of dark matter. Before discussing the results of direct and indirect searches, we have reviewed the present and future experiments on which they are based.

Among the most promising dark matter searches appears to be direct detection. The current situation is complicated by the claim of a positive detection by the DAMA experiment, which have been contradicted by several other experiments. It is unclear, but more and more improbable, whether it is possible to find a theoretical scenario that accommodates all the experimental findings. The much higher (several orders of magnitude) sensitivity of future experiments should be able to solve this controversy.

Indirect dark matter detection via annihilations in the Galactic Centerregion is also an exciting possibility, although the prospects for the observation of gamma-rays, neutrinos and synchrotron radiation from that direction strongly depend on astrophysical parameters, such as the profile of dark matter in the innermost regions, which unfortunately are poorly known. Nevertheless, the development of next-generation gamma-ray and neutrino telescopes will allow us to test many scenarios, especially if effects such as the adiabatic accretion onto the central black hole significantly enhance the dark matter density and corresponding annihilation signal. If the Galactic Centerturns out to contain less dark matter, observations of dwarf galaxies, external galaxies and local dark substructure may play an important role for indirect searches.

Indirect searches for dark matter through the observation of high energy neutrinos produced in dark matter annihilations in the Sun are also promising. These rates do not depend strongly on the dark matter halo distribution and are thus fairly unambiguous probes of particle dark matter models. Measurement of the positron and anti-proton spectra, which are soon to improve dramatically, can also provide an opportunity to observe products of dark matter annihilations in the galactic halo.

Collectively, the direct, indirect and collider searches for particle dark matter have incredible prospects for discovery in the coming years. We hope that this review can be a useful tool in guiding members of the scientific community closer to the goal of dark matter identification which has eluded us for so long.

8 Acknowledgements

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A Neutralino Mass Eigenstates

In the Minimal Supersymmetric Standard Model (MSSM), the neutral electroweak gauginos $(\tilde{B}, \tilde{W^3})$ and higgsinos $(\tilde{H}_1^0, \tilde{H}_2^0)$ have the same quantum numbers and, therefore, mix into four mass eigenstates called neutralinos. The neutralino mass matrix in the $\widetilde{B}$-$\widetilde{W}^3$-$\widetilde{H}_1^0$-$\widetilde{H}_2^0$ basis is given by

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -M_Z\cos\beta\sin\theta_W & M_Z\sin\beta\sin\theta_W \\ 0 & M_2 & M_Z\cos\beta\cos\theta_W & -M_Z\sin\beta\cos\theta_W \\ -M_Z\cos\beta\sin\theta_W & M_Z\cos\beta\cos\theta_W & 0 & -\mu \\ M_Z\sin\beta\sin\theta_W & -M_Z\sin\beta\cos\theta_W & -\mu & 0 \end{pmatrix}, \tag{146}$$

where M_1, M_2 and μ are the bino, wino and higgsino mass parameters, respectively, θ_W is the Weinberg angle and $\tan\beta$ is the ratio of the vacuum expectation values of the Higgs bosons. This matrix can be diagonalized by the matrix, N.

$$M_{\chi^0}^{\rm diag} = N^\dagger M_{\chi^0} N. \tag{147}$$

The masses of the four mass eigenstates are then given by [198, 58]

$$\epsilon_1 M_{\chi_1^0} = -(\frac{1}{2}a - \frac{1}{6}C_2)^{1/2} + \left[-\frac{1}{2}a - \frac{1}{3}C_2 + \frac{C_3}{(8a-\frac{8}{3}C_2)^{1/2}}\right]^{1/2} + \frac{1}{4}(M_1+M_2)\,, \tag{148}$$

$$\epsilon_2 M_{\chi_2^0} = +(\frac{1}{2}a - \frac{1}{6}C_2)^{1/2} - \left[-\frac{1}{2}a - \frac{1}{3}C_2 - \frac{C_3}{(8a-\frac{8}{3}C_2)^{1/2}}\right]^{1/2} + \frac{1}{4}(M_1+M_2)\,, \tag{149}$$

$$\epsilon_3 M_{\chi_3^0} = -(\frac{1}{2}a - \frac{1}{6}C_2)^{1/2} - \left[-\frac{1}{2}a - \frac{1}{3}C_2 + \frac{C_3}{(8a-\frac{8}{3}C_2)^{1/2}}\right]^{1/2} + \frac{1}{4}(M_1+M_2)\,, \tag{150}$$

$$\epsilon_4 M_{\chi_4^0} = +(\frac{1}{2}a - \frac{1}{6}C_2)^{1/2} + \left[-\frac{1}{2}a - \frac{1}{3}C_2 - \frac{C_3}{(8a-\frac{8}{3}C_2)^{1/2}}\right]^{1/2} + \frac{1}{4}(M_1+M_2)\,, \tag{151}$$

where ϵ_i is the sign of the ith eigenvalue of the neutralino mass matrix, and

$$C_2 = (M_1M_2 - M_Z^2 - \mu^2) - \frac{3}{8}(M_1+M_2)^2\,, \tag{152}$$

$$\begin{aligned} C_3 = & -\frac{1}{8}(M_1+M_2)^3 + \frac{1}{2}(M_1+M_2)(M_1M_2 - M_Z^2 - \mu^2) + (M_1+M_2)\mu^2 \\ & +(M_1\cos^2\theta_W + M_2\sin^2\theta_W)M_Z^2 + \mu M_Z^2 \sin 2\beta\,, \end{aligned} \tag{153}$$

$$\begin{aligned} C_4 = & -(M_1\cos^2\theta_W + M_2\sin^2\theta_W)M_Z^2\mu\sin 2\beta - M_1M_2\mu^2 \\ & +\frac{1}{4}(M_1+M_2)[(M_1+M_2)\mu^2 + (M_1\cos^2\theta_W + M_2\sin^2\theta_W)M_Z^2 + \mu M_Z^2\sin 2\beta] \\ & +\frac{1}{16}(M_1M_2 - M_Z^2 - \mu^2)(M_1+M_2)^2 - \frac{3}{256}(M_1+M_2)^4\,, \end{aligned} \tag{154}$$

$$a = \frac{1}{2^{1/3}} \mathrm{Re}\left[-S + i(D/27)^{1/2}\right]^{1/3} , \tag{155}$$

$$D = -4U^3 - 27S^2 , \quad U = -\frac{1}{3}C_2^2 - 4C_4, \quad S = -C_3^2 - \frac{2}{27}C_2^3 + \frac{8}{3}C_2C_4 . \tag{156}$$

The four masses above are not generally in the order $M_{\chi_1^0} < M_{\chi_2^0} < M_{\chi_3^0} < M_{\chi_4^0}$, although it is conventional to relabel the states, from lightest to heaviest.

