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Cold dark matter in the SE_6SSM

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The supersymmetric (SUSY) E_6 inspired $U(1)_N$ extension of the Standard Model includes exotic fermions to ensure anomaly cancellation. The cold dark matter in this SUSY model (SE_6SSM) can be formed by the gravitino with mass $m_{3/2} \lesssim 1$ GeV and the lightest neutral exotic fermion. We examine the spin-independent dark matter nucleon scattering cross-section within the SE_6SSM and argue that it can be substantially smaller than the current experimental bounds.

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1. Introduction

In supersymmetric (SUSY) Grand Unified Theories (GUTs) with E_6 gauge symmetry [1] E_6 may be broken so that

$$\begin{aligned} E_6 \to SO(10)\times U(1)_\psi \to SU(5)\times U(1)_\chi \times U(1)_\psi \to \\ \to SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_\psi \times U(1)_\chi \,, \end{aligned} \quad (1)$$

where $SU(3)_C \times SU(2)_W \times U(1)_Y$ is a gauge group of the Standard Model (SM). If $U(1)_\chi \times U(1)_\psi$ symmetry is broken to its discrete subgroup $Z_2^M = (-1)^{3(B-L)}$ called matter parity, where B and L are baryon and lepton numbers, then such SUSY GUTs may lead to the minimal supersymmetric standard model (MSSM), Next-to-minimal supersymmetric standard model (NMSSM) and its extensions at low energies. The E_6 GUTs can result in the $U(1)'$ extensions of the MSSM when $U(1)_\chi \times U(1)_\psi$ is reduced to $U(1)'$ where

$$U(1)' = U(1)_\chi \cos\theta_{E_6} + U(1)_\psi \sin\theta_{E_6} \,. \quad (2)$$

In these SUSY extensions of the SM the anomalies are automatically cancelled if the low energy particle spectrum involves three complete 27-plets of E_6 (27_i with $i = 1, 2, 3$). Each 27-plet contains one generation of quarks (Q_i, u_i^c, d_i^c) and leptons (L_i, e_i^c, N_i^c), two $SU(2)_W$–doublets (H_i^d and H_i^u) that have quantum numbers of the MSSM Higgs doublets, two colour triplets ($\overline{D}_i$ and D_i) associated with exotic quarks with electric charges $\pm 1/3$ as well as a SM singlet field S_i, which carries non-zero $U(1)'$ charge.

The exceptional supersymmetric standard model (E$_6$SSM) [2, 3] (for the review, see Ref. [4, 5]) is the $U(1)_N$ extension of the MSSM, that corresponds to $\theta_{E_6} = \arctan\sqrt{15}$. It implies that near the GUT scale $U(1)_\psi \times U(1)_\chi \to U(1)_N \times Z_2^M$. In the E$_6$SSM the right-handed neutrinos N_i^c do not participate in the $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$ interactions [2, 3]. Therefore, they may be superheavy shedding light on the origin of the mass hierarchy in the lepton sector. The heavy right-handed neutrinos decay into final states with lepton number $L = \pm 1$, inducing lepton and baryon asymmetries in the Universe [6–8].

As in other E_6 inspired $U(1)'$ extensions of the MSSM extra exotic matter in the E$_6$SSM may lead to flavor changing neutral currents (FCNCs) and rapid proton decay. Since 2006, several modifications of the E$_6$SSM, in which the corresponding operators are suppressed, have been explored [2, 3, 9–15]. Within the E$_6$SSM the upper bound on the mass of the lightest Higgs boson near the quasi-fixed point was examined in [16]. Such quasi-fixed point is an intersection of the quasi–fixed and invariant lines [17, 18]. The particle spectrum within the constrained E$_6$SSM and its modifications was analyzed in [19–21]. Extra exotic states in these models may result in distinctive LHC signatures [2, 3, 11, 12, 22–24]. They can also lead to non-standard Higgs decays [10, 25, 26].

Using the approach considered in [27], it was revealed that within the E$_6$SSM and its simplest extensions the lightest SUSY particles (LSPs) are linear combinations of the fermion components of the superfields S_i [26]. In the simplest phenomenologically viable scenarios, these states are considerably lighter than 1 eV. Such lightest exotic fermions may compose hot dark matter in our Universe. The existence of very light neutral fermions may also result in some interesting implications for the neutrino physics [28].

