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Weak Electric Dipole Moments of Heavy Fermions in the MSSM

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Abstract

A minimal supersymmetric version of the Standard Model with complex parameters allows contributions to the weak–electric dipole moments of fermions at the one–loop level. Assuming generation–diagonal trilinear soft–susy–breaking terms and the usual GUT constraint, a set of CP–violating physical phases can be introduced. In this paper the general expressions for the one–loop contribution to the WEDM in a generic renormalizable theory are given and the size of the WEDM of the τ lepton and the b quark in such a supersymmetric model is discussed.

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In the electroweak Standard Model (SM) there is only one source of CP violation, the $\delta_{\rm CKM}$ phase of the Cabibbo–Kobayashi–Maskawa mixing matrix for quarks [1].[1] The only place where CP violation has been currently measured, the neutral K system, fixes the value of this phase but does not constitute itself a test for the origin of CP violation [2]. On the other hand, if the baryon asymmetry of the universe has been dynamically generated, CP must be violated. The SM cannot account for the size of the observed asymmetry [3]. Many extensions of the SM contain new CP–violating phases, in particular, the supersymmetric models [4]. It has also been shown that the Minimal Supersymmetric Standard Model (MSSM) [5] can provide the correct size of baryon asymmetry in some range of parameters if the CP–violating phases are not suppressed [6].

One needs soft–breaking terms to introduce physical phases in the MSSM, different from the $\delta_{\rm CKM}$ [7]. We assume that the soft–breaking terms preserve R–parity. Other possibilities for CP violation can arise in R–parity violating models (cf. e.g. [8] in the context of R–parity violating scalar interactions). For simplicity, we restrict ourselves to generation–diagonal trilinear soft–breaking terms to prevent FCNC. Doing this we ignore CP–violating effects that have been already considered in the literature [9]. In our analysis we do not make any additional assumption, except for the unification of the soft–breaking gaugino masses at the GUT scale. However we do not assume unification of the scalar mass parameters or trilinear mass parameters. In such a constrained framework the following SUSY parameters can be complex: the Higgs–Higgsino mass parameter μ; the gaugino mass parameters M_1, M_2 and M_3; the bilinear mixing mass parameter m_{12}^2 and the trilinear soft supersymmetry breaking parameters A_τ, A_t and A_b (and accordingly for the other two generations). Not all of these phases are physical. Namely, the MSSM has two additional U(1) symmetries for vanishing μ and soft–breaking terms: the Peccei–Quinn and the R–symmetry. For non vanishing μ and soft–breaking terms these symmetries can be used to absorb two of the phases by redefinition of the fields [10]. In addition, the GUT relationship between M_1, M_2 and M_3 leads to only one common phase for the gaugino mass parameters. Hence, we are left with four independent CP–violating physical phases (only two assuming universality in the sfermion sector).

The most significant effect of the CP–violating phases in the phenomenology is their contribution to electric dipole moments (EDMs) [11]. Unlike the SM, where the contribution to the EDM of fermions arises beyond two loops [12], the MSSM can give a contribution already at the one–loop level. The measurement of the neutron EDM [13] constrains the phases and the supersymmetric spectrum in a way that may demand fine

[1] The QCD Lagrangian includes an additional source of CP violation, the $\theta_{\rm QCD}$, but we will not consider it here. Extreme fine tuning is needed in order that its contribution to the neutron EDM does not exceed the present experimental upper bound. Various mechanisms beyond the SM have been proposed to solve this problem [2].

tuning (supersymmetric CP problem): phases of $\mathcal{O}(10^{-2})$ [4], or SUSY particles very heavy (several TeV [14]). This problem could be solved if soft supersymmetry breaking terms are universal and the genuine SUSY CP phases vanish (the Yukawa matrices are then the only source of CP violation, like in the SM). It has been argued that one could still keep the SUSY phases of $\mathcal{O}(1)$ and the SUSY spectrum not very heavy and satisfy the experimental bounds due to cancellations among the different components of the neutron EDM [15]. Furthermore it has been recently shown [16] that, even without such cancellations and in the context of non universal soft–breaking terms, the current experimental limits on the neutron EDM can be met with almost no fine tunning on the CP–violating phases (even for the first generation ones), at the only price of $\arg(\mu)$ of $\mathcal{O}(10^{-2})$. In our analysis we do not assume universality and keep all the SUSY phases as free parameters.

