EXAMINING THE CANCELLATION MECHANISM OF NEUTRON EDM IN A MODEL WITH DILATON-DOMINATED SUSY BREAKING

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ABSTRACT

We examine the cancellation mechanism between the different contributions to the electric dipole moment of the neutron in a model with dilaton-dominated SUSY breaking. We find these accidental cancellations occur at few points in parameter space. For a wide region of this space we must constrain the phase of $\mu$ to be of order $10^{-1}$ and have the phases of $A$ and $\mu$ strongly correlated in order to have small neutron EDM. Moreover, we consider the indirect CP violation parameter $\varepsilon$ in this region where the electric dipole moment is less than the experimental limit and find that we can generate $\varepsilon$ of order $10^{-6}$.

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1. Introduction

It is well known that in supersymmetric theories there are new possibilities for CP violation. In particular, the soft SUSY breaking terms contain several parameters that may be complex, as can also be the $\mu$-parameter. These phases can cause at one loop level an electric dipole moment (EDM) for the quarks and leptons, and therefore also for the neutron. It has been known for a long time that in a generic SUSY model the contributions to the neutron electric dipole moment are larger than the experimental limit $1.1 \times 10^{-25}$ e cm unless either the new ‘SUSY phases’ are tuned to be of order $10^{-3}$, or the SUSY masses are of order a TeV. Recently it has been suggested that a natural cancellation mechanism exists whereby the electric dipole moment of the neutron may be made small without such fine-tuning. In this paper we examine this possibility in the context of a concrete model of SUSY breaking, namely dilaton-dominated breaking.

In the minimal supersymmetric standard model there are really only two new CP-violating phases. This can be seen as follows. The parameters $M, A$ and $B$ and $\mu$ can be complex. But of these four phases only two are physical. First, by an R-rotation with R-charge $Q_R = 1$ for lepton and quark superfields and $Q_R = 0$ for the vector and the Higgs superfields, the gaugino mass parameter $M$ can be made real. Second, $B\mu$ can be made real by a change of phase of the Higgs superfield. This ensures that the Higgs vacuum expectation values are real. The remaining phases cannot be defined away and violate CP. One is in $A = A_0 e^{i\phi_A}$ and the other in $B = B_0 e^{i\phi_B}$. The $\mu$ parameter then has a fixed phase $\mu = \mu_0 e^{-i\phi_B}$. In any phase convention

$$\phi_A = \arg(AM^*) \quad \phi_B = \arg(BM^*).$$

The fact that there are two phases can also be seen as follows. In the absence of the above soft SUSY breaking terms and the bilinear $\mu$-term of the superpotential there are two additional global $U(1)$ symmetries $[\mathbb{P}]$ an R-symmetry and a Peccei-Quinn symmetry. One finds that there are three independent combinations of the four parameters which are invariant under both $U(1)$’s, but only two of their phases are independent, and can be chosen to be $\phi_A$ and $\phi_B$.

It is these phases which must be of order $10^{-3}$ or less if the neutron EDM is not to be too large and the SUSY masses are not unnaturally heavy. The common choice is to
assume that $\phi_A$ and $\phi_B$ are identically zero. However, unless there is a symmetry which implies that these phases vanish or are very small at the unification scale, it would be unnatural to assume that they are. Clearly the SUSY phases problem is a problem of SUSY breaking since the relevant phases originate from SUSY breaking terms.

The new suggestion for suppressing the EDM of the neutron is made in Ref. [2] where chromoelectric and purely gluonic operator contributions are taken into account. It is proposed that there are internal cancellations among the various components contributing to the EDM which allows for the existence of large CP violating phases. However the EDM analysis for the minimal supergravity model [3] given in that paper actually shows that there is only a tiny region in the parameter space where these cancellations occur. Moreover, a heavy SUSY spectrum, of order TeV and small SUSY phases of order $10^{-1}$ are considered there. In fact, at these values the electric dipole moment is already suppressed without any cancellation. The region where large cancellations occur is bigger in Ref. [4] where a generic supersymmetric standard model is considered although all the superpartner masses are kept light. The reason for this is that there are more parameters to be adjusted to give a large cancellation between the different EDM contributions.

