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Higgs-Boson Two-Loop Contributions to Electric Dipole Moments in the MSSM

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ABSTRACT

The complete set of Higgs-boson two-loop contributions to electric dipole moments of the electron and neutron is calculated in the minimal supersymmetric standard model. The electric dipole moments are induced by CP-violating trilinear couplings of the ‘CP-odd’ and charged Higgs bosons to the scalar top and bottom quarks. Numerical estimates of the individual two-loop contributions to electric dipole moments are given.

Supersymmetric (SUSY) theories, including the minimal supersymmetric standard model (MSSM), face the difficulty of explaining naturally the apparent absence of electric dipole moments (EDMs) of the neutron and electron [1, 2, 3, 4, 5]. Several suggestions have been made to suppress the SUSY contributions to electron and neutron EDMs, at a level just below their present experimental 2σ upper bounds: $|d_e| < 0.5 \times 10^{-26} e \text{ cm}$ and $|d_n| < 1.12 \times 10^{-25} e \text{ cm}$ [6]. Apart from the obvious choice of suppressing the new CP-violating phases of the theory to the 10^{-3} level [1, 2], a more phenomenologically appealing possibility is to make the first two generations of scalar fermions as heavy as few TeV, but keep the soft-breaking mass parameters of the third generation relatively small, e.g. 0.5–0.7 TeV [3]. An interesting alternative is to arrange for partial cancellations among the different EDM contributions [4], within the framework of superstring-derived models [7].

Here, we shall focus our interest on studying additional two-loop contributions to electron and neutron EDMs in a SUSY scenario, in which the first two generations are rather heavy, e.g. of order few TeV, whereas the third generation is relatively light below the TeV scale [8]. In such a theoretical framework, the leading effects on the EDMs arise from CP-violating trilinear interactions related to the scalar top and bottom quarks through the three-gluon operator [2] and through the coupling of the ‘CP-odd’ Higgs boson, a , to the gauge bosons [5], as shown in Fig. 1. Despite their similarity to the graphs due to Barr and Zee [9], our reasoning of considering the two-loop contributions of Fig. 1 is completely different from [9], as the third-generation scalar quarks can have a significant impact by themselves on the electron and neutron EDMs, independently of the chirality-unsuppressed one-loop contributions. In our analysis, we shall assume rather heavy gluino masses $m_{\tilde{g}}$, e.g. $m_{\tilde{g}} > 0.5 \text{ TeV}$, such that their effect through the three-gluon operator, which scales as $1/m_{\tilde{g}}^3$ [2], will be much smaller than the present experimental upper bound [4].

In this paper, we shall complement the analysis of Ref. [5], and consider the complete set of two-loop EDM graphs, shown in Fig. 1, including those due to CP-violating $aA_\mu Z_\lambda$ and $H^+ A_\mu W_\lambda^-$ couplings. At this point, it is worth stressing that the EDM constraints we shall study here will have important consequences on Higgs-sector CP violation within the MSSM found recently [11] and on related phenomena in B -meson decays, dark-matter searches and collider experiments [12].

Our starting point is the scalar top and bottom mass matrices, which may conve-

*Analogous EDM contributions induced by charged-Higgs-boson two-loop graphs were studied in [10] within the context of three-Higgs-doublet models.

niently be expressed, in the weak basis $(\tilde{q}_L, \tilde{q}_R)$, as follows:

$$\tilde{\mathcal{M}}_q^2 = \begin{pmatrix} \tilde{M}_Q^2 + m_q^2 + \cos 2\beta M_Z^2 (T_z^q - Q_q \sin^2 \theta_w) & m_q(A_q^* - \mu R_q) \\ m_q(A_q - \mu^* R_q) & \tilde{M}_q^2 + m_q^2 + \cos 2\beta M_Z^2 Q_q \sin^2 \theta_w \end{pmatrix}, \quad (1)$$