As a generalization of the electromagnetic dipole moments of fermions, one can define weak dipole moments (WDMs), corresponding to couplings with a Z boson instead of a photon. The most general Lorentz structure of the vertex function that couples a Z boson and two on–shell fermions (with outgoing momenta q and $\bar{q}$) can be written in terms of form factors $F_i(s \equiv (q+\bar{q})^2)$ as

$$\begin{aligned}\Gamma^{Zff}_{\mu} &= ie\Big\{\gamma_\mu\left[\left(F_{\rm V}-\frac{v_f}{2s_Wc_W}\right)-\left(F_{\rm A}-\frac{a_f}{2s_Wc_W}\right)\gamma_5\right] \\ &\quad +(q-\bar{q})_\mu[F_{\rm M}+F_{\rm E}\gamma_5]-(q+\bar{q})_\mu[F_{\rm S}+F_{\rm P}\gamma_5]\Big\}\ , \qquad (1)\end{aligned}$$

where $v_f \equiv (I_3^f - 2s_W^2Q_f)$, $a_f \equiv I_3^f$. The form factors $F_{\rm M}$ and $F_{\rm E}$ are related to the anomalous weak magnetic and electric dipole moments of the fermion f with mass m_f as follows:

$$\begin{aligned}\text{AWMDM} \equiv a_f^W &= -2m_f\ F_{\rm M}(M_Z^2)\ , \\ \text{WEDM} \equiv d_f^W &= ie\ F_{\rm E}(M_Z^2)\ .\end{aligned}$$

The $F_{\rm M}$ ($F_{\rm E}$) form factors are the coefficients of the *chirality–flipping* term of the CP–conserving (CP–violating) effective Lagrangian describing Z–fermion couplings. Therefore, they are expected to get contributions proportional to some positive power of the mass of the fermions involved. This allows the construction of observables which can be probed experimentally most suitably by heavy fermions. Hence, for on–shell Z bosons, where the dipole form factors are gauge independent, the b quark and τ lepton are the most promising candidates.

In this work we concentrate on the analysis of the one–loop contribution of the MSSM with complex parameters and conserved R–parity to the WEDM of the τ lepton and the b quark.[2]

[2]The AWMDM has been considered in Refs. [18] where real supersymmetric couplings were used.

The WEDM

All the possible one–loop contributions to the WEDM can be classified in terms of six classes of triangle diagrams (Fig. 1). The vertices are labelled by generic couplings, according to the following interaction Lagrangian, for vectors $V_\mu^{(k)} = A_\mu,\ Z_\mu,\ W_\mu,\ W_\mu^\dagger$, general fermions Ψ_k and general scalars Φ_k:

$$\begin{aligned}\mathcal{L} &= ieJ(W^\dagger_{\mu\nu}W^\mu Z^\nu - W^{\mu\nu}W^\dagger_\mu Z_\nu + Z^{\mu\nu}W^\dagger_\mu W_\nu) + eV^{(k)}_\mu \bar\Psi_j\gamma^\mu(V^{(k)}_{jl} - A^{(k)}_{jl}\gamma_5)\Psi_l \\ &+ ieG_{jk}Z^\mu\Phi^\dagger_j \overleftrightarrow{\partial}_\mu \Phi_k + \Big\{e\bar\Psi_f(S_{jk} - P_{jk}\gamma_5)\Psi_k\Phi_j + eK_{jk}Z^\mu V^{(k)}_\mu\Phi_j + h.c.\Big\} \qquad (2)\end{aligned}$$

