объединенный
институт
ядерных
исследований
дубна

E2-93-448

V.A.Bednyakov, H.V.Klapdor-Kleingrothaus*,
S.G.Kovalenko

ON SUSY DARK MATTER DETECTION WITH SPINLESS NUCLEI[1]

Submitted to «Physics Review D»

*Max-Planck-Institut für Kernphysik, D-6900, Heidelberg, Germany
[1]This work was supported in part by the Russian Foundation for Fundamental Research (93-02-3744).

1993

1 Introduction

Analysis of the data on distribution and motion of astronomical objects within our galaxy and far beyond indicates presence of a large amount of non-luminous dark matter (DM). According to estimations, it constitutes more than 90% of the total mass of the universe if a mass density ρ of the universe close to the critical value ρ_{crit} is assumed. The exact equality $\Omega = \rho/\rho_{crit} = 1$, corresponding to a flat universe, is supported by naturalness arguments and by inflation scenarios. Also, in our galaxy most of the mass should be in a dark halo. Detailed models predict a spherical form for the galaxy halo and a Maxwellian distribution for DM particle velocities in the galactic frame. The mass density of DM in the Solar system should be about $\rho \approx 0.3\text{GeV}\cdot\text{cm}^{-3}$ and the DM particles should arrive at the earth's surface with mean velocities $v \approx 320$ km/sec, producing a substantial flux $\Phi = \rho \cdot v/M$ ($\Phi > 10^7\text{cm}^{-2}$ sec^{-1} for the particle mass M$\sim$ 1 GeV). Therefore one may hope to detect DM particles directly, for instance through the elastic scattering from nuclei inside a detector.

The theory of primordial nucleosynthesis restricts the amount of baryonic matter in the universe to 10%. Thus a dominant component of DM is non-baryonic. The recent data by the COBE satellite [1] on anisotropy in the cosmic background radiation and the theory of the formation of large scale structures of the universe lead to the conclusion that non-baryonic DM itself consists of a dominant (70%) "cold" DM (CDM) and smaller (30%) "hot" DM (HDM) component [2], [3].

The neutralino (χ) is a favorable candidate for CDM. This is a Majorana ($\chi^c = \chi$) spin-half particle predicted by supersymmetric (SUSY) models.

There are four neutralinos in the minimal SUSY extension of the standard model (MSSM). They are a mixture of gauginos ($\tilde{W}_3, \tilde{B}$) and Higgsinos ($\tilde{H}_{1,2}$), which are SUSY partners of gauge (W_3, B) and Higgs ($H_{1,2}$) bosons. The DM neutralino χ is the lightest of them. Moreover, χ is assumed to be the lightest SUSY particle (LSP) which is stable in SUSY models with R-parity conservation.

The problem of direct detection of the DM neutralino χ via elastic scattering off nuclei has been considered by many authors and remains a field of great experimental and theoretical activity [4]-[11].

The final goal of theoretical calculations in this problem is the event rate R for elastic χ-nucleus scattering. In general, the spin-dependent (R_{sd}) and spin-independent (R_{si}) neutralino-nucleus interactions contribute to the event rate:

$R = R_{sd} + R_{si}$. R_{sd} vanishes for spinless nuclei and this fact is often regarded as a reason to assert spinless nuclei to be irrelevant for the DM neutralino detection as giving a much smaller event rate. One can meet this statement in the literature. However, this is right only if the spin-dependent interaction dominates in elastic neutralino scattering off nuclei with non-zero spin.

In this paper we address the question on the role of nuclear spin in the DM-neutralino detection. We investigate this problem in the framework of the MSSM. We avoid using specific nuclear and nucleon structure models but rather base our consideration on the known experimental data about nuclei and nucleon. It allows us to free the consideration of theoretical uncertainties specific for the structure models. To restrict the MSSM parameter space we use experimental constraints on SUSY-particle masses, the cosmological bound on neutralino relic abundance and the proton life-time constraint.

We have found that R_{si} contribution dominates in the total event rate R for nuclei with atomic weight $A > 50$ in the region of the MSSM parameter space where $R = R_{sd} + R_{si} < 0.01$. The lower bound 0.01 is far below the sensitivity of realistic present and near future DM detectors. Therefore we can exclude the region where $R < 0.01$ as invisible for these detectors.

We do *not* expect a crucial dependence of the DM event rate on the nuclear spin for detectors with target nuclei having an atomic weight larger than 50. As a result, we expect equal chances for J = 0 and J $\neq$ 0 detectors to discover DM events. In particular, this conclusion supports the idea that presently operating $\beta\beta$-detectors with spinless nuclear target material can be successfully used for DM neutralino search. These highly developed set-ups (for a review see [12]), operating under extremely low background conditions, use detection technology which is suitable for the DM search.

2 General Properties of the Neutralino - Nucleus Interactions

A DM event is elastic neutralino-nucleus scattering causes the nuclear recoil detected by a detector. The event rate per unit mass of the target material depends on the distribution of the DM neutralinos in the solar vicinity and the cross section $\sigma_{el}(\chi A)$ of neutralino-nucleus elastic scattering. One can calculate $\sigma_{el}(\chi A)$ starting from the neutralino-quark effective Lagrangian. In

the most general form it can be given by the formula

$$L_{eff} = \sum_q \{\mathcal{B}(q) - \mathcal{A}(q)\} \cdot \bar{\chi}\gamma_\mu\gamma_5\chi \cdot \bar{q}\gamma^\mu\gamma_5 q + \frac{m_q}{M_W} \cdot \mathcal{C}_q \cdot \bar{\chi}\chi \cdot \bar{q}q, \quad (1)$$

where terms with the vector and pseudoscalar quark currents are omitted being negligible in the case of the non-relativistic DM neutralino with typical velocities $v_\chi \approx 10^{-3}c$. Following the tradition we retain in the first term the difference of two coefficients representing just one independent parameter of the Lagrangian. The coefficients $\mathcal{B}(q), \mathcal{A}(q), \mathcal{C}(q)$ depend on the SUSY model and will be considered in the next section. Here we survey general properties of χ-A scattering following from the Lagrangian (1).