Here, we investigate the dark matter-nucleon scattering cross section in the variant of the E$_6$SSM in which the most dangerous baryon number violating operators as well as non-diagonal tree-level flavor transitions are forbidden by a single discrete $\tilde{Z}_2^H$ symmetry. Below the GUT scale M_X the matter content of this model (SE$_6$SSM) involves three complete 27–plets, a set of supermultiplets M_l and $\overline{M}_l$ from additional $27'_l$ and $\overline{27}'_l$ representations as well as four E_6 singlet superfields (ϕ and ϕ_i). The set of supermultiplets M_l and $\overline{M}_l$ contains superfields S and $\overline{S}$. It also includes three pairs of $SU(2)_W$ doublets, i.e. H_u and $\overline{H}_u$, H_d and $\overline{H}_d$, L_4 and $\overline{L}_4$. Only supermultiplets H_d, H_u, L_4, $\overline{L}_4$, S, $\overline{S}$ and ϕ are even under the $\tilde{Z}_2^H$ symmetry [9, 10]. All other supermultiplets are odd under this symmetry. In the simplest scenario $\overline{H}_u$ and $\overline{H}_d$ get combined with the superposition of the appropriate components from the 27_i, forming vectorlike states with masses of order M_X. The components of L_4, $\overline{L}_4$, S, $\overline{S}$, ϕ and ϕ_i are expected to gain the TeV scale masses. Thus in the SE$_6$SSM the components of the supermultiplets

$$\begin{aligned}(Q_i,\, u^c_i,\, d^c_i,\, L_i,\, e^c_i,\, N^c_i) + (D_i,\, \bar{D}_i) + S_i + \phi_i + (H^u_\alpha,\, H^d_\alpha)\\ +L_4 + \overline{L}_4 + S + \overline{S} + H_u + H_d + \phi\,,\end{aligned} \tag{3}$$

have masses which are many orders of magnitude smaller than M_X ($i = 1, 2, 3$ and $\alpha = 1, 2$). The presence of L_4 and $\overline{L}_4$ with the TeV scale masses facilitates the generation of the baryon asymmetry of the Universe [7] and gauge coupling unification [9]. They also permit the lightest exotic quarks to decay within a reasonable time. The existence of the fermion components of the superfields ϕ_i at low energies allows to avoid the appearance of the exotic fermions with tiny masses in the particle spectrum.

The most general renormalisable superpotential of the SE$_6$SSM is given by [8, 15]

$$\begin{aligned}W_{\rm SE_6SSM} = \lambda S(H_u H_d) - \sigma\phi S\overline{S} + \frac{\kappa}{3}\phi^3 + \frac{\mu}{2}\phi^2 + \Lambda\phi + \mu_L L_4\overline{L}_4 + \tilde{\sigma}\phi L_4\overline{L}_4 + W_{IH}\\ +\kappa_{ij} S(D_i\overline{D}_j) + g^D_{ij}(Q_i L_4)\overline{D}_j + h^E_{i\alpha} e^c_i (H^d_\alpha L_4) + g_{ij}\phi_i \overline{L}_4 L_j + W_N + W_{\rm MSSM}(\mu = 0)\,,\end{aligned} \tag{4}$$

where $\alpha, \beta = 1, 2$ while $i, j = 1, 2, 3$ and

$$W_{IH} = \tilde{M}_{ij}\phi_i\phi_j + \tilde{\kappa}_{ij}\phi\phi_i\phi_j + \tilde{\lambda}_{ij}\overline{S}\phi_i S_j + \lambda_{\alpha\beta} S(H^d_\alpha H^u_\beta) + \tilde{f}_{i\alpha} S_i (H^d_\alpha H_u) + f_{i\alpha} S_i (H_d H^u_\alpha)\,, \tag{5}$$

$$W_N = \frac{1}{2} M_{ij} N^c_i N^c_j + \tilde{h}_{ij} N^c_i (H_u L_j) + h_{i\alpha} N^c_i (H^u_\alpha L_4)\,. \tag{6}$$

The breakdown of the $SU(2)_W \times U(1)_Y \times U(1)_N$ symmetry in the SE$_6$SSM is caused by the vacuum expectation values (VEVs) of ϕ, S, $\overline{S}$, H_u and H_d. The $U(1)_N$ symmetry forbids the term $\mu_0 H_d H_u$ in the superpotential (4). The sum of all other terms, which are present in the MSSM superpotential, is denoted as $W_{\rm MSSM}(\mu = 0)$. In the SE$_6$SSM S and $\overline{S}$ may develop VEVs along the D-flat direction, i.e. $\langle S\rangle \simeq \langle\overline{S}\rangle \simeq S_0$. If $\sigma \to 0$ the value of S_0 can be much larger than the sparticle mass scale M_S. As a consequence the Z' boson and all extra exotic states tend to be rather heavy in this case.

