The Electric Dipole Moment and CP Violation in $B \rightarrow X_s l^+ l^-$ in SUGRA Models with Nonuniversal Gaugino Masses

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Abstract

The constraints of electric dipole moments (EDMs) of electron and neutron on the parameter space in supergravity (SUGRA) models with nonuniversal gaugino masses are analyzed. It is shown that with a light sparticle spectrum, the sufficient cancellations in the calculation of EDMs can happen for all phases being order of one in the small $\tan\beta$ case and all phases but $\phi_\mu$ ($|\phi_\mu| \lesssim \pi/6$) order of one in the large $\tan\beta$ case. This is in contrast to the case of mSUGRA in which in the parameter space where cancellations among various SUSY contributions to EDMs happen $|\phi_\mu|$ must be less than $\pi/10$ for small $\tan\beta$ and $\mathcal{O}(10^{-2})$ for large $\tan\beta$. Direct CP asymmetries and the T-odd normal polarization of lepton in $B \rightarrow X_s l^+ l^-$ are investigated in the models. In the large $\tan\beta$ case, $A_{\text{CP}}^2$ and $P_N$ for $l=\mu$ ($\tau$) can be enhanced by about a factor of ten (ten) and ten (three) respectively compared to those of mSUGRA.
Recent observation of $\text{Re}(\epsilon'/\epsilon)$ by KTeV collaboration \cite{1} definitely confirms the earlier NA31 experiment \cite{2}. This direct CP violation measurement in the Kaon system can be accommodated by the CKM phase in standard model (SM) within the theoretical uncertainties. However, the CKM phase is not enough to explain the matter-antimatter asymmetry in the universe and gives the contribution to EDMs much smaller than the limits of EDMs of electron and neutron. One needs to have new sources of CP violation and examine their phenomenological effects.

There exist new sources of CP violation in SUSY theories which come from the phases of soft SUSY breaking parameters. It is well-known for a long time that in order to satisfy the current experimental limits on EDMs of electron and neutron SUSY CP-violating phases have to be much smaller ($\lesssim 10^{-2}$) unless sfermion masses of the first and second generations are very large ($> 1$ Tev) \cite{3}. Recently it has been shown that various contributions to EDMs cancel with each other in significant regions of the parameter space so that the current experimental limits on the EDM of electron (EDME) \cite{4} and the EDM of neutron (EDMN) \cite{5},

$$|d_e| < 4.3 \times 10^{-27} \text{ecm}$$

and

$$|d_n| < 6.3 \times 10^{-26} \text{ecm},$$

can be satisfied for SUSY models with SUSY phases of order one and relatively light sparticle spectrum ($< 1$ Tev) \cite{6,7}. In mSUGRA even in the parameter space where cancellation among various SUSY contributions to neutron EDM(EDMN) happens $|\phi_{\mu}|$ must be less than $\pi/10$ for small $\tan \beta$ \cite{6} and $O(10^{-2})$ for large $\tan \beta$ \cite{8} while the allowed range of $\phi_{A_0}$ is almost unconstrained. Brhlik et al. pointed out that more sufficient cancellations happen in MSSM if gaugino masses are complex \cite{9}. In the letter we consider cancellation phenomena in SUGRA with nonuniversal gaugino masses.

CP violation has so far only been observed in K system. It is one of the goals of the B factories presently under construction to discover and examine CP violation in the B system. Direct CP violation in $B \rightarrow X_s l^+ l^-$ in SM has been examined and the result is that it is unobservably small \cite{10}. In mSUGRA the CP asymmetry of branching ratio on $B \rightarrow X_s l^+ l^-$ has been given in \cite{22}. A detailed analysis of SUSY contributions to CP Violation in semileptonic B decays has been performed using the mass insertion approximation in \cite{23}. Direct CP asymmetries and the T-odd normal polarization of lepton in $B \rightarrow X_s l^+ l^-$ in mSUGRA with CP-violating phases are investigated in our previous paper \cite{8}. In the letter we extend the investigation to SUGRA models with nonuniversal gaugino masses after studying the allowed regions of the parameter space in the models by EDM data.

