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Interactions of exotic neutralino dark matter with nucleons in $U(1)$ extensions of the MSSM originating from E_6 GUTs

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Abstract

To ensure anomaly cancellation, the E_6 inspired $U(1)$ extensions of the Minimal Supersymmetric (SUSY) Standard Model (MSSM) involve extra exotic matter. The lightest exotic neutralino in these models can be stable contributing to the cold dark matter density. We consider the interactions of such neutralino with nucleons within a specific extension of the MSSM with an additional $U(1)_N$ gauge symmetry (SE$_6$SSM). The constraints on the couplings of this state, which are set by the present experimental bounds caused by the direct detection experiments, are examined. The obtained results can be generalised to other E_6 inspired SUSY models with extra $U(1)$ gauge symmetry.

Keywords: unified field theories and models, models beyond the Standard Model, supersymmetry, cold dark matter

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1. Introduction

Astrophysical and cosmological observations imply that non-luminous cold dark matter constitutes about 20–25% of the energy density of the Universe [1]. The nature of dark matter is unknown. It cannot consist of any elementary particles discovered so far. Therefore, the presence of dark matter is widely considered as one of the strongest indications for new physics beyond the Standard Model (SM) which describes well the experimental data measured in Earth-based experiments.

Models with softly broken supersymmetry (SUSY) remain the best motivated extensions of the SM. The discovered SM-like Higgs scalar with mass around 125 GeV is consistent with the Minimal and Next-to-Minimal Supersymmetric Standard Models (MSSM and NMSSM). In these models, the electroweak (EW) scale is almost stabilized and the lightest neutralino is absolutely stable if R-parity is preserved so that this state can play the role of dark matter. As neutralinos are heavy weakly interacting massive particles (WIMPs), they can explain the large-scale structure of the Universe [2, 3] providing the correct relic abundance of cold dark

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matter [4]. Within the MSSM, the EW and strong gauge couplings extrapolated to high energies using the renormalisation group equations (RGEs) converge to a common value near some high energy scale $M_X \sim 2\cdot 10^{16}$ GeV [5–8]. This permits embedding the EW and strong interactions into Grand Unified Theories (GUTs) [9], which are based on $SU(5)$, $SO(10)$ or E_6 gauge groups, as well as into a ten-dimensional superstring theory with $E_8 \times E_8'$ gauge symmetry [10].

Around the GUT scale M_X, the gauge group E_6 (or E_8) reduces to $SO(10)\times U(1)_\psi$ (maximal subgroup of E_6) with sequential breakdown of $SO(10)$ to $SU(5)\times U(1)_\chi$ and $SU(5)$ to the SM gauge group $SU(3)_C \times SU(2)_W \times U(1)_Y$ (for reviews, see Refs. [11–13]). If the SUSY breaking mechanism gives rise to the sparticle mass scale in the multi-TeV range, then the breakdown of the $U(1)_\chi \times U(1)_\psi$ symmetry to its discrete subgroup

$$U(1)_\psi \times U(1)_\chi \to P_M = (-1)^{3(B-L)}, \tag{1}$$

where B and L are the baryon and lepton numbers, can result in a variety of SUSY models, including MSSM, NMSSM and their different modifications. The conservation of matter parity P_M implies that R-parity

$$Z_2^R = (-1)^{3(B-L)+2s}, \tag{2}$$

where s is the spin of the state, is also preserved.

The $U(1)$ extensions of the MSSM may arise in this unification scheme when the $U(1)_\chi \times U(1)_\psi$ symmetry is reduced to

$$U(1)' = U(1)_\chi \cos\theta_{E_6} + U(1)_\psi \sin\theta_{E_6}. \tag{3}$$

Given that $\theta_{E_6} \neq 0$ and $\theta_{E_6} \neq \pi$, the anomalies in such $U(1)$ extensions of the MSSM get cancelled if the particle spectrum involves complete representations of E_6. Consequently, one needs to augment quarks and leptons of the SM by a number of exotics so that all these states form three complete fundamental 27 representations of E_6 at low energies. These 27-plets decompose under the $SU(5)\times U(1)_\psi \times U(1)_\chi$ subgroup of E_6 as follows:

$$\begin{aligned} 27_i \to \left(10,\, \frac{1}{\sqrt{24}},\, -\frac{1}{\sqrt{40}}\right)_i + \left(5^*,\, \frac{1}{\sqrt{24}},\, \frac{3}{\sqrt{40}}\right)_i + \left(5^*,\, -\frac{2}{\sqrt{24}},\, -\frac{2}{\sqrt{40}}\right)_i + \\ + \left(5,\, -\frac{2}{\sqrt{24}},\, \frac{2}{\sqrt{40}}\right)_i + \left(1,\, \frac{4}{\sqrt{24}},\, 0\right)_i + \left(1,\, \frac{1}{\sqrt{24}},\, -\frac{5}{\sqrt{40}}\right)_i . \end{aligned} \tag{4}$$

The quantities in brackets are the $SU(5)$ representation as well as extra $U(1)_\psi$ and $U(1)_\chi$ charges, while $i = 1,2,3$ is a family index. An ordinary SM family, which includes the doublets of left-handed leptons L_i and quarks Q_i, right-handed charged leptons, up- and down-quarks (e_i^c, u_i^c and d_i^c), is assigned to $\left(10,\, \frac{1}{\sqrt{24}},\, -\frac{1}{\sqrt{40}}\right)_i + \left(5^*,\, \frac{1}{\sqrt{24}},\, \frac{3}{\sqrt{40}}\right)_i$. The last term in Eq. (4) corresponds to the right-handed neutrinos N_i^c. The next-to-last term in Eq. (4) can be identified with extra SM-singlet fields S_i with non-zero $U(1)_\psi$ charges. The pairs of $SU(2)_W$ doublets (H_i^d and H_i^u), which are contained in $\left(5^*,\, -\frac{2}{\sqrt{24}},\, -\frac{2}{\sqrt{40}}\right)_i$ and $\left(5,\, -\frac{2}{\sqrt{24}},\, \frac{2}{\sqrt{40}}\right)_i$, have the quantum numbers of Higgs doublets (Higgs-like doublets) so that they form either Higgs or exotic Higgs $SU(2)_W$ multiplets. The colour triplet components of these $SU(5)$ multiplets are associated with the exotic quarks $\overline{D}_i$ and D_i with electric charges $+1/3$ and $-1/3$. The $B-L$ charges of these exotic quarks are twice larger than that of ordinary ones, i.e., $\left(\pm\frac{2}{3}\right)$. Thus, they can be either leptoquarks or diquarks.

