Phases in the MSSM

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Abstract

The effect of CP violating phases in the MSSM on the relic density of the lightest super-symmetric particle (LSP) is considered. In particular, the upper limits on the LSP mass are relaxed when phases in the MSSM are allowed to take non-zero values and when the LSP is predominantly a gaugino (bino). Previous limits of $m_{\tilde{B}} \lesssim 250$ GeV for $\Omega h^2 < 0.25$ can be relaxed to $m_{\tilde{B}} \lesssim 650$ GeV. The additional constraints imposed by the neutron and electron electric dipole moments induced by these phases are also considered. Though there is some restriction on the phases, the bino mass may still be as large as $\sim 350$ GeV and certain phases can be arbitrarily large.

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It is well known that by considering the cosmological relic density of stable particles, one can establish mass limits on these particles. In the minimal supersymmetric standard model (MSSM), because of the many unknown parameters it is a significantly complicated task to set limits on the mass of the lightest particle, the LSP. The annihilation cross section will depend on several parameters which determine the identity of the LSP as well as parameters such as scalar masses which mediate LSP annihilations. In addition, there are (at least) two CP violating phases in the MSSM which play a role in determining the LSP relic abundance. As will be shown, when the phases are allowed to take their maximal values, mass limits on the LSP are greatly relaxed. However, constraints from the neutron electric dipole moment will restrict the degree to which these mass limits can be relaxed.

The mass and relic density of the LSP depends on a relatively large number of parameters in the MSSM. With some assumptions on these parameters, their number can be whittled down to 6: 1) if the gaugino masses are assumed to obey a GUT relation, they can be characterized by a single mass parameter, $M_2$; 2) the Higgs mixing mass, $\mu$; 3) the ratio of Higgs vevs, $\tan \beta$; 4) the soft (diagonal) sfermion mass parameters, $m_{\tilde{f}}$ (assumed here to be degenerate at the weak scale); 5) the soft trilinear terms, $A_f$ (also assumed equal at the weak scale); and 6) the soft bilinear, $B$. In general, all but the sfermion masses, $m_{\tilde{f}}$, are complex. However not all of the phases associated with these are parameters are physical. It is common to rotate away the phase of the gaugino masses, and to make $B\mu$ real, which ensures that the vacuum expectation values of the Higgs fields are real. For simplicity, generation mixing in the sfermion sector will be ignored though left-right mixing will be included. If furthermore we assume that all of the $A_f$’s are equal, and simply label them by $A$, we are left with two independent phases, the phase of $A$, $\theta_A$, and the phase of $\mu$, $\theta_\mu$. The phase of $B$ is fixed by the condition that $B\mu$ is real.

There has been a considerable amount of work concerning phases in the MSSM. For the most part these phases are ignored because they tend to induce large electric dipole moments for the neutron. If one considers the direct contribution to the edm from gluino exchange, one finds a dipole moment

$$d_n \simeq e \frac{\alpha_s m_u m_{\tilde{g}} |\overline{m}_t|}{\pi m_{\tilde{f}}^4} \sin \gamma_t \sim \text{few} \times 10^{-23} \sin \gamma_t \text{ ecm}$$

