Cold Dark Matter Candidate in a Class of Supersymmetric Models with an extra U(1)

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Abstract

In supersymmetric models whose gauge group includes an additional U(1) factor at the TeV scale, broken by the VEV of an standard model singlet $S$, the parameter space can accommodate a very light neutralino not ruled out experimentally. This higgsino-like fermion, stable if $R$-parity is conserved, can make a good cold dark matter candidate. We examine the thermal relic density of this particle and discuss the prospects for its direct detection if it forms part of our galactic halo.

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1 Introduction

Supersymmetric models with conserved $R$-parity have in the lightest supersymmetric particle (LSP) a natural candidate for dark matter. This very appealing feature has motivated much work (see [1, 2] for review and references) and many increasingly sophisticated studies on the field. Most of these analyses have concentrated in the minimal supersymmetric standard model (MSSM) with or without theoretical constraints on its wide parameter space. It is important, however, to keep an open mind to the possibility that this simplest supersymmetric generalization of the standard model may not be the one realized in nature. One should be careful not to identify the predictions of the MSSM with those generic of low-energy supersymmetry. In this respect it is healthy to explore (well motivated) extensions of the MSSM in search of phenomenological (or cosmological) consequences that are different from those expected in the minimal model.

Perhaps among the best motivated extended models are those that include an additional $U(1)$ factor in the gauge group, broken radiatively at some scale below $\mathcal{O}(1)$ TeV by the VEV of a standard model singlet $S$. This type of models can be generically expected to arise as the low-energy limit of some string models [3] and have a number of interesting consequences both at the phenomenological (e.g. for $Z'$ and Higgs physics) and theoretical level (e.g. they can accommodate a natural solution to the $\mu$-problem). For a detailed study of this kind of scenarios we refer the reader to refs. [3, 4]. In this letter we would like to consider the lightest neutralino in this type of models as a possible good dark matter candidate. We do not attempt a complete exploration of the parameter space of these models (even wider than that of the MSSM) but rather focus on a particular region in which the LSP has properties completely different from those that could be expected in the MSSM. As is described in the next section, in the region we study, the LSP is a neutralino mainly composed of the fermionic superpartner of the singlet used to break the extra $U(1)$. The relic density of such particle, computed in section 3, turns out to be of the right order of magnitude for an interesting dark matter candidate, and the prospects for its laboratory detection are estimated in section 4. Finally, in section 5 we present some conclusions.
2 The Dark Matter Candidate

The symmetry breaking sector in these models includes a chiral multiplet $S$, singlet under the standard model gauge group but with a $U(1)'$ charge $Q_S$. This field couples in the superpotential

$$ W = hSH_1H_2 ,$$

(1)

to the usual Higgs doublets $H_1, H_2$ [with $U(1)'$ charges $Q_1$ and $Q_2$ respectively; gauge invariance requires $Q_1 + Q_2 + Q_S = 0$] and the $\mu$ parameter is dynamically generated by the VEV of $S$ as $\mu_s = h\langle S \rangle = hs/\sqrt{2}$.

The masses of the neutral gauge bosons $Z$ and $Z'$ are

$$ M^2_{Z-Z'} = \begin{bmatrix} \frac{1}{4}G^2(v_1^2 + v_2^2) & \frac{1}{2}g_1'(Q_1v_1^2 - Q_2v_2^2) \\ \frac{1}{2}g_1'(Q_1v_1^2 - Q_2v_2^2) & g_1'^2(Q_1v_1^2 + Q_2v_2^2 + 2v_3^2s^2) \end{bmatrix} ,$$

(2)

where $g, g', g_1'$ are the gauge couplings of $SU(2)_L, U(1)_Y$ and $U(1)'$ respectively, and $G^2 = g^2 + g'^2$. $v_{1,2}$ are the VEVs of $H^0_{1,2}$ with $v_1^2 + v_2^2 \equiv v_W^2 = (246$ GeV$)^2$. For numerical work we use $g_1'^2 = (5/3)g^2$. The $Z - Z'$ mixing angle is constrained experimentally to be less than a few times $10^{-3}$ (although larger values are allowed in some cases, e.g., if the $Z'$ has leptophobic couplings). We will assume in this paper $Q_1 = Q_2$ so that the requirement of a small $Z - Z'$ mixing will force $\tan \beta$ to be close to 1. This constraint is less stringent for larger values of $M_{Z'}$.

