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Neutralino-Stau Coannihilation and the Cosmological Upper Limit on the Mass of the Lightest Supersymmetric Particle

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Abstract

We consider the effects of neutralino-stau ($\chi - \tilde{\tau}$) coannihilations on the cosmological relic density of the lightest supersymmetric particle (LSP) $\widetilde{\chi}$ in the minimal supersymmetric extension of the Standard Model (MSSM), particularly in the constrained MSSM in which universal supergravity inputs at the GUT scale are assumed. For much of the parameter space in these models, $\widetilde{\chi}$ is approximately a $U(1)$ gaugino $\tilde{B}$, and constraints on the cosmological relic density $\Omega_{\widetilde{B}} h^2$ yield an upper bound on $m_{\widetilde{B}}$. We show that in regions of parameter space for which the cosmological bound is nearly saturated, coannihilations of the $\tilde{B}$ with the $\tilde{\tau}$, the next lightest sparticle, are important and may reduce significantly the $\tilde{B}$ relic density. Including also $\tilde{B}$ coannihilations with the $\tilde{e}$ and $\tilde{\mu}$, we find that the upper limit on $m_{\widetilde{\chi}}$ is increased from about 200 GeV to about 600 GeV in the constrained MSSM, with a similar new upper limit expected in the MSSM.

Supersymmetry is one of the most appealing options for possible physics beyond the Standard Model, motivated theoretically by the help it offers in stabilizing the gauge hierarchy, its successful prediction of $\sin^2\theta_W$ in the context of GUTs, and its consistency with the range of Higgs boson masses favoured by the precision electroweak data. Accordingly, supersymmetry has been the focus of intense phenomenological studies, particularly in the framework of the minimal supersymmetric extension of the Standard Model (MSSM) [1]. Many of these studies assume that supersymmetry is broken by unspecified dynamics in some hidden sector of the theory, which is communicated to the observable MSSM particles by gravitational interactions. One may further assume that the supersymmetry-breaking mass parameters $m_0, m_{1/2}, A$ and B are universal at the supergravity input scale, providing the constrained MSSM (CMSSM) framework that is privileged in this paper.

The lightest supersymmetric particle (LSP) is stable in the MSSM, unless additional R-violating interactions are postulated [2]. The LSP is generally thought to be the lightest neutralino $\widetilde{\chi}$ [3, 4] and is a favoured candidate for the Cold Dark Matter favoured by astrophysicists and theorists of structure formation. The relic LSP density can in principle be calculated reliably as a function of the parameters of the MSSM [5]. It is remarkable that there is a large generic domain of the parameter space in the MSSM and in the CMSSM, consistent with all experimental constraints [6], in which the $\widetilde{\chi}$ has a relic mass density $\Omega_{\widetilde{\chi}} h^2 \sim 0.1$ as favoured by astrophysical and cosmological arguments.

The phenomenological arguments for supersymmetry based on the gauge hierarchy, $\sin^2\theta_W$ and the Higgs mass m_h all suggest that supersymmetric particles should weigh ~ 1 TeV or less, but do not provide very precise upper limits on their masses. For example, the amount of fine tuning required to maintain the gauge hierarchy increases as the MSSM mass parameters are increased, but there is no objective criterion how much fine tuning is tolerable [7]. Moreover, $\sin^2\theta_W$ and m_h are only logarithmically sensitive to the sparticle masses. On the other hand, the LSP relic mass density is very sensitive to $m_{\widetilde{\chi}}$, since the annihilation cross section tends to decrease as $m_{\widetilde{\chi}}$ increases, increasing also its relic number number density. The relic density is also very sensitive over much of the parameter space to the scalar mass parameters, as these, too, control the annihilation cross section. The constraints on the general MSSM parameter space have been explored in some detail, particularly in the CMSSM framework. The possibility that the LSP might be mainly a photino $\tilde{\gamma}$ has been excluded by lower limits on sparticle masses from LEP and elsewhere [6]. When studying the possibility that the LSP might be largely a Higgsino $\tilde{H}$, coannihilations [8] between the Higgsino-like LSP and the next-to-lightest supersymmetric particle (NLSP) have to be taken into account [9, 10]. This Higgsino LSP possibility has also been tightly constrained by LEP and may be explored completely by upcoming runs [6]. This leaves us with the likelihood of a Bino- ($\tilde{B}$-) like LSP, and cosmology imposes an important upper limit on its mass, which has been given as $m_{\tilde{B}} \lesssim 300$ GeV in the MSSM [11] and $\lesssim 200$ GeV the CMSSM framework [12].