To calculate $\sigma_{el}(\chi A)$ one should average the χ-q interactions sequentially over the nucleon and the nuclear structure. The first and the second terms in L_{eff} (1) averaged over the nucleon states give the spin-dependent and the spin-independent matrix elements $\mathbf{M}_{sd}$ and $\mathbf{M}_{si}$, respectively. For the spin-dependent amplitude we have [4], [5]:

$$\mathbf{M}_{sd} = 4\vec{S}_\chi\vec{S}_{p(n)} \sum_{q\in p(n)} \{\mathcal{B}(q) - \mathcal{A}(q)\}\Delta q, \quad (2)$$

where $\vec{S}_\chi$ and $\vec{S}_{p(n)}$ are the neutralino and proton (neutron) spin operators; Δq are the fractions of the nucleon spin carried by the quark q. The standard definition is

$$< p(n)|\bar{q}\gamma^\mu\gamma_5 q|p(n) >= 2S^\mu_{p(n)}\Delta q, \quad (3)$$

where $S^\mu_{p(n)} = (0, \vec{S}_{p(n)})$ is the 4-spin of the nucleon. The parameters Δq (for the proton) can be extracted from the EMC [13] and hyperon data [14]:

$$\Delta u = 0.77 \pm 0.08, \quad \Delta d = -0.49 \pm 0.08, \quad \Delta s = -0.15 \pm 0.08. \quad (4)$$

The relevant values for the neutron can be found from (4) by the isospin symmetry substitution $u \to d$, $d \to u$.

The spin-independent matrix element has the form [7], [8][1]:

$$\mathbf{M}_{si} = \left[\hat{f}\frac{m_u\mathcal{C}(u) + m_d\mathcal{C}(d)}{m_u + m_d} + f\mathcal{C}(s) + \frac{2}{27}(1 - f - \hat{f})(\mathcal{C}(c) + \mathcal{C}(b) + \mathcal{C}(t))\right] \cdot \frac{M_{p(n)}}{M_W} \cdot \bar{\chi}\chi \cdot \bar{\Psi}\Psi, \quad (5)$$

[1] When this paper had been completed we received a paper ref.[10] with more refined treatment of the spin-independent matrix element.

where the parameters f and $\hat{f}$ are defined as follows:

$$
\begin{aligned}
< p(n)|(m_u + m_d)(\bar{u}u + \bar{d}d)|p(n) > &= 2\hat{f}M_{p(n)}\bar{\Psi}\Psi, \qquad (6)\\
< p(n)|m_s\bar{s}s|p(n) > &= fM_{p(n)}\bar{\Psi}\Psi.
\end{aligned}
$$

The values extracted from the data are [15],[16]: $\hat{f} = 0.05$ and $f = 0.2$.

Averaging (2), (5) over the nuclear states $|A>$ we deal with the following matrix elements at vanishing momentum transfer:

$$
\begin{aligned}
< A|M_{p(n)}\bar{\Psi}\Psi|A > &= M_A\bar{A}A, \qquad (7)\\
< A|\vec{S}_{p(n)}|A > &= \lambda < A|\vec{J}|A > .
\end{aligned}
$$

Here $\vec{J}$ is the nuclear spin. On the basis of the odd-group shell model [17] (essentially somewhat relaxed single particle shell model) the parameter λ can be related to the nuclear magnetic moment, μ, as follows

$$\lambda J = \frac{\mu - g^l J}{g^s - g^l}, \qquad (8)$$

where $g^l = 1(0)$ and $g^s = 5.586(-3.826)$ are orbital and spin proton (neutron) g-factors. Then one can extract values of λ for various nuclei from the experimental data on nuclear magnetic moments [2]. We use in this paper the values of λ as presented in ref.[17].

For large M_χ and M_A the momentum transfer may be comparable to the inverse radius of a nucleus and then we have to take into account the finite size effect (see also [18]). It can be done by introducing the coherence loss factor [19].

$$\zeta(r) = \frac{0.573}{b}\left(1 - \frac{\exp(-\frac{b}{1+b})}{\sqrt{1+b}}\frac{\mathrm{erf}(\sqrt{\frac{1}{1+b}})}{\mathrm{erf}(1)}\right), \qquad (9)$$

where

$$b = \frac{8}{9}\sigma^2 r^2 \frac{M_\chi^2 M_A^2}{(M_\chi + M_A)^2}.$$

Here σ^2 is the dispersion of the Maxwellian neutralino velocity distribution $\sigma = 0.9 \cdot 10^{-3}$. To obtain the coherence loss factor for spin-independent scattering

[2] A more direct way of calculation based on the theory of finite Fermi systems is presented in [18].

we take $r = r_{charge}$ in (9), where r_{charge} is the rms charge radius of the nucleus A [9]:

$$r_{charge} = (0.3 + 0.89M_A^{1/3}) \text{ fm}. \qquad (10)$$

The coherence loss factor for spin-dependent scattering is given by (9) with $r = r_{spin}$. The rms spin radius of the nucleus A can be estimated as $r_{spin} = \xi \cdot r_{charge}$ with $\xi \approx 1.25$ from harmonic well potential calculations [9].