The E_6 inspired $U(1)$ extensions of the MSSM were extensively studied over the years. Different aspects of phenomenology of the exotic quarks/squarks and Z' boson were considered in Refs. [14–16]. The implications of these models were explored for neutrino physics [17, 18] and models explaining the hierarchy of fermion masses [19], EW baryogenesis [20, 21], leptogenesis [22–25], muon anomalous magnetic moment [26, 27], CP-violation in the Higgs sector [28] and EW symmetry breaking [29–35], lepton flavour violating processes like $\mu \to e\gamma$ [36], electric dipole moment of electron [37] and tau lepton [38]. The neutralino sector in such models was studied in Refs. [35–46]. In the E_6 inspired $U(1)$-extended SUSY models, the upper bound on the SM-like Higgs mass and the Higgs sector were analysed in Refs. [34, 46–50].

Here we explore the interactions of dark matter particles with nucleons in the framework of specific E_6 inspired SUSY model with an additional $U(1)_N$ gauge symmetry associated with $\theta_{E_6} = \arctan\sqrt{15}$ in Eq. (3) [48, 49] (for recent review, see Ref. [51]). This exceptional Supersymmetric Standard Model (E$_6$SSM) implies that near the GUT scale, $U(1)_\psi$ and $U(1)_\chi$ are reduced to

$$U(1)_\psi \times U(1)_\chi \to U(1)_N \times P_M. \tag{5}$$

Only in this case, the right-handed neutrinos have zero charges so that they can be extremely heavy [48, 49]. Although we restrict our analysis here to a particular value of θ_{E_6}, our results can be easily generalised to other E_6 inspired $U(1)$ extensions of the MSSM in which matter parity is preserved.

In all E_6 inspired $U(1)$-extended SUSY models, extra exotic matter may lead to rapid proton decay and non-diagonal flavour transitions. A set of discrete symmetries permits suppressing the corresponding operators [48, 49]. In this paper, we consider a variant of the E$_6$SSM (SE$_6$SSM) [52–55] in which the most dangerous baryon and lepton number violating operators as well as tree-level flavour-changing transitions are forbidden by a single discrete $\tilde{Z}_2^H$ symmetry. Two discrete symmetries $\tilde{Z}_2^H$ and P_M give rise to at least two stable states, i.e., gravitino and the lightest exotic neutralino, that may compose cold dark matter density. In the previous article [56], it was pointed out that the dark matter-nucleon scattering cross section can be sufficiently strongly suppressed. Here we compare the computed values of this cross section with the corresponding experimental bounds and identify the restrictions on the parameter space of the SE$_6$SSM caused by the latest results of the LUX–ZEPLIN (LZ) experiment [57].

The layout of this article is as follows. In Section 2, we briefly review the SUSY model under consideration. In Section 3, the constraints on the SE$_6$SSM parameter space, which come from the direct detection experiments, are considered. Section 4 concludes the paper.

2. Exotic neutralino dark matter

Different modifications of the E$_6$SSM were considered in Refs. [48, 49, 52, 53, 58–68]. To ensure the cancellation of anomalies below the GUT scale M_X, the matter content of the simplest variant of the E$_6$SSM involves three 27-plets as well as a pair of $SU(2)_W$ doublets, i.e., L_4 and $\overline{L}_4$ [51]. As L_4 and $\overline{L}_4$ come from extra $27'$ and $\overline{27}'$, they carry opposite $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$ quantum numbers, and gauge anomalies still cancel. In the simplest variant of the E$_6$SSM, the supermultiplets L_4 and $\overline{L}_4$ facilitate the gauge coupling unification near the scale M_X [69]. In the vicinity of the quasi-fixed point, the upper bound on the mass of the lightest Higgs boson within this variant of the E$_6$SSM was analysed in Ref. [70]. Such quasi-fixed point appears as an intersection of the quasi-fixed and invariant lines [71–73]. Extra exotic states in the E$_6$SSM can give rise to distinctive LHC signatures [48–50, 59, 62, 74–76] and non-standard Higgs decays [53, 77–83].

Using approach proposed in Refs. [84, 85], it was found that in the E$_6$SSM the lightest supersymmetric particle (LSP), i.e., the lightest R-parity odd state, has to be lighter than 60−65 GeV [78]. Such LSPs may account for some of the cold dark matter density if their masses are close to half the Z-boson mass $M_Z/2$ [78]. Nevertheless, in this part of the E$_6$SSM parameter space, the SM-like Higgs scalar decays predominantly into a pair of LSPs, while other branching ratios of the Higgs boson are rather small. LHC experiments ruled out such scenario. The simplest phenomenologically viable scenario implies that LSPs in the E$_6$SSM are considerably lighter than 1 eV forming hot dark matter in our Universe. They give only a minor contribution to the dark matter density. The presence of very light neutral fermions may result in interesting implications for the neutrino physics [86].

Table 1. The $U(1)_Y$ and $U(1)_N$ charges of the SE$_6$SSM supermultiplets.