The mixing matrix, N, is then given by [198, 58]

$$\frac{N_{i2}}{N_{i1}} = -\frac{1}{\tan\theta_W} \frac{M_1 - \epsilon_i M_{\chi_i^0}}{M_2 - \epsilon_i M_{\chi_i^0}} , \tag{157}$$

$$\frac{N_{i3}}{N_{i1}} = \frac{-\mu[M_2 - \epsilon_i M_{\chi_i^0}][M_1 - \epsilon_i M_{\chi_i^0}] - M_Z^2 \sin\beta\cos\beta[(M_1 - M_2)\cos^2\theta_W + M_2 - \epsilon_i M_{\chi_i^0}]}{M_Z[M_2 - \epsilon_i M_{\chi_i^0}]\sin\theta_W[-\mu\cos\beta + \epsilon_i M_{\chi_i^0}\sin\beta]} , \tag{158}$$

$$\frac{N_{i4}}{N_{i1}} = \frac{-\epsilon_i M_{\chi_i^0}[M_2 - \epsilon_i M_{\chi_i^0}][M_1 - \epsilon_i M_{\chi_i^0}] - M_Z^2 \cos^2\beta[(M_1 - M_2)\cos^2\theta_W + M_2 - \epsilon_i M_{\chi_i^0}]}{M_Z[M_2 - \epsilon_i M_{\chi_i^0}]\sin\theta_W[-\mu\cos\beta + \epsilon_i M_{\chi_i^0}\sin\beta]} , \tag{159}$$

and

$$N_{i1} = \left[1 + \left(\frac{N_{i2}}{N_{i1}}\right)^2 + \left(\frac{N_{i3}}{N_{i1}}\right)^2 + \left(\frac{N_{i4}}{N_{i1}}\right)^2\right]^{-1/2} . \tag{160}$$

The lighest neutralino (χ_1^0) is a mixture of gauginos and higgsinos:

$$\chi_1^0 = N_{11}\tilde{B} + N_{12}\tilde{W}^3 + N_{13}\tilde{H}_1^0 + N_{14}\tilde{H}_2^0. \tag{161}$$

The gaugino fraction of χ_1^0 is defined as

$$f_G = N_{11}^2 + N_{12}^2 \tag{162}$$

and its higgisino fraction as

$$f_H = N_{13}^2 + N_{14}^2. \tag{163}$$

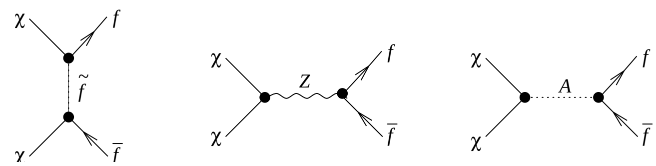


Figure 38: Tree level diagrams for neutralino annihilation into fermion pairs. From Ref. [305].

B Neutralino Annihilation Cross Sections in the Low Velocity Limit

In this appendix, we give the amplitudes and cross sections for the most important neutralino annihilation channels in the low velocity limit (the first term in the expansion $\sigma v = a + bv^2 + ...$). This is sufficient for indirect detection but generally insufficient for relic density calculations in which velocity dependent contributions are important. For a more complete list, with all S and P-wave tree level annihilation amplitudes, see Refs. [187, 305, 379, 101].

B.1 Annihilation Into Fermions

Neutralinos can annihilate to fermion pairs by three tree level diagrams [187, 204, 262, 263]. These processes consist of s-channel exchange of pseudoscalar Higgs and Z^0-bosons and t-channel exchange of sfermions (see Fig. 38).

The amplitude for pseudoscalar Higgs exchange is given by

$$\mathcal{A}_A = 4\sqrt{2}\, g\, T_{A\,11}\, h_{Aff}\ \frac{1}{4-(m_A/m_\chi)^2 + i\,\Gamma_A m_A/m_\chi^2}. \tag{164}$$

Here, m_A is the pseudoscalar Higgs mass and Γ_A is the pseudoscalar Higgs width. $T_{A\,11}$ is the A^0-neutralino-neutralino coupling and is given by

$$T_{A\,11} = -\sin\beta Q''_{1,1} + \cos\beta S''_{1,1}, \tag{165}$$

where $Q''_{1,1} = N_{3,1}(N_{2,1} - \tan\theta_W N_{1,1})$ and $S''_{1,1} = N_{4,1}(N_{2,1} - \tan\theta_W N_{1,1})$. N is the matrix which diagonalizes the neutralino mass matrix in the $\widetilde{B}$-$\widetilde{W}^3$-$\widetilde{H}_1^0$-$\widetilde{H}_2^0$ basis, $M_{\chi^0}^{\mathrm{diag}} = N^\dagger M_{\chi^0} N$ (see Appendix A). θ_W is the Weinberg angle and $\tan\beta$ is the ratio of the Higgs vacuum expectation values. h_{Aff} is the A^0-fermion-fermion Yukawa coupling. For up-type fermions, this is given by

$$h_{Aff} = -\frac{g m_f \cot\beta}{2 m_{W^\pm}}. \tag{166}$$

For down-type fermions, it is

$$h_{Aff} = -\frac{g m_f \tan\beta}{2 m_{W^\pm}}. \tag{167}$$

The amplitude for neutralino annihilation via sfermion exchange to a pair of fermions, $f_i \bar{f}_i$, is given by

$$\mathcal{A}_{\tilde{f}} = \sqrt{2} \sum_{j=1}^{6} \frac{1}{P_j} \Bigg(\Big[(X'_{ij1})^2 + (W'_{ij1})^2 \Big] \frac{m_{f_i}}{m_\chi} + 2 X'_{ij1} W'_{ij1} \Bigg), \tag{168}$$

where $P_j = 1 + (m_{\tilde{f}_j}/m_\chi)^2 - (m_{f_i}/m_\chi)^2$ and the sum is over the six sfermion states which couple to the final state fermion. The fermion-sfermion-neutralino couplings, X'_{ij1} and W'_{ij1}, are given by

$$X'_{ij1} = X_1 (\Pi_L \Theta_f)_{i,j} + Z_{i,k,1} (\Pi_R \Theta_f)_{k,j} \tag{169}$$

and

$$W'_{ij1} = Y_1 (\Pi_R \Theta_f)_{i,j} + Z_{i,k,1} (\Pi_L \Theta_f)_{k,j}, \tag{170}$$

where

$$X_1 = -g\sqrt{2} \, [\, T_3(f_i) N^*_{2,1} - \tan\theta_W \, (\, T_3(f_i) - e(f_i) \,) \, N^*_{1,1} \,], \tag{171}$$

and

$$Y_1 = g\sqrt{2} \, \tan\theta_W \, e(f_i) N^*_{1,1}. \tag{172}$$

For final state up-type quarks,

$$Z_{i,j,1} = -\frac{g}{\sqrt{2}\, m_{W^\pm} \, \sin\beta} \Theta_{i,j} \, N^*_{4,1}. \tag{173}$$

For final state down-type quarks,

$$Z_{i,j,1} = -\frac{g}{\sqrt{2}\, m_{W^\pm} \, \cos\beta} \Theta_{i,j} \, N^*_{3,1}. \tag{174}$$

And for final state leptons,

$$Z_{i,j,1} = -\frac{g}{\sqrt{2}\, m_{W^\pm} \, \cos\beta} \Theta_{i,j} \, N^*_{3,1}. \tag{175}$$