where the gluino mass, $m_{\tilde{g}}$ and $m_{\tilde{f}}$ as well as $|\overline{m}_t| \equiv A^* + \mu \cot \beta$ are all taken to be of order 100 GeV. The angle $\gamma_t \equiv \arg \overline{m}_t$ is clearly a function of both $\theta_\mu$ and $\theta_A$. To suppress the electric dipole moments, either large scalar masses (approaching 1 TeV) or small angles (of order $10^{-3}$, when all SUSY masses are of order 100 GeV) are required. For the most part, the community has opted for the latter, though the possibilities for large phases was recently considered in [4]. To reconcile large phases with small electric dipole moments, some of the...
sparticle masses are required to be heavy. In \cite{4}, it was first assumed that the SUSY phases were large (\(\pi/4\)), and by computing the various contributions to the edm of the neutron from neutralino, gluino, and chargino exchange (it was the latter that they found to be dominant) either large sfermion or neutralino masses (or both) were required. For example, Kizukuri and Oshimo \cite{4} found that for \(M_2 = \mu = 0.5\) TeV, and \(\tan \beta = 2\) the chargino contribution to the edm of the neutron exceeds the experimental upper bound \cite{5} of \(|d_n| < 1.1 \times 10^{25}\) emc unless, \(m_{\tilde{f}} \gtrsim 2\) TeV. When \(M_2 = \mu = 0.2\) TeV the limit on the sfermion masses becomes \(m_{\tilde{f}} \gtrsim 3\) TeV for \(\tan \beta = 2\) and \(m_{\tilde{f}} \gtrsim 6\) TeV for \(\tan \beta = 10\).

Figure 1: Typical contribution to the edm.

However, unless \(R\)-parity is broken and the LSP is not stable, one requires that sfermions be heavier than neutralinos, and if they are much heavier, this would result in an excessive relic density of neutralinos. Thus one must determine the relationship between potentially large phases in the MSSM and the relic density while remaining consistent with experimental bounds on the electric dipole moments. Here I will consider only neutralinos as the LSP. The mass limits will then depend sensitively on the masses of the scalars. For example when parameters are chosen so that the LSP is a gaugino (when the supersymmetry breaking gaugino masses \(M_{1,2}\) are taken to be smaller than the supersymmetric Higgs mixing mass, \(\mu\)), for a given LSP mass, the requirement that \(\Omega h^2 < 1/4\) places an upper bound on the sfermion mass (\(\Omega\) is the fraction of critical density and \(h\) is the hubble parameter scaled to \(100\) km/s/Mpc). The reason being that \(\Omega h^2 \propto 1/\langle \sigma v \rangle_A\), and \(\langle \sigma v \rangle_A \propto m_{\tilde{f}}^{-4}\) is the annihilation cross section. There is of course in addition, a dependence on the gaugino mass such that for \(m_{\text{gaugino}} < m_{\tilde{f}}\), a larger gaugino mass implies a larger annihilation cross section.

For small gaugino masses, the LSP is a photino, and one obtains a lower limit on the photino mass \cite{3}. At larger gaugino masses the LSP is a bino \cite{4} whenever \(M_2 < \mu\) and \(M_2 \gtrsim 200\) GeV; for smaller \(M_2\), the bino is still the LSP for large enough values of \(\mu\). The bino portion of the \(M_2-\mu\) LSP parameter plane is attractive, as it offers the largest possibility for a significant relic density \cite{4, 3, 9}. The complementary portion of the parameter plane, with \(\mu < M_2\), only gives a sizable density in a limited region, due to the large annihilation cross sections to \(W^+W^-\) and \(ZZ\) and due to co-annihilations \cite{10} with the next lightest neutralino (also a Higgsino in this case), which is nearly degenerate with the LSP \cite{11}. For the remainder of the discussion here, I will focus on binos as the LSP. In addition to resulting...
in a sizable relic density, the analysis is simplified by the fact that in the nearly pure bino region, the composition and mass of the LSP is not very sensitive to the new phases. However, as is shown, the relic density, which is determined primarily by annihilations mediated by sfermion exchange, is quite sensitive to the phases, $\gamma_t$ and $\gamma_b$.