The expressions for the WEDM are evaluated in the 't Hooft–Feynman gauge (all the would–be–Goldstone bosons must be included) and are written in terms of three–point one–loop tensor integrals $\bar C$ and vertex coefficients:[3]

$$\begin{aligned}\frac{d_f^W}{e}(\mathrm{I}) = \frac{\alpha}{4\pi}\Big\{&4m_f\sum_{jkl}\mathrm{Im}[V^Z_{jk}(V^{(l)}_{fj}A^{(l)*}_{fk} + A^{(l)}_{fj}V^{(l)*}_{fk}) + A^Z_{jk}(V^{(l)}_{fj}V^{(l)*}_{fk} + A^{(l)}_{fj}A^{(l)*}_{fk})] \\ &\times[2C_2^{+-} - C_1^-](k,j,l) \\ &-4\sum_{jkl}m_k\mathrm{Im}[V^Z_{jk}(V^{(l)}_{fj}A^{(l)*}_{fk} - A^{(l)}_{fj}V^{(l)*}_{fk}) - A^Z_{jk}(V^{(l)}_{fj}V^{(l)*}_{fk} - A^{(l)}_{fj}A^{(l)*}_{fk})] \\ &\times[2C_1^+ - C_0](k,j,l)\Big\} \qquad (3)\end{aligned}$$

$$\begin{aligned}\frac{d_f^W}{e}(\mathrm{II}) = -\frac{\alpha}{4\pi}\Big\{&2m_f\sum_{jkl}\mathrm{Im}[J(V^{(j)}_{fl}A^{(k)*}_{fl} + A^{(j)}_{fl}V^{(k)*}_{fl})][4C_2^{+-} - C_1^-](k,j,l) \\ &+6\sum_{jkl}m_l\mathrm{Im}[J(V^{(j)}_{fl}A^{(k)*}_{fl} - A^{(j)}_{fl}V^{(k)*}_{fl})]C_1^+(k,j,l)\Big\} \qquad (4)\end{aligned}$$

$$\begin{aligned}\frac{d_f^W}{e}(\mathrm{III}) = \frac{\alpha}{4\pi}\Big\{&-2m_f\sum_{jkl}\mathrm{Im}[V^Z_{jk}(P_{lj}S^*_{lk} + S_{lj}P^*_{lk}) + A^Z_{jk}(S_{lj}S^*_{lk} + P_{lj}P^*_{lk})] \\ &\times[2C_2^{+-} - C_1^-](k,j,l) \\ &+2\sum_{jkl}m_k\mathrm{Im}[V^Z_{jk}(P_{lj}S^*_{lk} - S_{lj}P^*_{lk}) + A^Z_{jk}(S_{lj}S^*_{lk} - P_{lj}P^*_{lk})] \\ &\times[C_1^+ + C_1^-](k,j,l)\Big\} \qquad (5)\end{aligned}$$

$$\begin{aligned}\frac{d_f^W}{e}(\mathrm{IV}) = \frac{\alpha}{4\pi}\Big\{&2m_f\sum_{jkl}\mathrm{Im}[G_{jk}(S_{jl}P^*_{kl} + P_{jl}S^*_{kl})][2C_2^{+-} - C_1^-](k,j,l) \\ &-2\sum_{jkl}m_l\mathrm{Im}(G_{jk}S_{jl}P^*_{kl})[2C_1^+ - C_0](k,j,l)\Big\} \qquad (6)\end{aligned}$$

[3] Equivalent expressions for classes III and IV can be found in Ref. [19] where a different set of generic couplings and tensor integrals is employed.

$$\frac{d_f^W}{e}(\mathrm{V}+\mathrm{VI}) = -\frac{\alpha}{4\pi}\sum_{jkl} 2\mathrm{Im}[K_{jk}(V_{fl}^{(k)}P_{jl}^{*}+A_{fl}^{(k)}S_{jl}^{*})][C_1^{+}+C_1^{-}](k,j,l) \tag{7}$$

