

Superlight neutralino as a dark matter particle candidate

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We address the question of how light the lightest supersymmetric particle neutralino can be to be a reliable cold dark matter (CDM) particle candidate. To this end we perform a combined analysis of the parameter space of the minimal supersymmetric standard model (MSSM) taking into account cosmological and accelerator constraints including those from the radiative $b \rightarrow s \gamma$ decay. Appropriate grand unification (GUT) scenarios are considered. We find that the relaxation of gaugino mass unification is sufficient to obtain a phenomenologically and cosmologically viable solution of the MSSM with a neutralino as light as 3 GeV. We find good prospects for the direct detection of these superlight CDM neutralinos via elastic scattering off various nuclei in the forthcoming experiments with low-threshold DM detectors. In a certain sense, these experiments can probe the gaugino mass unification giving constraints on the possible GUT scenarios within the MSSM.

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I. INTRODUCTION

The minimal supersymmetric standard model (MSSM) [1] is a leading candidate for a low energy theory consistent with the grand unified theory (GUT) idea. The gauge coupling constants precisely measured at the CERN e^+e^- collider LEP make unification in the standard model (SM) rather problematic while in the MSSM it occurs naturally with an excellent precision.

The MSSM, supplied *a priori* with a complete set of the grand unification conditions, possesses a remarkable predictive power. The complete set of the GUT conditions includes gauge coupling constants unification as well as unification of the “soft” supersymmetry- (SUSY-) breaking parameters at the same GUT scale $M_X \sim 10^{16}$ GeV. Instead of this ultimate GUT scenario one can consider less restrictive particular GUT scenarios relaxing some of the GUT conditions. We have no firm theoretical arguments yet in favor of one of these scenarios. Analyzing the prospects for discovering SUSY in various experiments it is more attractive to adopt a phenomenological low-energy approach and disregard certain GUT conditions. Following these arguments we will discuss several GUT scenarios in the present paper.

Another advantage of the MSSM is the prediction of a stable lightest supersymmetric particle (LSP)—neutralino

(χ). Now the neutralino is the best known cold dark matter (CDM) particle candidate [2].

There is a well-known lower limit on the neutralino mass $m_{\chi} \geq 18.4$ GeV [3,4]. This result is strongly connected to the unification scenario with the universal gaugino masses $M_i(M_X) = m_{1/2}$ at the GUT scale M_X . The renormalization group evolution of M_i , starting with the same $m_{1/2}$ value, leads to tight correlation between the neutralino m_{χ_i} (with $i=1-4$), chargino $m_{\chi_{1,2}^\pm}$ and gluino $m_{\tilde{g}}$ masses (see Sec. II). As a result, direct and indirect SUSY searches at Fermilab and at LEP [4,5] strongly disfavor the neutralinos lighter than 18.4 GeV within the universal gaugino mass scenario [4]. As discussed in the following, the nonuniversal gaugino mass scenario with nonequal gaugino masses at the GUT scale allows essentially lighter neutralinos. In a certain sense, direct searches for the “superlight” neutralinos in the mass range $m_{\chi} < 18.4$ GeV could be a test of the gaugino mass unification.

In the present paper we will consider the discovering potential of DM experiments searching for the superlight DM neutralinos via neutralino-nucleus elastic scattering. Since the nuclear recoil energy in this collision is $E_r \sim 10^{-6} m_{\chi}$, a DM detector should have the low threshold $E_r^{\text{thr}} \sim \text{few keV}$ to be sensitive to the DM neutralinos as light as $m_{\chi} \sim \text{few GeV}$.

There are several projects of DM experiments with low-threshold detectors which are able to probe this mass region [6]. Some of them are expected to run in the near future. These experiments will use either a new generation of cryogenic calorimeters [7–9] or germanium detectors of special configuration [10].

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We will argue that these experiments have good prospects for verification of the gaugino mass unification at the GUT scale. The basic reason follows from the results of our analysis of the MSSM parameter space within the nonuniversal gaugino mass unification scenarios. In these scenarios we have found superlight neutralinos with masses as small as 3 GeV which produce cosmologically viable relic density $0.025 \leq \Omega_{\tilde{\chi}} \leq 1$ and a substantial total event rate $R \sim 1$ event/kg day of elastic scattering from nuclei in a DM detector. These values of R are within the expected sensitivity of the above-mentioned low-threshold DM detectors. Therefore, the superlight DM neutralinos appearing in the nonuniversal gaugino mass GUT scenarios can be observed with these setups. Negative results of DM neutralino searches in the mass region $3 \text{ GeV} \leq m_{\tilde{\chi}} \leq 18.4 \text{ GeV}$ would discriminate these scenarios.

The paper is organized as follows. In Sec. II we specify the MSSM and give formulas used in the subsequent sections. In Sec. III we summarize the experimental and cosmological inputs for our analysis. Section IV is devoted to calculation of the event rate of the elastic neutralino-nucleus

scattering. In Sec. V we discuss the results of our numerical analysis and Sec. VI contains the conclusion.

II. MINIMAL SUPERSYMMETRIC STANDARD MODEL

The MSSM is completely specified by the standard $SU(3) \times SU(2) \times U(1)$ gauge couplings as well as by the low-energy superpotential and ‘‘soft’’ SUSY-breaking terms [1]. The most general gauge-invariant form of the R -parity-conserving superpotential is

$$\mathcal{W} = h_E L^j E^c H_1^i \epsilon_{ij} + h_D Q^j D^c H_1^i \epsilon_{ij} + h_U Q^j U^c H_2^i \epsilon_{ij} + \mu H_1^i H_2^i \epsilon_{ij} \quad (1)$$

($\epsilon_{12} = +1$). The following notation is used for the quark $Q(\bar{3}, 2, 1/6)$, $D^c(\bar{3}, 1, 1/3)$, $U^c(\bar{3}, 1, -2/3)$, lepton $L(1, 2, -1/2)$, $E^c(1, 1, 1)$ and Higgs boson $H_1(1, 2, -1/2)$, $H_2(1, 2, 1/2)$ chiral superfields with the $SU(3)_c \times SU(2)_L \times U(1)_Y$ assignment given in parentheses. Yukawa coupling constants $h_{E,D,U}$ are matrices in the generation space and nondiagonal in the general case. For simplicity we suppressed generation indices.

In general, the ‘‘soft’’ SUSY-breaking terms are given by [11]

$$\begin{aligned} \mathcal{L}_{\text{SB}} = & -\frac{1}{2} \sum_A M_A \bar{\lambda}_A \lambda_A - m_{H_1}^2 |H_1|^2 - m_{H_2}^2 |H_2|^2 - m_{\tilde{Q}}^2 |\tilde{Q}|^2 - m_{\tilde{D}}^2 |\tilde{D}^c|^2 - m_{\tilde{U}}^2 |\tilde{U}^c|^2 - m_{\tilde{L}}^2 |\tilde{L}|^2 - m_{\tilde{E}}^2 |\tilde{E}^c|^2 - (h_E A_E \tilde{L}^j \tilde{E}^c H_1^i \epsilon_{ij} \\ & + h_D A_D \tilde{Q}^j \tilde{D}^c H_1^i \epsilon_{ij} + h_U A_U \tilde{Q}^j \tilde{U}^c H_2^i \epsilon_{ij} + \text{H.c.}) - (B \mu H_1^i H_2^i \epsilon_{ij} + \text{H.c.}). \end{aligned} \quad (2)$$

As usual, $M_{3,2,1}$ are the masses of the $SU(3) \times U(2) \times U(1)$ gauginos, \tilde{g} , \tilde{W} , \tilde{B} and m_i are the masses of scalar fields. A_L , A_D , A_U , and B are trilinear and bilinear couplings.

