Prospects for Neutralino Detection with a $^{73}\text{Ge} + ^{76}\text{Ge}$ Detector

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Abstract

We survey the elastic scattering rates of neutralinos in a proposed $500\text{g}^{73}\text{Ge} + 500\text{g}^{76}\text{Ge}$ cryogenic detector, choosing model parameters that give a critical density of the lightest neutralino and that are currently allowed by all LEP and other accelerator constraints. We work with the full one-loop radiatively-corrected Higgs mass matrix and use grand unification relations among sfermion and gaugino masses, and we use the odd-group model to estimate nuclear matrix elements incorporating a simple model for the momentum dependence of the elastic spin and mass form factors. We find that the rates are above the $0.1 (\text{kg}\cdot\text{day})^{-1}$ level of sensitivity of such a detector over a significant fraction of the parameter space for low values of m_A , the mass of the pseudoscalar Higgs boson.

Astronomical observations have indicated for some time that the bulk of the matter in the universe is invisible. Some of the most complete and convincing evidence comes from the study of the rotation of spiral galaxies, including that of our own galaxy. Whilst the dark matter in galactic halos could conceivably be baryonic, there is increasing evidence that considerably more Dark Matter is present on larger distance scales, quite possibly sufficient to match the critical density required to close the universe: $\Omega \equiv \rho/\rho_{crit} = 1$. This is the density expected on the basis of naturalness arguments, and specifically predicted by inflationary models of cosmology. Theories of galaxy formation also favour a high-density universe, which is moreover very consistent with the recent COBE observations of structure in the microwave background radiation [1]

Constraints from the cosmological synthesis of light elements exclude the possibility that so much Dark Matter could be baryonic [2], and it is natural to hypothesize that the dark matter in galactic halos is also non-baryonic. There is a plethora of candidates for this Dark Matter, ranging in mass from macroscopic black holes down to elementary particles weighing down to a few eV in the case of neutrinos, or even less in the case of axions. Theories of structure formation seem to favor Cold Dark Matter candidates, such as axions or particles weighing more than 1 keV. Detection strategies have been devised for many of these candidates if they do indeed constitute our galactic halo, which we will not review here [3]. In particular, weakly-interacting Cold Dark Matter particles weighing a GeV or more can in principle be detected via their elastic scattering off massive nuclei.

The elastic scattering rates depend not only on the mass and coupling strength of the conjectured Dark Matter particle, but also on its spin and the spin of

the target nucleus. For example, spin-zero or Dirac spin-half particles scatter coherently off large spinless nuclei, whilst Majorana spin-half particles prefer to scatter off nuclei with spin, although they do also have some coherent interactions with spinless nuclei. Thus the category of Cold Dark Matter particle that can be detected depends on the species of target nucleus chosen. This choice is affected by practical and technological considerations, and the nuclei used in experiments so far (*e.g.* ^{76}Ge and ^{29}Si) have not been able to yield very stringent constraints on Majorana Dark Matter candidates such as the lightest supersymmetric particle (LSP), which is expected to be some combination of electromagnetically-neutral and non-strongly-interacting gauginos and higgsinos, commonly called neutralinos.

A potentially important step forward in the search for Cold Dark Matter particles is now being taken with the preparation of an experiment using 500g of Germanium 73 and 500 g of Germanium 76, a mixture of nuclei with and without spin [4]. In this paper we present predictions for the scattering rate of the LSP in this experiment, and discuss the prospects that it could either detect or exclude this Cold Dark Matter candidate in at least a domain of parameter space.

We use the full gaugino/higgsino composition of the LSP in the minimal supersymmetric extension of the Standard Model (MSSM), and incorporate all the experimental constraints on its parameter space coming from LEP and other accelerators. We choose the masses of the supersymmetric partners of the quarks and leptons, the sfermions, in such a way that the relic density of the LSP is critical, $\Omega_{LSP} = 1$. We incorporate the supersymmetric Grand Unified Theory (GUT) relations between the sfermion masses, and also parameters of the Higgs

mass matrix. We work with the full one-top-quark-loop radiatively-corrected MSSM Higgs mass matrix and mixing angle. To estimate the nucleonic matrix elements of the effective four-fermion LSP-quark interaction operator, we use the latest EMC estimates for the spin-dependent parts, and the standard sigma-term estimates for the spin-independent parts. To estimate the nuclear matrix elements, we use the odd group model, which has been shown in the past to agree well with more detailed shell model calculations wherever they are available. Finally, we model the mass and spin form factors in the manner we discussed in ref. [5].