with $q = t, b$, $Q_t = 2/3$, $Q_b = -1/3$, $T_z^t = -T_z^b = 1/2$, $R_b = \tan \beta = v_2/v_1$, $R_t = \cot \beta$, and $\sin \theta_w = (1 - M_W^2/M_Z^2)^{1/2}$. Moreover, \tilde{M}_Q^2 and \tilde{M}_q^2 are soft-SUSY-breaking masses for the left-handed and right-handed scalar top and bottom quarks. The matrix $\tilde{\mathcal{M}}_q^2$ may be diagonalized through a unitary rotation, which relates the weak $(\tilde{q}_L, \tilde{q}_R)$ to the mass eigenstates $(\tilde{q}_1, \tilde{q}_2)$:

$$\begin{pmatrix} \tilde{q}_L \\ \tilde{q}_R \end{pmatrix} = \begin{pmatrix} 1 & 0 \\ 0 & e^{i\delta_q} \end{pmatrix} \begin{pmatrix} \cos \theta_q & \sin \theta_q \\ -\sin \theta_q & \cos \theta_q \end{pmatrix} \begin{pmatrix} \tilde{q}_1 \\ \tilde{q}_2 \end{pmatrix}, \quad (2)$$

where $\delta_q = \arg(A_q - R_q \mu^*)$ and θ_q are mixing angles determined by

$$\begin{aligned} \cos \theta_q &= \frac{m_q |A_q - R_q \mu^*|}{\sqrt{m_q^2 |A_q - R_q \mu^*|^2 + [(\tilde{\mathcal{M}}_q^2)_{LL} - M_{\tilde{q}_1}^2]^2}}, \\ \sin \theta_q &= \frac{|(\tilde{\mathcal{M}}_q^2)_{LL} - M_{\tilde{q}_1}^2|}{\sqrt{m_q^2 |A_q - R_q \mu^*|^2 + [(\tilde{\mathcal{M}}_q^2)_{LL} - M_{\tilde{q}_1}^2]^2}}. \end{aligned} \quad (3)$$

In Eq. (3), the quantity $(\tilde{\mathcal{M}}_q^2)_{LL}$ is the $\{11\}$ -matrix element of the scalar quark mass matrix $\tilde{\mathcal{M}}_q^2$. In addition, the mass eigenvalues of $\tilde{\mathcal{M}}_q^2$ are given by

$$\begin{aligned} M_{\tilde{q}_1(\tilde{q}_2)}^2 &= \frac{1}{2} \left\{ \tilde{M}_Q^2 + \tilde{M}_q^2 + 2m_q^2 + T_z^q \cos 2\beta M_Z^2 \right. \\ &\quad \left. - (+) \sqrt{[\tilde{M}_Q^2 - \tilde{M}_q^2 + \cos 2\beta M_Z^2 (T_z^q - 2Q_q \sin^2 \theta_w)]^2 + 4m_q^2 |A_q^* - \mu R_q|^2} \right\}. \end{aligned} \quad (4)$$

As usual, we consider the convention in which $\arg \mu$ and $\arg A_{t,b}$ are the only physical SUSY CP-violating angles in the MSSM. Then, CP violation originates from the interaction Lagrangian

$$\mathcal{L}_{\text{CP}} = a \sum_{\tilde{q}=\tilde{t},\tilde{b}} \sum_{i,j=1,2} (\tilde{q}_i^* \text{Re} \Gamma^{a\tilde{q}_i^* \tilde{q}_j} \tilde{q}_j) + (H^+ \sum_{i,j=1,2} \tilde{t}_i^* i \text{Im} \Gamma^{H^+ \tilde{t}_i^* \tilde{b}_j} \tilde{b}_j + \text{H.c.}). \quad (5)$$

The real and imaginary parts of the couplings $a\tilde{q}^* \tilde{q}$ and $H^+ \tilde{t}_i^* \tilde{b}_j$, which are denoted by $\Gamma^{a\tilde{q}_i^* \tilde{q}_j}$ and $\Gamma^{H^+ \tilde{t}_i^* \tilde{b}_j}$, respectively, are given by

$$\text{Re} \Gamma^{a\tilde{q}_i^* \tilde{q}_j} = v \xi_q \begin{pmatrix} 1 & -\cot 2\theta_q \\ -\cot 2\theta_q & -1 \end{pmatrix}, \quad (6)$$