The arguments of the tensor integrals refer to $\bar{C}(k,j,l) \equiv \bar{C}(-\bar{q},q,M_k,M_j,M_l)$ of Ref. [20]. The tensor integrals are defined in such a way that for equal external fermion masses C_1^- and C_2^{+-} are antisymmetric under the interchange of k and j, whereas C_0 and C_1^+ are symmetric. The contribution of diagrams of class I and II vanishes as they can only involve SM fermions in the loop (MSSM preserves R–parity) whose couplings to gauge bosons are either real (Z–exchange) or self–conjugated (W–exchange). The gluonic contribution in class I contains only real couplings. For the class V and VI diagrams, the only contribution to the WEDM might occur when a pseudoscalar Higgs boson is involved in the loop, but there is no coupling of two neutral gauge bosons to a pseudoscalar and hence they also vanish. One can easily check that the Higgs sector for both the SM and the MSSM to class III and IV diagrams does not contribute to the WEDM, consistently with the CP–conserving character of both the SM and MSSM Higgs sectors. In a general 2HDM a CP–violating contribution is possible [17]. These considerations lead to the well known result that the SM one–loop contribution to the WEDM is zero. The MSSM contribution comes from charginos, neutralinos, gluinos and sfermions via diagrams of class III and IV. Finally, notice that all the contributions are proportional to one of the fermion masses involved, as expected from the chirality flipping character of the dipole moments.[4]

The WEDM of τ lepton and the b quark

The conventions for couplings and mixings in the MSSM are the ones in Ref. [5, 21] except for the complex character of the μ parameter and the trilinear soft supersymmetry breaking parameters A_τ, A_t and A_b. For convenience, we deal with the following CP–violating phases:[5] $\varphi_\mu \equiv \arg(\mu)$, $\varphi_{\tilde{f}} \equiv \arg(m_{\mathrm{LR}}^f)$ ($f=\tau,\ t,\ b$) with $m_{\mathrm{LR}}^t \equiv A_t - \mu^*\cot\beta$ and $m_{\mathrm{LR}}^{\tau,b} \equiv A_{\tau,b} - \mu^*\tan\beta$. We assume a common squark mass parameter $m_{\tilde{q}} \equiv m_{\tilde{Q}} = m_{\tilde{U}} = m_{\tilde{D}}$ as well as a common slepton mass parameter $m_{\tilde{l}} \equiv m_{\tilde{L}} = m_{\tilde{E}}$. We take real gaugino mass parameters constrained by the GUT relations:

$$M_1 = \frac{5}{3}\tan^2\theta_W M_2\ , \qquad M_3 = \frac{\alpha_s}{\alpha}s_W^2 M_2\ . \tag{8}$$

A "natural" scale for the EDMs is a "magneton" defined by $\mu_f \equiv e/2m_f = 1.7\times 10^{-15}$ (0.7×10^{-15}) $e$cm for the τ lepton (b quark). In the plots the dimensionless quantity

[4]For the diagrams of class V and VI the chirality flipping occurs at the scalar–fermion vertex and the fermion mass is embedded in S and P, which are in this case Yukawa couplings.

[5]Such a choice leads to a dependence on φ_μ of chargino and neutralinos masses. Conversely the sfermion masses are independent on $\varphi_{\tilde{f}}$.

d^W_f/μ_f is displayed.

We make a full scan of the SUSY parameter space and determine the values of the CP–violating phases that yield the maximum effect on the WEDM. The result of this analysis is described below.

Chargino and scalar neutrino contribution to d^W_τ

There is only one phase, φ_μ, involved in the chargino contribution as there is no mixing in the scalar neutrino sector. The result grows with $\tan\beta$. It also depends on the common slepton mass (whose effect consists of dumping the result through the tensor integrals) and the $|\mu|$ and M_2 mass parameters. A value $\varphi_\mu = \pi/2$ enhances the WEDM. Taking $M_2 = |\mu| = 250$ GeV, $\mathrm{Re}(d^W_\tau[\text{charginos}]) = 0.18\ (5.52) \times 10^{-6}\ \mu_\tau$ for $\tan\beta = 1.6\ (50)$ and $m_{\tilde{l}} = 250$ GeV. There is no contribution to the imaginary part assuming the present bounds on the chargino masses.