The neutralino sector has an extra $U(1)'$ bino and an additional higgsino $\tilde{S}$ (the singlino) besides the four MSSM neutralinos. The $6 \times 6$ mass matrix reads (in the basis $\{\tilde{B}', \tilde{B}, \tilde{W}_3, \tilde{H}_1^0, \tilde{H}_2^0, \tilde{S}\}$):

$$ M^\phi = \begin{pmatrix} M_1' & 0 & 0 & g_1'Q_1v_1 & g_1'Q_2v_2 & g_1'Q_{SS} \\ 0 & M_1 & 0 & -\frac{1}{2}g'v_1 & \frac{1}{2}g'v_2 & 0 \\ 0 & 0 & M_2 & \frac{1}{2}gv_1 & -\frac{1}{2}gv_2 & 0 \\ g_1'Q_1v_1 & -\frac{1}{2}gv_1 & \frac{1}{2}gv_1 & 0 & -\mu_s & -\mu_s \frac{v_2}{s} \\ g_1'Q_2v_2 & \frac{1}{2}g'v_2 & -\frac{1}{2}gv_2 & -\mu_s & 0 & -\mu_s \frac{v_1}{s} \\ g_1'Q_{SS} & 0 & 0 & -\mu_s \frac{v_2}{s} & -\mu_s \frac{v_1}{s} & 0 \end{pmatrix} ,$$

(3)

where $M_1', M_1$ and $M_2$ are the gaugino masses associated with $U(1)', U(1)_Y$ and $SU(2)_L$ respectively. If we assume unification of the gaugino masses at
the gauge unification scale, $M_1', M_1, M_2$ are in the proportion $g'^2k'_1 : \frac{2}{3}g'^2 : g^2$
where $k'_1$ is a normalization constant. With this assumption, the neutralino mass matrix depends (for fixed charges) on two unknown mass parameters, $M_1'$ and $M_{Z'}$ (through $\mu_s$) and the dimensionless $\tan \beta$. In fig. 1 we give a contour plot of the lightest neutralino mass $m_\chi$ in the plane $M_1' - M_{Z'}$ for $\tan \beta = 1$ and $Q_1 = Q_2 = 1$. This mass is zero along a line in the half-plane $M_1' > 0$. For the study of the region around this line (where $\chi$ will be quite light) it is useful to define the normalized neutralino state $\tilde{N}$:

$$\tilde{N} = \frac{1}{N} \left[ \frac{h}{\sqrt{2}g'_1}r^2 \sin 2\beta \hat{B}' + (\overline{Q}_H r^2 - Q_S)\tilde{S} + r_1 (Q_S - \overline{Q}_- r^2)\hat{H}_1^0 + r_2 (Q_S + \overline{Q}_- r^2)\hat{H}_2^0 \right], \quad (4)$$

where $\overline{Q}_- = Q_1c_\beta - Q_2s_\beta$ (when the $Z - Z'$ mixing angle is small due to cancellations, $\overline{Q}_- \simeq 0$); $\overline{Q}_H = Q_1c_\beta + Q_2s_\beta$; $r_{1,2} = v_{1,2}/s$ and $r^2 = r_1^2 + r_2^2$. This state gives a precise analytic description of the $\chi$ composition near the $m_\chi = 0$ line. This is also shown in fig. 1 where the thick solid line delimits the area in which the lightest neutralino has a $\tilde{N}$ purity $P \equiv \tilde{\chi} \cdot \tilde{N} \geq 0.99$. In this region $M_{Z'}$ is typically large ($M_{Z'} > 300$ GeV) so that $r$ is expected to be small. To first order in $r$ we can approximate:

$$\tilde{N} \simeq \frac{1}{\sqrt{1 + r^2}} \left[ r (\cos \beta \hat{H}_1^0 + \sin \beta \hat{H}_2^0) - \tilde{S} \right]. \quad (5)$$

