New Constraints on Neutralino Dark Matter in the Supersymmetric Standard Model

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ABSTRACT

We investigate the prospects for neutralino dark matter within the Supersymmetric Standard Model (SSM) including the constraints from universal soft supersymmetry breaking and radiative breaking of the electroweak symmetry. The latter is enforced by using the one-loop Higgs effective potential which automatically gives the one-loop corrected Higgs boson masses. We perform an exhaustive search of the allowed five-dimensional parameter space and find that the neutralino relic abundance $\Omega_\chi h^2_0$ depends most strongly on the ratio $\xi_0 \equiv m_0/m_{1/2}$. For $\xi_0 \gg 1$ the relic abundance is almost always much too large, whereas for $\xi_0 \ll 1$ the opposite occurs. For $\xi_0 \sim 1$ there are wide ranges of the remaining parameters for which $\Omega_\chi \sim 1$. We also determine that $m_{\tilde{q}} \gtrsim 250\text{ GeV}$ and $m_{\tilde{\ell}} \gtrsim 100\text{ GeV}$ are necessary in order to possibly achieve $\Omega_\chi \sim 1$. These lower bounds are much weaker than the corresponding ones derived previously when radiative breaking was not enforced.
1. Introduction

The fundamental observation by the COBE satellite of minute anisotropies in the cosmic microwave background radiation [1], has once again made evident the synergism between particle physics and cosmology [2], since the most compelling explanation for these observations [3] appears to be found in inflationary theories of the early universe [4]. Theoretically, the observed large-scale structure of the universe requires the presence of a dominant cold dark matter (CDM) component in the energy density [5]. However, CDM may not be enough to produce the observed small-scale structure [6], and recent re-evaluations of this problem [7] suggest the presence of a sub-dominant hot dark matter (HDM) component as well, in the form of $\sim$ few eV neutrinos.

Even though the CDM component could have various origins [8], from the particle physics point of view it is customary to ascribe it to an $\mathcal{O}(100\,\text{GeV})$ stable particle, i.e., the lightest neutralino ($\chi$) (and also the lightest supersymmetric particle) of the minimal supersymmetric Standard Model (MSSM) [9]. The standard calculational procedure consists of searching through the vast parameter space of the model for ‘cosmologically favored’ regions, that is, regions with near critical relic abundance of neutralinos, i.e., where $\Omega_{\chi} h_0^2 \lesssim h_0^2$, with $\Omega_{\chi} = \rho_{\chi}/\rho_0$ the fraction of the present total energy density (assumed to be the critical density) in the form of neutralinos, and $h_0$ the Hubble parameter in units of $100\,\text{km}\,\text{s}^{-1}\,\text{Mpc}^{-1}$ (current observations indicate that $0.5 \leq h_0 \leq 1$). A more conservative (i.e., cosmological-model–independent) outcome of this procedure is the identification of regions of parameter space which are cosmologically disfavored, i.e., regions where $\Omega_{\chi} h_0^2 \gtrsim h_0^2$.

Since $\Omega_{\chi}$ depends on the annihilation cross section $\chi\chi \rightarrow \text{all}$, the whole set of couplings and masses in the model need to be specified before $\Omega_{\chi}$ can be accurately computed. In the MSSM twenty-one parameters are needed to specify the relevant quantities. To make things tractable, several ad-hoc assumptions about the model parameters are usually made [9,10], resulting in approximate results of some interest but of limited scope. Things simplify dramatically when, following well-established theoretical prejudices, the MSSM is embedded in a generic supergravity model with universal soft supersymmetry breaking [11,12,13,14,15], since then the parameter count drops to just eight. The final step in this theoretical round-up is to enforce the requirement of radiative electroweak symmetry breaking, which cuts down the parameter count to five variables, introduces a correlation between the usual relic density variables ($M_2, \mu$), and determines the Higgs boson masses...
in terms of the basic model parameters. (Hereafter we refer to this unified model as the Supersymmetric Standard Model (SSM).) The latter two novel effects have important consequences in the determination of the allowed region of parameter space and in the value of $\Omega_\chi h_0^2$ throughout this region. This ultimate calculation has been performed before for a ‘no-scale’-inspired supersymmetry breaking scenario (where $m_{1/2} \gg m_0, A \approx 0$) imposing the radiative breaking constraint at tree-level \cite{16} and one-loop \cite{11}, although in both cases with the inaccurate approximation of tree-level Higgs boson masses. The purpose of this note is to present the results of this calculation in the context of supergravity models with universal soft supersymmetry breaking and radiative electroweak breaking in the one-loop approximation \cite{18}, including the ensuing one-loop corrected Higgs boson masses \cite{19,20}.