The paper is organised as follows. In the next section, we focus on the SE$_6$SSM scenario in which the cold dark matter is formed by the gravitino with mass $m_{3/2} \lesssim 1$ GeV and the lightest exotic fermion. In section 3 we explore the interactions of the dark matter states with the nucleons. Our results are summarised in section 4.

2. Dark matter in the SE$_6$SSM

For the analysis of the SE$_6$SSM it is convenient to introduce the Z_2^E symmetry defined as $\tilde{Z}_2^H = Z_2^M \times Z_2^E$ [9]. The components of the supermultiplets L_4, $\overline{L}_4$, S_i, ϕ_i, H^d_α, H^u_α, $\overline{D}_i$ and D_i are odd under the Z_2^E symmetry whereas all other supermultiplets are Z_2^E even. The invariance of the SE$_6$SSM Lagrangian with respect to Z_2^M and $\tilde{Z}_2^H$ symmetries implies that Z_2^E symmetry and R-parity are also preserved. As a result the lightest R-parity odd state is stable. Thus it can contribute to the dark matter density. Here we consider the SE$_6$SSM scenario in which the superpartner of graviton (gravitino) is the lightest R-parity odd state that has a mass which is substantially smaller than the masses of all other new particles in the model under consideration. Because the Z_2^E symmetry is also conserved and gravitino is a Z_2^E even state, the lightest exotic particle, i.e. lightest Z_2^E odd state, has to be stable as well [9].

The fermion components of the supermultiplets H^u_α, H^d_α, S_i and ϕ_i, that do not acquire VEVs, compose the exotic (inert) chargino and neutralino states. If all components of ϕ_i are significantly heavier than the bosons and fermions from H^u_α, H^d_α and S_i, the superfields ϕ_i can be integrated out. As a consequence the part of the SE$_6$SSM superpotential (5) reduces to

$$W_{IH} \to \widetilde{W}_{IH} \simeq -\widetilde{\mu}_{ij} S_i S_j + \lambda_{\alpha\beta} S (H^d_\alpha H^u_\beta) + \tilde{f}_{i\alpha} S_i (H^d_\alpha H_u) + f_{i\alpha} S_i (H_d H^u_\alpha) + \ldots \qquad (7)$$

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

In this article we study the scenarios with the lightest exotic states formed by the fermion components of H^u_1 and H^d_1. It is expected that all sparticles except gravitino as well as all other exotic states have masses which are substantially larger than 1 TeV. It is also assumed that H^u_1 and H^d_1 mostly interact with S_1, H_u and H_d. To simplify our analysis we ignore all other couplings of H^u_1 and H^d_1. Then the mass of the lightest charged exotic fermion $\chi^\pm_1$ is given by

$$m_{\chi^\pm_1} \simeq |\mu_{11}|, \qquad\qquad \mu_{11} \simeq \lambda_{11} \langle S \rangle\,.$$

If $|\widetilde{\mu}_1|$ is much larger than the EW scale and $|\mu_{11}|$ the mass matrix of the lightest exotic neutralino can be diagonalised using the perturbation theory method (see, for example, [29–32]) that yields

$$\begin{array}{lll} m_{\chi_1} \simeq m_{\chi^\pm_1} - \Delta_1\,, & m_{\chi_2} \simeq m_{\chi^\pm_1} + \Delta_2\,, & m_{\chi_3} \simeq \widetilde{\mu}_1 + \Delta_1 + \Delta_2\,, \\ \Delta_1 \simeq \dfrac{(\tilde{f}_{11} v \sin\beta + f_{11} v \cos\beta)^2}{4(\widetilde{\mu}_1 - m_{\chi^\pm_1})}\,, & & \Delta_2 \simeq \dfrac{(\tilde{f}_{11} v \sin\beta - f_{11} v \cos\beta)^2}{4(\widetilde{\mu}_1 + m_{\chi^\pm_1})}\,. \end{array} \qquad (8)$$

where $\tan\beta = v_2/v_1$ and $v = \sqrt{v_1^2 + v_2^2} \approx 246\,\text{GeV}$, while v_1 and v_2 are the VEVs of the Higgs doublets H_d and H_u, i.e. $\langle H_d \rangle = v_1/\sqrt{2}$ and $\langle H_u \rangle = v_2/\sqrt{2}$. In the leading approximation the masses of the lightest exotic neutralino (χ_1 and χ_2) and chargino $\chi^\pm_1$ states are determined by μ_{11}. Here we require that $m_{\chi_2} - m_{\chi_1} > 200\,\text{MeV}$. In this case the second lightest exotic neutralino χ_2 decays before Big Bang Nucleosynthesis (BBN), i.e. the lifetime of χ_2 is shorter than 1 sec.

