Scenarios for gluino coannihilation

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Abstract: We study supersymmetric scenarios in which the gluino is the next-to-lightest supersymmetric particle (NLSP), with a mass sufficiently close to that of the lightest supersymmetric particle (LSP) that gluino coannihilation becomes important. One of these scenarios is the MSSM with soft supersymmetry-breaking squark and slepton masses that are universal at an input GUT renormalization scale, but with non-universal gaugino masses. The other scenario is an extension of the MSSM to include vector-like supermultiplets. In both scenarios, we identify the regions of parameter space where gluino coannihilation is important, and discuss their relations to other regions of parameter space where other mechanisms bring the dark matter density into the range allowed by cosmology. In the case of the non-universal MSSM scenario, we find that the allowed range of parameter space is constrained by the requirement of electroweak symmetry breaking, the avoidance of a charged LSP and the measured mass of the Higgs boson, in particular, as well as the appearance of other dark matter (co)annihilation processes. Nevertheless, LSP masses $m_\chi \lesssim 8\text{ TeV}$ with the correct dark matter density are quite possible. In the case of pure gravity mediation with additional vector-like supermultiplets, changes to the anomaly-mediated gluino mass and the threshold effects associated with these states can make the gluino almost degenerate with the LSP, and we find a similar upper bound.

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1 Introduction

The absence of supersymmetry so far, at LHC Run I [1–3] and elsewhere, raises the question where, if anywhere, is it hiding. There are scenarios for which, at least, some supersymmetric particles were produced in LHC Run I, but have been overlooked. Examples include models where $R$-parity is violated [4, 5], or the spectra are compressed [6–9]. Alternatively, sparticles might be too heavy to have been detected at LHC Run I, but might be within range of future LHC runs [10]. It is also possible that supersymmetric particles may lie beyond the reach of the LHC altogether, and require a future higher-energy $pp$ collider for their detection.

If one assumes that $R$-parity is conserved, the lightest supersymmetric particle (LSP) must be stable, and hence makes at least a contribution to the cosmological cold dark matter density [11, 12]. The total density of cold dark matter is very tightly constrained by measurements of the cosmic microwave background radiation [13]. It is clear, therefore, that the parameters of generic models are constrained in very specific ways in order to realize the correct dark matter density [14–31]. Moreover, this parameter space with the correct density is likely to be found in a region of parameter space where the density varies rapidly with the parameters. In these cases, regions where the LSP contributes only a fraction of the cold dark matter density will have parameters similar to those regions yielding the correct total density.

This sensitivity of the dark matter density to parameters are particularly relevant for models with compressed and/or very heavy spectra that have survived LHC searches. Examples of specific choices of heavy spectra that yield the correct cosmological dark matter density include scenarios in which the LSP, $\chi$, would have annihilated with itself.

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1 Full ATLAS Run 1 results can be found at https://twiki.cern.ch/twiki/bin/view/AtlasPublic/SupersymmetryPublicResults.
2 Full CMS Run 1 results can be found at https://twiki.cern.ch/twiki/bin/view/CMSPublic/PhysicsResultsSUS.
through a direct-channel boson such as the heavier neutral Higgs bosons $A$ and $H$ \cite{32-36}. Alternatively, there might be one or more heavier supersymmetric particles that are nearly degenerate with the LSP, $\chi$, and would have coannihilated with it in the early Universe \cite{37}. There are several examples of possible coannihilating sparticles, including the lighter stau, or possibly some other slepton \cite{38-45}, the lighter stop squark \cite{46-56}, the lighter chargino \cite{55,57-60} and the gluino \cite{56,61-72}.

In most cases, coannihilation with a sparticle having stronger interactions extends the allowed mass range of the LSP. The possibility of gluino coannihilation is therefore particularly interesting since it interacts strongly suggesting it can accommodate a heavier LSP than is possible from coannihilation with a stau or slepton. In fact, it has been shown that a dark matter density realized by an LSP coannihilating with the gluino could lie well beyond the reach of the LHC, with a mass as heavy as $m_\chi \lesssim 8 \text{ TeV}$ \cite{72}.

The possibility of gluino coannihilation does not arise in the minimal supersymmetric extension of the Standard Model (MSSM) with the soft supersymmetry-breaking parameters constrained to be universal at the input GUT scale (the CMSSM) \cite{14,32-36,73-92}, nor in related models with non-universal Higgs masses \cite{83,84,93-106}. However, as we discuss in this paper, gluino coannihilation can become important in variants of the MSSM with non-universal gaugino masses, and in variations of pure gravity mediation (PGM) with non-minimal matter content such as additional vector-like supermultiplets \cite{67-69}.

We use the SSARD (SuperSymmetry And Relic Density) code\textsuperscript{3} to calculate the particle spectrum and relic density. SSARD first calculates the supersymmetric particle spectrum for a given set of boundary conditions defined by the model — the modified CMSSM or PGM. Coupled renormalization-group equations (RGEs) are then run back and forth between the weak scale and the GUT scale, which is defined by the renormalization scale where the two electroweak gauge couplings are equal. The gauge and Yukawa couplings are run at two loops, whereas the soft supersymmetry-breaking masses are run at one loop. The strong gauge coupling is fixed at the weak scale. Once convergence of the RGEs is obtained, the sfermion mass matrices are evaluated to obtain physical masses. SSARD determines $\mu$ and $B\mu$ at the weak scale by minimizing the Higgs tadpole equations. We calculate the Higgs mass using the procedure outlined in \cite{107,108}.\textsuperscript{4} The neutralino and chargino mass matrices are then diagonalized with one-loop corrections applied. With the sparticle spectrum determined, the cross sections for annihilation and co-annihilation are computed and input into a routine that integrates numerically the Boltzmann equation to determine the cosmological relic density. Finally, branching fractions for rare decays are computed, along with the value of $g_\mu - 2$ and the neutralino-nucleon scattering cross section.

\textsuperscript{3}Information about this code is available from K.A. Olive: it contains important contributions from J. Evans, T. Falk, A. Ferstl, G. Ganis, F. Luo, A. Mustafayev, J. McDonald, K.A. Olive, P. Sandick, Y. Santos, V. Spanos, and M. Srednicki.