If we write

$$\chi = N_{10} \hat{B}' + N_{30} \hat{H}_1^0 + N_{40} \hat{H}_2^0 + N_{50} \tilde{S}, \quad (6)$$

we find numerically $N_{10} \sim 0.1 - 0.2$, $N_{30} \sim N_{40} \sim 0.2 - 0.3$ and $|N_{50}| \sim 0.9 - 0.95$ in the region of interest. This is our dark matter candidate and it is basically singlino dominated, with a small higgsino doublet component and an even smaller $\hat{B}'$ part. This result is independent of the assumption of unification of the gaugino masses and holds as long as they are large enough. More precisely, we would have obtained the same eq. (4) provided $M_1, M_2 \gg M_Z$ at the electroweak scale.

The rest of neutralinos (and charginos) have masses controlled by the large gaugino masses and $\mu_s$ and are thus much heavier than $\chi$. In these models the typical soft mass scale is of order $M_{Z'}$ so that the spectrum of superpartners is expected to have masses of similar magnitude (a possible
Figure 1: Contour plot of the lightest neutralino mass \(m_\chi/\text{GeV}\) in the plane \(M'_1 - M_{Z'}\) with \(Q_1 = Q_2 = 1\) and \(\tan \beta = 1\). In the region delimited by the thick solid lines, the light neutralino has the composition with a purity \(\geq 0.99\).

exception could be the lightest stop because of a sizeable \(\tilde{t}_R - \tilde{t}_L\) mixing). The fact that electroweak symmetry occurs at a lower scale (and thus \(M_Z \ll M_{Z'}\)) is a result of accidental cancellations among soft masses (see [4]). The lightest scalar Higgs boson, \(h^0\), remains also at the electroweak scale (roughly given by \(M_Z\)) while the rest of Higgs states (two more scalars \(H^0_2, H^0_3\), one pseudoscalar \(A^0\) and a charged pair \(H^\pm\)) have heavy masses comparable to \(M_{Z'}\).

When \(\chi\) is so light that the decay \(Z \to \chi \chi\) is kinematically allowed it gives an extra contribution to the invisible \(Z\) width. The LEP constraint \(\delta \Gamma_{inv} < 4 \text{ MeV}\) is however easily satisfied; the coupling of \(\chi\) to the \(Z\) boson is proportional to \(r^2 \cos 2\beta\) and this is small, first, because \(\tan \beta\) is close to 1 as required to suppress the \(Z - Z'\) mixing, and second, because \(r\) is also small.

In the region of parameters described above, \(\chi\) is the LSP and, with the assumption of conserved \(R\)-parity, becomes a possible candidate for cold dark matter.
3 Relic Abundance

The present relic abundance of $\chi$’s ($\Omega_\chi h^2$) is determined by their annihilation cross section at freeze-out, which, in the non-relativistic expansion applicable in this case, we write as

$$\sigma_{\text{ann}} v \simeq a + \frac{b}{6} v^2,$$

where $v$ is the relative velocity of $\chi$’s in the c.m. frame. The freeze-out temperature $T_F$ (and thus $v$) can be iteratively computed from the condition (obtained from the equality of annihilation and expansion rates):

$$\frac{m_\chi}{T_F} \equiv x_F = \ln \frac{0.1 M_{Pl}(\sigma_{\text{ann}}v)m_\chi}{\sqrt{g_*x_F}},$$

where $M_{Pl} = 1.22 \times 10^{19}$ GeV is the Planck mass, $g_*$ the number of relativistic degrees of freedom at $T_F$ ($\sqrt{g_*} \sim 8 - 9$) and

$$\langle \sigma_{\text{ann}} v \rangle = a + \left( b - \frac{3}{2} a \right) \frac{1}{x_F},$$