In order to concentrate on the effects of the phases arising from complex gaugino masses we limit ourself to a class of SUGRA models with nonuniversal gaugino masses in which gaugino masses at high energy scale (GUT scale) are nonuniversal but scalar masses and trilinear couplings at GUT scale are still universal. Such a class of effective SUGRA models can naturally arise from string models \cite{11}. In this class of models, compared to the mSUGRA, there are two more new independent phases \cite{9} which we choose to be $\phi_1$ and $\phi_3$, the phases of gaugino masses $M_1$ and $M_3$, in addition to the phases $\phi_{\mu}$ and $\phi_{A_0}$. From the one loop renormalization
group equations (RGEs) for $M_i$ $(i=1,2,3)$:

$$\frac{dM_i}{dt} = \frac{1}{4\pi} b_i \alpha_i M_i \quad i = 1, 2, 3 \tag{3}$$

where $\alpha_i = \frac{g_i^2}{4\pi}$, $t = \ln(Q^2/M_{GUT}^2)$, we know that the phases of $M_i$ do not run, like the phase of $\mu$.

Let us recall the cancellation mechanism for EDME. There are only two contributions, the chargino (-sfermion loop) and neutrino (-sfermion loop) contributions, to the EDME. The chargino contribution involves gaugino-Higgsino (g-h) mixing. The neutrino contribution involves not only gaugino-Higgsino mixing but also gaugino-gaugino (g-g) mixing. The chargino contribution and the part of the neutrino contribution which comes from g-h mixing have automatically opposite sign because of the opposite sign of $\mu$ in chargino and neutrino mass matrices. In general, the chargino contribution in magnitude is significantly larger than the part of neutrino contribution. Therefore, as pointed out in ref. [9], a cancellation can happen only if the another part of the neutrino contribution which comes from the g-g mixing can balance some of the difference between the two contributions. For EDME the neutrino contribution which comes from the g-g mixing is proportional to

$$\frac{1}{m_e^2} m_e [A_e sin(\phi_A - \phi_1) + |\mu| tan\beta sin(\phi_\mu + \phi_1)] \tag{4}$$

Therefore, given $\phi_\mu$, the sign of the contribution can be controlled by choosing $\phi_1$ and $\phi_A$ and the magnitude of the contribution can increase by increasing $\mu$ and/or $A_e$. Because the chargino contribution is dependent on $\mu$ and independent on $\phi_A$, it is sufficient to have a cancellation that the magnitude of the $A_e$ term (first term) in eq.(4) is comparable to that of the $\mu tan\beta$ term (second term) in eq.(4). This is easy to be down in MSSM in which $A_e$ and $\mu$ are free parameters. Thus, an almost exact cancellation can occur for the whole range of $\phi_\mu$. That is exactly what happens in MSSM [9].

However, in SUGRA models low energy properties are determined by running RGEs from the high energy scale to the electroweak (EW) scale and the radiative breaking mechanism of the EW symmetry puts constraints on CP-violating phases. As long as we limit our discussion to mass spectra less than than 1 Tev, $M_3$ and $A_0$ (hence $A_e$) can not be too large. For small tan$\beta$ (say $\sim 2$), the sufficient condition (i.e., the two terms in eq.(4) have size of the same order) can easily be realized in the almost whole range of $\phi_\mu$ by choosing $\phi_A$ and $\phi_1$. For moderate and large tan$\beta$, it is difficult for the condition to be realized due to the limited values of $A_0$ (hence $A_e$) so that only for some limited ranges of $\phi_\mu$ the EDM constraint can be satisfied. The similar (but more complicated) situation occurs for EDMN with appropriate $\phi_3$ as well as $\phi_A$ chosen.