We argue that this is an accidental cancellation and is as much a fine tuning as the previously known mechanisms for suppression the neutron EDM that use small phases or a heavy SUSY spectrum. In this paper we examine this cancellation mechanism in a string-inspired model where the supersymmetry is broken in the hidden sector by the vacuum expectation values of the dilaton and/or moduli fields. The soft SUSY breaking terms of this scenario are very constrained and they are given in terms of two parameters only. Indeed, we find that such a constrained model does not allow for a large cancellation, the electric contribution to EDM is the dominant one, and the only way to suppress it is either by making the phase small or making the masses of order a TeV.

The paper is organized as follows. In section 2 we review the formulae of the soft SUSY breaking terms in the dilaton-dominated SUSY breaking model. In section 3 we examine the cancellation of the electric dipole moment contributions. It is found that this constrained model allows for this accidental cancellation only at few points in the parameter space. Section 4 is devoted to studying the indirect CP violation parameter $\varepsilon$ in the region where the electric dipole moment is smaller than $10^{-25}e\text{ cm}$. Finally, we give our conclusions in section 5.
2. Dilaton Dominated SUSY Breaking

We consider an example of the string inspired model where the dilaton $S$ and overall modulus field $T$ contribute to SUSY breaking and the vacuum energy vanishes. In this scenario, the soft masses $m_i$ and the gaugino masses $M_a$ are written as

\[ m_i^2 = m_{3/2}^2 (1 + n_i \cos^2 \theta) , \]
\[ M_a = \sqrt{3} m_{3/2} \sin \theta e^{-i\alpha_S} , \]

where $m_{3/2}$ is the gravitino mass, $n_i$ is the modular weight of the chiral multiplet and $\sin \theta$ corresponds to a ratio between the $F$-terms of $S$ and $T$. For example, the limit, $\sin \theta \to 1$, corresponds to the dilaton-dominated SUSY breaking. Here the phase $\alpha_S$ originates from the $F$-term of $S$. Similarly the $A$-parameters are also written as

\[ A_{ijk} = -\sqrt{3} m_{3/2} \sin \theta e^{-i\alpha_S} - m_{3/2} \cos \theta (3 + n_i + n_j + n_k) e^{-i\alpha_T} , \]

where $n_i, n_j$ and $n_k$ are the modular weights of the fields that are coupled by this $A$ term. One needs a correction term in eq (3) when the corresponding Yukawa couplings depend on moduli fields. This correction depends on the derivative of the Yukawa couplings with respect to the moduli field, and is therefore small since the Yukawa couplings are constants or tend exponentially to constants. So we ignore them. Here the phase $\alpha_T$ originates from the $F$-term of $T$.

Finally, the magnitude of the scalar bilinear soft breaking term $B \mu H_1 H_2$ depends on how the $\mu$-term is generated. Therefore here we take $\mu$ and $B$ as free parameters and we fix them by requiring successful electroweak symmetry breaking.

Thus, gaugino masses and $A$-terms as well as the $B$-term are, in general, complex. We have a degree of freedom to rotate $M_a$ and $A_{ijk}$ at the same time. Here we use the basis where $M_a$ is real. Similarly we rotate the phase of $B$ so that $B \mu$ itself is real. In this basis, the phases of $B$ and $\mu$ satisfy $\phi_B = -\phi_\mu = \arg(BM^*)$. In $A$-terms in the above basis, there remains only one independent phase, namely, $\alpha' \equiv \alpha_T - \alpha_S$. i.e., $A_{ijk} = A_{ijk}(\alpha')$. We assume the following modular weights for quark and lepton fields

\[ n_Q = n_U = n_{H_1} = -1 \]

and

\[ n_D = n_L = n_E = n_{H_2} = -2 . \]
As will be seen later, this assumption is favorable for electroweak breaking. Under this assumption we have $A_t = A_b = A$, where $A$ is given by

$$A = -\sqrt{3}m_{3/2}\sin\theta + m_{3/2}\cos\theta e^{-i\alpha'}.$$  \hfill (4)