Besides the particle physics model used to calculate the annihilation cross section, various approximations have been used regarding the composition of the $\chi$ state, the dominant annihilation channels, and the solution of the Boltzmann equation needed to determine $\Omega_\chi h_0^2$. See Refs. \cite{13,21,17} for a discussion of these matters. The actual relic density calculations in this paper have been carried out following the methods previously described in Ref. \cite{13}.

2. The favored regions

The Supersymmetric Standard Model (SSM) is an $SU(3) \times SU(2) \times U(1)$ model with the minimal three generations and two Higgs doublets of matter representations, and which is assumed to unify into a larger gauge group at a unification mass of $M_U \approx 10^{16}$ GeV. The parameter space of this model can be described in terms of three universal soft-supersymmetry breaking parameters: $m_{1/2}, m_0, A$, the top-quark mass $m_t$, and the ratio of Higgs vacuum expectation values $\tan \beta = v_2/v_1$; the sign of the superpotential Higgs mixing term $\mu$ is also undetermined \cite{16}. Several consistency and phenomenological constraints restrict the range of the model parameters \cite{16,18}. For ease of comparison with earlier work we calculate $\Omega_\chi h_0^2$ and plot it in the traditional relic density $(M_2, \mu)$ plane for a given $\tan \beta$, where $M_2 = (\alpha_2/\alpha_U)m_{1/2} \approx 0.83m_{1/2} \approx 0.30m_\tilde{\chi}$ is the $SU(2)$ gaugino mass. This is accomplished by numerically inverting the relation $\mu = \mu(m_t)$. We also take

\footnote{Concurrently with our calculation, there has appeared a new calculation \cite{17} imposing radiative breaking at tree-level but including the one-loop corrections to the Higgs boson masses, although in a more restrictive supersymmetry breaking scenario where $A$ is determined, and also assuming $\lambda_b(M_U) = \lambda_\tau(M_U)$.}
\[ \tan \beta = 2, 8, 15, 30, \xi_0 \equiv m_0/m_{1/2} = 0.0, 0.2, 0.4, 0.6, 0.8, 1.0 \text{ and } 2, 4, 6, 8; \xi_A \equiv A/m_{1/2} = -\xi_0, 0, +\xi_0; \text{ and vary } 0 \leq M_2 \leq 300 \text{ GeV and } 0 \leq |\mu| \leq 800 \text{ GeV so that } m_\tilde{g} \lesssim 1 \text{ TeV and fine-tuning of the parameters occurs at two-orders-of-magnitude or less} \ [18]. \]

The resulting set of figures gives a good sampling of the whole parameter space. In Fig. 1 we present a representative set of plots for \( \Omega_\chi h^2 \) in the \((M_2, \mu)\) plane for \( \tan \beta = 2, 8, 15, 30, \xi_0 = 0.0, 0.4, 1.0, \xi_A = 0, \) and both signs of \( \mu \). Note that since we have restricted \( M_2 < 300 \text{ GeV}, \) then \( m_\chi \) is also restricted: for large \( M_2, m_\chi < M_1 = \frac{5}{3} \tan^2 \theta_w M_2 \approx \frac{1}{2} M_2 \lesssim 150 \text{ GeV}. \) This upper bound on \( M_2 \) also eliminates the possibility of nearly pure higgsino \( \chi \)'s.

