JHEP HYPER VERSION

Received: *August 7, 2002*, Accepted: *December 10, 2002*
Revised: *November 20, 2002*

Upper and lower limits on neutralino WIMP mass and spin-independent scattering cross section, and impact of new $(g-2)_\mu$ measurement

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Abstract: We derive the allowed ranges of the spin-independent interaction cross section σ_p^{SI} for the elastic scattering of neutralinos on proton for wide ranges of parameters of the general Minimal Supersymmetric Standard Model. We investigate the effects of the lower limits on Higgs and superpartner masses from colliders, as well as the impact of constraints from $b \to s\gamma$ and the new measurement of $(g-2)_\mu$ on the upper and lower limits on σ_p^{SI}. We further explore the impact of the neutralino relic density, including coannihilation, and of theoretical assumptions about the largest allowed values of the supersymmetric parameters. For $\mu > 0$, requiring the latter to lie below 1 TeV leads to $\sigma_p^{SI} \gtrsim 10^{-11}$ pb at $m_\chi \sim 100$ GeV and $\sigma_p^{SI} \gtrsim 10^{-8}$ pb at $m_\chi \sim 1$ TeV. When the supersymmetric parameters are allowed above 1 TeV, for 440 GeV $\lesssim m_\chi \lesssim$ 1020 GeV we derive a *parameter-independent lower limit* of $\sigma_p^{SI} \gtrsim 2 \times 10^{-12}$ pb. (No similar lower limits can be set for $\mu < 0$ nor for 1020 GeV $\lesssim m_\chi \lesssim$ 2.6 TeV.) Requiring $\Omega_\chi h^2 < 0.3$ implies a *parameter-independent upper limit* $m_\chi \lesssim 2.6$ TeV. The new e^+e^--based measurement of $(g-2)_\mu$ restricts $m_\chi \lesssim 350$ GeV at $1\,\sigma$ CL and $m_\chi \lesssim 515$ GeV at $2\,\sigma$ CL, and implies $\mu > 0$. The largest allowed values of σ_p^{SI} have already become accessible to recent experimental searches.

Keywords: Supersymmetric Effective Theories, Cosmology of Theories beyond the SM, Dark Matter.

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1. Introduction

The hypothesis of the lightest neutralino χ, as the lightest supersymmetric particle (LSP), providing the dominant contribution to cold dark matter (CDM) in the Universe, has inspired much activity in the overlap of today's particle physics and cosmology. It is well-known that the relic density of the neutralinos is often comparable with the critical density [1, 2]. The expectation that the Galactic dark matter (DM) halo is mostly made of weakly-interacting massive particles (WIMPs) has further led to much experimental activity. In particular, the experiments looking for CDM WIMPs elastically scattering off underground targets have recently set limits on spin-independent (SI), or scalar, cross section of the order of 10^{-6} pb [3, 4, 5]. They have also nearly ruled out the region of (m_χ, σ_p^{SI}) that has been claimed by the DAMA experiment to be consistent with an annual modulation effect [6]. Initial and early studies [7, 8, 9, 10] were followed by more recent work [11], where it was concluded that current experimental sensitivity is generally comparable with the ranges expected from the neutralino WIMP in the Minimal Supersymmetric Standard Model (MSSM). It is, however, still at least on order of magnitude, or so, above the ranges predicted by recent analyses of the Constrained MSSM (CMSSM) [12]–[16].

For comparison, cross sections for spin-dependent (SD) interactions are in the case of the neutralino generally some two or three orders of magnitude larger than the SI ones. On the other hand, at present detectors are still not sensitive enough to explore the parameter space of the MSSM, despite recent progress [17].

In light of the ongoing and planned experimental activities, it is timely to conduct a thorough and careful re-analysis of the predicted cross sections for SI scattering of neutralino WIMPs. Such a study is rather challenging because resulting ranges often strongly depend on a given SUSY model and on related theoretical assumptions. They are further

affected by experimental limits on SUSY, both from colliders and from indirect searches, as well as by cosmological input, where the relic abundance of the CDM has been measured with better accuracy both directly and in CMBR studies [18]. Over the last few years and months there have been also new results for LEP lower bounds on the masses of the lightest Higgs and electroweakly-interacting superpartners, Tevatron lower limits on strongly-interacting superpartners, as well as limits on allowed SUSY contributions to $b \to s\gamma$, and especially to the anomalous magnetic moment of the muon $(g-2)_\mu$ [19]. The new result for $(g-2)_\mu$ indicates a sizable deviation from the Standard Model prediction, whose value is still a subject of much discussion. As we will show, when interpreted in terms of SUSY, the required extra contribution to $(g-2)_\mu$ plays a unique role in implying a stringent *upper* bound $m_\chi \lesssim 350\,\mathrm{GeV}$ ($1\,\sigma$ CL) and $m_\chi \lesssim 512\,\mathrm{GeV}$ ($2\,\sigma$), but it does *not* affect much the allowed ranges of the SI scattering cross section σ_p^{SI}.

In this paper we carefully study the impact of the above constraints. In an attempt to minimize theoretical bias, we work here in the context of the general MSSM, which will be defined below, with an additional assumption of R-parity conservation. We focus here on the SI cross section case. Other recent studies of the general MSSM include [20, 21, 22]. The results presented here show that the level of experimental sensitivity that has recently been reached [3, 4, 5] has now indeed allowed one to start exploring cosmologically favored ranges of the neutralino WIMP mass and SI cross section. However, we point out a number of caveats and relations and further discuss the origin of, and robustness of, the upper and lower limits on σ_p^{SI}. In particular, for a big range of the heavy neutralino mass $440\,\mathrm{GeV} \lesssim m_\chi \lesssim 1020\,\mathrm{GeV}$ we are able to derive *parameter-independent* lower bounds on σ_p^{SI}. Collider lower limits on Higgs and superpartner masses plus requiring $\Omega_\chi h^2 < 0.3$ alone leads to a *parameter-independent* upper bound $m_\chi \lesssim 2.6\,\mathrm{TeV}$. This is true except near Higgs resonances for which we can have even larger m_χ consistent with $\Omega_\chi h^2$.

2. The MSSM

We start by reviewing relevant features of the general MSSM [23]. (We follow the convention of [24].) By the MSSM we mean a supersymmetrized version of the Standard Model, with Yukawa and soft SUSY-breaking terms consistent with R-parity. We neglect CP-violating phases in the Higgs and SUSY sectors and assume no mixings among different generations of squarks and sleptons, since both are probably small. In the MSSM, all the slepton and squark masses can be considered as a priori free parameters set at the electroweak scale. In the same way we treat the trilinear parameters A_i $(i = t, b, \tau)$ of the third generation while neglecting the ones of the first two. The Higgs sector is determined at the tree level by the usual ratio of the neutral Higgs VEV's $\tan\beta = v_t/v_b$ and the mass of the pseudoscalar m_A. In computing full one-loop and leading two-loop radiative corrections to the lightest scalar Higgs we use the package FeynHiggsFast (FHF) [25]. As we will see, Higgs masses will play an important role in the analysis.