Observable quantities can be calculated in terms of the gauge and the Yukawa coupling constants as well as the soft SUSY-breaking parameters and the Higgs mass parameter μ introduced in Eqs. (1) and (2).

Under the renormalization they depend on the energy scale Q according to the renormalization group equations (RGEs).

It is a common practice to implement the GUT conditions at the GUT scale M_X . It allows one to reduce the number of free parameters of the MSSM. A complete set of GUT conditions is

$$\begin{aligned} m_{\tilde{L}}(M_X) = m_{H_{1,2}}(M_X) = m_{\tilde{E}}(M_X) = m_{\tilde{Q}}(M_X) \\ = m_{\tilde{U}}(M_X) = m_{\tilde{D}}(M_X) = m_0, \end{aligned} \quad (3)$$

$$A_U(M_X) = A_D(M_X) = A_L(M_X) = A_0,$$

$$M_i(M_X) = m_{1/2}, \quad (4)$$

$$\begin{aligned} \alpha_i(M_X) = \alpha_{\text{GUT}}, \quad \text{where } \alpha_1 = \frac{5}{3} \frac{g'^2}{4\pi}, \quad \alpha_2 = \frac{g^2}{4\pi}, \\ \alpha_3 = \frac{g_s^2}{4\pi}, \end{aligned} \quad (5)$$

g' , g , and g_s are the $U(1)$, $SU(2)$, and $SU(3)$ gauge coupling constants.

The above GUT conditions put very stringent constraints on the weak scale particle spectrum and couplings. In this scenario the neutralino is heavier than 18.4 GeV, as mentioned in the Introduction. Because of specific correlations between the MSSM parameters at the weak scale, stemming from Eqs. (3)–(5), the detection rate of the DM neutralinos is very small, being typically beyond the realistic abilities of the present and future DM detectors.

However, at present we have no strong motivation to impose a complete set of the GUT conditions in Eqs. (3)–(5) on the MSSM.

Therefore, in our analysis we use relaxed versions of the above discussed GUT conditions and consider two scenarios with the nonuniversal gaugino masses $M_3 = M_2 \neq M_1$: (a) with the universal

$$m_{H_{1,2}}(M_X) = m_{\tilde{f}}(M_X) = m_0,$$

and (b) nonuniversal

$$m_{H_1}(M_X) = m_{H_2}(M_X) \neq m_0 \quad (6)$$

GUT scale scalar masses. Deviation from the GUT scale universality of the mass parameters in the scalar sector was previously considered in Refs. [12] and [13].

Accepting the GUT conditions, we end up with the following free MSSM parameters: the common gauge cou-

pling α_{GUT} ; the matrices of the Yukawa couplings h_i^{ab} , where $i=E, U, D$; the soft supersymmetry-breaking parameters $m_0, m_{1/2}, A_0, B$, the Higgs field mixing parameter μ , and an additional parameter of the Higgs sector m_A being the mass of the CP -odd neutral Higgs boson. Since the masses of the third generation are much larger than the masses of the first two ones, we consider only the Yukawa coupling of the third generation and drop the indices a, b .

Additional constraints follow from the minimization conditions of the scalar Higgs potential. Under these conditions the bilinear coupling B can be replaced in the given list of free parameters by the ratio $\tan \beta = v_2/v_1$ of the vacuum expectation values of the two Higgs doublets.

We calculate the weak-scale parameters in Eqs. (1) and (2) in terms of the above-listed free parameters on the basis of two-loop RGE's following the iteration algorithm developed in [14].

The Higgs potential V including the one-loop corrections ΔV can be written as

$$V(H_1^0, H_2^0) = m_1^2 |H_1^0|^2 + m_2^2 |H_2^0|^2 - m_3^2 (H_1^0 H_2^0 + \text{H.c.}) + \frac{g^2 + g'^2}{8} (|H_1^0|^2 - |H_2^0|^2)^2 + \Delta V,$$

$$\text{with } \Delta V = \frac{1}{64\pi^2} \sum_i (-1)^{2J_i} (2J_i + 1) C_i m_i^4 \left[\ln \frac{m_i^2}{Q^2} - \frac{3}{2} \right], \quad (7)$$

where the sum is taken over all possible particles with the spin J_i and with the color degrees of freedom C_i . The mass parameters of the potential are introduced in the usual way as

$$m_{1,2}^2 = m_{H_{1,2}}^2 + \mu^2, \quad m_3^2 = B\mu. \quad (8)$$

They are running parameters $m_i(Q)$ with the Q -scale dependence determined by the RGE. The one-loop potential (7) itself is Q independent up to, irrelevant for the symmetry breaking, field-independent term depending on Q . At the minimum of this potential the neutral components of the Higgs field acquire nonzero vacuum expectation values $\langle H_{1,2}^0 \rangle = v_{1,2}$ triggering the electroweak symmetry breaking with $g^2(v_1^2 + v_2^2) = 2M_W^2$.

The minimization conditions read

$$2m_1^2 = 2m_3^2 \tan \beta - M_Z^2 \cos 2\beta - 2\Sigma_1, \quad (9)$$

$$2m_2^2 = 2m_3^2 \cot \beta + M_Z^2 \cos 2\beta - 2\Sigma_2, \quad (10)$$

where $\Sigma_k \equiv \partial \Delta V / \partial \psi_k$, with $\psi_{1,2} = \text{Re } H_{1,2}^0$, are the one-loop corrections [15]

$$\Sigma_k = -\frac{1}{32\pi^2} \sum_i (-1)^{2J_i} (2J_i + 1) \frac{1}{\psi_k} \frac{\partial m_i^2}{\partial \psi_k} m_i^2 \left(\ln \frac{m_i^2}{Q^2} - 1 \right). \quad (11)$$

As a remnant of two Higgs doublets $H_{1,2}$ after the electroweak symmetry breaking there occur five physical Higgs particles: a CP -odd neutral Higgs boson A , CP -even neutral Higgs bosons H, h , and a pair of charged Higgs bosons H^\pm . Their masses $m_A, m_{h,H}, m_{H^\pm}$ can be calculated including all one-loop corrections as second derivatives of the Higgs potential in Eq. (7) with respect to the corresponding fields evaluated at the minimum [16,17].

The neutralino mass matrix written in the basis $(\tilde{B}, \tilde{W}^3, \tilde{H}_1^0, \tilde{H}_2^0)$ has the form

$$\mathcal{M}_{\tilde{\chi}} = \begin{pmatrix} M_1 & 0 & -M_Z \cos \beta \sin \theta_W & M_Z \sin \beta \sin \theta_W \\ 0 & M_2 & M_Z \cos \beta \cos \theta_W & -M_Z \sin \beta \cos \theta_W \\ -M_Z \cos \beta \sin \theta_W & M_Z \cos \beta \cos \theta_W & 0 & -\mu \\ M_Z \sin \beta \sin \theta_W & -M_Z \sin \beta \cos \theta_W & -\mu & 0 \end{pmatrix}. \quad (12)$$

The universal gaugino mass unification scenarios with the GUT condition in Eq. (4) imply the relation

$$M_1 = \frac{5}{3} \tan^2 \theta_W M_2. \quad (13)$$

Diagonalizing the mass matrix by virtue of the orthogonal matrix \mathcal{N} one can obtain four physical neutralinos χ_i with the field content

$$\chi_i = \mathcal{N}_{i1} \tilde{B} + \mathcal{N}_{i2} \tilde{W}^3 + \mathcal{N}_{i3} \tilde{H}_1^0 + \mathcal{N}_{i4} \tilde{H}_2^0 \quad (14)$$

and with masses m_{χ_i} being eigenvalues of the mass matrix (12). We denote the lightest neutralino χ_1 by χ . In our analysis χ is the lightest SUSY particle (LSP).

