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The electric dipole moment of the neutron in the constrained minimally supersymmetric standard model

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Abstract

We analyse the electric dipole moment of the neutron in the MSSM, induced by the renormalisation of the soft-susy breaking terms. We run the RGEs using two-loop expressions for gauge and Yukawa couplings and retaining family dependence. The μ and B parameters were determined by minimising the full one-loop Higgs potential, and we find that the neutron EDM lies in the range $10^{-33} < |d_n| < 10^{-29}$ e cm.

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Supersymmetric unified theories are the most promising candidates for physics beyond the standard model in that they resolve the crucial gauge hierarchy problem of widely separated electroweak and grand unified scales, are the consequence of string theories and are favoured over non-supersymmetric unified theories by recent high precision measurements at LEP. In addition to the usual signatures of grand unification such as proton decay, neutrino masses, fermion mass relations and weak mixing angle prediction, supersymmetric unification is characterised by the resultant mass spectrum of the supersymmetric particles (squarks, sleptons, charginos, neutralinos) and the flavour-changing and CP-violating processes which arise as the renormalisation group equations (RGE) scale the physics from the unification scale $M_U \sim 10^{16}$ GeV down to the electroweak scale. Of particular interest are the flavour changing neutral current transitions involving the quark-squark-gluino vertex, with their implications for rare B-decays and mixings [1], and the non-removable CP-violating phases, with their implications for quark electric dipole moments (EDM) [2, 3, 4] and for non-standard-model patterns of CP-violation in neutral B decays [5], which result from the RGE scaling of the soft supersymmetry (SUSY) breaking scalar interactions in these models.

CP-violation in the standard model (SM) arises from the single phase δ_{CKM} in the Cabbibo-Kobayashi-Maskawa (CKM) mixing matrix relating the quark weak interaction and mass eigenstates and, in principle, this source of CP violation can accommodate the known CP properties of the kaon system. Non-zero quark (and lepton) electric dipole moments are very sensitive probes of CP violation beyond the standard model [6] because, unlike the other observables of CP violation which are small because of the intergenerational mixing angles of the CKM matrix, electric dipole moments are *particularly* suppressed by the chiral nature of the weak interaction and vanish at both one- and two-loop order in the Standard Model, resulting in quark EDMs of

$$d_{u,d}^{SM} \sim O(10^{-34}) \quad e \text{ cm}. \tag{1}$$

At present the experimental bounds on quark EDMs are obtained indirectly from measurements of the neutron EDM. In the non-relativistic quark model the neutron EDM is

$$d_n(quarks) = \frac{4}{3}d_d - \frac{1}{3}d_u \tag{2}$$

and is of the same order as the u and d-quark EDMs. However the neutron EDM is expected to be dominated by long distance effects such that

$$d_n^{SM}(LD) = O(10^{-32}) \quad e \text{ cm}. \tag{3}$$

By contrast, non-zero supersymmetric phases, collectively denoted δ_{soft}^{SUSY}, arising from the complexity of the soft SUSY breaking terms can generate quark EDMs at one-loop order, irrespective of generation mixing [2] from diagrams involving gluino, chargino (or neutralino) exchange and mixing of right- and left-handed virtual squarks. However, for squark, gluino and chargino masses of order 100 GeV, these induced one-loop quark EDMs yield a neutron EDM which exceeds the experimental upper bound of

$$d_n^{EXP} < 12 \times 10^{-26} \quad e \text{ cm} \tag{4}$$

unless these soft phases are constrained to satisfy

$$\delta_{soft}^{SUSY} < 0.01. \tag{5}$$

The alternative scenario of not imposing any condition of smallness on these phases but instead making the supersymmetric scalar masses heavy enough (of the order of 1 TeV) to suppress the EDMs has been considered by Kizukuri and Oshimo[7]. This scenario also has consequences for the relic density of the lightest supersymmetric particle[8].