Given the boundary conditions in eqs. (1-3) at the compactification scale, we determine the evolution of the couplings and the mass parameters according to their one loop renormalization group equation in order to estimate the mass spectrum of the SUSY particles at the weak scale. The radiative electroweak symmetry breaking imposes the following conditions on the renormalized quantities:

$$m_{H_1}^2 + m_{H_2}^2 + 2\mu^2 > 2B\mu, \hfill (5)$$

$$(m_{H_1}^2 + \mu^2)(m_{H_2}^2 + \mu^2) < (B\mu)^2, \hfill (6)$$

$$\mu^2 = \frac{m_{H_1}^2 - m_{H_2}^2\tan^2\beta}{\tan^2\beta - 1} - \frac{M_Z^2}{2}, \hfill (7)$$

and

$$\sin 2\beta = \frac{-2B\mu}{m_{H_1}^2 + m_{H_2}^2 + 2\mu^2}; \hfill (8)$$

where $\tan\beta = \langle H_2^0 \rangle / \langle H_1^0 \rangle$ is the ratio of the two Higgs VEVs that gives masses to the up and down type quarks and $m_{H_1}, m_{H_2}$ are the two soft Higgs masses at the electroweak scale. We take here $\tan\beta = 3$, using the above equations we can determine $|\mu|$ and $B$ in terms of $m_{3/2}, \theta$ and $\alpha'$. The phase of $\mu$ ($\phi_\mu$) remains undetermined.

3. The Calculation of Electric Dipole Moment

In supersymmetric theories, the EDM of the quark receives contributions at the one loop level from diagrams in which the charginos, neutralinos or gluinos are exchanged together with the squarks. The EDM operator changes the chirality of the quark. The gaugino couples the quark to the squark with the same chirality via the gauge interactions, while the Higgsino couples the quark to the squark with the opposite chirality via the Yukawa interactions. In a supersymmetric model, this can happen in two ways, either through the gaugino-Higgsino mixing or through the mixing of $\tilde{q}_L$ and $\tilde{q}_R$. The chargino loop diagrams involve both possibilities, while gluino-loop diagram involves the second one only.
Moreover, besides the quark electric dipole, there are additional operators that contribute to the EDM. They are the gluonic operator $O_G = -\frac{1}{6}f^{abc} G_a G_b \tilde{G}_c$ and the quark chromoelectric dipole moment $O_q = \frac{i}{4} q \sigma_{\mu \nu} \gamma_5 T^a G^{\mu \nu}$, where $T^a$ are the generators of $SU(3)$. In Ref. [9], it was estimated that the gluonic operator is the smallest contribution when all the mass scales are taken to be equal. However, in Ref. [3], it was claimed that the contributions of the chromoelectric and the purely gluonic operators can be comparable to the contribution of the electric dipole operator. The calculation of the gluino contributions to the quark EDM in that reference shows that they are given by

$$d_{q - \text{gluino}}^E/e = -\frac{2\alpha_s}{3\pi} \sum_{k=1}^{2} \text{Im}(\Gamma_q) \frac{m_{\tilde{g}_k}}{M_{\tilde{q}_k}^2} Q_d B\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right), \quad (9)$$

and

$$d_{q - \text{gluino}}^C = \frac{g_s \alpha_s}{4\pi} \sum_{k=1}^{2} \text{Im}(\Gamma_q) \frac{m_{\tilde{g}_k}}{M_{\tilde{q}_k}^2} C\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right). \quad (10)$$