	Q_i	u_i^c	d_i^c	L_i, L_4	e_i^c	N_i^c	S_i, S	H_α^u, H_u	H_α^d, H_d	D_i	$\overline{D}_i$	$\overline{L}_4$	$\overline{S}$	ϕ_i, ϕ
$\sqrt{\frac{5}{3}}Q_i^Y$	$\frac{1}{6}$	$-\frac{2}{3}$	$\frac{1}{3}$	$-\frac{1}{2}$	1	0	0	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{1}{3}$	$\frac{1}{2}$	0	0
$\sqrt{40}Q_i^N$	1	1	2	2	1	0	5	-2	-3	-2	-3	-2	-5	0

In addition to three 27-plets, L_4 and $\overline{L}_4$, the field content of the SE$_6$SSM also includes a pair of superfields S and $\overline{S}$ arising from another $\widetilde{27}+\overline{\widetilde{27}}$ representations as well as four E_6 singlet superfields ϕ and ϕ_i $(i=1,2,3)$. Extra superfields permit avoiding the appearance of extremely light neutral fermions in the particle spectrum. The SE$_6$SSM matter content can originate from the E_6 orbifold GUT model in six dimensions [52]. The $U(1)_Y$ and $U(1)_N$ charges of all matter supermultiplets in the SE$_6$SSM are given in Table 1. The supermultiplets ϕ, L_4, $\overline{L}_4$, S, $\overline{S}$ and one pair of the Higgs-like doublets (say, $H_d \equiv H_3^d$ and $H_u \equiv H_3^u$) are required to be even under the discrete $\tilde{Z}_2^H$ symmetry, while the remaining supermultiplets have to be $\tilde{Z}_2^H$ odd. Neglecting all suppressed non-renormalisable interactions, the SE$_6$SSM superpotential can be written as

$$\begin{aligned} W_{\mathrm{SE_6SSM}} = \lambda S(H_u H_d) - \sigma\phi S\overline{S} + \frac{\kappa}{3}\phi^3 + \frac{\mu}{2}\phi^2 + \Lambda\phi + \mu_L L_4\overline{L}_4 + \tilde{\sigma}\phi L_4\overline{L}_4 + W_{IH} + \\ + \kappa_{ij}S(D_i\overline{D}_j) + g_{ij}^D(Q_iL_4)\overline{D}_j + h_{i\alpha}^E e_i^c(H_\alpha^d L_4) + g_{ij}\phi_i\overline{L}_4 L_j + W_N + W_{\mathrm{MSSM}}, \end{aligned} \tag{6}$$

$$W_{IH} = \tilde{M}_{ij}\phi_i\phi_j + \tilde{\kappa}_{ij}\phi\phi_i\phi_j + \tilde{\lambda}_{ij}\overline{S}\phi_i S_j + \lambda_{\alpha\beta}S(H_\alpha^d H_\beta^u) + \tilde{f}_{i\alpha}S_i(H_\alpha^d H_u) + f_{i\alpha}S_i(H_d H_\alpha^u), \tag{7}$$

$$W_N = \frac{1}{2}M_{ij}N_i^cN_j^c + \frac{1}{2}\eta_{ij}\phi N_i^c N_j^c + \tilde{h}_{ij}N_i^c(H_uL_j) + h_{i\alpha}N_i^c(H_\alpha^u L_4), \tag{8}$$

$$W_{\mathrm{MSSM}} = y_{ij}^U(Q_iH_u)u_j^c + y_{ij}^D(Q_iH_d)d_j^c + y_{ij}^L(L_iH_d)e_j^c. \tag{9}$$

In Eqs. (6)–(9), $\alpha,\beta = 1,2$ and $i,j = 1,2,3$. When $M_{ij} = \eta_{ij} = 0$, the expressions (6)–(9) are valid for any E_6 inspired $U(1)$ extension of the MSSM (with $\theta_{E_6} \neq 0$ and $\theta_{E_6} \neq \pi$) if $\tilde{Z}_2^H$ and P_M are preserved. The $\tilde{Z}_2^H$ symmetry ensures that the SE$_6$SSM superpotential does not contain any operators that lead to rapid proton decay. The $SU(2)_W$ doublets H_d and H_u play the role of the MSSM Higgs supermultiplets. Because of the $\tilde{Z}_2^H$ symmetry the down-type quarks and charged leptons couple to just H_d, whereas the up-type quarks couple to H_u only. As a consequence, the non-diagonal flavour transitions are suppressed at tree level. The Yukawa couplings of L_4 to Q_i and $\overline{D}_j$, which are allowed by the $\tilde{Z}_2^H$ symmetry, permit the lightest exotic quarks to decay within a reasonable time.

The $U(1)_N$ gauge symmetry is spontaneously broken in the SE$_6$SSM by the vacuum expectation values (VEVs) of the superfields S and $\overline{S}$. In the limit $\sigma \to 0$, the corresponding part of the scalar potential is given by

$$V_S(S,\,\overline{S}) = m_S^2|S|^2 + m_{\overline{S}}^2|\overline{S}|^2 + \frac{Q_S^2 g_1'^2}{2}\left(|S|^2 - |\overline{S}|^2\right)^2, \tag{10}$$

where m_S^2 and $m_{\overline{S}}^2$ are the soft SUSY breaking masses, while g_1' and Q_S are the $U(1)_N$ gauge coupling and the $U(1)_N$ charge of S. When $(m_S^2 + m_{\overline{S}}^2) < 0$, the minimum of the potential (10) is attained for $\langle S\rangle = \langle \overline{S}\rangle \to \infty$. This run-away direction is stabilized for non-zero values of coupling σ. If $\sigma \ll 1$, then the VEVs

$$\langle \phi\rangle \sim \langle S\rangle \simeq \langle \overline{S}\rangle \sim \frac{M_S}{\sigma} \tag{11}$$

tend to be much larger than the sparticle mass scale M_S. The exotic states and Z' boson are rather heavy in this case.

For practical purposes, it is convenient to introduce the Z_2^E symmetry which is defined as $\tilde{Z}_2^H = P_M \times Z_2^E$ [52]. The transformation properties of the SE$_6$SSM supermultiplets under the P_M, Z_2^E and $\tilde{Z}_2^H$ symmetries are presented in Table 2. Since $\tilde{Z}_2^H$ and P_M are exact, the Z_2^E symmetry and R-parity are conserved. Therefore, the lightest R-parity odd state should be stable. Here we assume that gravitino is the lightest SUSY particle in the spectrum and can give a substantial contribution to the cold dark matter density. As gravitino is even under the discrete Z_2^E symmetry, the lightest Z_2^E odd state has to be stable as well [52].