Typically, bino annihilation is dominated by the $\tilde{B}\tilde{B} \rightarrow f\bar{f}$ channel. When the thermally averaged cross-section is expanded in powers of $T/m_{\tilde{B}}$ (see e.g. [12]), the cross section can be written as

$$\langle \sigma v \rangle_A = a + bT/m_{\tilde{B}}$$

where

$$a \propto \frac{g_4^4}{128\pi} \left( Y_L^2 + Y_R^2 \right)^2 \frac{m_f^2}{\Delta_f^2}$$

$$b \sim \frac{g_4^4}{128\pi} \left( Y_L^4 + Y_R^4 \right) \frac{m_{\tilde{B}}^2}{\Delta_f^2}$$

and $\Delta_f^2 = m_f^2 + m_{\tilde{B}}^2 - m_f^2$. In the above equations, $Y_{L(R)}$ is the hypercharge of $f_{L(R)}$. The expression for $b$ is significantly more complicated but what is important to note is that $b$ contains a piece which is proportional to $m_{\tilde{B}}^2$ whereas $a$ does not. This is a manifestation of the p-wave suppression of the annihilation cross-section for majorana fermions.

When common sfermion masses are assumed and when sfermion (left-right) mixing is ignored, there is an upper limit $m_{\tilde{B}} \lesssim 250$–300 GeV [7, 8] in order to ensure that $\Omega_{\tilde{B}}h^2 < 1/4$. (This upper limit is somewhat dependent on the value of the top quark mass. In [7, 13], we found that the upper limit was $\sim 250$ GeV for $m_t \sim 100$ GeV. When $m_t = 174$ GeV, the upper limit is 260 GeV. For $m_t \sim 200$ GeV, this limit is increased to $m_{\tilde{B}} \lesssim 300$ GeV. Furthermore, one should be aware that there is an upward correction of about 15% when three-body final states are included [14] which raises the bino mass limit to about 350 GeV. This latter correction would apply to the limits discussed below though it has not been included.) As the bino mass is increased, the sfermion masses, which must be larger than $m_{\tilde{B}}$, are also increased, resulting in a smaller annihilation cross section and thus a higher relic density. At $m_{\tilde{B}} \simeq 250$ GeV, even when the sfermion masses are equal to the bino mass, $\Omega h^2 \sim 1/4$. Note that this provides an upper bound on the sfermion masses as well, since the mass of the lightest sfermion is equal to the mass of the bino, when the bino mass takes its maximum value. One should also note that in the range $20$ GeV $\lesssim m_{\tilde{B}} \lesssim 300$ GeV, it is always possible to adjust $m_f$ in order to obtain $\Omega_{\tilde{B}}h^2 = 1/4$.

In [13], it was shown that the upper limit to the bino mass is sensitive to the level of sfermion mixing. The general form of the sfermion mass squared matrix is [13]

$$
\begin{pmatrix}
M_L^2 + m_f^2 + \cos 2\beta (T_3f - Q_f \sin^2 \theta_W) M_Z^2 & m_f \overline{m_f} e^{i\gamma_f} \\
m_f \overline{m_f} e^{-i\gamma_f} & M_R^2 + m_f^2 + \cos 2\beta Q_f \sin^2 \theta_W M_Z^2
\end{pmatrix}
$$
where $M_{L(R)}$ are the soft supersymmetry breaking sfermion mass which are assumed to be generation independent and generation diagonal and hence real. I also assume $M_L = M_R$. As discussed above, there is a non-trivial phase associated with the off-diagonal entries, which is denoted by $m_f(|m_f|e^{i\gamma_f})$, of the sfermion mass\(^2\) matrix, and

$$
|m_f|e^{i\gamma_f} = R_f \mu + A^* = R_f |\mu| e^{i\theta_f} + |A| e^{-i\theta_A}
$$

(6)

where $R_f = \cot \beta$ ($\tan \beta$) for weak isospin +1/2 (-1/2) fermions.