The lightest neutralino χ is lightest of the four mass eigenstates of the linear combinations of the bino $\widetilde{B}$, the wino $\widetilde{W}_3^0$ and the two higgsinos $\widetilde{H}_b^0$ and $\widetilde{H}_t^0$

$$\chi \equiv \chi_1^0 = N_{11}\widetilde{B} + N_{12}\widetilde{W}_3^0 + N_{13}\widetilde{H}_b^0 + N_{14}\widetilde{H}_t^0\,. \tag{2.1}$$

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The neutralino mass matrix $\mathcal{M}$ [24] is determined by the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ gaugino mass parameters M_1 and M_2, respectively (and we impose the relation $M_1 = \frac{5}{3}\tan^2\theta_W M_2$, which comes from assuming gaugino mass unification at GUT scale), the Higgs/higgsino mass parameter μ, as well as $\tan\beta$. In the region $|\mu| \gg M_1$, the lightest neutralino is mostly a bino with mass $m_\chi \simeq M_1$. In the other extreme, it is mostly higgsino-dominated and $m_\chi \simeq |\mu|$.

For the purpose of this analysis, we take as independent parameters: $\tan\beta$, μ, M_2, m_A, $A_{t,b}$ as well as the soft masses of the sleptons and of the squarks. In order to make our analysis manageable, we make an additional assumption that, at the electroweak scale, the soft mass parameters of the sleptons are all equal to some common value $\widetilde{m}_{\tilde{l}}$, and analogously $\widetilde{m}_{\tilde{q}}$ for all the squarks. One normally expects certain relations among the physical masses of the sleptons and the squarks since they in addition receive well-defined D-term and F-term contributions to their mass matrices. Assuming common soft mass terms, at either GUT or electroweak scale, this normally leads to the sleptons being lighter than the squarks. Furthermore, for the sfermions of the 3rd generation it is natural to expect large mass splittings. However, we believe that, for our purpose, introducing just two separate common soft mass scales $\widetilde{m}_{\tilde{l}}$ and $\widetilde{m}_{\tilde{q}}$, while greatly simplifying the analysis, will not play much role in our overall conclusions for the SI cross sections. (This is in contrast to often assumed *full* degeneracy of soft sfermion masses which leads in our opinion to an unnecessary limitation on the allowed SUSY parameter space.)

What we do find important is to disentangle squark and slepton masses. This is mainly because experimental limits on slepton masses are significantly weaker than in the case of squarks. For the case of the bino-like neutralino, which is the most natural case for providing $\Omega_\chi h^2 \sim 1$ [26], it is therefore the mass of the lightest slepton which often predominantly determines $\Omega_\chi h^2$. Assuming common soft masses for sleptons and squarks at the electroweak scale would therefore generally lead to overestimating $\Omega_\chi h^2$ by missing the cases where relatively light sleptons (below current squark mass bounds) would otherwise provide acceptable $\Omega_\chi h^2$. An additional effect is that of coannihilation. When slepton mass is only somewhat larger than that of the LSP, the neutralino relic abundance is strongly reduced and otherwise forbidden cases become allowed [27].

3. The spin-independent cross section

For non-relativistic Majorana particles, like the neutralino WIMP, the elastic scattering off constituent quarks and gluons of some nucleon ${}^A_Z X$ is given by an effective differential cross section [7, 8, 10]

$$\frac{d\sigma}{d|\vec{q}|^2} = \frac{d\sigma^{SI}}{d|\vec{q}|^2} + \frac{d\sigma^{SD}}{d|\vec{q}|^2}\,, \tag{3.1}$$

where the transferred momentum $\vec{q} = \mu_A \vec{v}$ depends on the velocity $\vec{v}$ of the incident WIMP, and $\mu_A = m_A m_\chi/(m_A + m_\chi)$ is the reduced mass of the system. The effective WIMP-nucleon cross sections σ^{SI} and σ^{SD} are computed by evaluating nucleonic matrix elements of corresponding WIMP-quark and WIMP-gluon interaction operators.

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In the SI part, contributions from individual nucleons in the nucleus add coherently and the finite size effects are accounted for by including the SI nuclear form factor $F(q)$. The differential cross section for the scalar part then takes the form [2]

$$\frac{d\sigma^{SI}}{d|\vec{q}|^2} = \frac{1}{\pi v^2}\left[Zf_p + (A-Z)f_n\right]^2 F^2(q)\,, \tag{3.2}$$

where f_p and f_n are the effective neutralino couplings to protons and neutrons, respectively. Explicit expressions for the case of the supersymmetric neutralino can be found, e.g., in [28]. The formalism we follow has been reviewed in several recent papers [2, 28, 12]. We have re-done the original complete calculation of Drees and Nojiri [10] and agreed with their results.

A convenient quantity which is customarily used in comparing theory and experimental results for SI interactions is the cross section σ_p^{SI} for WIMP elastic scattering of free proton in the limit of zero momentum transfer:

$$\sigma_p^{SI} = \frac{4}{\pi}\mu_p^2 f_p^2 \tag{3.3}$$

where μ_p is defined similarly to μ_A above. The analogous quantity for a target with nuclei with mass number A can then be expressed in terms of σ_p^{SI} as

$$\sigma_A^{SI} = \frac{4}{\pi}\mu_A^2\left[Zf_p + (A-Z)f_n\right]^2 = \left(\frac{\mu_A}{\mu_p}\right)^2 A^2\sigma_p^{SI}\,. \tag{3.4}$$

One can do so because, for Majorana WIMPs, $f_p \simeq f_n$.

The coefficients $f_{p,n}$ can be expressed as [10]

$$\frac{f_p}{m_p} = \sum_{q=u,d,s}\frac{f_{Tq}^{(p)}}{m_q}f_q + \frac{2}{27}f_{TG}^{(p)}\sum_{c,b,t}\frac{f_q}{m_q} + \cdots$$

where $f_{TG}^{(p)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p)}$, and $f_{Tq}^{(p)}$ is given by $\langle p|m_q\bar{q}q|p\rangle = m_p f_{Tq}^{(p)}$ $(q = u, d, s)$, and analogously for the neutron. The masses and ratios $B_q = \langle p|\bar{q}q|p\rangle$ of light constituent quarks in a nucleon come with some uncertainties. For definiteness, we follow a recent re-evaluation [12] and assume $m_u/m_d = 0.553 \pm 0.043$, $m_s/m_d = 18.9 \pm 0.8$, and $B_d/B_u = 0.73 \pm 0.02$, as well as

$$\begin{aligned}
f_{Tu}^{(p)} &= 0.020 \pm 0.004\,, & f_{Td}^{(p)} &= 0.026 \pm 0.005\,, & f_{Ts}^{(p)} &= 0.118 \pm 0.062 \\
f_{Tu}^{(n)} &= 0.014 \pm 0.003\,, & f_{Td}^{(n)} &= 0.036 \pm 0.008\,, & f_{Ts}^{(n)} &= 0.118 \pm 0.062\,,
\end{aligned} \tag{3.5}$$

which numerically gives SI cross section values very similar to using the set of [2]. Some other recent studies use a new determination of $\sigma_{\pi N}$ to derive a much larger value for $f_{Ts}^{(p)} \simeq 0.37$ [29]. We find such values somewhat questionable since they imply that the strange quark component of the nucleon would be larger than the up and down ones. We have numerically checked that using the set of input parameters of [29] gives typically SI cross section values a factor of six higher than in our case.