Here, $T_3(f_i)$ and $e(f_i)$ are the weak hypercharge and electric charge of the final state fermion. N, again, is the matrix which diagonalizes the neutralino mass matrix. Θ_f's are the appropriate 6 x 6 sfermion mass matrices and $\Pi_{L,R}$ are left and right projection operators:

$$\Pi_L = \begin{pmatrix} 1 & 0 & 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 & 0 & 0 \end{pmatrix}, \tag{176}$$

$$\Pi_R = \begin{pmatrix} 0 & 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 0 & 0 & 1 \end{pmatrix}. \tag{177}$$

Lastly, the amplitude for neutralino annihilation to fermions via Z exchange is given by

$$\mathcal{A}_Z = 2\sqrt{2}\frac{g^2}{\cos^2\theta_{\rm W}} O''^{L}_{1,1} T_3(f_i)\frac{m_{f_i} m_\chi}{m_Z^2}, \tag{178}$$

where $T_3(f_i)$ is the weak hypercharge of the fermion. The coupling $O''^{L}_{1,1}$ is given by $\frac{1}{2}(-N_{3,1}N^*_{3,1} + N_{4,1}N^*_{4,1})$.

Summing these three contributions to the amplitude, we can calculate the cross section for this process:

$$\sigma v(\chi\chi \to \bar{f}_i f_i)_{v\to 0} = \frac{c_f \beta_f}{128\pi m_\chi^2}|\mathcal{A}_A(\chi\chi \to \bar{f}_i f_i) + \mathcal{A}_{\tilde{f}}(\chi\chi \to \bar{f}_i f_i) + \mathcal{A}_Z(\chi\chi \to \bar{f}_i f_i)|^2, \tag{179}$$

where $\beta_f = \sqrt{1 - m_f^2/m_\chi^2}$. c_f is a color factor which is equal to three for quark final states and one otherwise.

We emphasize that all tree level (low velocity) neutralino annihilation diagrams to fermion pairs have amplitudes which are proportional to the final state fermion mass. For sfermion and Z^0 exhange, this is because the Z^0-fermion-fermion and neutralino-fermion-fermion couplings preserve chirality. For psuedoscalar Higgs exchange, the amplitude introduces an explicit factor of the fermion mass in the Yukawa coupling. We also note that the Yukawa coupling which appears in the psuedoscalar Higgs exchange amplitude is proportional to $\tan\beta$ for down-type quarks and $\cot\beta$ for up-type quarks. The net result of these observations is that neutralino annihilation into fermions will be dominated by heavy final states, $b\bar{b}$, $\tau^-\tau^+$ and, if kinematically allowed, $t\bar{t}$. Furthermore, if $\tan\beta$ is large, bottom-type fermions may dominate over up-type fermions, even if less massive. For example, annihilations to $b\bar{b}$ may dominate over $t\bar{t}$, even for heavy neutralinos.

B.2 Annihilation Into Gauge Bosons

Generally, neutralinos can annihilate into gauge boson pairs via several processes (see Fig. 39) [187, 265, 384, 385]. In the low velocity limit, however, only t-channel processes via chargino or neutralino exchange are non-vanishing.

In the low velocity limit, the amplitude for neutralino annihilation to $W^\pm$-pairs is given by

$$\mathcal{A}(\chi\chi \to W^+W^-)_{v\to 0} = 2\sqrt{2}\,\beta_W g^2 \sum_{n=1}^{2}\left[(O^L_{1,n})^2 + (O^R_{1,n})^2\right]\frac{1}{P_n}, \tag{180}$$

where $\beta_W = \sqrt{1 - m_W^2/m_\chi^2}$ and $P_n = 1 + (m_{\chi^\pm_n}/m_\chi)^2 - (m_W/m_\chi)^2$. The sum is over chargino states. $O^L_{1,n}$ and $O^R_{1,n}$ are the neutralino couplings to charginos given by $\frac{-1}{\sqrt{2}}N_{4,1}V^*_{2,n} + N_{2,1}V^*_{1,n}$ and $\frac{1}{\sqrt{2}}N^*_{3,1}U_{2,n} + N^*_{2,1}U_{1,n}$, respectively, where N, again, is the matrix which diagonalizes the neutralino mass matrix. The V's and U's are components of the chargino mass matrix, in the basis

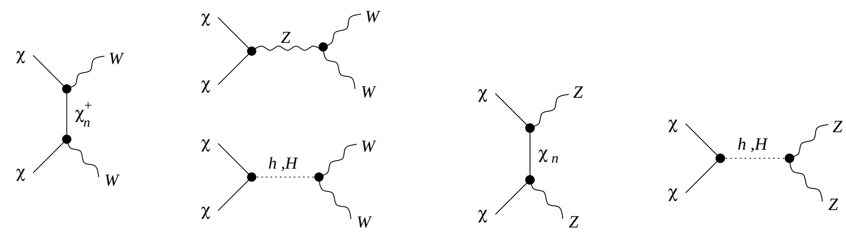


Figure 39: Tree level diagrams for neutralino annihilation into gauge boson pairs. From Ref. [305].

$$U = \begin{pmatrix} \cos\phi_- & -\sin\phi_- \\ \sin\phi_- & \cos\phi_+ \end{pmatrix} \tag{181}$$

and

$$V = \begin{pmatrix} \cos\phi_+ & -\sin\phi_+ \\ \sin\phi_+ & \cos\phi_- \end{pmatrix}, \tag{182}$$

where

$$\tan 2\phi_- = 2\sqrt{2} m_W \frac{(\mu \sin\beta + M_2 \cos\beta)}{(M_2^2 - \mu^2 + 2m_W^2 \cos 2\beta)} \tag{183}$$

and

$$\tan 2\phi_+ = 2\sqrt{2} m_W \frac{(\mu \cos\beta + M_2 \sin\beta)}{(M_2^2 - \mu^2 - 2m_W^2 \cos\beta)}. \tag{184}$$

The amplitude for annihilations to Z^0-pairs is similar:

$$\mathcal{A}(\chi\chi \to Z^0 Z^0)_{v\to 0} = 4\sqrt{2}\,\beta_Z \frac{g^2}{\cos^2\theta_W} \sum_{n=1}^{4} \left(O''^{L}_{1,n}\right)^2 \frac{1}{P_n}. \tag{185}$$

Here, $\beta_Z = \sqrt{1 - m_Z^2/m_\chi^2}$, and $P_n = 1 + (m_{\chi_n}/m_\chi)^2 - (m_Z/m_\chi)^2$. The sum is over neutralino states. The coupling $O''^{L}_{1,n}$ is given by $\frac{1}{2}(-N_{3,1}N^*_{3,n} + N_{4,1}N^*_{4,n})$.

The low velocity annihilation cross section for this mode is then given by

$$\sigma v(\chi\chi \to GG)_{v\to 0} = \frac{1}{S_G} \frac{\beta_G}{128\pi m_\chi^2} |\mathcal{A}(\chi\chi \to GG)|^2, \tag{186}$$

where G indicates which gauge boson is being considered. S_G is a statistical factor equal to one for W^+W^- and two for Z^0Z^0.

