Calculations of Rates for Direct Detection of Neutralino Dark Matter

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(Received 18 March 1988)

The detection rates in cryogenic detectors of neutralinos, the most well motivated supersymmetric dark-matter candidate, are calculated. These rates can differ greatly from the special case of pure photinos and pure Higgsinos which are usually considered. In addition, a new term is found in the elastic-scattering cross section proportional to the Z-ino component which is "spin independent," even for these Majorana particles. As a result, substantial detection rates exist for previously disfavored, mostly spinless materials such as germanium and mercury.

PACS numbers: 98.60.Ln, 14.80.Ly

It is quite possible that the dark matter (DM) known to exist in galactic halos consists of some, as yet undiscovered, elementary particle. In part because it is stable in most models, one of the most well motivated particle DM candidates is the lights supersymmetric particle (LSP). Of the candidates for LSP, the lightest neutralino ($\tilde{\chi}$), a linear combination of photino, Zino, and neutral Higgsinos, is probably the most likely. These were considered in detail by Ellis *et al.*,[1] who showed that over a wide range of parameters a relic density of neutralinos equal to critical density exists. Realizing that particles in the galaxy's halo might be detectable in laboratory experiments, many authors[2-4] have published predictions for event rates of neutralinos in cryogenic detectors. Most groups have, however, considered only the pure photino and pure Higgsino, two special cases of the more general neutralino. If the neutralino is very light then one might expect a reasonably pure photino or Higgsino; but, there are no strong theoretical or experimental reasons to expect such a light LSP and as the mass increases a pure photino or pure Higgsino becomes more unlikely.

In this Letter, I reconsider the elastic-scattering cross section for the general neutralino and apply the result to cryogenic-detection estimates. I find a relative sign difference with respect to Ref. 1 and also a new term which can be written as an additional scalar-scalar interaction in the effective Lagrangean. Applying a technique of Shifman, Vainshstein, and Zakharov,[5] I find that the new terms results in a piece of the elastic-scattering cross sections which is proportional to the mass of the nucleus. This differs from previous work which considered cross sections for Majorana particles to be "spin dependent" and means that neutralinos might be detectable even with mostly spinless materials such as germanium or mercury. The size of this scalar term is not large, but it does eliminate cancellations which occur otherwise and can dominate for heavy materials. Enhancements over the naive rate of several orders of magnitude are possible. This new term may also change the capture rate of neutralinos into the body of the Earth. Both the sign change and the new term do not contribute to pure photino or pure Higgsino elastic scattering. For clarity, in this Letter I make several simplifying assumptions; the full details will be reported elsewhere.[6]

Throughout we use the minimal supersymmetric extension of the standard model described by Haber and Kane[7] (see especially the appendices) and Gunion and Haber.[8] In these models there exists four neutralinos which are mixtures of the supersymmetric partners of the neutral W, the B, and two neutral Higgs bosons. These can also be characterized as the photino, Z-ino, and two neutral Higgsinos. Only the lightest will be stable (I assume a conserved R parity) and I denote it as $\tilde{\chi}=Z_{11}\tilde{B}+Z_{12}\tilde{W}^3+Z_{13}\tilde{H}_1+Z_{14}\tilde{H}_2$, where the Z_{1i} are the elements of the real orthogonal matrix which diagonalizes the neutralino mass matrix; that is, if $Z_{11}=Z_{12}=0$, $\tilde{\chi}$ is a pure Higgsino, if $Z_{11}=\cos\theta_W$, $Z_{12}=\sin\theta_W$, $\tilde{\chi}$ is a pure photino, and if $Z_{11}=-\sin\theta_W$, $Z_{12}=\cos\theta_W$, $\tilde{\chi}$ is a pure Z-ino.

The neutralino masses and the Z_{ij}'s (and also the chargino parameters) are fully determined by four parameters: $\tan\beta$, μ, M, and M', where $\tan\beta=v_2/v_1$ is the ratio of Higgs vacuum expectation values,[9] M and M' are soft supersymmetry breaking parameters, and μ is a supersymmetric Higgs mass. Throughout, we make the standard[7] simplification $M'=\frac{5}{3}M\tan^2\theta_W$ to reduce the parameter space. Overall then we have three undetermined parameters $\tan\beta$, M, and μ, and it is this parameter space that we explore.

