Fermi National Accelerator Laboratory

FERMILAB-Pub-88/52-A
April 1988

Calculations of Rates for Direct Detection of Neutralino Dark Matter

KIM GRIEST

Astronomy and Astrophysics Center, Enrico Fermi Institute, The University of Chicago, Chicago, Illinois 60637

and

NASA/Fermilab Astrophysics Center, Fermi National Accelerator Laboratory, Batavia, Illinois 60510

ABSTRACT

We calculate the detection rates in cryogenic detectors of neutralinos, the most well motivated supersymmetric Dark Matter candidate. These rates can differ greatly from the special cases of pure photinos and pure higgsinos which are usually considered. In addition, we find a new term in the elastic scattering cross section proportional to the zino 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.

Operated by Universities Research Association Inc. under contract with the United States Department of Energy

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 lightest supersymmetric particle (LSP). Of the candidates for LSP the lightest neutralino ($\widetilde{\chi}$), 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 our 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, we reconsider the elastic scattering cross section for the general neutralino and apply the result to cryogenic detection estimates. We 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 Lagrangian. Applying a technique of Shifman, *et al.*,[5] we find that the new term 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 we make several simplifying assumptions; the full details will be reported elsewhere.[6]

Throughout we use the minimal supersymmetric extension of the Standard Model described in Ref. 7 (see especially the appendices) and Ref. 8. In these models there exist 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, zino and two neutral higgsinos. Only the lightest will be stable (we assume a conserved R parity) and we denote it as $\widetilde{\chi} = Z_{11}\widetilde{B} + Z_{12}\widetilde{W}^3 + Z_{13}\widetilde{H}_1 + Z_{14}\widetilde{H}_2$, where the Z_{1i} are the elements of the real orthogonal matrix which diagonalizes the neutralino mass matrix; that is, if $Z_{11} = Z_{12} = 0$, $\widetilde{\chi}$ is a pure higgsino, if $Z_{11} = \cos\theta_w, Z_{12} = \sin\theta_w$, $\widetilde{\chi}$ is a pure photino, and if $Z_{11} = -\sin\theta_w, Z_{12} = \cos\theta_w$, $\widetilde{\chi}$ is a pure zino.

The neutralino masses, 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] and the rest are soft supersymmetry breaking parameters. Throughout, we make the standard[7] simplification $M' = \frac{5}{3}M\tan\theta_w$ to reduce the parameter space. Overall then we have three undetermined parameters $\tan\beta$, M, and μ, and it is this parameter space we explore.

For a neutralino of mass m_χ less than the Z^0 mass m_Z, both elastic scattering ($\widetilde{\chi}q \to \widetilde{\chi}q$) and annihilation ($\widetilde{\chi}\bar{\widetilde{\chi}} \to q\bar{q}$) 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, using an effective Lagrangian technique

$$\begin{aligned}L_{eff} = &\frac{-4g^2}{M^2_{\tilde{q}L}}\bar{\widetilde{\chi}}(aP_R + bP_L)q\bar{q}(aP_L + bP_R)\widetilde{\chi} \\ &- \frac{4g^2}{M^2_{\tilde{q}R}}\bar{\widetilde{\chi}}(cP_R - aP_L)q\bar{q}(cP_L - aP_R)\widetilde{\chi} \\ &- \frac{g^2}{2m^2_W}(Z^2_{13} - Z^2_{14})\bar{q}\gamma^\mu(c_L P_L + c_R P_R)q\bar{\widetilde{\chi}}\gamma_\mu\gamma_5\widetilde{\chi},\end{aligned} \tag{1}$$

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, $\sin^2\theta_w = .23$ and $d_f = Z_{13}/\cos\beta$ for down type quarks and $Z_{14}/\sin\beta$ for up types. Finally, we 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 $\bar{q}P_L\widetilde{\chi}\bar{\widetilde{\chi}}P_R q = -\frac{1}{2}\bar{q}P_L\gamma^\mu q\bar{\widetilde{\chi}}P_R\gamma_\mu\widetilde{\chi}$ terms such as $\bar{q}P_L\widetilde{\chi}\bar{\widetilde{\chi}}P_L q = -\frac{1}{2}\bar{q}P_L q\bar{\widetilde{\chi}}P_L\widetilde{\chi}$ 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