$$\begin{aligned} \text{Im } \Gamma^{H^+ \tilde{t}_i^* \tilde{b}_j} &= \frac{v}{\sqrt{2}} \begin{pmatrix} -\frac{\cos \theta_t}{\cos \theta_b} \xi_b + \frac{\cos \theta_b}{\cos \theta_t} \xi_t & \frac{\cos \theta_t}{\sin \theta_b} \xi_b + \frac{\sin \theta_b}{\cos \theta_t} \xi_t \\ -\frac{\sin \theta_t}{\cos \theta_b} \xi_b - \frac{\cos \theta_b}{\sin \theta_t} \xi_t & \frac{\sin \theta_t}{\sin \theta_b} \xi_b - \frac{\sin \theta_b}{\sin \theta_t} \xi_t \end{pmatrix} \\ &+ \frac{\sqrt{2} m_b m_t \sin(\delta_b - \delta_t)}{\sin \beta \cos \beta v^2} \begin{pmatrix} \sin \theta_t \sin \theta_b & -\sin \theta_t \cos \theta_b \\ -\cos \theta_t \sin \theta_b & \cos \theta_t \cos \theta_b \end{pmatrix}, \end{aligned} \quad (7)$$

where $v = \sqrt{v_1^2 + v_2^2} = 2g_w/M_W$, and ξ_t and ξ_b are the CP-violating quantities

$$\xi_q = R_q \frac{\sin 2\theta_q m_q \text{Im}(\mu e^{i\delta_q})}{\sin \beta \cos \beta v^2}. \quad (8)$$

We shall first calculate the couplings $aA_\mu(k)A_\lambda(q)$, $ag_\mu(k)g_\lambda(q)$, $aA_\mu(k)Z_\lambda(q)$ and $H^-A_\mu(k)W_\lambda^+(q)$, which are induced by \tilde{t} - and \tilde{b} -mediated one-loop graphs shown in Fig. 1. We adopt the Feynman-'t Hooft gauge, and neglect Feynman diagrams that lead to suppressed EDM contributions proportional to m_f^2 . With the convention that the momenta q and k flow into the vertex, the analytic result of the one-loop couplings is found to have the gauge-invariant form

$$i\Gamma_{\mu\lambda}^V(k, q) = iA^V(q^2) [(q \cdot k) g_{\mu\lambda} - q_\mu k_\lambda], \quad (9)$$

where the superscript V denotes the gauge boson propagating in the loop, and

$$A^\gamma(q^2) = \sum_{q=t,b} \frac{N_c e^2 Q_q^2 v}{8\pi^2} \sum_{i=1,2} (-1)^{i+1} \xi_q \int_0^1 dx \frac{x(1-x)}{M_{\tilde{q}_i}^2 - q^2 x(1-x)}, \quad (10)$$

$$\begin{aligned} A^Z(q^2) &= \sum_{q=t,b} \frac{N_c g_w e Q_q}{16 \cos \theta_w \pi^2} \sum_{i,j=1,2} \text{Re} \Gamma^{a\tilde{q}_i^* \tilde{q}_j} K_{ij}^q \\ &\times \int_0^1 dx \frac{x(1-x)}{M_{\tilde{q}_i}^2 x + M_{\tilde{q}_j}^2 (1-x) - q^2 x(1-x)}, \end{aligned} \quad (11)$$

$$A^W(q^2) = \frac{N_c g_w e}{8\sqrt{2}\pi^2} \sum_{i,j=1,2} (\Gamma^{H^+ \tilde{t}_i^* \tilde{b}_j})^* K_{ij}^{tb} \int_0^1 dx \frac{x(1-x) [Q_t x + Q_b (1-x)]}{M_{\tilde{t}_i}^2 x + M_{\tilde{b}_j}^2 (1-x) - q^2 x(1-x)}. \quad (12)$$