We find that there are significant domains of the LSP parameter space where the planned Germanium experiment has a chance of detecting or excluding the LSP as the main constituent of the galactic halo. Only parts of these domains could be explored in parallel by the continuation of direct searches at LEP II. The extents of these domains depend significantly on the values assumed for the ratio $\tan\beta$ of Higgs vacuum expectation values in the MSSM, and on m_A , the mass of the MSSM's CP-odd neutral Higgs boson. We find in particular that a larger fraction of the domain of LSP parameter space where $\Omega_{LSP} = 1$ is attainable can be explored if m_A is small, although this is not a large fraction of the total LSP parameter space allowed by present experimental constraints.

There are four neutralinos in the MSSM that are mixtures of the standard gauginos (W_3, B) and Higgsinos (H_1, H_2). Their well-known mixing matrix [6] is characterized by the ratio of the hypercharge +1 to the hypercharge -1 Higgs vacuum expectation values $v_1/v_2 \equiv \tan\beta$, a superpotential $H_1 H_2$ mixing coefficient μ , and SU(2) and U(1) gaugino masses M_2 and M_1 respectively, which we assume as usual to be related by grand unification to a common gaugino mass,

$m_{1/2} : M_i = (\alpha_i / \alpha_G \Gamma) m_{1/2}$ ($i=1,2$), so that $M_1 = \frac{5}{3} \tan^2 \theta_W M_2$. The masses of other particles (squarks, sleptons and Higgses) also enter into the elastic scattering amplitudes and the calculation of the relic LSP abundance. For the squarks and sleptons we assume conventional grand unification relations (neglecting mixing):

$$m_i^2 = m_0^2(1 + C_i \xi^2) + m_{1/2}^2 \pm \frac{1}{2}(v_1^2 - v_2^2) \left\{ \frac{1}{2} \xi_i g_1^2 - T_{3i} g_2^2 \right\}, \quad (1)$$

where the + (-) sign is for $\tilde{\nu}_L$ ($\tilde{\nu}_R$), the scalar partner of the left- (right-) handed component of the fermion f , $\xi \equiv m_{1/2}/m_0$ and [7] $C_{\tilde{q}L} = C_{\tilde{u}R} = 10C_{\tilde{t}L} \approx 30C_{\tilde{t}R} \approx 5$. Sfermion mixing is potentially very important for the partners of the lightest quarks, if present, because it can lead to spin-independent interactions that are present even if the LSP is a pure gaugino [8], and it also affects the Higgs spectrum and mixing if present for the partners of the heaviest quarks. Here we shall include sfermion mixing for the stop and sbottom, characterized by μ and soft-supersymmetry-breaking parameters A_t and A_b respectively. We have not included sfermion mixing for the partners of the lightest quarks, as we find it to have a negligible effect on the elastic scattering rates in the parameter space we survey (see the discussion below). An extra parameter is required to specify the Higgs spectrum and mixing angle α_H at the one-loop level, in addition to the top quark mass m_t : the mass of the pseudoscalar Higgs boson, m_A . The one-loop Higgs mass matrix [9] is then specified by $\tan\beta$, m_A , μ , A_t , A_b , and the top quark mass m_t . Our results are not very sensitive to the choice of parameters A_t and A_b . Current constraints [10] place the value of m_t in the range 90–200 GeV, and we choose $m_t = 140$ GeV as a representative value.

The principal experimental constraints on the parameters μ , M_2 and $\tan\beta$ entering the neutralino mass matrix are: (1) $M_{\text{chargino}} > 45 \text{ GeV}$ (LEP) [11]; (2) an upper limit on non-standard contributions to the Z decay width [12], which we implement assuming that the processes $Z \rightarrow \chi\chi'$ and $Z \rightarrow \chi'\chi'$ contribute to the hadronic decay width, while only $Z \rightarrow \chi\chi$ contributes to the invisible width (here χ' is a generic heavier neutralino, and we note that the resulting bounds are independent of any assumption on the chargino and neutralino decay modes); (3) an upper limit on the branching ratio for $Z \rightarrow \chi\chi'$ and $Z \rightarrow \chi'\chi'$ decays, assuming that $\chi' \rightarrow \chi\gamma$ predominantly [13]; (4) a lower bound $M_{\text{gluino}} > 70 \text{ GeV}$ (CDF) [14]. In Figure 1 we show the envelope of these constraints in the (μ, M_2) plane for two representative choices $\tan\beta = 2$ and δ . There are in addition constraints on the parameters $\tan\beta$ and m_A that enter the one-loop scalar Higgs mass matrix set by the unsuccessful searches at LEP for the pseudoscalar (A) and lightest scalar (h) Higgs bosons in the decays $Z \rightarrow Z^*h$ and $Z \rightarrow hA$ [15]. We implement these constraints, which depend on the sfermion masses and μ via α_H , by finding the sfermion mass scale m_o that gives $\Omega_{LSP} = 1$, and then finding whether α_H is allowed by the LEP constraints.