The lifetime of unstable exotic and R-parity odd particles Y that decay mainly into gravitino can be estimated as [33]

$$\tau_Y \sim 48\pi \frac{m^2_{3/2} M^2_{Pl}}{m^5_Y}\,, \qquad (9)$$

where m_Y is the mass of the state Y and $M_{Pl} = (8\pi G_N)^{-1/2} \simeq 2.4 \cdot 10^{18}$ GeV is the reduced Planck mass. The lifetime (9) should be also less than 1 sec. Otherwise the products of the decays of the corresponding particles may alter the abundances of light elements predicted by the BBN in the SM. For $m_Y \simeq 1$ TeV this requirement is satisfied if $m_{3/2} \lesssim 1$ GeV. Gravitinos originating from the scattering of particles in the thermal bath give rise to the dark matter density [34, 35]

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

Here $M_{\tilde{g}}$ is a gluino mass and T_R is a reheating temperature. When $m_{3/2} \simeq 1$ GeV and $M_{\tilde{g}} \gtrsim 3$ TeV the appropriate gravitino contribution to the dark matter density can be obtained for $T_R \lesssim 10^{6-7}$ GeV [36]. Nevertheless even for so low T_R the baryon asymmetry in the SE$_6$SSM can be induced via the lightest right-handed neutrino (sneutrino) decays into exotic states [8].

Using the approximate formula [37, 38]:

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

one can estimate the contribution of the lightest neutral exotic neutralino χ_1 to the cold dark matter density. At the same time from the Planck observations it follows $(\Omega h^2)_{\text{exp}} = 0.1188 \pm 0.0010$ [39]. This leads to the upper bound $\mu_{11} < 1.1$ TeV. If μ_{11} is considerably smaller than 1.1 TeV, gravitino may account for some part of the observed density of the cold dark matter.

3. Spin-independent interaction of the dark matter with nucleons

In the SUSY model under consideration the interactions of the dark matter with baryons are defined by the couplings of χ_1 since the gravitino couplings to the SM particles are extremely small. The low energy Lagrangian $\mathcal{L}_{\chi_1 q}$, that describes the spin-independent interactions of χ_1 with quarks, and the corresponding χ_1-nucleon scattering cross section σ_{SI} can be written as

$$\mathcal{L}_{\chi_1 q} = \sum_q a_q \bar{\chi}_1 \chi_1 \bar{q} q, \qquad \sigma_{SI} = \frac{4 m_r^2}{\pi} \frac{(Z f^p + (A-Z) f^n)^2}{A^2}, \qquad m_r = \frac{m_{\chi_1} m_N}{m_{\chi_1} + m_N}, \qquad (12)$$

where m_N is a nucleon mass, Z and A are the charge and nucleon number of the target nucleus, $f^p \approx f^n \approx f^N$ and

$$\frac{f^N}{m_N} = \sum_{q=u,d,s} \frac{a_q}{m_q} f^N_{Tq} + \frac{2}{27} \sum_{Q=c,b,t} \frac{a_Q}{m_Q} f^N_{TQ}, \qquad (13)$$

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

Here the hadronic matrix elements are set to be equal to the default values used in micrOMEGAs, i.e. $f^N_{Tu} \simeq 0.0153$, $f^N_{Td} \simeq 0.0191$ and $f^N_{Ts} \simeq 0.0447$ [40].

