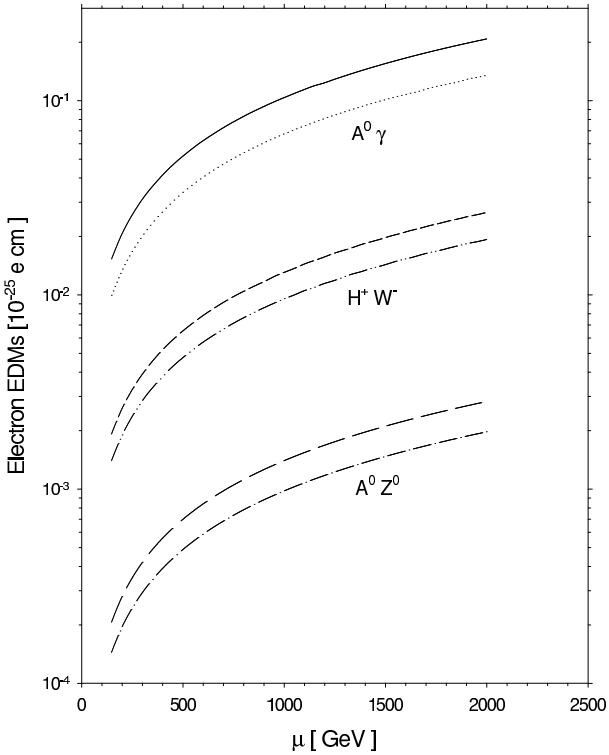

Fig. 3c

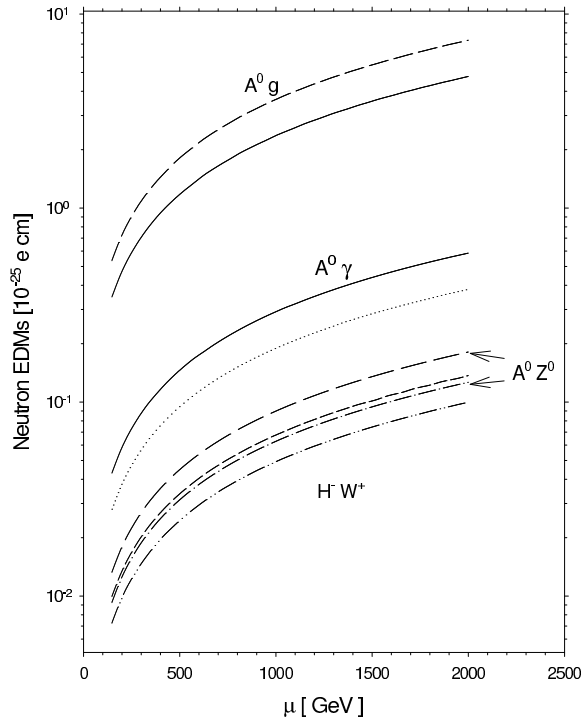

Fig. 3d

Fig. 3c Numerical estimates of various contributions to the EDM of the electron as a function of μ, for the case $\tan\beta = 20$, $M^2_{\tilde{q}_L} = M^2_{\tilde{q}_R} = 0.6$ TeV, $M_A = 150$ GeV and 300 GeV, real μ and imaginary $A_t = A_b = i(1\ \text{TeV})$.

Fig. 3d Numerical estimates of various contributions to the EDM of the neutron as a function of μ, for the case $\tan\beta = 20$, $M^2_{\tilde{q}_L} = M^2_{\tilde{q}_R} = 0.6$ TeV, $M_A = 150$ GeV and 300 GeV, real μ and imaginary $A_t = A_b = i(1\ \text{TeV})$.

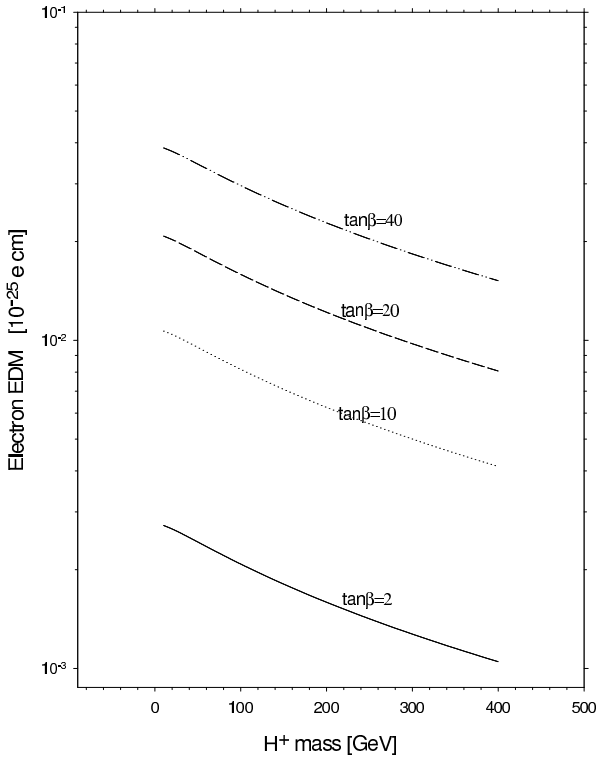

Fig. 3e

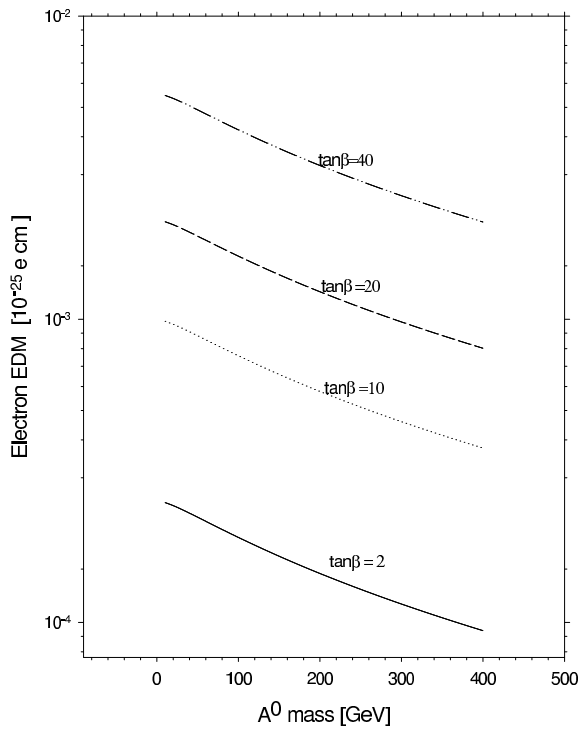

Fig. 3f

Fig. 3e Numerical estimates of H^+ contributions to the EDM of the electron as a function of $M_{H^\pm}$, for the case $M^2_{\tilde{q}_L} = M^2_{\tilde{q}_R} = 0.6$ TeV, real $\mu = 1$ TeV and imaginary $A_t = A_b = i(1 \text{ TeV})$.

Fig. 3f Numerical estimates of A^0 contributions to the EDM of the electron as a function of M_A, for the case $M^2_{\tilde{q}_L} = M^2_{\tilde{q}_R} = 0.6$ TeV, real $\mu = 1$ TeV and imaginary $A_t = A_b = i(1 \text{ TeV})$.