We assume, as usual, that most of the dark matter halo surrounding our galaxy is made of elementary particle dark matter, so that the typical dark matter particle velocities are $\approx 300 \text{ km}\cdot\text{s}^{-1}$. At these very small velocities the energy transfer involved is very small, $\Delta E < M_{\chi^2}^2 = 10 \text{ keV}(M_{\chi}/10\text{GeV})$, which is much smaller than any elementary particle mass scale in the problem. We can therefore use an effective Lagrangian to describe low-energy LSP-quark interactions. Neglecting kinetic energy terms:

$$L_{\text{eff}} = \bar{\chi}\gamma_\mu\gamma_5\chi\bar{q}\gamma^\mu(AP_L + BP_R)q + \bar{\chi}\chi C m_q\bar{q}q, \quad (2)$$

where $P_{RL} \equiv \frac{1}{2}(1 \pm \gamma_5)$ and χ is the LSP mixture of the gauginos and higgsinos, $\chi = \alpha W_3 + \beta B + \gamma H_1 + \delta H_2$. The first term causes a nuclear-spin-dependent interaction and is due to sfermion and Z boson exchange. The second term contributes a spin-independent term to the cross-section that is proportional to the square of the nuclear mass, and is due to Higgs [16] and sfermion [17] exchange as well as sfermion mixing [8,18]. The coefficients A, B and C are given by (neglecting sfermion mixing)

$$A = -\frac{(\gamma^2 - \delta^2)}{4M_Z^2}(g_I\sin\theta_W + g_2\cos\theta_W)\left(\frac{\gamma_L}{2}g_I\sin\theta_W - T_3g_2\cos\theta_W\right) + \frac{(\alpha g_2T_3 + \beta g_I\frac{Y_L}{2})^2}{2(m_L^2 - M_Z^2)} + \frac{m_q^2}{4(m_R^2 - M_Z^2)} \begin{cases} \gamma^2/v_I^2; T_3 = +\frac{1}{2} \\ \delta^2/v_2^2; T_3 = -\frac{1}{2} \end{cases} \quad (2a)$$

$$B = -\frac{(\gamma^2 - \delta^2)}{4M_Z^2}(g_I\sin\theta_W + g_2\cos\theta_W)\left(\frac{\gamma_R}{2}g_I\sin\theta_W\right) - \frac{(\beta g_I\frac{Y_R}{2})^2}{2(m_R^2 - M_Z^2)} - \frac{m_q^2}{4(m_L^2 - M_Z^2)} \begin{cases} \gamma^2/v_I^2; T_3 = +\frac{1}{2} \\ \delta^2/v_2^2; T_3 = -\frac{1}{2} \end{cases} \quad (2b)$$

$$C = -\frac{g_2^2}{4M_W} \left\{ \frac{-\cos\alpha_H/\sin\beta}{\sin\alpha_H/\cos\beta} \right\} \frac{F}{m_h^2} + \frac{g_2}{4M_W} \left\{ \frac{\alpha g_2T_3 + \beta g_I\frac{Y_L}{2}}{m_L^2 - M_Z^2} - \frac{\beta g_I\frac{Y_R}{2}}{m_R^2 - M_Z^2} \right\} \begin{cases} \gamma/v_I\sin\beta; T_3 = +\frac{1}{2} \\ \delta/\cos\beta; T_3 = -\frac{1}{2} \end{cases} \quad (2c)$$

Here $\langle p(n) | (m_u + m_d) (\bar{u}u + \bar{d}d) | p(n) \rangle \equiv 2\hat{f} M_{p(n)} \bar{\Psi} \Psi$ and $\langle p(n) | m_s \bar{s} s | p(n) \rangle \equiv f M_{p(n)} \bar{\Psi} \Psi$; $\hat{f} \approx .05$ and $f \approx .2$ [24].