Finally we arrive at the formula for the event rate of elastic neutralino-nucleus scattering in the detector per day per unit mass of the target material:

$$R = R_{si} + R_{sd}, \qquad (11)$$

where the spin-dependent and spin-independent parts are:

$$
\begin{aligned}
R_{sd} &= 5.8 \cdot 10^{10} \cdot \lambda^2 J(J+1)\zeta(r_{spin})\mathcal{M}_{sd}^2 \cdot \mathcal{D}, \qquad (12)\\
R_{si} &= 1.44 \cdot 10^{10} \cdot \left(\frac{M_A}{M_W}\right)^2 \zeta(r_{charge}) \cdot \mathcal{M}_{si}^2 \cdot \mathcal{D}. \qquad (13)
\end{aligned}
$$

The common kinematic factor $\mathcal{D}$ and properly normalized nucleon matrix elements $\mathcal{M}_{si}, \mathcal{M}_{sd}$ are defined as:

$$
\begin{aligned}
\mathcal{D} &= \left[\frac{4M_\chi M_A}{(M_\chi + M_A)^2}\right]\left[\frac{\rho}{.3GeV \cdot cm^{-3}}\right]\left[\frac{<|\vec{v_E}|>}{320km/s}\right]\frac{events}{kg \cdot day} \qquad (14)\\
\mathbf{M}_{si} &= \mathcal{M}_{si} \cdot \frac{M_{p(n)}}{M_W}\bar{\chi}\chi \cdot \bar{\Psi}\Psi, \qquad (15)\\
\mathbf{M}_{sd} &= 4 \cdot \mathcal{M}_{sd} \cdot \vec{S}_\chi \vec{S}_{p(n)}. \qquad (16)
\end{aligned}
$$

For the definition of $\mathbf{M}_{si}$, $\mathbf{M}_{sd}$ see formulae (2), (5). Here $\rho \approx 0.3$ GeV$\cdot$cm^{-3} is the DM neutralino density in the solar vicinity and $<|\vec{v_E}|> \approx 320$ km/s is DM neutralino averaged velocity at the earth's surface.

To study the role of nuclear spin in elastic χ-nucleus scattering we introduce the ratio

$$\kappa = R_{sd}/R_{si}. \qquad (17)$$

characterizing the relative contribution of spin-dependent and spin-independent interactions. From the practical point of view it determines the expected relative sensitivity of DM detectors with spin-non-zero (J $\neq$ 0) and spin-zero (J = 0) nuclei as target material. If $\kappa < 1$, then detectors with spin-non-zero and spin-zero target materials have approximately equal sensitivities to the DM signal, whereas if $\kappa > 1$ then, the spin-non-zero detectors are more sensitive than the spin-zero ones.

Let us consider separately the dependence of κ on the nuclear structure and the choice of a specific SUSY model. We may write:

$$\kappa = \eta_A \eta_{susy}^{p(n)}, \tag{18}$$

where

$$\eta_A = 4.03\lambda^2 J(J+1)\cdot\frac{\zeta(r_{spin})M_W^2}{\zeta(r_{charge})M_A^2}, \tag{19}$$

$$\eta_{susy}^{p(n)} = \left(\frac{\mathcal{M}_{sd}^{p(n)}}{\mathcal{M}_{si}}\right)^2. \tag{20}$$

Here η_A is a factor depending on the properties of the nucleus A; $\eta_{susy}^{p(n)}$ is defined by the SUSY-model which specifies the neutralino composition and the interactions with matter. The SUSY-factor also depends on the shell-model class to which nucleus A belongs, being η_{SUSY}^n for the shell-model "neutron" (^{3}He, ^{29}Si, ^{73}Ge,...) and η_{SUSY}^p for the shell-model "proton" (^{19}F, ^{35}Cl, ^{205}Tl,...).

Fig.1 depicts the nuclear factor η_A versus the atomic weight A. The error bars represent the interval of the neutralino masses 20 GeV$< M_\chi <$ 200 GeV. The lower bound corresponds to the present experimental constraints [20]-[22]. The upper bound is taken to include recent estimations for the mass of the cosmologically favorable neutralino [23]. It follows from fig.1 that $\eta_A < 1$ for $A > 50$. Thus at $A > 50$ there is no nuclear structure enhancement of the spin-dependent event rate as compared to the spin-independent one.

The next is an estimation of the SUSY-factor $\eta_{susy}^{p(n)}$.

3 Specific SUSY-model Predictions

To estimate the factor η_{SUSY} in (20) one should calculate the parameters $\mathcal{A}(q)$, $\mathcal{B}(q)$ and $\mathcal{C}(q)$ of the effective Lagrangian (1) in the specific SUSY model. We will follow the MSSM. This model is specified by the superpotential and "soft" SUSY breaking terms [24].

The effective low energy superpotential is:

$$\tilde{W} = \sum_{generations}(h_U\hat{H}_2\hat{Q}\hat{U} + h_D\hat{H}_1\hat{Q}\hat{D} + h_L\hat{H}_1\hat{L}\hat{E}) + \mu\hat{H}_1\hat{H}_2. \tag{21}$$

H_1 and H_2 are the Higgs fields with a weak hypercharge $Y = -1, +1$ respectively.

SUSY breaking in the "hidden" sector of N=1 supergravity produces "soft" supersymmetry breaking terms in the scalar potential:

$$V_{soft} = \sum_{i=scalars} m_i^2|\phi_i|^2 + h_U A_U H_2\tilde{Q}\tilde{U} + h_D A_D H_1\tilde{Q}\tilde{D} + h_L A_L H_1\tilde{L}\tilde{E} \tag{22}$$
$$+\mu B H_1 H_2 + \text{ h.c.}$$

and a "soft" gaugino mass term

$$\mathcal{L}_{FM} = -\frac{1}{2}\left[M_1\tilde{B}\tilde{B} + M_2\tilde{W}^3\tilde{W}^3 + M_3\tilde{g}^a\tilde{g}^a\right] - \text{ h.c.} \tag{23}$$

The model is also characterized by the set of boundary conditions at the unification scale M_X:

$$A_U = A_D = A_L = A_0, \tag{24}$$

$$m_{H_1} = m_{H_2} = m_L = m_E = m_Q = m_U = m_D = m_0, \tag{25}$$

$$M_1 = M_2 = M_3 = m_{1/2}, \tag{26}$$

$$g_1(M_X) = g_2(M_X) = g_2(M_X) = g_{GUT}, \tag{27}$$

where g_3, g_2, g_1 are the $SU(3)\times SU(2)\times U(1)$ gauge coupling constants equal to g_{GUT} at the unification scale M_X. At the Fermi scale $Q \sim M_W$ these parameters can be evaluated on the basis of the 1-loop renormalization group equations (RGE) [25],[26].