In order to show the important role of $\phi_A$ played in the cancellation mechanism, in fig.1a and 1b we display EDME as function of $\phi_1$ for $\phi_A=0$ and different $\phi_\mu$ for both small tan$\beta$ (2) and large tan$\beta$ (30) cases. We can see from the fig.1 that most of the range of $\phi_\mu$ is excluded by EDME in both cases. For EDMN as function of $\phi_3$, similar results are obtained. That is, like $\phi_1$, for positive $\phi_\mu$ cancellations happen in some narrow ranges within $[\pi,2\pi]$ of $\phi_3$ and within $[0,\pi]$ for negative $\phi_\mu$. When we vary the values of $\phi_A$ we achieve the above mentioned results: almost whole range of $\phi_\mu$ is allowed by EDME and EDMN for small tan$\beta$ and $|\phi_\mu| \sim \pi/6$ for
large $\tan\beta$ (see fig. 2). Moreover, because $\phi_1(\phi_3)$ is correlated with $\phi_\mu$, we find that with varying $\phi_\mu$ the whole range of $\phi_1$ and $\phi_3$ can be allowed by EDM constraints. For large $\tan\beta$ (30) case largest $|\phi_\mu|$ (about $\pi/6$) correspond to $\phi_1$ and $\phi_3$ around $\mp\pi/2 \mp\pi/6$, while $\phi_1$ and $\phi_3$ are around $\pm\pi/2$ when $\phi_\mu$ about $\mp0.4$, and when $\phi_\mu$ is about $\pm0.2$ they are around $\mp\pi/4$. The correlated values of $\phi_3$ and $\phi_\mu$ are needed in analyses of $B \to X_l l^+ l^-$ and $B \to X_s l^+ l^-$ or $\bar{B} \to \bar{X}_s l^+ l^-$. Correlation between $\phi_\mu$ and $\tan\beta$, with the absolute value of soft breaking terms chosen as those in fig1 and appropriate phases chosen, is shown in fig. 2 where all of the points are allowed by the experimental bounds on EDM's and EDMN. One can see in the figure that $\phi_\mu$ becomes more constrained as $\tan\beta$ is increased. Nevertheless, for $\tan\beta$ larger than 6 the allowed regions of $\phi_\mu$ are almost unchanged, which means that effects of the $A_e$ term (and $A_d$ term in the case of EDMN) in eq.(4) are relatively small and the balance is provided by the $\mu\tan\beta$ term in eq.(4). Since we also consider the large $\tan\beta$ case, we include the two loop contribution given by D. Chang et.al[13]. But the numerical calculations in the regions of the parameter space in which one loop EDMs satisfy the current experimental limits due to the cancellation mechanism show that it is very small compared to one loop contributions.

We now turn to the calculations of the CP violation in $B \to X_l l^+ l^-$. The direct CP asymmetries in decay rate and backward-forward asymmetry for $B \to X_s l^+ l^-$ and $\bar{B} \to \bar{X}_s l^+ l^-$ are defined by \[8, 14\]

$$A_{CP}^1(\hat{s}) = \frac{d\Gamma/\hat{d}s - d\Gamma/\hat{d}s}{d\Gamma/\hat{d}s + d\Gamma/\hat{d}s} = \frac{D(\hat{s}) - \bar{D}(\hat{s})}{D(\hat{s}) + \bar{D}(\hat{s})},$$

$$A_{CP}^2(\hat{s}) = A(\hat{s}) - \bar{A}(\hat{s}) A(\hat{s}) + \bar{A}(\hat{s})$$

where

$$A(\hat{s}) = 3\sqrt{1 - \frac{4t^2}{s}} E(\hat{s}) \frac{D(\hat{s})}{D(\hat{s})},$$

$$D(\hat{s}) = 4|C_7|^2(1 + \frac{2t^2}{s})(1 + \frac{2t^2}{s}) + |C_8^{eff}|^2(2\hat{s} + 1)(1 + \frac{2t^2}{s}) + |C_9|^2[1 + 2\hat{s} + (1 - 4s)2t^2]$$

$$+ 12Re(C_8^{eff}C_7^*)(1 + \frac{2t^2}{s}) + \frac{3}{2}|C_{Q_1}|^2(1 - \frac{4t^2}{s})\hat{s} + \frac{3}{2}|C_{Q_2}|^2\hat{s} + 6Re(C_9C_{Q_2}^*)t$$

$$E(\hat{s}) = Re(C_8^{eff}C_7^*\hat{s} + 2C_7C_{Q_3}^* + C_8^{eff}C_{Q_2}^*t + 2C_7C_{Q_1}^*t)$$

(5)

Another observable related to CP violating effects in $B \to X_s l^+ l^-$ is the normal polarization of the lepton in the decay, $P_N$, which is the T-odd projection of the lepton spin onto the normal of the decay plane, i.e $P_N \sim \hat{s}_l \cdot (\vec{p}_e \times \vec{p}_l^-)$ \[15\]. A straightforward calculation leads to \[8, 16\]

$$P_N = \frac{3\pi}{4} \sqrt{1 - \frac{4t^2}{s}} \hat{s}^2 Im\left[2C_8^{eff}C_9 t + 4C_9 C_7^* \frac{t}{\hat{s}} + C_8^{eff}C_{Q_1} + 2C_7^* C_{Q_1} + C_9^* C_{Q_2}\right]/D(\hat{s})$$

(6)