The chargino mass term is

$$(\tilde{W}^-, \tilde{H}_1^-) \mathcal{M}_{\tilde{\chi}^\pm} \begin{pmatrix} \tilde{W}^+ \\ \tilde{H}_2^+ \end{pmatrix} + \text{H.c.} \quad (15)$$

with the mass matrix

$$\mathcal{M}_{\tilde{\chi}^\pm} = \begin{pmatrix} M_2 & \sqrt{2} M_W \sin \beta \\ \sqrt{2} M_W \cos \beta & \mu \end{pmatrix}, \quad (16)$$

which can be diagonalized by the transformation

$$\begin{aligned} \tilde{\chi}^- &= U_{i1} \tilde{W}^- + U_{i2} \tilde{H}_1^-, \\ \tilde{\chi}^+ &= V_{i1} \tilde{W}^+ + V_{i2} \tilde{H}_2^+, \end{aligned} \quad (17)$$

with $U^* \mathcal{M}_{\tilde{\chi}^\pm} V^\dagger = \text{diag}(M_{\tilde{\chi}_1^\pm}, M_{\tilde{\chi}_2^\pm})$, where the chargino masses are

$$M_{\tilde{\chi}_{1,2}^\pm}^2 = \frac{1}{2} [M_2^2 + \mu^2 + 2M_W^2 \mp \sqrt{(M_2^2 - \mu^2)^2 + 4M_W^4 \cos^2 2\beta + 4M_W^2(M_2^2 + \mu^2 + 2M_2\mu \sin 2\beta)}].$$

It is seen from Eqs. (12) and (16) that in the universal gaugino mass scenarios, leading to relation (13), the neutralino and chargino sectors of the MSSM are strongly correlated since they are described by the same set of three parameters $M_2, \mu, \tan \beta$. Relaxing condition (13) we get one independent parameter in the neutralino sector, M_1 . In this case, taking place within the nonuniversal gaugino mass GUT scenarios, the neutralino mass can be appreciably smaller than 18.4 GeV. The latter is the lower neutralino mass bound following from the experimental data in the case of the universal gaugino mass.

The mass matrices for the three-generation sfermions \tilde{t} , \tilde{b} , and $\tilde{\tau}$ in the $\tilde{f}_L - \tilde{f}_R$ basis are

$$\mathcal{M}_{\tilde{t}}^2 = \begin{pmatrix} m_{\tilde{Q}}^2 + m_t^2 + \frac{1}{6} (4M_W^2 - M_Z^2) \cos 2\beta & m_t(A_t - \mu \cot \beta) \\ m_t(A_t - \mu \cot \beta) & m_{\tilde{U}}^2 + m_t^2 - \frac{2}{3} (M_W^2 - M_Z^2) \cos 2\beta \end{pmatrix},$$

$$\mathcal{M}_{\tilde{b}}^2 = \begin{pmatrix} m_{\tilde{Q}}^2 + m_b^2 - \frac{1}{6} (2M_W^2 + M_Z^2) \cos 2\beta & m_b(A_b - \mu \tan \beta) \\ m_b(A_b - \mu \tan \beta) & m_{\tilde{D}}^2 + m_b^2 + \frac{1}{3} (M_W^2 - M_Z^2) \cos 2\beta \end{pmatrix},$$

$$\mathcal{M}_{\tilde{\tau}}^2 = \begin{pmatrix} m_{\tilde{L}}^2 + m_\tau^2 - \frac{1}{2} (2M_W^2 - M_Z^2) \cos 2\beta & m_\tau(A_\tau - \mu \tan \beta) \\ m_\tau(A_\tau - \mu \tan \beta) & m_{\tilde{E}}^2 + m_\tau^2 + (M_W^2 - M_Z^2) \cos 2\beta \end{pmatrix}.$$

For simplicity, in the sfermion mass matrices we ignored nondiagonality in the generation space, which is important only for the $b \rightarrow s \gamma$ decay.

III. CONSTRAINED MSSM PARAMETER SPACE

In this section we summarize the theoretical and experimental constraints used in our analysis.

The solution of the gauge coupling constants unification [see Eq. (5)] allows us to define the unification scale M_X . Our numerical procedure is based on the two-loop RGE's. We use the world averaged values of the gauge couplings at the Z^0 energy obtained from a fit to the LEP data [18], M_W [4] and m_t [19,20]:

$$\alpha^{-1}(M_Z) = 128.0 \pm 0.1, \quad (18)$$

$$\sin^2 \theta_{\overline{MS}} = 0.2319 \pm 0.0004, \quad (19)$$

$$\alpha_3 = 0.125 \pm 0.005, \quad (20)$$

where \overline{MS} denotes the modified minimal subtraction scheme. The value of the fine structure constant $\alpha^{-1}(M_Z)$ was updated from [21] by using new data on the hadronic vacuum polarization [22]. The standard relations are implied:

$$\alpha_1 = \frac{5}{3} \frac{\alpha}{\cos^2 \theta_W}, \quad \alpha_2 = \frac{\alpha}{\sin^2 \theta_W}. \quad (21)$$

SUSY particles have not been found so far and from the searches at LEP one knows that the lower limit on the charged sleptons is half the Z^0 mass (45 GeV) [4] and the Higgs boson mass has to be above 60 GeV [23,24], while the sneutrinos have to be above 41 GeV [4]. For the charginos

the preliminary lower limit of 65 GeV was obtained from the LEP 140 GeV run [5]. The above mass limits are incorporated in our analysis.

Radiative corrections trigger spontaneous symmetry breaking in the electroweak sector. In this case the Higgs potential has its minimum for the nonzero vacuum expectation values of the Higgs fields $H_{1,2}^0$. Solving M_Z from Eqs. (9) and (10) yields

$$\frac{M_Z^2}{2} = \frac{m_1^2 + \Sigma_1 - (m_2^2 + \Sigma_2) \tan^2 \beta}{\tan^2 \beta - 1}, \quad (22)$$

where Σ_1 and Σ_2 are defined in Eq. (11). This is an important constraint which relates the true vacuum to the physical Z-boson mass $M_Z = 91.187 \pm 0.007$ GeV.

Another stringent constraint is imposed by the branching ratio $B(b \rightarrow s \gamma)$, measured by the CLEO Collaboration [25] to be $B(b \rightarrow s \gamma) = (2.32 \pm 0.67) \times 10^{-4}$.

In the MSSM this flavor-changing neutral current (FCNC) process receives contributions from $H^\pm - t$, $\tilde{\chi}^\pm - \tilde{t}$, and $\tilde{g} - \tilde{q}$ loops in addition to the SM $W - t$ loop. The $\tilde{\chi} - \tilde{t}$ loops, which are expected to be much smaller, have been neglected [26,27]. The $\tilde{g} - \tilde{q}$ loops are proportional to $\tan \beta$. It was found [14] that this contribution should be small, even in the case of large $\tan \beta$ and therefore it can be neglected. The chargino contribution, which becomes large for large $\tan \beta$ and small chargino masses, depends sensitively on the splitting of the two stop mass eigenstates $\tilde{t}_{1,2}$.