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It is worth noting that, despite several different diagrams and complicated expressions, it is the exchange of the heavy scalar Higgs that in most cases comes out to be numerically dominant. It is further enhanced when the neutralino is a mixed gaugino-higgsino state [30]. As $m_H \simeq m_A$ increases, σ_p^{SI} drops as m_H^{-4} because of the t-channel propagator effect in $\chi q \to \chi q$ elastic scattering. Eventually, at smaller σ_p^{SI}, squark exchange becomes important, and even dominant, instead.

4. Details of the scan

As outlined above, we use seven parameters $\tan\beta$, M_2, μ, m_A, $\widetilde{m}_{\tilde{q}}$, $\widetilde{m}_{\tilde{l}}$ and $A_t = A_b$ to conduct a careful scan of the general MSSM parameter space. For their allowed ranges we take:

$$\begin{aligned}
50\,\text{GeV} \leq\ & M_2 & \leq 2\,\text{TeV} \\
50\,\text{GeV} \leq\ & |\mu| & \leq 2\,\text{TeV}\ (4\,\text{TeV}) \\
50\,\text{GeV} \leq\ & \widetilde{m}_{\tilde{l}} & \leq 2\,\text{TeV}\ (4\,\text{TeV}) \\
200\,\text{GeV} \leq\ & \widetilde{m}_{\tilde{q}} & \leq 2\,\text{TeV}\ (4\,\text{TeV}) \\
90\,\text{GeV} \leq\ & m_A & \leq 2\,\text{TeV} \\
0 \leq\ & |A_{t,b}| & \leq 1\,\text{TeV} \\
5 \leq\ & \tan\beta & \leq 65
\end{aligned} \tag{4.1}$$

while we set $A_\tau = 0$ since we treat the masses of the slepton as independent parameters anyway. In addition to a general scan of the parameter space, in many cases we do several focused scans and explore the effect of extremely large values of μ, $\widetilde{m}_{\tilde{l}}$ and $\widetilde{m}_{\tilde{q}}$ beyond 2 TeV (given in brackets above) in order to derive *parameter-independent* lower limits on σ_p^{SI} for a big range of large m_χ, as described below. The minimum values are set so that the resulting physical masses of Higgs and superpartners are limited from below by collider bounds. For the lighter chargino we take $m_{\chi_1^\pm} > 104\,\text{GeV}$ [31], for sleptons the lower limit of 90 GeV [31] and for squarks 200 GeV [32]. The lower limit on m_χ depends not only on a model but also on a number of additional assumptions [33]. For this reason, in our analysis we conservatively do not impose a direct experimental limit on m_χ, but instead infer it from the other limits, especially the one on the chargino mass. As regards the lightest Higgs mass m_h, in much of the parameter space ($m_A > 120\,\text{GeV}$) the lower limit on the Standard Model Higgs of 114.1 GeV [34] applies. However there are two important points to note. Firstly, theoretical uncertainties in computing m_h in the MSSM are estimated at $2-3\,\text{GeV}$. Conservatively, we thus require only $m_h > 111\,\text{GeV}$. Secondly, for $90\,\text{GeV} < m_A < 120\,\text{GeV}$, there still remains a sizable range of the (m_h, m_A)-plane where the lightest Higgs mass given roughly by $m_h > 0.78(m_A + 21.7\,\text{GeV})$ is allowed [34]. We will comment below on the effect of this low-mass range on increasing the largest allowed values of σ_p^{SI}.

Among indirect limits on SUSY, $b \to s\gamma$ often places an important additional constraint on the allowed parameter space. We calculate the SM contribution to $\text{BR}(B \to X_s\gamma)$ at

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the full NLO level and include dominant $\tan\beta$-enhanced NLO SUSY [35], and also include the c-quark mass effect on the SM value [37]. At this level of accuracy the SM prediction is $\mathrm{BR}(B \to X_s\gamma) = (3.70 \pm 0.30) \times 10^{-4}$ [37]. This range is partially overlapping with the new world-average [39]

$$\mathrm{BR}(B \to X_s\gamma) = (3.41 \pm 0.36) \times 10^{-4} \tag{4.2}$$

which has gone up from the previous range of $(3.23 \pm 0.72) \times 10^{-4}$ following the new result from BaBAR [36]. As described in more detail in [14], with an update in [38], we accordingly allow the full SM+SUSY contribution to be in the range $\mathrm{BR}(B \to X_s\gamma) = (3.41 \pm 0.67) \times 10^{-4}$. It is important, however, to stress here an important salient point. In computing the SUSY contribution to $\mathrm{BR}(B \to X_s\gamma)$ one usually makes an implicit assumption of minimal flavor violation in the down-type squark sector, which is theoretically poorly justified. Even a slight modification of the assumption often leads to a significant relaxation of the bound from $b \to s\gamma$ to the point of even allowing $\mu < 0$ and relatively light superpartner masses [39].

A very recent measurement of the anomalous magnetic moment of the muon $a_\mu = (g_\mu - 2)/2$ [19] has confirmed a previous value [40] but with twice-increased precision [19]

$$a_\mu^{\mathrm{expt}} - 11659000 \times 10^{-10} = (203 \pm 8) \times 10^{-10}\,. \tag{4.3}$$

The LO hadronic vacuum polarization contribution has recently been re-evaluated in [41, 42] and the light-by-light corrections in [43]. In [41] the updated SM prediction of [44] has been found to be $a_\mu^{\mathrm{SM}} - 11659000 \times 10^{-10} = (169.1 \pm 7.8) \times 10^{-10}$ when applying data from e^+e^- annihilation cross sections and $(186.3 \pm 7.1) \times 10^{-10}$ when applying τ-decay data. This leads to a $3\,\sigma$ discrepancy

$$\Delta a_\mu = a_\mu^{\mathrm{expt}} - a_\mu^{\mathrm{SM}} = (33.9 \pm 11.2) \times 10^{-10} \tag{4.4}$$

when using the e^+e^--based data, or to a $1.6\,\sigma$ deviation

$$\Delta a_\mu = a_\mu^{\mathrm{expt}} - a_\mu^{\mathrm{SM}} = (16.7 \pm 10.7) \times 10^{-10} \tag{4.5}$$

when applying the τ-based data.

If interpreted in terms of SUSY, eq. (4.4) restricts the allowed SUSY contribution to

$$22.7 \times 10^{-10} < a_\mu^{\mathrm{SUSY}} < 45.1 \times 10^{-10} \qquad (1\,\sigma) \tag{4.6}$$

$$11.5 \times 10^{-10} < a_\mu^{\mathrm{SUSY}} < 56.3 \times 10^{-10} \qquad (2\,\sigma)\,. \tag{4.7}$$

Similar ranges are obtained by using the results of [42, 45]. The eqs. (4.4) and (4.5) further imply that $\mu > 0$ (the sign of the SUSY contribution is the same as that of μ).