The Wilson coefficients $C_i$ and $C_{Q_i}$ in eqs.(6) and (7) have been given in ref.\[8, 17, 18\]. Since only $C_8^{eff}$ contains the non-trivial strong phase, $A_{CP}^1$ is determined by $Im C_7$ and $A_{CP}^2$ by $Im C_{Q_1}$ and $Im C_7$. Although $P_N$ depends on all the relevant Wilson coefficients a large $P_N$ does require relatively large values of $Im C_{Q_i}(i = 1, 2)$ \[8\]. With the main contributions coming from exchanging chargino-stop loop with neutral Higgs coupled to external b quark \[17\], imaginary
parts of $C_{Q_i}$s come mainly from terms proportional to (unitarity condition for stop mixing matrix has been used)

$$\frac{m_{\chi_i}m_t}{m_{WZ}\sin\beta \cos\beta} U(i, 2)V(i, 2)D_{i21}D_{i11}^*, \quad i = 1, 2 \tag{8}$$

i.e CP violating phases enter into the imaginary parts of $C_{Q_i}$ through g-h mixings (U, V) and chiral mixing ($D_t$) of stops. From the chargino mass matrix

$$M_C = \begin{pmatrix} M_2 & \sqrt{2}m_W\sin\beta \\ \sqrt{2}m_W\cos\beta & \mu \end{pmatrix} \tag{9}$$

and stop mass matrix

$$M^2_t = \begin{pmatrix} M^2_0 + m_t^2 + M^2_z(\frac{1}{2} - Q_u \sin^2\theta_W) \cos 2\beta & m_t(A_t^* - \mu \cot\beta) \\ m_t(A_t - \mu^* \cot\beta) & M^2_0 + m_t^2 + M^2_u Q_u \sin^2\theta_W \cos 2\beta \end{pmatrix} \tag{10}$$

we know that $\sum_{i=1}^2 m_{\chi_i} U(i, 2)V(i, 2) = \mu$ and $D_{i21}D_{i11}^* = \frac{m_\tau}{m_{t1} - m_{t2}} (A_t - \mu^* \cot\beta)$. Therefore, $A_t$ itself is as important as $\mu$ for providing imaginary contributions to $C_{Q_i}$, in particular, for large $\tan\beta$. The similar conclusion holds also for $C_7$.

We have known well that $A_t$ at the EW scale mainly depends on $M_3$ at the GUT scale through RGE effects\cite{19, 20}. In fact there exists a quasi fixed point which shows the ratio of $A_t$ at $m_Z$ scale to $M_3$ at GUT scale to be about $-1.6$ provided the Yukawa couplings of the third generation are large enough \cite{19}. Hence it is possible for $A_t$ to achieve a very large imaginary parts in the non-universal gaugino mass models, in contrast to the case of mSUGRA. Especially in large $\tan\beta$ case, $|\mu|$ is limited by EDM data to be less than $\pi/6$, so $A_t$ plays a more important role in CP Violation than $\mu$.

To study the effect of large $\phi_3$ (hence large $\phi_{A_t}$) on $C_7$, we notice that essentially $A_t$ is multiplied by $\mu$. Changing the sign of $A_t$ has the same effects of switching the sign of $\mu$ and switching the sign of $\mu$ results in a sign change in $C_7$ (because in most of the parameter space $\mu$ are much larger than the non-diagonal terms in eq.(9)), so if $\phi_3$ is in $[\pi/2, 3\pi/2]$ (hence $\phi_{A_t}$ in $[-\pi/2, \pi/2]$) and $\phi_\mu$ in $[-\pi/2, \pi/2]$), supersymmetry contributions give enhancement to $ReC_7$ so that it is hard to satisfy the $B\to X_s\gamma$ constraints: $2.4\times10^{-4} < Br(B\to X_s\gamma) < 4.5\times10^{-4}$ \cite{21}. Similar situation occurs for $\phi_3$ being in $[\pi/2, 2\pi/2]$) and $\phi_\mu$ in $[\pi/2, 2\pi/2]$. Since supersymmetry gives large contributions to $C_7$ only when $\tan\beta$ is large we shall focus on large $\tan\beta$ case in the following. From the above analysis of EDMN we know that for somewhat large $\phi_\mu$(say around $\mp0.4$) cancellations happen for $\phi_3$ near $\pm\pi/2$. With this kind of phases of $M_3$ and $\mu$ SUSY contributions to $C_7$ are almost totally imaginary. So the value of $\Im C_7$ is constrained to be very small by the branching ratio of $B\to X_s\gamma$. One way to avoid the $B\to X_s\gamma$ constraint is to make use of the cancellation happened at $\phi_3$ around $\pm\pi/4$ and $\phi_\mu$ about $\mp0.2$. With such kind of phases and a low mass spectrum ($|M_2|$ and $|M_3|$ around 150 GeV) the real part of SUSY contributions to $C_7$ cancels those from W-top and charged Higgs-top loops. The real part can even be cancelled to be near zero and only a large imaginary part of $C_7$ remains. Another way is to just suppress the total SUSY contributions to $C_7$, i.e., to make the mass spectrum heavier but still less than 1 Tev ($|M_2|$ and $|M_3|$ larger than about 300 GeV and less than about 500 GeV). In the former case $A_{CP}^0$ can reach order 1% for $B\to X_s\ell^+\ell^-$ and $B\to X_s\tau^+\tau^-$, because of their larger Yukawa coupling there are great
enhancement of branching ratio$^{[17]}$, so that $A_{CP}^1$ are smaller than that for $B \to X_s e^+ e^-$. In the later case (we shall call it as region A hereafter) it can only be a few thousandth at most, i.e., the same order as that in SM.