The first evident feature of these figures is that the allowed parameter space is considerably constrained. For \( \xi_0 \sim 1 \) the radiative breaking constraints exclude the lower diagonal portion of the figures due to the determination of \( \mu \). For example, at tree-level \[18],

\[
\mu^2_{\text{tree}} = m^2_{1/2} \left[ X_{1/2} + \xi_0^2 X_0 - \frac{1}{2} (M_Z/m_{1/2})^2 \right], \tag{2.1}
\]

and clearly for \( \xi_0 \lesssim 1, \mu \lesssim m_{1/2} \) since \( X_{1/2}, X_0 \) are \( \mathcal{O}(1) \) functions \[18]. This correlation between \( \mu \) and \( M_2 \) is a distinctive feature of supergravity models with radiative electroweak breaking and (for \( \xi_0 \lesssim 1 \)) eliminates a large portion of the otherwise allowed parameter space (for example, compare Fig. 1c with Fig. 2 in Ref. \[14\]). Another consequence of the radiative breaking constraints is the determination of the Higgs boson masses in terms of the model parameters. This implies that as one moves around the \((M_2, \mu)\) plane in Fig. 1, the Higgs masses are varying continuously. In fact, with growing \( M_2 \) \( \text{(i.e., } m_{1/2}) \), the lightest Higgs boson \( (h) \) quickly approaches its upper bound (which is variable at one-loop but does not exceed \( \approx 150 \text{ GeV for } m_\tilde{g} < 1 \text{ TeV} \)) and the other three Higgs bosons become quite massive. This means that: (i) the important Higgs-mediated \( s \)-channel annihilation into fermion pairs becomes suppressed due to massive propagators as \( M_2 \) grows; and (ii) the \( \chi\chi \to hh \) channel is open throughout a fixed region in the \((M_2, \mu)\) plane \( \text{(i.e., where } m_\chi > m_h) \). These two effects suppress the annihilation rate and hence enhance the relic density relative to the usual calculations (with \[12,13,14,15\] or without \[10\] the supergravity mass relations) where the radiative breaking constraints are not imposed and all the Higgs masses are kept fixed (at moderate values) throughout the \((M_2, \mu)\) plane (for example, compare Fig. 1c with Fig. 2 in Ref. \[14\]).

\[2 \text{ For contours of } m_\chi \text{ and } \chi \text{-composition in the } (M_2, \mu) \text{ plane see } e.g., \text{ Fig. 1 in Ref. } [13].\]

\[3 \text{ We should point out that in Refs. } [13,14] \mu \text{ followed an opposite sign convention than in the present paper and Refs. } [16,20,18].\]
For lower values of $\xi_0$ the radiative breaking constraints apply as well but the phenomenological constraints bite more into the allowed region. The left edge of the allowed area is determined primarily by the LEP lower bound on the chargino mass ($m_{\chi^+} > 45 \text{ GeV}$) and the constraint $m_t \gtrsim 90 \text{ GeV}$. For $\xi_0 \approx 0$ the allowed area is further suppressed \[13\] by demanding an electrically neutral lightest supersymmetric particle.

Within the allowed area the value of $\Omega_\chi h_0^2$ grows steadily with $\xi_0$ since larger values of $m_{\tilde{q}, \tilde{l}}^2 \approx m_{1/2}^2 (c_\chi + \xi_0^2)$ suppress the important $t$- and $u$-channel sfermion-mediated $\chi \chi \to f \bar{f}$ ($f = q, l$) annihilation channels. On the other hand, $\xi_A$ does not affect the relic density in any significant way since $A$ mainly determines the degree of left-right stop mixing. However, $\xi_A$ affects the calculated value of $\mu$ (mainly for small $\mu$) \[18\], and therefore it affects the left edge of the allowed area; $\xi_A > 0$ ($\xi_A < 0$) moves the left edge to the left (right).