\textsuperscript{4}This differs from most implementations of SSARD, which often uses FeynHiggs \cite{109-114} to calculate the Higgs mass. However, FeynHiggs loses stability at mass scales significantly above 10 TeV let alone the order 100–1000 TeV mass scales considered here. In the implementation used here, the Higgs mass is computed from its effective quartic coupling as discussed in \cite{107,108}. For more information on how the Higgs mass calculation is used in SSARD see \cite{115}.
The layout of this paper is as follows. In section 2 we set out the coupled set of Boltzmann equations that we use to calculate the relic LSP density, discussing the circumstances under which the analysis can be reduced to a single Boltzmann equation for a particular combination of sparticle abundances [72]. Then, in section 3 we discuss various scenarios with non-universal gaugino masses in which gluino coannihilation can become important, delineating the corresponding strips in parameter space and comparing their extents with the results of [72]. We find that this scenario is constrained by the requirement of consistent electroweak symmetry breaking (EWSB), by the measurement of $m_H$, and by avoidance of a stop or chargino LSP. We give examples showing that the correct dark matter density is possible with LSP masses as large as 8 TeV. Section 4 contains a similar analysis of PGM models with vector-like supermultiplets, focusing on an example with a single extra pair of $10$ and $\bar{10}$ representations of SU(5). Because the anomaly-mediated contribution to the gluino mass is zero in this case, threshold effects due to these additional states generate almost the entire gluino mass. This suppresses the gluino mass relative to those of the other gauginos, leading to near-degeneracy between the gluino and the LSP. Neutralino dark matter candidates with similarly large values of $m_\chi$ are again possible. Finally, section 5 summarizes our conclusions and discusses the prospects for discovering supersymmetry in these gluino coannihilation scenarios.

2 Calculations of gluino coannihilation

In this section, we present general formulae for calculating the dark matter thermal relic density, and then specialize it to the case of the gluino coannihilation scenarios we consider in this paper, taking into account the effects of gluino-gluino bound states.

We consider $N$ $R$-odd species in the thermal bath in the early Universe. We assume that the rates for interconverting the LSP (which is labeled as the first species, with mass $m_1$) and the first $l$ species ($1 \leq l \leq N$) are sufficiently large, compared to the Hubble expansion rate, that to a very good approximation the ratios of densities are equal to the equilibrium ratios: $n_i/n_1 = n_i^{eq}/n_1^{eq}$ for ($i = 1, \ldots, l$).

For any of the $N$ species, the evolution of its number density is governed by the Boltzmann equation

$$\frac{dn_i}{dt} + 3H(T)n_i = -\sum_{j=1}^{N} \langle \sigma v \rangle_{ij \rightarrow SM} \left( n_i n_j - n_i^{eq} n_j^{eq} \right) - \sum_{j=1}^{N} \langle \Gamma \rangle_{i \rightarrow j} \left( n_i - n_i^{eq} \frac{n_j}{n_j^{eq}} \right),$$  \hspace{1cm} (2.1)

where $\langle \sigma v \rangle_{ij \rightarrow SM}$ is the product of the thermally-averaged relative velocity with the total cross section for the channels of $i$ and $j$ (co)annihilating into Standard Model particles, and $\langle \Gamma \rangle_{i \rightarrow j}$ is the sum of all the thermally-averaged decay and conversion rates for decay and conversion processes for which there is one particle $i$ in the initial state and one particle $j$ in the final state, with all other particles involved in these processes being Standard Model particles. The relations between the thermally-averaged forward and backward reactions are used in (2.1), and we assume for all the Standard Model particles involved
that $n_{SM} = n_{SM}^{eq}$, so that, for example,
\[
\langle \Gamma \rangle_{i \rightarrow j} n_{i}^{eq} = \langle \Gamma \rangle_{j \rightarrow i} n_{j}^{eq}.
\] (2.2)

Written in terms of the yields, $Y_{i} \equiv n_{i}/s$, (2.1) becomes
\[
\frac{dY_{i}}{dx} = - \frac{x}{H(m_{1})} \left( 1 + \frac{T}{3g_{ss}} \right) \left[ \sum_{j=1}^{N} \langle \sigma v \rangle_{ij} \langle Y_{i}Y_{j} \rangle + \sum_{j=1}^{N} \sum_{j_{n_{eq}^{i}}}^{N} \frac{\langle \Gamma \rangle_{i \rightarrow j}}{s} \left( Y_{i} - Y_{i}^{eq} \frac{Y_{j}}{Y_{j}^{eq}} \right) \right],
\] (2.3)
where
\[
x \equiv \frac{m_{1}}{T}, \quad s = \frac{2\pi^{2}}{45} g_{ss} T^{3}, \quad H(m_{1}) \equiv H(T) x^{2} = \left( \frac{4\pi^{3} G_{N} g_{s}}{45} \right)^{\frac{1}{2}} m_{1}^{2},
\] (2.4)
and $g_{ss}$ and $g_{s}$ are the total numbers of effectively massless degrees of freedom associated with the entropy density and the energy density, respectively.

Defining $\bar{Y} \equiv \sum_{i=1}^{l} Y_{i}$, $\Delta_{i} \equiv (m_{i} - m_{1})/m_{1}$ and $\tilde{g}_{\text{eff}} \equiv \sum_{i=1}^{l} g_{i} (1 + \Delta_{i})^{3/2} e^{-\Delta_{i} x}$, we have
\[
\frac{Y_{i}^{eq}}{\bar{Y}^{eq}} = \frac{g_{i} \left( \frac{m_{1} T}{2\pi} \right)^{3/2} e^{-m_{i}/T}}{\sum_{j=1}^{l} g_{j} \left( \frac{m_{1} T}{2\pi} \right)^{3/2} e^{-m_{j}/T}} = \frac{g_{i} (1 + \Delta_{i})^{3/2} e^{-\Delta_{i} x}}{\tilde{g}_{\text{eff}}},
\] (2.5)
Using $Y_{i}/\bar{Y} = Y_{i}^{eq}/\bar{Y}^{eq}$ for $i = 1, \ldots, l$ and summing over (2.3) for the first $l$ species, we find
\[
\frac{d\bar{Y}}{dx} = - \frac{x}{H(m_{1})} \left( 1 + \frac{T}{3g_{ss}} \right) \left\{ \langle \bar{\sigma}_{\text{eff}} v \rangle \left( \bar{Y}^{2} - \bar{Y}_{\text{eq}}^{2} \right) \right. \\
+ \left. \sum_{k=l+1}^{N} \left( \langle \sigma k_{\text{eff}} v \rangle \left( \bar{Y} Y_{k} - \bar{Y}_{\text{eq}} Y_{k}^{eq} \right) + \frac{\langle \Gamma k_{\text{eff}} \rangle}{s} \left( \bar{Y} - \bar{Y}_{\text{eq}} \frac{Y_{k}}{Y_{k}^{eq}} \right) \right) \right\},
\] (2.6)
while for each of the species $k$ ($l < k \leq N$), we get
\[
\frac{dY_{k}}{dx} = - \frac{x}{H(m_{1})} \left( 1 + \frac{T}{3g_{ss}} \right) \left\{ \langle \sigma k_{\text{eff}} v \rangle \left( \bar{Y} Y_{k} - \bar{Y}_{\text{eq}} Y_{k}^{eq} \right) \right. \\
+ \left. \sum_{j=l+1}^{N} \langle \sigma v \rangle_{kj} \langle Y_{k} Y_{j} \rangle - \langle \Gamma k \rangle_{\text{eff}} \langle Y_{k} \rangle \left( \bar{Y} - \bar{Y}_{\text{eq}} \frac{Y_{k}}{Y_{k}^{eq}} \right) \right) \right\},
\] (2.7)
where

