Direct dark matter detection in the next-to-minimal supersymmetric standard model

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Direct dark matter detection is considered in the next-to-minimal supersymmetric standard model (NMSSM). The effective neutralino-quark Lagrangian is obtained and event rates are calculated for the ^{73}Ge isotope. Accelerator and cosmological constraints on the NMSSM parameter space are included. By means of scanning the parameter space at the Fermi scale we show that the lightest neutralino could be detected in dark matter experiments with a sizable event rate.
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I. INTRODUCTION

In not too far in the future new very sensitive dark matter (DM) detectors [1–3] may start to operate, and one expects new, very important data from these experiments. The future experimental progress forces investigators to better understand the variety and property of the dark matter particles. The lightest supersymmetric particle (LSP), the neutralino, is considered now as a most promising candidate, which may compose the main fraction of the so-called cold dark matter. The prospects of direct and indirect detection of the LSP have comprehensively been investigated [4] in the various versions of the minimal supersymmetric standard model (MSSM) [5].

In this paper we consider direct detection of this relic LSP in the next-to-minimal supersymmetric standard model (NMSSM) [6,7]. The Higgs sector of the NMSSM contains five physical neutral Higgs bosons, three Higgs scalars, two pseudoscalars, and two degenerate physical charged Higgs particles $C^{\pm}$. The neutralino sector is extended to five neutralinos instead of four in the MSSM. The remaining particle content is identical with that of the MSSM.

The NMSSM is mainly motivated by its potential to eliminate the so-called μ problem of the MSSM [8], where the origin of the μ parameter in the superpotential $W_{\rm MSSM}=\mu H_1 H_2$ is not understood. For phenomenological reasons it has to be of the order of the electroweak scale, while the ''natural'' mass scale would be of the order of the grand unified theory (GUT) or Planck scale. This problem is evaded in the NMSSM where the μ term in the superpotential is dynamically generated through $\mu=\lambda x$ with a dimensionless coupling λ and the vacuum expectation value x of the Higgs singlet. Another essential feature of the NMSSM is the fact that the mass bounds for the Higgs bosons and neutralinos are weakened. While in the MSSM experimental data imply a lower mass bound of about 20 GeV for the LSP [9], very light or massless neutralinos and Higgs bosons are not excluded in the NMSSM [10,11]. Furthermore the upper tree level mass bound for the lightest Higgs scalar of the MSSM

$$m_h^2 \leqslant m_Z^2 \cos^2 2\beta \tag{1}$$

is increased to $m_{S_1}^2 \leqslant m_Z^2 \cos^2 2\beta + \lambda^2 (v_1^2+v_2^2)\sin^2 2\beta$. Taking into account the weak coupling of the Higgs singlet the NMSSM may still remain a viable model when the MSSM can be ruled out due to Eq. (1).

The above arguments make an intensive study of the NMSSM phenomenology very desirable. Previously the Higgs and neutralino sectors of the NMSSM were carefully studied in [10–15]. The calculation of the LSP relic abundance in the NMSSM was performed for the first time in [16] and recently in [17].

The outline of this paper is as follows. In Sec. II we describe the Lagrangian of the NMSSM. Since the additional singlet superfield of the NMSSM leads to extended Higgs and neutralino sectors, we present the Higgs and neutralino mixings. Section III collects formulas relevant for calculation of the event rate for direct dark matter detection in the framework of the NMSSM. In Sec. IV we discuss the constraints on the NMSSM parameter space which are used in our analysis. In Sec. V we shortly describe our numerical procedure and discuss the results obtained. Section VI contains a conclusion.

II. THE LAGRANGIAN OF THE NMSSM

The NMSSM superpotential is [12] ($\varepsilon_{12}=-\varepsilon_{21}=1$)

$$W=\lambda \varepsilon_{ij} H_1^i H_2^j N - \frac{1}{3} k N^3 + h_u \varepsilon_{ij} \tilde{Q}^i \tilde{U} H_2^j - h_d \varepsilon_{ij} \tilde{Q}^i \tilde{D} H_1^j - h_e \varepsilon_{ij} \tilde{L}^i \tilde{R} H_1^j, \tag{2}$$

where $H_1=(H_1^0,H^-)$ and $H_2=(H^+,H_2^0)$ are the SU(2) Higgs doublets with hypercharge $-1/2$ and 1/2 and N is the Higgs singlet with hypercharge 0. The notation of the fermion doublets and singlets is conventional, generation indices are omitted. Contrary to the MSSM, the superpotential of the NMSSM consists only of trilinear terms with dimensionless couplings.

The electroweak gauge-symmetry $SU(2)_I \times U(1)_Y$ is spontaneously broken to the electromagnetic gauge-symmetry $U(1)_{em}$ by the Higgs VEVs $\langle H_i^0 \rangle = v_i$ with $i=1,2$ and $\langle N \rangle = x$, where $v=\sqrt{v_1^2+v_2^2}=174$ GeV, $\tan\beta = v_2/v_1$.