The purpose of this paper is re-evaluate this upper limit, including for the first time detailed calculations of coannihilations between the $\tilde{B}$ LSP and the lighter supersymmetric partner of the τ, the right-handed stau $\tilde{\tau}_R$, which is the NLSP in much of the $\tilde{B}$ LSP region. We find that $\tilde{B} - \tilde{\tau}_R$ coannihilation is particularly important when the relic mass density is close to the cosmological upper limit, which we take to be $\Omega_{\widetilde{\chi}} h^2 < 0.3$, resulting in a considerable relaxation of the previous upper bound on $m_{\tilde{B}}$. Including also $\tilde{B}$ coannihilation with the $\tilde{e}_R$ and $\tilde{\mu}_R$, we now find that $m_{\tilde{B}}$ may be as large as 600 GeV in the CMSSM or MSSM.

As already commented, the LSP is a $\tilde{B}$ in much of parameter space that leads to an interesting relic density, both in the generic MSSM and in the CMSSM [6]. Indeed, this is a prediction of the CMSSM. Unless the $\tilde{B}$ mass happens to lie near $m_Z/2$ or $m_h/2$, in which case there are large contributions to the annihilation through direct s-channel resonance exchange, the dominant contribution to the $\tilde{B}\tilde{B}$ annihilation cross section comes from crossed t-channel sfermion exchange. The resonant case is anyway not relevant for the upper bounds on $m_{\tilde{B}}$ to be discussed here. In the absence of such a resonance, the thermally-averaged cross section for $\tilde{B}\tilde{B} \to f\bar{f}$ takes the generic form

$$\langle \sigma v \rangle = (1 - \frac{m_f^2}{m_{\widetilde{B}}^2})^{1/2} \frac{g_1^4}{128\pi} \left[(Y_L^2 + Y_R^2)^2 (\frac{m_f^2}{\Delta_f^2}) \; + \; (Y_L^4 + Y_R^4)(\frac{4m_{\widetilde{B}}^2}{\Delta_f^2})(1 + ...) \, x \right] \tag{1}$$

where $Y_{L(R)}$ is the hypercharge of $f_{L(R)}$, $\Delta_f \equiv m_{\tilde{f}}^2 + m_{\widetilde{B}}^2 - m_f^2$, and we have shown only the leading P-wave contribution proportional to $x \equiv T/m_{\widetilde{B}}$.

The upper limit on $m_{\tilde{B}}$ due to the cosmological relic density comes about as follows [11]. The assumption that the $\tilde{B}$ is the LSP requires, in particular, that $m_{\widetilde{B}} < m_{\tilde{f}}$. In order to minimize the relic density, we must maximize the cross section, which is done by setting $m_{\tilde{f}} = m_{\widetilde{B}}$. The cross section is then approximately inversely proportional to $m_{\widetilde{B}}^2$. The cosmological upper limit on $\Omega_{\tilde{B}} h^2$ translates into a lower limit on $\langle \sigma v \rangle$ which then, in turn, yields an upper limit to $m_{\widetilde{B}}$. In the MSSM, this limit is $m_{\widetilde{B}} \lesssim 300$ GeV, when all sfermion masses are taken to be equal at the weak scale.