\[
\langle \sigma_{\text{eff}} v \rangle \equiv \sum_{i=1}^{l} \langle \sigma v \rangle_{i \rightarrow \text{SM}} \frac{Y_{i}^{\text{SM}}}{Y_{\text{eq}}} = \sum_{i=1}^{l} \langle \sigma v \rangle_{i \rightarrow \text{SM}} \frac{g_{i}(1+\Delta_{i})^{3/2}e^{-\Delta_{i}x}}{\tilde{g}_{\text{eff}}},
\]

(2.8)

\[
\langle \Gamma_{k} \rangle_{\text{eff}} \equiv \sum_{i=1}^{l} \langle \Gamma \rangle_{i \rightarrow k} \frac{Y_{i}^{\text{eq}}}{Y_{\text{eq}}} = \sum_{i=1}^{l} \langle \Gamma \rangle_{i \rightarrow k} \frac{g_{i}(1+\Delta_{i})^{3/2}e^{-\Delta_{i}x}}{\tilde{g}_{\text{eff}}},
\]

(2.9)

\[
\langle \tilde{\sigma}_{\text{eff}} v \rangle \equiv \sum_{i,j=1}^{l} \langle \sigma v \rangle_{i \rightarrow \text{SM}} \frac{Y_{i}^{\text{eq}} Y_{j}^{\text{eq}}}{Y_{\text{eq}}^{2}} = \sum_{i,j=1}^{l} \langle \sigma v \rangle_{i \rightarrow \text{SM}} \frac{g_{i}g_{j}(1+\Delta_{i})^{3/2}(1+\Delta_{j})^{3/2}e^{-(\Delta_{i}+\Delta_{j})x}}{\tilde{g}_{\text{eff}}^{2}}.
\]

(2.10)

We note that in the case $N = l + 1$, the final term in (2.7) does not appear. In the case $N = l$, (2.6) does not have the two terms in the squared bracket, and reverts to the familiar form for coannihilations when all the $N$ species are sufficiently coupled to the LSP.

We now specialize the above general formulae to the gluino coannihilation scenarios we consider in this paper. First of all, following the discussion in [72], the effect of gluino-gluino bound states on the calculation of the dark matter relic density can be taken into account simply by modifying the Boltzmann equation by including the Sommerfeld-enhanced thermal-averaged velocity-weighted gluino pair annihilation cross section [70], which includes gluino-pair annihilation to two gluons and to all the quark anti-quark pair channels:

\[
\langle \sigma v \rangle_{g\bar{g} \rightarrow g\bar{g},q\bar{q}} \rightarrow \langle \sigma v \rangle_{g\bar{g} \rightarrow g\bar{g},q\bar{q}}^{\text{incl.}, R} \equiv \langle \sigma v \rangle_{g\bar{g} \rightarrow g\bar{g},q\bar{q}} + \langle \sigma v \rangle_{bsf} \frac{\langle \Gamma \rangle_{R}}{\langle \Gamma \rangle_{R} + \langle \Gamma \rangle_{\text{dis}}},
\]

(2.11)

where $\langle \sigma v \rangle_{bsf}$, $\langle \Gamma \rangle_{R}$ and $\langle \Gamma \rangle_{\text{dis}}$ are the thermally-averaged formation cross section, decay rate and dissociation rate for the bound state $R$, respectively. The details of these quantities and the derivation of eq. (2.11) can be found in section 3, 5 and appendix B of [72].

When the rate for interconverting the neutralino LSP and the gluino is sufficiently large, compared to the Hubble rate, so that to a good approximation the relation $Y_{\tilde{g}}(T)/Y_{1}(T) = Y_{\tilde{g}}^{\text{eq}}(T)/Y_{1}^{\text{eq}}(T)$ holds at all temperatures during which the sum of $Y_{\tilde{g}}(T)$ and $Y_{1}(T)$ changes non-negligibly, we can use a single Boltzmann equation to solve for the dark matter relic abundance, including the gluino species in $\tilde{Y}$ and $\langle \tilde{\sigma}_{\text{eff}} v \rangle$ and using (2.6) without the two terms in the squared bracket. Otherwise, one should use a coupled set of Boltzmann equations, namely (2.6) and (2.7), to solve for the dark matter relic abundance. For the scenarios considered in this paper, any of the $R$-odd species apart from the gluino is either sufficiently coupled to the LSP by having a Standard Model particle in the propagator of a tree-level Feynman diagram describing its interconversion with the LSP, or is so heavy compared to the LSP that it is effectively not participating coannihilations. Therefore, when using a coupled set of Boltzmann equations, we have $N = l + 1$ in (2.6) and (2.7), and the index $k$ is for the gluino.

We end this section by emphasizing that, in principle, the coupled set of Boltzmann equations can always be used to solve for the dark matter relic density, whether the rate for interconverting the gluino and the LSP is sufficiently large compared to the Hubble
rate or not. However, for the former case, solving a single Boltzmann equation is usually easier than solving the coupled ones and requires less computing time.

3 The non-universal MSSM scenario

It was assumed in [72] that the squarks were all degenerate with a common mass \( m_{\tilde{q}} \), and the effects of sparticles with only electroweak interactions were neglected. It was found in [72] that in the presence of gluino coannihilation, a Bino LSP, \( \chi \), could be the dark matter of the universe if it weighed \( \lesssim 8 \) TeV, the exact value depending on the ratio \( m_{\tilde{g}}/m_{\chi} \), with smaller values of \( m_{\chi} \) being found for \( m_{\tilde{g}}/m_{\chi} \lesssim 5 \) and \( \gtrsim 100 \). Here we make a more complete study in a variant of the MSSM with universal soft supersymmetry-breaking scalar masses \( m_0 \) and trilinear couplings \( A_0 \), allowing a restricted form of non-universality in the gaugino sector with \( M_1 = M_2 \neq M_3 \) at the input GUT scale. The results therefore depend on \( M_1/M_3 \) as well as the usual CMSSM parameters \( m_0, A_0 \) and \( \tan \beta \) (the ratio of MSSM Higgs vev’s). This is therefore a one-parameter extension of the CMSSM (with the new free parameter being \( M_3 \)) as is the NUHM1 (with the soft Higgs masses \( m_1 = m_2 \neq m_0 \)) [104–106]. We consider in this section various \((M_1, M_3)\) planes for various choices of the other parameters which illustrate the range of possibilities.