The exotic neutralino states in the SE$_6$SSM do not interact with quarks and squarks. Therefore the only contributions to a_q come from the t–channel exchange of the Higgs states. Because in the SE$_6$SSM all extra Higgs scalars are expected to be substantially heavier than 1 TeV, all contributions which are caused by the so heavy Higgs exchange can be ignored. Since in the limit, when the

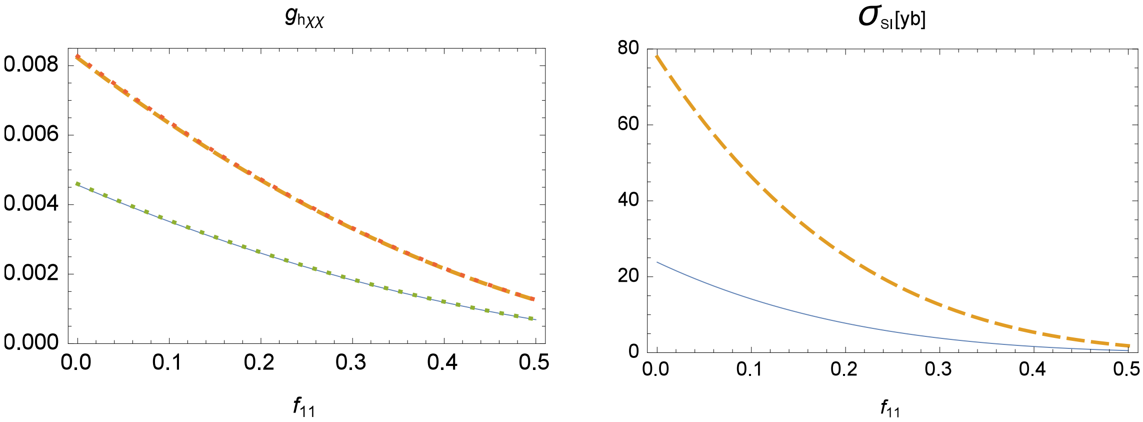

Figure 1: The coupling $g_{h\chi\chi}$ (left) and the cross–section σ_{SI} (right) as a function of f_{11} for $\tilde{f}_{11} = -0.41$, $\tan\beta = 2$, $\tilde{\mu}_1 = 2\,\text{TeV}$, $\mu_{11} = 200\,\text{GeV}$ (solid lines) and $\mu_{11} = 1\,\text{TeV}$ (dashed lines). The dotted lines correspond to the approximate expression for $g_{h\chi\chi}$ (15).

sparticle mass scale is much larger than v, the lightest Higgs state h_1 with mass $m_{h_1} \approx 125\,\text{GeV}$ has couplings which are almost identical to the ones in the SM, the coefficients a_q are given by

$$\frac{a_q}{m_q} \simeq \frac{a_Q}{m_Q} \simeq \frac{g_{h\chi\chi}}{v m_{h_1}^2}\,, \tag{14}$$

where $g_{h\chi\chi}$ is the coupling of χ_1 to h_1. If $\Delta_1 \ll \mu_{11} \ll \tilde{\mu}_1$ then

$$|g_{h\chi\chi}| \simeq \frac{\Delta_1}{v}\,. \tag{15}$$

The mass and coupling $g_{h\chi\chi}$ of χ_1 are determined by μ_{11}, $\tilde{\mu}_1$, f_{11}, $\tilde{f}_{11}$ and $\tan\beta$. The interval of variation of the parameter μ_{11} is fixed so that $\mu_{11} \lesssim 1\,\text{TeV}$ and $\mu_{11} \gtrsim 200\,\text{GeV}$. This permits us to avoid the lower experimental bound on $\chi^{\pm}_1$. It also ensures that χ_1 results in the phenomenologically acceptable density of the dark matter. The range of variations of $\tilde{f}_{11}$ and f_{11} at low energies is constrained by the requirement of the applicability of perturbation theory up to the scale M_X. To simplify our analysis here we set $\tan\beta \simeq 2$ and $\tilde{\mu}_1 \simeq 2\,\text{TeV}$. For moderate values of $\tan\beta$ the 125 GeV Higgs state can be obtained only if $\lambda \gtrsim 0.5$. When λ is so large all additional Higgs bosons tend to have masses beyond the multi-TeV range [2, 3, 22, 23]. Therefore they cannot be detected at the LHC.

The results of our analysis presented in Figure 1 indicate that the approximate expression (15) describes reasonably well the dependence of $g_{h\chi\chi}$ on the SE$_6$SSM parameters. Indeed, $g_{h\chi\chi}$ grows when μ_{11} increases from 200 GeV to 1 TeV. From Eq. (15) it follows that $g_{h\chi\chi}$ vanishes for $f_{11} \approx -\tilde{f}_{11}\tan\beta$. Thus in some part of the SE$_6$SSM parameter space the interactions of χ_1 with nucleons tend to be extremely weak. In our analysis we set f_{11} to be positive while $\tilde{f}_{11}$ is chosen to be negative. As a result $g_{h\chi\chi}$ and σ_{SI} diminish when f_{11} grows and approaches $-\tilde{f}_{11}\tan\beta$. If $|g_{h\chi\chi}|$ becomes smaller than 10^{-3} the contributions to σ_{SI} induced by the t-channel exchange of additional heavy Higgs states cannot be ignored. In this case one also needs to take into account the quantum corrections that stem from the diagrams involving the SM gauge bosons [41, 42]. The inclusion of such corrections lead to $\sigma_{SI} \sim 0.1$ yb even if $|g_{h\chi\chi}| \ll 10^{-3}$ [43].