While the chargino loop contributions are given by

$$d_{u - \text{chargino}}^E/e = -\frac{\alpha_{EM}}{4\pi \sin^2 \theta_W} \sum_{k=1}^{2} \sum_{i=1}^{2} \text{Im}(\Gamma) \frac{m_{\chi_i}}{M_{\chi_i}^2} [Q_d B\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right) + (Q_u - Q_d) A\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right)], \quad (11)$$

$$d_{d - \text{chargino}}^E/e = -\frac{\alpha_{EM}}{4\pi \sin^2 \theta_W} \sum_{k=1}^{2} \sum_{i=1}^{2} \text{Im}(\Gamma) \frac{m_{\chi_i}}{M_{\chi_i}^2} [Q_d B\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right) + (Q_d - Q_u) A\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right)], \quad (12)$$

and

$$d_{\tilde{q} - \text{chargino}}^C = \frac{-g_s^2 g_s}{16\pi^2} \sum_{k=1}^{2} \sum_{i=1}^{2} \text{Im}(\Gamma) \frac{m_{\chi_i}}{M_{\chi_i}^2} B\left(\frac{m_{\tilde{q}_k}^2}{M_{\tilde{q}_k}^2}\right). \quad (13)$$

Finally, the coefficient of the CP violating dimension six operator is given by

$$d^C = -3\alpha_s m_t \left(\frac{g_s}{4\pi}\right)^3 \text{Im}(\Gamma^{12}_t) \frac{m_{t_1}^2 - m_{t_2}^2}{m_{\bar{t}_1}^2} H\left(\frac{m_{t_1}^2}{m_{\bar{t}_1}^2}, \frac{m_{t_2}^2}{m_{\bar{t}_2}^2}, \frac{m_t^2}{m_{\bar{t}_1}^2} \right). \quad (14)$$

In these equations, $M_{\tilde{q}_k}$ are the masses of the corresponding scalar particle running in the loop. The functions $A, B, C, H$ and $\Gamma$ can be found in Ref. [3]. Another contribution to the quark electric dipole moment comes from the neutralino loops. However, as explained in detail in Ref. [4], this contribution is small and can be neglected in the electric dipole moment analysis. There are several reasons for that. For instance, the elements of the neutralino diagonalizing matrix yield smaller imaginary parts than the elements of the chargino matrices. Also the chargino contribution is enhanced due to large values of the loop function compared to the function in the neutralino expression.
The analysis of the EDM above is at the electro-weak scale and it must be evolved down
to the hadronic scale via the renormalization group evolution as explained in Ref [3]. The
contribution to the neutron EDM coming from the $d^E$ is obtained by using the $SU(6)$ quark model which gives
\[ d_n = \frac{1}{3}(4d_d - d_u). \]
The contributions to $d_n$ coming from the $d^C$ and $d^G$ are estimated using naive dimensional
analysis [9]. As stated above, we determine $\mu$ and $B$ from the electroweak breaking
conditions. Hence the parameter space of this model consists of only $m_3/2$, $\theta$, $\phi_A$ and $\phi_{\mu}$.
In this scenario we take $\phi_{\mu}$ to be unconstrained while $\phi_A$ is given in terms of $\alpha'$ at GUT
scale. It is worth mentioning that $\phi_A$ at electroweak scale is less than $\alpha'$ since the real
part of $A$ is running faster than the imaginary part. In figure (1) we show the values of the
total EDM of the neutron as function of $m_3/2$, for $\cos^2 \theta \simeq 1/2$ and $\phi_{\mu} = \alpha' = \pi/2$.

![Figure 1. The total EDM as function of $m_3/2$, for $\phi_{\mu} = \alpha' = \pi/2$.](image)

From this figure we note that with these large phases we still have to have $m_3/2 > \text{TeV}$
to obtain value for the EDM less than the experimental limit $1.1 \times 10^{-25}$e cm. (i.e., all
the scalar masses are of order TeV). The experimental limit on the chargino mass puts
a lower bound on $m_3/2$, namely $m_3/2 > 100$ GeV. In the region of $m_3/2$ between 100 and
115 GeV the chromoelectric contribution to the EDM exceeds the electric contribution
and at one point in the parameter space these two contributions are approximately equal
so that we find the total EDM of neutron is less than the experimental limit. However, as
we have said, this is happening in a very tiny region in the parameter space. Now, to see
the effect of each of the phases on the EDM values we plot in figure (2) the EDM versus
the $\phi_\mu$ for two values of $\alpha'$. The solid line corresponds to the EDM where $\alpha' = 0$, while
the dashed line corresponds to the EDM where $\alpha'$ is of order $\pi/2$.