We first consider the example with \( m_0 = 1000 \) TeV, \( A_0/m_0 = 1.5 \) and \( \tan \beta = 2.5 \) shown in figure 1. In the left panel and in subsequent figures, the regions where the relic LSP density \( \Omega_h \) falls within the range allowed by Planck and other data are shown as dark blue strips, and the regions where the lightest neutralino is no longer the LSP are shaded brick-red. In this case, the gluino is the LSP in the shaded region. Because of the scale of the plot, it is difficult to discern the relic density strip, which lies very close to the boundary of this region. However, we note that it lies to the left of the red shaded region only when \( M_3 \) is between \( \sim 400 \) and \( \sim 1200 \) GeV, as shown in the right panel of figure 1 by the left axis and blue curve, which shows the mass difference \( \Delta M \equiv m_{\tilde{g}} - m_{\chi} \) between the gluino and the neutralino along the coannihilation strip as a function of the input gluino mass. Also shown in the right panel (as a red line) is the neutralino mass as a function of \( M_3 \).

The shapes of the blue and red curves in the right panel of figure 1 can be understood by comparing with figure 6 of [72], which shows the value of \( m_{\chi} \) at the end-point of the gluino coannihilation strip when \( \Delta M = 0 \) for a pure Bino LSP, as a function of the ratio of the assumed common squark mass, \( m_{\tilde{g}} \), and \( m_{\chi} \). We reproduce this plot here (figure 2) for the convenience of the reader. In this plot, the green band corresponds to the 3-\( \sigma \) range allowed by Planck and other data. The drop at small \( m_{\tilde{g}}/m_{\chi} \) is due to the cancellations between the \( s-, t\)- and \( u\)-channel diagrams for gluino-pair annihilation into quark and anti-quark, which results in a smaller annihilation cross section. The very rapid drop at large \( m_{\tilde{g}}/m_{\chi} \) is due to the decoupling of the gluino and neutralino densities. In between there is a plateau with \( m_{\chi} \sim 8 \) TeV along the green band. As seen in figure 2, the choice of \( m_0 \) in figure 1 corresponds to values of \( m_{\tilde{g}}/m_{\chi} \) extending from beyond the plateau at small \( M_3 \) to values along the plateau at large \( M_3 \). The gluino coannihilation strip therefore has two end-points where \( \Delta M \to 0 \), corresponding to the limiting values \( m_{\chi} \sim 6 \) TeV and
$m_\chi \sim 8 \text{ TeV}$ seen in the right panel of figure 1: for larger and smaller $M_3$, $\Delta M < 0$ and the gluino is the LSP.

We list in the first column of table 1 some details of a sample parameter set from the model plane in figure 1. The first few lines display the values of the input parameters $m_0, M_1, M_3, A_0/m_0$ and $\tan \beta$, followed by the Higgs mixing parameter $\mu$ derived from the electroweak vacuum conditions. The next few lines exhibit the masses of the gluino, neutralino LSP, lighter stop squark and Higgs boson, and the last two lines exhibit the spin-independent and -dependent LSP-proton scattering cross sections $\sigma_{SI,SD}$. All the mass parameters are given in TeV units, except for the Higgs mass, which is expressed in GeV. The succeeding columns in table 1 show the corresponding numbers for parameter sets from the model planes shown in later figures. In the cases of the pure gravity mediation (PGM) models in the last two columns, the values of $c_{10}$ are quoted in parentheses (\ldots).

We note that the Higgs mass is relatively insensitive to the choice of $M_1$ and $M_3$, and therefore varies very little across the plane with $\tan \beta, A_0$, and $m_0$ fixed. For the case shown in figure 1, we calculate $m_H \approx 126.3$ GeV, which is compatible with the experimental measurement, within the theoretical uncertainties. We do not show any other $(M_1, M_3)$ planes for $m_0 = 1000$ TeV, since the possibilities are quite limited: there are no consistent solutions of the electroweak symmetry-breaking conditions for much smaller values of $A_0/m_0 \lesssim 1$ and/or larger values of $\tan \beta$, and $m_H$ is too large for larger values of $\tan \beta$ and/or $A_0/m_0$ (though it increases quite slowly with $A_0$).

We consider next an example of a $(M_1, M_3)$ plane for $m_0 = 200$ TeV, which corresponds to values of $m_\tilde{q}/m_\chi$ along the plateau in figure 2. The left panel of figure 3 shows the $(M_1, M_3)$ plane for $\tan \beta = 3$ and $A_0/m_0 = 1.5$. In this case there is a longer gluino
Table 1. For the sample parameter sets taken from the figures noted in the top line, the values of the input parameters $m_0, M_1, M_3, A_0/m_0$ (or $c_{10}$) and $\tan \beta$ are shown in the next five lines, followed by the Higgs mixing parameter $\mu$ derived from the electroweak vacuum conditions. Points 10 and 11 come from the PGM models with vector matter and do not have gaugino mass inputs. The masses of the gluino, neutralino LSP, lighter stop squark and Higgs boson are shown in the next four lines, and the spin-independent and -dependent LSP-proton scattering cross sections $\sigma_{SI, SD}$ are shown in the last two lines. All the mass parameters are expressed in TeV units, except for the Higgs mass, which is expressed in GeV.
Figure 2. This figure is adapted from figure 6 of [72]. It shows the value of $m_\chi$ at the end-point of the gluino coannihilation strip when $\Delta M = 0$ for a pure Bino LSP, as a function of $m_{\tilde{q}}/m_\chi$, where a common mass $m_{\tilde{q}}$ is assumed for all the squarks. The drop at small $m_{\tilde{q}}/m_\chi$ is due to the cancellations between the s-, t- and u-channel diagrams for gluino-pair annihilation into quark and anti-quark, and that at large $m_{\tilde{q}}/m_\chi$ is due to the decoupling of the gluino and neutralino densities. The green band corresponds to the 3-σ range of the dark matter density: $\Omega_\chi h^2 = 0.1193 \pm 0.0042$.