Within the MSSM the following ratio has been calculated [26]:

$$\frac{B(b \rightarrow s \gamma)}{B(b \rightarrow c e \bar{\nu})} = \frac{|V_{ts}^* V_{tb}|^2}{|V_{cb}|^2} K_{\text{NLO}}^{\text{QCD}} \frac{6\alpha}{\pi} \times \frac{[\eta^{16/23} A_\gamma + \frac{8}{3} (\eta^{14/23} - \eta^{16/23}) A_g + C]^2}{I(x_{cb}) [1 - (2/3\pi) \alpha_s(m_b) f(x_{cb})]},$$

where

$$C \approx 0.175, \quad I(x_{cb}) = 0.4847,$$

$$\eta = \alpha_s(M_W) / \alpha_s(m_b), \quad f(x_{cb}) = 2.41.$$

Here $f(x_{cb})$ represents corrections from leading order QCD to the known semileptonic $b \rightarrow c e \bar{\nu}$ decay rate; $I(x_{cb})$ is a phase space factor; $x_{cb} = m_c/m_b = 0.316$; $K_{\text{NLO}}^{\text{QCD}}$ describes the next-to-leading-order QCD corrections [28]. $A_{\gamma,g}$ are the coefficients of the effective operators for $bs - \gamma$ and $bs - g$ interactions, respectively; C describes mixing with the four-quark operators. The ratio of Cabibbo-Kobayashi-Maskawa (CKM) matrix elements $|V_{ts}^* V_{tb}|^2 / |V_{cb}|^2 = 0.95$ was taken from [29].

Assuming that the neutralinos form a dominant part of the DM in the universe one obtains a cosmological constraint on the neutralino relic density.

The present lifetime of the universe is at least 10^{10} years, which implies an upper limit on the expansion rate and correspondingly on the total relic abundance. Assuming $h_0 > 0.4$ one finds that the contribution of each relic particle species χ has to obey [30]

$$\Omega_\chi h_0^2 < 1, \quad (23)$$

where the relic density parameter $\Omega_\chi = \rho_\chi / \rho_c$ is the ratio of the relic neutralino mass density ρ_χ to the critical one $\rho_c = 1.88 \times 10^{-29} h_0^2 \text{g cm}^{-3}$.

We calculate $\Omega_\chi h_0^2$ following the standard procedure on the basis of the approximate formula

$$\Omega_\chi h_0^2 = 2.13 \times 10^{-11} \left(\frac{T_\chi}{T_\gamma} \right)^3 \left(\frac{T_\gamma}{2.7K} \right)^3 N_F^{1/2} \left(\frac{\text{GeV}^{-2}}{ax_F + bx_F^2/2} \right). \quad (24)$$

Here T_γ is the present day photon temperature, T_χ/T_γ is the reheating factor, $x_F = T_F/m_\chi \approx 1/20$, T_F is the neutralino freeze-out temperature, and N_F is the total number of degrees of freedom at T_F . The coefficients a, b are determined from the nonrelativistic expansion

$$\langle \sigma_{\text{annih}} v \rangle \approx a + bx \quad (25)$$

of the thermally averaged cross section of neutralino annihilation. We adopt an approximate treatment ignoring complications, which appear when the expansion (25) fails [31]. In our analysis all possible channels of the $\chi - \chi$ annihilation are taken into account. The complete list of the formulas, which we used, for the coefficients a, b and numerical values for the other parameters in Eqs. (24) and (25) can be found in [32].

Since the neutralinos are mixtures of gauginos and Higgsinos, the annihilation can occur via s -channel exchange of the Z^0 , Higgs bosons, and t -channel exchange of a scalar particle, like a selectron [33]. Therefore, the cosmological constraint in Eq. (23) substantially restricts the MSSM parameter space, as discussed by many groups [2,23,34,35].

In the analysis we ignore possible rescaling of the local neutralino density ρ which may occur in the region of the MSSM parameter space where $\Omega_\chi h^2 < 0.025$ [36,37]. At lower relic densities DM neutralinos cannot saturate even galactic halos in the universe and the presence of additional DM components should be taken into account. One may assume that it can be done by virtue of the above-mentioned rescaling ansatz. Let us note that the halo density is a very uncertain quantity. Its actual value can be one order of magnitude smaller (or larger) than the quoted value 0.025 [38]. The SUSY solution of the DM problem with such low neutralino density becomes questionable. Therefore, we simply skip the corresponding domains of the MSSM parameter space as cosmologically uninteresting.

Thus we assume neutralinos to be a dominant component of the DM halo of our galaxy with a density $\rho_\chi = 0.3 \text{ GeV cm}^{-3}$ in the solar vicinity.

IV. NEUTRALINO-NUCLEUS ELASTIC SCATTERING

A dark matter event is elastic scattering of a DM neutralino from a target nucleus producing a nuclear recoil which can be detected by a detector. The corresponding event rate depends on the distribution of the DM neutralinos in the solar vicinity and the cross section $\sigma_{\text{el}}(\chi^A)$ of neutralino-nucleus elastic scattering. In order to calculate $\sigma_{\text{el}}(\chi^A)$ one should specify neutralino-quark interactions. The relevant low-energy effective Lagrangian can be written in a general form as

$$L_{\text{eff}} = \sum_q \left(\mathcal{A}_q \cdot \bar{\chi} \gamma_\mu \gamma_5 \chi \cdot \bar{q} \gamma^\mu \gamma_5 q + \frac{m_q}{M_W} \cdot \mathcal{C}_q \cdot \bar{\chi} \chi \cdot \bar{q} q \right) + \mathcal{O}\left(\frac{1}{m_q^4}\right), \quad (26)$$

where terms with vector and pseudoscalar quark currents are omitted being negligible in the case of nonrelativistic DM neutralinos with typical velocities $v_\chi \approx 10^{-3} c$.

In the Lagrangian (26) we also neglect terms which appear in supersymmetric models at the order of $1/m_q^4$ and higher, where m_q^4 is the mass of the scalar superpartner \tilde{q} of the quark q . These terms, as pointed out in Ref. [32], are potentially important in the spin-independent neutralino-nucleon scattering, especially in the domains of the MSSM parameter space where $m_{\tilde{q}}$ is close to the neutralino mass m_χ . Below we adopt the approximate treatment of these terms proposed in [32], which allows ‘‘effectively’’ absorbing them into the coefficients \mathcal{C}_q in a wide region of the SUSY model parameter space. Our formulas for the coefficients \mathcal{A}_q and \mathcal{C}_q of the effective Lagrangian take into account squark mixing $\tilde{q}_L - \tilde{q}_R$ and the contribution of both CP -even Higgs bosons h, H :

$$\begin{aligned}
\mathcal{A}_q = & -\frac{g^2}{4M_W^2} \left[\frac{\mathcal{N}_{14}^2 - \mathcal{N}_{13}^2}{2} T_3 - \frac{M_W^2}{m_{\tilde{q}_1}^2 - (m_\chi + m_q)^2} (\cos^2 \theta_q \phi_{qL}^2 + \sin^2 \theta_q \phi_{qR}^2) - \frac{M_W^2}{m_{\tilde{q}_2}^2 - (m_\chi + m_q)^2} (\sin^2 \theta_q \phi_{qL}^2 \right. \\
& + \cos^2 \theta_q \phi_{qR}^2) - \frac{m_q^2}{4} P_q^2 \left(\frac{1}{m_{\tilde{q}_1}^2 - (m_\chi + m_q)^2} + \frac{1}{m_{\tilde{q}_2}^2 - (m_\chi + m_q)^2} \right) - \frac{m_q}{2} M_W P_q \sin 2\theta_q T_3 (\mathcal{N}_{12} - \tan \theta_W \mathcal{N}_{11}) \\
& \left. \times \left(\frac{1}{m_{\tilde{q}_1}^2 - (m_\chi + m_q)^2} - \frac{1}{m_{\tilde{q}_2}^2 - (m_\chi + m_q)^2} \right) \right], \tag{27}
\end{aligned}$$