For a neutralino of mass m_X less than the Z^0 mass m_Z, both elastic scattering ($\tilde{\chi}q\rightarrow\tilde{\chi}q$) and annihilation ($\tilde{\chi}\bar{\tilde{\chi}}\rightarrow q\bar{q}$) processes are found from the same five Feynman diagrams: one involving Z^0 exchange, two involving left-chiral-squark (or slepton) exchange, and two involving right-chiral-squark exchange. The complete matrix elements including different left- and right-chiral-squark masses and propagator momenta are quite cumbersome and will be presented elsewhere.[6] In the limit of heavy squarks (and heavy Z^0) the elastic-scattering cross section can be found, however, with an effective-Lagrangean technique

$$L_{\rm eff} = (-4g^2/M^2_{\tilde{q}L})\bar{\tilde{\chi}}(aP_R+bP_L)q\,\bar{q}(aP_L+bP_R)\tilde{\chi} - (4g^2/M^2_{\tilde{q}R})\bar{\tilde{\chi}}(cP_R-aP_L)q\,\bar{q}(cP_L-aP_R)\tilde{\chi} - (g^2/2m_W^2)(Z_{13}^2-Z_{14}^2)\bar{q}\gamma^\mu(c_LP_L+c_RP_R)q\,\bar{\tilde{\chi}}\gamma_\mu\gamma_5\tilde{\chi}, \quad (1)$$

where q is the quark field, g is the weak coupling constant, and $P_L = \frac{1}{2}(1-\gamma_5)$, etc., while $a = m_q d_q/2m_W$, $b = T_{3L}Z_{12} - \tan\theta_W(T_{3L}-e_q)Z_{11}$, and $c = \tan\theta_W e_q Z_{11}$. Here m_W is the W-boson mass, T_{3L} is the third component of the weak isospin, e_q is the charge of the quark or lepton, $\sin^2\theta_W = 0.23$, and $d_q = Z_{13}/\cos\beta$ for down-type quarks and $Z_{14}/\sin\beta$ for up-type quarks. Finally, I define $c_L = T_{3L} - e_q\sin^2\theta_W$ and $c_R = -e_q\sin^2\theta_W$.

To get Eq. (1) in a more useful form we perform Fierz transformations on the first two terms. For clarity, we set $M_{\tilde{q}L} = M_{\tilde{q}R}$ although this is not necessary. The important feature is that while $\bar{q}P_L\tilde{\chi}\bar{\tilde{\chi}}P_Rq = -\frac{1}{2}\bar{q}P_L\gamma^\mu q\,\bar{\tilde{\chi}}P_R\gamma_\mu\tilde{\chi}$, terms such as $\bar{q}P_L\tilde{\chi}\bar{\tilde{\chi}}P_Lq = -\frac{1}{2}\bar{q}P_Lq\,\bar{\tilde{\chi}}P_L\tilde{\chi}$ are not of the form of an axial-vector coupling. Using the fact[7] that for Majorana fermions $\bar{\tilde{\chi}}\gamma_\mu\tilde{\chi} = 0$ we find

$$L_{\rm eff} = (g^2/2m_W^2)[\bar{\tilde{\chi}}\gamma_\mu\gamma_5\tilde{\chi}\,\bar{q}\gamma^\mu(V'+A'\gamma_5)q + 2a(b-c)x_q^2(\bar{\tilde{\chi}}\tilde{\chi}\,\bar{q}q + \bar{\tilde{\chi}}\gamma_5\tilde{\chi}\,\bar{q}\gamma_5q)], \quad (2)$$

where $V' = -\frac{1}{2}(c_R+c_L)(Z_{13}^2-Z_{14}^2) + x_q^2(b^2-c^2)$, $A' = \frac{1}{2}(c_L-c_R)(Z_{13}^2-Z_{14}^2) - x_q^2(2a^2+b^2+c^2)$, and $x_q = m_W/M\tilde{q}$. Apart from the scalar and pseudoscalar interaction terms Eq. (2) differs from the corresponding equation in Ref. 1 only by the signs of a^2 and b^2. Since the complete calculation[6] involving the five diagrams gives the same relative signs I believe that the present signs are the correct ones. The complete calculation also shows that for elastic scattering (extreme nonrelativistic limit) the pseudoscalar term vanishes and also that there is no interference between the axial-vector and scalar pieces. Note that the scalar piece is proportional to $b-c$, the Z-ino component of the neutralino. The axial-vector piece of the elastic cross section can be evaluated as in Goodman and Witten[2] and I find the elastic-scattering cross section off a nucleus of mass m_N to be