is the thermally averaged cross section $\mathbb{1}$. Typically $x_F \simeq 20$ and $v \simeq 1/3$. For the range of masses we consider, $5 - 10 \lesssim m_\chi/GeV \lesssim 70$, $T_F$ is above the QCD quark-hadron phase transition and below the electroweak phase transition. The neutralino relic density is then

$$\Omega_\chi h^2 \equiv \frac{\rho_\chi}{\rho_c/h^2} = \frac{8.77 \times 10^{-11} x_F \text{GeV}^{-2}}{\sqrt{g_*} \left[ a + \frac{1}{2}(b - \frac{3}{2} a) \frac{1}{x_F} \right]},$$

where $\rho_c$ is the critical density and $h$ is the Hubble constant in units of 100 km $Mpc^{-1}$ sec$^{-1}$. $\chi$’s in that mass range can annihilate only into standard model fermion-antifermion pairs $\mathbb{2}$ (top quarks excluded). The relevant processes are mediated by $Z$, $Z'$ and neutral Higgs bosons in the s-channel or by sfermions in the t-channel. The cross section formulae for the MSSM $\mathbb{3}$ can be easily generalized to our model. For the particle spectrum and neutralino composition described previously, the cross section is dominated by $Z'$ exchange with ($\sigma v = \sum_f [a_f + (b_f/6)v^2]$):

$$a_f = \frac{2c_f}{\pi} \beta_f g^4 \frac{m_f^2}{M_{Z'}^2} [Q_f Q_\chi]^2,$$

1Annihilation into gluons or photon pairs is a one-loop process and can be ignored.

5
and

\[ b_f = b_{Z^0Z'}^{(f)} + \left( -\frac{3}{2} + \frac{3}{41 - \xi_f} \right) a_f, \]  

with

\[ b_{Z^0Z'}^{(f)} = \frac{2c_f}{\pi} \beta_f \left( \frac{g_1^2 Q_\chi m_\chi}{4m_\chi^2 - M_{Z'}^2} \right)^2 \left[ (Q_f^V)^2 (2 + \xi_f) + 2(Q_f^A)^2 (1 - \xi_f) \right]. \]

In this formulas, \( c_f = 1(3) \) for leptons (quarks) in the final state, \( \xi_f = m_f^2 / m_\chi^2, \beta_f = \sqrt{1 - \xi_f} \) and \( Q_\chi = Q_1 N_{30}^2 + Q_2 N_{40}^2 + Q_3 N_{50}^2 \) for a \( \chi \) with composition as given by eq. (6). The axial and vector \( U(1)' \) charges of the final state fermions are

\[ Q_f^A = \frac{1}{2} \left[ Q'(f_L) - Q'(f_R) \right], \quad Q_f^V = \frac{1}{2} \left[ Q'(f_L) + Q'(f_R) \right]. \]

The dependence of \( \sigma_{\text{ann}}v \) on these charges introduces some model dependence in the results. However, annihilation into \( b\bar{b} \) usually dominates so that \( \Omega_{\chi h^2} \) depends basically on \( Q_f^A \), which is equal to \(-Q_1 \) if the bottom mass is generated by \( \langle H_1 \rangle \).

Other subdominant annihilation channels are discussed below:

- The amplitude for \( Z^0 \)-mediated annihilation, which is usually dominant for higgsino LSP’s, is proportional to \((N_{30}^2 - N_{40}^2)\), and therefore is doubly suppressed in our case: \( \tan \beta \simeq 1 \) implies \( N_{30} \simeq N_{40} \) and, furthermore, \( N_{30}^2 \) and \( N_{40}^2 \) are very small.

- The t-channel contribution from sfermions, which are expected to have masses comparable to \( M_{Z'} \), is suppressed by extra powers of small Yukawa couplings and/or the smallness of \( N_{10}, N_{30}, N_{40} \).