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The most general supersymmetry breaking potential can be written as [12]

$$\begin{aligned}V_{\rm soft}&=m_1^2|H_1|^2+m_2^2|H_2|^2+m_3^2|N|^2\\&\quad+m_Q^2|\tilde{Q}|^2+m_U^2|\tilde{U}|^2+m_D^2|\tilde{D}|^2+m_L^2|\tilde{L}|^2+m_E^2|\tilde{R}|^2\\&\quad-(\lambda A_\lambda\varepsilon_{ij}H_1^iH_2^jN+\text{H.c.})-\left(\frac{1}{3}kA_kN^3+\text{H.c.}\right)\\&\quad+(h_uA_U\varepsilon_{ij}\tilde{Q}^i\tilde{U}H_2^j-h_dA_D\varepsilon_{ij}\tilde{Q}^i\tilde{D}H_1^j\\&\quad-h_eA_E\varepsilon_{ij}\tilde{L}^i\tilde{R}H_1^j+\text{H.c.})+\frac{1}{2}M\lambda^a\lambda^a+\frac{1}{2}M'\lambda'\lambda'.\end{aligned} \tag{3}$$

As free parameters appear the ratio of the doublet vacuum expectation values, $\tan\beta$, the singlet vacuum expectation value x, the couplings in the superpotential λ and k, the parameters A_λ, A_k, as well as A_U, A_D, A_E (for three generations) in the supersymmetry breaking potential, the gaugino mass parameters M and M', and the scalar mass parameters for the Higgs bosons $m_{1,2,3}$, squarks $m_{Q,U,D}$ and sleptons $m_{L,E}$.

The minimization conditions for the scalar potential $\partial V/\partial v_{1,2}=0$, $\partial V/\partial x=0$ eliminate three parameters of the Higgs sector which are normally chosen to be m_1^2, m_2^2, and m_3^2. Then at the tree level the elements of the symmetric CP-even mass squared matrix $\mathcal{M}_S^2=(M_{ij}^{S^2})$ become, in the basis (H_1,H_2,N),

$$M_{11}^{S^2}=\frac{1}{2}v_1^2(g'^2+g^2)+\lambda x\tan\beta(A_\lambda+kx),$$

$$M_{12}^{S^2}=-\lambda x(A_\lambda+kx)+v_1v_2\left(2\lambda^2-\frac{1}{2}g'^2-\frac{1}{2}g^2\right)$$

$$M_{13}^{S^2}=2\lambda^2v_1x-2\lambda kxv_2-\lambda A_\lambda v_2,$$

$$M_{22}^{S^2}=\frac{1}{2}v_2^2(g'^2+g^2)+\lambda x\cot\beta(A_\lambda+kx),$$

$$M_{23}^{S^2}=2\lambda^2v_2x-2\lambda kxv_1-\lambda A_\lambda v_1,$$

$$M_{33}^{S^2}=4k^2x^2-kA_kx+\frac{\lambda A_\lambda v_1v_2}{x}.$$

In the same way one finds, for the elements of the CP-odd matrix $\mathcal{M}_P^2$,

$$M_{11}^{P^2}=\lambda x(A_\lambda+kx)\tan\beta,\quad M_{12}^{P^2}=\lambda x(A_\lambda+kx),$$

$$M_{13}^{P^2}=\lambda v_2(A_\lambda-2kx),\quad M_{22}^{P^2}=\lambda x(A_\lambda+kx)\cot\beta,$$

$$M_{23}^{P^2}=\lambda v_1(A_\lambda-2kx),$$

$$M_{33}^{P^2}=\lambda A_\lambda\frac{v_1v_2}{x}+4\lambda kv_1v_2+3kA_kx,$$

and for the charged Higgs matrix one obtains

$$\mathcal{M}_c^2=\left(\lambda A_\lambda x+\lambda kx^2-v_1v_2\left(\lambda^2-\frac{g^2}{2}\right)\right)\begin{pmatrix}\tan\beta & 1\\ 1 & \cot\beta\end{pmatrix}.$$

In our numerical analysis we have included 1-loop radiative corrections to Higgs mass matrices following [14,18].

Assuming CP conservation in the Higgs sector, the Higgs matrices are diagonalized by the real orthogonal 3×3 matrices U^S and U^P, respectively,

$$\text{Diag}(m_{S_1}^2,m_{S_2}^2,m_{S_3}^2)=U^{S^T}\mathcal{M}_S^2U^S,$$

$$\text{Diag}(m_{P_1}^2,m_{P_2}^2,0)=U^{P^T}\mathcal{M}_P^2U^P,$$

where $m_{S_1}<m_{S_2}<m_{S_3}$ and $m_{P_1}<m_{P_2}$ denote the masses of the mass eigenstates of the neutral scalar Higgs bosons S_a $(a=1,2,3)$ and neutral pseudoscalar Higgs bosons P_α $(\alpha=1,2)$ [12].