The corresponding one-loop form factor A^g may be obtained by replacing the colour factor $N_c = 3$ by $1/2$, and Q_q^2 by 1 in Eq. (10). In Eqs. (11) and (12), K^q ($q = t, b$) and K^{tb} are (2×2) -non-unitary matrices describing the mixing in the $Z\tilde{q}_i^* \tilde{q}_j$ and $W^+ \tilde{t}_i^* \tilde{b}_j$ sectors, respectively:

$$K^q = \begin{pmatrix} 2T_z^q \cos^2 \theta_q - 2Q_q \sin^2 \theta_w & T_z^q \sin 2\theta_q \\ T_z^q \sin 2\theta_q & 2T_z^q \sin^2 \theta_q - 2Q_q \sin^2 \theta_w \end{pmatrix}, \quad (13)$$

$$K^{tb} = \begin{pmatrix} \cos \theta_t \cos \theta_b & \cos \theta_t \sin \theta_b \\ \sin \theta_t \cos \theta_b & \sin \theta_t \sin \theta_b \end{pmatrix}. \quad (14)$$

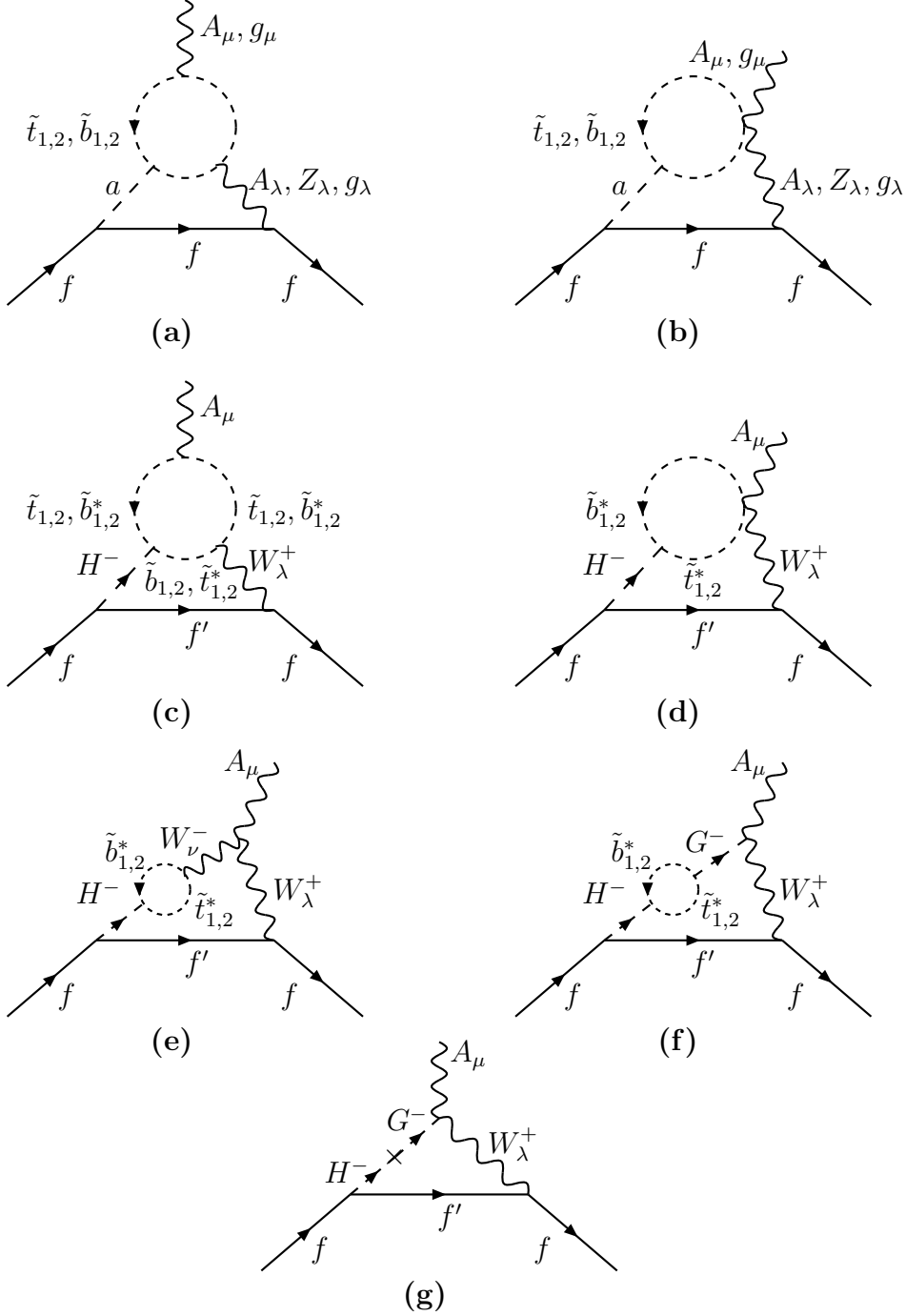


Figure 1: Higgs-boson two-loop contributions to EDM and CEDM of a fermion in the MSSM in the Feynman-'t Hooft gauge (mirror graphs are not shown); f' represents the conjugate fermion of f under T_z^f .

At this stage, we should remark that there are also contributions originating from chargino

loops. However, these contributions are proportional to $\arg\mu$, on which strict constraints exist from one-loop graphs that contribute to the electron EDM, and may therefore be neglected. Finally, we should notice that the one-loop H^-AW^+ -coupling receives its gauge-invariant form of Eq. (9) in the Feynman-'t Hooft gauge, after including the H^-W^+ and H^-G^+ wave functions and the respective tadpole contribution related to the H^-G^+ transition, as shown in Figs. 1(e)–(g).

It is now straightforward to compute the individual contributions to the EDM of a fermion that come from quantum corrections involving γ , Z and W^\pm bosons in the loop. These individual EDM contributions shown in Fig. 1 may conveniently be cast into the form:

$$\left(\frac{d_f}{e}\right)^\gamma = Q_f \frac{N_c \alpha \tan \beta m_f}{32\pi^3 M_a^2} \sum_{q=t,b} \xi_q Q_q^2 \left[F\left(\frac{M_{\tilde{q}_1}^2}{M_a^2}\right) - F\left(\frac{M_{\tilde{q}_2}^2}{M_a^2}\right) \right], \quad (15)$$

$$\begin{aligned} \left(\frac{d_f}{e}\right)^Z &= -(T_z^f - 2Q_f \sin^2 \theta_w) \frac{N_c \alpha_w \tan \beta m_f}{64 \cos^2 \theta_w \pi^3 M_a^2} \\ &\times \sum_{q=t,b} \sum_{i,j=1,2} \frac{1}{v} \operatorname{Re} \Gamma^{a\tilde{q}_i^* \tilde{q}_j} K_{ij}^q Q_q G\left(\frac{M_Z^2}{M_a^2}, \frac{M_{\tilde{q}_i}^2}{M_a^2}, \frac{M_{\tilde{q}_j}^2}{M_a^2}\right), \end{aligned} \quad (16)$$

$$\begin{aligned} \left(\frac{d_f}{e}\right)^W &= \frac{N_c \alpha_w \tan \beta m_f}{128\sqrt{2}\pi^3 M_{H^+}^2} \sum_{i,j=1,2} \frac{1}{v} \operatorname{Im} \Gamma^{H^+ \tilde{t}_i^* \tilde{b}_j} K_{ij}^{tb} \\ &\times \left[Q_t G\left(\frac{M_W^2}{M_{H^+}^2}, \frac{M_{\tilde{t}_i}^2}{M_{H^+}^2}, \frac{M_{\tilde{b}_j}^2}{M_{H^+}^2}\right) + Q_b G\left(\frac{M_W^2}{M_{H^+}^2}, \frac{M_{\tilde{b}_i}^2}{M_{H^+}^2}, \frac{M_{\tilde{t}_j}^2}{M_{H^+}^2}\right) \right], \end{aligned} \quad (17)$$

where $\alpha = e^2/(4\pi)$ and $\alpha_w = g_w^2/(4\pi)$ are the electromagnetic and weak fine structure constants, respectively, and $F(z)$ and $G(a, b, c)$ are two-loop functions given by

$$F(z) = \int_0^1 dx \frac{x(1-x)}{z-x(1-x)} \ln \left[\frac{x(1-x)}{z} \right], \quad (18)$$

$$\begin{aligned} G(a, b, c) &= \int_0^1 dx x \left[\frac{ax(1-x) \ln a}{(a-1)[ax(1-x) - bx - c(1-x)]} \right. \\ &+ \left. \frac{x(1-x)[bx + c(1-x)]}{[ax(1-x) - bx - c(1-x)][x(1-x) - bx - c(1-x)]} \ln \left(\frac{bx + c(1-x)}{x(1-x)} \right) \right]. \end{aligned} \quad (19)$$