The effect of $M_2 \approx 0.83 m_{1/2} \approx 0.30 m_{\tilde{g}}$ is to scale up all Higgs and sparticle masses and therefore one would expect $\Omega_\chi h_0^2$ to grow steadily with $m_{1/2}$. Even though this is generally the case, the steady growth can be locally (in $m_{1/2}$) depleted due to the presence of poles and thresholds of the annihilation cross section. The most prominent of these are the $Z$- and $h$-poles ($\chi \chi \to Z, h \to f \bar{f}$) which occur for $m_\chi \approx \frac{1}{2} M_Z$ and $m_\chi \approx \frac{1}{2} m_h$ respectively, and the $\chi \chi \to hh$ threshold for $m_\chi \approx m_h$. In Fig. 2 (top row) we show $\Omega_\chi h_0^2$ versus $m_{\tilde{g}}$ for $\xi_0 = 1, \xi_A = 0$, $m_t = 150 \text{ GeV}$, $\tan \beta = 2, 8$, and both signs of $\mu$ (solid $\mu > 0$, dashed $\mu < 0$). These ‘rays’ correspond to a path of increasing $M_2$ and $|\mu|$ (and therefore $m_\chi$) through the plots in Fig. 1c (with $M_2/|\mu| \approx 0.5$ (0.7) for $\tan \beta = 2$ (8)) such that $m_t = 150 \text{ GeV}$. For $\tan \beta = 2$ and $\mu > 0$ the $Z$-pole occurs for $m_{\tilde{g}} \approx 270 \text{ GeV}$ and the $h$-pole does not occur (since $m_\chi > \frac{1}{2} m_h$ always). For $\mu < 0$ the $h$- and $Z$-poles are close: $m_\chi \approx \frac{1}{2} m_h \approx 37 \text{ GeV}$ for $m_{\tilde{g}} \approx 340 \text{ GeV}$ and $m_\chi \approx \frac{1}{2} M_Z$ for $m_{\tilde{g}} \approx 380 \text{ GeV}$. For $\tan \beta = 8$ the $h$- and $Z$-poles are even closer: $m_\chi \approx \frac{1}{2} M_Z$ for $m_{\tilde{g}} \approx 310$ ($370$) GeV and $m_\chi \approx \frac{1}{2} m_h \approx 50 \text{ GeV}$ for $m_{\tilde{g}} \approx 320$ ($400$) GeV for $\mu > 0$ ($\mu < 0$). These features are evident in the figures. The $\chi \chi \to hh$ threshold is also noticeable as a drop in the rate of increase of $\Omega_\chi h_0^2$. The degree of effectiveness of this new annihilation channel depends on the $\chi$-composition. Pure bino $\chi$‘s do not couple to the $h$ field. This is the case for the $\tan \beta = 2$, $\mu > 0$ ray where at $m_{\tilde{g}} \approx 440 \text{ GeV}$, $m_\chi \approx m_h \approx 70 \text{ GeV}$ but only an almost imperceptible drop occurs. The other rays encounter the threshold when $\chi$ is of mixed composition and therefore the expected drop is clearly observable. We remark that perhaps the most important effect of the one-loop corrections to the Higgs boson masses is
to shift the position of the poles and thresholds discussed above, relative to the tree-level case.