This is clearly an intriguing hint for "new physics". However, since some other recent evaluations tend to give a larger value of a_μ^{SM} [46] and a larger error bar [47], at this point we will not strictly impose the $(g-2)_\mu$ constraint on the parameter space that is otherwise allowed by all other constraints. Nevertheless, below we will discuss the important impact it has on the upper bounds on m_χ.

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As regards the WIMP relic abundance, a lower limit on the age of the Universe conservatively gives $\Omega_\chi h^2 < 0.3$, while "direct" measurements of the CDM lead to $0.1 < \Omega_\chi h^2 < 0.2$ which we will treat as a preferred range. Recent studies of the CMBR seem to imply even more restrictive ranges; for example, $\Omega_\chi h^2 < 0.12 \pm 0.04$ in [48]. We will not apply this narrower range yet but will comment below on its impact on the upper and/or lower limits on σ_p^{SI}. We compute $\Omega_\chi h^2$ as accurately as one reasonably can, at the level of a few per cent both near and further away from poles and thresholds, by applying our recently derived exact analytic expressions for neutralino pair-annihilation [49] and neutralino-slepton coannihilation [50], and by using an exact procedure for the neutralino coannihilation with chargino and next-to-lightest neutralino [51, 52].

5. Results

The allowed ranges of the SI cross section that result from our scans are illustrated in figure 1 for $\tan\beta = 35$, $A_t = A_b = 1\,\mathrm{TeV}$ and $\mu > 0$. Since σ_p^{SI} generally grows with $\tan\beta$ due to an enhancement in the heavy scalar Higgs coupling to down-type quarks, the above choice is a reasonable compromise between the low and the very large values of $\tan\beta$. In deriving the allowed ranges of σ_p^{SI} we have imposed the bounds from colliders and from $b \to s\gamma$, and have further required $0.1 < \Omega_\chi h^2 < 0.2$, as described earlier, but *not* yet the bound from $(g-2)_\mu$. We also mark with a thick solid line a *parameter-independent* lower bound on σ_p^{SI} which will be explained below.

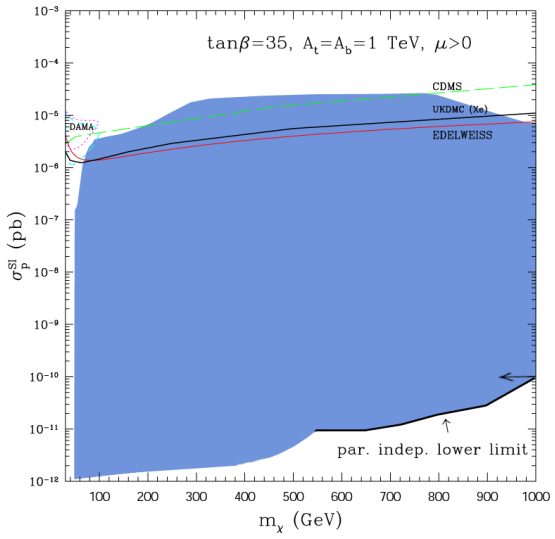

Figure 1: Ranges of σ_p^{SI} in the general MSSM vs. m_χ for $\tan\beta = 35$, $A_t = A_b = 1\,\mathrm{TeV}$ and $\mu > 0$, which are allowed by the bounds from colliders, $b \to s\gamma$ and $0.1 < \Omega_\chi h^2 < 0.2$, but not from $(g-2)_\mu$. Also marked are some results of recent experimental WIMP searches. The thick black line (and a left-pointing arrow) indicates a *parameter-independent* lower bound on σ_p^{SI} for $550\,\mathrm{GeV} < m_\chi < 1020\,\mathrm{GeV}$. No similar bound can be set for lower m_χ because of the neutralino-slepton coannihilation effect, as explained in the text.

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We can see a big spread of σ_p^{SI} over some seven (four) orders of magnitude at small (large) m_χ. In fact, the upper values of σ_p^{SI} exceed the latest experimental limits, including the recent result from Edelweiss [4] and the new limit from the UKDMC Zeplin I detector [5]. Also shown is the CDMS bound and the so-called DAMA region. It is clear that today's experiments have already started probing the most favored ranges of σ_p^{SI} that come from SUSY predictions for neutralino cold dark matter.

In the four windows of figure 2 we show the effect of the most important constraints on the upper and lower limits on the allowed ranges of σ_p^{SI}. Firstly, in the upper left window we show the effect of relaxing the cosmological bound by allowing $\Omega_\chi h^2 < 0.1$. Obviously, larger ranges of σ_p^{SI} now become allowed since an enhancement in the neutralino pair-annihilation cross section often, by crossing symmetry, implies an increase in σ_p^{SI}. Note, however, that a combination of all the other constraints, most notably a lower limit on m_h from LEP and the constraint from $b \to s\gamma$, prevents σ_p^{SI} from rising by more than about one order of magnitude and only for not very large values of m_χ. On the other hand, imposing a narrower range $0.08 < \Omega_\chi h^2 < 0.16$ has almost no effect on the upper and lower limits on σ_p^{SI}, although it does remove a number of points from the allowed ranges of σ_p^{SI}.

In the same window we also show an important effect, already pointed out in [21], of including neutralino coannihilation with sleptons (predominantly with the lighter stau) on allowing very low ranges of σ_p^{SI} at smaller m_χ. At lowest $m_\chi \lesssim 120\,\mathrm{GeV}$ the relic abundance can be reduced to the favored range by choosing in the scan light enough sleptons, even without coannihilation. By simultaneously choosing large enough heavy Higgs and squark masses, one can reduce σ_p^{SI} to very low values of a few $\times 10^{-12}\,\mathrm{pb}$. As m_χ increases, $\Omega_\chi h^2$ would normally increase as well, and become too large, but it is there that neutralino-slepton coannihilation kicks in. Since σ_p^{SI} is independent of the slepton masses, by carefully scanning the parameter space, one can always find $\widetilde{m}_{\widetilde{l}}$ not much above m_χ, in which case $\Omega_\chi h^2$ can be sufficiently reduced again to fall into the favored range. The effect is very strong for smaller m_χ, thus explaining a sharp rise of the left side of the dark-red region allowed by neutralino-slepton coannihilation, but, as the process becomes increasingly inefficient at larger m_χ [27, 50], it gradually fades away.

In the upper right window of figure 2 we present the effect of imposing the constraint from $(g-2)_\mu$. We can see that, for this case, $m_\chi \lesssim 245\,\mathrm{GeV}$ ($1\,\sigma$ CL) and $m_\chi \lesssim 420\,\mathrm{GeV}$ ($2\,\sigma$ CL). This upper limit comes from the fact that, as m_χ increases, the SUSY contribution from the $\chi - \widetilde{\mu}$ and $\chi^- - \widetilde{\nu}_\mu$ loops become suppressed and at some point becomes too small to explain the apparent discrepancy between the SM and the experimental measurement [53, 54]. On the other hand, the upper and lower limits on σ_p^{SI} are not really affected.