We now turn to the nuclear matrix element of the scalar ($\bar{\Psi}\Psi$) and axial-vector ($\bar{\Psi}\gamma_\mu\gamma_5\Psi$) nucleon operators at vanishing momentum transfer. The scalar operator simply gives the mass of the nucleus N , $\langle N | M_{p(n)} \bar{\Psi}\Psi | N \rangle = M_N \bar{N}N$. Secondly, $\langle N | \bar{\Psi} \vec{\gamma} \gamma_5 \Psi | N \rangle = \langle N | \vec{S} | N \rangle = \lambda \langle N | \vec{J} | N \rangle$, where $\vec{J}$ is the total nuclear angular momentum. In the single particle shell model, $\lambda = \pm(2L + 1)^{-1}$ for a nucleon in an orbital L , and $J = L \pm \frac{1}{2}$. However, this model is known to be a good approximation only near closed shells, and to be rather poor otherwise. Engel and Vogel [25] have noted that a good approximation is to relax the assumption that a single nucleon of the odd group of nucleons contributes all the nuclear angular momentum, while still assuming that the even group of nucleons has vanishingly small angular momentum, as indicated by nuclear transfer data. Then λ is directly related to the measured magnetic moment, μ_N , by [25]

$$J\lambda = \frac{\mu_N - g_{odd} \frac{L}{2}}{g_{odd} - g_{odd} \frac{L}{2}}, \quad (5)$$

where $g_{odd} \frac{L}{2}$ is the orbital (spin) nucleon g-factor: 1(5.586) for protons and 0(-3.826) for neutrons. For ^{73}Ge ($J = \frac{9}{2}$, $\mu_N = -0.879$), $\lambda^2 J(J+1) = 0.065$.

For large M_χ the typical momentum transfer, $|\vec{q}| = M_\chi |\Delta\vec{v}| = 10^{-3} M_\chi$, gets comparable to the inverse size of a large nucleus and it is necessary to correct the nuclear matrix elements by the appropriate form factors [5,26]. Since we shall be interested here only in the total elastic scattering rate for the LSP mass range M_χ

Here $-ig_2 F = -ig_2(\alpha - \beta \tan\theta_W) \chi \cos\alpha_H + \delta \sin\alpha_H$ is the $h\chi\chi$ coupling. The Higgs-exchange contribution in (2c) agrees with [16]; the sfermion exchange contribution disagrees by a factor of 2 with [17], but agrees with [8]. The sign of the Z-exchange contribution to the axial-vector-axial-vector term in A and B agrees with [17].

In order to obtain the $\chi N_{\text{nucleus}} \rightarrow \chi N_{\text{nucleus}}$ elastic scattering amplitude we first evaluate $\langle \chi, \text{nucleon} | L_{eff} \chi, \text{nucleon} \rangle$. For a general neutralino, the spin-dependent invariant amplitude for scattering off a proton (neutron) p (n) is [18,19]

$$M_{SD}(\chi p(n) \rightarrow \chi p(n)) = 4 \vec{S}_\chi \cdot \vec{S}_{p(n)} \sum_{q \in p(n)} (B - A) \Delta q, \quad (3)$$

where $\langle p(n) | \bar{q} \gamma_\mu \gamma_5 q | p(n) \rangle \equiv 2S_{\mu}^{p(n)} \Delta q$ and $S_{\mu}^{p(n)} = \langle 0, \vec{S}_{p(n)} \rangle$ is the nucleon's four-spin. We use the most recent analysis of the EMC [20] combined with the most recent analyses of hyperon data [21] to determine the Δq 's: $\Delta u = .77 \pm .08$, $\Delta d = -.49 \pm .08$ and $\Delta s = -.15 \pm .08$.