$$\sigma_{\rm el} = \frac{24m_\chi^2 m_N^2 G_F^2}{\pi(m_\chi+m_N)^2}\left[\tfrac{4}{3}\lambda^2 J(J+1)\Big(\sum_{u,d,s}A'\Delta q\Big)^2 + \left(\frac{2m_N}{27m_W}\right)^2\Big(\sum_{c,b,t}(b-c)x_q^2 d_q\Big)^2\right], \quad (3)$$

where J is the total spin of the nucleus and the sums are over the indicated quarks. The first term agrees with Ref. 2 in the photino limit (see also Kane and Kani and Campbell *et al.*[10]) and the second term is new and requires some explanation. In the above, I followed Ref. 2 and Ellis and co-workers[11,12] in defining $\lambda = \frac{1}{2}\{1 + [s_p(s_p+1) - l(l+1)]/J(J+1)\}$ from the one-particle nuclear shell model[13] and the Landé formula, where l is the shell-model angular momentum and s_p is the proton (or neutron) spin. I also follow Refs. 11, 12, and Ashman *et al.*[14] in defining $\langle p|\bar{q}\gamma_\mu\gamma_5 q|p\rangle = 2\Delta q\,\mathbf{s}_q$, where $\mathbf{s}_q$ is the spin of quark q and Δg measures the fraction of the proton spin carried by quark q. Ashman *et al.* (European Muon Collaboration)[14] give $\Delta u = 0.746$, $\Delta d = -0.508$, and $\Delta s = -0.226$ and I use these values in the numerical work. [In Ref. 6 the values of Δq from flavor SU(3) are also considered.] Note that since we set $M_{\tilde{q}L} = M_{\tilde{q}R}$ no vector pieces can appear in Eq. (2). As discussed in Refs. 2 and 12 these terms can be important if there is significant left- and right-chiral-squark mixing; however, this is expected to be small.

In deriving the second term of Eq. (3) I modified slightly a technique described in Ref. 5 and used recently by Raby and West.[15] For coherent scattering of a neutralino off a nucleus we need to find

$$\langle N|\textstyle\sum_q 2a(b-c)x_q^2\bar{q}q|N\rangle \propto \langle N|\sum_q T_{3L}x_q^2 d_q m_q\bar{q}q|N\rangle,$$

where $\langle N|$ is the nucleus state and the sum is over all the quarks, both valence and sea. Using the "heavy-quark expansion" for the charm, bottom, and top quarks, Shifman, Vainshtein, and Zakharov write

$$m_q\bar{q}q \simeq -\tfrac{2}{3}(\alpha_s/8\pi)G^a_{\mu\nu}G^a_{\mu\nu} + O(\alpha_s^2/m_q^2)$$

and by including the anomaly in the trace of the quark energy-momentum tensor $\theta_{\mu\mu}$ they find

$$m_N\bar{\Psi}_N\Psi_N = \langle N|\theta_{\mu\mu}|N\rangle \simeq -(9\alpha_s/8\pi)\langle N|G^a_{\mu\nu}G^a_{\mu\nu}|N\rangle.$$

Physically, this last equation says that the mass of the nucleon (and therefore the nucleus) comes from the light-quark anomaly. Since the light quarks in the sum above are quite light I follow Ref. 5 in ignoring them and find

$$\langle N|\sum_q 4aT_{3L}x_q^2\bar{q}q|N\rangle \simeq \left(\frac{4m_N}{27m_W}\right)\sum_{c,b,t}T_{3L}x_q^2 d_q \simeq \frac{2m_N x_q^2}{27m_W}\left(\frac{2Z_{14}}{\sin\beta} - \frac{Z_{13}}{\cos\beta}\right), \quad (4)$$

where in the last step I made the simplifying assumption that all squarks have the same mass. Using Eq. (4), one finds Eq. (3) in a straightforward manner. The Higgsino coupling to the nucleon, like the Higgs-boson coupling, is via a loop diagram involving heavy quarks, in which the heavy-quark masses cancel out. I do not claim that the above cross section is exact, but is shows that "spin-independent" cross sections exist for Majorana particles. Uncertainties include the extent to which the charm, bottom, and top quarks contribute equally, the extent to which the strange quark contributes, the possibility of additional generations of quarks, and higher-order contributions, both in the heavy-quark and the heavy-squark expansions.