In the same window we also show the effect of relaxing the constraints from $b \to s\gamma$. We can see that if it were not imposed, at large m_χ the upper limit on σ_p^{SI} would significantly increase. In this region, the mass of the pseudoscalar Higgs, and therefore also the heavier scalar, is rather small, thus giving larger σ_p^{SI}. However, the mass of the charged Higgs is then also on the lower side, and a cancellation between a charged Higgs-top quark loop and chargino-stop loop contribution is not sufficient to reduce $\mathrm{BR}(B \to X_s\gamma)$ to agree with the experimental limit.

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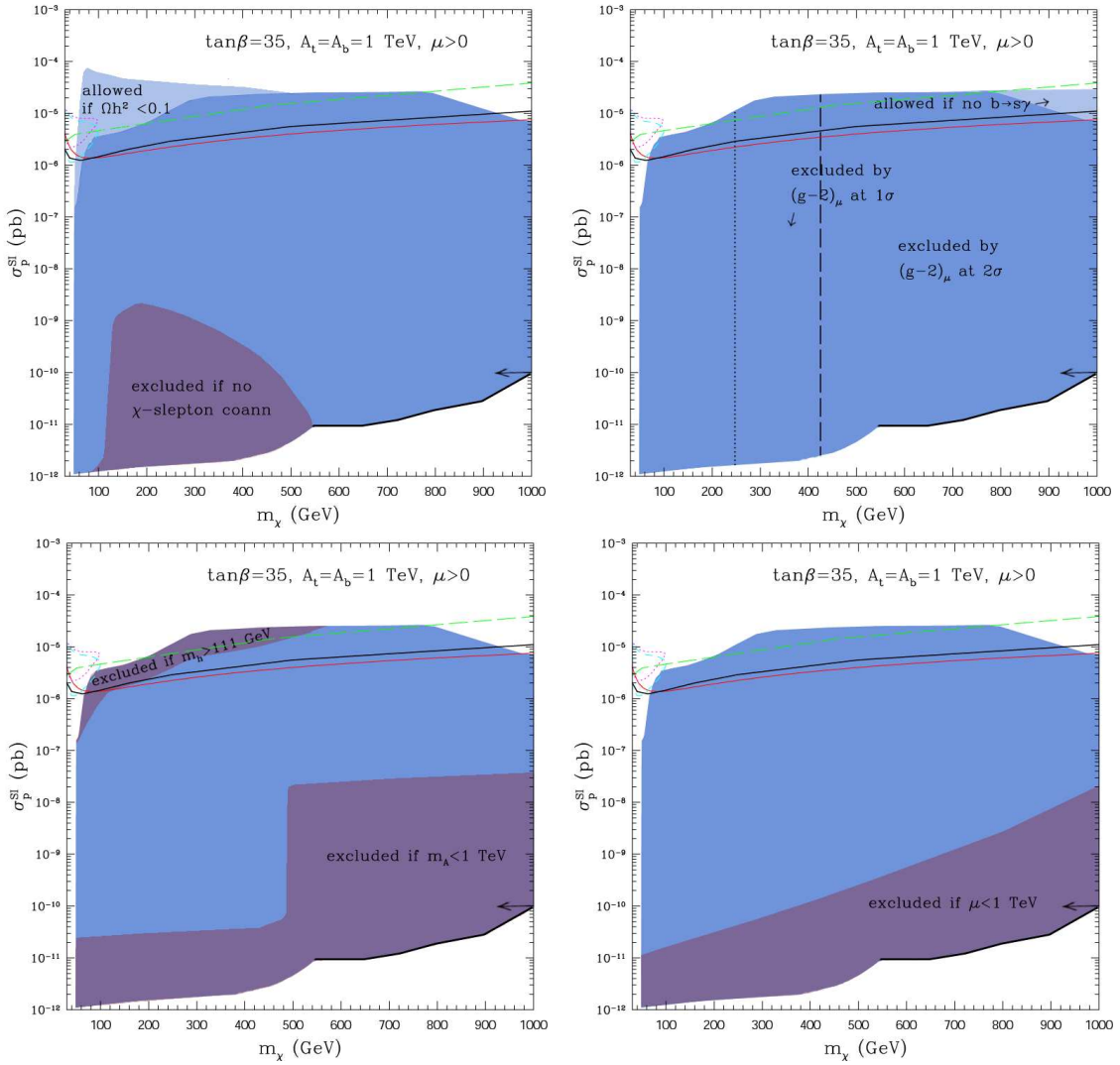

Figure 2: Sensitivity of the upper and lower limits on σ_p^{SI} in figure 1 to various assumptions and constraints. Upper left window: the light blue region would be allowed if $\Omega_\chi h^2 < 0.1$. The dark-red region would be excluded if one neglected the effect of neutralino-slepton coannihilation. Upper right window: the light blue region would be allowed if one lifted the constraint from $b \to s\gamma$. The regions to the right of the vertical dotted (dashed) lines are excluded by imposing current $1\,\sigma$ ($2\,\sigma$) CL bound from $(g-2)_\mu$. Lower left window: the upper dark-red region would be excluded by assuming $m_h > 111\,\text{GeV}$ for all m_A (i.e., by neglecting a window of lighter m_h which is still allowed for $m_A < 120\,\text{GeV}$). Also shown in this window is the effect of restricting $m_A < 1\,\text{TeV}$. Lower right window: the same as for m_A but for μ.

However, we remind the reader that a slight relaxation of the underlying assumption of minimal flavor violation in the squark sector often leads to a significant weakening of the bound from $b \to s\gamma$ to the point of even allowing $\mu < 0$ [39]. We would therefore be cautious in applying the $b \to s\gamma$ constraint rigidly.

In the lower left window of figure 2 we present the sensitivity of the *upper* limit on σ_p^{SI} to the lower limit on the light Higgs mass. We can see that it would sizably decrease if we neglected the region of small $90\,\text{GeV} \lesssim m_h \lesssim 111\,\text{GeV}$ which is still allowed at $m_A \lesssim 120\,\text{GeV}$,

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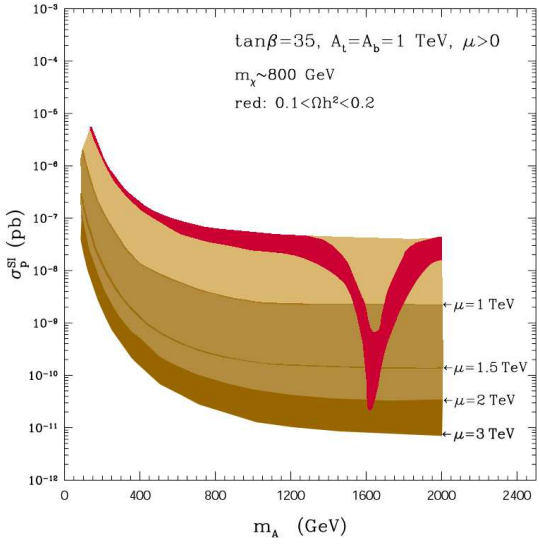