Table 2. Transformation properties of the SE$_6$SSM supermultiplets under the discrete symmetries.

	$Q_i, u_i^c, d_i^c, L_i, e_i^c, N_i^c$	$\overline{D}_i, D_i, H_\alpha^d, H_\alpha^u, S_i, \phi_i$	$H_d, H_u, S, \overline{S}, \phi$	$L_4, \overline{L}_4$
$\tilde{Z}_2^H$	$-$	$-$	$+$	$+$
P_M	$-$	$+$	$+$	$-$
Z_2^E	$+$	$-$	$+$	$-$

The phenomenologically acceptable scenarios with gravitino LSP imply that all unstable particles decay before Big Bang Nucleosynthesis (BBN). For the sparticle mass scale M_S in the multi-TeV range, this requirement can be satisfied if gravitino mass $m_{3/2}$ is lower than 10 GeV [56]. The contribution of such gravitinos to the cold dark matter density is determined by the reheating temperature T_R [87, 88]:

$$\Omega_{3/2}h^2 \sim 0.27\left(\frac{T_R}{10^8\ \text{GeV}}\right)\left(\frac{1\ \text{GeV}}{m_{3/2}}\right)\left(\frac{M_{\tilde{g}}}{1\ \text{TeV}}\right)^2, \tag{12}$$

where $M_{\tilde{g}}$ is a gluino mass. As $\Omega_{3/2}h^2 \leqslant 0.12$ [1], for $M_{\tilde{g}} \gtrsim 3$ TeV and $m_{3/2} \simeq 1$ GeV, one finds an upper limit $T_R \lesssim 10^{6-7}$ GeV [89]. Nevertheless, even for so low T_R, the appropriate baryon asymmetry can be generated within the SE$_6$SSM via the out-of-equilibrium decays of the lightest right-handed neutrino/sneutrino [22].

The part of the SE$_6$SSM superpotential W_{IH} describes the interactions of the $\tilde{Z}_2^H$ even supermultiplets ϕ, S, $\overline{S}$, H_u and H_d with ϕ_i, S_i, H_α^u and H_α^d which do not develop VEVs. The fermion components of these $\tilde{Z}_2^H$ odd supermultiplets compose the exotic neutralino and chargino states. The signatures associated with such exotic neutralino states were studied in

Refs. [90, 91]. If all components of ϕ_i are much heavier than the fermions and bosons from H^u_α, H^d_α and S_i, the superfields ϕ_i can be integrated out and W_{IH} reduces to

$$W_{IH} \to \widetilde{W}_{IH} \simeq -\widetilde{\mu}_{ij} S_i S_j + \lambda_{\alpha\beta} S(H^d_\alpha H^u_\beta) + \tilde{f}_{i\alpha} S_i (H^d_\alpha H_u) + f_{i\alpha} S_i (H_d H^u_\alpha) + \ldots \tag{13}$$

Here and further we use the basis in which $\lambda_{\alpha\beta} = \lambda_{\alpha\alpha}\,\delta_{\alpha\beta}$ and $\widetilde{\mu}_{ij} = \widetilde{\mu}_i\,\delta_{ij}$.

As mentioned before, the lightest exotic state with $Z^E_2 = -1$ has to be stable. In this paper, it is assumed that such state is mostly formed by the neutral fermion components $\widetilde{H}^{u0}_1$ and $\widetilde{H}^{d0}_1$ of the supermultiplets H^u_1 and H^d_1. We explore the scenarios in which all sparticles except gravitino as well as all other exotic states are rather heavy, i.e., they are considerably heavier than 1 TeV. It is also expected that the supermultiplets H^u_1 and H^d_1 mainly interact with S_1, H_u and H_d. To simplify the analysis, all other couplings of these $SU(2)_W$ doublets are set to be negligibly small. In this approximation, the masses of the lightest exotic neutralinos are eigenvalues of the mass matrix [56]:

$$M^{ab} = -\begin{pmatrix} 0 & \mu_{11} & \dfrac{\tilde{f}_{11}}{\sqrt{2}} v_2 \\ \mu_{11} & 0 & \dfrac{f_{11}}{\sqrt{2}} v_1 \\ \dfrac{\tilde{f}_{11}}{\sqrt{2}} v_2 & \dfrac{f_{11}}{\sqrt{2}} v_1 & \widetilde{\mu}_1 \end{pmatrix}, \tag{14}$$

where v_1 and v_2 are the VEVs of H_d and H_u, i.e., $\langle H_u \rangle = v_2/\sqrt{2}$, $\langle H_d \rangle = v_1/\sqrt{2}$ and $v = \sqrt{v_1^2 + v_2^2} \approx 246$ GeV. In Eq. (14), $\mu_{11} = \lambda_{11}\langle S \rangle$.

The perturbation theory method allows one to diagonalise the mass matrix (14) if $|\widetilde{\mu}_1|$ is much larger than v and $|\mu_{11}|$ (see, for example, Refs. [92–94]). For $\mu_1 > 0$ and $\mu_{11} > 0$, one obtains

$$m_{\chi_1} \simeq \mu_{11} - \Delta_1, \qquad m_{\chi_2} \simeq \mu_{11} + \Delta_2, \qquad m_{\chi_3} \simeq \widetilde{\mu}_1 + \Delta_1 + \Delta_2, \tag{15}$$
$$\Delta_1 \simeq \frac{(\tilde{f}_{11} v \sin\beta + f_{11} v \cos\beta)^2}{4(\widetilde{\mu}_1 - \mu_{11})}, \qquad \Delta_2 \simeq \frac{(\tilde{f}_{11} v \sin\beta - f_{11} v \cos\beta)^2}{4(\widetilde{\mu}_1 + \mu_{11})},$$

where $\tan\beta = v_2/v_1$. Equations (15) indicate that the masses of the lightest and next-to-lightest exotic neutralinos (m_{χ_2} and m_{χ_1}) are determined by μ_{11}. In our analysis, we require that $m_{\chi_2} - m_{\chi_1} > 200$ MeV so that the next-to-lightest exotic neutralino χ_2 decays before BBN.