The spin-independent contribution in (2) involves the matrix elements $\langle p(n) | m_q \bar{q} q | p(n) \rangle$. Because of the m_q factor, heavy quarks do not decouple [22], and $\langle p(n) | \sum_{\text{heavy}} m_q \bar{q} q | p(n) \rangle = \frac{2}{9} M_{p(n)} \bar{\Psi} \Psi$. The spin-independent invariant amplitude is then [23]:

$$M_{SI}(\chi p(n) \rightarrow \chi p(n)) = \left(\hat{f} \frac{m_u C_u + m_d C_d}{m_u + m_d} + f C_s \right. \\ \left. + \frac{2}{27} (1 - f - \hat{f}) (C_c + C_b + C_l) \right) \vec{S}_\chi \cdot \vec{S}_{p(n)} \bar{\Psi}. \quad (4)$$

$\leq M_W$, we can approximate both form factors by a gaussian fit [5] $|F(\vec{q})|^2 = \exp(-\alpha(q, \sqrt{\text{sup3}(2)})\alpha(q, \sqrt{\text{sup3}(2)})\alpha(q, \sqrt{\text{sup3}(2)}))$, where $\vec{q}_0^2 \equiv 3(r_{\text{mass}}^2 (3r_{\text{spin}}^2) for the mass (spin) form factor and $r_{\text{mass}} (r_{\text{spin}})$ is the rms mass (spin) radius. We approximate $r_{\text{mass}} \approx r_{\text{charge}}$ [1], and we take the rms charge radius from a fit to the data [27]: $r_{\text{charge}} = (0.30 + 0.89A^{1/3}) \text{ fm}$. Integrating over a Maxwellian neutralino velocity distribution of dispersion σ^2 one gets the following correction factor ζ for the spin-independent cross section (assuming no energy threshold) [28]:$

$$\zeta = \frac{0.573}{b} \left(1 - \frac{\exp\left(\frac{-b}{1+b}\right) \text{erf}\left(\sqrt{\frac{1}{1+b}}\right)}{\text{erf}(1)} \right) \quad (6)$$

where $b \equiv (M_{\tilde{\chi}}^2 M_N^2 / (M_{\tilde{\chi}} + M_N)^2) \frac{8}{9} \sigma^2 r_{\text{charge}}^2$. Likewise, the spin-dependent cross section is corrected by $\zeta(r_{\text{spin}}) = \zeta(r_{\text{charge}}) \times (r_{\text{spin}} / r_{\text{charge}})$. We have calculated previously the ratios $r_{\text{spin}} / r_{\text{charge}}$ in the infinite square-well shell model for a variety of nuclei [5]. It has been noted [29], however, that since spin-dependent elastic scattering is dominated by nucleons at the Fermi surface, whose wave functions are more spread out than that of core nucleons, it is a deficiency of this model that all nucleons are confined within a fixed radius. A more realistic model in this regard is, for example, the square well potential. For ^{73}Ge we find that the ratio does not vary significantly between the two models: $r_{\text{spin}} / r_{\text{charge}} =$

[1] The mass and charge rms radii should not differ significantly since $Z \sim A - Z$. We find explicitly that in the shell models discussed below these radii differ by $< 1\%$, and we shall not distinguish them in our discussion.

1.25 (1.13) for a harmonic (square) well potential. We shall use the harmonic well value in our calculations.

We now combine the above results to obtain the rate of neutralino-nucleus elastic scattering per unit mass of target material. From (3) and (4) we get (for 100% isotopic abundance)

$$R = (R_{S.D.} + R_{S.I.}) \left(\frac{4 M_{\tilde{\chi}} M_N}{(M_{\tilde{\chi}} + M_N)^2} \right) \left(\frac{\rho}{3 \text{ GeV} \cdot \text{cm}^{-3}} \right) \left(\frac{\langle \vec{v}_E \rangle}{320 \text{ km} \cdot \text{s}^{-1}} \right) \frac{\text{events}}{\text{kg} \cdot \text{day}} \quad (7)$$

where ρ is the density of dark matter in the solar neighborhood and $\langle \vec{v}_E \rangle$ is the average velocity with which the particles hit the detector. Here

$$R_{S.D.} = 580 \lambda^2 J(J+1) \zeta(r_{\text{spin}}) \left(10^4 \text{ GeV}^2 \sum_{q \in P(n)} (B - A) \Delta q \right)^2; \quad (7a)$$

$$R_{S.I.} = 210 M_N^2 M_{\tilde{\chi}}^4 \zeta(r_{\text{charge}}) \left(\frac{m_u C_u + m_d C_d}{m_u + m_d} + f_C \right) + \frac{2}{27} (1 - f - \frac{1}{3})(C_c + C_b + C_t)^2. \quad (7b)$$