For supersymmetric unified models such as spontaneously broken $N = 1$ supergravity with flat Kähler metrics [9] the resultant explicit soft SUSY breaking terms at the scale $M_{SUSY} \sim M_U$ of local SUSY breaking are quite simple and these generic phases δ_{soft}^{SUSY} are reduced to just two phases $\delta_{A,B}^{SUSY}$ in addition to the usual δ_{CKM}. The most natural way of satisfying the experimental bound (4) and ensuring that the squark, gluino and chargino masses are not much above the electroweak scale is to assume that these phases $\delta_{A,B}^{SUSY}$ vanish identically at the unification scale because of CP conservation in the SUSY breaking sector. Under these conditions the only explicit CP violation at the unification scale is in the flavour-dependent Yukawa coupling matrices which are required to have the structure necessary to reproduce the CKM mixing matrix at the electroweak scale under RGE scaling. However, as the RGEs for the soft SUSY breaking trilinear couplings $A_{u,d,e}$ and bilinear coupling B depend on the Yukawa couplings, inclusion of flavour mixing in the RGEs can lead to large RGE-induced CP-violating phases in the off-diagonal components of the couplings triggered, in particular, by the complexity of the large t-quark Yukawa coupling.

The implications of such large phases for EDMs of quarks have been studied recently by Bertolini and Vissani [3] and Inui *et al* [4] within a $N = 1$ supergravity inspired minimally supersymmetric standard model (MSSM) in which the spontaneous breaking of the electroweak $SU(2) \times U(1)$ symmetry is driven by radiative corrections. Bertolini and Vissani argue that the dominant induced EDM is that of the d-quark arising from the one-loop diagram involving chargino exchange and find

$$d_d^{SUSY}(BV) \sim \mathcal{O}(10^{-30}) \quad e \text{ cm}, \tag{6}$$

four orders of magnitude greater than the standard model prediction (1) but still satisfying the experimental upper bound (4). Inui *et al* also find that the d-quark EDM from chargino exchange is dominant but obtain the much larger value

$$d_d^{SUSY}(INUI) \sim O(10^{-27} - 10^{-29}) \quad e \text{ cm} \tag{7}$$

which they ascribe to the inclusion of gaugino masses in their RGEs.

Recently Dimopoulos and Hall [6] have considered quark and lepton EDMs in a class of supersymmetric unified theories based on the gauge group SO(10) where the unification of all quarks and leptons of a particular generation into a single **16** spinorial representation leads to non-removable CKM-like phases in the Yukawa couplings which, under RGE scalings induced by a large t-quark Yukawa coupling, give rise to EDMs close to the experimental limits such that some regions of the parameter space of the minimal SO(10) theory are excluded.

The sensitivity of the quark EDMs to CP violation, and the fact that the EDMs of the Standard Model, the MSSM and GUT theories are nicely separated makes this an important window to physics at the unification scale. In addition the dipole moments in the MSSM are a minimum prediction of supersymmetry. Because of this, and also because of the numerical discrepancies between the calculations of Bertolini and Vissani and Inui *et al* for the EDM of the d-quark in the constrained MSSM, and the limited nature of the free parameters chosen in both sets of calculations, we have undertaken a more detailed study of quark EDMs.

We have used two loop evaluation of gauge and Yukawa couplings, rather than the one-loop RGEs as used for the existing EDM calculations, and minimised the full one-loop Higgs potential, including contributions from matter and gauge sectors. We do this following the very complete analyses of Kane *et al* [10] and Barger *et al* [11], but retaining the full flavour dependence in the RGEs as we run them. We do not consider it necessary to describe the entire procedure since this is outlined in some detail in refs.[11, 10], but we shall briefly consider some details of our analysis.

The superpotential of the MSSM is given by

$$W = h_u Q_L^\dagger H_2 U_R + h_d Q_L^\dagger H_1 D_R + h_e L^\dagger H_1 E_R + \mu H_1 \epsilon H_2, \tag{8}$$

where generation indices are implied, and we define the VEVs of the Higgs fields (v_1 and v_2) such that $m_u = h_u v_2$, $m_d = h_d v_1$ and $m_e = h_e v_1$. Supersymmetry may be broken softly by generic mass-squared scalar terms, gaugino masses, and by 'trilinear' couplings of the form,