$$\begin{aligned}
\mathcal{C}_q = & -\frac{g^2}{4} \left[\frac{F_h}{m_h^2} h_q + \frac{F_H}{m_H^2} H_q + P_q \left(\frac{\cos^2 \theta_q \phi_{qL} - \sin^2 \theta_q \phi_{qR}}{m_{\tilde{q}_1}^2 - (m_\chi + m_q)^2} - \frac{\cos^2 \theta_q \phi_{qR} - \sin^2 \theta_q \phi_{qL}}{m_{\tilde{q}_2}^2 - (m_\chi + m_q)^2} \right) + \sin 2\theta_q \left(\frac{m_q}{4M_W} P_q^2 - \frac{M_W}{m_q} \phi_{qL} \phi_{qR} \right) \right] \\
& \times \left(\frac{1}{m_{\tilde{q}_1}^2 - (m_\chi + m_q)^2} - \frac{1}{m_{\tilde{q}_2}^2 - (m_\chi + m_q)^2} \right). \tag{28}
\end{aligned}$$

Here

$$F_h = (\mathcal{N}_{12} - \mathcal{N}_{11} \tan \theta_W) (\mathcal{N}_{14} \cos \alpha_H + \mathcal{N}_{13} \sin \alpha_H),$$

$$F_H = (\mathcal{N}_{12} - \mathcal{N}_{11} \tan \theta_W) (\mathcal{N}_{14} \sin \alpha_H - \mathcal{N}_{13} \cos \alpha_H),$$

$$h_q = \left(\frac{1}{2} + T_3 \right) \frac{\cos \alpha_H}{\sin \beta} - \left(\frac{1}{2} - T_3 \right) \frac{\sin \alpha_H}{\cos \beta},$$

$$H_q = \left(\frac{1}{2} + T_3 \right) \frac{\sin \alpha_H}{\sin \beta} + \left(\frac{1}{2} - T_3 \right) \frac{\cos \alpha_H}{\cos \beta},$$

$$\phi_{qL} = \mathcal{N}_{12} T_3 + \mathcal{N}_{11} (Q - T_3) \tan \theta_W,$$

$$\phi_{qR} = \tan \theta_W Q \mathcal{N}_{11},$$

$$P_q = \left(\frac{1}{2} + T_3 \right) \frac{\mathcal{N}_{14}}{\sin \beta} + \left(\frac{1}{2} - T_3 \right) \frac{\mathcal{N}_{13}}{\cos \beta}.$$

The above formulas coincide with the relevant formulas in [32] neglecting the terms $\sim 1/m_{\tilde{q}}^4$ and higher. These terms are taken into account ‘‘effectively’’ by introducing an ‘‘effective’’ stop quark \tilde{t} propagator.

A general representation of the differential cross section of neutralino-nucleus scattering can be given in terms of three spin-dependent $\mathcal{F}_{ij}(q^2)$ and one spin-independent $\mathcal{F}_S(q^2)$ form factors [39]:

$$\begin{aligned}
\frac{d\sigma}{dq^2}(v, q^2) = & \frac{8G_F}{v^2} [a_0^2 \cdot \mathcal{F}_{00}^2(q^2) + a_0 a_1 \cdot \mathcal{F}_{10}^2(q^2) \\
& + a_1^2 \cdot \mathcal{F}_{11}^2(q^2) + c_0^2 \cdot A^2 \mathcal{F}_S^2(q^2)], \tag{29}
\end{aligned}$$

where v is a projectile neutralino velocity and q is the momentum transferred to the nucleus. The last term corresponding to the spin-independent scalar interaction gains coherent enhancement A^2 (A is the atomic weight of a nucleus in the reaction). The coefficients a_0, c_0 do not depend on the nuclear structure and relate to the parameters $\mathcal{A}_q, \mathcal{C}_q$ of the effective Lagrangian (26) as well as to the parameters $\Delta q, f_s, \hat{f}$ characterizing the nucleon structure. One has the relationships

$$a_0 = (\mathcal{A}_u + \mathcal{A}_d)(\Delta u + \Delta d) + 2\Delta s \mathcal{A}_s,$$

$$a_1 = (\mathcal{A}_u - \mathcal{A}_d)(\Delta u - \Delta d),$$

$$c_0 = \hat{f} \frac{m_u \mathcal{C}_u + m_d \mathcal{C}_d}{m_u + m_d} + f_s \mathcal{C}_s + \frac{2}{27} (1 - f_s - \hat{f})(\mathcal{C}_c + \mathcal{C}_b + \mathcal{C}_t). \tag{30}$$

Here $\Delta q^{p(n)}$ are the fractions of the proton (neutron) spin carried by the quark q . The standard definition is

$$\langle p(n) | \bar{q} \gamma^\mu \gamma_5 q | p(n) \rangle = 2S_{p(n)}^\mu \Delta q^{p(n)}, \tag{31}$$

where $S_{p(n)}^\mu = (0, \vec{S}_{p(n)})$ is the four-spin of the nucleon. The parameters $\Delta q^{p(n)}$ can be extracted from the data on polarized nucleon structure functions [40,41] and the hyperon semileptonic decay data [42].

We use Δq values extracted both from the EMC [40] and SMC [41] data. The other nuclear structure parameters f_s and \hat{f} in formula (30) are defined as

$$\langle p(n) | (m_u + m_d)(\bar{u}u + \bar{d}d) | p(n) \rangle = 2\hat{f} M_{p(n)} \bar{\Psi} \Psi,$$

$$\langle p(n) | m_s \bar{s}s | p(n) \rangle = f_s M_{p(n)} \bar{\Psi} \Psi. \tag{32}$$

The values extracted from the data under certain theoretical assumptions are [43]

$$\hat{f} = 0.05 \quad \text{and} \quad f_s = 0.14. \tag{33}$$

The strange quark contribution f_s is known to be uncertain to about a factor of 2. Therefore we take its value within the interval $0.07 < f_s < 0.3$ [43,44].

The nuclear structure comes into play via the form factors $\mathcal{F}_{ij}(q^2), \mathcal{F}_S(q^2)$ in Eq. (29). The spin-independent form factor $\mathcal{F}_S(q^2)$ can be represented as the normalized Fourier transform of a spherical nuclear ground state density distribution $\rho(\mathbf{r})$. We use the standard Woods-Saxon inspired distribution [45]. It leads to the form factor

$$\mathcal{F}_S(q^2) = \int d^3\mathbf{r} \rho(\mathbf{r}) e^{i\mathbf{r}\mathbf{q}} = 3 \frac{j_1(qR_0)}{qR_0} e^{-(1/2)(qs)^2}, \tag{34}$$

TABLE I. The MSSM parameters within the GUT scenarios (a), (b) and intervals of their variations in the present numerical analysis.