The neutralino mass matrix in this model has the form [24]:

$$M_\chi = \begin{pmatrix} M_2 & 0 & -M_Z c_W s_\beta & M_Z c_W c_\beta \\ 0 & M_2\frac{5}{3}t_W & M_Z s_W s_\beta & -M_Z c_W s_\beta \\ -M_Z c_W s_\beta & M_Z s_W s_\beta & 0 & -\mu \\ M_Z c_W c_\beta & -M_Z s_W c_\beta & -\mu & 0 \end{pmatrix}. \tag{28}$$

where $c_W = \cos\theta_W$, $s_W = \sin\theta_W$, $t_W = \tan\theta_W$, $s_\beta = \sin\beta$, $c_\beta = \cos\beta$. The matrix is written in the basis of fields $(\tilde{W}^3, \tilde{B}, \tilde{H}_2^0, \tilde{H}_1^0)$. As usual, M_2, μ are the gaugino mass and the Higgs mixing parameter; the angle β is defined by the vacuum expectation values of the neutral components of the Higgs fields: $\tan\beta = < H_2^0 > / < H_1^0 > = v_2/v_1$. Diagonalizing the mass matrix (28) we obtain the lightest neutralino of the mass M_χ with the field content

$$\chi = \alpha\tilde{W}^3 + \beta\tilde{B} + \gamma\tilde{H}_2^0 + \delta\tilde{H}_1^0.$$

To calculate the low-energy neutralino-quark interactions we also need to have the spectrum of squarks $\tilde{q}$ and Higgs particles at the Fermi scale. Their mass matrices depend on the soft SUSY breaking parameters. We obtain them from the boundary conditions at the GUT scale M_X (24)-(26) as a solution of the 1-loop RGE.

We analyze the Higgs sector of the MSSM at the 1-loop level [27]. In the analysis we take into account $\tilde{t}_L - \tilde{t}_R$, $\tilde{b}_L - \tilde{b}_R$ mixing between the third-generation squarks. Diagonalization of the Higgs mass matrix gives three neutral mass-eigenstates. There are two CP-even states, H, h, with the masses m_H, m_h and the relevant mixing angle α_H and one CP-odd A with the mass m_A. We take the mass m_A as an independent free parameter of the MSSM.

A complete list of essential free parameters of the MSSM is

$$\tan\beta, A_0, B, \mu, m_{1/2}, m_A, m_0, m_t. \tag{29}$$

Having a particle spectrum one can derive the effective Lagrangian L_{eff} of low-energy neutralino-quark interactions. As discussed in the previous section, its general form is given by eqn. (1). In the MSSM the first term of L_{eff} is induced by the Z-boson and $\tilde{q}$ exchange [28] whereas the second one is due to the Higgs particle [29] and $\tilde{q}$ exchange [6] as well as $\tilde{q}_L - \tilde{q}_R$ mixing [4],[30]. The coefficients of L_{eff} are

$$\begin{aligned}\mathcal{A}(q) &= -\frac{|\gamma|^2-|\delta|^2}{4M_Z^2}(g_1\sin\theta_W + g_2\cos\theta_W)(\frac{Y_L}{2}g_1\sin\theta_W - T_3 g_2\cos\theta_W)\\ &+ \frac{1}{2}\frac{|\alpha g_2 T_3 + \beta g_1 \frac{Y_L}{2}|^2}{m_{\tilde{q}L}^2 - M_\chi^2} + \frac{1}{2}\frac{m_q^2}{m_{\tilde{q}R}^2 - M_\chi^2}[\frac{\frac{1}{2}+T_3}{v_2^2}|\gamma|^2 + \frac{\frac{1}{2}-T_3}{v_1^2}|\delta|^2], \end{aligned}\tag{30}$$

$$\begin{aligned}\mathcal{B}(q) &= -\frac{|\gamma|^2-|\delta|^2}{4M_Z^2}(g_1\sin\theta_W + g_2\cos\theta_W)(\frac{Y_R}{2}g_1\sin\theta_W)\\ &- \frac{1}{2}\frac{|\beta g_1 \frac{Y_R}{2}|^2}{m_{\tilde{q}R}^2 - M_\chi^2} - \frac{1}{2}\frac{m_q^2}{m_{\tilde{q}L}^2 - M_\chi^2}[\frac{\frac{1}{2}+T_3}{v_2^2}|\gamma|^2 + \frac{\frac{1}{2}-T_3}{v_1^2}|\delta|^2], \end{aligned}\tag{31}$$

$$\begin{aligned}\mathcal{C}(q) &= \frac{g_2^2}{4m_{h1}^2}F_q[(\frac{1}{2}+T_3)\frac{\cos\alpha_H}{\sin\beta} - (\frac{1}{2}-T_3)\frac{\sin\alpha_H}{\cos\beta}]\\ &+ \frac{g_2}{4}[\frac{\alpha g_2 T_3 + \frac{Y_L}{2}\beta g_1}{m_{\tilde{q}L}^2 - M_\chi^2} - \frac{\frac{Y_R}{2}g_1\beta}{m_{\tilde{q}R}^2 - M_\chi^2}][(\frac{1}{2}+T_3)\frac{\gamma}{\sin\beta} - (\frac{1}{2}-T_3)\frac{\delta}{\cos\beta}]. \end{aligned}\tag{32}$$

Here

$$F_q = (\alpha - \beta\tan\theta_W)(\gamma\cos\alpha_H + \delta\sin\alpha_H). \tag{33}$$

In these formulae we ignore $\tilde{q}_L - \tilde{q}_R$ mixing since they give a small contribution according to the estimation of ref. [9].