Note that $2G(0, a, a) = -F(a)$. Among the EDM terms given by Eqs. (15)–(17), the fermion EDM induced by the photon-exchange graphs $(d_f/e)^\gamma$ represents the dominant contribution [5]. For completeness, we give the two-loop contribution to the CEDM of a coloured fermion [5]

$$\left(\frac{d_f^C}{e}\right)^g = \frac{\alpha_s \tan \beta m_f}{64\pi^3 M_a^2} \sum_{q=t,b} \xi_q \left[F\left(\frac{M_{\tilde{q}_1}^2}{M_a^2}\right) - F\left(\frac{M_{\tilde{q}_2}^2}{M_a^2}\right) \right]. \quad (20)$$

We can now estimate the neutron EDM d_n induced by d_u and d_d in the valence quark model, by including QCD renormalization effects [5], i.e.

$$\frac{d_n}{e} \approx \left(\frac{g_s(M_Z)}{g_s(\Lambda)} \right)^{32/23} \left[\frac{4}{3} \left(\frac{d_d}{e} \right)_\Lambda - \frac{1}{3} \left(\frac{d_u}{e} \right)_\Lambda \right]. \quad (21)$$

The u - and d -quark masses occurring in Eq. (15) are running masses evaluated at the low-energy hadronic scale Λ . In Eq. (21), we have assumed that the renormalization-group running factors of the strong coupling constant g_s from m_b to Λ is almost of order 1. To be specific, we consider the values: $m_u(\Lambda) = 7$ MeV, $m_d(\Lambda) = 10$ MeV, $\alpha_s(M_Z) = 0.12$, and $g_s(\Lambda)/(4\pi) = 1/\sqrt{6}$ [2]. Likewise, the light-quark CEDMs d_u^C and d_d^C lead to a neutron EDM

$$\frac{d_n^C}{e} \approx \left(\frac{g_s(M_Z)}{g_s(\Lambda)} \right)^{74/23} \left[\frac{4}{9} \left(\frac{d_d^C}{e} \right)_\Lambda + \frac{2}{9} \left(\frac{d_u^C}{e} \right)_\Lambda \right], \quad (22)$$

where in turn the strong coupling constant g_s and the u - and d -quark masses in d_u^C and d_d^C are calculated at the scale Λ .

In the following, we shall give a more quantitative discussion of the individual two-loop contributions to electron and neutron EDMs for $M_a = 150$ and 300 GeV. At the tree level, the charged-Higgs-boson mass M_{H^+} is related to the would-be CP-odd mass M_a by

$$M_{H^+}^2 = M_a^2 + M_W^2. \quad (23)$$

Even though this very last relation receives appreciable radiative corrections in the MSSM with Higgs-sector CP violation [13], we shall still make use of Eq. (23), as required by a consistent expansion in perturbation theory. As we have explicitly demonstrated in [5], the EDMs crucially depend on μ and $\tan\beta$ through $\xi_{t,b}$ in Eq. (8) and through the couplings of a and H^+ to electron and d quark. For the purpose of illustration, we therefore plot in Figs. 2 and 3 the numerical predictions for electron and neutron EDMs as functions of $\tan\beta$ and μ , respectively. Specifically, we consider the following values for the SUSY parameters:

$$\text{Fig. 2 : } \quad M_0 = \tilde{M}_Q = \tilde{M}_u = \tilde{M}_d = 0.6 \text{ TeV}, \quad A = |A_t| = |A_b| = 1 \text{ TeV}, \\ A = \mu = 1 \text{ TeV}, \quad \arg A = 90^\circ$$

$$\text{Fig. 3 : } \quad M_0 = \tilde{M}_Q = \tilde{M}_u = \tilde{M}_d = 0.6 \text{ TeV}, \quad A = |A_t| = |A_b| = 1 \text{ TeV}, \\ \tan\beta = 20, \quad \arg A = 90^\circ \quad (24)$$