In general, one finds for the s-wave annihilation cross-section \([1]\),

$$
\tilde{a}_f = \frac{g'^4}{128\pi} \left| \frac{\Delta_1 (m_f w^2 - 2m_f w z + m_f z^2) + \Delta_2 (m_f x^2 + 2m_f x y + m_f y^2)}{\Delta_1 \Delta_2} \right|^2
$$

(7)

where $\Delta_i = m_f^2 + m_B^2 - m_f^2$ and the factors $x, y, w, z$ are determined from the bino-fermion-sfermion interaction Lagrangian,

$$
\mathcal{L}_{ff\bar{B}} = \frac{1}{\sqrt{2}} g' \left( Y_R \bar{f}_L \bar{f}_R \bar{F} + Y_L \bar{f}_L \bar{f}_R \bar{F} \right) + \text{h.c.}
$$

$$
= \frac{1}{\sqrt{2}} g' \left( \bar{f}_1 x P_L + \bar{f}_2 w P_L + \bar{f}_1 y P_R - \bar{f}_2 z P_R \right) \bar{B} + \text{h.c.}
$$

(8)

When phases are ignored, $x = Y_R \sin \theta_f$, $y = Y_L \cos \theta_f$, $w = Y_R \cos \theta_f$, $z = Y_L \sin \theta_f$. The angle $\theta_f$ relates the sfermion mass eigenstates to the interaction eigenstates by $\tilde{f}_1 = \tilde{f}_L \cos \theta_f + \tilde{f}_R \sin \theta_f$ and $\tilde{f}_2 = -\tilde{f}_L \sin \theta_f + \tilde{f}_R \cos \theta_f$.

When the off diagonal elements of the sfermion mass matrix vanish ($|m_f| = 0$), $\Delta_1 = \Delta_2$ and $x y = w z$ and the pieces proportional to $m_B$ in (7) cancel yielding the suppressed form of the s-wave cross-section given in (3). When sfermion mixing is included \([13]\), the limits, which now depend on the magnitude of the off-diagonal elements $m_f m_f$, are modified. In Figure 2, the upper and lower limits to the bino mass is shown as a function of $A_f$, which for fixed

Figure 2: Upper and lower limits to the bino mass as a function of $A_f$ for $\mu = -3000$ GeV.
\( \mu \) is equivalent to showing the limits as a function of the off diagonal element \( |\mathbf{m}| \). Because it is the combination \( m_j|\mathbf{m}| \) that appears in the off diagonal element, sfermion mixing really only affects the stops. Thus when mixing is important one of the stop masses is much less than the diagonal elements \( M_{L(R)} \). In order to ensure that the bino is lighter than the light stop, the diagonal elements must be pushed up relative to the no mixing case. This results in a reduction in the annihilation cross-section away from the no stop mixing region near \( A = 1500 \) GeV. The upper limit is asymmetric about this point as the couplings of the bino to \( \tilde{t}_1 \) and \( \tilde{t}_2 \) differ, and the lightest stop is either \( \tilde{t}_1 \) or \( \tilde{t}_2 \), depending on whether or not \( A \) is greater than or less than 1500 GeV. The lower limit on \( m_\tilde{B} \) is easily understood: when there is a lot of stop mixing and the diagonal elements are large, the annihilation into channels other than tops, is suppressed. If \( m_\tilde{B} < m_t \), the annihilation cross-section is too small and \( \Omega h^2 \) is too large. Thus we must require the top channel to be available and \( m_\tilde{B} > m_t \). Near the no stop mixing region, we have assumed that \( m_{\tilde{f}} \) is much less than \( 74 \) GeV \[16\]. In addition, we are able to obtain \( \Omega h^2 = 1/4 \) for a wide range of bino masses (as was the case when sfermion mixing was ignored by varying the the sfermion mass) as well as for a wide range of sfermion masses, by varying the magnitude of the off diagonal elements in the sfermion mass matrix. For fixed \( \mu \) this corresponds to varying \( A \). The upper limit to the bino mass is relaxed considerably when the phases are allowed to take non-zero values.