With fixed parameters of the Higgs sector the masses and mixings of the neutralinos are determined by the two furthest parameters M and M' of the Lagrangian

$$\mathcal{L}=-\frac{1}{2}\Psi^TM\Psi+\text{H.c.},$$

$$\Psi^T=(-i\lambda_1,-i\lambda_2^3,\Psi_{H_1}^0,\Psi_{H_2}^0,\Psi_N).$$

In this basis the symmetric mass matrix M of the neutralinos has the form

$$\begin{pmatrix}M' & 0 & -m_Z\sin\theta_W\cos\beta & m_Z\sin\theta_W\sin\beta & 0\\ 0 & M & m_Z\cos\theta_W\cos\beta & -m_Z\cos\theta_W\sin\beta & 0\\ -m_Z\sin\theta_W\cos\beta & m_Z\cos\theta_W\cos\beta & 0 & \lambda x & \lambda v_2\\ m_Z\sin\theta_W\sin\beta & -m_Z\cos\theta_W\sin\beta & \lambda x & 0 & \lambda v_1\\ 0 & 0 & \lambda v_2 & \lambda v_1 & -2kx\end{pmatrix}.$$

The mass of the neutralinos is obtained by diagonalizing the mass matrix M with the orthogonal matrix N:

$$\mathcal{L}=-\frac{1}{2}m_i\overline{\tilde{\chi}_i^0}\tilde{\chi}_i^0, \quad \tilde{\chi}_i^0=\begin{pmatrix}\chi_i^0\\ \overline{\chi_i^0}\end{pmatrix}$$

$$\text{with } \chi_i^0=\mathcal{N}_{ij}\Psi_j \quad \text{and} \quad M_{\rm diag}=NMN^T.$$

The neutralinos $\tilde{\chi}_i^0$ $(i=1-5)$ are ordered with increasing mass $|m_i|$, thus $\chi\equiv\tilde{\chi}_1^0$ is the LSP neutralino. The matrix elements $\mathcal{N}_{ij}$ $(i,j=1-5)$ describe the composition of the neutralino $\tilde{\chi}_i^0$ in the basis Ψ_j. For example the bino fraction of the lightest neutralino is given by $\mathcal{N}_{11}^2$ and the singlino fraction of this neutralino by $\mathcal{N}_{15}^2$.

III. 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 suitable detector. The corresponding event rate depends on the distribution of the DM neutralinos in the solar vicinity and the cross section of neutralino-nucleus elastic scattering.

The relevant low-energy effective neutralino-quark Lagrangian can be written in a general form as [4,19–21]

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

where terms with vector and pseudoscalar quark currents are omitted being negligible in the case of non-relativistic DM neutralinos with typical velocities $v\approx10^{-3}c$.

The coefficients in the effective Lagrangian (4) have the form

$$\mathcal{A}_q=-\frac{g_2^2}{4M_W^2}\Bigg[\frac{\mathcal{N}_{14}^2-\mathcal{N}_{13}^2}{2}T_3-\frac{M_W^2}{m_{\tilde{q}1}^2-(m_\chi+m_q)^2}(\cos^2\theta_q\phi_{qL}^2+\sin^2\theta_q\phi_{qR}^2)-\frac{M_W^2}{m_{\tilde{q}2}^2-(m_\chi+m_q)^2}(\sin^2\theta_q\phi_{qL}^2+\cos^2\theta_q\phi_{qR}^2)$$
$$-\frac{m_q^2}{4}P_q^2\left(\frac{1}{m_{\tilde{q}1}^2-(m_\chi+m_q)^2}+\frac{1}{m_{\tilde{q}2}^2-(m_\chi+m_q)^2}\right)-\frac{m_q}{2}M_WP_q\sin2\theta_qT_3(\mathcal{N}_{12}-\tan\theta_W\mathcal{N}_{11})$$
$$\times\left(\frac{1}{m_{\tilde{q}1}^2-(m_\chi+m_q)^2}-\frac{1}{m_{\tilde{q}2}^2-(m_\chi+m_q)^2}\right)\Bigg] \tag{5}$$

$$\mathcal{C}_q=-\frac{g_2^2}{4}\Bigg[-\sum_{a=1,2,3}Q_{a11}^{L\prime\prime}\frac{1}{m_a^2}\mathcal{V}_{aq}+P_q\left(\frac{\cos^2\theta_q\phi_{qL}-\sin^2\theta_q\phi_{qR}}{m_{\tilde{q}1}^2-(m_\chi+m_q)^2}-\frac{\cos^2\theta_q\phi_{qR}-\sin^2\theta_q\phi_{qL}}{m_{\tilde{q}2}^2-(m_\chi+m_q)^2}\right)$$
$$+\sin2\theta_q\left(\frac{m_q}{4M_W}P_q^2-\frac{M_W}{m_q}\phi_{qL}\phi_{qR}\right)\times\left(\frac{1}{m_{\tilde{q}1}^2-(m_\chi+m_q)^2}-\frac{1}{m_{\tilde{q}2}^2-(m_\chi+m_q)^2}\right)\Bigg]. \tag{6}$$