Neutralino and $\tilde{\tau}$ slepton contribution to d^W_τ

Now both φ_μ and $\varphi_{\tilde{\tau}}$ contribute. Assuming that $|m^\tau_{LR}|$ is of the order of $|\mu|\tan\beta$ or below we get that there is no large influence of $\varphi_{\tilde{\tau}}$ on the neutralino contribution,[6] for both low and high $\tan\beta$. Thus we observe that this contribution is roughly proportional to $\sin\varphi_\mu$ regardless the value of $\varphi_{\tilde{\tau}}$. Taking $\varphi_\mu = \pi/2$, $M_2 = |\mu| = 250$ GeV and $m_{\tilde{l}} = 250$ GeV we find that the $\mathrm{Re}(d^W_\tau[\text{neutralinos}]) = -0.01\ (-0.25) \times 10^{-6}\ \mu_\tau$ for $\tan\beta = 1.6\ (50)$. For presently non excluded masses of the neutralinos there can be a contribution to the imaginary part, of the order of $10^{-6}\ \mu_\tau$.

Chargino and $\tilde{t}$ squark contribution to d^W_b

Two CP–violating phases are involved in this contribution: φ_μ and $\varphi_{\tilde{t}}$. In Fig. 2(a) the dependence on these phases is shown, for $M_2 = |\mu| = m_{\tilde{q}} = 250$ GeV, $|m^t_{LR}| = |\mu|\cot\beta$ and both low and high $\tan\beta$ scenarios. The maximum effect on the WEDM is obtained for $\varphi_\mu = \pi/2$ and $\varphi_{\tilde{t}} = \pi$. For example, one gets $\mathrm{Re}(d^W_b[\text{charginos}]) = 1.17\ (27.1) \times 10^{-6}\ \mu_b$ for low (high) $\tan\beta$. As expected, in the high $\tan\beta$ scenario our assumed $|m^t_{LR}|$ takes a small value and the dependence on $\varphi_{\tilde{t}}$ tends to dissappear. To have an idea of the maximum value achievable for the chargino contribution, we show in Fig. 3 the dependence on M_2 and $|\mu|$ for $\varphi_\mu = \pm\pi/2$ and $m^t_{LR} = 0$ with the previous value for the common squark mass parameter.

[6]The size of $|m^\tau_{LR}|$ is critical for the τ sleptons to have a physical mass.

Neutralino and $\tilde{b}$ squark contribution to d_b^W

The two relevant CP–violating phases for this case are: φ_μ and $\varphi_{\tilde{b}}$. As before, the most important effect from $\varphi_{\tilde{b}}$ arises when the off–diagonal term is larger, which in this case corresponds to high $\tan\beta$, as the trilinear soft breaking parameter is taken to be of the order of $|\mu|\tan\beta$. The maximum value for the neutralino contribution occurs for $\varphi_\mu = \varphi_{\tilde{b}} = \pi/2$ (Fig. 2(b)). The total contribution increases with $\tan\beta$. Thus one gets $\mathrm{Re}(d_b^W[\text{neutralinos}]) = -0.29\ (-12.6) \times 10^{-6}\ \mu_b$ for low (high) $\tan\beta$ with $M_2 = |\mu| = m_{\tilde{q}} = 250$ GeV and $|m_{LR}^b| = |\mu|\tan\beta$.