As we have just seen that though sfermion mixing has a large effect on the bino mass limits, the mass limit was not greatly increased. Because of the mass splitting among the sfermions, \( (\Delta_2 \gg \Delta_1) \) there is a reduction in the cross-section which compensates somewhat for the presence of the term proportional to \( m_\tilde{B} \) in \( a \). With non-zero phases, the enhancement in the cross-section is much greater. Even with \( \Delta_2 \simeq \Delta_1 \), the terms of interest in \( a \) combine to \( 2m_\tilde{B}(xy - wz) \neq 0 \) which is in fact proportional to \( \sin \gamma_f \). The effect of the phases on the bino mass limits is shown in Figures 3, 4 as a function of the magnitude of the off-diagonal term in the top-squark mass matrix, \( \mathbf{m}_t \), given the conditions: 1) \( \Omega_\tilde{B} h^2 < 1/4 \), 2) the lightest sfermion is heavier than the bino, and 3) the lightest sfermion is heavier than 74 GeV. In both figures we have taken \( \tan \beta = 2 \) and \( m_{\text{top}} = 174 \) GeV. In Figure 3, \( |\mu| = 3000 \) GeV and in Figure 4, \( |\mu| = 1000 \) GeV. The various curves are labeled by the value of \( \gamma_b \) assumed, and in addition, \( m_\tilde{B} \) has been maximized for all allowed values of \( \theta_\mu \).

The lower limit on \( m_\tilde{B} \) assumes \( \gamma_b = \pi/2 \). As one can see, when \( \gamma_b \) is allowed to take its maximal value of \( \pi/2 \), the upper limits are greatly relaxed to \( m_\tilde{B} \approx 650 \) GeV. With \( |\mu| \) and \( \gamma_b \) fixed, for a given value of \( \mathbf{m}_t \) and \( \theta_\mu \) all of the remaining quantities such as \( |A|, \theta_A, \gamma_t, \) and \( \mathbf{m}_b \) are determined.

For large \( \mathbf{m}_t \), the diagonal mass terms \( M^2_L \) must be taken large to ensure that the mass of the lightest stop is \( \gtrsim m_\tilde{B} \). This drives up the masses of the other sfermions and suppresses their contribution to the annihilation. As \( \mathbf{m}_t \) is decreased, \( M^2_L \) must drop and the other sfermions begin to contribute and the upper bound on \( m_\tilde{B} \) is increased. In particular,
annihilation to $\mu$'s and $\tau$'s becomes important, since

$$Y_L^2 Y_R^2 \bigg|_{\mu,\tau} : Y_L^2 Y_R^2 \bigg|_{c,t} : Y_L^2 Y_R^2 \bigg|_{s,b} = 81 : 4 : 1$$  \hspace{1cm} (9)

Decreasing $\gamma_b$ reduces the effect of $\mu$'s and $\tau$'s, and this can be seen as a decrease in the upper bounds in Figures 3 and 4. For $\overline{m}_t$ sufficiently small, the stops become unmixed, diminishing somewhat their contribution and slightly decreasing the upper bound on $m_{\tilde{B}}$.

Figure 3: Upper limits on the bino mass as a function of the off-diagonal element $\overline{m}_t$ in the top squark mass matrix, for various values of $\gamma_b$, the argument of the off-diagonal element of the $T_3 = -1/2$ sfermion mass matrix. Also shown is the lower bound (lowest curve) on the bino mass assuming $\gamma_b = \pi/2$. The value of $|\mu|$ was chosen to be 3000 GeV.

Figure 4: As in Fig. 3, with $|\mu| = 1000$ GeV.