$$\delta\mathcal{L} = A_u h_u \tilde{q}_L h_2 \tilde{u}_R + A_d h_d \tilde{q}_L h_1 \tilde{d}_R + A_e h_e \tilde{l} h_1 \tilde{e}_R + B\mu h_1 \epsilon h_2, \tag{9}$$

where again, generation indices are suppressed. In order to determine the Yukawa couplings at the weak scale, we first ran the Standard Model down to 1 GeV using the two-loop QED and three-loop QCD RGEs of Arason *et al* [12]. We then ran the full supersymmetry RGEs up to the GUT scale (i.e. where the $SU(2)$ and $U(1)$ gauge couplings unified). For this we used two-loop RGEs for the gauge and Yukawa couplings [13], and one-loop RGEs for everything else, in the $\overline{DR}$ scheme. Here we 'unify' by setting the strong coupling equal to the unified $SU(2)$ and $U(1)$ couplings. This neglects the effects of thresholds at the GUT scale which depend on the precise details of the GUT theory, and tends to give values of α_s (≈ 0.126 for a top mass of 174 GeV) at the weak scale which are a little on the high side [14]. At the GUT scale there are three parameters which we set by hand, the common scalar mass m_0, the common gaugino mass $m_{1/2}$, and the common trilinear coupling A. Some of this degeneracy (for example of the gaugino masses) is motivated by the presumed existence of a GUT theory, and some by minimal supergravity (together with the assumption that the effects of renormalisation between the GUT and Planck scales is small). From the point of view of determining the effects of renormalisation of the pure MSSM on the electric dipole moments, degeneracy of these parameters is the natural assumption. We then run the entire theory back down to the weak scale. We determined the mass eigenvalues and eigenstates at the relevant physical scale, $Q = m(Q)$. At the same time we retained the full supersymmetric spectrum in the *running* theory, and only decoupled states in the running of the gauge and Yukawa couplings. Given the value of $\tan\beta$ and the sign of μ, it is possible to minimise the effective potential using the tadpole equations of ref.[11] in the running theory. This is a valid procedure if we wish to minimise the one-loop effective potential only, and avoids the need to match Lagrangians at each particle threshold [10, 11]. The minimisation was done at $Q = m_{top}$ (which was taken to be 174 GeV throughout), leaving the canonical, hybrid, four-dimensional, parameter space (m_0, $m_{1/2}$, A, $\tan\beta$) in addition to sign(μ).

However, since all the contributions from matter and gauge sectors were included in the minimisation, the vacuum expectation values of the two neutral Higgs fields (or equivalently the values obtained for B and μ) should be insensitive to the momentum scale at which they are evaluated [15]. This was indeed found to be the case. The whole process was then iterated a number of times. Generally, the procedure converges very rapidly (within a few iterations), and we accurately recover the entire spectrum given in ref.[11] for a wide range of parameters[1].

The mass matrices (in the super-KM basis) were diagonalised numerically to yield the required mass eigenstates and diagonalisation matrices as follows,

$$\begin{aligned} \text{squarks}: &\quad V_{\tilde{q}}^{\dagger} M_{\tilde{q}}^{2} V_{\tilde{q}} = m_{\tilde{q}}^{2} \\ \text{neutralinos}: &\quad V_{N}^{\dagger} M_{N} V_{N} = m_{\chi^{0}} \\ \text{charginos}: &\quad U_{C}^{\dagger} M_{C} V_{C} = m_{\chi^{\pm}}, \end{aligned} \tag{10}$$

and the following constraints applied.

In addition to experimental constraints (we adopt those used in ref.[10], $m_{\chi^0} > 18$ GeV, $m_{\chi^\pm} > 47$ GeV, $m_{h^0} > 44$ GeV, $m_{h^\pm} > 44$ GeV, $m_A > 21$ GeV, $m_{\tilde{g}} > 141$ GeV, $m_{\tilde{\nu}} > 43$ GeV, $m_{\tilde{q}} > 45$ GeV) we insisted that the minimum was stable in the sense that the Higgs and squark mass-squareds were positive. We also required that the minimum which we obtained was global, and that there were no other minima which may have broken colour or charge. The constraints