In the CMSSM, the argument is somewhat similar, although $m_{\widetilde{B}}$ and the sfermion masses are no longer entirely independent, because it is assumed in the CMSSM that there is a common scalar mass m_0 at the GUT scale. For a given value of the common gaugino mass

$m_{1/2}$ at the GUT scale, the relic $\tilde{B}$ density falls with m_0, since $m^2_{\tilde{f}} = m_0^2 + C_f m^2_{1/2} + O(m_Z^2)$, where C_f is a positive numerical coefficient that is calculable via the renormalization-group evolution of the sfermion masses. Therefore, the cosmological upper limit on $\Omega_{\tilde{B}} h^2$ translates at fixed $m_{1/2}$ into an upper limit on m_0. Typically, this upper limit is not larger than $m_0 \lesssim 150$ GeV, unless one is sitting on a direct-channel pole, i.e., when $m_{\widetilde{B}} \sim m_Z/2$ or $m_h/2$, in which case s-channel annihilation is dominant and there is no upper limit to m_0. However, as already mentioned, this is not our case, as we are interested in an upper bound on $m_{\widetilde{B}}$. We recall that $m_{\widetilde{B}}$ scales with $m_{1/2}$, and it transpires for $m_{1/2} \gtrsim 400$ GeV that $m_{\widetilde{B}}$ exceeds mass of the lightest sfermion, which is typically the $\tilde{\tau}_R$, for m_0 small enough to satisfy the cosmological bound [12]. Thus, the LSP is no longer a neutralino for such large values of $m_{1/2}$, and hence an upper bound $m_{\widetilde{B}} \lesssim 200$ GeV [12] can be established. [1]

When $m_{\widetilde{B}}$ attains this upper bound, the $\tilde{B}$ is degenerate in mass with the $\tilde{\tau}_R$, and quite close in mass to the $\tilde{e}_R$ and $\tilde{\mu}_R$. It is well known [8] that, in such circumstances, the neutralinos can be maintained in equilibrium by scatterings with a slightly heavier particle, and the number density of neutralinos can be significantly reduced by such coannihilations. The case of heavy Higgsinos is a well studied example [10]. Analogously to that case, the $\tilde{B}$ relic density can be reduced through coannihilation with slightly heavier $\widetilde{\tau}_{\mathrm{R}}$'s or other sleptons, as we now discuss in detail.

To derive a thermally-averaged cross section, we use the technique of [5]. Thus, we expand $\langle \sigma v_{\rm rel} \rangle$ in a Taylor expansion in powers of $x = T/m_{\widetilde{B}}$:

$$\langle \sigma v_{\rm rel} \rangle = a + bx + O(x^2) \; . \tag{2}$$

Repeating the analysis [5] for initial particles with different masses m_1 and m_2 yields

$$\langle \sigma v_{\rm rel} \rangle = \frac{1}{m_1 m_2} \left(1 - \frac{3(m_1+m_2)T}{2 m_1 m_2} \right) w(s)|_{s \to (m_1+m_2)^2 + 3(m_1+m_2)T} \; + \; O(T^2). \tag{3}$$

where

$$\begin{aligned} w(s) \; &\equiv \; \frac{1}{4} \int \, d\mathrm{LIPS} \; |\mathcal{M}|^2 &(4) \\ &= \; \frac{1}{32\pi} \frac{p(s)}{s^{1/2}} \int_{-1}^{+1} d\cos\theta_{\rm CM} \; |\mathcal{M}|^2, &(5) \end{aligned}$$

Here dLIPS is the Lorentz-Invariant phase-space element, $p(s)$ is the magnitude of the three momentum of one of the initial particles in the center-of-mass frame, as a function of the

[1] This upper bound can be strengthened by requiring that the global minimum of the effective potential of the MSSM conserve electric charge and color [13].

total center-of-mass energy-squared s, $\theta_{\rm CM}$ is the center-of-mass scattering angle, and $|\mathcal{M}|^2$ is the absolute square of the reduced matrix element for the annihilation, summed over final spins and averaged over initial spins wherever appropriate. The a and b coefficients in (2) may be read off the right-hand side of (3), after expanding in powers of x. [2]