Using the approximate formula

$$\Omega_{\tilde{H}} h^2 \simeq 0.1 \left(\frac{\mu_{11}}{1\ \text{TeV}}\right)^2, \tag{16}$$

one can estimate the contribution of the lightest exotic neutralino χ_1 to the dark matter density. Equation (16) was derived within the MSSM in the case when dark matter is formed by the Higgsino states [95, 96]. Planck observations indicate that $(\Omega h^2)_{\text{exp}} = 0.1188 \pm 0.0010$ [1]. Thus, in the phenomenologically viable scenarios, μ_{11} has to be lower than 1.1 TeV. For $|\mu_{11}| < 1.1$ TeV, gravitino may account for some part of the dark matter density.

3. Direct detection constraints

Now let us consider the interactions of dark matter states with the baryons within the SE_6SSM. It is well known that gravitino with masses of a few GeV has negligibly small couplings to the SM particles. Therefore, the interactions of dark matter particles with nucleons are defined by the couplings of the stable lightest exotic neutralino χ_1. The two lightest exotic neutralinos are linear combinations of $\tilde{H}_1^{d0}$, $\tilde{H}_1^{u0}$ and the fermion component $\tilde{S}_1$ of the superfield S_1, i.e.,

$$\chi_\alpha = N_\alpha^1 \tilde{H}_1^{d0} + N_\alpha^2 \tilde{H}_1^{u0} + N_\alpha^3 \tilde{S}_1, \tag{17}$$

where $\alpha = 1, 2$. The exotic neutralino mixing matrix N_i^a $(a, i = 1, 2, 3)$ is defined by

$$N_i^a M^{ab} N_j^b = m_i \delta_{ij}, \qquad \text{no sum on } i. \tag{18}$$

The two lightest exotic neutralinos mostly interact with the Z-boson and Higgs states. The masses of all Higgs bosons except the SM-like one are set to be equal to 10 TeV. As a consequence, their contribution to the scattering of the lightest exotic neutralino on nuclei can be ignored. The part of the Lagrangian describing the interactions of the SM-like Higgs scalar h and Z boson with χ_1 and χ_2 is given by

$$\mathcal{L}_{Zh\chi} = \sum_{\alpha,\sigma} \frac{M_Z}{2v} R_{Z\alpha\sigma} Z_\mu \Bigl(\chi_\alpha^T \gamma_\mu \gamma_5 \chi_\sigma \Bigr) + \sum_{\alpha,\sigma} g_{h\alpha\sigma} h \Bigl(\chi_\alpha^T \chi_\sigma \Bigr), \tag{19}$$

$$R_{Z\alpha\sigma} = N_\alpha^1 N_\sigma^1 - N_\alpha^2 N_\sigma^2, \qquad g_{h\alpha\sigma} = -\frac{1}{\sqrt{2}} \Bigl(f_{11} N_\alpha^3 N_\sigma^2 \cos\beta + \tilde{f}_{11} N_\alpha^3 N_\sigma^1 \sin\beta \Bigr),$$

where $\alpha, \sigma = 1, 2$. The entries of the mixing matrix N_i^a can be computed using the perturbation theory method. Substituting the corresponding analytical expressions into Eqs. (19), one obtains the approximate formulae

$$|g_{h11}| \simeq \frac{\Delta_1}{v}, \qquad R_{Z11} \simeq \frac{v^2 (f_{11}^2 \cos^2\beta - \tilde{f}_{11}^2 \sin^2\beta)}{4\mu_{11}(\widetilde{\mu}_1 - \mu_{11})}, \tag{20}$$

which are valid when $\widetilde{\mu}_1 \gg \mu_{11} > 0$ and μ_{11} is considerably larger than $\tilde{f}_{11} v \sin\beta$ and $f_{11} v \cos\beta$.

The couplings of the SM-like Higgs scalar h with mass $m_h \simeq 125$ GeV to nucleons g_{hNN} are set by

$$g_{hNN} = a_S^N \frac{m_N}{v}, \qquad a_S^N = \sum_{q=u,d,s} f_{Tq}^N + \frac{2}{27} \sum_{Q=c,b,t} f_{TQ}^N, \tag{21}$$

where $N = p, n$, m_N is a nucleon mass and

$$\langle N | m_q \bar{q} q | N \rangle = m_N f_{Tq}^N, \qquad f_{TQ}^N = 1 - \sum_{q=u,d,s} f_{Tq}^N. \tag{22}$$

From Eqs. (21) and (22) it follows that the couplings g_{hNN} are determined by the hadronic matrix elements, i.e., the coefficients f_{Tq}^N. These coefficients are related to the π-nucleon σ term and the spin content of the nucleon. In our analysis, we set $f_{Tq}^p \simeq f_{Tq}^n \simeq f_{Tq}$ so that $a_S^p \simeq a_S^n \simeq a_S$. Here we also fix $f_{Tu} \simeq 0.0153$, $f_{Td} \simeq 0.0191$ and $f_{Ts} \simeq 0.0447$ which are the default values used in micrOMEGAs [97]. The t-channel exchange of the lightest CP-even Higgs scalar gives rise to the spin-independent part of χ_1-nucleon cross section which is given by

$$\sigma_{SI} = \frac{4 m_r^2 m_N^2}{\pi v^2 m_{h_1}^4} |g_{h\chi\chi} a_S|^2, \qquad m_r = \frac{m_{\chi_1} m_N}{m_{\chi_1} + m_N}. \tag{23}$$