We have evaluated the rate (7) for an equal mixture of pure ^{73}Ge ($J = \frac{9}{2}$) + ^{76}Ge ($J = 0$), assuming a closure density of neutralinos for a present-day Hubble expansion rate of $50 \text{ kms}^{-1} \text{ Mpc}^{-1}$. To estimate this, we have calculated the density of relics for a general neutralino composition, including all channels kinematically allowed for $M_{\tilde{\chi}} \leq M_W$, and fully expanding the annihilation-cross sections to first

order in $\tilde{p}_Z^2/M_Z^2$ [2]. In Figure 1 we present our results for the elastic scattering rate in the (μ, M_2) plane for two representative values of $\tan\beta$ and m_A . Here we have neglected sfermion mixing for the light quarks, which should be present for $\mu \neq 0$ unless canceled by soft-supersymmetry-breaking terms. For large $\tan\beta$ and μ we find that the contribution to the elastic scattering rate from squark mixing (in the limit $\chi = B$, the most favorable case) is given by $R_{\text{smixing}} \approx 0.05 \cdot \left(\frac{100 \text{ GeV}}{M_{sf}} \right)^8 \left(\frac{\tan\beta}{10} \right)^2 \left(\frac{\mu}{\text{TeV}} \right)^2 \left(\frac{M_N}{M_W} \right)^2 \frac{4M_Z M_N}{(0.2(M_Z + M_N))^2} (kg\text{-day})^{-1}$ for a common sfermion mass scale M_{sf} . Thus, even for $\tan\beta = 10$ and $\mu = 1 \text{ TeV}$, the contribution to the rate of scattering off $^{73}\text{Ge} + ^{76}\text{Ge}$ is $R_{\text{smixing}} \approx 0.06\text{--}0.003 (kg\text{-day})^{-1}$ for $M_Z = 20\text{--}80 \text{ GeV}$ and $\Omega_{LSP} = 1$, and we have neglected it in our survey.

As can be seen in Figure 1, for low values of m_A the predicted elastic scattering rate would exceed $0.1 (kg\text{-day})^{-1}$, the expected level of sensitivity of the detector, over a substantial fraction of the parameter space where a critical LSP density is attainable, making detection possible. In order to quantify this better, in Figure 2 we plot the fraction of the $\Omega_{LSP} = 1$ parameter space where the elastic scattering rate $> 0.1 (kg\text{-day})^{-1}$, as a function of $\tan\beta$ and m_A . Note that a larger fraction of the domain of LSP parameter space where $\Omega_{LSP} = 1$ is possible can be explored if m_A is small, although this is not a large fraction of the total LSP parameter space currently allowed by experimental constraints. In the dark region

[2] Although radiative corrections to the charged Higgs mass now make it possible for it to be lighter than the W boson [30], we have not included $\chi\chi \rightarrow H^+H^-$ annihilations since this is possible in only a very small region of the parameter space we survey. We have also determined that our relic density calculation would fail for the reasons discussed in ref [31] only in regions already excluded by experimental constraints and regions where the elastic scattering rate is already exceedingly small.

this fraction is 100%, making possible the detection or exclusion of the LSP as the main constituent of the galactic halo.

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Figure Captions

Figure 1: Elastic scattering rate, eq. (7), in the (μ, M_2) plane. Here $m_A = 70$ GeV (a) and 200 GeV (b), $m_{top} = 140$ GeV, and $A_t = A_b = 100$ GeV. For the left (right) panels $\mu < 0$ ($\mu > 0$) and for the upper (lower) panels $\tan\beta = 2(8)$. The elastic scattering rate $R(\chi\bar{\chi}N \rightarrow \chi\bar{\chi}N)/(kg\cdot day) > 1$ ($0.1, 0.01$) in the dark (cross-hatched, hatched) regions. The light-shaded region is excluded by experiments at LEP and CDF, and would be extended up to the dashed line by an unsuccessful search for charginos lighter than the Z boson at LEP II.

Figure 2: Fraction of the $\Omega_{LSP} = 1$ parameter space where the elastic scattering rate $R(\chi\bar{\chi}N \rightarrow \chi\bar{\chi}N)/(kg\cdot day) > 0.1$, as a function of $\tan\beta$ and m_A . In the dark region the fraction is 100%, and in the progressively lighter-shaded regions the fraction is at least 50%, 25% and 12.5%. In the dark region, the $\Omega_{LSP} = 1$ parameter space is $\approx 10\%$ of the parameter space currently allowed by experimental constraints for $M_\chi \leq M_W$, and increases to 30%-40% in the lightest shaded region.

- Fig. 1a -

- Fig. 1b -

- Fig. 2 -