The procedure we use for the neutralino mass matrix diagonalization always leads to positive mass eigenvalues and to either real or pure imaginary values of the coefficients $\alpha, \beta, \gamma, \delta$. Therefore in formulae (30),(31) the absolute values of these coefficients appear.

Now we are ready to calculate the η_{susy}-factor (20) substituting the definitions (30)-(32) in formula (20).

To get complete information about possible values of the η_{susy}-factor we scan the MSSM parameter space within the constraints imposed by the experimental data and some general theoretical principles. The well known experimental constraints [22] are summarized in the Table.

Table: Present Limits on Supersymmetric Particles

(Table is taken from ref.[31].)

Particle	Bound on Particle Mass (GeV)	Source
$\tilde{\chi}_1^\circ$	18.4	Based on the LEP non-observation of $\tilde{\chi}_i^\circ$
$\tilde{\chi}_2^\circ$	45	and $\tilde{\chi}_1^+$, CDF non-observation of $\tilde{g}$, and
$\tilde{\chi}_3^\circ$	70	the assumption of gaugino mass unifica-
$\tilde{\chi}_4^\circ$	108	tion.
$\tilde{\chi}_1^\pm$	45.2	LEP
$\tilde{\chi}_2^\pm$	99	See neutralino mass limits above.
$\tilde{\nu}$	41	Assumes the $\tilde{\nu}$ decays are invisible (otherwise $M_{\tilde{\nu}} < 32$ GeV). Based on LEP measurement of $\Gamma(Z \to$ invisible final states).
$\tilde{e}$	45	LEP; assumes $M_{\tilde{\chi}_1^\circ} < 41$ GeV
$\tilde{\mu}$	45	LEP; assumes $M_{\tilde{\chi}_1^\circ} < 41$ GeV
$\tilde{\tau}$	45	LEP; assumes $M_{\tilde{\chi}_1^\circ} < 38$ GeV
$\tilde{q}$	45	LEP; assumes $M_{\tilde{\chi}_1^\circ} < 20$ GeV
	74	UA2 (any $M_{\tilde{g}}$)
	~ 95	CDF ($M_{\tilde{q}} < M_{\tilde{g}}$)
$\tilde{g}$	79	UA2 ($M_{\tilde{g}} < M_{\tilde{q}}$)
	~ 95	CDF

The constraints are given for masses of squarks $\tilde{q}$, the gluino $\tilde{g}$, charginos $\chi^\pm$, the neutralino χ, charged sleptons $\tilde{l}^\pm$, the sneutrino $\tilde{\nu}$, lightest CP-even h

and CP-odd A Higgs bosons. We also include the constraints

$$1.12 < \tan\beta < 4.7. \tag{34}$$

The lower limit follows from the finiteness condition for the top Yukawa coupling Y_t. If

$$\sin\beta > (m_t/200\ GeV), \tag{35}$$

then Y_t is finite up to the unification scale M_X. For $m_t = 150$ GeV we get the lower limit in (34). The upper limit in this formula is expected from proton stability considerations [32].

From the "naturalness" arguments [33] we may choose:

$$m_{\tilde{q}}, m_{\tilde{g}} < 1\ TeV, \tag{36}$$

where $m_{\tilde{f}}$ is the mass of any sfermion $\tilde{f}$. The choice of the interval for the neutralino mass

$$20\ GeV < M_\chi < 200\ GeV$$

was already explained at the end of *section 2.*

The additional constraint we use in the analysis of the role of nuclear spin is the constraint on the realistic sensitivity of the DM detector. In terms of the total event rate R we choose the sensitivity to be not better than:

$$R > 0.01 \frac{event}{kg \cdot days}. \tag{37}$$

We do not expect the DM detectors to go below this lower bound in near future. Therefore the constraint (37) reflects the realistic capacities of the present and near-future set-ups. It excludes the region in the parameter space corresponding to the low-level DM signals inaccessible to these detectors.

We have performed a numerical analysis of the MSSM parameter space within the above-defined constraints. In fig.2 the typical behavior of the η_{susy}-factor in particular domains of the MSSM parameter space is presented. The following upper bound for the SUSY-factor in eqn. (18) was found:

$$\eta_{susy} \leq 1.2. \tag{38}$$

Combining this result with the values of the nuclear factor η_A represented in fig.1 we conclude that

$$\kappa = R_{sd}/R_{si} = \eta_A \eta_{susy}^{p(n)} \leq 1 \quad \text{for nuclei with } A > 50 \tag{39}$$

at a detector sensitivity up to $R > 0.01$. The tendency is that at higher sensitivities (lower R accessible) we get $\kappa \leq 1$ for heavier nuclei.

As a by-product of our analysis in fig.3 we also give the event rate for some nuclei of special interest in DM search.

We do not take into account possible rescaling of the local neutralino density ρ which may occur in the region of the MSSM parameter space where $\Omega h^2 < 0.05$ [7]. This effect, if it took place, would essentially reduce the event rate R [11]. Of course, it has no influence on the ratio κ in the formula (17) and on our conclusion about the role of nuclear spin. Plots in fig.3 correspond to a situation when neutralinos constitute a dominant component of the DM halo of our galaxy with the density $\rho = 0.3$ GeV$\cdot$cm^{-3} in the solar vicinity.

4 Conclusions

In the framework of general assumptions about the nuclear and nucleon structure considering the MSSM as the basis for description of the neutralino properties we have drawn the following basic conclusions.

For sufficiently heavy nuclei with atomic weights $A > 50$ the spin-independent event rate R_{si} is larger than the spin-dependent one R_{sd} if low-level signals with total event rates $R = R_{sd} + R_{si} < 0.01$ are ignored. This cut condition reflects the realistic sensitivities of the present and the near-future DM detectors.