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

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 we believe our signs are the correct one. The complete calculation also shows that for elastic scattering (extreme non-relativistic 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 zino component of the neutralino. The axial vector piece of the elastic cross section can be evaluated as in Goodman and Witten[2] and we find the elastic scattering cross section off a nucleus of mass m_N to be

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

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 Ref. 10) and the second term is new and requires some explanation. In the above, we followed Goodman and Witten,[2] the EMC group[11] and Refs. 12 and 13 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[14] and the Lande formula, where l is the shell model angular momentum and s_p is the proton (or neutron) spin. We also follow Refs. 11, 12 and 13 in defining $\langle p|\bar{q}\gamma_\mu\gamma_5 q|p\rangle = 2\Delta q\vec{s}_q$, where $\vec{s}_q$ is the spin of quark q and Δq measures the fraction of the proton spin carried by quark q. The EMC group[11] gives $\Delta u = .746$, $\Delta d = -.508$ and $\Delta s = -.226$ and we use these values in our numerical work. (In Ref. 6 the values of Δq from flavor SU(3) are also considered.) Note that for simplicity we left out vector pieces in Eq. (2) which can be important if there is significant left and right chiral squark mixing. These terms have been discussed in Refs. 2 and 13 and are expected to be small.

In deriving the second term of Eq. (3) we modified slightly a technique described in Shifman, *et al.*[5] and used recently by Raby and West.[15] For coherent scattering of a neutralino off a nucleus we need to find $\langle N|\sum_q 2a(b-c)x_q^2\bar{q}q|N\rangle \propto \langle N|\sum_q T_{3L}x_q^2 d_q m_q\bar{q}q|N\rangle$, 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, *et al.* write $m_q\bar{q}q \simeq -\frac{2}{3}\frac{\alpha_s}{8\pi}G^a_{\mu\nu}G^a_{\mu\nu} + O\left(\frac{\alpha_s^2}{m_q^2}\right)$ and by including the anomaly in the trace of the quark energy-momentum tensor $\theta_{\mu\mu}$ they find $m_N\bar{\Psi}_N\Psi_N = \langle N|\theta_{\mu\mu}|N\rangle \simeq -\frac{9\alpha_s}{8\pi}\left\langle N|G^a_{\mu\nu}G^a_{\mu\nu}|N\right\rangle$. 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 very light we can ignore them and find

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

where in the last step we 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 coupling, is via a loop diagram involving heavy quarks, in which the heavy quark masses cancels out. We do not claim that the above cross section is exact, but it 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 possibility of additional generations of quarks, and higher order contributions, both in the heavy quark and the heavy squark expansion.

We now turn to numerical results. Displaying event rates in a cryogenic detector is problematic since there are many free parameters. Figs. (1a) and (1b) show event 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, 3/4, and .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 we used the complete annihilation cross section[6] with a Hubble parameter of 50 km s^{-1} Mpc^{-1}. For simplicity, we set all the slepton and squark masses equal. The event rate is given by Ref. 4, $R = (8/3)^{\frac{1}{2}}\sigma_{el}\rho_{halo}\bar{v}\eta_v\eta_c/(\pi m_N m_X)$, where $\rho_{halo} \approx .4$ 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} we 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 one as either mass gets small.[4] Actually, in the one particle nuclear shell model coherence loss is not important, so we have not corrected the "spin dependent" piece, but we have applied this correction to the scalar piece. For Mercury $\eta_c \simeq .8$ for $m_X \simeq 20$

GeV, .5 for $m_X \simeq 40$, dropping to .2 for $m_X \simeq 90$ GeV. Fig. (1a) shows the total event rate while Fig. (1b) shows the result of leaving out the new scalar-scalar term. The rates in Fig. (1b) are smaller overall since $\lambda^2 J(J+1) = 1/12$ 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. (1a) and (1b). Fig. (1b) 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. (1a) since the scalar term gives a minimum cross section, but the Z^0 pole suppression remains. The almost pure photino blobs do not move from Fig. (1a) to Fig. (1b) 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 the table.