$$\begin{aligned} A_{\tau}^{2} &< 3h_{\tau}^{2}(m_{\tilde{\tau}}^{2} + m_{L}^{2} + m_{1}^{2}) \\ A_{b}^{2} &< 3h_{b}^{2}(m_{\tilde{b}}^{2} + m_{Q}^{2} + m_{1}^{2}) \\ A_{t}^{2} &< 3h_{t}^{2}(m_{\tilde{t}}^{2} + m_{Q}^{2} + m_{2}^{2}) \end{aligned} \tag{11}$$

provide a coarse indication of this [16]. We also insisted that the lightest supersymmetric partner was the neutralino, and finally we required that the process converged (i.e. that our choice of parameters was not too close to any fixed points).

The diagonalisation matrices appear in trilinear couplings between the quarks and the heavy supersymmetric scalar bosons and fermions, in particular squarks and charginos or gluinos. It is CP violation (i.e. non-zero phases) in these matrices, at the interaction vertices of the diagrams shown in figs(1a,1b,2), which may induce a non-zero EDM. As discussed in ref.[4], we find that such phases are indeed induced into the A-terms by the running of the RGEs, and hence into the diagonalisation matrices.

Having established this fact, let us consider how the EDM arises. For completeness, we wish to include the u-quark contribution (which is usually neglected), and so we shall briefly re-examine the EDM calculation. In doing so we also hope to gain a little insight into the CP violating nature of the EDM. First focus on a gaugeless Lagrangian with two fermionic fields, one scalar field, and a single cubic coupling,

$$\begin{aligned} \mathcal{L} &= \overline{\psi}_1 \left(i\gamma^\mu \partial_\mu - m_1\right) \psi_1 + \overline{\psi}_2 \left(i\gamma^\mu \partial_\mu - m_2\right) \psi_2 \\ &\quad + \left|\partial_\mu \phi\right|^2 - m_\phi^2 \left|\phi\right|^2 \\ &\quad + \left(a\overline{\psi}_{2L}\psi_{1R} + b\overline{\psi}_{2R}\psi_{1L}\right)\phi^* + \left(a^*\overline{\psi}_{1R}\psi_{2L} + b^*\overline{\psi}_{1L}\psi_{2R}\right)\phi. \end{aligned} \tag{12}$$

[1]We would like to thank P. Ohmann for discussions

The ψ_1 field is the light quark whose EDM we would like to calculate, the scalar field, ϕ, represents the squark fields, and ψ_2 is the heavy fermion field (not another quark). For any particular quark, we may choose that basis in which the mass parameters are real. A CP transformation on this Lagrangian shows that, in order to have CP violation, the phases of a and b must be different (a common phase may be absorbed into the definition of ϕ). Now consider the self energy graphs in fig(1a,b). In addition to giving mass and wave function renormalisation, these diagrams also induce non-local terms which may be obtained by performing a derivative expansion

$$\begin{aligned}\Delta\mathcal{L} &= D\left(\partial_\mu\overline{\psi}_{1R}\partial^\mu\psi_{1L} - m_1^2\overline{\psi}_{1R}\psi_{1L}\right)G(x) \\ &+ D^*\left(\partial_\mu\overline{\psi}_{1L}\partial^\mu\psi_{1R} - m_1^2\overline{\psi}_{1L}\psi_{1R}\right)G(x) + \ldots\end{aligned} \tag{13}$$

where the dots represent terms which are higher order in momentum, and where

$$\begin{aligned}G(x) &= \frac{1}{(1-x)^3}\left(1 - x^2 + 2x\log x\right) \\ D &= a^*b/\left(32\pi^2 m_2\right) \\ x &= m_\phi^2/m_2^2.\end{aligned} \tag{14}$$