Figure 3. As for figure 1, but for $m_0 = 200 \text{ TeV}$ and $A_0/m_0 = 1.5$, with $\tan \beta = 3$. 
coannihilation strip extending nearly all the way to $M_3 \sim 3$ TeV. The panel on the right again shows the gluino-neutralino mass difference $\Delta M$ (blue line) which in this case peaks at approximately 170 GeV, which is consistent with the results of [72] for intermediate squark-to-gluino mass ratios. Also shown is the neutralino mass as a function of $M_3$ (red line): it again rises to $m_\chi \sim 8$ TeV at the tip of the coannihilation strip, which has $M_3$ slightly $> 3$ TeV. We list in the second column of table 1 some details of a sample parameter set from the model plane in figure 3.

The Higgs mass in this case is again very slowly varying across the plane and takes the value $m_H \sim 125$ GeV for this choice of $\tan \beta = 3$. In comparison, had we chosen $\tan \beta = 5$, the $(M_1, M_3)$ plane would look almost identical but with $m_H \sim 131$ GeV. This and larger values of $\tan \beta$ are therefore excluded for this value of $A_0/m_0$. We have also studied smaller values of $A_0/m_0$ and found no consistent solutions of the electroweak symmetry-breaking conditions for $M_3 \lesssim 500$ GeV for $A_0/m_0 = 1$ and no consistent solutions across the plane at somewhat lower $A_0/m_0$. This is also the case for $\tan \beta = 10$ and $A_0/m_0 = 1.5$, for which $m_H \sim 134$ GeV. Larger values of $A_0/m_0$ also give values that tend to increase $m_H$ and, if increased too much, the stop becomes the LSP and eventually tachyonic.

Next we consider some sample $(M_1, M_3)$ planes with $m_0 = 20$ TeV and $\tan \beta = 5$, corresponding to the lower end of the $m_\chi$ plateau in figure 2. Figure 4 is for the case $A_0/m_0 = 1$, where we see in the left panel that electroweak symmetry breaking is possible up to values of $M_1 \lesssim 14$ TeV. There is a gluino coannihilation strip close to the colored LSP boundary for $M_1 \lesssim 9$ TeV. This is terminated by a spur extending to large $M_3$ when $9$ TeV $\lesssim M_1 \lesssim 10$ TeV, where the lighter chargino is the LSP. There is no chargino coannihilation strip along the boundary of this region at large $M_3$, because the relic density is too high: for these values of $M_1$ and $M_3$, the Higgsino mass is too large and other coannihilations are not sufficient to bring the relic density down. At larger values, $10$ TeV $\lesssim M_1 \lesssim 11$ TeV, there is a Higgsino-gluino coannihilation strip, which is followed at larger $M_1$ by a focus-point
strip [116–121] hugging the electroweak symmetry breaking boundary where the neutralino is well-tempered [122]. In this case we see both the 124 and 125 GeV Higgs mass contours and, as in the previous example, $m_H$ is compatible with experiment whenever the dark matter density falls within the allowed range.

Because the relic density strip is a multi-valued function of $M_3$, the structure of the gluino-neutralino mass difference $\Delta M$ (blue curve) and the neutralino mass (red curve) shown in the right panel of figure 4 are more complicated than in the previous cases. After growing to a local maximum $\sim 170$ GeV when $M_3 \sim 1$ TeV, $\Delta M$ starts to fall at larger $M_3$. We then see a change in behaviour at $M_3 \approx 1800$ GeV along the gluino coannihilation strip. Here, the neutralino becomes Higgsino-like and, as $M_1$ is increased, the coannihilation strip tends toward lower $M_3$ with an increasing mass difference, as seen in the lower branch of the blue curve. A Higgsino LSP emerges for larger $M_1$ because it gives a positive contribution to the up Higgs soft mass from renormalization group running. As the up Higgs soft mass goes to zero so does $\mu$ and the Higgsino becomes the LSP. Once $\mu$ is small enough, the Higgsino can be a thermal relic without any assistance in setting the relic density from other particles. In this focus-point-like region, the mass difference increases beyond the range displayed. This behaviour is correlated with the value of $m_\chi$ (red curve), which increases monotonically to $\sim 4$ TeV. When the Bino/Higgsino transition occurs at $M_3 \approx 1800$ GeV, $m_\chi$ doubles back down to $M_3 \approx 500$ GeV. Then, on the focus-point branch of the relic density strip, the LSP is mostly Higgsino, the value of $M_3$ grows, and the lightest neutralino mass takes the characteristic value $m_\chi \sim 1$ TeV. We list in the third, fourth and fifth columns of table 1 some details of sample parameter sets from the model plane in figure 4.

In figure 5 we choose a larger value of $A_0/m_0 = 1.5$, and we see in the left panel a gluino coannihilation strip that extends to $M_1 \sim 14$ TeV, along which $m_H$ varies between 124 and 126 GeV as seen by the three Higgs mass contours. The focus-point Higgsino dark matter region has disappeared, due to the large $A_0$ driving the Higgs mass to large negative values. The end-point of the gluino coannihilation strip is clearly seen in the right panel of figure 5, where $\Delta M \to 0$ (blue curve) at $M_3 \approx 3300$ GeV. Qualitatively, this case is similar to that shown in figure 3, rather than to figure 4 with its truncated gluino coannihilation strip. In this case, the LSP mass (red curve) rises monotonically to $m_\chi \sim 7.5$ TeV at the end-point of the strip. We list in the sixth column of table 1 some details of a sample parameter set from the model plane in figure 5.

In the left panel of figure 6 we display the $(M_1, M_3)$ plane for $m_0 = 20$ TeV and $A_0/m_0 = 2$. We see again a gluino coannihilation strip, but extending only to $M_1 \sim 7.5$ TeV. It is terminated by a stop LSP region that extends to larger values of $M_3$ than those displayed. In principle, one might have expected to see a stop coannihilation strip running up along the boundary of the stop LSP region. However, in this case the relic density is too high along the boundary shown in this figure: as in the chargino case mentioned earlier, the would-be end-point of the stop coannihilation strip lies within the gluino LSP region.

\footnote{A one-parameter extension of the CMSSM with lighter gluinos was considered in [123] in the context of a mixed Higgsino-Bino neutralino.}
The value of $m_H$ is generally higher than in the previous case, though compatible with experiment along all the dark matter strip. In the right panel of figure 6, the curves for $\Delta M$ and $m_\chi$ terminate when the stop becomes the LSP, with $m_\chi < 3.5$ TeV. We list in the seventh column of table 1 some details of a sample parameter set from the model plane in figure 6.