The interactions of the Z boson with quarks are described by

$$\mathcal{L}_{Zq} = \sum_q \frac{\bar{g}}{2}\,\overline{q}(a_V^q\gamma^\mu + a_{PV}^q\gamma^\mu\gamma^5)q\,Z_\mu = \frac{\bar{g}}{2}J^\mu_{NC}Z_\mu, \tag{24}$$

$$a_V^q = T_{3q} - 2s_W^2 Q_q, \qquad a_{PV}^q = T_{3q}.$$

Using the Lagrangian (24), the following hadronic matrix elements can be computed:

$$\langle N'|J^\mu_{NC}|N\rangle = \bar{\psi}'_N\gamma^\mu\left(a_V^N + a_{PV}^N\gamma_5\right)\psi_N, \tag{25}$$

$$a_V^p \simeq \left(\frac{1}{2} - 2s_W^2\right), \qquad a_{PV}^p \simeq \left(\frac{1}{2}\Delta_u^{(p)} - \frac{1}{2}\Delta_d^{(p)} - \frac{1}{2}\Delta_s^{(p)}\right), \tag{26}$$

$$a_V^n \simeq -\frac{1}{2}, \qquad a_{PV}^n \simeq \left(\frac{1}{2}\Delta_u^{(n)} - \frac{1}{2}\Delta_d^{(n)} - \frac{1}{2}\Delta_s^{(n)}\right), \tag{27}$$

where $\Delta_u^p = \Delta_d^n = 0.777$, $\Delta_d^p = \Delta_u^n = -0.438$ and $\Delta_s^p = \Delta_s^n = -0.053$ [98]. From Eqs. (26) and (27) one finds $a_{PV}^p \simeq 0.63$ and $a_{PV}^n \simeq -0.58$. The t-channel exchange of the Z boson results in the spin-dependent parts of χ_1-proton and χ_1-neutron cross sections

$$\sigma^p = \frac{3m_r^2}{\pi v^4}|R_{Z11}a_{PV}^p|^2, \qquad \sigma^n = \frac{3m_r^2}{\pi v^4}|R_{Z11}a_{PV}^n|^2. \tag{28}$$

For each set of parameters $\tan\beta$, μ_{11}, $\widetilde{\mu}_1$, f_{11} and $\tilde{f}_{11}$, the χ_1-nucleon cross sections σ_{SI}, σ^p and σ^n may be calculated. The value of $\tan\beta$ in the SE$_6$SSM is less constrained as compared with the MSSM. In particular, in the MSSM, the scenarios with moderate values of $\tan\beta$, i.e., $\tan\beta \lesssim 4$, are ruled out because the lightest Higgs scalar has a mass which is considerably smaller than 125 GeV. In the SE$_6$SSM with $\lambda \gtrsim \sqrt{2}(M_Z/v) \simeq 0.5$, one can find the scenarios with 125 GeV SM-like Higgs for any $\tan\beta \gtrsim 2$. For so large values of λ, all Higgs bosons except the lightest CP-even Higgs scalar tend to have masses beyond the multi-TeV range so that they cannot be discovered at the LHC [48–50, 74]. To simplify our analysis, we set $\tan\beta \simeq 3$ and $\widetilde{\mu}_1 \simeq 2$ TeV.

The parameter space in the E_6 inspired $U(1)$ extensions of the MSSM is strongly constrained by the LHC experimental lower bounds on the Z' masses. Such gauge bosons are required to be heavier than 4.5 TeV [99, 100]. For $\langle S\rangle \simeq \langle\overline{S}\rangle$, the Z' mass in the SE$_6$SSM is given by

$$M_{Z'} \approx 2g'_1 Q_S\,\langle S\rangle. \tag{29}$$

Assuming gauge coupling unification, one can compute the low energy value of g'_1. With this value of g'_1, the VEVs $\langle S\rangle \simeq \langle\overline{S}\rangle$ have to be larger than 6 TeV to ensure that $M_{Z'} \gtrsim 4.5$ TeV.

The evaluated χ_1-nucleon cross sections σ_{SI}, σ^n and σ^p have to be compared with the most stringent experimental bounds set by the LZ [57] and IceCube [101] recently. However, one needs to take into account that for $m_{\chi_1} \approx \mu_{11} \lesssim 1$ TeV, the lightest exotic neutralino states compose only some fraction of the cold dark matter density. Thereby, the experimental limits become weaker, i.e.,

$$\sigma_{SI} < \sigma_{SI}^0 = \frac{(\Omega h^2)_{\rm exp}}{\Omega_{\tilde{H}}h^2}(\sigma_{SI})_{\rm exp}, \qquad \sigma^{p,n} < \sigma_0^{p,n} = \frac{(\Omega h^2)_{\rm exp}}{\Omega_{\tilde{H}}h^2}(\sigma^{p,n})_{\rm exp}. \tag{30}$$

Here $\Omega_{\tilde{H}}h^2$, σ_{SI} and $\sigma^{p,n}$ are calculated values of the cold dark matter density and χ_1-nucleon cross sections for each set of μ_{11}, f_{11} and $\tilde{f}_{11}$. At the same time, $(\sigma_{SI})_{\rm exp}$ and $(\sigma^{p,n})_{\rm exp}$ are the experimental bounds on the spin-independent χ_1-nucleon cross section as well as spin-dependent

χ_1-proton and χ_1-neutron cross sections at the given mass m_{χ_1}. Any set of parameters, which does not satisfy the conditions (30), is ruled out.

The results of our analysis are presented in Figures 1 and 2. We consider two scenarios associated with $\mu_{11} \simeq 500$ GeV and $\mu_{11} \simeq 1$ TeV. Equations (15), (20), (23) and (28) indicate that for fixed values of $\tan\beta$, $\tilde{\mu}_1$ and μ_{11}, the conditions (30) set limits on the Yukawa couplings f_{11} and $\tilde{f}_{11}$. Therefore, in Figures 1 and 2, the allowed regions in the $\tilde{f}_{11} - f_{11}$ plane are shown. To guarantee the validity of perturbation theory up to the GUT scale M_X, we require the absolute values of f_{11} and $\tilde{f}_{11}$ to be smaller than 0.4.