Scenario	Parameter	Lower and upper bounds	
(a), (b)	$\tan \beta$	1	50
(a), (b)	m_0	0 TeV	1 TeV
(a), (b)	A_0	-1 TeV	1 TeV
(a), (b)	μ	-1 TeV	1 TeV
(a), (b)	M_2	0 TeV	1 TeV
(a), (b)	M_1	-1 TeV	1 TeV
(b)	m_A	50 GeV	1 TeV

where $R_0 = (R^2 - 5s^2)^{1/2}$ and $s \approx 1$ fm are the radius and the thickness of a spherical nuclear surface; j_1 is the spherical Bessel function of index 1.

Spin-dependent form factors $\mathcal{F}_{ij}(q^2)$ are much more nuclear-model-dependent quantities. In the past few years noticeable progress in detailed nuclear-model calculations of these form factors has been achieved. For many nuclei of interest in the DM search they have been calculated within various models of nuclear structure [46,47]. Unfortunately, these calculations do not cover all isotopes which we are going to consider in the present paper. Therefore, we use a simple parameterization of the q^2 dependence of $\mathcal{F}_{ij}(q^2)$ in the form of a Gaussian with the rms spin radius of the nucleus calculated in the harmonic well potential [48]. For our purposes this semiempirical scheme is sufficient.

An experimentally observable quantity is the differential event rate per unit mass of the target material:

$$\frac{dR}{dE_r} = \left[N \frac{\rho_\chi}{m_\chi} \right] \int_{v_{\min}}^{v_{\max}} dv f(v) v \frac{d\sigma}{dq^2}(v, E_r), \quad (35)$$

where E_r is the nuclear recoil energy. The function $f(v)$ is the velocity distribution of neutralinos in the Earth's frame. In the galactic frame it is usually assumed to have an approximate Maxwellian form. $v_{\max} = v_{\text{esc}} \approx 600$ km/s and $\rho_\chi = 0.3$ GeV cm $^{-3}$ are the escape velocity and the mass density of the relic neutralinos in the solar vicinity; $v_{\min} = (M_A E_r / 2M_{\text{red}}^2)^{1/2}$ with M_A and M_{red} being the mass of nucleus A and the reduced mass of the neutralino-nucleus system, respectively. Note that $q^2 = 2M_A E_r$.

The differential event rate is the most appropriate quantity for comparing with the observed recoil spectrum. It allows one to properly take into account the spectral characteristics of a specific DM detector and to discriminate a background. However, discussing general problems and prospects of DM detection it is enough to consider the total event rate R integrated over the whole kinematical domain of recoil energy. Notice that this quantity is less sensitive to details of the nuclear structure than the differential event rate in Eq. (35). The q^2 shape of the form factors $\mathcal{F}_{ij}(q^2)$, $\mathcal{F}_S(q^2)$ in Eq. (29) may substantially change from one nuclear model to another. Integration over q^2 , as in the case of the total event rate R , reduces this model dependence.

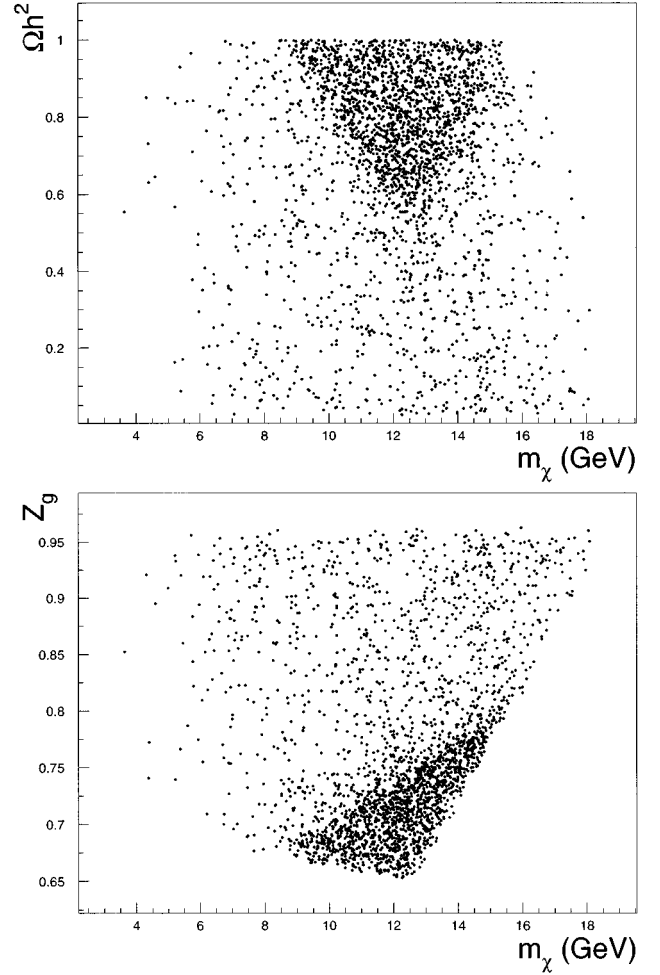


FIG. 1. The neutralino relic density Ωh^2 and the gaugino fraction $Z_g = \mathcal{N}_{11}^2 + \mathcal{N}_{12}^2$ versus the neutralino mass m_χ in GUT scenario (a).

The present paper aims at the general investigation of detectability of the cosmologically viable ‘‘superlight’’ DM neutralinos independently of a specific DM detector. Therefore, we may use the total event rate R as a characteristic of the DM signal.

V. NUMERICAL ANALYSIS AND DISCUSSION

In our numerical analysis we randomly scanned the MSSM parameter space including all experimental and cosmological constraints discussed in Sec. III. Two GUT scenarios with the nonuniversal gaugino mass $M_3 = M_2 \neq M_1$ [see Eqs. (3)–(5) and the comments below them] have been considered: (a) with the universal $m_{H_{1,2}}(M_X) = m_{\bar{f}}(M_X) = m_0$ and (b) nonuniversal $m_{H_1}(M_X) = m_{H_2}(M_X) \neq m_0$ GUT scale scalar masses.

The free parameters of both scenarios and intervals of their variation in our analysis are given in Table I.

After the scan procedure had been finished we found the superlight neutralinos in the mass range $3 \text{ GeV} \leq m_\chi \leq 18.4 \text{ GeV}$ in both GUT scenarios (a) and (b). The criterion we