- The amplitude for \( \chi\chi \rightarrow A^0 \rightarrow f\bar{f} \) has a gauge part suppressed by the smallness of \( N_{10} N_{30}, N_{10} N_{40} \). There is also a new part proportional to \( hN_{50}(N_{30} \cos \beta + N_{40} \sin \beta) \) due to the new Yukawa coupling in the superpotential (7) and, the sizeable value expected for \( h \) \( (h \simeq 0.7 \) from renormalization group analyses (6)), can in principle compensate for the smallness of \( N_{30}, N_{40} \). However, assuming \( m_A \sim M_{Z'} \), the \( A^0 \) contribution is suppressed with respect to the \( Z' \) contribution by an extra factor \( m_\chi/v_W \).

- The same suppression factor appears for \( \chi\chi \rightarrow H_2^0 \rightarrow f\bar{f} \). In addition, the \( \chi\chi H_2^0 \) coupling goes like \( \cos 2\beta \) in the limit \( m_A \simeq m_{H^0} \gg M_Z \) and thus is unimportant for \( \tan \beta \simeq 1 \).
In the limit $M'_Z \gg M_Z$, the third scalar, $H^0_3$, is singlet dominated and does not mediate the annihilation into fermions.

More important can be the $h^0$ mediated annihilation; the suppression factor $m_\chi/v_W$ can be compensated by the fact that $m_{h^0}$ cannot be as heavy as $M_{Z'}$ and, in addition, there can be a resonant enhancement of the cross section for $2m_\chi \simeq m_h$. Nevertheless, we have checked numerically that this channel can also be neglected. First, the CP odd nature of the $\chi\chi$ s-wave initial state forces $h^0$ to contribute to $\sigma_{\text{ann}}v$ only to order $v^2$. Second, in these models the mass of $h^0$ receives extra contributions and is given by

$$m^2_{h^0} = \left[ \frac{1}{4} G^2 \cos^2 2\beta + h^2 \sin^2 2\beta + g'^2 Q_H^2 \right] v^2_W + \Delta_{\text{rad}} m^2_{h^0},$$

where the last piece, coming from loop corrections can also be sizeable. One sees that, even in the case $\tan \beta = 1$, the tree-level mass can easily be as large as 170 GeV, and $2m_\chi$ is never close to the pole for light $m_\chi$.

Finally, as we showed above, the rest of neutralinos and the charginos are much heavier than $\chi$ and “co-annihilation” effects (such as $\chi\chi' \rightarrow f \bar{f}$ or $\chi\chi^\pm \rightarrow f \bar{f}'$) play no role.

Figure 2 shows the relic density of $\chi$’s vs. its mass for different values of $M_{Z'}$ from 300 GeV (lower curve) to 800 GeV (upper). For concreteness, we have fixed $N_{30} = N_{40} = 0.25$ and $N_{50} = -0.9$ in eq. (4) and included Higgs subdominant annihilation channels to compute the abundance. We use $Q_1 = Q_2 = 1$ and fix $m_h = 170$ GeV, and $m_A = M_{Z'}$. In the range of $m_\chi$ shown, $Z'$ gives the dominant annihilation channel, as explained above, so that larger $M_{Z'}$ reduces the annihilation rate making the relic abundance grow. For $m_\chi \gtrsim 70$ GeV, $h^0$-exchange starts to be important and reduces significantly the relic density. Smaller values of the $U(1)'$ charges would also tend to increase the abundance but, on the other hand, $m_h$ is lighter in that case and will become important for lighter values of $m_\chi$.

In the figure, we show only the region $\Omega_\chi h^2 < 1$, conservative limit which follows from a lower bound on the age of the Universe of $\sim 10$ Gyr. Taking the uncertainty in the Hubble constant to be $0.4 \lesssim h \lesssim 1$, the cosmologically interesting range is $0.01 \lesssim \Omega_\chi h^2 \lesssim 0.5$ if $\chi$’s form a significant fraction of the total dark matter in the Universe. The upper bound $\Omega_\chi h^2 < 1$ sets a lower limit on the value of $m_\chi$ for which $\chi$ is a good CDM candidate, given fixed values of the charges and $M_{Z'}$. This lower bound increases with $M_{Z'}$ and
can be very small for low $M_{Z'}$ without conflicting with experimental bounds from $Z^0$ decays.