It is useful to note that pure-gaugino neutralinos have a no S-wave annihilation amplitude to gauge bosons. Pure-higgsinos or mixed higgsino-gauginos, however, can annihilate efficiently via these channels, even at low velocities.

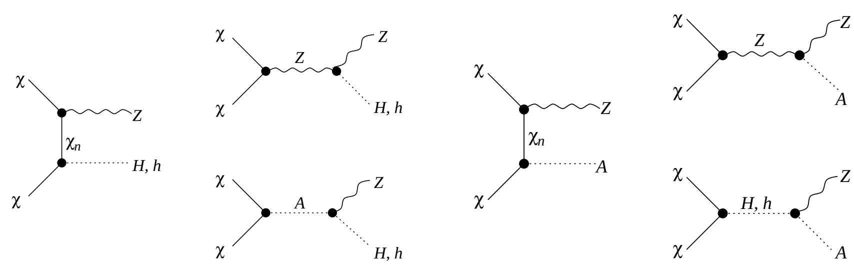


Figure 40: Tree level diagrams for neutralino annihilation into a Z and a Higgs boson. From Ref. [305].

B.3 Annihilation Into Higgs Bosons

There are many tree level diagrams which contribute to neutralino annihilation into Higgs boson pairs or a Higgs boson and a gauge boson (see Figs. 40, 41 and 42)[187, 265, 384, 385, 309, 405, 361].

In the low velocity limit, the amplitude for neutralino annihilation to a Z^0 and a light neutral Higgs, h^0, is given by

$$\mathcal{A}(\chi\chi \rightarrow Z^0 h^0)_{v\rightarrow 0} = -2\sqrt{2}\,\beta_{Zh}\,\frac{m_\chi}{m_Z}\,\frac{g^2}{\cos\theta_W}\Bigg[-2\sum_{n=1}^{4} O''^{L}_{1,n} T_{h\,1,n} \tag{187}$$
$$\times \frac{m_{\chi_n}-m_\chi}{m_\chi P_n} + O''^{L}_{1,1}\frac{m_\chi \sin(\beta-\alpha)}{m_Z\cos\theta_W} - \frac{2\cos(\alpha-\beta)T_{A\,1,1}}{4-m_A^2/m_\chi^2 + i\,\Gamma_A m_A/m_\chi^2}\Bigg],$$

where Γ_A is the pseudoscalar Higgs width and $T_{h\,1,n}$ is the $h^0-\chi_0-\chi_n$ Yukawa coupling (see below). The couplings, $O''^{L}_{1,n}$, is given by $N_{3,1}(N_{2,n}-\tan\theta_W N_{1,n})/2+N_{3,n}(N_{2,1}-\tan\theta_W N_{1,1})/2$ and $P_n = 1+(m_{\chi_n}/m_\chi)^2-\frac{1}{2}(m_Z/m_\chi)^2-\frac{1}{2}(m_h/m_\chi)^2$. $\tan\beta$ is the ratio of the Higgs vacuum expectation values and the mixing angle, α, is related to *beta* by

$$\sin 2\alpha = -\sin 2\beta \left(\frac{m_H^2+m_h^2}{m_H^2-m_h^2}\right) \tag{188}$$

and

$$\cos 2\alpha = -\cos 2\beta \left(\frac{m_A^2-m_Z^2}{m_H^2-m_h^2}\right). \tag{189}$$

The three terms of Eq. 187 correspond to neutralino, Z^0 and A^0 exchange, from first to last.

The expression for neutralino annihilations to a Z^0 and a heavy Higgs boson, H^0, is the same, but with $\sin(\beta-\alpha)$ and $\cos(\alpha-\beta)$ exchanged, and the couplings and masses of h^0 replaced by the couplings and masses of H^0. These Yukawa couplings are given by

$$T_{h\,1,n} = \sin\alpha\, Q''_{1,n} + \cos\alpha\, S''_{1,n} \tag{190}$$

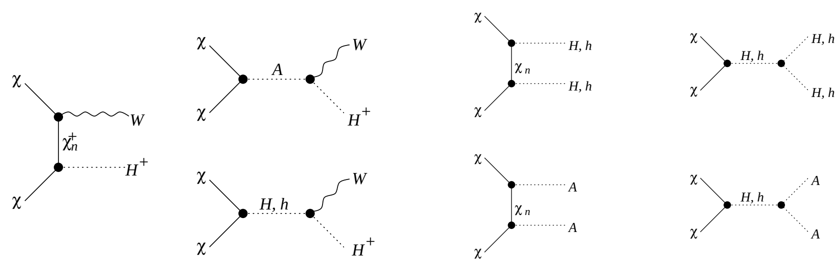


Figure 41: Tree level diagrams for neutralino annihilation into a $W^{\pm}$ and a Higgs boson or a pair of Higgs bosons. From Ref. [305].

and

$$T_{H\,1,n} = -\cos\alpha\, Q''_{1,n} + \sin\alpha\, S''_{1,n}. \tag{191}$$

Here, $S''_{1,n} = N_{4,1}(N_{2,n} - \tan\theta_W N_{1,n})/2 + N_{4,n}(N_{2,1} - \tan\theta_W N_{1,1})/2$. $Q''_{1,n}$ is defined above.

The amplitude for annihilations to a $W^{\pm}$ and a charged Higgs boson is given by

$$\mathcal{A}(\chi\chi \to W^{\pm} H^{\mp})_{v\to 0} = 4\sqrt{2}\,\beta_{WH}\, g^2 \left[-\frac{1}{2}\sum_{n=1}^{2} \frac{m_\chi}{m_W} \frac{O^R_{1,n} Q'^R_{1,n} - O^L_{1,n} Q'^L_{1,n}}{P_n} \right. \tag{192}$$
$$\left. + \frac{1}{2}\sum_{n=1}^{2} \frac{m_{\chi^+_n}}{m_W} \frac{O^R_{1,n} Q'^L_{1,n} - O^L_{1,n} Q'^R_{1,n}}{P_n} - \frac{m_\chi\, T_{A\,11}}{m_W(4 - m_A^2/m_\chi^2)} \right],$$

where $P_n = 1 + (m_{\chi^{\pm}_n}/m_\chi)^2 - \frac{1}{2}(m_{H^{\pm}}/m_\chi)^2 - \frac{1}{2}(m_W/m_\chi)^2$. $O^R_{1,n}$ and $O^L_{1,n}$ are couplings given earlier in this appendix. $Q'^L_{n,m}$ and $Q'^R_{n,m}$ are the chargino-neutralino-charged Higgs couplings, given by

$$Q'^{\,L}_{nm} = \cos\beta \left[N^*_{4n} V^*_{1m} + \sqrt{\frac{1}{2}}(N^*_{2n} + \tan\theta_W N^*_{1n}) V^*_{2m} \right] \tag{193}$$

and

$$Q'^{\,R}_{nm} = \sin\beta \left[N_{3n} U_{1m} - \sqrt{\frac{1}{2}}(N_{2n} + \tan\theta_W N_{1n}) U_{2m} \right], \tag{194}$$

where each of the quantities used have been defined earlier in this appendix. The first and second terms of Eq. 192 correspond to chargino exchange. The third term corresponds to pseudoscalar Higgs exchange (there is no low velocity contribution from scalar Higgs exchange).