Figure 3: Sensitivity of σ_p^{SI} to m_A and μ for the case of figure 1. We concentrate on the region of parameter space where $m_\chi \sim 800\,\mathrm{GeV}$. The whole marked region is consistent with all the constraints from colliders and $b \to s\gamma$. By further imposing the constraint $0.1 < \Omega_\chi h^2 < 0.2$ one selects only the red region. Near the resonance $m_A \simeq 2m_\chi$, significantly smaller values of σ_p^{SI} become allowed by $0.1 < \Omega_\chi h^2 < 0.2$, (mostly in the increasingly pure bino region), but become eventually limited from below, independently of increasing the maximum allowed value of μ.

and instead required $m_h > 111\,\mathrm{GeV}$ for all m_A. Note also that the new experimental limits on σ_p^{SI} are for the most part inconsistent with the possibility of the light Higgs scalar.

In the same window and in the lower right window we explore the existence of the *lower* limit on σ_p^{SI} and its dependence on the assumed *upper* limit on m_A and μ, respectively. As we can see, the lowest values of σ_p^{SI} are often to a large extent determined by a somewhat subjective restrictions from above on these parameters. As one allows either μ or m_A above 1 TeV the lower limit on σ_p^{SI} relaxes considerably.

However, we argue that, by requiring sizable enough $\Omega_\chi h^2$ (e.g., $\Omega_\chi h^2 > 0.1$), it is possible to set a *parameter-independent* lower bound on σ_p^{SI} for a considerable range of large $m_\chi \lesssim 1020\,\mathrm{GeV}$ (marked with a thick solid line and left-pointing arrow in figure 2). The limit holds for an arbitrary case of the neutralino (gaugino, higgsino or mixed), independently of how large μ and other SUSY parameters are taken. Let us first consider the gaugino limit $\mu \gg M_2$. In this case the origin of the parameter-independent bound is displayed in figure 3 where we plot σ_p^{SI} as a function of the pseudoscalar Higgs m_A for $m_\chi \simeq 800\,\mathrm{GeV}$ and scan over all the other parameters. As m_A and other parameters are varied, large ranges of σ_p^{SI} remain allowed by collider and indirect constraints, depending on the maximum allowed value of μ. For each fixed μ, σ_p^{SI} decreases proportionally to the fourth power of $m_H \simeq m_A$ because of the (typically dominant) t-channel exchange of the heavy Higgs, as the marked cases of μ in figure 3 clearly demonstrate. As μ increases, at fixed m_χ one moves deeper into the gaugino region and typically finds large $\Omega_\chi h^2 > 0.2$. By imposing $0.1 < \Omega_\chi h^2 < 0.2$ one selects only a narrow red (dark) range with a large bino component. The bino purity ($p_{\widetilde{B}} = N_{11}^2$) increases with increasing μ, which normally quickly gives too large $\Omega_\chi h^2$. However, for each m_χ one can choose $m_A \simeq 2m_\chi$

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(roughly $m_A = 1600\,\text{GeV}$ in figure 3) in which case $\Omega_\chi h^2$ is reduced to an allowed level by a wide resonance due to A-exchange. This leads to allowing a much reduced σ_p^{SI}, while still being consistent with $0.1 < \Omega_\chi h^2 < 0.2$. However, because of the finite width of the A-resonance, at large enough μ one reaches the *lowest* value of σ_p^{SI} (in this case $2 \times 10^{-11}\,\text{pb}$) which is *parameter-independent*. Away from the resonance (for example, if one imposed $m_A < 1\,\text{TeV}$), one would obtain the lower bound $\sigma_p^{SI} \gtrsim 3 \times 10^{-8}\,\text{pb}$, basically independently of whether $\mu < 1\,\text{TeV}$ is imposed or not. In fact, one can see this effect in the lower left window of figure 2 where imposing $m_A < 1\,\text{TeV}$ causes the lower limit on σ_p^{SI} to suddenly jump up at around $m_\chi \simeq m_A/2 \simeq 500\,\text{GeV}$.

The above parameter-independent lower limit arises in the gaugino case and applies to arbitrarily large m_χ. The case of the gaugino-like LSP can be argued to be more attractive as being less fine-tuned than the higgsino-like one [26]. In the MSSM the LSP is a nearly-pure higgsino for $M_2 \gtrsim 300\,\text{GeV}$ which, by applying the assumption of gaugino mass-unification (see below (2.1)) to the gluino mass, implies $m_{\tilde{g}} = \frac{\alpha_s}{\alpha_2} M_2 \gtrsim 1\,\text{TeV}$ and therefore large and less "natural" soft SUSY-breaking scale. Nevertheless, in the spirit of generality, we need to extend the analysis to the case of the higgsino-like neutralino.

For the higgsino-like LSP ($M_2 \gg \mu \simeq m_\chi$), if we remain within the ranges of parameters given in (4.1), we find σ_p^{SI} some two orders of magnitude larger than in the gaugino case presented above. However, one can in principle reduce σ_p^{SI} to arbitrarily small values by going to the limit of pure enough higgsino (M_2 in the multi- TeV range), in which case the neutralino-Higgs coupling would be arbitrarily reduced, and by further suppressing the squark contribution by making them extremely heavy. It is therefore reasonable to question the existence of the lower bound on σ_p^{SI}. However, in the multi-hundred GeV range ($m_t < m_\chi \lesssim 1020\,\text{GeV}$) $\Omega_\chi h^2$ remains typically very small $\Omega_\chi h^2 \ll 0.1$ although it does increase with m_χ. For $\tan\beta = 35$ it reaches 0.1 for $m_\chi \lesssim 1020\,\text{GeV}$, 0.2 for $m_\chi \lesssim 1.6\,\text{TeV}$ and 0.3 for $m_\chi \lesssim 2.5\,\text{TeV}$. (These values decrease somewhat with increasing $\tan\beta$.) Since, as mentioned above, we impose $0.1 < \Omega_\chi h^2$ (or a similar sizable lower limit on $\Omega_\chi h^2$), in the range of $m_\chi < 1\,\text{TeV}$ displayed in figures 1 and 2 the parameter-independent lower limit on σ_p^{SI} holds. (In order to be clear that the parameter-independent lower bound on σ_p^{SI} applies to a general neutralino so long as $m_\chi \lesssim 1020\,\text{GeV}$, in figures 1 and 2 we have put a left-pointing arrow at $m_\chi = 1\,\text{TeV}$.)