Gluino contribution to d_b^W

The gluino contribution is affected only by m_{LR}^b, $m_{\tilde{q}}$ and the gaugino mass M_3. Therefore the maximum value occurs for $\varphi_{\tilde{b}} = \pi/2$. The mixing in the $\tilde{b}$ sector is determined by m_{LR}^b and intervenes in the contribution due to chirality flipping in the gluino internal line (the contribution proportional to M_3). The contribution to the AWMDM is enhanced by the largest values of $|m_{LR}^b|$ compatible with an experimentally not excluded mass for the lightest $\tilde{b}$ squark. For zero gluino mass, only the term proportional to the mass of the b quark provides a contribution. As we increase the gluino mass, the term proportional to M_3 dominates, especially for large $|m_{LR}^b|$, being again suppressed at high M_3 due to the gluino decoupling. Thus for $\varphi_{\tilde{b}} = \pi/2$ one gets $\mathrm{Re}(d_b^W[\text{gluinos}]) = 0.26\ (9.31) \times 10^{-6}\ \mu_b$ for low (high) $\tan\beta$ and $|m_{LR}^b| = |\mu|\tan\beta$, $M_2 = |\mu| = m_{\tilde{q}} = 250$ GeV and M_3 fulfilling the GUT relation (8).

Conclusions

Unlike the SM, an R–parity preserving MSSM version with complex parameters contains enough freedom to provide a contribution to the (W)EDMs to one loop: considering generation–diagonal trilinear soft–susy–breaking terms, to reduce undesired FCNC, and the GUT constraint between the gauginos mass parameters, at most two (three) CP–violating physical phases are available for lepton (quark) WEDMs, apart from the δ_{CKM}.

In this work, the one–loop analytical expressions for the (W)EDM of fermions in any renormalizable theory are given in terms of a set of generic couplings. Moreover, a full scan of the MSSM parameter space has been performed in search for the maximum effect on the WEDM of the τ lepton and the b quark. The Higgs sector does not contribute and chargino diagrams are more important than neutralino ones. Gluinos are also involved in the b case and compete in importance with charginos. In the most favourable configuration of CP–violating phases and for values of the rest of the parameters still not excluded by experiments, these WEDMs can be as much as twelve orders of magnitude larger than the SM predictions, although still far from experimental reach:

$$\begin{aligned}|\mathrm{Re}(d_\tau^W)| &\lesssim 0.3\ (10) \times 10^{-21}\ e\mathrm{cm} \\ |\mathrm{Re}(d_b^W)| &\lesssim 1\ (20) \times 10^{-21}\ e\mathrm{cm}\end{aligned}$$

There may be a contribution to the imaginary part if the neutralinos are light. The current experimental bound on the τ WEDM [23] is $|\mathrm{Re}(d_\tau^W)| < 5.6 \times 10^{-18}$ $e$cm and $|\mathrm{Im}(d_\tau^W)| < 1.5 \times 10^{-17}$ $e$cm at 95% confidence level. There does not exist any similar analysis for the b which might be not possible due to hadronization. For comparison with other theoretical predictions: the electron EDM in the SM [11] ($d_f \propto m_f$) can be estimated to be $d_e \sim 10^{-37}$ $e$cm; in multi–Higgs models [24] $d_\tau^W \sim 3 \times 10^{-22}$ $e$cm; in leptoquark models [25] $d_\tau^W \sim 10^{-19}$ $e$cm.