Here

$$\mathcal{V}_{aq}=\left[\left(\frac{1}{2}+T_{3q}\right)\frac{U_{a2}^S}{\sin\beta}+\left(\frac{1}{2}-T_{3q}\right)\frac{U_{a1}^S}{\cos\beta}\right],$$

$$Q_{a11}^{L\prime\prime}=(\mathcal{N}_{12}-\tan\theta_W\mathcal{N}_{11})[U_{a1}^S\mathcal{N}_{13}-U_{a2}^S\mathcal{N}_{14}]+\sqrt{2}\lambda\mathcal{N}_{15}[U_{a1}^S\mathcal{N}_{14}+U_{a2}^S\mathcal{N}_{13}]-2\sqrt{2}kU_{a3}^S\mathcal{N}_{15}^2,$$

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

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

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

The coefficients $\mathcal{A}_q$ and $\mathcal{C}_q$ take into account squark mixing $\tilde{q}_L - \tilde{q}_R$ and the contributions of all CP-even Higgs bosons. Under the assumption $\lambda = k = 0$ these formulas coincide with the relevant formulas in the MSSM [19].

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 as follows [22]:

$$\frac{d\sigma}{dq^2}(v,q^2) = \frac{8G_F}{v^2}(a_0^2 \cdot \mathcal{F}_{00}^2(q^2) + a_0 a_1 \cdot \mathcal{F}_{10}^2(q^2) + a_1^2 \cdot \mathcal{F}_{11}^2(q^2) + c_0^2 \cdot A^2 \mathcal{F}_S^2(q^2)). \quad (7)$$

The last term corresponding to the spin-independent scalar interaction gains coherent enhancement A^2 (A is the atomic weight of the nucleus in the reaction). The coefficients $a_{0,1}, c_0$ do not depend on nuclear structure and relate to the parameters $\mathcal{A}_q, \mathcal{C}_q$ of the effective Lagrangian (4) and to parameters characterizing the nucleon structure. In what follows we use notations and definitions of our paper [23].

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

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

Here $f(v)$ is the velocity distribution of neutralinos in the earth's frame which is usually assumed to be a Maxwellian distribution in the galactic frame. N is the number density of the target nuclei. $v_{max} = v_{esc} \approx 600$ 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; $v_{min} = (M_A E_r / 2M_{red}^2)^{1/2}$ with M_A and M_{red} being the mass of nucleus A and the reduced mass of the neutralino-nucleus system, respectively.

The differential event rate is the most appropriate quantity for comparing with the observed recoil spectrum and allows one to take properly into account spectral characteristics of a specific detector and to separate the background. However, in many cases the total event rate R integrated over the whole kinematic domain of the recoil energy is sufficient. It is widely employed in theoretical papers for estimating the prospects for DM detection, ignoring experimental complications which may occur on the way. In the present paper we are going to perform a general analysis aimed at searching for domains with large values of the event rate R like those reported in [24]. This is the reason why we use in the analysis the total event rate R.

IV. CONSTRAINTS ON THE NMSSM PARAMETER SPACE

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 [25] $\Omega_\chi h_0^2 < 1$, 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 $\rho_c = 1.88 \times 10^{-29} h_0^2$g cm^{-3}.

We calculate $\Omega_\chi h_0^2$ following the standard procedure on the basis of the approximate formula [26,27]

$$\Omega_\chi h_0^2 = 2.13 \times 10^{-11} \left(\frac{T_\chi}{T_\gamma} \right)^3 \left(\frac{T_\gamma}{2.7K^o} \right)^3 \times N_F^{1/2} \left(\frac{\text{GeV}^{-2}}{a x_F + b x_F^2/2} \right). \quad (8)$$

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 non-relativistic expansion $\langle \sigma_{ann.} v \rangle \approx a + bx$ of the thermally averaged cross section of neutralino annihilation in the NMSSM. We adopt an approximate treatment not taking into account complications, which occur when the expansion fails [28]. We take into account all possible channels of the χ-χ annihilation. The complete list of the relevant formulas in the NMSSM can be found in [17].

Since the neutralinos are mixtures of gauginos, higgsinos, and singlino the annihilation can occur both, via s-channel exchange of the Z^0 and Higgs bosons and t-channel exchange of a scalar particle, like a selectron. This constrains the parameter space, as discussed by many groups [27,29,30].

In the analysis we ignore possible rescaling of the local neutralino density ρ which may occur in the region of the NMSSM parameter space where $\Omega_\chi h_0^2 < 0.025$ [31–33]. If the neutralino is accepted as a dominant part of the DM its density has to exceed the quoted limiting value 0.025. Otherwise the presence of additional DM components should be taken into account, for instance, by the mentioned rescaling ansatz. However, the halo density is known to be very uncertain. Therefore, one can expect that the rescaling takes place in a small domain of the parameter space. Another point is that the SUSY solution of the DM problem with such low neutralino density becomes questionable. We assume neutralinos to be a dominant component of the DM halo of our galaxy with a density $\rho_\chi = 0.3$ GeV cm^{-3} in the solar vicinity and disregard in the analysis points with $\Omega_\chi h_0^2 < 0.025$.