Finally, the low velocity amplitude for neutralino annihilation into one neutral Higgs boson and one pseudoscalar Higgs boson is given by

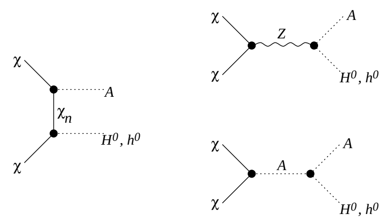


Figure 42: Tree level diagrams for neutralino annihilation into a neutral Higgs boson and a pseudoscalar Higgs boson. From Ref. [305].

$$\mathcal{A}(\chi\chi \to h^0 A^0)_{v\to 0} = \sqrt{2}\, g^2 \Bigg\{ -4 \sum_{n=1}^{4} T_{h0n} T_{A0n} \left[\frac{m_{\chi_n}}{m_\chi P_n} - \frac{m_A^2 - m_h^2}{m_\chi^2} \right] -2 \frac{m_Z}{m_\chi} \frac{\sin(\alpha+\beta)\cos 2\beta}{2\cos\theta_W} \frac{T_{A00}}{4 - m_A^2/m_\chi^2} - \frac{\cos(\alpha-\beta) O''_{L00}}{2\cos^2\theta_W} \frac{m_A^2 - m_a^2}{m_Z^2} \Bigg\}. \tag{195}$$

Here, $P_n = 1 + (m_{\chi_n}/m_\chi)^2 - \frac{1}{2}(m_A/m_\chi)^2 - \frac{1}{2}(m_h/m_\chi)^2$. The other quantities have been defined earlier in this appendix. Again, the amplitude for the analogous process with a heavy rather than light Higgs boson in the final state is the same, but with $\sin(\alpha+\beta)$ and $\cos(\alpha-\beta)$ exchanged and the light Higgs couplings and masses replaced with those for the heavy Higgs boson.

In the low velocity limit, there is no amplitude for neutralino annihilations to H^+H^-, h^0h^0, H^0H^0, A^0A^0 or Z^0A^0.

The low velocity cross section for neutralino annihilation via any of these modes is

$$\sigma v(\chi\chi \to XY)_{v\to 0} = \frac{\beta_{XY}}{128\pi m_\chi^2} |\mathcal{A}(\chi\chi \to XY)_{v\to 0}|^2, \tag{196}$$

where X and Y are labels referring to the final state particles.

B.4 Annihilation Into Photons

Although there are no tree level processes for neutralino annihilation into photons, loop level processes to $\gamma\gamma$ and γZ^0 are very interesting, as they may provide a spectral line feature observable in indirect detection experiments.

In Fig. 43, all of the one-loop diagrams are shown for neutralino annihilation to a pair of photons. In Fig. 44, the corresponding diagrams to a photon and a Z^0 are shown. We do not include the corresponding amplitudes or cross sections here. For those results, see Ref. [74] and Ref. [453] for $\gamma\gamma$ and γZ^0 final states, respectively.

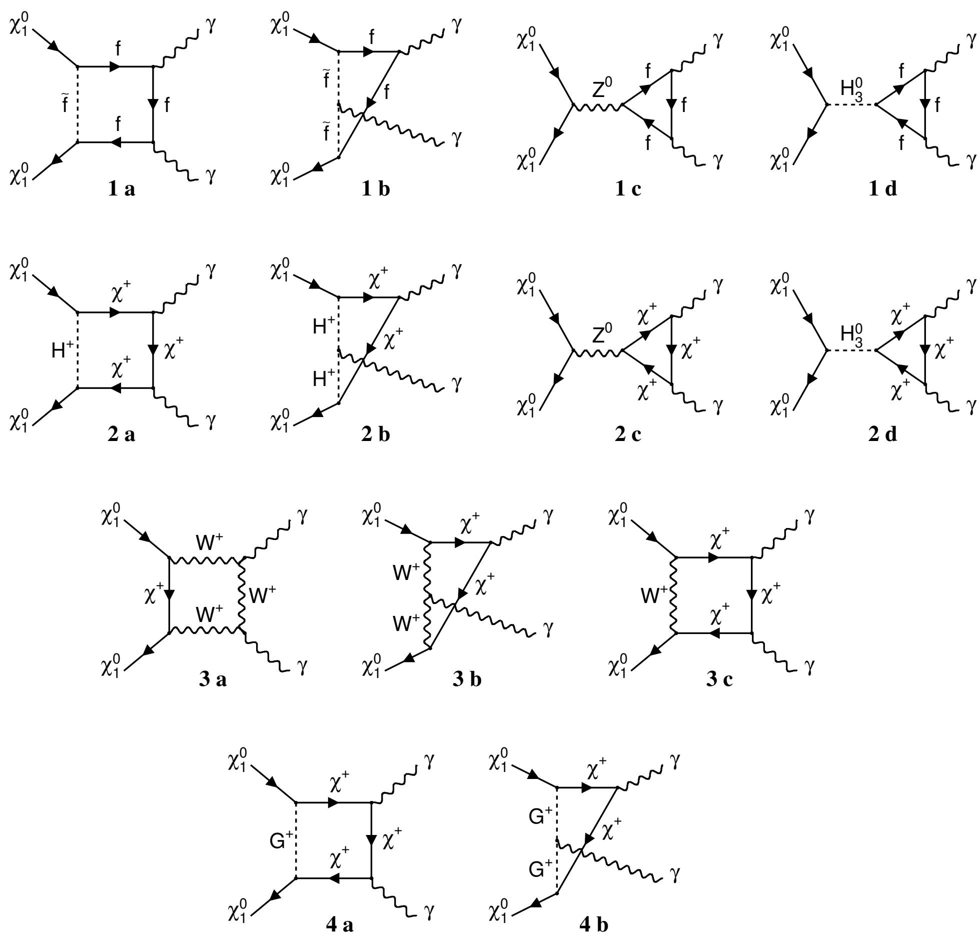
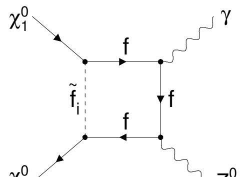
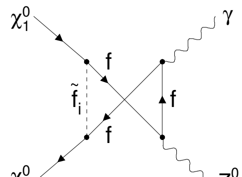


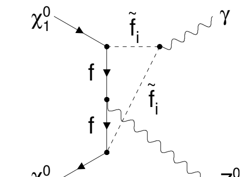
Figure 43: Diagrams contributing, at one loop level, to neutralino annihilation into two photons. From Ref. [74].