As pointed out above, the effects of large $φ_3$ on $C_{Q_i}$ are similar to those on $C_7$. In large tan$β$ case, for $φ_μ \sim ±0.2$ and $φ_3 \sim ±π/4$ or $φ_μ \sim π ± 0.2$ and $φ_3 \sim π ± π/4$ and with $|M_2|$ and $|M_3|$ around 150 GeV (which we shall call as region B for simplicity), which are allowed by the EDME and EDMN limits, Im$C_{Q_i}$ reaches maxima. In small tan$β$ case, although the constraints of EDMs on $φ_μ$ and $φ_3$ are relaxed the magnitude of Im$C_{Q_i}$ is very small since $C_{Q_i}$ is proportional to $m_l tan^2 β$ (even $m_l tan^3 β$ in some regions of the parameter space). Therefore, we expect the significant CP violation in large tan$β$ case.

From eqs.(5) and (6), $A_{CP}^2$ can be rewritten as

$$A_{CP}^2 = \frac{E(\hat{s})D(\hat{s}) - \overline{E}(\hat{s})D(\hat{s})}{E(\hat{s})\overline{D}(\hat{s}) + \overline{E}(\hat{s})D(\hat{s})}$$

For $l=e$, the difference between $E(\hat{s})$ and $\overline{E}(\hat{s})$ can be neglected (as it is proportional to lepton mass square, see eq.(6)). So $A_{CP}^2$ for $l=e$ can be reduced to

$$A_{CP}^2 = \frac{D(\hat{s}) - D(\hat{s})}{\overline{D}(\hat{s}) + D(\hat{s})}$$

that is exactly the opposite of $A_{CP}^1$. The same conclusion can be drawn for $l=μ, τ$ in small tan$β$ case due to smallness of $C_{Q_i}$. On the other hand, for $l=τ$ in large tan$β$ case, $|E(\hat{s}) - \overline{E}(\hat{s})|$ can be more important than $|D - \overline{D}|$ and consequently one has approximately

$$A_{CP}^2 = \frac{E(\hat{s}) - \overline{E}(\hat{s})}{E(\hat{s}) + \overline{E}(\hat{s})}$$

Thus, it is proportional to Im$C_{Q_1}$. Therefore, in region B where $C_{Q_i}$ reach the maxima, $A_{CP}^2$ can be over 50%. In region A, $C_{Q_i}$ are less important and $A_{CP}^2$ can reach about 5% at most.

The correlation between $A_{CP}^2$ and EDME (or EDMN) in region A is plotted in fig.3 (note that we choose $\hat{s} = 0.76$ as representative in the figure). For $l=μ$, the magnitude of $A_{CP}^2$ can be estimated to be of order 1% at most in the large tan$β$ case. Numerical calculations prove this estimate.