The main practical issue is that two different DM detectors with (J = 0, A_1) and with (J $\neq$ 0, A_2) nuclei as target material have equal chances to discover the DM event if $A_1 \sim A_2 > 50$.

Another aspect of the DM search is the investigation of the SUSY-model parameter space from nonobservation of DM events. Apparently, in this case experiments with J $\neq$ 0 nuclei are important since they provide new information about the SUSY model parameters from R_{sd} which is inaccessible in J = 0 experiments.

The results presented above were obtained in a specific SUSY-model. Therefore it is a natural question whether our basic conclusions hold for other popular SUSY-models. We plan to investigate this question in a subsequent paper.

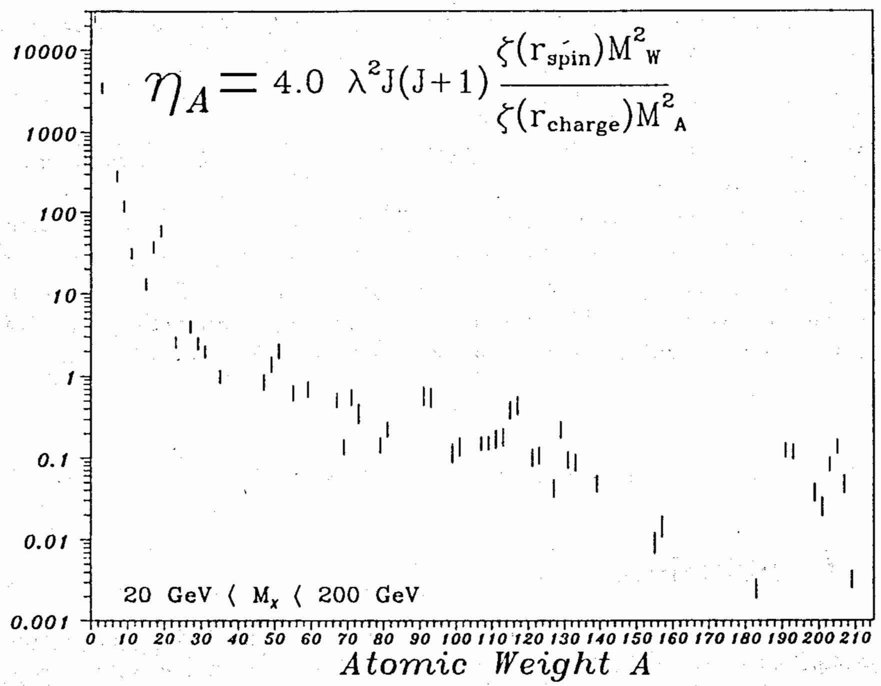
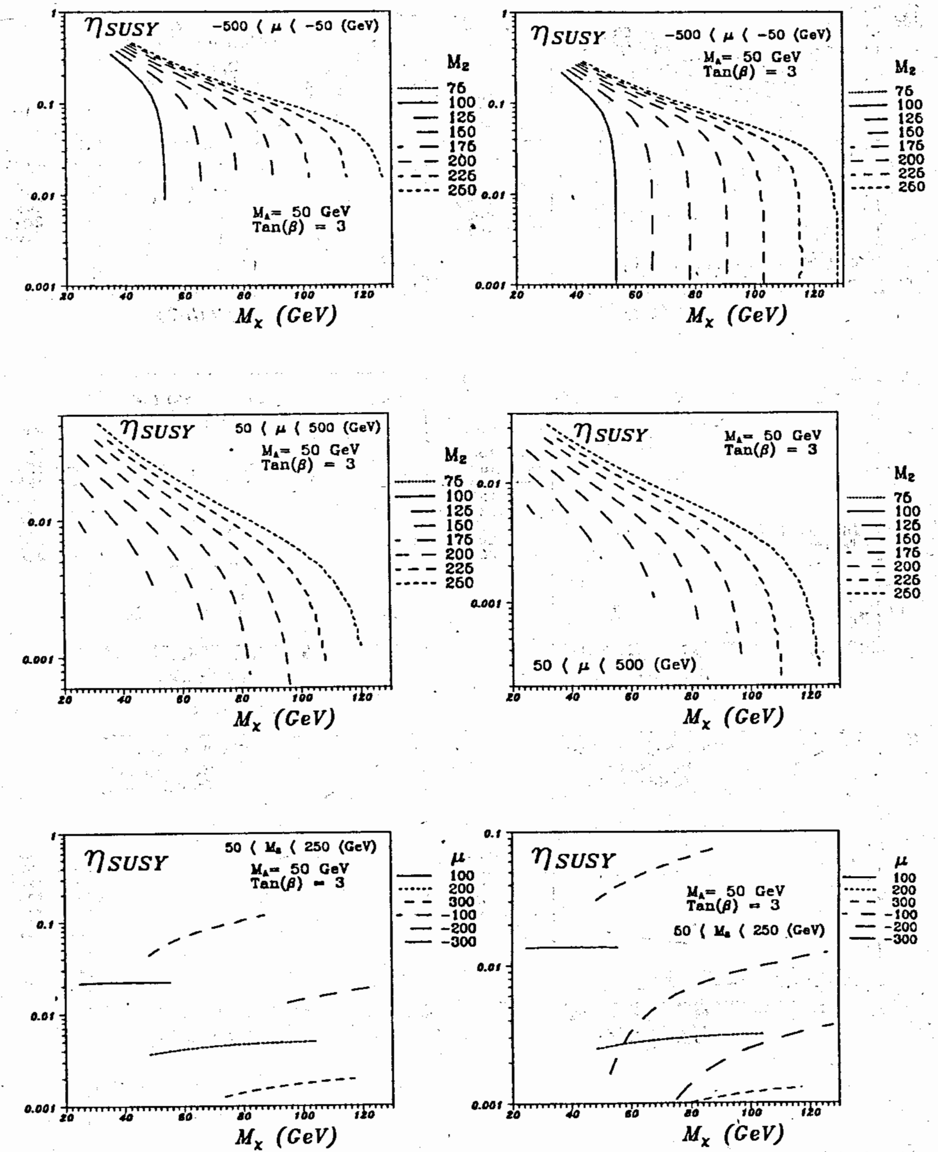