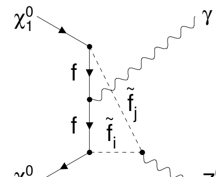
1 a



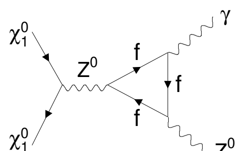
1 b



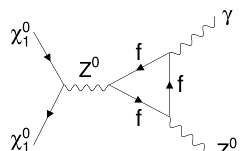
1 c



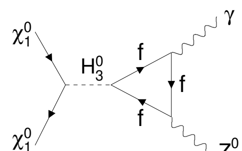
1 d



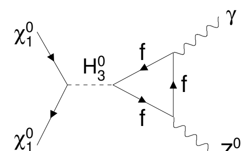
1 e



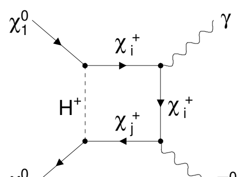
1 f



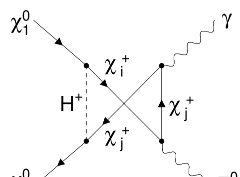
1 g



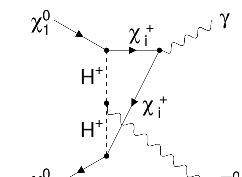
1 h



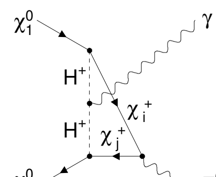
2 a



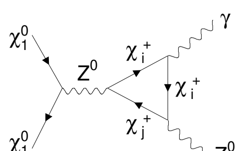
2 b



2 c



2 d



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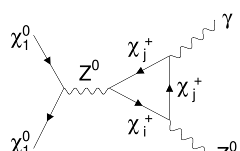

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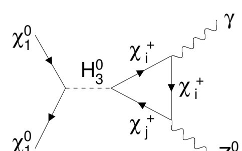

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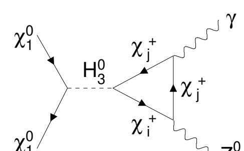

2 h

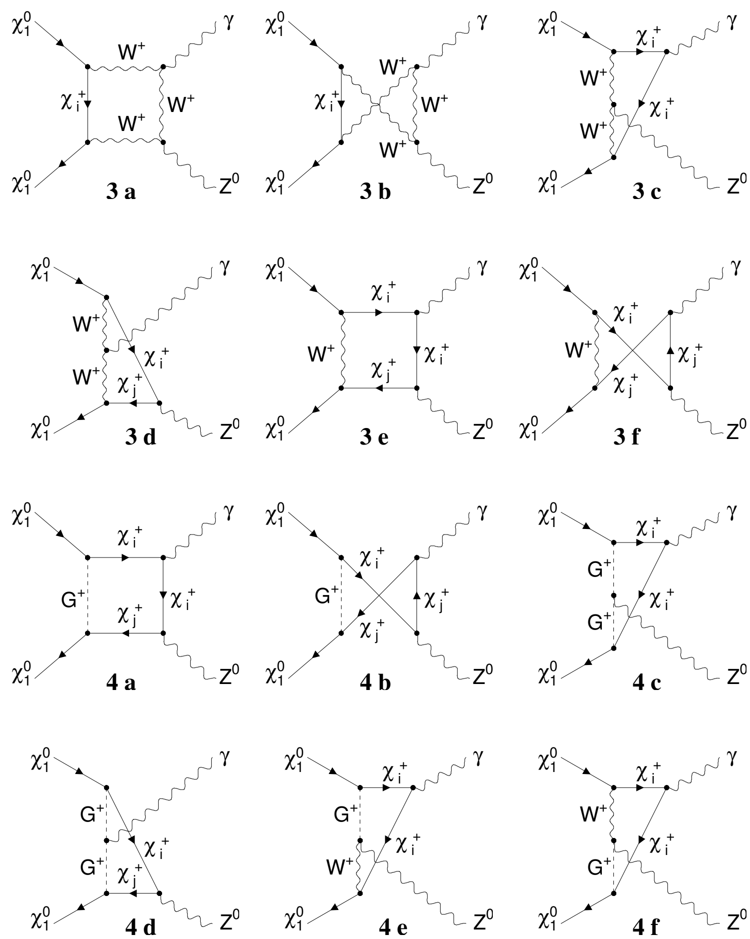

Figure 44: Diagrams contributing, at one loop level, to neutralino annihilation into a photon and a Z^0. From Ref. [453].

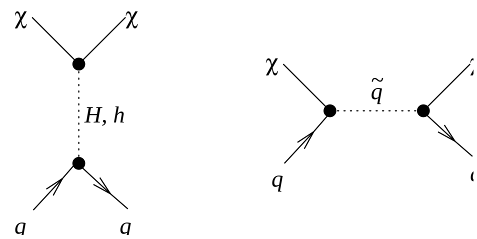

Figure 45: Tree level Feynman diagrams for neutralino-quark scalar (spin-independent) elastic scattering. From Ref. [305].

C Elastic Scattering Processes

C.1 Scalar Interactions

Consider a WIMP with scalar interactions with quarks given by

$$\mathcal{L}_{scalar} = a_q \bar{\chi}\chi\bar{q}q, \tag{197}$$

where a_q is the WIMP-quark coupling. Then the scattering cross section for the WIMP off of a proton or neutron is given by

$$\sigma_{scalar} = \int_0^{4m_r^2 v^2} \frac{d\sigma(v=0)}{d|\vec{v}|^2} = \frac{4m_r^2}{\pi} f_{p,n}^2, \tag{198}$$

where v is the relative velocity of the WIMP, m_r is the reduced mass of the nucleon ($m_r \simeq m_{p,n}$ for WIMPs heavier than ~ 10 GeV) and $f_{p,n}$ is the WIMP coupling to protons or neutrons, given by

$$f_{p,n} = \sum_{q=u,d,s} f_{Tq}^{(p,n)} a_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{(p,n)} \sum_{q=c,b,t} a_q \frac{m_{p,n}}{m_q}, \tag{199}$$

where $f_{Tu}^{(p)} = 0.020 \pm 0.004$, $f_{Td}^{(p)} = 0.026 \pm 0.005$, $f_{Ts}^{(p)} = 0.118 \pm 0.062$, $f_{Tu}^{(n)} = 0.014 \pm 0.003$, $f_{Td}^{(n)} = 0.036 \pm 0.008$ and $f_{Ts}^{(n)} = 0.118 \pm 0.062$ [200]. $f_{TG}^{(p,n)}$ is related to these values by

$$f_{TG}^{(p,n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p,n)}. \tag{200}$$

The term in Eq. 199 which includes $f_{TG}^{(p,n)}$ results from the coupling of the WIMP to gluons in the target nuclei through a heavy quark loop. The couplings of squarks and Higgs bosons to heavy quarks leads to a loop level coupling of the WIMP to gluons [263, 57, 309]. Such diagrams are shown in Fig. 46.