In addition, we assume that the μ -parameter is real. In agreement with [5], we find that the dominant EDM effects originate from the $aA_\mu A_\lambda$ -coupling in d_e , $(d_e)^\gamma$, and from the $ag_\mu g_\lambda$ -coupling in d_n , $(d_n)^C$. In Figs. 2(a) and 3(a), we see that the charged-Higgs-boson two-loop contribution to EDM, $(d_e)^W$, is smaller than $(d_e)^\gamma$ by a factor 8. On the other

hand, Figs. 2(b) and 3(b) show that the corresponding EDM contributions, $(d_n)^W$, $(d_n)^\gamma$, and $(d_n)^Z$, are all of comparable size. They are roughly one order of magnitude smaller than $(d_n)^C$, and of opposite sign.

In conclusion, we have shown that EDM constraints on the CP-violating parameters related to the sectors of scalar top and bottom quarks can be significant for $\tan\beta \gtrsim 10$, and $\mu, A_{t,b} \gtrsim 0.5$ TeV. In particular, we have studied additional two-loop contributions mediated by W - and Z -boson interactions, which are found to be sub-dominant but non-negligible, and are therefore expected to play an important role in future phenomenological analyses of Higgs-sector CP violation in the MSSM.

Note added

While revising the paper, I became aware of Ref. [14] which addresses the same topic. After the final revisions, the results obtained by the two groups agree both analytically and numerically. The author wishes to thank Darwin Chang and Wai-Yee Keung for discussions.

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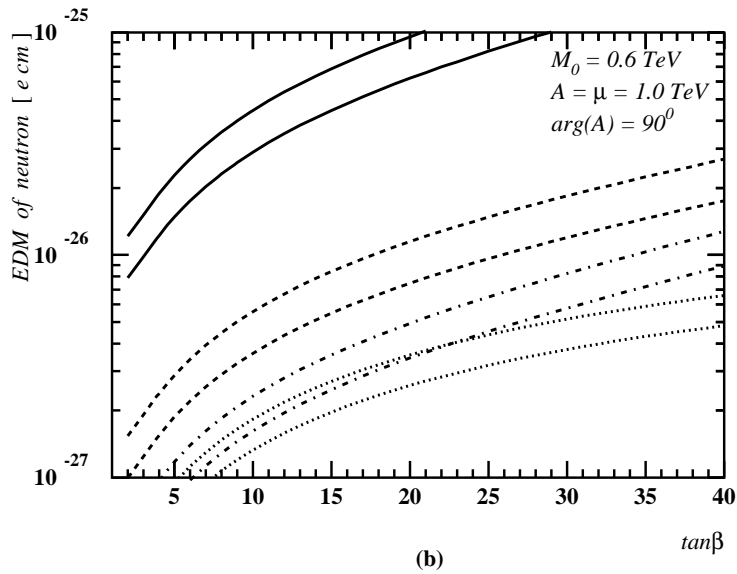
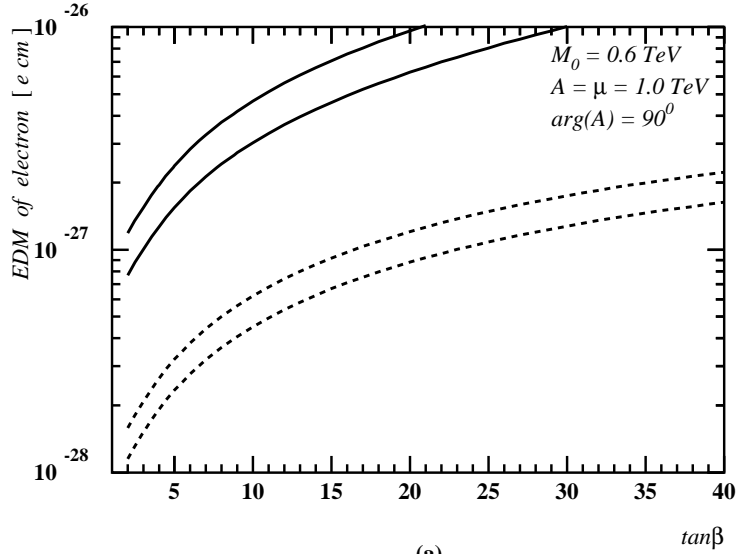