The parameter space of the NMSSM and the masses of the supersymmetric particles are constrained by the results from the high energy colliders LEP at CERN and Tevatron at Fermilab [10,11]. A key role for the production of Higgs bosons at e^+e^- colliders plays the Higgs coupling to Z bosons, while neutralino production at LEP crucially depends on the $Z\tilde{\chi}^0\tilde{\chi}^0$ coupling which is formally identical in NMSSM and MSSM and differs only by the neutralino mixing. All those couplings are suppressed in the NMSSM if the respective neutralinos or Higgs bosons have significant singlet components. Therefore NMSSM neutralino and Higgs mass bounds are much weaker than in the minimal model [12]. The consequences from the negative neutralino search at the CERN e^+e^- collider LEP for the parameter space and the neutralino masses have been studied in [10]. In [12] it is shown that a very light NMSSM neutralino cannot even be ruled out at LEP2.

We used the following constraints from LEP. For new physics contributing to the total Z width $\Delta\Gamma(Z\to\tilde{\chi}^+\tilde{\chi}^- + Z\to\tilde{\chi}_i^0\tilde{\chi}_j^0)<23$ MeV. For new physics contributing to the invisible Z width $\Delta\Gamma(Z\to\tilde{\chi}_i^0\tilde{\chi}_j^0)<8$ MeV. From the direct neutralino search $B(Z\to\tilde{\chi}_1^0\tilde{\chi}_j^0)<2\times10^{-5}$ for $j=2,\ldots,5$, and $B(Z\to\tilde{\chi}_i^0\tilde{\chi}_j^0)<5\times10^{-5}$, for $i,j=2,\ldots,5$. The results of LEP searches for S_aZ and S_aZ^* productions [9], which impose restrictions on the S_aZZ couplings were included in our analysis. We have also included the experimental bounds from the direct search for pseudoscalar Higgs bosons produced together with a Higgs scalar at LEP [9], but this in accordance with [12] does not significantly affect the excluded parameter domain.

In our numerical analysis we use the following experimental restrictions for the SUSY particle spectrum in the NMSSM: $m_{\tilde{\chi}_1^+}\geq 90$ GeV for charginos, $m_{\tilde{\nu}}\geq 80$ GeV for sneutrinos, $m_{\tilde{e}_R}\geq 80$ GeV for selectrons, $m_{\tilde{q}}\geq 150$ GeV for squarks, $m_{\tilde{t}_1}\geq 60$ GeV for light top squark, $m_{H^+}\geq 65$ GeV for charged Higgs bosons and $m_{S_1}\geq 1$ GeV for the light scalar neutral Higgs boson. In fact, it appeared that all above-mentioned constraints do not allow m_{S_1} to be smaller than 20 GeV.

V. NUMERICAL ANALYSIS

Randomly scanned parameters of the NMSSM at the Fermi scale are the following: the gaugino mass parameters M' and M, the ratio of the doublet vacuum expectation values, $\tan\beta$, the singlet vacuum expectation value x, the couplings in the superpotential λ and k, squared squark mass parameters $m^2_{Q_{1,2}}$ for the first two generations and $m^2_{Q_3}$ for the third one, the parameters A_λ, A_k, as well as A_t for the third generation. The parameters are varied in the intervals given below:

$$-1000 \text{ GeV} < M' < 1000 \text{ GeV}$$

$$-2000 \text{ GeV} < M < 2000 \text{ GeV}$$

$$1 < \tan\beta < 50$$

$$0 \text{ GeV} < x < 10000 \text{ GeV}$$

$$-0.87 < \lambda < 0.87$$

$$-0.63 < k < 0.63$$

$$100 \text{ GeV}^2 < m^2_{Q_{1,2}} < 1\,000\,000 \text{ GeV}^2$$

$$100 \text{ GeV}^2 < m^2_{Q_3} < 1\,000\,000 \text{ GeV}^2$$

$$-2000 \text{ GeV} < A_t < 2000 \text{ GeV}$$

$$-2000 \text{ GeV} < A_\lambda < 2000 \text{ GeV}$$

$$-2000 \text{ GeV} < A_k < 2000 \text{ GeV}.$$

For simplicity the other sfermion mass parameters $m^2_{U_{1,2}}$, $m^2_{D_{1,2}}$, $m^2_{L_{1,2}}$, $m^2_{E_{1,2}}$, and m^2_T, m^2_B, $m^2_{L_3}$, $m^2_{E_3}$ are chosen to be equal to $m^2_{Q_2}$ and $m^2_{Q_3}$, respectively. Therefore masses of the sfermions in the same generation differ only due to the D-term contribution. Other parameters (except A_t) of the supersymmetry breaking potential A_U, A_D, A_E (for all three generations) are fixed to be zero.