As follows from Eqs. (28), the computed values of σ^p and σ^n are always relatively close. On the other hand, the experimental limits $\sigma^p_{\rm exp}$ and $\sigma^n_{\rm exp}$ are very different. Taking into account that the lightest exotic neutralino states annihilate mainly into a pair of gauge bosons, the experimental bounds on the spin-dependent WIMP-proton scattering cross section $\sigma^p_{\rm exp} \approx 6 \cdot 10^3$ zb $(10^4$ zb) for $m_{\chi_1} \approx 500$ GeV(1 TeV), where 1 zb $= 10^{-45}$ cm^2 [101]. The spin-dependent WIMP-neutron scattering cross section is more tightly constrained, i.e., $\sigma^n_{\rm exp} \approx 2300$ zb (5000 zb) for $m_{\chi_1} \approx 500$ GeV(1 TeV) [57]. These experimental bounds set rather weak constraints on the Yukawa couplings f_{11} and $\tilde{f}_{11}$. In Figures 1 and 2, the regions of the parameter space, where $\sigma^n < \sigma^n_0/10$, are limited by the solid lines.

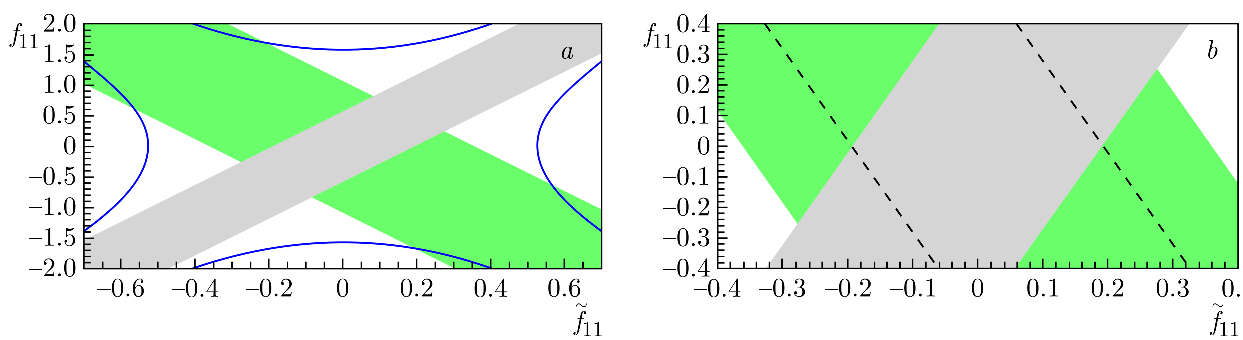

Figure 1. The constraints on the SE$_6$SSM parameter space in the $\tilde{f}_{11} - f_{11}$ plane for $\tan\beta = 3$, $\tilde{\mu}_1 = 2$ TeV and $\mu_{11} = 500$ GeV. Green region marks the part of the parameter space, where the experimental restriction on σ_{SI} is satisfied. In the grey area, $\Delta_2 \leqslant 200$ MeV. *a*) The area limited by the solid lines corresponds to the region, where $\sigma^n < \sigma^n_0/10$. *b*) The dashed-dotted lines limit the part of the parameter space, where $\sigma_{SI} < \sigma^0_{SI}/10$.

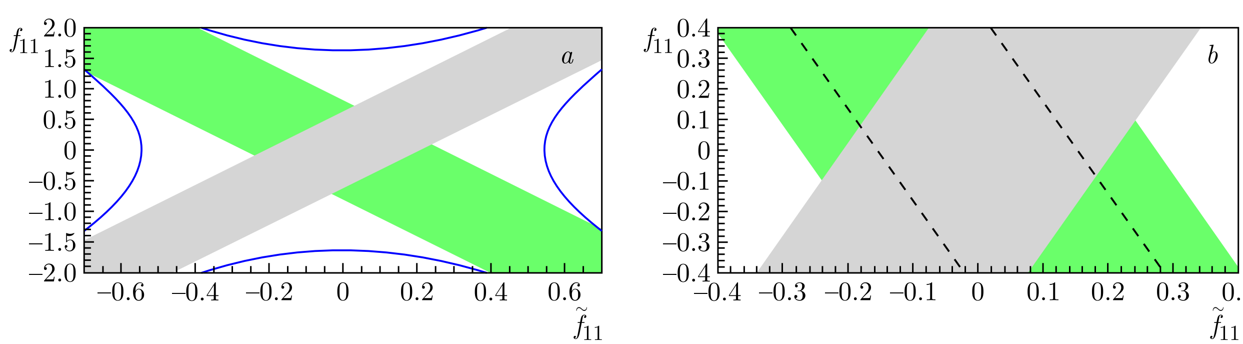

Figure 2. The constraints on the SE$_6$SSM parameter space in the $\tilde{f}_{11} - f_{11}$ plane for $\tan\beta = 3$, $\tilde{\mu}_1 = 2$ TeV and $\mu_{11} = 1000$ GeV. Green region marks the part of the parameter space, where the experimental restriction on σ_{SI} is satisfied. In the grey area, $\Delta_2 \leqslant 200$ MeV. *a*) The area limited by the solid lines corresponds to the region, where $\sigma^n < \sigma^n_0/10$. *b*) The dashed-dotted lines limit the part of the parameter space, where $\sigma_{SI} < \sigma^0_{SI}/10$.

The most stringent restrictions on the Yukawa couplings f_{11} and $\tilde{f}_{11}$ come from experimental bounds on the spin-independent WIMP-nucleon scattering cross section. The corresponding experimental limits are 12 yb for $m_{\chi_1} \approx 500$ GeV and 30 yb for $m_{\chi_1} \approx 1$ TeV (1 yb $= 10^{-48}$ cm^2) which are the values of $(\sigma_{SI})_{\rm exp}$ obtained by the LZ experiment [57]. The green areas in Figures 1 and 2 represent the allowed ranges of the Yukawa couplings f_{11} and $\tilde{f}_{11}$, where $\sigma_{SI} < \sigma^0_{SI}$. The requirement $\Delta_2 \leqslant 200$ MeV disfavours the grey regions. The areas between dashed-dotted lines correspond to the regions of the parameter space, where $\sigma_{SI} < \sigma^0_{SI}/10$. Thus, even if the experimental bounds become considerably more stringent, the scenarios under consideration are not going to be ruled out.