The cross section σ_{SI} shown in Figure 1 is always smaller than 300 yb for $m_{\chi_1} \approx \mu_{11} \approx 1$ TeV and 60 yb for $m_{\chi_1} \approx \mu_{11} \approx 200$ GeV (1 yb $= 10^{-48}$ cm^2) which are current experimental limits [44, 45]. The maximal values of σ_{SI} in the SE$_6$SSM are substantially larger than the spin–independent χ_1–nucleon scattering cross section presented in Figure 1. The corresponding cross section attains its maximal value for $\tilde{f}_{11} \sim f_{11} \sim 1$ and $\mu_{11} \simeq \widetilde{\mu}_1$. It can reach $20 - 30$ zb. In the part of the SE$_6$SSM parameter space explored in Figure 1 the suppression of σ_{SI} is caused by the large $\widetilde{\mu}_1$ and by the partial cancellation of different contributions to $g_{h\chi\chi}$.

4. Conclusions

The particle spectrum of the exceptional supersymmetric standard model (E$_6$SSM) involves at least three fundamental representations of E_6 to ensure anomaly cancellation. These three 27-plets, in particular, include three SM singlet superfields S_i with non-zero $U(1)_N$ charges as well as three families of Higgs-like doublets H^u_i and H^d_i. Here we considered the variant of the E$_6$SSM (SE$_6$SSM) in which two pairs of Higgs-like doublets (H^u_α and H^d_α, where $\alpha = 1, 2$) and three superfields S_i do not acquire VEVs. The fermion components of these supermultiplets compose the exotic neutralino and chargino states.

In this paper we explored the scenarios in which the cold dark matter is formed by the lightest exotic neutralino and gravitino with mass $m_{3/2} \lesssim 1$ GeV. We assumed that the lightest exotic neutralino state χ_1 is composed of the neutral fermion components of H^u_1 and H^d_1. As a consequence χ_1, the second lightest exotic neutralino χ_2 and the lightest exotic chargino $\chi^\pm_1$ are nearly degenerate. If the masses of these states are smaller than 1.1 TeV they give rise to the phenomenologically acceptable value of the dark matter density.

Several experiments have searched for almost degenerate neutralino and chargino states. When the mass splitting between such states is rather small the decay products of $\chi^\pm_1$ and χ_2 tend to be quite soft and may escape detection. The results of the searches for these states depend on $\Delta_0 = m_{\chi_2} - m_{\chi_1}$ and $\Delta = m_{\chi^\pm_1} - m_{\chi_1}$. In our analysis the parameters of the SE$_6$SSM are chosen so that $\Delta_1 > 200$ MeV whereas Δ is always larger than 300 MeV [43]. Thus $\chi^\pm_1$ and χ_2 are not long–lived particles. At the LHC χ_1, χ_2 and $\chi^\pm_1$ are pair produced via off-shell Z and W–bosons. Then χ_2 and $\chi^\pm_1$ decays mainly into hadrons and χ_1. ATLAS ruled out such exotic chargino with masses below 140 GeV (193 GeV) for $\Delta \simeq 2$ GeV (4.7 GeV) [46]. CMS excluded the corresponding chargino with masses below 112 GeV if $\Delta \simeq 1$ GeV [47].

Because of the weakness of the interactions of gravitino with the SM particles we only studied the dependence of the χ_1-nucleon coupling on the SE$_6$SSM parameters. We focused on the scenarios in which the masses of all extra states beyond the SM except χ_1, χ_2, $\chi^\pm_1$ and gravitino are of the order of a few TeV. Therefore, the spin-independent χ_1-nucleon scattering cross-sections σ_{SI} is dominated by the t-channel exchange of the lightest Higgs boson. The results of our analysis indicate that there is a part of the SE$_6$SSM parameter space where σ_{SI} is considerably smaller than the current experimental bounds.

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