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Figures

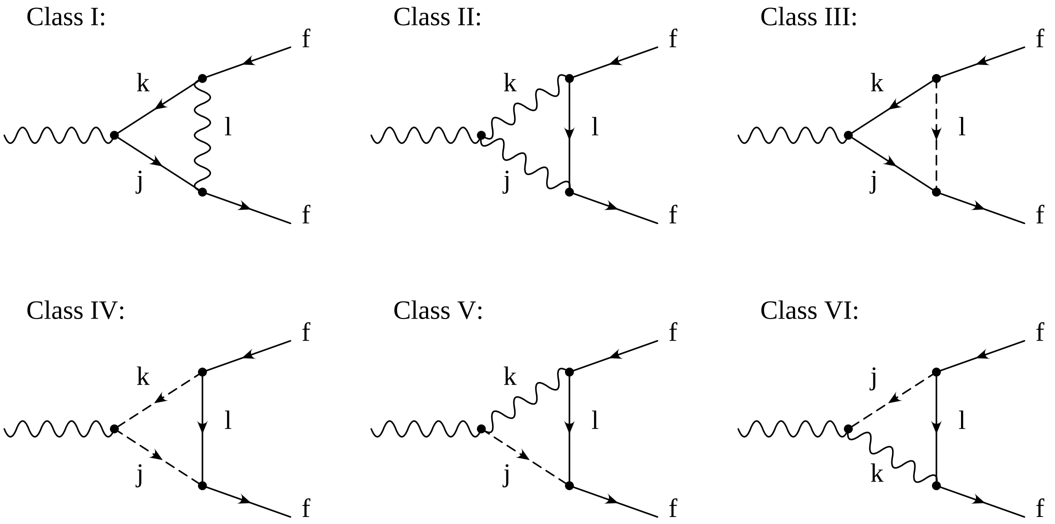

Figure 1: The one–loop Zff diagrams.

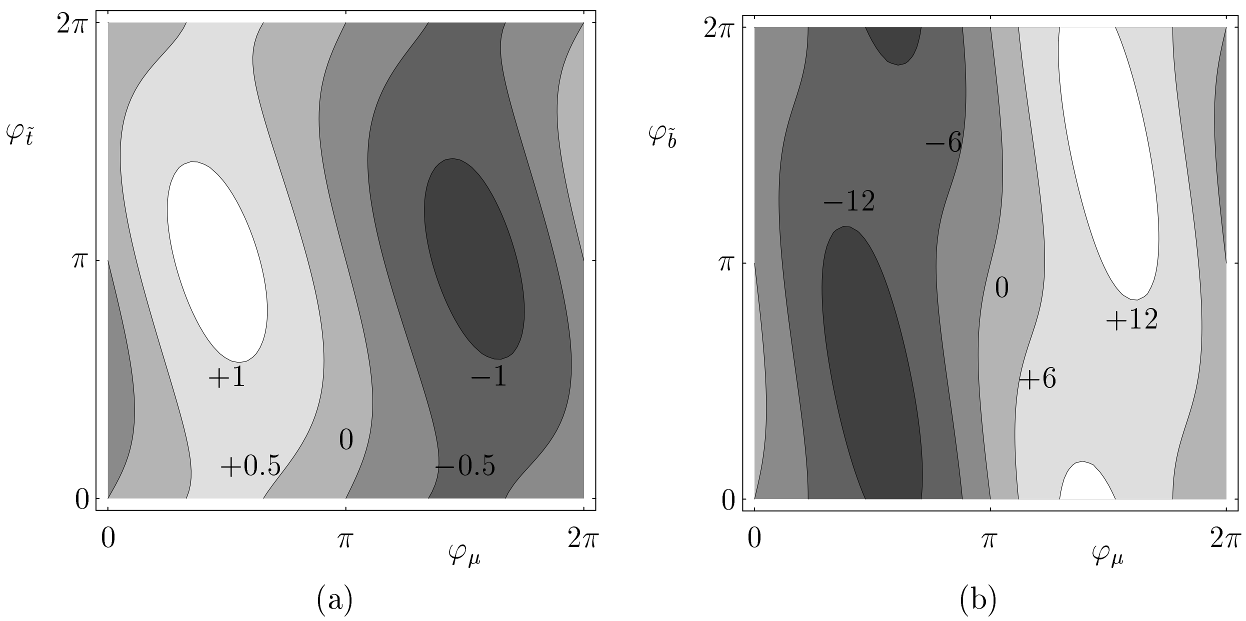

Figure 2: The boundaries of the different shaded areas are contour lines in the plane $\varphi_{\tilde{f}} - \varphi_\mu$ showing (a) $\mathrm{Re}(d_b^W[\text{charginos}])$ in units of 10^{-6} μ_b for low $\tan\beta$ and (b) $\mathrm{Re}(d_b^W[\text{neutralinos}])$ in units of 10^{-6} μ_b for high $\tan\beta$. $M_2 = |\mu| = m_{\tilde{q}} = 250$ GeV and $|m_{LR}^t| = |\mu|\cot\beta$ for (a) and $|m_{LR}^b| = |\mu|\tan\beta$ for (b).