For the higgsino-like LSP in the mass range $1020\,\text{GeV} \lesssim m_\chi \lesssim 1600\,\text{GeV}$ the relic abundance is $0.1 < \Omega_\chi h^2 < 0.2$, basically indpendently of how large M_2 is, and no lower bound on σ_p^{SI} can in principle be set. The range of m_χ increases to 2.5 TeV if we allow $\Omega_\chi h^2 < 0.3$. Larger values of m_χ for an arbitrary neutralino gaugino/higgsino composition are inconsistent with $\Omega_\chi h^2 < 0.3$. This upper limit on m_χ relaxes to 2.6 TeV for $\tan\beta = 10$.This is true except near Higgs resonances (m_H, $m_A \sim 2m_\chi$) for which we can have even larger m_χ consistent with $\Omega_\chi h^2$.

On the other side, at low enough m_χ coannihilation with sleptons prevents one from deriving a firm lower limit on σ_p^{SI}. Indeed, by suitably choosing the slepton mass not too much above m_χ, we can always reduce $\Omega_\chi h^2$ below 0.2. Thus, it is possible to set firm lower limits on σ_p^{SI} but only for large enough m_χ and even though they correspond to extremely large values of μ and accordingly involve much fine-tunning.

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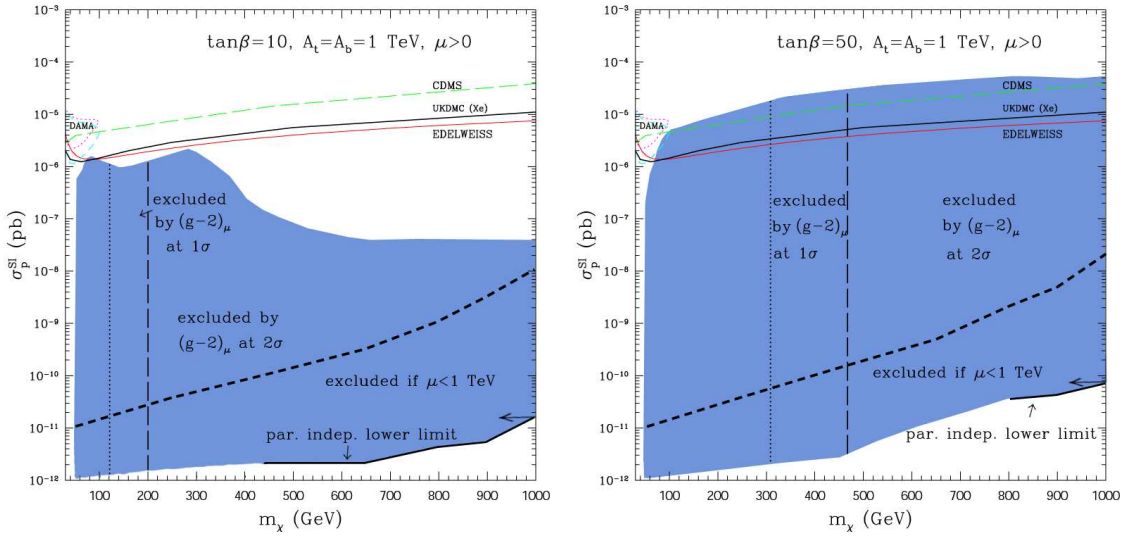

Figure 4: The same as in figure 1 but for $\tan\beta = 10$ (left window) and 50 (right window). Also marked is the effect of imposing $\mu < 1$ TeV.

The above discussion of the parameter-independent lower limit on σ_p^{SI} and upper limit on m_χ has been presented in the case of $\mu > 0$ but it obviously applies also to the case $\mu < 0$.

The dependence of σ_p^{SI} on A_t and A_b is rather weak. As the tri-linear terms deviate from zero, the mass of the lightest Higgs generally increases due to somewhat larger mass splittings among the stops and sbottoms. As a result, the normally subdominant contribution to σ_p^{SI} from the t-channel h-exchange is slightly reduced.

A far more important effect is that the number of SUSY configurations satisfying all experimental constraints, especially that from $b \to s\gamma$, decreases significantly. This is because the cancellation between charged Higgs loop and chargino loop contribution to the $b \to s\gamma$ decay rate becomes more inefficient as A_t decreases. For example, for $A_t = -1$ TeV only a handful of points remain allowed. Generally, we have concluded that the regions allowed by lower values of A_t and A_b fall into the regions allowed by the choice $A_t = 1$ TeV.

In order to display the dependence on $\tan\beta$, in the left and right window of figure 4 we present the cases of $\tan\beta = 10$ and 50, respectively. Note that, for small $\tan\beta = 10$ the largest allowed values of σ_p^{SI} are roughly an order of magnitude smaller than at $\tan\beta = 50$ because of the $\tan\beta$-dependence of the heavy scalar coupling to down-type quarks, as mentioned earlier. Notice a significant decrease in the upper ranges of σ_p^{SI} at large m_χ for $\tan\beta = 10$, which is caused by exceeding the upper limit (4.08×10^{-4}) of the allowed range of $\mathrm{BR}(B \to X_s\gamma)$. It is clear that the constraint is more severe in the case $\tan\beta = 10$ rather than at larger $\tan\beta$. This may sound somewhat counter-intuitive since, for example, in the Constrained MSSM, the constraint from $b \to s\gamma$ on the (CMSSM) parameter space becomes more pronounced at larger $\tan\beta$. This is because, in the CMSSM the pseudoscalar Higgs mass, hence also the charged Higgs mass, is typically large, and the corresponding charged Higgs-top quark loop contribution to $\mathrm{BR}(B \to X_s\gamma)$ becomes suppressed. At smaller $m_{1/2}$ (thus also m_χ and $m_{\chi_1^\pm}$) and m_0 (thus also light enough stop) the $\tan\beta$-dependent

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(negative) contribution from the chargino-stop loop gives too small BR($B \to X_s\gamma$), below the lower experimental limit, thus producing a strong lower bound on m_χ at not too large m_0. In contrast, in the general MSSM case, we can choose small values of m_A which is a free parameter. This small m_A implies a big positive charged Higgs loop contribution to BR($B \to X_s\gamma$). For smaller $\tan\beta$ and at smaller m_χ this is reduced to an acceptable range by the chargino-stop loop contribution. However, at large m_χ, the chargino-stop loop cannot cancel the charged Higgs contribution anymore and one exceeds the upper experimental limit on BR($B \to X_s\gamma$). Since the chargino loop contribution is proportional to $\tan\beta$, at large $\tan\beta$ the cancellation can be achieved even at large m_χ. For this reason, in the right window of figure 4 there is no analogous decrease in the largest allowed σ_p^{SI} at large m_χ, in contrast to the left window. Clearly, the constraint from $b \to s\gamma$ is more severe for smaller $\tan\beta$.

A similar effect of can be observed in the case of the $(g-2)_\mu$ constraint. By comparing figure 4 with the upper right window of figure 2 one can see that it produces as stronger upper bound on m_χ at smaller $\tan\beta$. For example, imposing a $1\,\sigma$ bound implies $m_\chi \lesssim 120\,\text{GeV}$ for $\tan\beta = 10$, while for $\tan\beta = 50$ the bound moves up to $m_\chi \lesssim 310\,\text{GeV}$, as denoted in the two windows of figure 4. Also marked is the effect of imposing $\mu < 1\,\text{TeV}$. While not being a firm constraint, it does, in our opinion, indicate the region which may be considered as somewhat less fine-tuned.