We now turn to numerical results. Displaying even rates in a cryogenic detector is problematic since there are many free parameters. Figures 1(a) and 1(b) show even rates versus the neutralino mass for one set of parameter values. These were all chosen so that $\Omega_X = 1$. Values of $\tan\beta$ of 2, $\frac{3}{4}$, and 0.2 and values of $M_{\tilde{q}}$ of 50, 100, and 200 GeV were chosen and then all values of $\mu > 0$ and M which satisfied $\Omega = 1$ found. Any values which resulted in $\tilde{\chi}$ being heavier than the chargino or squark were removed. In deriving the relic abundance of neutralinos I used the complete annihilation cross section[6] with a Hubble parameter of 50 km s^{-1} Mpc^{-1}. For simplicity, I set all the slepton and squark masses equal. The even rate is given by Ref. 4,

$$R = (8/3\pi)^{1/2} \sigma_{el} \rho_{halo} \bar{v} \eta_v \eta_c / m_N m_X,$$

where $\rho_{halo} \approx 0.4$ GeV cm^{-3} is the local halo density, $\bar{v} \approx 270$ km s^{-1} is the dispersion velocity in the halo, and η_v is a correction for the motion of the Earth through the halo. Considering only the circular velocity of the Sun around the galactic center $v_\odot \approx 220$ km s^{-1}, I find $\eta_v \approx 1.3$. The factor η_c is a correction for loss of coherence at large m_X and large m_N which goes to 1 as either mass gets small.[4] For mercury $\eta_c \simeq 0.8$ for $m_X \simeq 20$ GeV, 0.5 for $m_X \simeq 40$, dropping to 0.2 for $m_X \simeq 90$ GeV. Figure 1(a) shows the total event rate, while Fig. 1(b) shows the result of leaving out the new scalar-scalar term. The rates in Fig. 1(b) are smaller overall since $\lambda^2 J(J+1) = \frac{1}{12}$ for the 17% of mercury which is not spinless.

The first thing to notice is the rather large variation in event rate which comes from considering the general neutralino rather than just the photino. The almost pure photino is seen as the two dark blobs (corresponding to $M_{\tilde{q}} = 50$ and $M_{\tilde{q}} = 100$ GeV) in both Figs. 1(a) and 1(b). Figure 1(b) shows very low event rates for $m_X \approx 5$ GeV and $m_X \approx 17$ GeV, which result from cancellations among the $A'\Delta q$'s due to negative Z^0-squark exchange interference. As $M_{\tilde{q}}$ and $\tan\beta$ are varied, these cancellations actually occur for every value of m_X. There are also low rates at $m_X \approx m_Z/2$ due to the Z^0 pole.[16] The cancellations are not as pronounced in Fig. 1(a) since the