With CP violation D is complex. The EDM appears when we now introduce electromagnetic interactions whilst keeping this expansion gauge invariant by introducing covariant derivatives,

$$\begin{aligned}\partial_\mu\psi_1 &\to (\partial_\mu + iq_1A_\mu)\psi_1 \\ \partial_\mu\psi_2 &\to (\partial_\mu + iq_2A_\mu)\psi_2 \\ \partial_\mu\phi &\to (\partial_\mu + iq_\phi A_\mu)\phi\end{aligned} \tag{15}$$

with $q_1 = q_2 + q_\phi$. When D is complex, one can anticipate an electric dipole moment from $\Delta\mathcal{L}$, of

$$d = q_1 G(x)\mathrm{Im}(D). \tag{16}$$

In fact, when the heavy fermion is electrically neutral (like the gluino of fig(1a)), this *is* the one-loop contribution to the dipole moment. In general there is an additional gauge invariant contribution to the Lagrangian, coming from the heavy fermion charge. To determine this we must resort to the usual one-loop diagram shown in fig(2). It is found to be of the form

$$\Delta\mathcal{L}' = -q_2F_{\mu\nu}H(x)\left(\mathrm{Re}(D)\overline{\psi}_1\sigma^{\mu\nu}\psi_1 - i\mathrm{Im}(D)\overline{\psi}_1\sigma^{\mu\nu}\gamma^5\psi_1\right), \tag{17}$$

where,

$$\begin{aligned}F_{\mu\nu} &= \partial_\mu A_\nu - \partial_\nu A_\mu, \\ \sigma^{\mu\nu} &= \frac{i}{2}[\gamma^\mu, \gamma^\nu],\end{aligned} \tag{18}$$

and where,

$$H(x) = \frac{2}{(1-x)^2}\left(1 - x + x\log x\right). \tag{19}$$

The total EDM of a quark coming from chargino/squark loops is then,

$$d = \mathrm{Im}(D)\left[q_1 G(x) - q_2 H(x)\right]. \tag{20}$$

For each of the quarks this gives,

$$\begin{aligned} d_d &= \frac{e}{3}\mathrm{Im}(D_d)F_d(x) \\ d_u &= -\frac{2e}{3}\mathrm{Im}(D_u)F_u(x), \end{aligned} \tag{21}$$

where we have defined the functions,

$$\begin{aligned} F_d &= \frac{1}{(1-x)^3}\left[5 - 12x + 7x^2 + 2x(2-3x)\log x\right] \\ F_u &= \frac{1}{(1-x)^3}\left[2 - 6x + 4x^2 + x(1-3x)\log x\right]. \end{aligned} \tag{22}$$

The above analysis generalises in a straightforward manner. The coupling constants a and b become matrices a_{ij} and b_{ij}, with i, j running over the appropriate mass eigenstates. For the chargino contributions we find

$$\begin{aligned} d_d &= \frac{e}{32\pi^2}\sum_i^2 \frac{(V_C)_{2i}(U_C)_{2i}}{m_{\chi_i^\pm}}\mathrm{Im}\left(h_u\left[V_{\tilde{u}}F_u\left(\frac{m_{\tilde{u}}^2}{m_{\chi_i^\pm}^2}\right)V_{\tilde{u}}^\dagger\right]_{RL}^T K^\dagger h_d\right)_{11} \\ d_u &= -\frac{e}{16\pi^2}\sum_i^2 \frac{(V_C)_{2i}(U_C)_{2i}}{m_{\chi_i^\pm}}\mathrm{Im}\left(h_d\left[V_{\tilde{d}}F_d\left(\frac{m_{\tilde{d}}^2}{m_{\chi_i^\pm}^2}\right)V_{\tilde{d}}^\dagger\right]_{RL}^T h_u\right)_{11}, \end{aligned} \tag{23}$$

where we are using the down-quark diagonal basis, and where K is the CKM matrix. For the gluino contributions we find,

$$\begin{aligned} d_d &= \frac{e\alpha_s}{9\pi m_{\tilde{g}}}\mathrm{Im}\left(\left[V_{\tilde{d}}G\left(\frac{m_{\tilde{d}}^2}{m_{\tilde{g}}^2}\right)V_{\tilde{d}}^\dagger\right]_{LR}\right)_{11} \\ d_u &= \frac{e\alpha_s}{9\pi m_{\tilde{g}}}\mathrm{Im}\left(\left[V_{\tilde{u}}G\left(\frac{m_{\tilde{u}}^2}{m_{\tilde{g}}^2}\right)V_{\tilde{u}}^\dagger\right]_{LR}\right)_{11}. \end{aligned} \tag{24}$$