The main results of our scan are presented in Fig. 1 in the form of scatter plots. Given in Fig. 1 are the total event rates R for ^{73}Ge, and the LSP gaugino fraction ($\mathcal{N}^2_{11}+\mathcal{N}^2_{12}$), singlino fraction ($\mathcal{N}^2_{15}$), and finally relic density parameter $\Omega_\chi h_0^2$ versus the LSP mass. The left panel in Fig. 1 presents the above-mentioned observables obtained without taking into account the cosmological relic density constraint.

In this case the total expected event rate R reaches values up to about 50 events per day and per 1 kg of the ^{73}Ge isotope. As on can see from Fig. 1 the small-mass LSP (less than about 100 GeV) is mostly gauginos, with a very small admixture of the singlino component. Large masses of the LSP (larger than 100 GeV) correspond to sizable gaugino and singlino components together perhaps with some higgsino fraction.

The results of implementation of the cosmological constraint

$$0.025 < \Omega_\chi h_0^2 < 1$$

can be seen in the right panel of Fig. 1. There is approximately a 5-fold reduction of the number of the points which fulfill all restrictions in this case. Nevertheless quite large values of event rate R (above 1 event/day/kg) still survive the cosmological constraint. The lower bound for the mass of the LSP now becomes about 3–5 GeV. The gaugino component becomes more significant, but the singlino fraction cannot be completely ruled out especially for large masses of the LSP. The higgsino component of the LSP remains still possible only for LSP masses in the vicinity of M_Z.

For illustration in Fig. 2 we present the calculated event rate R as function of the mass of the lightest scalar Higgs boson, m_{S_1}. The largest values of R are concentrated mostly in the region of quite large masses m_{S_1}, where LEP constraints are not very significant. The upper bound for m_{S_1} is also clearly seen.

VI. CONCLUSION

In the paper we address the question whether the next-to-minimal supersymmetric standard model can be attractive from the point of view of the direct detection of neutralinos provided the neutralino is the stable LSP.

To answer the question we derived the effective low energy neutralino-quark Lagrangian, which takes into account the contributions of extra scalar Higgs bosons and extra neutralinos. On this basis we calculated the total direct-dark-matter-detection event rate in ^{73}Ge as a representative isotope which is interesting for construction of a realistic dark