A non-negligible $\chi\chi \rightarrow Z' \rightarrow f\bar{f}$ rate usually implies a sizeable $Z'$ production cross section at hadron colliders and one should check that the interesting region for $\Omega_{\chi}h^2$ is not in conflict with the non-observation of $Z'$ events (basically $p\bar{p} \rightarrow Z' \rightarrow e^+e^-$) at the Tevatron CDF and D0 experiments. The lower limits on $M_{Z'}$ obtained by these experiments are model-dependent (all the $U(1)'$ charges and masses of the possible decay products of $Z'$ enter the computation) and we do not attempt such detailed analysis. We remark, however, that the supersymmetric decays of the $Z'$ (in particular the invisible decay $Z' \rightarrow \chi\chi$ can be very important) reduce significantly the $Z' \rightarrow e^+e^-$ branching ratio, so that the usual limits on $M_{Z'}$ are relaxed and can be easily evaded in our model.

It is illustrative to compare our dark matter candidate with similar light higgsino candidates proposed in other supersymmetric models. In the MSSM one such light neutralino, with composition $\chi \simeq \sin \beta H^0_1 + \cos \beta H^0_2$, was
studied in [8] motivated by the SUSY interpretation of the CDF $e e \gamma\gamma + \not{E}_T$ event. The $\chi\chi Z^0$ coupling, proportional to $\cos 2\beta$, is reduced for $\tan \beta \approx 1$. This suppresses the contribution to the invisible $Z^0$ width and the $\chi\chi \rightarrow Z \rightarrow f\bar{f}$ cross-section, producing relic abundances in the interesting range.

In the NMSSM, the minimal model extended by an extra singlet but no additional $U(1)$, another light higgsino dark matter candidate was studied in [10] (for more general analyses of NMSSM neutralino dark matter see [11]). In this case $\chi \simeq (\cos \beta \tilde{H}_1^0 + \sin \beta \tilde{H}_2^0 + \alpha \tilde{S})/N$ with all three components of similar magnitude. The contribution to $\Gamma_{Z'}^{\text{inv}}$, also proportional to $\cos^2 2\beta$ in this case, is further reduced by the non-negligible $\tilde{S}$ component. If $\chi \simeq \tilde{S}$, the relic abundance would be too large because $\tilde{S}$’s do not annihilate efficiently.

Our model is similar to the previous case but with $\tilde{S}$ dominant in the $\chi$ composition. Now, however, $\chi$’s annihilate through $\chi\chi \rightarrow Z' \rightarrow f\bar{f}$. The rate is not too large due to the smallness of $g_1' Q S$ and the large mass of the $Z'$ boson, resulting in a relic abundance of the right order of magnitude.

Finally, some comments from the model building point of view are in order. The region of parameter space we have examined corresponds to $M_1' \gg M_{Z'}$ and $M_1, M_2 \gg M_Z$. Such hierarchy of masses cannot be easily accommodated in models with universal boundary conditions at a high energy scale (a GUT scale or the string scale) and would rather point to models in which soft breaking is dominated by gaugino masses. Models of this kind (in the context of the MSSM) have been considered in the past [12]. A renormalization group analysis of the evolution of parameters from low-energy to the string scale, with particular attention to symmetry breaking constraints, would be required.

4 Detection Rates

If $\chi$’s form the bulk of the dark matter in our galactic halo (with density $\rho_0 \sim 0.3 \text{ GeV/cm}^3$ measured by its gravitational effects), we would like to estimate the prospects for its detection. We focus on direct detection experiments, which are the best suited for small mass WIMPs. In these experiments one hopes to detect calorimetrically nuclear recoils in specialized materials after elastic scattering with the flux of $\chi$’s. In the limit $N_{50} = 1$, the (spin-dependent) $\chi$-nucleus scattering proceeds by $Z'$ exchange and the
rate of events per day and kg of material is \[ R = \frac{\sigma_{sd} \rho_0}{m_\chi m_N} \left( \xi v \right)_{sd} \left[ \frac{7.3 \times 10^{55}}{\text{kg} \cdot \text{day}} \right], \tag{16} \]