Fig.4 shows the correlation of EDM constraints and $P_N$ of $B \to X_s τ^+ τ^−$ in region A. We can see in this figure that $P_N$ can reach more than 15 percent. In region B, as $C_{Q_i}$ are much larger the numerator in eq.(7) is increased a lot. But the denominator in eq.(7) (and consequently the branching ratio of $B \to X_s τ^+ τ^−$) is also greatly enhanced in this region, so $P_N$ is just about 15%, i.e., not larger than the magnitude that can be achieved in region A. Situations for muon are similar and because of its much smaller Yukawa coupling the magnitude of $P_N$ can only reach about 6%. For electron $P_N$ is negligibly small, due to its negligibly small mass. An important feature that can be seen from fig.3 and fig.4 is that the magnitudes of $A_{CP}^2$ and $P_N$ will not be reduced if EDM constraints improved. That is because the regions of parameter space where EDM constraints are satisfied are of width about $π/20$ for $φ_3$ and about $π/4$ for $φ_{A_0}$ (adjustment needed), while $C_{Q_i}$ do not change sharply within these regions.

In summary, we have analyzed the constraints of electric dipole moments of electron and neutron on the parameter space in supergravity models with nonuniversal gaugino masses. It
is shown that with a light sparticle spectrum, the sufficient cancellations in the calculation of EDMs can happen due to the presence of the two new phases arising from complex gaugino masses, in addition to the phases $\phi_{\mu}$ and $\phi_{A_{0}}$. With appropriate correlation between $\phi_{\mu}$ and $\phi_{1}$ (for EDME) or $\phi_{3}$ (for EDMN) as well as an appropriate choice of $\phi_{A_{0}}$, cancellations can occur and all phases can be order of one in the small $\tan\beta$ case and all phases but $\phi_{\mu}$ ($|\phi_{\mu}| \sim \pi/6$) order of one in the large $\tan\beta$ case. This is in contrast to the case of mSUGRA where in the parameter space where cancellations among various SUSY contributions to EDMs happen $\phi_{\mu}$ must be less than $\pi/10$ for small $\tan\beta$ and $\mathcal{O}(10^{-2})$ for large $\tan\beta$. And our analysis show that the branching ratio of $B \rightarrow X_s \gamma$ gives an extra constraint on the phases for large $\tan\beta$ case with light mass spectrum. We have calculated direct CP asymmetries and the T-odd normal polarization of lepton in $B \rightarrow X_s l^{+} l^{-}$ in the regions of the parameter space in the models where the constraints from EDMs as well as $B \rightarrow X_s \gamma$ are satisfied. It is shown that the results for $A_{1 \text{CP}}$ are similar to those in mSUGRA if the mass spectrum is relatively heavier($|M_2|,|M_3| \sim 300\text{GeV}$) and it is also true for $A_{2 \text{CP}}$ in the small $\tan\beta$ case. The former is due to the constraint from $B \rightarrow X_s \gamma$ and the latter is due to smallness of the contributions from exchanging neutral Higgs bosons in the small $\tan\beta$ case. However, in the large $\tan\beta$ case, $A_{2 \text{CP}}$ can reach 1% for $l = \mu$ and is a few percent in most of allowed regions and can reach 50% in some allowed regions for $l=\tau$. $A_{2 \text{CP}}$ for $l = e$ is approximately equal to $A_{1 \text{CP}}$ even in the large $\tan\beta$ case. $P_N$ can reach 6% for $l = \mu$ and is in the range from 1% to 15% in most of the allowed regions for $l=\tau$ in the large $\tan\beta$ case. That is, there is a significant enhancement compared to the mSUGRA in which $P_N$ only can reach about 0.5% for $l=\mu$ and about 5% for $l=\tau$. In the small $\tan\beta$ case the results are similar to those in mSUGRA. For $l=e$, $P_N$ is negligibly small, as it should be.

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Fig 1. EDME($10^{-27}$) as functions of $\phi_1$ from 0 to $2\pi$. (a) $\tan\beta=2$ and lines from below refer to $\phi_\mu=1.4, 1.0, 0.6, 0.2, -0.2, 0.6, 1.0, 1.4$. (b) $\tan\beta=30$ and six lines from below refer to $\phi_\mu=0.5, 0.3, 0.1, -0.1, 0.3, 0.5$. Phases of $M_2, M_3$ and $A_0$ are chosen to be zero. Other parameters are chosen such that $M_0=|M_2|=400\text{Gev}$, $|M_1|=|M_3|=500\text{Gev}$, $|A_0|=800\text{Gev}$. 
Fig. 2 Correlation between the allowed regions of $\phi_\mu$ and $\tan \beta$ with $M_0 = |M_2| = 400$ GeV, $|M_1| = |M_3| = 500$ GeV, $|A_0| = 800$ GeV.
Fig. 3 Correlation of EDMs and $A_{CP}^2$ in region A.
Fig. 4 Correlation of EDMs and $P_N$ in region A.