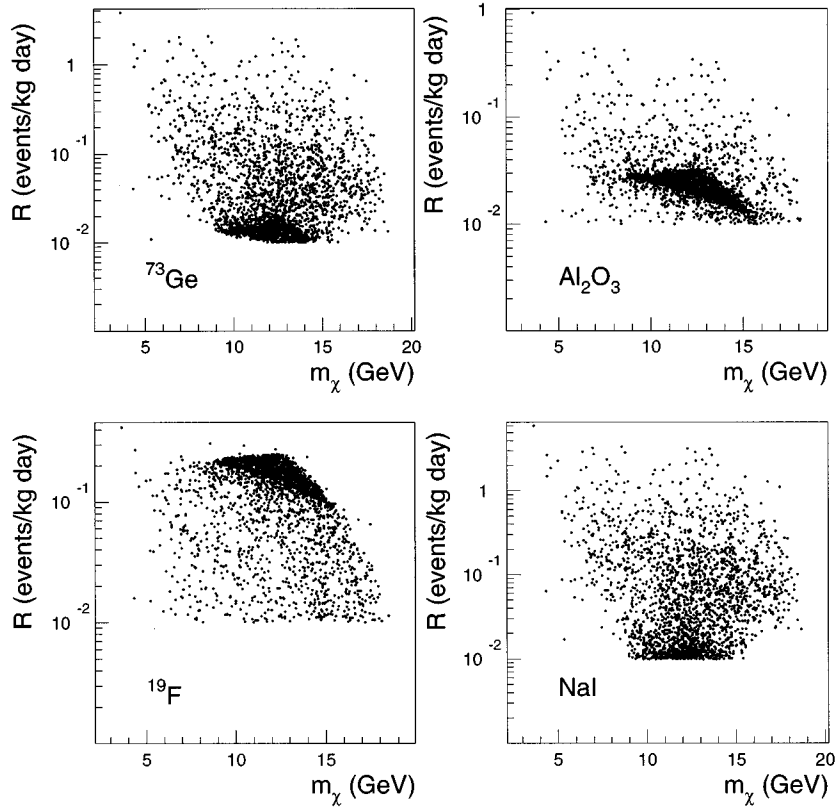


FIG. 2. The expected DM neutralino event rate R versus the neutralino mass m_χ for various isotopes of experimental interest. GUT scenario (a). All points with $R \leq 0.01$ events/kg day are cut off.

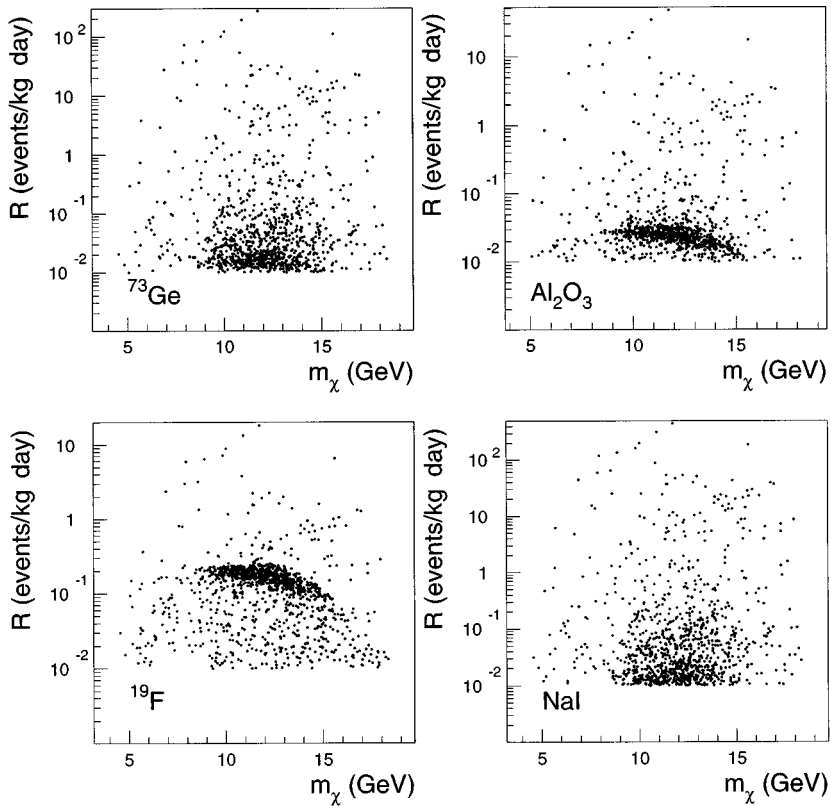


FIG. 3. The same as in Fig. 2 but for large event rate GUT scenario (b).

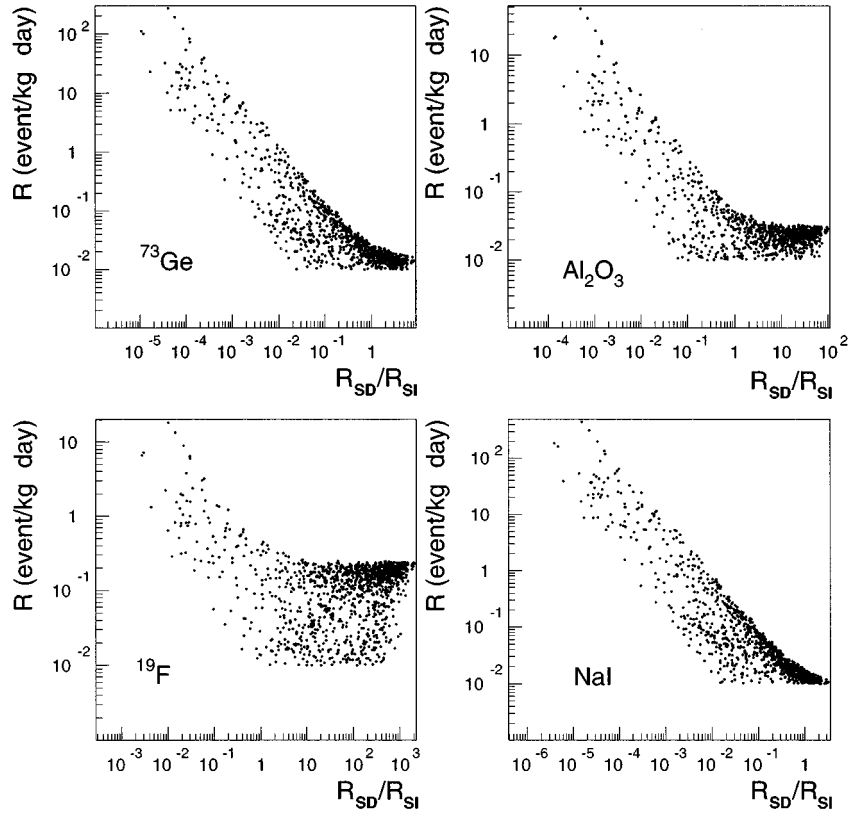


FIG. 4. Spin sensitivities of various isotopes of experimental interest. Given is the total event rate R in GUT scenario (b) versus the ratio R_{SD}/R_{SI} , R_{SD} and R_{SI} are the spin-dependent and spin-independent components of R ($R=R_{SD}+R_{SI}$). All points with $R \leq 0.01$ events/kg day are cut off.

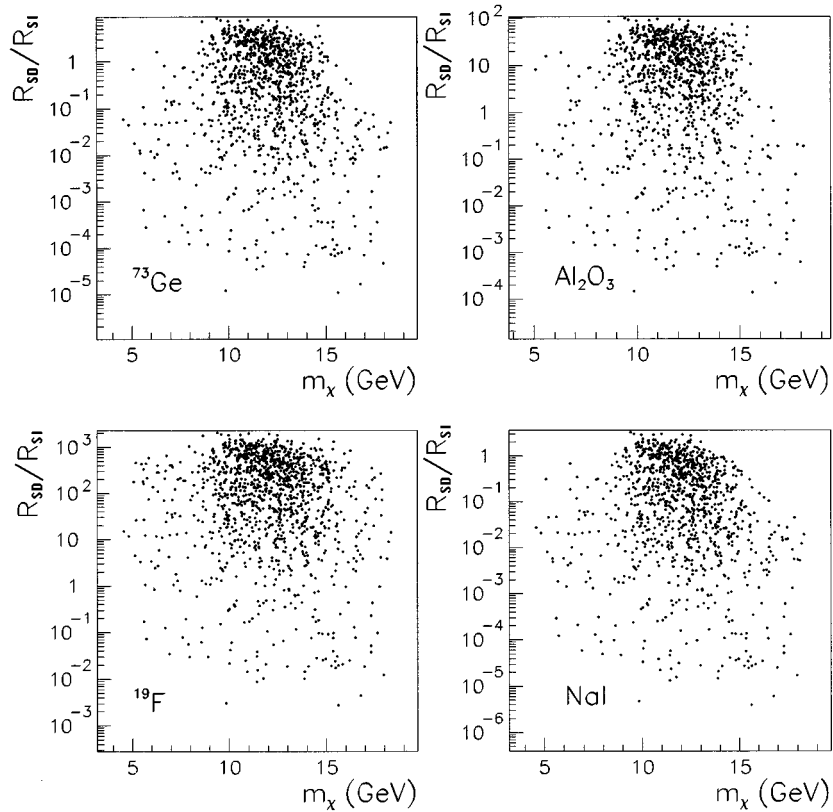


FIG. 5. The ratio R_{SD}/R_{SI} versus the neutralino mass m_χ . R_{SD} and R_{SI} are the spin-dependent and spin-independent components of R ($R=R_{SD}+R_{SI}$). Large event rate GUT scenario (b).

used to stop the running numerical code was stabilization of the m_χ lower bound.