where \( m_N \) is the nucleus mass, \( (\xi v)_{sd} \) takes into account the nuclear form factor suppression and velocity distribution of \( \chi \)'s, and \[ \sigma_{sd} = \frac{16}{\pi} m_r^2 J(J + 1) \left[ \sum_{q=u,d,s} d_q (\Delta q_q^{(p)} \langle S_p \rangle + \Delta q_q^{(n)} \langle S_n \rangle) \right]^2. \tag{17} \]

In this formula, \( J \) is the spin of the nucleus, \( m_r = m_N m_\chi / (m_N + m_\chi) \) is the reduced mass of the \( \chi \)-nucleon system, \( \Delta q_q^{(p)} (\Delta q_q^{(n)}) \) is the quark spin content of the proton (neutron), \( \langle S_{p,n} \rangle \) is the expectation value of the spin content of the proton (or neutron) group in the nucleus and finally \[ d_q = g_1^2 \frac{Q_q^4 Q_\chi}{M_{Z'}^2}, \tag{18} \]
gives the effective neutralino-quark axial coupling.

For non-zero \( N_{30}, N_{40} \) there is also a scalar interaction mediated by \( h^0 \) exchange. The corresponding rate is of the form (16) with a different form factor \( (\xi v)_{sc} \) and \( \sigma_{sd} \) replaced by \[ \sigma_{sc} = \frac{4}{\pi} m_r^2 [Z f_p + (A - Z) f_n]^2. \tag{19} \]

\( Z \) and \( A \) are the nuclear charge and atomic number respectively. \( f_p \) and \( f_n \), the effective scalar couplings of \( \chi \) to protons and neutrons, are proportional to \[ f_h^{(p,n)} = \frac{h m_{p,n}}{\sqrt{2} v_W} N_{50} (N_{30} \sin \beta + N_{40} \cos \beta) \frac{1}{m_h^2}, \tag{20} \]
and \( m_{p,n} \) are the proton and neutron masses, so that \( f_p \simeq f_n \). For heavy nuclei, the \( A^2 \) factor can compensate for the smallness of \( N_{30}, N_{40} \), and \( \sigma_{sc} \) eventually dominates over the spin-dependent rate.

We have computed the total rate (spin-dependent plus scalar) for two typical materials used in dark matter detectors, \(^{73}\)Ge and \(^{29}\)Si, both with \( J \neq 0 \). The results, as a function of \( m_\chi \), are presented in fig. 3 for the same choice of parameters used in fig. 2. In each set of curves, the upper one corresponds to \( M_{Z'} = 300 \text{ GeV} \) and the lowest to \( M_{Z'} = 800 \text{ GeV} \). For \( M_{Z'} > 400 \text{ GeV} \) the curves are nearly indistinguishable: the spin-dependent rate is negligible and
the scalar rate (independent of $M_{Z'}$) gives all the effect. The small difference between the two limiting curves in each set shows that the scalar interaction provides the dominant effect. The rates obtained are below the current experimental sensitivity [1]. For instance the CDMS experiment, using germanium, is planning on obtaining a sensitivity of $\sim 10^{-1}$ event/kg/day after a year of exposure. However, the next generation of experiments is expected to achieve sensitivities $\sim 0.01$ event/kg/day and therefore, we conclude that $\chi$’s could well be on the reach of near future experiments.

5 Conclusions

In supersymmetric models with an extra $U(1)$ (broken at the TeV scale by the VEV of an standard model singlet $S$) the LSP neutralino can be singlino dominated. We have showed that the thermal relic density of this fermion, stable if $R$-parity is conserved, can be of the right order of magnitude to be a good cold dark matter candidate. We have also estimated direct detection
rates at typical cryogenic devices in search of halo dark matter and found that they are typically small but may be reachable by the next generation of experiments.

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