Since \( m_{1/2} \) and \( \xi_0 \) determine to a large extent the magnitude of \( \Omega \chi h_0^2 \), it is possible to scan the parameter space and determine the lowest values of these parameters which give \( \Omega \chi h_0^2 \sim h_0^2 \). We have done this for \( h = 1/2 \) and find that it is necessary to have \( m_{\tilde{q}} \gtrsim 250 \text{ GeV} \) and \( m_{\tilde{l}} \gtrsim 100 \text{ GeV} \) to possibly achieve this. Lower values of \( m_{\tilde{q},\tilde{l}} \) do not suppress the annihilation cross section sufficiently to give a large enough relic density. The bounds are substantially weaker than previously obtained (c.f. \( m_{\tilde{q}} \gtrsim 600 \text{ GeV} \) and \( m_{\tilde{l}} \gtrsim 200 \text{ GeV} \) \[13\]) without imposing the radiative breaking constraints since (as discussed above) now the important Higgs-dependent annihilation channels also fade away with increasing \( m_{1/2} \).

The effect of \( \tan \beta \) on \( \Omega \chi h_0^2 \) is felt in the \( \chi \)-composition for fixed \( M_2 \) and \( \mu \), which affects the annihilation rate directly. Heuristically we find that \( \Omega \chi h_0^2 \) generally decreases for large values of \( \tan \beta \). This variable also affects the size of the allowed area in the \((M_2, \mu)\) plane. The allowed area first increases with \( \tan \beta \), it peaks in size, and then it decreases for large \( \tan \beta \) until it disappears completely \[18\].

From Fig. 1 and the preceding discussion it is clear that for \( \xi_0 \lesssim 1 \) it is quite possible to obtain \( \Omega \chi h_0^2 \sim h_0^2 \), i.e., pluses, squares, diamonds for \( h = 1, 1/\sqrt{2}, 1/2 \). As \( \xi_0 \) decreases the fraction of the allowed area occupied by ‘preferred’ points decreases and for \( \xi_0 \lesssim 0.2 \) it is no longer possible to obtain \( \Omega \chi h_0^2 \sim h_0^2 \) for \( 0.5 \leq h_0 \leq 1 \). However, \( \xi_0 \ll 1 \) may still provide enough neutralino relic density for it to be the major component of the galactic halo.

3. The excluded regions

We have determined that for \( \xi_0 \lesssim 1 \) and \( m_{\tilde{g}} < 1 \text{ TeV} \) there are no points in parameter space where \( \Omega \chi h_0^2 > 1 \) and therefore one cannot conclusively exclude any of these points on cosmological grounds. Note however that if \( h = 1/\sqrt{2} (1/2) \) then \( \Omega \chi h_0^2 > 1/2 (1/4) \) will become disfavored. On the other hand, for \( \xi_0 \gtrsim 2 \) there are always regions of parameter space where \( \Omega \chi h_0^2 > 1 \) for almost all allowed values of \( M_2 \). The fraction of the allowed area which becomes cosmologically excluded increases with \( \xi_0 \).

To explore this dependence in a quantitative way it is better to change variables from \((M_2, \mu)\) to \((M_2, m_t)\) for \( \xi_0 \gtrsim 1 \). This is because for large \( \xi_0 \) values, the range of \( \mu \) is much enlarged, i.e., \( \mu \lesssim \xi_0 m_{1/2} \) (see Eqn. \(2.1\)), whereas \( m_t \) is always restricted to be below
≈ 190 GeV \cite{16,18}. We thus interchange \( \mu \) with the basic parameter \( m_t \). In Fig. 3 we show \( \Omega_\chi h_0^2 \) in the \((M_2, m_t)\) plane for \( \xi_0 = 2, \xi_A = 0, \tan \beta = 2, 8 \), and both signs of \( \mu \). Note the large number of cosmologically excluded points (at least the stars).