If the masses of the next-to-lightest sparticles (NLSPs) are close to the LSP mass: $\Delta M = \mathcal{O}(x_f) M$, where $x_f \sim (1/20 - 1/25)$ is the value of $T/m_{\tilde{\chi}}$ at the time of neutralino decoupling, the number densities of the NLSPs have only slight Boltzmann suppressions with respect to the LSP number density when the LSP freezes out of chemical equilibrium with the thermal bath. In such a case, coannihilations of NLSPs with the LSP, along with NLSP-NLSP annihilations, may play an important rôle in keeping the LSPs in chemical equilibrium with the bath [8]. These processes can be particularly important when the LSP annihilation rate itself is suppressed, as is the case for neutralinos. Gaugino-like neutralinos typically annihilate predominantly into fermion pairs, and such processes exhibit P-wave suppressions [3], so that $a \ll b$ in (2). This effect can be seen from (1) where the a-term is suppressed relative to b by $m_f^2/m_{\widetilde{B}}^2$, reducing the neutralino annihilation cross section by a factor of $\mathcal{O}(x_f)$. We also emphasize that 2-2 scatterings with particles in the thermal bath keep the NLSPs, in this case the $\widetilde{\tau}_{\rm R}$, $\widetilde{e}_{\rm R}$ and $\widetilde{\mu}_{\rm R}$, in chemical equilibrium with each other and with $\widetilde{\chi}$, down to temperatures well below the temperature at which the comoving LSP number density freezes out.

We consider the total density $n \equiv \sum_i n_i$, where the index i runs over $\widetilde{\tau}_{\rm R}, \widetilde{\tau}^*_{\rm R}, \widetilde{e}_{\rm R}, \widetilde{e}^*_{\rm R}, \widetilde{\mu}_{\rm R}$ and $\widetilde{\mu}^*_{\rm R}$ as well as $\widetilde{\chi}$, and write the rate equation for n:

$$\frac{dn}{dt} = -3Hn - \langle \sigma_{\rm eff} v_{\rm rel} \rangle (n^2 - n_{\rm eq}^2), \tag{6}$$

where H is the Hubble parameter, and

$$\sigma_{\rm eff} = \sum_{ij} \sigma_{ij} r_i r_j. \tag{7}$$

Here $r_i \equiv n_i^{\rm eq}/n^{\rm eq}$ where $n_i^{\rm eq}$ is the equilibrium density of particle species i, and σ_{ij} is the total cross section for particle i to annihilate with particle j. Since the sleptons decay into neutralinos after freeze-out, the number density of neutralinos becomes n. Many of the σ_{ij} are related, and if we take the $\widetilde{e}_{\rm R}$ and $\widetilde{\mu}_{\rm R}$ to be degenerate in mass and ignore the electron and muon mass, we can write

$$\begin{aligned} \sigma_{\rm eff} \;=\; & \sigma_{\chi\chi} r_\chi r_\chi + 4\,\sigma_{\chi\tau} r_\chi r_\tau + 8\,\sigma_{\chi e} r_\chi r_e + 2\,(\sigma_{\tau\tau} + \sigma_{\tau\tau^*}) r_\tau r_\tau + 8\,(\sigma_{\tau e} + \sigma_{\tau e^*}) r_\tau r_e + \\ & 4\,(\sigma_{ee} + \sigma_{ee^*}) r_e r_e + 4\,(\sigma_{e\mu} + \sigma_{e\mu^*}) r_e r_e \end{aligned} \tag{8}$$

[2] A similar calculation is necessary when the LSP is assumed to be a sneutrino [14].