To attain the scalar cross section for a WIMP scattering off of a target nucleus, one should sum over the protons and neutrons in the target:

$$\sigma = \frac{4m_r^2}{\pi}\left(Zf_p + (A-Z)f_n\right)^2, \tag{201}$$

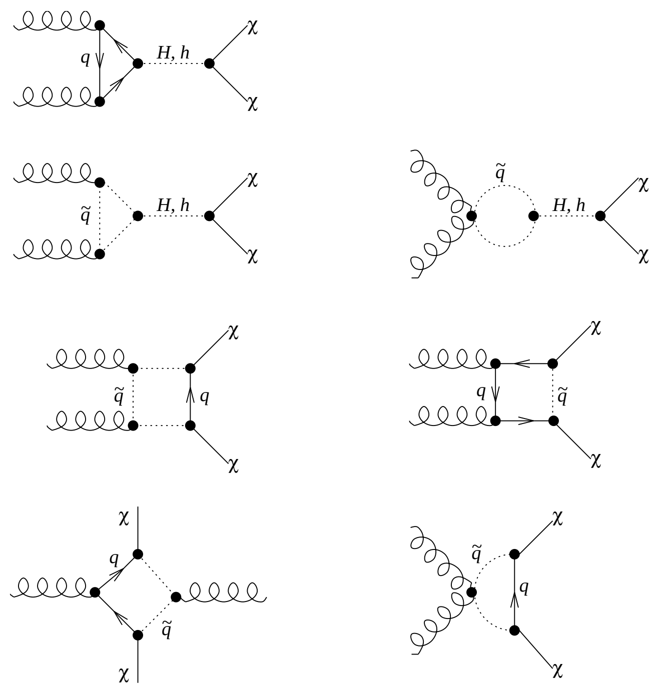

Figure 46: Feynman diagrams for neutralino-gluon scalar (spin-independent) elastic scattering. Notice that no tree level processes exist. From Ref. [305].

where Z and $A-Z$ are the numbers of protons and neutrons in the nucleus, respectively.

The above expression is valid only at zero momentum transfer between the WIMP and the nucleon. For finite momentum transfer, the differential cross section must be multiplied by a nuclear form factor. The appropriate factor, called the *Woods-Saxon* form factor, is given by [211]

$$F(Q) = \left(\frac{3j_1(qR_1)}{qR_1}\right)^2 \exp[-(\mathrm{qs})^2], \tag{202}$$

where j_1 is the first spherical bessel function and the momentum transferred is $q=\sqrt{sm_N Q}$. R_1 is given by $\sqrt{R^2-5s^2}$, where R and s are approximately equal to $1.2\,\mathrm{fm}\,\mathrm{A}^{1/3}$ and $1\,\mathrm{fm}$, respectively.

Although less accurate than the Woods-Saxon form factor, the following simple form factor is sometimes used in its place [16, 229]:

$$F(Q) = \exp[-\mathrm{Q}/2\mathrm{Q}_0]. \tag{203}$$

Here, Q is the energy tranferred from the WIMP to the target and $Q_0 = 1.5/(m_N R_0^2)$ where $R_0 = 10^{-13}\,\mathrm{cm}\,[0.3+0.91(\mathrm{m_N/GeV})^{1/3}]$.

In the context of neutralino scattering, the value of a_q can be calculated from the parameters of the MSSM [236, 434, 189, 188]. Following Ref. [200], a_q

is in this case given by:

$$
\begin{aligned}
a_q \;\; = \;\; & -\frac{1}{2(m_{1i}^2 - m_\chi^2)} Re\left[(X_i)(Y_i)^*\right] - \frac{1}{2(m_{2i}^2 - m_\chi^2)} Re\left[(W_i)(V_i)^*\right] \\
& - \frac{g m_q}{4 m_W B}\left[Re\left(\delta_1[gN_{12} - g' N_{11}]\right) DC\left(-\frac{1}{m_H^2} + \frac{1}{m_h^2}\right)\right. \\
& \left. + Re\left(\delta_2[gN_{12} - g' N_{11}]\right)\left(\frac{D^2}{m_H^2} + \frac{C^2}{m_h^2}\right)\right], \qquad (204)
\end{aligned}
$$

where

$$
\begin{aligned}
X_i \;\; & \equiv \;\; \eta_{11}^* \frac{g m_q N_{1,5-i}^*}{2 m_W B} - \eta_{12}^* e_i g' N_{11}^*, \\
Y_i \;\; & \equiv \;\; \eta_{11}^* \left(\frac{y_i}{2} g' N_{11} + g T_{3i} N_{12}\right) + \eta_{12}^* \frac{g m_q N_{1,5-i}}{2 m_W B}, \\
W_i \;\; & \equiv \;\; \eta_{21}^* \frac{g m_q N_{1,5-i}^*}{2 m_W B} - \eta_{22}^* e_i g' N_{11}^*, \\
V_i \;\; & \equiv \;\; \eta_{22}^* \frac{g m_q N_{1,5-i}}{2 m_W B} + \eta_{21}^* \left(\frac{y_i}{2} g' N_{11}, + g T_{3i} N_{12}\right) \qquad (205)
\end{aligned}
$$

where y_i and T_{3i} denote hypercharge and isospin, and

$$
\begin{gathered}
\delta_1 = N_{13}(N_{14}), \quad \delta_2 = N_{14}(-N_{13}), \\
B = \sin\beta(\cos\beta), \quad C = \sin\alpha(\cos\alpha), \\
D = \cos\alpha(-\sin\alpha). \qquad (206)
\end{gathered}
$$

Here, $i = 1$ for up-type and 2 for down-type quarks. α is the Higgs mixing angle. m_{1i}, m_{2i} denote elements of the appropriate 2 x 2 squark mass matrix. N_{1n} are elements of the matrix which diagonalizes the neutralino mass matrix in the $\widetilde{B}$-$\widetilde{W}^3$-$\widetilde{H}_1^0$-$\widetilde{H}_2^0$ basis (see Appendix A). η is the matrix which diagonalizes the appropriate squark mass matrices.

To crudely estimate what scale we expect for the scalar cross section between a neutralino and nucleon, we can carry out a back-of-the-envelope estimate. Considering a gaugino-like neutralino, we see that as δ_1 and δ_2 both vanish, so do most of the terms in Eq. 204. We are left with a neutralino-quark coupling of $a_q \sim A/m_{\tilde{q}}^2$, where A is the product of the various order 1 couplings, mixing matrix parameters, etc. which contribute. For a typical case, A might be $\sim 10^{-3}$ or so, although it can vary a great deal. Inserting this coupling into Eqs. 198 and 199, we estimate a neutralino-nucleon scalar cross section of $\sim A^2 m_p/m_{\tilde{q}}^4$, which is roughly 10^{-9} picobarns, for TeV mass squarks. These results can vary dramatically, however, depending on the characteristics of the model being considered (see Figs. 22 and 23).