I turn now to the calculation of the electric dipole moments of the neutron and the electron. As discussed above, the edm's of the electron and quarks receive contributions from one-loop diagrams involving the exchange of sfermions and either neutralinos, charginos, or (for the quarks) gluinos. In the case of the neutron edm, there are additional operators besides the quark electric dipole operator, $O_{\gamma} = \frac{1}{4} \bar{q} \sigma_{\mu\nu} q \tilde{F}^{\mu\nu}$ which contribute. They are the gluonic operator $O_G = -\frac{1}{6} f^{abc} G_a G_b G_c$ and the quark color dipole operator.
The gluonic operator is the smallest \[ O_q : O_g = 21 : 4.5 : 1. \] when all mass scales are taken to be equal. These three operators are conveniently compared to one another in [20] and relative to the gluino exchange contribution to the \( O_\gamma \) operator, it was found that \( O_\gamma : O_q : O_g = 21 : 4.5 : 1. \).

Because of the reduced importance of the additional operators contributing to the neutron edm, it is sufficient to only include the three contributions to the quark electric dipole moment. The necessary CP violation in these contributions comes from either \( \gamma_f \) in the sfermion mass matrices or \( \theta_\mu \) in the neutralino and chargino mass matrices. Full expressions for the chargino, neutralino and gluino exchange contributions are found in [4]. The dependencies of the various contributions on the CP violating phases can be neatly summarized: for the chargino contribution

\[
d_C \sim \sin \theta_\mu ,
\]

with essentially no dependence on \( \gamma_f \); whilst for the gluino contribution

\[
d_G \sim \overline{m}_f \sin \gamma_f ,
\]

independent of \( \theta_\mu \), and the neutralino contribution has pieces that depend on both \( \sin \theta_\mu \) and \( \overline{m}_f \sin \gamma_f \). All three contributions can be important (including the neutralino contribution, in the case of the electron edm), and depending on \( \sin \theta_\mu \) and \( \sin \gamma_f \), they can come in with either the same or opposite signs. In particular, \( \text{sign}[d_C / d_G] = \text{sign}[\sin \theta_\mu / \sin \gamma_f] \). For the mass ranges we consider, the dipole moments fall as the sfermion masses are increased, and sfermion masses in the TeV range can bring these contributions to the neutron and electron electric dipole moments below the experimental bounds of \( |d_n| < 1.1 \times 10^{-25} \text{e cm} \) [5] and \( |d_e| < 1.9 \times 10^{-26} \text{e cm} \) [21], even for large values of the CP violating phases [4]. However, these large sfermion masses are inconsistent with the cosmological bounds mentioned above, where sfermion masses must be relatively close to the bino mass in order to keep the relic density in check.

To determine the allowed parameter ranges [1], we first fix the value of \( \gamma_b \) and take \( |\mu| = 3000 \text{ GeV} \). Then for several values of \( \overline{m}_t \) between 0 and 1500 GeV, we determine the upper bound on \( m_\tilde{B} \), as a function of \( \theta_\mu \). As we vary \( \theta_\mu \) across its full range, \( \overline{m}_t \) and \( \gamma_t \) change, and this affects the annihilation rate and consequently the bound on \( m_\tilde{B} \). Taking \( m_\tilde{B} \) at its maximum value allows us to take \( M_L^2 \) as large as possible; although the electric dipole moments depend on \( m_\tilde{B} \) as well, the dependence on \( M_L^2 \) is sufficiently strong that the edm’s take their minimum values for the maximum values of \( m_\tilde{B} \) and \( M_L^2 \). The quark and electron edm’s can then be computed as a function of \( \theta_\mu \) and \( \overline{m}_t \), and one can use the nonrelativistic quark model to relate the neutron edm to the up and down-quark edm’s via

\[
d_n = \frac{1}{3}(4d_d - d_u).
\]

If there is no region of the \( \theta_\mu, \overline{m}_t \) parameter space which satisfies both the neutron and electron edm bounds, we can decrease \( \gamma_b \) and repeat the procedure. In practice, the bound
on the neutron edm the more difficult of the two to satisfy, and every region of the parameter space we show which produces an acceptable neutron edm also produces a sufficiently small electron edm.