In figure 7, we choose a lower value of $A_0/m_0 = 0.75$ and keep $\tan \beta = 5$. We see, in the left panel, that consistent electroweak symmetry breaking is possible only for relatively large $M_3$ and small $M_1$, and that there is a strip hugging the curved electroweak symmetry-breaking boundary where the LSP has an enhanced Higgsino component. Its relic density is brought into the allowed range by the same mechanism as we discussed in the case of a
well-tempered neutralino. As one can see in the right panel of figure 7, once $M_3$ is large enough for EWSB solutions to exist, the mass of the lightest neutralino (which is mainly a Higgsino) is $m_\chi \sim 1.1$ TeV (red line), almost independent of $M_3$ for values $\gtrsim 1.1$ TeV. The gluino-neutralino mass difference does not play a role in the relic density determination and is not shown here. The red dot-dashed contour shows where $m_H = 125$ GeV: $m_H$ is smaller (larger) above (below) this contour. The Higgs mass is highly compatible with the LHC measurement all along the displayed part of the relic density strip. We list in the eighth column of table 1 some details of a sample parameter set from the model plane in figure 7.

At larger values of $\tan \beta$, the planes would look similar, though $m_H$ would be larger. The $(M_1, M_3)$ plane for $m_0 = 20$ TeV with $\tan \beta = 10$ and $A_0/m_0 = 1$ resembles that in the left panel of figure 7 for $\tan \beta = 5$ and $A_0/m_0 = 0.75$, with a focus-point strip following closely the curved electroweak symmetry breaking boundary. The most notable difference is the Higgs mass $m_H$, which is around 128 GeV and only marginally compatible with experiment after allowing for the theoretical uncertainties. For the same values of $m_0$ and $\tan \beta = 10$, we find no consistent electroweak symmetry breaking for smaller values of $A_0/m_0$, and for larger values we find that $m_H$ is too high. Thus we find no interesting examples of gluino coannihilation for $m_0 = 20$ TeV and $\tan \beta = 10$.

Finally, we consider in figures 8 and 9 two examples for $m_0 = 10$ TeV, corresponding to values of $m_\tilde{q}/m_\chi$ below the $m_\chi$ plateau in figure 2. Figure 8 is for $\tan \beta = 10$ and $A_0/m_0 = 1$ and displays a truncated gluino coannihilation strip extending to $M_1 \sim 4$ TeV, followed by a Higgsino coannihilation strip extending to $M_1 \sim 5$ TeV, and then a focus-point strip extending beyond the limits of the plot. We find that $m_H$ is always compatible with the experimental measurement. This example resembles that of figure 4 for $m_0 = 20$ TeV and $A_0/m_0 = 1$, the main difference being that the chargino spur has disappeared: we see instead a chargino LSP island at $M_1 \sim 6$ TeV and $M_3 \gtrsim 2$ TeV. As in figure 4, we see in
Figure 8. As for figure 1, but for \( m_0 = 10 \text{ TeV}, \tan \beta = 10 \) and \( A_0/m_0 = 1.0 \).

the right panel that the gluino-neutralino mass difference (blue curve) has a multivalued form, and becomes larger as the neutralino becomes more Higgsino-like along the gluino coannihilation strip. At larger \( M_1, M_3 \) drops along the Higgsino-gluino coannihilation strip and the mass difference increases slightly as \( M_3 \) decreases, increasing beyond the displayed range as one moves on to the focus point strip. Similarly, the neutralino mass rises as \( M_1 \) is increased, then falls back to about 1.1 TeV when the LSP is mostly a Higgsino. We list in the ninth, tenth and eleventh columns of table 1 some details of sample parameter sets from the model plane in figure 8.

Figure 9 for \( A_0/m_0 = 1.5 \) displays a more extended gluino coannihilation strip reaching \( M_1 \sim 7 \text{ TeV} \) and \( m_\chi \sim 3.5 \text{ TeV} \), where it is terminated by a stop LSP region. This stop LSP region would dominate for larger values of \( A_0/m_0 \), and the range of \( m_H \) would also become too high. At lower \( \tan \beta \), the figures would look similar, but with a smaller Higgs mass. For example, for \( \tan \beta = 5 \), with \( A_0/m_0 = 1.5 \) (as in figure 9), the Higgs mass would drop by roughly 3 GeV. Lower values of \( \tan \beta \) would have \( m_H \) too small, and lower values of \( A_0/m_0 \) but the same value of \( \tan \beta \) would have no electroweak symmetry breaking solutions, while the stop LSP region would dominate for larger \( A_0/m_0 \). We list in the twelfth column of table 1 some details of a sample parameter set from the model plane in figure 9.

Our analysis of gluino coannihilation in non-universal MSSM scenarios with \( M_1 \neq M_3 \) has shown that large values of \( m_\chi \lesssim 8 \text{ TeV} \) are certainly possible, though restricted by competing mechanisms. This possibility occurs when \( m_0 = 1000 \text{ TeV} \), but only for \( A_0/m_0 \sim 1.5 \) and low values of \( \tan \beta \lesssim 3 \). The possibility of gluino coannihilation becomes more prominent for \( m_0 = 200 \text{ TeV} \), appearing for an extended range of \( \tan \beta \) though still only for \( A_0/m_0 \sim 1.5 \). Gluino coannihilation is also prominent for \( m_0 = 20 \text{ TeV} \) and \( \tan \beta = 5 \), but the focus point may also be important (it becomes dominant when \( \tan \beta = 10 \)), as is Higgsino coannihilation. These features also appear when \( m_0 = 10 \text{ TeV} \). In general, this scenario is constrained at small \( A_0/m_0 \) by the absence of electroweak symmetry breaking, and at large \( A_0/m_0 \) by \( m_H \) and the appearance of a stop LSP.
4 Pure gravity mediation with vector multiplets

Another possible way of realizing a spectrum that can lead to gluino coannihilation is in models with pure gravity mediation [124–131] of supersymmetry breaking with additional vector multiplets that are charged under the SM gauge symmetries [67–69]. Since the mass spectrum of the gaugino sector of PGM is quite constrained, one of the few viable ways to alter this spectrum is through the addition of vector-like matter with all of its mass coming from a Giudice-Masiero (GM) term [132–134] in the Kähler potential. The vector-like matter alters the gaugino mass spectrum in two ways. First, it alters the beta functions of the SM gauge couplings that set the overall coefficient of the anomaly-mediated contribution to the gaugino masses. Secondly, when the additional vector-like matter is integrated out, it generates gauge-mediation-like contributions to the gaugino masses.\footnote{If the masses of the vector-like states come from a term in the superpotential, when the vector multiplets are integrated out the gauge-mediated contribution to the mass exactly cancels the additional anomaly-mediated contribution, so that the anomaly-mediated relations are still valid below the mass of the additional vector states. If the mass comes from a GM term in the Kähler potential, this cancellation no longer occurs.}

Below we consider a PGM model with an additional pair of SU(5) $10$ and $\overline{10}$ multiplets, since this simple extension of PGM automatically generates a spectrum with gluino-Bino coannihilation.