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, UA(1) claims limits on squark masses, PETRA has placed limits on the chargino masses, and ASP has 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, we note that some results presented in this Letter, and some of the above issues have been considered recently by Ellis and Flores.[13]

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.

We 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 DoE (at Chicago) and for computing by the DoE and NASA (at Fermilab).

REFERENCES

1. J. Ellis, J. S. Hagelin, D. V. Nanopoulos, K. A. Olive and M. Srednicki, *Nucl. Phys.* **B238** (1984) 453.
2. M. W. Goodman and E. Witten, *Phys. Rev* **D31** (1986) 3059.
3. A. K. Drukier, K. Freese and D. N. Spergel, *Phys. Rev.* **D33** (1986) 3495.
4. K. Freese, J. Frieman and A. Gould, SLAC Report No. SLAC-PUB-4427, 1987 (to be published).
5. M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, *Phys. Lett.* **78B** (1978) 443; A. I. Vainshtein, V. I. Zakharov and M. A. Shifman, *Sov. Phys. Usp.* **23** (1980) 429.
6. K. Griest, in preparation, to be submitted to *Phys. Rev.* **D** (1988).

7. H. E. Haber and G. L. Kane, *Phys. Rep.* **117** (1985) 75.

8. J. F. Gunion and H. E. Haber, *Nucl. Phys.* **B272** (1986) 1.

9. We follow the convention in which H_1 gives mass to the down type quarks and H_2 gives mass to the up type. This is the opposite of the Ellis, *et al.*[1] convention, implying $Z_{13} \leftrightarrow Z_{14}$ and $\tan\beta \to v_1/v_2$. Also, our μ corresponds to their $-\epsilon$.

10. G. L. Kane and I. Kani, *Nucl. Phys.* **B277** (1986) 525; B. A. Campbell, *et al.*, *Phys. Lett.* **173B** (1986) 270.

11. J. Ashman, *et al.*, CERN Report No. CERN-EP/87-230, 1987 (to be published).

12. J. Ellis, R. A. Flores and S. Ritz, *Phys. Lett.* **198B** (1987) 393.

13. J. Ellis and R. A. Flores, CERN Report No. CERN-TH.4911/87 (Draft), 1988 (to be published).

14. A. de-Shalit and I. Talmi, *Nuclear Shell Theory* (Academic Press, New York, 1963).

15. S. Raby and G. B. West, Los Alamos Report No. LA-UR-87-3664 1987 (to be published).

16. Annihilation due to Z^0 exchange is very efficient near $m_X = m_Z/2$, implying that weak couplings are needed for $\Omega_X = 1$. These weak touplings remain in the elastic scattering cross section, although the Z^0 pole is no longer there, causing the event rates to be small.