For the large value of $|\mu| = 3000 \text{ GeV}$, the largest contribution to the neutron edm comes either from gluino exchange (for the more negative values of $\theta_\mu$) or chargino exchange (for the more positive values of $\theta_\mu$), and the value for $|d_n|$ is too large unless $\gamma_b$ takes a relatively small value. In particular, non-negligible experimentally acceptable regions of the parameter space are found only for $\gamma_b \lesssim \pi/25$. In Figure 5, a contour plot of the neutron edm as a function of $\theta_\mu$ and $\overline{m}_t$ for $\gamma_b = \pi/40$ is shown \[\text{[1]}\]. The shaded regions demarcate the range of $\theta_\mu$ for this choice of $\gamma_b$. Much of this range produces a sufficiently small $|d_n|$. As we increase $m_t$, the $\tilde{d}$ and $\tilde{u}$ masses become large and $|d_n|$ falls. As we move to values of $\overline{m}_t$ greater than $\sim 1500 \text{ GeV}$, we begin to require a significant tuning of $M_L^2$ to produce $\Omega_{\tilde{B}}h^2 < 1/4$.

This procedure can be repeated for $|\mu| = 1000 \text{ GeV}$ and the result is shown in Figure 6 \[\text{[1]}\]. For lower $|\mu|$, the chargino exchange contribution is enhanced relative to the gluino exchange contribution. In this case, one can allow larger values of $\gamma_b$ up to about $\pi/6$. For the same range in $\overline{m}_t$, $\theta_\mu$ can takes values from -0.3 to 0.2. For either value of $\mu$, the angles $\gamma_t$ and $\theta_A$ are unconstrained.

In summary, we have found that CP violating phases in the MSSM can significantly affect the cosmological upper bound on the mass of an LSP bino. In particular, taking the maximal value $\pi/2$ for the phase $\gamma_b$ of the off-diagonal component of the $T_3 = -1/2$ sfermion mass matrices pushes the upper bound on $m_{\tilde{B}}$ up from $\sim 250 \text{ GeV}$ to $\sim 650 \text{ GeV}$. When we additionally consider constraints on neutron and electron electric dipole moments, we find the upper bound on $m_{\tilde{B}}$ is reduced to $\sim 350 \text{ GeV}$. Various combinations of the CP violating phases are constrained as well: $|\theta_\mu| \lesssim 0.3$ and $|\gamma_b| \lesssim \pi/6$ for $|\mu| \gtrsim 1000 \text{ GeV}$, while $\gamma_t$ and $\theta_A$ are essentially unconstrained. We note that although the bounds on $\theta_\mu$ and $\gamma_b$ are small, they are much larger than the values of order $10^{-3}$ typically considered.

Figure 5: Contours of the neutron electric dipole moment, $d_n$, in the $\theta_\mu - \overline{m}_t$ plane in units of $10^{-25} e \text{ cm}$. The value of $|\mu|$ was chosen to be $3000 \text{ GeV}$ and $\gamma_b = \pi/40$. The shaded region corresponds to values of $\theta_\mu$ and $\overline{m}_t$ which are not allowed algebraically for this value of $\mu$ and $\gamma_b$. The value of $|\mu|$ was chosen to be $3000 \text{ GeV}$ and $\gamma_b = \pi/40$. The shaded region corresponds to values of $\theta_\mu$ and $\overline{m}_t$ which are not allowed algebraically for this value of $\mu$ and $\gamma_b$.
Figure 6: As in Figure 5, with $|\mu|$ chosen to be 1000 GeV and $\gamma_b = \pi/8$.

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\[ |\mu| = 3000 \text{ GeV} \]

\[ M_{\tilde{B}} (\text{GeV}) \]

\[ m_t (\text{GeV}) \]

- \( \gamma_b = \pi/2 \)
- \( = \pi/4 \)
- \( = \pi/8 \)
- \( = \pi/20,40 \)
$|\mu| = 1000 \text{ GeV}$

$\gamma_b = \pi/2$

$\gamma_b = \pi/8$