In order to present our results, we choose points generated at random in the parameter space given by

$$\begin{aligned} 0 < \ & m_0 & < 1\ \mathrm{TeV} \\ 0 < \ & m_{1/2} & < 1\ \mathrm{TeV} \\ -1 < \ & A & < 1\ \mathrm{TeV} \\ 0 < \ & \tan\beta & < 20. \end{aligned} \tag{25}$$

In practice, values higher than these seldom satisfy all the criteria detailed above (i.e. they imply fine-tuning). The region below the low $\tan\beta$ fixed point is excluded. For a top-quark of mass 174 GeV, $A = 0$ GeV and $m_{1/2} = m_0 = 150$ GeV, this was found to be at $\tan\beta = 2.2$.

The modulus of the neutron EDM is plotted against $\tan\beta$ in fig(3). There is a slight tendency for it to be positive, and the largest values occur for negative μ and positive A. Clearly the value of $\tan\beta$ dominates the EDM of the neutron, and we see the approximately linear behaviour for large $\tan\beta$ coming from the increased down-quark Yukawa couplings. Phases feed into the off-diagonal elements of A_d, especially into A_{d13}. The EDM becomes smaller as we approach the low $\tan\beta$ fixed point, since the top Yukawa coupling dominates the running. Here the gluino contribution to the up-quark can be the dominant contribution. Elsewhere however, the down-quark, chargino diagram is nearly always dominant. No obvious pattern emerges with the other three parameters. The value of the neutron EDM in the MSSM is much less than the value of 10^{-27} indicated in [4], because values of the A parameter as large as those used in ref.[4], give problems with colour or charge breaking minima, or do not lead to a solution for μ and B on minimisation. We find that the expected range for the EDM is therefore

$$10^{-33} < |d_n| < 10^{-29} e \text{ cm}. \tag{26}$$

It is expected that future developments will push the experimental bound on the neutron EDM down from eq.(4) to $O(10^{-28})e$ cm but, as the present calculations indicate, will still not provide a test of the constrained MSSM[2]. A more remote possibility is the direct measurement of the t-quark EDM using either $t\bar{t}$ decay correlations in $e^+e^- \to t\bar{t}$ [17] or $t\bar{t}$ production via photon-photon fusion using linearly polarised photons generated by Compton back-scattering of laser light on electron or positron beams of linear e^+e^- or e^-e^- colliders [18]. In this analysis, we found that the t-quark EDM is usually larger than the neutron EDM by a factor of 3–5, a slight improvement but still unlikely to be measured in the forseeable future, at least in the MSSM.

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Figure Captions

Figure 1 : Quark self energy diagrams involving (a) gluino and (b) chargino exchange.

Figure 2 : SUSY contribution to quark EDM from chargino-photon coupling.

Figure 3 : The modulus of Neutron EDM (in units of 10^{-33} e cm) as a function of $\tan\beta$ for random choices of m_0, $m_{1/2}$ and $|A|$ in the range (0, 1) TeV.

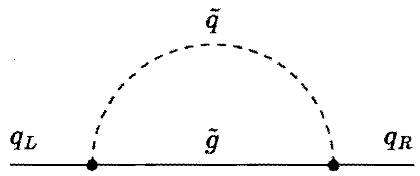

Figure 1a

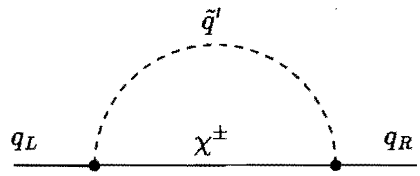

Figure 1b

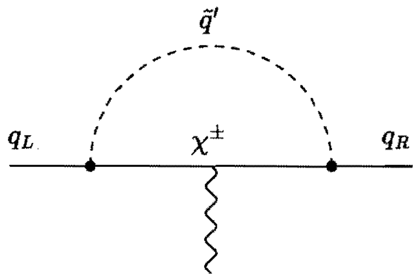

Figure 2

Fig(3): Electric Dipole Moment