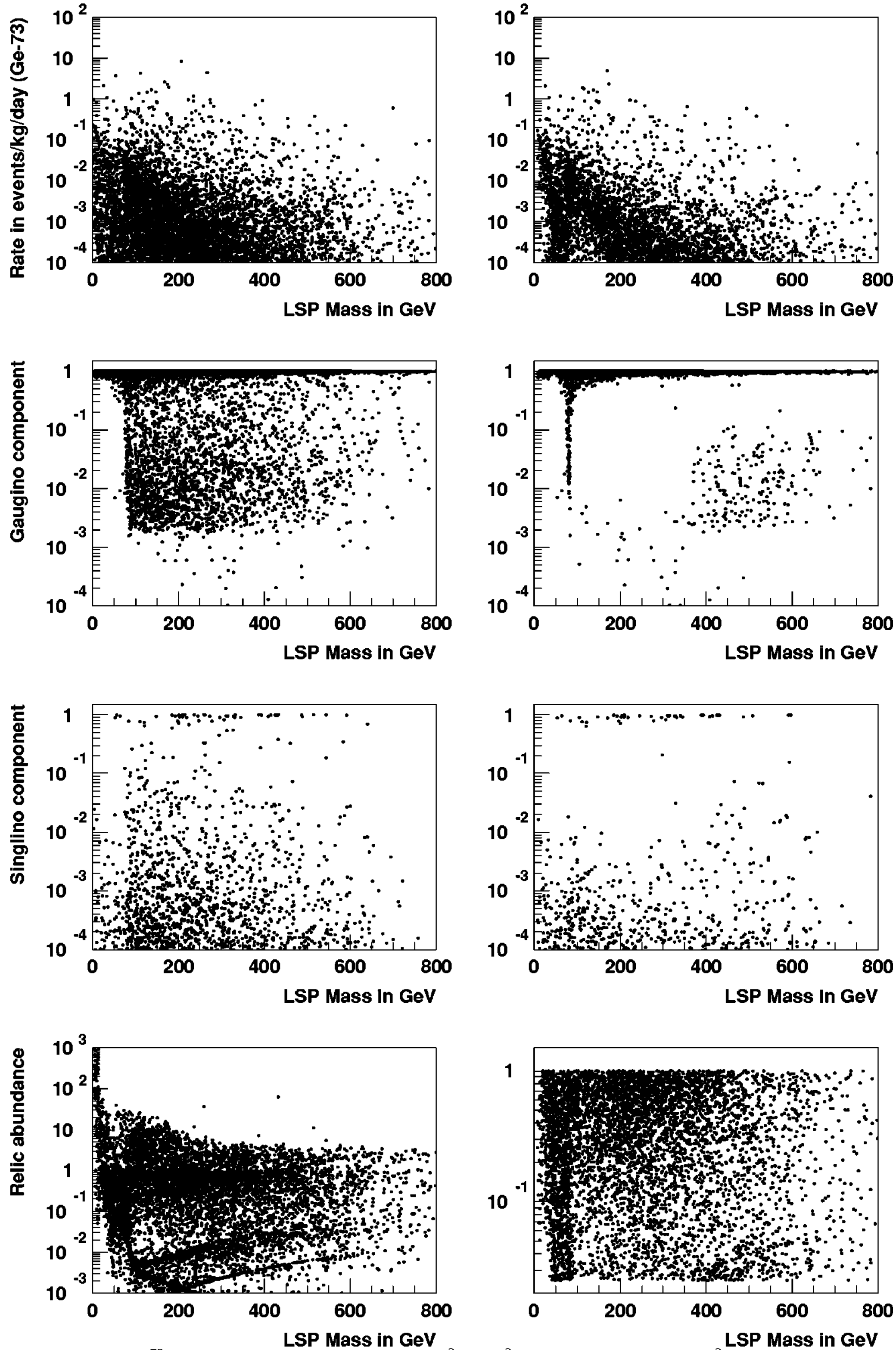

FIG. 1. Total event rate R for ^{73}Ge, the LSP gaugino fraction ($\mathcal{N}_{11}^2+\mathcal{N}_{12}^2$), singlino fraction ($\mathcal{N}_{15}^2$), and the relic abundance parameter $\Omega_\chi h_0^2$ versus the LSP mass (from up to down). The left (right) panel presents results obtained without (with) taking into account the cosmological relic density constraint.

matter detector. We analyzed the NMSSM taking into account the available accelerator and cosmological constraints by means of a random scan of the NMSSM parameter space at the Fermi scale. We demonstrated that the cosmological constraint does not rule out domains in the parameter space which correspond to quite sizable event rates in a germanium detector.

Due to relaxation of the gaugino unification condition,

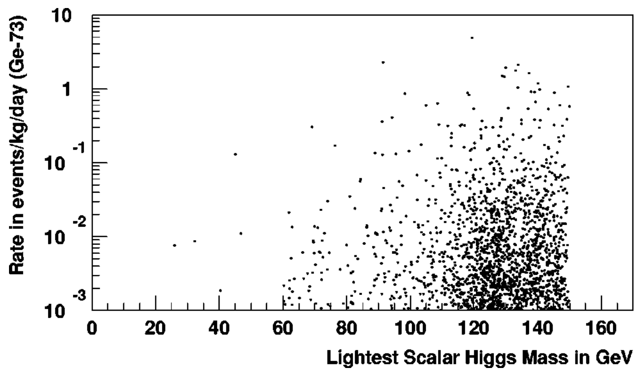

FIG. 2. Total event rate R for ^{73}Ge as function of the mass of the lightest scalar Higgs boson m_{S_1}.

contrary to previous consideration [17] we found domains in the parameter space where the lightest neutralinos have quite small masses (about 3 GeV), acceptable relic abundance and a sufficiently large expected event rate for direct detection with a ^{73}Ge-detector.

Therefore the NMSSM looks no worse than the MSSM from the point of view of direct dark matter detection. The question arises: Is it possible to distinguish MSSM and NMSSM by means of direct dark matter detection of LSP? It is a problem to be solved in the future. The question can disappear by itself if a negative search for light Higgs boson with the CERN Large Hadron Collider (LHC) rules out the MSSM. As already mentioned in the Introduction the NMSSM can bypass the most crucial constraint for the MSSM with the upper bound for the light Higgs boson (1). Therefore the NMSSM might remain a viable theoretical background for direct dark matter search for relic neutralinos in the post-MSSM epoch.

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[1] H. V. Klapdor-Kleingrothaus, in *Proceedings of the International Conference on Beyond the Desert*, edited by H. V. Klapdor-Kleingrothaus and H. Pais (IOP, Bristol, 1998), p. 485; in *Proceedings of the International Workshop on Non-Accelerator New Physics*, Dubna, 1997, edited by V. Brudanin, V. Bednyakov, and S. Kovalenko [Phys. At. Nucl. **61**, 967 (1998)]; H. V. Klapdor-Kleingrothaus, J. Hellmig, and M. Hirsch, J. Phys. G **24**, 483 (1998).

[2] H. V. Klapdor-Kleingrothaus and Y. Ramachers, in *Proceedings of the International Conference on Beyond the Desert* [1], p. 786.

[3] L. Baudis *et al.*, Nucl. Instrum. Methods Phys. Res. A **385**, 265 (1997); L. Baudis *et al.*, in *Proceedings of the International Conference on Dark98,* Heidelberg, Germany, 1998, edited by H. V. Klapdor-Kleingrothaus and L. Baudis (IOP, Bristol, in press).

[4] G. Jungman, M. Kamionkowski, and K. Griest, Phys. Rep. **267**, 195 (1996), and references therein.