![Figure 2](image)

**Figure 2.** The total EDM versus $\sin \phi_\mu$. The solid line corresponds $\alpha' = 0$
and the dashed one corresponds to $\alpha' = \pi/2$.

This figure shows that the values of the EDM of the neutron in the case of non-vanishing
$\phi_A$ is less than in the case of $\phi_A = 0$. This is an interesting result since by adjusting the
phase of $\mu$ we can leave the phase of $A$ large and have the EDM of the neutron less
than the experimental bound. This could generate, as we examine in the next section, a
sizable value for the CP violation $\varepsilon$ parameter. For fixed value of $m_{3/2} \simeq \mathcal{O}(100)$ GeV
and $\cos^2 \theta \simeq 1/2$, the behaviour of the neutron EDM in the phase plane $\phi_\mu - \alpha'$ can
be understood as follows. The interesting region of this plane (i.e., where the neutron
EDM is less than the experimental limit) is where the phases are opposite in sign and are
strongly correlated. For instance, the constraint obtained on the phase $\phi_\mu$ is more severe
for $\alpha' = 0$ than the one obtained for non vanishing $\alpha'$, as Figure (2) confirms.

In fact, in our model we have left the $B$ parameter free and determined it at electroweak
scale from the radiative breaking conditions so that the phase of $\mu$ in this case is a free
parameter. However in Ref. 3 three sources for the $B$ parameter were considered, labeled
by $B_Z$, $B_\mu$ and $B_\lambda$. The source of $B_Z$ is the presence of certain bilinear term in the
Kähler potential which can naturally induce a $\mu$-term of order $m_{3/2}$ after SUSY breaking.
An alternative mechanism to generate a $B$-term in a scalar potential is to assume that superpotential $W$ includes a SUSY mass term $\mu(S, T)H_1H_2$ induced by a non-perturbative effect, then a $B$-term is automatically generated and it is called $B_\mu$. Also, it has been pointed out that the presence of a non-renormalizable term in the superpotential $\lambda W H_1H_2$ yields dynamically a $\mu$ parameter when $W$ acquires VEV. The corresponding $B$-term is denoted by $B_\lambda$. Furthermore, it was shown in Ref. [10] that $B_\mu$ is the favorable choice to realize the radiative electroweak symmetry breaking. In mixed dilaton/modulus SUSY breaking the $B_\mu$ is given by

$$B_\mu = m_{3/2}[-e^{i\alpha_S} - \sqrt{3}\sin \theta - \sqrt{3}\cos \theta(3 + n_{H_1} + n_{H_2})e^{-i\alpha'}],$$

with the values of the modular weights we gave in section 2 this formula leads to

$$B_\mu = m_{3/2}[-e^{i\alpha_S} - \sqrt{3}\sin \theta],$$

and therefore

$$\tan(\phi_\mu) = \frac{\sin \alpha_S}{\cos \alpha_S + \sqrt{3}\sin \theta}.$$  

Thus, the phase of $\mu$ in this case is constrained such that $|\phi_\mu| \leq \pi/4$ (from $m_{H_2}$ with the modular weights we are assuming (see eq. (1)) one has $\cos^2 \theta < 1/2$. This implies that $\sin \theta > 0$ and even close to one). In other cases of generating the $B$-term, as well as for other choices for the modular weights, we could get a stronger bound on the phase of $\mu$. So it may be natural to have such a small phase of $\mu$. Some further comments are in order. It is important to note that in this model a large value of $\alpha'$ leads to a large value of $|A|$; and this is required to reduce the EDM value, as was also found in Ref. [3]. Moreover, the sign of the gluino contribution is reversed for large value of $\alpha'$, and a destructive interference between it and the chargino contribution occurs. This can be seen by looking to the lowest approximation of the gluino contribution to the electric dipole moment

$$d_{q-glino}^E/e \simeq \frac{2\alpha_s}{3\pi} m_{\bar{q}} Q_{\bar{q}} \frac{m_{\mu}}{M_{\mu}^2}(|A_q| \sin \phi_{A_q} + |\mu| \sin \phi_\mu R_q)B(\frac{m_{\mu}^2}{M_{\mu}^2}).$$