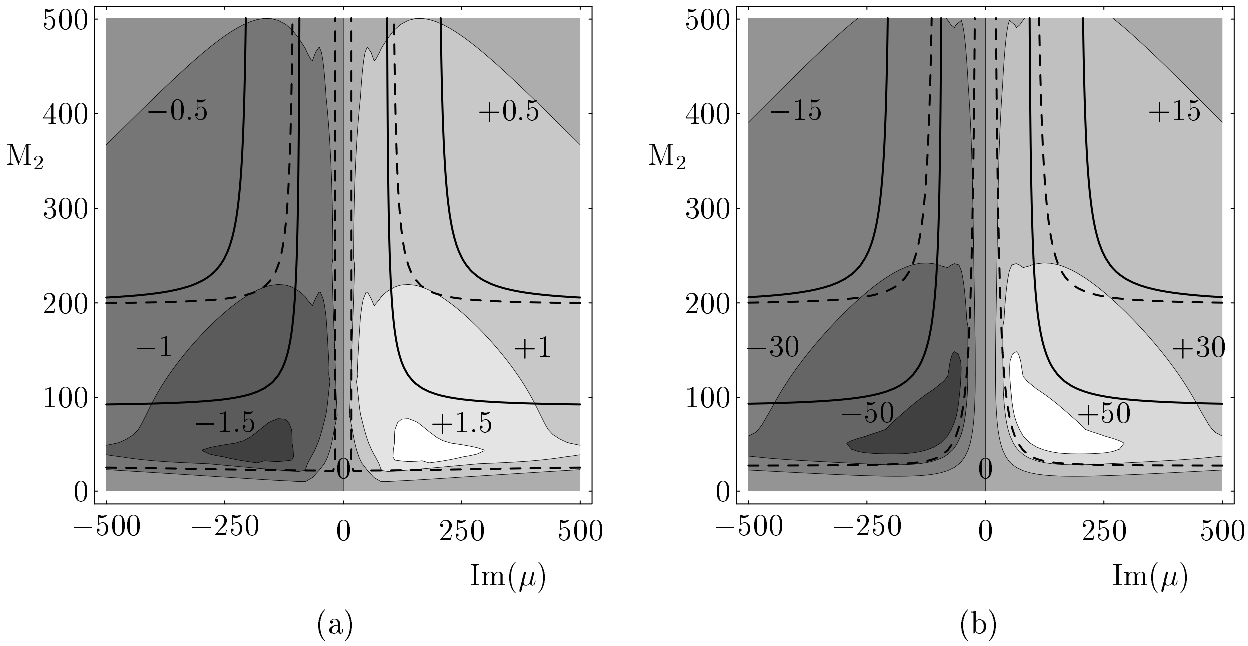

Figure 3: The boundaries of the different shaded areas are contour lines in the plane M_2−Im(μ) showing Re(d_b^W[charginos]) in units of 10^{-6} μ_b for (a) low and (b) high $\tan\beta$ with $|\sin\varphi_\mu| = 1$, $m_{\tilde{q}} = 250$ GeV and $|m_{LR}^t| = 0$. Also indicated are the isocurves corresponding to lightest chargino masses $m_{\tilde{\chi}_1^\pm} = 91$ and 200 GeV (solid) and lightest neutralino masses $m_{\tilde{\chi}_1^0} = 14$ and 100 GeV (dashed). The present LEP limits at $\sqrt{s} = 183$ GeV are $m_{\tilde{\chi}_1^\pm} > 91$ GeV and $m_{\tilde{\chi}_1^0} > 14$ GeV [22].