Figure 2: Numerical estimates of the individual two-loop EDM contributions as a function of $\tan\beta$: **(a)** $(d_e)^\gamma$ (solid line), $(d_e)^W$ (dashed line); **(b)** $-(d_n)^C$ (solid line), $(d_n)^\gamma$ (dashed line), $(d_n)^W$ (dotted line), $(d_n)^Z$ (dash-dotted line). Lines of the same type from the upper to the lower one correspond to $M_a = 150$ and 300 GeV, respectively.

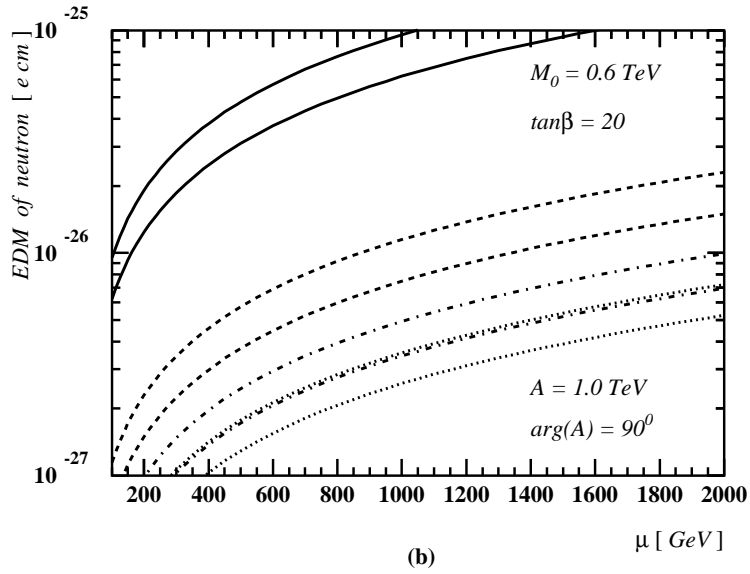
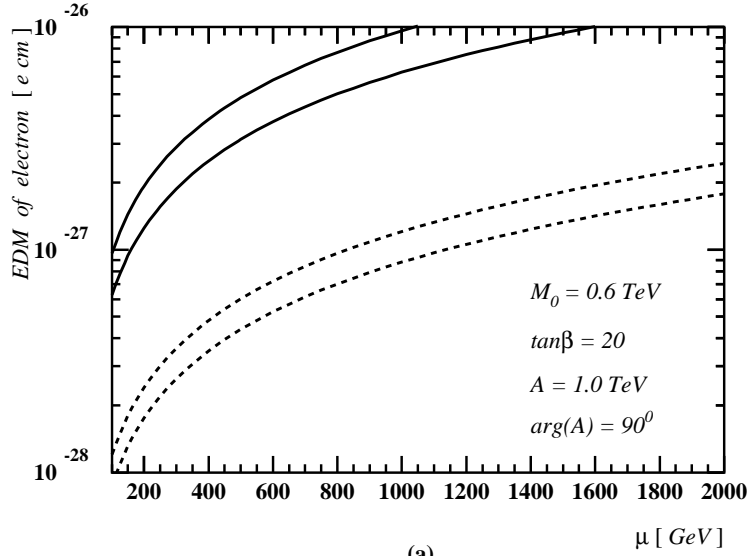


Figure 3: Numerical estimates of the individual two-loop EDM contributions as a function of μ : **(a)** $(d_e)^\gamma$ (solid line), $(d_e)^W$ (dashed line); **(b)** $-(d_n)^C$ (solid line), $(d_n)^\gamma$ (dashed line), $(d_n)^W$ (dotted line), $(d_n)^Z$ (dash-dotted line). Lines of the same type from the upper to the lower one correspond to $M_a = 150$ and 300 GeV, respectively.