Thus, for $|\phi_\mu| < 10^{-1}$ and $|\alpha'| = \pi/2$, and $\phi_\mu$ and $\alpha'$ having opposite signs, the sign of the gluino contribution is opposite to the sign of the chargino contribution, and consequently there is some cancellation at that point. It is important to note that this will not be the case if the phase of $\mu$ is larger than $10^{-1}$. Finally, with the phase $\alpha'$ of order $\pi/2$ and the phase of $\mu$ of order $10^{-1}$ we find the limit on the EDM of the neutron is satisfied for a large region of the parameter space. We have to mention that the cancellation between
the electric and chromoelectric contributions helps but it is not the reason for reaching this region. The reason is that the value of the electric contribution to the EDM of the neutron by itself is small.

In this connection it is very important to mention that in case of $\alpha' = \pi/2$, the phase of $A$ at GUT scale from eq(3) is of order $\pi/6$. This is the maximum phase of the $A$-term in this model. Moreover, due to the running from the GUT scale to weak scale this phase is reduced more and becomes of order $\pi/20$, i.e., it is of the same order as $\phi_\mu$. In fact, it is a feature of all supersymmetric models that the phase of the trilinear couplings at the weak scale is less than at the GUT scale due to the difference in the running between the real and imaginary parts of $A$. On the other hand, the phase of $\mu$ does not run, so that it has to be of order $10^{-1}$ at the GUT scale too. It is the smallness of these phases seperately, and not some cancellation between phases, that is chiefly responsible for the smallness of the EDM.

We have not discussed the electron EDM, since the cancellation mechanisms we have been considering (between gluino and chargino diagrams, and between chromoelectric and electric dipole operators) obviously do not apply to it. Even in the region of parameter space where the neutron EDM is suppressed, there is no reason to expect the electron EDM to be suppressed, and therefore the electron EDM would remain a problem, especially as the sleptons are lighter than the squarks.
4. The parameter $\varepsilon$

It is very important to show if it is possible to generate a sufficiently large value of $\varepsilon$ in the region of the parameter space that we found in the last section leads to values for the EDM of the neutron less than the experimental limit. It is not important whether the values of the phases at the large scale or at the low scale are of order one or $10^{-1}$. What is really important is that these phases alone can generate a sizable contribution to the CP violation processes.

The value of the indirect CP violation in the Kaon decays, $\varepsilon$, is defined as

$$\varepsilon = \frac{e^{i\frac{\pi}{4}}\text{Im}M_{12}}{\sqrt{2}\Delta m_K}, \quad (17)$$

where $\Delta m_K = 2\text{Re}(K^0|H_{eff}|\bar{K}^0) = 3.52 \times 10^{-15}$ GeV. The amplitude $M_{12} = \langle K^0|H_{eff}|\bar{K}^0 \rangle$. The relevant supersymmetric contributions to $K^0 - \bar{K}^0$ are the gluino and the chargino contributions, (i.e., the transition proceeds through box diagrams exchanging gluino-squarks and chargino-squarks). It is usually expected that the gluino is the dominant contribution. However, as we will show, it is impossible in the case of degenerate $A$-terms that the gluino gives any significant contribution to $\varepsilon$ when the CKM matrix is taken to be real even with large phase of $A$.