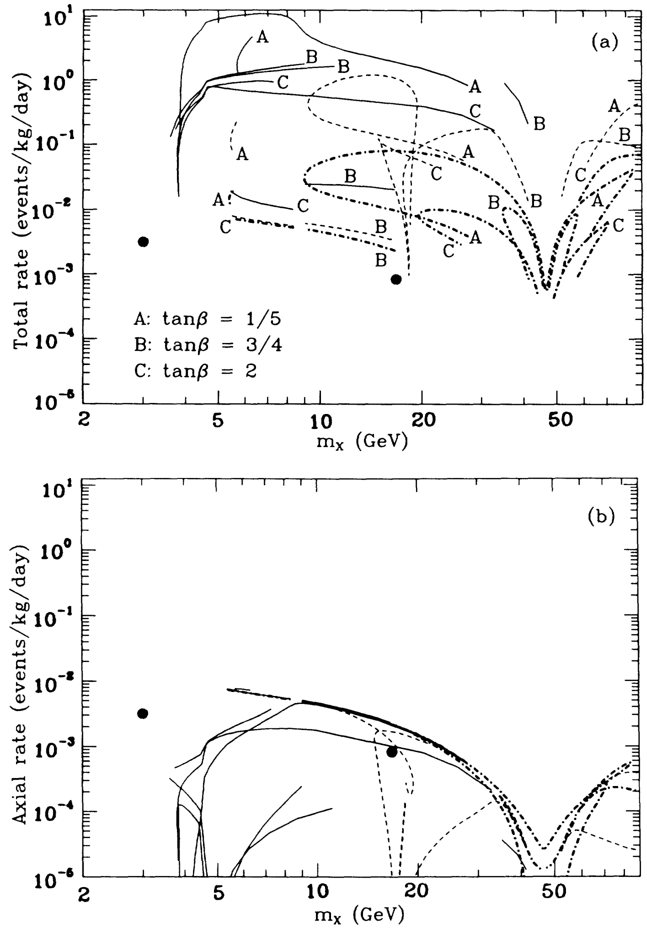

FIG. 1. Event rates in a mercury detector (natural abundance) for values of the model parameters chosen so that $\Omega = 1$. Solid lines indicate $M_{\tilde{q}} = 50$ GeV, dashed lines $M_{\tilde{q}} = 100$ GeV, and dot-dashed lines $M_{\tilde{q}} = 200$ GeV. Values of $\tan\beta$ of 2, 0.75, and 0.2 are included. (a) Total rate. (b) The rate including only the axial-vector term.

scalar term gives a minimum cross section, but the Z^0 pole suppression remains. The almost-pure-photino blobs do not move from Fig. 1(a) to Fig. 1(b) showing that the new term does not contribute to pure-photino scattering. Finally, these and examples of event rates for germanium and fluorine are shown in Table I.

The results presented here are illustrative only. Mercury, lead, or fluorine may not be ideal elements for DM detection. It is also likely that sleptons are lighter than squarks which will reduce event rates. This adds to the parameter space which needs to be explored. In addition, areas of parameter space can be eliminated by requiring consistency with experimental results. For example, Albajar *et al.* (the UA1 collaboration)[17] claims limits on square masses, experiments at the DESY e^+e^- collider PETRA[18] have placed limits on the chargino masses, and Hearty *et al.* and Bartha *et al.* (the

TABLE I. Total even rates (R_{tot}) and event rates without the scalar term (R_{ax}) of neutralinos scattering off various elements (in events kg^{-1} d^{-1}). Also shown are the model parameters and the photino ($\tilde{\gamma}$) and Z-ino ($\tilde{Z}$) components of the neutralino. The first three entries show the range possible at $m_X \approx 5$, and the second three at $m_X \approx 17$. All masses are in gigaelectronvolts.

m_X	M	μ	$\tan\beta$	$\tilde{M}$	$\tilde{\gamma}$	$\tilde{Z}$	R_{tot} (Hg)	R_{ax} (Hg)	R_{tot} (Ge)	R_{ax} (Ge)	R_{tot} (F)	R_{ax} (F)
4.9	53	100	0.2	50	0.28	−0.70	8.88	6×10^{-4}	3.0	3×10^{-3}	1.1	0.55
4.9	139	71	2	50	0.21	−0.57	0.78	1×10^{-3}	0.27	6×10^{-3}	0.99	0.94
5.0	90	100	0.75	50	0.23	−0.73	1.07	2×10^{-7}	0.36	8×10^{-7}	0.07	1×10^{-3}
17	67	132	0.2	100	0.42	−0.72	1.03	1×10^{-3}	0.32	4×10^{-3}	0.52	0.48
17	51	376	2	100	0.60	−0.78	0.01	6×10^{-7}	5×10^{-3}	2×10^{-6}	5×10^{-4}	2×10^{-4}
17	28	43	2	100	1.0	0.02	8×10^{-4}	8×10^{-4}	3×10^{-3}	3×10^{-3}	0.36	0.36

ASP collaboration)[19] have put limits on the process $e^+e^- \rightarrow \bar{\tilde{\chi}}\tilde{\chi}\gamma$. A more complete exploration is in progress and will be presented elsewhere.[6] Finally, I note that some results presented in this Letter, and some of the above issues have been considered recently by Ellis and Flores.[12]

In conclusion, it is seen that a general neutralino can give event rates rather different from those of the usually considered photino and Higgsino. Since the LSP is probably the most well motivated particle DM candidate and there is no strong reason to expect other than a combination state, experiments should aim for neutralino rather than photino or Higgsino detection. In this case, new terms in the cross sections can be important and should be included.

I would like to acknowledge very useful conversations with P. Arnold, W. Bardeen, H. Haber, D. Seckel, M. Turner, and especially R. Flores. This work was supported in part by the U.S. Department of Energy (at Chicago) and for computing by the U.S. Department of Energy and NASA (at Fermilab).

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