Recall that 18.4 GeV is the widely cited lower neutralino mass bound valid within the complete GUT scenario with relations (3)–(5) at the GUT scale. Scenario (a) implies a minimal relaxation of the complete GUT scenario necessary to obtain the superlight neutralinos. Scenario (b) does not change the lower m_χ bound 3 GeV, but allows a much larger DM event rate R than scenario (a). We have calculated the total event rate for germanium (^{73}Ge), sapphire (Al_2O_3), fluorine (^{19}F), and sodium iodide (NaI).

The main results of our analysis are presented in Figs. 1–5 in the form of scatter plots. In Figs. 2–4 we cut off all points with $R \leq 0.01$ events/kg day since they lie beyond the sensitivity of the present and the near-future DM detectors.

Figure 1 shows the relic density, $\Omega_\chi h^2$, produced by the superlight neutralinos and their gaugino fraction, $Z_g = \mathcal{N}_{11}^2 + \mathcal{N}_{12}^2$ [see Eq. (14)]. It is seen that the majority of the points are concentrated in the region of a large relic density, especially within the interval $8 \text{ GeV} \leq m_\chi \leq 16 \text{ GeV}$. However, the neutralinos with a marginal mass of 3 GeV can also produce cosmologically interesting values of $\Omega_\chi h^2$. Therefore we conclude that the superlight neutralinos can comprise a dominant part of the CDM. The field composition of these neutralinos characterized by the gaugino fraction Z_g shows specific tendencies. A large domain of the MSSM parameter space (the volume is proportional to the number of points in the scatter plots) corresponds to the mixed $Z_g \sim 0.7$ gaugino-Higgsino states around the mass value 12 GeV. For larger m_χ neutralinos lose their Higgsino component and become mostly gauginos.

In Figs. 2 and 3 we present the calculated event rate, R , for ^{73}Ge , Al_2O_3 , ^{19}F , and NaI versus the neutralino mass m_χ in two GUT scenarios (a) and (b). These figures demonstrate that in both scenarios there are a lot of points within the reach of the near-future low-threshold DM detectors mentioned in the Introduction. Their sensitivities are expected to be at the level of $R \geq 0.1$ events/kg day. In the scenario (b) points extend up to $R \sim 300$ –500 events/kg day for Ge, NaI, and $R \sim 20$ –50 events/kg day for ^{19}F , Al_2O_3 . In scenario (a) these values are reduced by approximately a factor of 10.

In unification scenario (b) with nonuniversal Higgs boson mass parameters [see Eq. (6)] the Higgs boson and sfermion masses are not strongly correlated. As discussed in [13], this relaxation of the complete unification in the scalar sector makes it possible to avoid one of the most stringent theoretical limitations on the allowed values of the DM neutralino event rate. In this case the masses of the CP -even and charged Higgs bosons $m_{H,h}$, m_{H^\pm} are calculated in terms of the CP -odd Higgs mass, m_A , as an extra free parameter.

An important question touches upon the fraction of the spin-dependent part R_{SD} of the total event rate $R = R_{\text{SD}} + R_{\text{SI}}$. From Eqs. (27)–(30) it is seen that measurement of R_{SD} and R_{SI} would give us complementary information about the MSSM parameters. In the previously investigated mass region $m_\chi \geq 18.4 \text{ GeV}$ most of the experimentally interesting isotopes are more sensitive to the spin-independent part R_{SI} due to the coherent enhancement effect (see [13] and references therein). The spin-dependent part R_{SD} originates from the neutralino-nucleus interaction via spin-spin coupling. It is not a coherent interaction with a whole nucleus since only

a few nucleons contribute to the nuclear spin.

In Figs. 4 and 5 we give the scatter plots in $R - R_{\text{SD}}/R_{\text{SI}}$ and $R_{\text{SD}}/R_{\text{SI}} - m_\chi$ planes for large event rate scenario (b). The following conclusions are valid in scenario (a) as well. One can see that in the mass region of the superlight neutralinos $3 \text{ GeV} \leq m_\chi \leq 20 \text{ GeV}$ the coherent R_{SI} component of the event rate dominates for all isotopes in question, except the region of very small total event rates R . As seen from Fig. 5, the spin-dependent component R_{SD} dominates for the majority of the points in the mass range $8 \text{ GeV} \leq m_\chi \leq 14 \text{ GeV}$. It reflects the fact illustrated in Fig. 1 that in this mass region the neutralino can acquire the largest Higgsino admixture $1 - Z_g \approx 0.35$. As a result, the Z -boson contribution to R_{SD} [via the coefficient \mathcal{A}_q in Eq. (27)] is enhanced.

We conclude this discussion with the following remark. Nuclear spin does not play an important role in the direct searches for the DM neutralinos in the mass region $3 \text{ GeV} \leq m_\chi \leq 18.4 \text{ GeV}$. Previously a similar conclusion was obtained concerning the mass region $18.4 \text{ GeV} \leq m_\chi$ [13]. Nevertheless, only a DM detector with a spin-non-zero target nuclei can provide us with information about R_{SD} and the corresponding MSSM parameters in \mathcal{A}_q [see Eq. (27)].

VI. CONCLUSION

We have analyzed the MSSM parameter space taking into account cosmological and accelerator constraints including those from the radiative $b \rightarrow s \gamma$ decay.

It is well known that the MSSM with the universal gaugino mass at the GUT scale disfavors neutralinos lighter than 18.4 GeV [3,4] if the known experimental data are taken into account.

A central result of the present paper is the conclusion about the existence of a substantial domain of the MSSM parameter space corresponding to the superlight neutralinos in the mass range $3 \text{ GeV} \leq m_\chi \leq 18.4 \text{ GeV}$ within the GUT scenarios with the nonuniversal gaugino mass. In this domain neutralinos are cosmologically viable and produce an event rate which is detectable in the near-future experiments with the low-threshold DM detectors.

When our paper was being prepared, we found Ref. [49] where prospects for the superlight DM neutralino detection are also considered but within a more phenomenological approach ignoring all GUT conditions. Our analysis shows that we need not disregard all GUT conditions in Eqs. (3)–(5) to get a window for the superlight neutralinos. To this end it is necessary and sufficient to relax only gaugino mass unification condition. Since we know a generic root of the superlight neutralino we can conclude that in a certain sense the experiments searching for the DM neutralinos in the mass range $m_\chi \geq 18.4 \text{ GeV}$ probe the gaugino mass unification at the GUT scale.

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