[5] H. E. Haber and G. L. Kane, Phys. Rep. **117**, 75 (1985); J. F. Gunion and H. E. Haber, Nucl. Phys. **B272**, 1 (1986).

[6] M. Drees, Int. J. Mod. Phys. A **4**, 3635 (1989).

[7] J. Ellis, J. F. Gunion, H. E. Haber, L. Roszkowski, and F. Zwirner, Phys. Rev. D **39**, 844 (1989).

[8] J. E. Kim and H. P. Nilles, Phys. Lett. **138B**, 150 (1984).

[9] ALEPH Collaboration, D. Decamp *et al.*, Phys. Rep. **216**, 253 (1992); L3 Collaboration, O. Adriani *et al.*, *ibid.* **236**, 1 (1993); ALEPH Collaboration, D. Buskulic *et al.*, Phys. Lett. B **313**, 312 (1993); OPAL Collaboration, R. Akers *et al.*, Z. Phys. C **64**, 1 (1994).

[10] F. Franke, H. Fraas, and A. Bartl, Phys. Lett. B **336**, 415 (1994).

[11] F. Franke and H. Fraas, Phys. Lett. B **353**, 234 (1995).

[12] F. Franke and H. Fraas, Z. Phys. C **72**, 309 (1996); Int. J. Mod. Phys. A **12**, 479 (1997).

[13] U. Ellwanger, Phys. Lett. B **303**, 271 (1993).

[14] T. Elliott, S. F. King, and P. L. White, Phys. Lett. B **314**, 56 (1993); Phys. Rev. D **49**, 2435 (1994).

[15] P. N. Pandita, Z. Phys. C **59**, 575 (1993).

[16] S. A. Abel, S. Sarkar, and I. B. Whittingham, Nucl. Phys. **B392**, 83 (1993).

[17] A. Stephan, Phys. Lett. B **411**, 97 (1997); Phys. Rev. D **58**, 035011 (1998).

[18] S. F. King and P. L. White, Phys. Rev. D **52**, 4183 (1995).

[19] M. Drees and M. M. Nojiri, Phys. Rev. D **48**, 3483 (1993).

[20] J. Ellis and R. Flores, Phys. Lett. B **263**, 259 (1991); **300**, 175 (1993); Nucl. Phys. **B400**, 2536 (1993).

[21] P. Nath and R. Arnowitt, Phys. Rev. D **56**, 2820 (1997).

[22] J. Engel and P. Vogel, Phys. Rev. D **40**, 3132 (1989); J. Engel, S. Pitel, and P. Vogel, Int. J. Mod. Phys. E **1**, 1 (1992).

[23] V. A. Bednyakov, H. V. Klapdor-Kleingrothaus, and S. G. Kovalenko, Phys. Rev. D **50**, 7128 (1994); **55**, 503 (1997); V. A. Bednyakov, H. V. Klapdor-Kleingrothaus, S. G. Kovalenko, and Y. Ramachers, Z. Phys. A **357**, 339 (1997).

[24] L. Bergström and P. Gondolo, Astropart. Phys. **5**, 263 (1996).

[25] E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, Reading, MA, 1990).

[26] J. Ellis *et al.*, Nucl. Phys. **B238**, 453 (1984); J. Ellis, L. Roszkowski, and Z. Lalak, Phys. Lett. B **245**, 545 (1990).

[27] M. Drees and M. M. Nojiri, Phys. Rev. D **47**, 376 (1993).

[28] K. Griest and D. Seckel, Phys. Rev. D **43**, 3191 (1991); P. Nath and R. Arnowitt, Phys. Rev. Lett. **70**, 3696 (1993); G. Gelmini and P. Gondolo, Nucl. Phys. **B360**, 145 (1991).

[29] G. Steigman, K. A. Olive, D. N. Schramm, and M. S. Turner, Phys. Lett. B **176**, 33 (1986); J. Ellis, K. Enquist, D. V. Nanopoulos, and S. Sarkar, Phys. Lett. **167B**, 457 (1986); L. Roszkowski, Phys. Rev. D **50**, 4842 (1994).

[30] J. L. Lopez, D. V. Nanopoulos, and H. Pois, Phys. Rev. D **47**,

2468 (1993); P. Nath and R. Arnowitt, Phys. Rev. Lett. **70**, 3696 (1993); J. L. Lopez, D. V. Nanopoulos, and K. Yuan, Phys. Rev. D **48**, 2766 (1993); G. L. Kane, C. Kolda, L. Roszkowski, and J. D. Wells, *ibid.* **49**, 6173 (1994).

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

[32] A. Bottino *et al.*, Astropart. Phys. **2**, 77 (1994); A. Bottino *et al.*, in *Proceedings of the Workshop on Double Beta Decay and Related Topics*, edited by H. V. Klapdor-Kleingrothaus and S. Stoica (World Scientific, Singapore, 1996), p. 281.

[33] V. Berezinsky *et al.*, Astropart. Phys. **5**, 1 (1996).