The model we consider here is based on that in [69], whose setup we briefly review here. The effective potential is similar to that for the CMSSM:

$$V = \left| \frac{\partial W}{\partial \phi^i} \right|^2 + \left( A_0 W^{(3)} + B_0 W^{(2)} + \text{h.c.} \right) + m^2_{\phi_i} \phi_i^* \phi_i, \quad (4.1)$$

where $W^{(3)}$ corresponds to the trilinear terms of the superpotential, $W^{(2)}$ contains the bilinear terms of the superpotential, and the $\phi_i$ signify the MSSM fields, with

$$W = (y_e H_1 L e^c + y_d H_1 Q d^c + y_u H_2 Q u^c) + \mu H_1 H_2. \quad (4.2)$$

---

**Figure 9.** As for figure 1, but for $m_0 = 10$ TeV, $\tan \beta = 10$ and $A_0/m_0 = 1.5$. 

The scalar masses are generated through gravity mediation with a minimal Kähler potential, in an identical manner to mSUGRA [135], and hence are equal to \( m_{3/2} \) at the GUT scale. In general, as seen in (4.1), the form of \( W^{(3)} \) dictates the pattern of the trilinear supersymmetry-breaking terms. However, the trilinear couplings are suppressed in PGM models, because the supersymmetry-breaking field is charged. Thus, we take \( A_0 = 0 \).

Since \( A_0 = 0 \) and the Kähler potential is minimal in this model, we have \( B_0 = -m_{3/2} \) for the Higgs fields. This leaves three free parameters, two of which are determined by enforcing the electroweak symmetry-breaking conditions:

\[
\mu^2 = \frac{m_1^2 - m_2^2 \tan^2 \beta + \frac{1}{2} m_Z^2 (1 - \tan^2 \beta) + \Delta_{\mu}^{(1)}}{\tan^2 \beta - 1 + \Delta_{\mu}^{(2)}} ,
\]

and

\[
B\mu = -\frac{1}{2}(m_1^2 + m_2^2 + 2\mu^2) \sin 2\beta + \Delta_B ,
\]

where \( \Delta_{\mu}^{(1)} , \Delta_{\mu}^{(2)} \) and \( \Delta_B \) are loop corrections to the relationships [136–138].

This scenario has a very restricted parameter space that does not, in general, realize electroweak symmetry breaking [131]. Therefore, we add a Giudice-Masiero (GM) term [132–134] for the Higgs fields, which modifies the GUT-scale values of both \( B_0 \) and \( \mu \):

\[
\mu = \mu_0 + c_H m_{3/2} ,
\]

\[
B\mu = -\mu_0 m_{3/2} + 2c_H m_{3/2}^2 .
\]

This additional degree of freedom in the EWSB sector allows us to choose \( \tan \beta \) (in addition to \( m_{3/2} \)) as a free parameter, and one finds viable parameter space as long as \( \tan \beta \lesssim 3 \) [131].

As already mentioned, the extension of this simplest viable version of the PGM scenario that we consider includes an additional \( 10 \) and \( \overline{10} \) of the SU(5) grand-unification group. Because these states are vector-like, the most general form of the Kähler potential is

\[
K = |10|^2 + |\overline{10}|^2 + (c_{10}(10 \cdot \overline{10}) + \text{h.c.}) ,
\]

which includes a GM-like coupling \( c_{10} \) that generates a supersymmetric mixing mass term, \( \mu_{10} \), and a supersymmetry-breaking \( B \) term for the additional vector-like fields. As can be seen from the minimal form of the kinetic terms in the limit \( c_{10} \to 0 \), the additional fields also have a gravity-mediated tree-level soft supersymmetry-breaking mass equal to \( m_{3/2} \). Since some of the SM fields are also contained in a \( 10 \) of SU(5), the SM like fields contained in the additional \( 10 \) can be combined with \( H_2 \), the SU(2) charged piece of the \( 5_H \), into gauge-invariant operators in the superpotential just as is done in the SM. Since \( H_1 \) comes from the \( 5_H \), it can likewise be combined with the SM-charged components of the additional \( \overline{10} \) to form gauge-invariant Yukawa interactions. If we impose only gauge symmetries, the most generic contribution to the superpotential is

\[
W = y_t H_u Q U' + y_d H_d Q U ,
\]
where \( Q' \) and \( U' \) are from the 10 and \( \bar{Q} \) and \( \bar{U} \) are from the \( \bar{10} \). However, to preserve \( R \) symmetry\(^7\) we must take either \( y_l^t = 0 \) or \( y_l^b = 0 \). Here, we take \( y_l^b = 0 \). The interactions proportional to \( y_l' \) contribute to the beta function of the up Higgs soft mass in a similar way to those controlled by \( y_l \). Specifying a comparable value of \( y_l' \) helps drive radiative electroweak symmetry breaking, which in turn allows larger values of \( \tan \beta > 3 \). The extended theory now has four parameters: \( m_{3/2}, \tan \beta, c_{10}, \) and \( y_l' \).

The gaugino masses in these models are generated by anomalies [140–144]. Because the contributions to gaugino masses are proportional to \( m_{3/2} \) times the gauge coupling beta function, the scalar masses tend to be much heavier than the gaugino masses, reminiscent of split supersymmetry [145–149]. With the addition of the 10 and \( \bar{10} \), the anomaly-mediated contributions to the gaugino masses are\(^8\)

\[
M_1 = \frac{48}{5} \frac{g_1^2}{16\pi^2} m_{3/2}, \quad (4.9) \\
M_2 = \frac{g_2^2}{4\pi^2} m_{3/2}, \quad (4.10) \\
M_3 = 0. \quad (4.11)
\]

In addition, the gauginos then get rather large threshold corrections from the 10 and \( \bar{10} \) when they are integrated out, which is in addition to the large threshold correction coming from integrating out the Higgsinos: for more details see [69]. Since the only contribution to the mass of the gluino comes from the threshold corrections, it tends to be lighter than in typical PGM models. Hence there are regions where the gluino can coannihilate with the Bino, yielding the possibility of a relatively heavy Bino dark matter candidate.