In Table 1 we list the sets of initial and final states for which we compute the annihilation cross sections. The cross sections for other reactions, such as $\widetilde{\ell}_{\rm R}\widetilde{\ell}^*_{\rm R} \to hH, hA, hZ, H^+H^-$, $W^+H^-, \ldots$ are either suppressed or kinematically unavailable in the regions of CMSSM parameter space relevant to our analysis. In practice, we find that the dominant contributions to $\sigma_{\rm eff}$ come from annihilations of $\widetilde{\ell}^i_{\rm R}\,\widetilde{\ell}^{i*}_{\rm R}$ to gauge bosons, $\widetilde{\ell}^i_{\rm R}\,\widetilde{\ell}^j_{\rm R}$ to lepton pairs, and $\widetilde{\ell}^i_{\rm R}\,\widetilde{\chi}$ to $\ell^i +$ gauge boson.

Table 1: Initial and Final States for Coannihilation: $\{i, j = \tau, e, \mu\}$

Initial State	Final States
$\widetilde{\ell}^i_{\rm R}\,\widetilde{\ell}^{i*}_{\rm R}$	$\gamma\gamma,\ ZZ\,,\ \gamma Z,\ W^+W^-,\ hh,\ \ell^i\bar{\ell}^i$
$\widetilde{\ell}^i_{\rm R}\,\widetilde{\ell}^j_{\rm R}$	$\ell^i\ell^j$
$\widetilde{\ell}^i_{\rm R}\,\widetilde{\ell}^{j*}_{\rm R},\ i \neq j$	$\ell^i\bar{\ell}^j$
$\widetilde{\ell}^i_{\rm R}\,\widetilde{\chi}$	$\ell^i\gamma, \ell^i Z, \ell^i h$

To get a simple estimate for the size of the effect of including the next-to-lightest states, we first assume degenerate LSP and NLSPs, and consider a model in which the NLSP-NLSP and NLSP-LSP annihilations are all unsuppressed. Thus, we take

$$\{a_{ij} \approx a_{1j} \approx b_{11}, a_{11} = 0; i, j > 1\}, \tag{9}$$

where the subscripts i, j refer to the NLSPs and 1 to the LSP. Denoting with superscripts 0 quantities that are computed ignoring the NLSP states, we estimate the following ratio of relic densities without and with coannihilation:

$$R \equiv \frac{\Omega^0}{\Omega} \approx (\frac{2}{x^0_f})\,(\frac{a_{\rm eff}}{b_{11}})\,(\frac{x_f}{x^0_f}), \tag{10}$$

where $x^0_f/x_f \approx 1 + x^0_f \ln(g_{\rm tot}/g_1 x^0_f) \approx 1.2$, $g_{\rm tot} = \sum_i g_i$, and $a_{\rm eff}/b_{11} \approx 1 - g_1^2/g_{\rm tot}^2 = 15/16$ for the case of three degenerate slepton NLSPs. Thus, in this crude approximation we find a factor ~ 35 reduction in the relic density. Ignoring the (heavier) left-handed sleptons, we may reduce (1) to $\langle\sigma v\rangle \approx 3g_1^4 x/(16\pi m^2_{\widetilde{\chi}})$, yielding $\Omega_{\widetilde{\chi}} h^2 \sim 8 \times 10^{-6}\ m^2_{\widetilde{\chi}}$. Thus, in this simple approximation, $\Omega_{\widetilde{\chi}} h^2 = (\Omega_{\widetilde{\chi}} h^2)^0/R < 0.3$ gives an upper bound on the $\tilde{B}$ mass of $m_{\widetilde{B}} \lesssim 1.2\,{\rm TeV}$.