Fig.1. The nuclear factor η_A versus the atomic weight A. The error bars represent the interval of the neutralino masses 20 GeV< M_χ < 200 GeV

Fig.2. The SUSY factor η_{susy} versus neutralino mass M_χ at various values of the MSSM free parameters. (a) and (b) for nuclei with proton and neutron shell model structure, respectively



(a) (b)

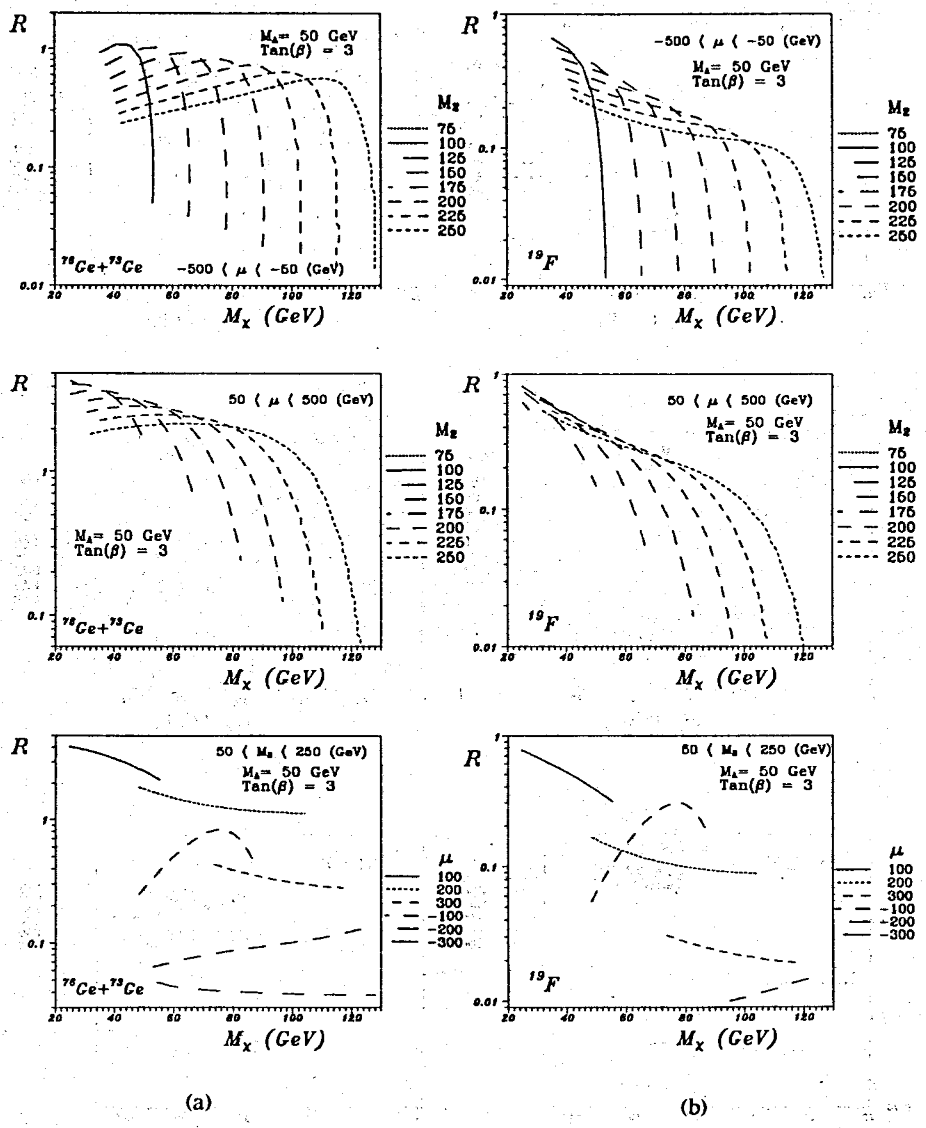


Fig.3. Samples of plots for the total event rate for neutralino elastic scattering off (a) equal parts mixture of $^{73}Ge +^{76} Ge$ and (b) ^{19}F nuclei

References

[1] G.F.Smoot et al.,Astrophys.J. 396 (1992) L1.

[2] A.N.Taylor, M.Rowan-Robinson, Nature 359 (1992) 336.

[3] M.Davis, F.J.Summers and D.Schlegel, Nature 359 (1992) 393.

[4] M.W.Googman, E.Witten, Phys.Rev. D31 (1985) 3059.

[5] J.Ellis, R.Flores, Nucl.Phys. B307 (1988) 883.

[6] K.Griest, Phys.Rev. D38 (1988) 2357; D39 (1989) 3802.

[7] G.B.Gelmini, P.Gondolo and E.Roulet, Nucl.Phys. B351 (1991) 623.

[8] R.Flores, K.A.Olive and D.Thomas, Phys.Lett. B263 (1991) 425.

[9] J.Ellis, R.Flores, Phys.Lett., B263 (1991) 259; Phys.Lett. B300 (1993) 175; CERN-TH-6588/92.

[10] M. Drees, M.M. Nojiri, Univ. Wisconsin - Madison, 1993, prepr. MAD/PH/768.

[11] A.Bottino, V.de Alfaro, N.Fornengo, G.Mignola and S. Scopel, Astropart.Phys.J. 1 (1992) 61.