The dependences of \( \Omega_\chi h_0^2 \) on the radiative breaking constraints, and on \( \xi_0, \xi_A, m_{1/2} \), and \( \tan \beta \) are qualitatively as discussed for the \( \xi_0 < 1 \) case above. The relic density increases steadily with \( \xi_0 \), except for values of \( M_2 \) where poles and/or thresholds occur. In Fig. 2 (bottom row) we show \( \Omega_\chi h_0^2 \) along a ray with \( m_t = 150 \text{ GeV} \) in Fig. 3. Compared to the top row, the bottom row has very similar features, except for the overall scaling of the vertical axis. The reason for this behaviour (i.e., the little change in \( m_\chi \)) is that \( \mu \) does not change too much for \( \xi_0 = 1 \rightarrow 2 \) and the \( m_\chi \) contours in the \((M_2, \mu)\) plane (see e.g., Fig. 1 in Ref. \cite{13}) tend to change little with \( \mu \) (for large enough \( \mu \)); the change (and the corresponding shifts in \( m_\chi \)) are larger for \( \mu < 0 \). Note that in contrast with the \( \tan \beta = 2, \mu > 0, \xi_0 = 1 \) case, for \( \xi_0 = 2 \) the \( h \)-pole does occur in this case for \( m_\tilde{g} \approx 150 \text{ GeV} \) and \( m_\chi \approx \frac{1}{2} m_h \approx 27 \text{ GeV} \). Regarding the \( \chi \chi \rightarrow hh \) thresholds in this case (see Fig. 2 bottom row), it is interesting to note that for \( \tan \beta = 2, \mu > 0 \), the threshold effect, even though small it is nevertheless noticeable (relative to the corresponding top row case in Fig. 2). This is because the relative size of the drop in the annihilation cross section due to the new channel is much larger here (\( \xi_0 = 2 \), small total cross section) than in the previous case (\( \xi_0 = 1 \), large total cross section).

To quantify the statement made above about larger values of \( \xi_0 \) giving larger cosmologically excluded fractions of the allowed region, we define \( f_h \equiv n(\Omega_\chi h_0^2 > h_0^2)/n_{tot} \), where \( n(\Omega_\chi h_0^2 > h_0^2) \) is the number of points inside the allowed region which have \( \Omega_\chi h_0^2 > h_0^2 \) and \( n_{tot} \) is the total number of points for a fixed discrete gridding of the \((M_2, m_t)\) space (with \( \Delta m_t = \Delta M_2 = 5 \text{ GeV} \)). In Fig. 4 we plot \( f_{h=1/2} \) versus \( \xi_0 \) for \( \xi_A = -\xi_0, 0, +\xi_0 \), \( \tan \beta = 2, 8, 15, 30 \), and both signs of \( \mu \). Due to the large extent of the computations to produce Fig. 4, these were performed in the tree-level approximation to radiative breaking and the Higgs boson masses (a full one-loop treatment would have required \( O(\text{days}) \) of NEC SX-3 supercomputer CPU time). However, in this particular instance we do not expect the full one-loop treatment to modify significantly the results in Fig. 4 (as verified explicitly in some cases) because its main effect is to shift the positions of poles and thresholds, which basically do not affect the value of \( f_h \). As anticipated, as \( \xi_0 \) grows the fraction of excluded points grows, and would probably approach \( \approx 100\% \) sooner were it not for the poles. In fact, from Fig. 2 one can see that increasing \( \xi_0 \) reduces the size of the interval in \( m_\tilde{g} \) where \( \Omega_\chi h_0^2 < h_0^2 \) (= 0.25 in this case). Because these regions can only
shrink but not disappear completely, it is not possible to obtain an absolute upper bound on $\xi_0$ from cosmological considerations alone. However, large specific regions of parameter space can be ruled out in this way.

4. Conclusions

We have presented an exhaustive exploration of the parameter space of the SSM which due to the constraints of universal soft supersymmetry breaking and radiative breaking of the electroweak symmetry needs only five variables to be described. The latter constraint imposes novel relations between the traditional relic density variables and allows the determination of the Higgs boson masses in terms of the basic parameters of the model. These two new effects affect significantly the variation of $\Omega_\chi h_0^2$ throughout the allowed parameter space. Furthermore, our calculations include the state-of-the-art one-loop Higgs effective potential to enforce the radiative breaking constraints and obtain automatically the one-loop corrected Higgs boson masses. As such we believe this to be the ultimate calculational framework upon which more specific supergravity models would need to be investigated (see e.g., [22]).