We have gone beyond the above crude approximations to make a detailed numerical analysis of coannihilation effects on the neutralino relic density, including light sleptons, some of whose results are displayed in Fig. 1. The light shaded region corresponds to $0.1 < \Omega_{\widetilde{\chi}} h^2 < 0.3$, and the dark shaded region to $m_{\widetilde{\tau}_R} < m_{\widetilde{\chi}}$. We have chosen the representative points $\tan\beta = 3$ and 10, and present results for both $\mu < 0$ and $\mu > 0$. In practice, the relationship (9) is not exact, not all of the a_{ij} are unsuppressed, there are contributions from the b_{ij}, and the $\widetilde{e}_R$ and $\widetilde{\mu}_R$ can be slightly heavier than the $\widetilde{\tau}_R$. These corrections have the net effect of reducing $\sigma_{\rm eff}$ by a factor of $\sim (3-4)$. Numerically, we find that $R \sim 10$ along the line where $m_{\widetilde{\chi}} = m_{\widetilde{\tau}_R}$, at the top of the dark shaded region in Fig. 1. As m_0 increases, the sleptons become heavier relative to the neutralino, and their number density rapidly falls, reducing their contribution to $\sigma_{\rm eff}$. The relic density rises rapidly in this region, leaving an allowed band in m_0 which is about 30-50 GeV wide for $m_{1/2} < 800\,{\rm GeV}$. In Fig. 2 we extend the coverage of Fig. 1a,c over a larger scale, to show the cross-over point between the regions with $\Omega_{\widetilde{\chi}} h^2 < 0.3$ and $m_{\widetilde{\tau}_R} < m_{\widetilde{\chi}}$, where there would be an unacceptable abundance of charged dark matter [4]. The two constraints together require $m_{1/2} \lesssim 1450$, corresponding to an upper bound on the neutralino mass of $m_{\widetilde{\chi}} \lesssim 600\,{\rm GeV}$. The results for $\mu > 0$ and $\mu < 0$ are very similar, so we do not display the latter. The width of the allowed region is insensitive to A_0, though the position of the line $m_{\widetilde{\chi}} = m_{\widetilde{\tau}_R}$ can vary somewhat. As already commented, the requirement that the electroweak vacuum conserve electric charge and color constrains significantly the CMSSM parameter space [13]. We find that the large-$m_{1/2}$ tail of the region newly allowed by coannihilation obeys this requirement, for $\tan\beta \gtrsim 3$ for some values of A_0.

We expect the corresponding bound in the MSSM to be very similar. In the general case, one must take all the squarks and sleptons degenerate with the neutralino and compute the annihilation and coannihilation cross-sections for all possible combination of sfermions. However, if the rates are the same as for the sleptons, the effect is about a 15% decrease in $(\Omega_{\widetilde{\chi}} h^2)^0/R$, leading to a similar bound on $m_{\widetilde{\chi}}$ as in the CMSSM.

The potential significance of coannihilation effects had been emphasized previously [8, 9, 10], particularly in the Higgsino LSP region. We have shown that this can also be an important effect in the $\tilde{B}$ region, where the NLSPs include the $\tilde{\tau}_R$ and other right-handed sleptons. In particular, we find in the context of the CMSSM that the upper bound on the LSP mass quoted previously [12] should be increased by a factor of about two, to $m_{\widetilde{\chi}} \lesssim 600\,{\rm GeV}$.

This observation has many potential ramifications, in particular for searches for neutralino dark matter. Generically, in regions of parameter space where coannihilation is important, a relic density of astrophysical interest is now obtained for a smaller annihilation cross section. This is likely to reduce the typical rates for signatures of annihilations in the galactic halo. The corresponding elastic scattering cross section is also likely to be reduced generically, with a consequent suppression of signatures for scattering on nuclei and capture followed by annihilation in the Sun or Earth. These points are worthy of further study.

On the basis of the previously-quoted upper bound on the LSP mass, it has been argued that the LHC is guaranteed to detect some supersymmetric particles [15], since it could reach out to $m_{1/2} \sim 1200$ GeV when $m_{\widetilde{B}} \sim m_{\tilde{\tau}_R}$. It is still true that the overwhelming majority of the CMSSM parameter space allowed by cosmology can be explored by the LHC, but there may be a narrow region of m_0 extending up to $m_{1/2} \sim 1500$ GeV that is problematical. This point is also worthy of further study.