4. Conclusions

The cancellation of anomalies in the E_6 inspired $U(1)$-extended SUSY models requires that the low energy matter content of these models involves three fundamental 27 representations of E_6. Three 27-plets, in particular, includes three SM singlet superfields S_i and three families of Higgs-like doublets H^u_i and H^d_i. One pair of such Higgs-like supermultiplets ($H_d \equiv H^d_3$ and $H_u \equiv H^u_3$) acquires VEVs breaking the EW symmetry. The fermion components of two other families H^u_α and H^d_α $(\alpha = 1, 2)$ as well as the fermion components of the SM singlet superfields S_i, which do not gain VEVs, compose exotic neutralino and chargino states. The lightest exotic neutralino in these $U(1)$ extensions of the MSSM can be stable forming some part of the cold dark matter density.

In this article, we examined the constraints on the couplings of such lightest exotic neutralino within a specific extension of the MSSM with extra $U(1)_N$ gauge symmetry (SE$_6$SSM), in which the single discrete $\tilde{Z}^H_2$ symmetry forbids the most dangerous baryon and lepton number violating operators as well as tree-level non-diagonal flavour transitions. In addition to three 27-plets, the low energy spectrum of the SE$_6$SSM is supplemented by four E_6 singlet superfields, a pair of $SU(2)_W$ lepton doublets L_4 and $\overline{L}_4$ with opposite $SU(2)_W \times U(1)_Y \times U(1)_N$ quantum numbers and a pair of the SM singlet superfields S and $\overline{S}$ with opposite $U(1)_N$ charges. The superfields S and $\overline{S}$ can acquire very large VEVs ($\langle S \rangle \simeq \langle \bar{S} \rangle \gg 10$ TeV) resulting in the breakdown of the $U(1)_N$ gauge symmetry and inducing sufficiently large masses of all extra exotic particles. The supermultiplets L_4 and $\overline{L}_4$ facilitate the gauge coupling unification and permit the lightest exotic quark/squark to decay before BBN.

The conservation of the $\tilde{Z}^H_2$ symmetry and R-parity in the SE$_6$SSM leads to two stable neutral states which can form cold dark matter density. Here we assume that one of these stable states is gravitino with mass $m_{3/2} \sim 1$ GeV, while another stable particle is the lightest exotic neutralino χ_1 composed of the fermion components of H^d_1 and H^u_1. In this case, the lightest exotic chargino $\chi^\pm_1$ as well as the lightest and second lightest exotic neutralinos (χ_1 and χ_2) are nearly degenerate. If the masses of these exotic particles are smaller than 1.1 TeV, they may result in the phenomenologically acceptable dark matter density.

The interactions of dark matter with the nucleons in the SE$_6$SSM are determined by the couplings of χ_1. In this article, we focused on the scenarios in which all exotic fermions and all scalars except the SM-like Higgs, χ_1, χ_2, $\chi^\pm_1$ and gravitino are heavier than 10 TeV. As a consequence, the main contributions to the spin-dependent χ_1-proton and χ_1-neutron cross sections (σ^p and σ^n) come from the diagrams with t-channel exchange of the Z boson, whereas the spin-independent χ_1-nucleon cross section σ_{SI} is dominated by the t-channel exchange of the SM-like Higgs scalar. Our analysis revealed that there is a significant part of the SE$_6$SSM parameter space, where σ_{SI} and $\sigma^{p,n}$ are considerably lower than the present experimental bounds.

The phenomenological viability of the SE$_6$SSM requires $\chi_1^{\pm}$, χ_2 and χ_1 to be lighter than 1.1 TeV. Otherwise, the contribution of the lightest exotic neutralino to the total dark matter density becomes larger than its measured value. If χ_1 is considerably lighter than 1.1 TeV, its annihilation cross section is sufficiently large giving rise to the relatively small density of χ_1. In this case, one can naively expect that the indirect signal from dark matter annihilation may be enhanced. However, the analysis of the similar scenarios performed within the MSSM indicates that the corresponding indirect signal gets weaker with diminishing of the lightest neutralino mass because the density of χ_1 decreases [102].

Since scenarios under consideration imply that $\chi_1^{\pm}$, χ_2 and χ_1 are lighter than 1.1 TeV, they could be produced at the LHC. Nevertheless, when the masses of the lightest exotic chargino $m_{\chi_1^{\pm}}$, lightest and second lightest exotic neutralino (m_{χ_1} and m_{χ_2}) are rather close, the decay products of χ_2 and $\chi_1^{\pm}$ may escape detection. This also occurs within natural SUSY if the mass splitting between the lightest chargino and neutralino states is a few GeV or even smaller [103–105]. The results of the searches in this case depend on the values $\Delta_0 = m_{\chi_2} - m_{\chi_1}$ and $\Delta = m_{\chi_1^{\pm}} - m_{\chi_1}$. When $\Delta \simeq 2$ GeV(4.7 GeV), ATLAS ruled out the lightest chargino with mass below 140 GeV(193 GeV) [106], whereas CMS excluded the lightest chargino masses below 112 GeV for $\Delta = 1$ GeV [107]. For $\Delta \lesssim 150$ MeV, the lightest chargino may be long-lived, and LHC experiments ruled out such charginos if they are lighter than 1090 GeV [108].

The last most stringent experimental limit is not applicable because in the SE$_6$SSM scenarios under consideration $\Delta \gtrsim 300$ MeV [109, 110]. The second lightest exotic neutralino χ_2 cannot be long-lived as well. Indeed, at the tree level $\Delta_0 = \Delta_1 + \Delta_2 \gtrsim \Delta_2$, whereas in our analysis we require $\Delta_2 \geqslant 200$ MeV. Thus, the lightest exotic chargino and the second lightest exotic neutralino decay into χ_1 and hadrons. Since $\chi_1^{\pm}$, χ_2 and χ_1 are nearly degenerate, it seems to be rather problematic to discover this set of states at hadron colliders. The discovery prospects for such exotic fermions look more promising at future International Linear Collider.

Conflicts of interest

The authors declare no conflicts of interest.

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