[12] H.V.Klapdor-Kleingrothaus in Proc. Weak and Electromagnetic Int. in Nuclei, Dubna, Russia 1992, World Scientific, Singapore 1993, p.201; Nucl.Phys. (proc. Suppl.) B31 (1993) 72; H.V.Klapdor-Kleingrothaus in Proc. Neutrinos in Cosmology, Astro, Particle and Nuclear Physics, Erice, Sicily, Italy, 8-17 Sept.,1993, Plenum press, in print.

[13] J.Ashman et al., EMC collaboration, Nucl.Phys. B328 (1989) 1.

[14] A.Manohar, R.Jaffe, Nucl.Phys. B337 (1990) 509.

[15] T.P.Cheng, Phys.Rev. D38 (1988) 2869; H.-Y.Cheng, Phys.Lett. B219 (1989) 347.

[16] J.Gasser, H.Leutwyler and M.E. Sainio, Phys.Lett. B253 (1991) 252.

[17] J.Engel, P.Vogel, Phys.Rev. D40 (1989) 3132.

[18] M.A.Nikolaev, H.V.Klapdor-Kleingrothaus, Z.Phys. A345 (1993) 183; 373.

[19] K.Freese, J.Frieman and A.Gould, Phys.Rev. D37 (1988) 3388; A.Gould, Astrophys.J. 321 (1987) 571.

[20] L.Roszkowski, Phys.Lett. B262 (1991) 59; Phys.Lett. B252 (1990) 471.

[21] K.Hidaka, Phys.Rev. D44 (1991) 927.

[22] Review of Particle Properties, Phys.Rev. D45, No.11 (1992) S1.

[23] R.G.Roberts, L.Roszkowski, Phys.Lett. B309(1993) 329;

[24] H.E.Haber, G.L.Kane, Phys.Rep. 117 (1985) 75.

[25] L.E.Ibañez, C.Lopez, Phys.Lett. B126 (1983) 54; L.E.Ibañez, C. Lopez and C.Muñoz, Nucl.Phys. B256 (1985) 218.

[26] K.Inoue, A.Kakuto, H.Komatsu and S. Takeshita, Progr.Theor.Phys. 68 (1982) 927; 71 (1984) 348.

[27] Y.Okada, M.Yamaguchi and T.Yanagida, Progr.Theor.Phys. 85 (1991) 1; Phys.Lett. B262 (1991) 54; J.Ellis, G.Ridolfi and F.Zwirner, Phys.Lett. B257 (1991) 83; B262 (1991) 477; H.E.Haber, R.Hempfling, Phys.Rev.Lett. 66 (1991) 1815; R.Barbieri, M.Frigeni, Phys.Lett. B258 (1991) 395.

[28] J.Ellis, J.S.Hagelin, D.V.Nanopoulos, K.Olive and M.Srednicki, Nucl.Phys. B238 (1984) 453.

[29] R.Barbieri, M.Frigeni and G.F. Guidice, Nucl.Phys. B313 (1989) 725.

[30] M.Srednicki, R.Watkins, Phys.Lett. B225 (1989) 140.

[31] H.E. Haber, in Proceedings of the Theoretical Advanced Study Institute, Colorado 1992 (to be published).

[32] P.Nath, R.Arnowitt, Phys.Rev.Lett. 69 (1992) 725; Phys.Lett. B287 (1992) 89; Phys.Lett. b289 (1992) 368.

[33] R.Barbieri, G.F.Guidice, Nucl.Phys. B306 (1988) 63; G.G.Ross, R.G.Roberts, Nucl.Phys. B377 (1992) 571.

Received by Publishing Department
on December 27, 1993.

Бедняков В.А., Клапдор-Кляйнгротхауз Х.В., Коваленко С.Г. E2-93-448
К вопросу о детектировании SUSY-темной материи
с помощью бесспиновых ядер

На основе общепринятых предположений о нуклонной и ядерной структуре в рамках минимального суперсимметричного расширения стандартной модели (MSSM) проанализированы возможности детектирования космических нейтралино при помощи детекторов с различными ядрами мишени, в том числе ^{73}Ge–^{76}Ge. Показано, что для ядер с атомным весом A > 50 скорость счета событий, обусловленная спин-независимым взаимодействием R_{si}, вследствие когерентного усиления заметно превышает скорость счета событий R_{sd} от спин-спинового взаимодействия. Поэтому детекторы, где в качестве мишени применяются изотопы с атомным весом A > 50, имеют примерно равные возможности зарегистрировать сигнал от частиц темной материи, при близких значениях атомных весов этих изотопов и независимо от наличия спина у ядра.

Работа выполнена в Лаборатории ядерных проблем ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна, 1993

Bednyakov V.A., Klapdor-Kleingrothaus H.V., Kovalenko S.G. E2-93-448
On SUSY Dark Matter Detection with Spinless Nuclei

We investigate the role of nuclear spin in elastic scattering of Dark Matter (DM) neutralinos from nuclei in the framework of the Minimal SUSY standard model (MSSM). The relative contribution of spin-dependent axial-vector and spin-independent scalar interactions to the event rate in a DM detector has been analyzed for various nuclei. Within general assumptions about the nuclear and nucleon structure we find that for nuclei with atomic weights $A > 50$ the spin-independent part of the event rate R_{si} is larger than the spin-dependent one R_{sd} in the domain of the MSSM parameter space allowed by the known experimental data and where the additional constraint for the total event rate $R=R_{sd}+R_{si} >$ >0.01 is satisfied. The latter reflects realistic sensitivities of present and near future DM detectors. Therefore we expect equal chances for discovering the DM event either with spin-zero or with spin-non-zero isotopes if their atomic weights are $A_1 \sim A_2 > 50$.

The investigation has been performed at the Laboratory of Nuclear Problems, JINR.

Preprint of the Joint Institute for Nuclear Research. Dubna, 1993