We can contrast this with the much larger neutralino annihilation cross sections. Considering again a gaugino-like neutralino, its amplitude for annihilations into $b\bar{b}$ via psuedoscaler Higgs exchange (see Eq. 164) is roughly $\mathcal{A}_A \sim$

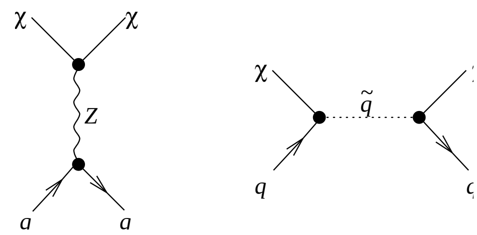

Figure 47: Tree level Feynman diagrams for neutralino-quark axial-vector (spin-dependent) elastic scattering. From Ref. [305].

$m_b \tan\beta\sqrt{f_h}/m_{W^\pm}$ where f_h is the higgsino fraction of the WIMP. The annihilation cross section (Eq. 179) is then roughly $\sigma \sim 3\, m_b^2 \tan^2\beta f_h/128\pi m_\chi^2 m_{W^\pm}^2$. For even a very small higgsino fraction, say 1%, and a 200 GeV neutralino, we find a cross section of $\sim 10^{-3}$ picobarns for small values of $\tan\beta$ and a few picobarns for $\tan\beta = 30$.

C.2 Axial-Vector Interactions

Next, we consider a WIMP with axial-vector interactions with quarks given by

$$\mathcal{L}_A = d_q\ \bar{\chi}\gamma^\mu\gamma_5\chi\ \ \bar{q}\gamma_\mu\gamma_5 q, \tag{207}$$

where d_q is the generic coupling.

For such a WIMP, the spin-dependent scattering cross section can be written as [247]

$$\frac{d\sigma}{d|\vec{v}|^2} = \frac{1}{2\pi v^2}\overline{|T(v^2)|^2}, \tag{208}$$

where v, again, is the relative velocity of the WIMP, and $T(v^2)$ is the scattering matrix element. This expression can be integrated over the Boltzman velocity distribution of halo WIMPs to arrive at an average elastic scattering cross section. At zero momentum, the matrix element, $T(v^2)$, is given by

$$\overline{|T(0)|^2} = \frac{4(J+1)}{J}|(d_u\Delta_u^p + d_d\Delta_d^p + d_s\Delta_s^p) < S_p > + (d_u\Delta_u^n + d_d\Delta_d^n + d_s\Delta_s^n) < S_n > |^2, \tag{209}$$

where J is the nuclear spin and the Δ's are the fraction of the nucleon spin carried by a given quark. Their values are measured to be $\Delta_u^p = \Delta_d^n = 0.78 \pm 0.02$, $\Delta_d^p = \Delta_u^n = -0.48 \pm 0.02$ and $\Delta_s^p = \Delta_s^n = -0.15 \pm 0.02$. $< S_p >$ and $< S_n >$ are the expectation values of the total spin of protons and neutrons, respectively. Notice that for target nuclei with even numbers of protons and neutrons, there is zero total spin, and the cross section vanishes.

The values of $< S_p >$ and $< S_n >$ depend on the nucleus being considered. For ^{73}Ge, the interacting shell model finds $< S_p >$ and $< S_n >$ to be 0.011 and 0.468, respectively. For ^{28}Si, they are given by -0.0019 and 0.133. For ^{27}A, they are 0.3430 and 0.269. And for ^{39}K, they are -0.184 and 0.054 [351].

For non-zero momenta, a more complex form of equation 209 is needed. This equation is given by

$$\begin{aligned}\overline{|T(v^2)|^2} = \frac{(J+1)}{J}|(d_u\Delta_u^p + d_d\Delta_d^p + d_s\Delta_s^p + d_u\Delta_u^n + d_d\Delta_d^n + d_s\Delta_s^n) \\ < S_p + S_n > F^0(v^2) + (d_u\Delta_u^p + d_d\Delta_d^p + d_s\Delta_s^p - d_u\Delta_u^n + d_d\Delta_d^n + d_s\Delta_s^n) \\ < S_p - S_n > F^1(v^2)|^2, \quad (210)\end{aligned}$$

where the F's are nuclear form factors given by

$$F^0(v^2) \simeq \exp(-r_0^2 v^2/\hbar^2) \tag{211}$$

and

$$F^1(v^2) \simeq \exp(-r_1^2 v^2/\hbar^2 + icv/\hbar), \tag{212}$$

where r_0 and r_1 are parameters which depend on the nucleus being considered, with typical values of $1.3 - 2.1\ \mathrm{fm}^{-1}$.

Again, within the context of neutralino scattering, the value of d_2 can be calculated from the parameters of the MSSM [254, 213, 262, 263, 202, 403]. Following Ref. [200], d_2 is in this case given by:

$$\begin{aligned} d_2 &= \frac{1}{4(m_{1i}^2 - m_\chi^2)}\left[|Y_i|^2 + |X_i|^2\right] + \frac{1}{4(m_{2i}^2 - m_\chi^2)}\left[|V_i|^2 + |W_i|^2\right] \\ &\quad - \frac{g^2}{4m_Z^2\cos^2\theta_W}\left[|N_{13}|^2 - |N_{14}|^2\right]\frac{T_{3i}}{2}, \end{aligned} \tag{213}$$

where the quantities used are defined in Appendix C.1.

C.3 Vector Interactions

As a third case, consider a WIMP with vector interactions with quarks, given by

$$\mathcal{L}_{vec}^q = b_q\ \bar{\chi}\gamma_\mu\chi\ \bar{q}\gamma^\mu q. \tag{214}$$

Here, b_q is the WIMP-quark vector coupling. In this case, the contributions of each quark in the nucleus add coherently and large cross sections result for large nuclei. The WIMP-nucleus cross section in this case is straight forward [254]

$$\sigma_{0\,vec} = \frac{m_\chi^2 m_N^2 b_N^2}{64\pi(m_\chi + m_N)^2}, \tag{215}$$

where b_N is simply $b_N = 2Zb_p + (A-Z)b_n$.

As an example of a WIMP with vector interactions, consider a Dirac neutrino. In this case, $b_q = G_F(T_q^3 - 2e_q\sin^2\theta_{\mathrm{W}})/\sqrt{2}$, where G_F is the Fermi constant, T_q^3 and e_q are the weak isospin and electric change of the quark q, respectively, and θ_{W} is the Weinberg angle. Summing over the quarks in a proton or neutron, we get $b_p = G_F(1 - 4\sin^2\theta_{\mathrm{W}})/(2\sqrt{2})$ and $b_n = -G_F/(2\sqrt{2})$. Since

$4\sin^2\theta_{\rm W} \cong 1$, the neutron-neutrino cross section is much larger than the analogous proton-neutrino interaction. The Dirac neutrino-neutron cross section is then given by $\sigma_{\nu,n} = G_F^2 m_\nu^2 m_n^2/(512\pi(m_\nu + m_n)^2)$. A cross section of this size has been ruled out by direct scattering experiments, except perhaps in the case of a very light WIMP. Similar calculations show that other similar particles, such as sneutrinos, are also excluded by this method [216].

Neutralinos, being Majorana fermions, do not have vector interactions with quarks.

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