Acknowledgments

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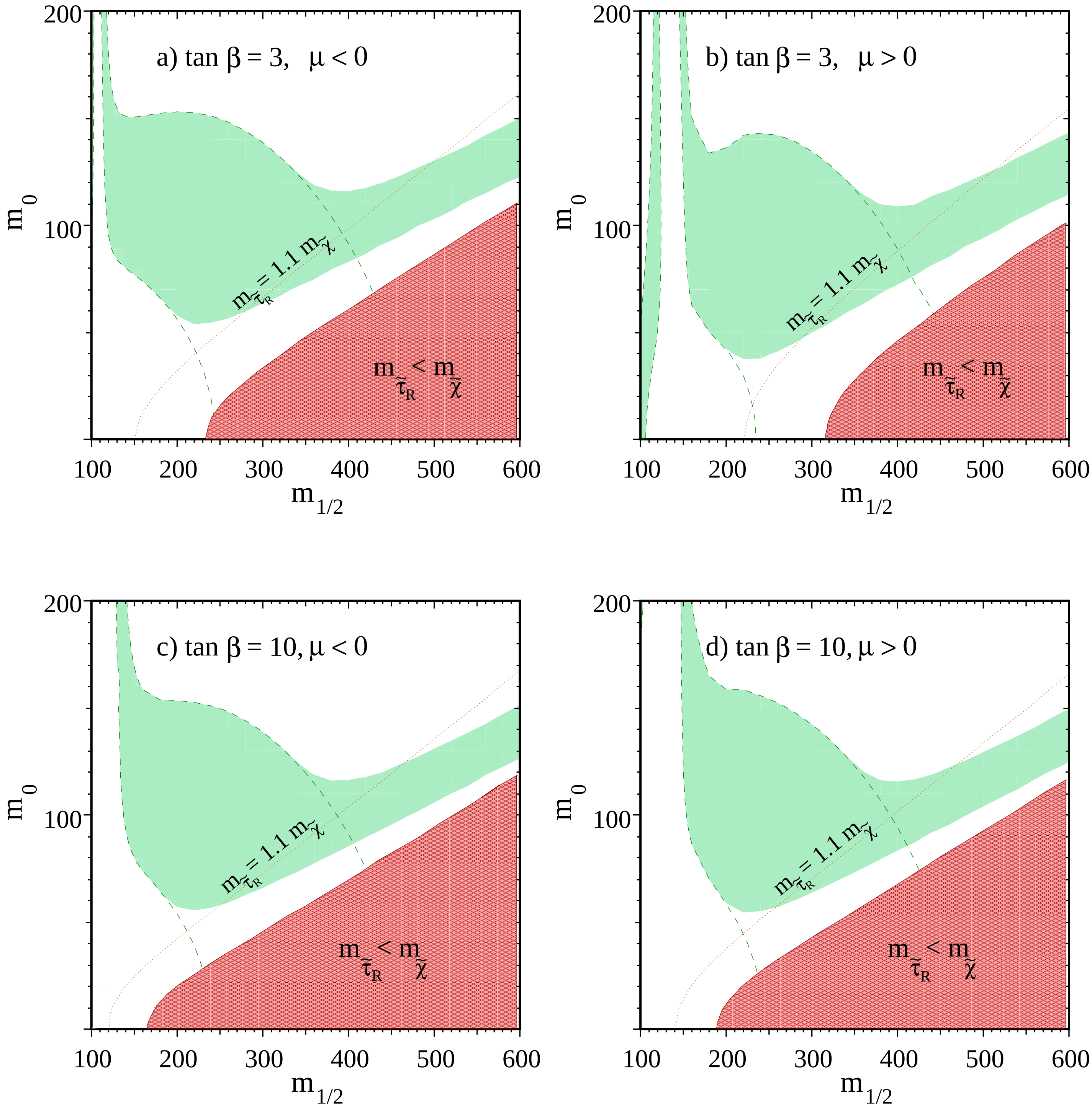

Figure 1: The light-shaded area is the cosmologically preferred region with $0.1 \leq \Omega_{\tilde{\chi}}\, h^2 \leq 0.3$. The dashed line shows the location of the cosmologically preferred region if one ignores the light sleptons. In the dark shaded region in the bottom right of each panel, the LSP is the $\tilde{\tau}_R$, leading to an unacceptable abundance of charged dark matter. Also shown as a dotted line is the contour $m_{\tilde{\tau}_{\rm R}} = 1.1\, m_{\tilde{\chi}}$.