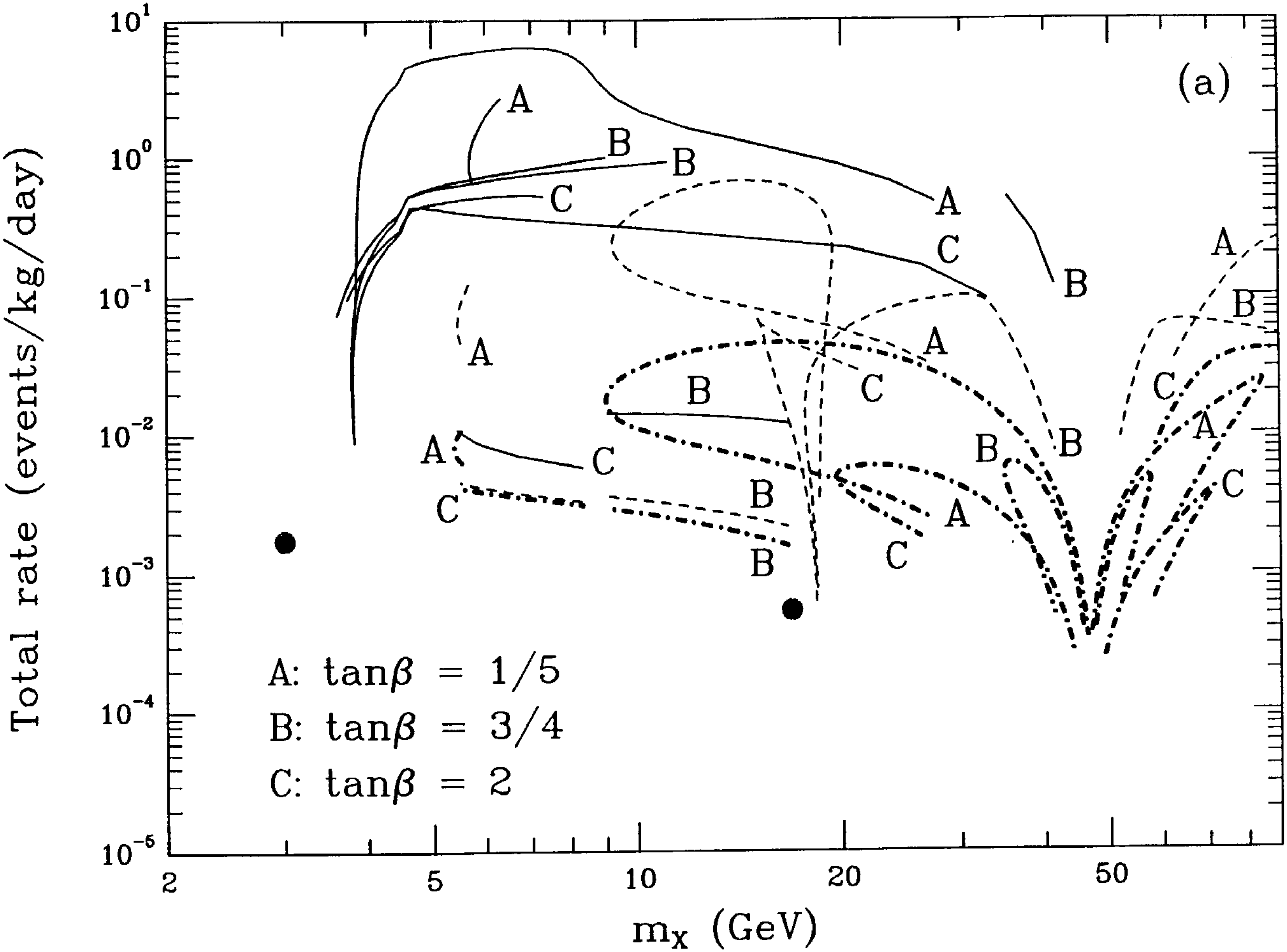
TABLE I

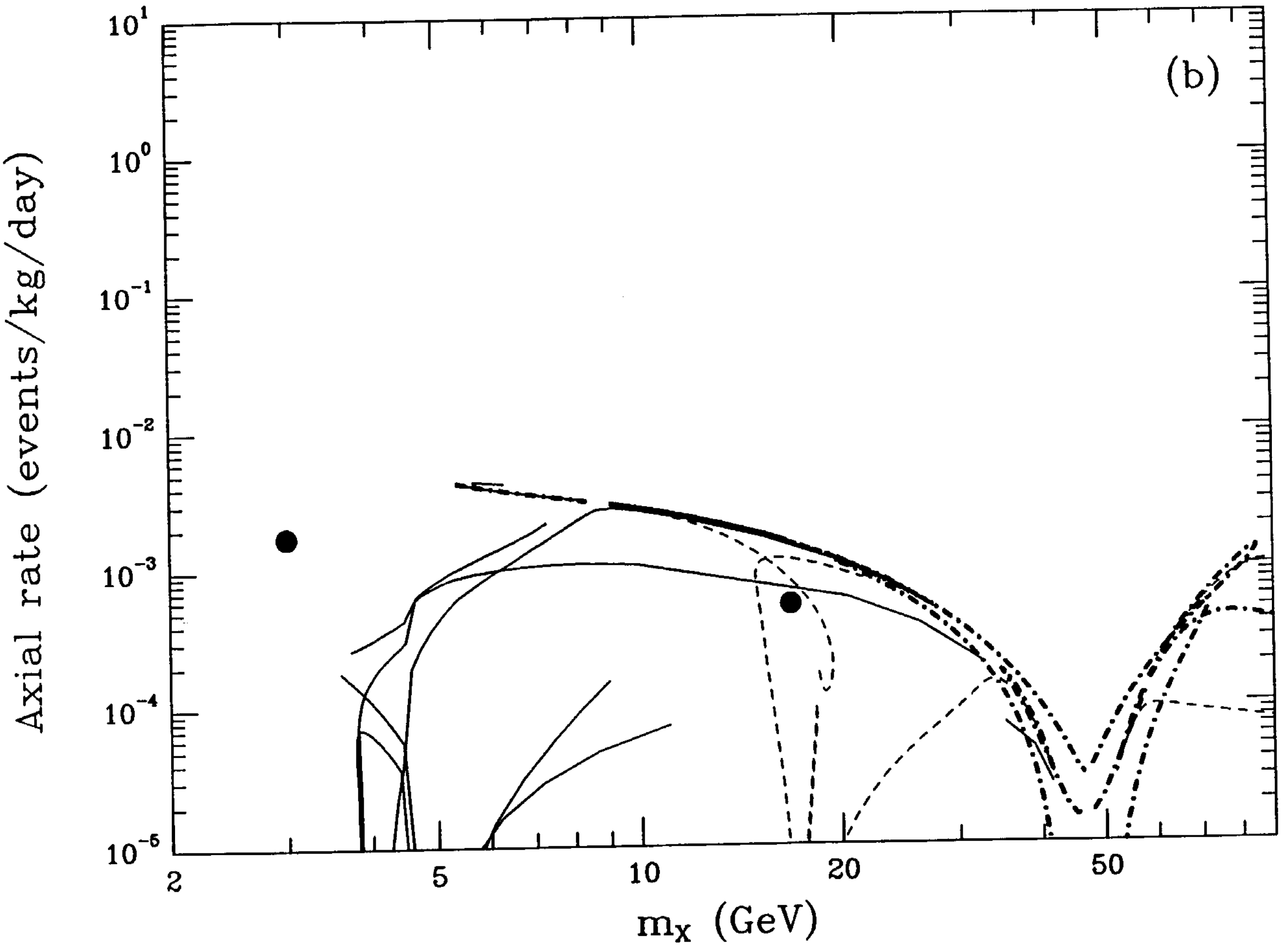
Total event rates (R_{tot}) and event rates without the scalar term (R_{ax}) of neutralinos scattering off various elements (in events kg^{-1} day^{-1}). Also shown are the model parameters and the photino ($\widetilde{\gamma}$) and zino ($\widetilde{Z}$) components of the neutralino. The first three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

m_X	M	μ	$\tan\beta$	$\widetilde{M}$	$\widetilde{\gamma}$	$\widetilde{Z}$	R_{tot} (Hg)	R_{ax} (Hg)	R_{tot} (Ge)	R_{ax} (Ge)	R_{tot} (F)	R_{ax} (F)
4.9	53	100	.2	50	.28	-.70	5.00	4E-4	1.7	2E-3	.62	.31
4.9	139	71	2	50	.21	-.57	.44	8E-4	.15	3E-3	.56	.53
5.0	90	100	.75	50	.23	-.73	.61	1E-7	.20	5E-7	.04	8E-4
17	67	132	.2	100	.42	-.72	.58	7E-4	.18	3E-3	.30	.28
17	51	376	2	100	.60	-.78	8E-3	4E-7	3E-3	1E-6	3E-4	1E-4
17	28	43	2	100	1.0	.02	5E-4	5E-4	2E-3	2E-3	.20	.20

Figure Caption

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, .75, and .2 are included. Fig. 1a is the total rate, while Fig. 1b is the rate including only the axial vector term.


(a)
A: tanβ = 1/5
B: tanβ = 3/4
C: tanβ = 2
Total rate (events/kg/day)
m_X (GeV)



(b)
Axial rate (events/kg/day)
10^{1}
10^{0}
10^{-1}
10^{-2}
10^{-3}
10^{-4}
10^{-5}
2
5
10
20
50
m_X (GeV)