The amplitude of the gluino contribution is given in Ref. [11] in terms of the mass insertion $\delta_{AB}$ defined by $\delta_{AB} = \frac{\Delta_{AB}}{\bar{m}}$ where $\bar{m}$ is an average sfermion mass and $\Delta$ is off-diagonal terms in the sfermion mass matrices. The mass insertion to accomplish the transition from $\bar{d}_{iL}$ to $\bar{d}_{jL}$ ($i, j$ are flavor indices) is given by [12]

$$\begin{align*}
(\Delta_{LL}^d)_{ij} &\approx -\frac{1}{8\pi^2} \left[ \frac{K^\dagger(M_u^{\text{diag}})^2K}{v^2\sin^2\beta} \ln\left(\frac{M_{GUT}}{M_W}\right) \right] (3\bar{m}^2 + |X|^2), \quad (18) \\
(\Delta_{LR}^d)_{ij} &\approx -\frac{1}{8\pi^2} \left[ \frac{K^\dagger(M_u^{\text{diag}})^2K M_d}{v^2\sin^2\beta\cos\beta} \ln\left(\frac{M_{GUT}}{M_W}\right) X, \quad (19) \\
(\Delta_{RL}^d)_{ij} &\approx -\frac{1}{8\pi^2} \left[ \frac{M_d K^\dagger(M_u^{\text{diag}})^2K}{v^2\sin^2\beta\cos\beta} \ln\left(\frac{M_{GUT}}{M_W}\right) X, \quad (20) \\
(\Delta_{RR}^d)_{ij} &\approx 0, \quad (21)
\end{align*}$$

where $X = A_d - \mu \tan\beta$. It is clear that $\Delta_{ij}$ in general are complex due to the complexity of the CKM matrix, the trilinear coupling $A$ and $\mu$ parameter. Here we assume the vanishing of $\delta_{CKM}$ to analyze the effect of the SUSY phases. We notice that $(\Delta_{LL}^d)_{12}$ is proportional to $|X|^2$ i.e., it is real and does not contribute to $\varepsilon$ whatever the phase of $A$ is. Moreover, the values of the $(\Delta_{LR}^d)_{12}$ and $(\Delta_{RL}^d)_{12}$ are proportional to $m_s$ and $m_d$, respectively.
hence they are quite small. Indeed in this case we find the gluino contribution to $\varepsilon$ is of order $10^{-6}$. The contribution is enhanced in case of non-degenerate $A$-terms [13].

For the chargino contribution the amplitude is given by [14]

$$\langle K^0|H_{\text{eff}}|\bar{K}^0\rangle = -\frac{G_F^2 M_W^2}{(2 \pi)^2} (V_{td}^* V_{ts})^2 f_K M_k \left( \frac{1}{3} C_1(\mu) B_1(\mu) + \frac{1}{24} C_2(\mu) B_2(\mu) + \frac{1}{24} C_3(\mu) B_3(\mu) \right).$$

(22)

The complete expression for these function can be found in Ref. [14]. Since we have $\tan \beta \simeq 3$ the value of $C_3$ is much smaller than $C_1$ since it is suppressed by the ratio of $m_s$ to $M_W$. However, by neglecting the flavor mixing in the squark mass matrix $C_1$ turned out to be exactly real [15]. The imaginary part of $C_1$ is associated to the size of the intergenerational sfermion mixings, thus it is maximal for large $\tan \beta$. In low $\tan \beta$ case, that we consider, the imaginary part of $C_1$ is very small, and the gluino contribution is still the dominant contribution $\varepsilon$. As explained, the EDM of the neutron constrains the phase of $\phi_\mu$ to be of order $10^{-1}$, where $\alpha' \simeq \pi/2$ and $\cos^2 \theta \simeq 1/2$. In this region we estimate the value of $\varepsilon$. It turns out that it is of order $10^{-6}$, which is less than the experimental value $2.26 \times 10^{-3}$.

5. Conclusions

We have examined the cancellation mechanism between the different contributions for the electric dipole moment of the neutron in a model with dilaton-dominated SUSY breaking. We found that this is an accidental cancellation and is as much a fine tuning as the previously known mechanisms for the suppression the EDM of the neutron. It occurs only at a few point in the parameter space of this model.

Furthermore, we studied the indirect CP violation parameter $\varepsilon$ and we showed that it is of order $10^{-6}$ only in the region where the EDM of the neutron is smaller than $10^{-25}$e cm.

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