Our results for the PGM model with vector 10 and \( \bar{10} \) multiplets can be displayed in \((c_{10}, m_{3/2})\) planes for fixed values of the Yukawa coupling, \( y_l' \) and \( \tan \beta \). Two examples of these planes are shown in figures 10 and 11. In the former, we have fixed \( \tan \beta = 3 \) (mainly to get an acceptable value for the Higgs mass, \( m_H \), over the range of \( m_{3/2} \leq 600 \text{ TeV} \) shown) and \( y_l'^2 = 0.15 \). In the left panel, we see a large red shaded region at small \( c_{10} \) where the gluino is the LSP. To the right of this boundary, we see the gluino coannihilation strip.\(^9\) In the lower right corner, the pink shaded region is excluded because one or more of the scalar components of the new vector matter have become tachyonic. As in previous figures, the Higgs mass contours are shown as red dot-dashed curves as labelled. Within the theoretical uncertainties, the Higgs mass agrees with experiment over the part of the

\(^7\)We recall that \( R \) symmetry is an important part of dynamical supersymmetry breaking [139]. Since we wish ultimately to have supersymmetry broken dynamically, we consider a theory that preserves \( R \) symmetry. We note furthermore that this \( R \) symmetry can also play an important part in generating the PGM spectrum, since \( R \) symmetry can forbid the tree-level gaugino masses generated in supergravity.

\(^8\)The anomaly-mediated gaugino masses are proportional to the gauge coupling beta functions. The addition of a 10 and \( \bar{10} \) alters the gauge coupling beta functions and so alters the gaugino masses. As it turns out, the additional contribution of the 10 and \( \bar{10} \) to the SU(3) beta function completely cancels the SM contribution at one-loop. Therefore, the total anomaly mediated contribution to the gluino mass is zero at one-loop.

\(^9\)The strip becomes less well defined at \( m_{3/2} \geq 350 \text{ TeV} \) due to inaccuracies of the relic density calculation at such large masses.
plane that is shown. In the right panel we see, as before, the gluino-neutralino mass difference $\Delta M$ along the gluino coannihilation strip (blue) and the neutralino mass along the strip (red). We see that the curve for $\Delta M$ has the same characteristic shape due to strong coannihilations involving the gluino and peaks at $\simeq 170$ GeV at $m_{3/2} \simeq 200$ TeV when $m_{\chi} \simeq 3$ TeV. The end-point of the coannihilation strip occurs at $m_{3/2} \simeq 500$ TeV where $m_{\chi} \simeq 8.3$ TeV. We list in the last two columns of table 1 some details of sample parameter sets from the model planes in figures 10 and 11. We recall that $A_0 = 0$ in these models, and display in parentheses (…) the model values of $c_{10}$.

In figure 11, we show in the left panel the corresponding $(c_{10}, m_{3/2})$ plane for fixed $y_{t'}^2 = 0.65$ and the same value of $\tan \beta = 3$. In this case with a higher Yukawa coupling, slightly higher $c_{10}$ is needed to obtain a neutralino LSP. As in the previous case, we see a brick-red shaded gluino LSP for low $c_{10}$ and, at slightly larger $c_{10}$, a gluino coannihilation strip. As previously, $m_H$ is acceptable along all the displayed portion of the strip where $m_{3/2} \lesssim 400$ TeV. In this case, the region at larger $c_{10}$ where one or more of the new vector scalars becomes tachyonic also extends to low $c_{10}$ for small $m_{3/2}$. The right panel of figure 11 shows the values of $\Delta M$ and $m_{\chi}$ along the gluino coannihilation strip. In this case, we see that $\Delta M$ is maximized at $\sim 160$ GeV for $m_{3/2} \sim 150$ TeV. The end-point of the gluino coannihilation strip also occurs at $m_{3/2} \simeq 500$ TeV with $m_{\chi}$ around $8.3$ TeV.

Finally, we show in figure 12 an example of a $(y_{t'}^2, c_{10})$ plane with $\tan \beta = 5$ for four choices of the gravitino mass, namely $m_{3/2} = 30, 50, 100$ and $250$ TeV. The red shaded

\footnote{The maximal value of the LSP mass that is compatible with it being a viable dark matter candidate in this case is similar to the CMSSM case, even though there are additional squarks to mediate $q q' \rightarrow q g$. This is because the limiting reaction is gluino-gluino annihilation, which is the same in the two cases.}
Figure 11. As in figure 10, the PGM \((c_{10}, m_{3/2})\) plane for fixed \(\tan \beta = 3\) and \(y_t^2 = 0.65\).

The right panel shows that the gluino-neutralino mass difference is almost independent of \(y_t^2\). We do not show the neutralino mass for these cases, as it is largely independent of \(y_t^2\) and is determined from the gravitino mass and can be read from either of the two previous figures.

5 Summary and conclusions

We have presented in this paper a couple of representative MSSM scenarios in which the gluino may be nearly degenerate with the neutralino LSP \(\chi\), whose relic dark matter density is brought into the range favoured by Planck and other data by gluino coannihilation. It had been shown previously that values of \(m_\chi \lesssim 8\) TeV are in principle possible when gluino coannihilation is operative \([72]\), and we have shown in this paper how such a possibility can be embedded within a scenario for non-universal soft supersymmetry breaking within the MSSM, on the one hand, and within a simple extension of the MSSM with pure gravity mediation of soft supersymmetry breaking that includes a vector-like \(10 + \overline{10}\) multiplet pair, on the other hand.

In both scenarios, the upper bound on \(m_\chi\) depends on the details of the models. In particular, in the non-universal MSSM scenario there is competition from other mechanisms for bringing the dark matter density into the Planck range. These, together with other phenomenological constraints such as the mass of the Higgs boson and the requirement to

\[11\] We recall also that the Higgs mass is sensitive to the choice of \(\tan \beta\).
ensure electroweak symmetry breaking, restrict the parameter region where gluino coannihilation is dominant. In the PGM scenario with extra vector-like multiplets, the allowed range of $m_\chi$ depends on the gravitino mass as well as a vector-like Yukawa coupling. In both cases, values of $m_\chi \sim 8$ TeV are quite possible.

Gluino coannihilation therefore offers the possibility that the LSP, and hence the rest of the supersymmetric spectrum, may lie in the multi-TeV range, beyond the reach of the LHC. Of course, we sincerely hope, if not expect, that supersymmetry will be discovered during future LHC runs. That said, the scenarios discussed here illustrate one way in which detection at the LHC could be evaded. An interesting and important question that lies beyond the scope of this paper is how to detect supersymmetry in a gluino coannihilation scenario with a multi-TeV LSP (see, e.g., the discussion in [150]). As we have discussed in this paper, the gluino-neutralino mass difference in such a scenario is typically $\mathcal{O}(100)$ GeV, resulting in a suppressed missing-energy signature whose detection at a future 100-TeV proton-proton collider might be challenging.

Une affaire à suivre.

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Figure 12. The PGM ($y_{t^2}, c_{10}$) plane with $\tan \beta = 5$ for four choices of the gravitino mass, $m_{3/2} = 30, 50, 100$ and 250 TeV. The right panel shows the gluino-neutralino mass difference as a function of $y_{t^2}$.
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