We find that $\Omega_\chi h_0^2$ depends most strongly on the parameter $\xi_0 = m_0/m_{1/2}$, as previously observed [14]. For $\xi_0 \gg 1$ the relic density is almost always (barring accidental depletion due to $Z$- and/or $h$-poles) in conflict with current cosmological observations. For $\xi_0 \ll 1$ the opposite occurs, although $\Omega_\chi h_0^2$ may still account for the dark matter in the galactic halo. For $\xi_0 \sim 1$ there is a wide range of the other parameters for which $\Omega_\chi \sim 1$. Besides these qualitative observations, one can draw more precise conclusions for specific points or regions of parameter space. It is also possible to obtain an approximate lower bound on the squark and slepton masses below which it is not possible to get $\Omega_\chi \sim 1$; we find $m_{\tilde{q}} \gtrsim 250\,\text{GeV}$ and $m_{\tilde{l}} \gtrsim 100\,\text{GeV}$. These lower bounds are much weaker than those obtained previously [13] without enforcing the radiative breaking constraints. Since on naturalness grounds [13] we have limited $m_{\tilde{g}}$ to be below 1 TeV, the value of $m_\chi$ is also restricted to $m_\chi \lesssim 150\,\text{GeV}$. These new bounds offer added hope that the supersymmetric spectrum may be discovered at the next generation of $ee$ and $pp$ machines.

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Figure Captions

Fig. 1. The neutralino relic density $\Omega_{\chi} h^2_0$ in the $(M_2, \mu)$ plane for $\tan \beta = 2, 8$, $\xi_A = 0$, both signs of $\mu$, and (a) $\xi_0 = 0$, (b) $\xi_0 = 0.4$, (c) $\xi_0 = 1.0$. The various symbols denote the following ranges of $\Omega_{\chi} h^2_0$: $< 0.025$ (dots), $0.025 - 0.10$ (crosses), $0.10 - 0.25$ (diamonds), $0.25 - 0.50$ (squares), $0.50 - 1.0$ (pluses), $> 1.0$ (stars). For $h = 1, 1/\sqrt{2}, 0.5$ stars, stars+pluses, stars+pluses+squares are excluded on cosmological grounds.

Fig. 2. The neutralino relic density $\Omega_{\chi} h^2_0$ as a function of the gluino mass $m_{\tilde{g}}$ along rays with $m_t = 150$ GeV in Fig. 1c ($\xi_0 = 1, \xi_A = 0$) and Fig. 3 ($\xi_0 = 2, \xi_A = 0$). (In Fig. 1c these rays correspond to $M_2/|\mu| \approx 0.5 (0.7)$ for $\tan \beta = 2 (8)$.) Solid (dashed) lines correspond to $\mu > 0$ ($\mu < 0$). Note the various poles and thresholds of the neutralino annihilation cross section.

Fig. 3. The neutralino relic density $\Omega_{\chi} h^2_0$ in the $(M_2, m_t)$ plane for $\xi_0 = 2, \xi_A = 0$, for $\tan \beta = 2, 8$ and both signs of $\mu$. The symbols are as in Fig. 1. For $h = 1, 1/\sqrt{2}, 0.5$ stars, stars+pluses, stars+pluses+squares are excluded on cosmological grounds.

Fig. 4. The fraction of the allowed parameter space in the $(M_2, m_t)$ plane which is cosmologically excluded for $h = 1/2$ ($f_{h=1/2}$) versus $\xi_0$ for $\xi_A = -\xi_0, 0, +\xi_0$, both signs of $\mu$ and $\tan \beta = 2$ (solid), 8 (dashed), 15 (dotdashed), 30 (dotted).