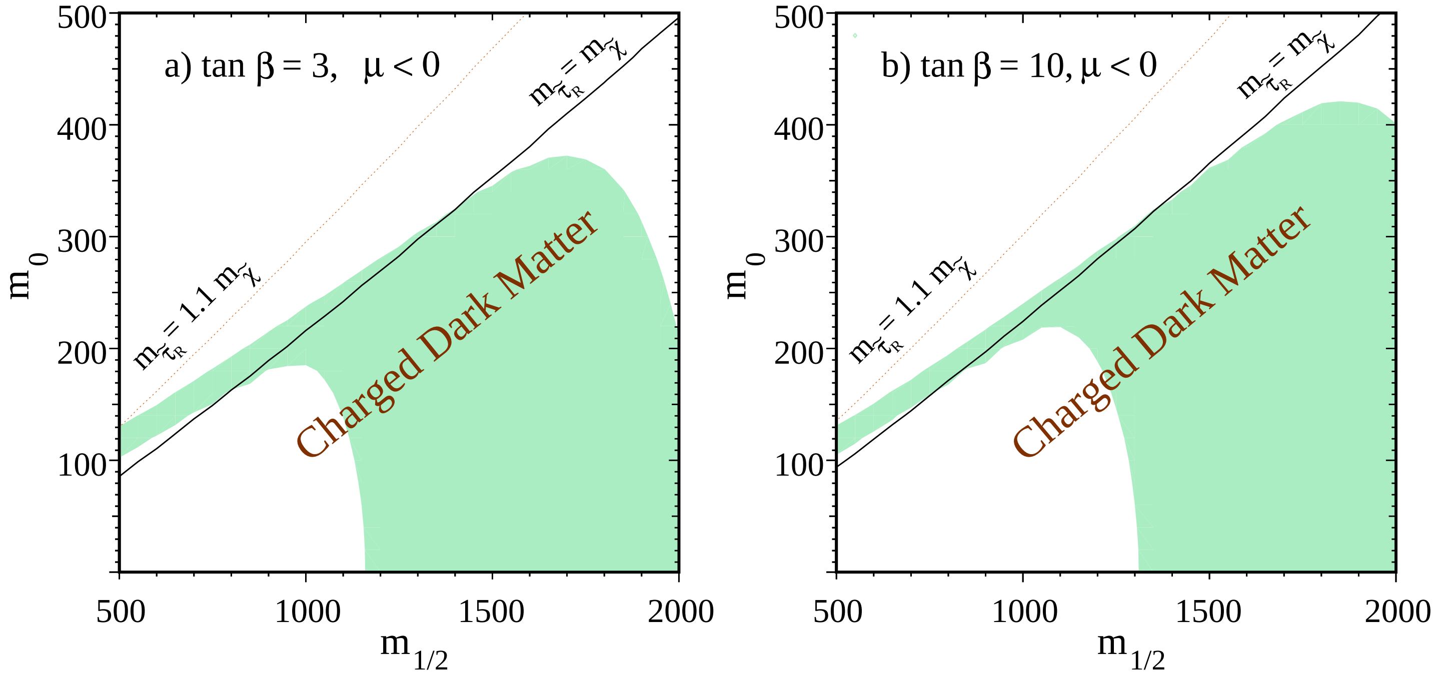

Figure 2: Same as Fig. 1(a,c), extended to larger values of $m_{1/2}$.