As discussed above, for a considerable range of large $m_\chi \lesssim 1020\,\text{GeV}$ (marked with a thick solid line and left-pointing arrow in figure 4), we can again set up the parameter-independent lower limit on σ_p^{SI}, in analogy with the case $\tan\beta = 35$ of figure 2. In order to do so, we had to explore extremely large ranges of μ up to some 4 TeV at smaller $\tan\beta$ in order to saturate the bound $\Omega_\chi h^2 < 0.2$. On the other hand, at smaller $m_\chi \lesssim 440(800)\,\text{GeV}$ for $\tan\beta = 10(50)$ the lower limit on σ_p^{SI} remains basically independent of $\tan\beta$ since $\Omega_\chi h^2$ at lower m_χ is determined mostly by the coannihilation with sleptons, as discussed in detail in the case $\tan\beta = 35$. In this case the lower limit on σ_p^{SI} arises from restricting μ below the value for which the parameter-independent lower limit arises at larger m_χ. Notice that the thick line extends to lower m_χ at smaller $\tan\beta$ because the efficiency of the neutralino-slepton coannihilation increases with $\tan\beta$.

Finally, in figure 5 we summarize the results for the full scan conducted so far for $\tan\beta = 5, 10, 35, 50, 55, 60, 65$ and for $\mu > 0$. We repeat that, in determining the allowed (blue) region we applied the constraints from collider searches, $b \to s\gamma$ and $0.1 < \Omega_\chi h^2 < 0.2$. The lower limit on σ_p^{SI} mostly comes from low $\tan\beta$ and, where applicable, we mark with the solid line where it is parameter-independent. The effect of restricting $\mu < 1\,\text{TeV}$ is marked with a dashed line. Also indicated is the impact of the new measurement of $(g-2)_\mu$ on the mass of the neutralino. If confirmed, the e^+e^--based range (4.4) will imply rather stringent upper limits on m_χ

$$m_\chi \lesssim 350\,\text{GeV} \qquad (1\,\sigma\ \text{CL}) \tag{5.1}$$

$$m_\chi \lesssim 510\,\text{GeV} \qquad (2\,\sigma\ \text{CL})\,. \tag{5.2}$$

If the τ-based numbers (4.5) are applied instead, one obtains $m_\chi \lesssim 800\,\text{GeV}$ ($1\,\sigma$ CL) and no upper bound at $2\,\sigma$ CL (for $\mu > 0$).

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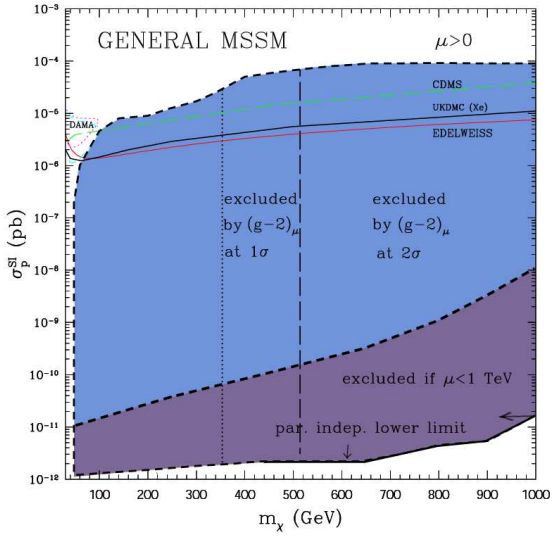

Figure 5: Ranges of σ_p^{SI} in the general MSSM vs. m_χ for $\mu > 0$, which are allowed by collider bounds, $b \to s\gamma$ and $0.1 < \Omega_\chi h^2 < 0.2$. Also marked are some results of recent experimental WIMP searches. The thick black line and a left-pointing arrow indicate a *parameter-independent* lower bound. The region below the dashed line is excluded if one imposes the constraint $\mu < 1\,\text{TeV}$. The ranges of m_χ to the vertical lines are excluded at $1\,\sigma$ and $2\,\sigma$ CL by the current discrepancy between the experimental value of $(g-2)_\mu$ and the Standard Model prediction.

The vast ranges of σ_p^{SI} predicted in the framework of the general MSSM may be somewhat discouraging to DM WIMP hunters. It is worth noting, however, that it is the region of smaller m_χ, below a few hundred GeV, that not only is implied by the new result for $(g-2)_\mu$, but is also theoretically more favored as corresponding to less fine-tuning. Furthermore, ranges of very small $10^{-12}\,\text{pb} \lesssim \sigma_p^{SI} \lesssim 10^{-8}\,\text{pb}$ generally correspond either to very large (and therefore perhaps somewhat less natural) values of μ and/or m_A, or become allowed by selecting slepton masses on the light side, and in the χ-slepton coannihilation region, within some 20 GeV of m_χ, which again can be be considered as a finely-tuned case.

At the end, we comment again on the case of the Constrained MSSM. Because the model is much more restrictive, the ranges of σ_p^{SI} that one obtains in the parameter space allowed by all constraints, are very much narrower [12]–[16]. They are also typically somewhat lower than the largest ones allowed in the general MSSM. For example, at $\tan\beta = 50$ we find $\sigma_p^{SI} \sim 10^{-7}\,\text{pb}$ at $m_\chi = 100\,\text{GeV}$ and $\sigma_p^{SI} \sim 7 \times 10^{-11}\,\text{pb}$ at the largest (neglecting $(g-2)_\mu$) allowed value of $m_\chi = 800\,\text{GeV}$. On the other hand, because the model is defined at the grand-unified, and not electroweak, scale, in the case of large $\tan\beta \gtrsim 50$ and/or large scalar masses, theoretical uncertainties involved in the running of parameters are substantial and have much impact on the resulting ranges of both m_χ and σ_p^{SI} [15]. We will explore the case of the Constrained MSSM in an oncoming publication.

6. Conclusions

We have delineated the ranges of the SI cross section σ_p^{SI} in the general MSSM, which are consistent with current experimental bounds and for which one finds the expected amount

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of dark matter. We have further discussed the dependence of our results on the experimental constraints and on the underlying theoretical assumptions. While the ranges which we have obtaine extend over more than six orders of magnitude, we find it encouraging that the experimental sensitivity that has recently been reached, now allows one to explore our theoretical predictions for the MSSM. As we have argued above, smaller values of the WIMP mass and also larger values of σ_p^{SI} may be considered as more natural, which will hopefully be confirmed by a measuring a positive WIMP detection signal in the near future.

Acknowledgments

We thank M. Brhlik for his help in testing our numerical code for computing the cross section for WIMP-proton scattering. L.R. is grateful to A. Czarnecki for helpful comments about theoretical and experimental aspects of $(g-2)_\mu$. This work was supported in part by the EU Fifth Framework network "Supersymmetry and the Early Universe" (HPRN-CT-2000-00152).

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