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Leptogenesis and Dark Matter–Nucleon Scattering Cross Section in the SE$_6$SSM

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Abstract: The E_6-inspired extension of the minimal supersymmetric (SUSY) standard model (MSSM) with an extra $U(1)_N$ gauge symmetry, under which right-handed neutrinos have zero charge, involves exotic matter beyond the MSSM to ensure anomaly cancellation. We consider the variant of this extension (SE$_6$SSM) in which the cold dark matter is composed of the lightest neutral exotic fermion and gravitino. The observed baryon asymmetry can be induced in this case via the decays of the lightest right-handed neutrino/sneutrino into exotic states even for relatively low reheating temperatures $T_R \lesssim 10^{6-7}$ GeV. We argue that there are some regions of the SE$_6$SSM parameter space, which are safe from all current constraints, and discuss the implications of this model for collider phenomenology.

Keywords: unified field theories and models; models beyond the standard model; supersymmetry; cold dark matter; leptogenesis

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1. Introduction

The observed baryon asymmetry and the presence of cold dark matter in the Universe stimulates the investigation of extensions of the Standard Model (SM). New physics beyond the SM permits to induce the baryon asymmetry if the Sakharov conditions are fulfilled [1]. The proposed new physics scenarios include baryogenesis in Grand Unified theories (GUTs) [2–8], the Affleck–Dine mechanism [9,10], baryogenesis via leptogenesis [11], electroweak baryogenesis [12,13], etc . In the case of thermal leptogenesis [11] lepton asymmetry is generated due to the decays of the lightest right-handed neutrino. The realisation of this mechanism within the type I seesaw models [14], in which CP and lepton number are violated, allows for understanding of the mass hierarchy in the lepton sector if the right-handed neutrinos are superheavy. In this scenario the induced lepton asymmetry is partially converted into baryon asymmetry via sphaleron processes [15,16].

After inflation in the reheating epoch, which is characterized by a reheat temperature T_R, the right-handed neutrinos can be produced by thermal scattering if $T_R > M_1$. In the SM and minimal supersymmetric (SUSY) standard model (MSSM) such production process results in the appropriate baryon asymmetry only when the lightest right-handed neutrino mass M_1 is larger than 10^9 GeV [17,18]. Therefore thermal leptogenesis in the MSSM and its extensions may take place when $T_R \gtrsim 10^9$ GeV. This lower bound on the reheat temperature leads to the gravitino problem [19,20] in the supergravity (SUGRA) models, that lead to the sparticle mass scale below 10 TeV. Indeed, such a high T_R gives rise to an overproduction of gravitinos. Since gravitinos are sufficiently long-lived they tend to decay after Big Bang Nucleosynthesis (BBN). Such decays destroy the agreement between the predicted and observed light element abundances. To preserve the success of BBN the relic abundance of gravitinos has to be relatively small. It becomes low enough if reheat temperature is lower than 10^{6-7} GeV [21–23].

In this context it seems to be interesting to study the generation of matter–antimatter asymmetry and formation of cold dark matter in the framework of well-motivated E_6-inspired extensions of the SM. In the E_6-inspired composite Higgs model (E$_6$CHM) [24,25] the process of the baryon asymmetry generation was explored in [26,27]. The E_6-inspired $U(1)$ extensions of the MSSM implies that near the GUT scale M_X the E_6 gauge group is broken down to $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)'$ (for review see [28,29]) where $SU(3)_C \times SU(2)_W \times U(1)_Y$ is the SM gauge group and

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

In Equation (1) the $U(1)_\psi$ and $U(1)_\chi$ symmetries are associated with the subgroups $E_6 \supset SO(10) \times U(1)_\psi$ and $SO(10) \supset SU(5) \times U(1)_\chi$. The E_6-inspired $U(1)$ extensions of the MSSM can originate from the heterotic superstring theory with $E_8 \times E_8'$-gauge symmetry. Some phenomenological consequences of the heterotic string model were considered in [30,31]

Within the SUSY models with extra $U(1)'$ the anomalies are canceled if the particle spectrum contains complete representations of E_6. Because of this the particle spectrum in the models under consideration is usually extended by the supermultiplets of exotics so that it consists of three 27-dimensional representations of E_6 (27_i with $i = 1, 2, 3$) at low energies. These 27-plets decompose under $SU(5) \times U(1)_\psi \times U(1)_\chi$ as follows:

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

Here the first, second and third quantities in brackets are the $SU(5)$ representation as well as $U(1)_\psi$ and $U(1)_\chi$ charges. The SM family, which consists of the doublets of left-handed quarks Q_i and leptons L_i, right-handed up- and down-quarks (u_i^c and d_i^c) as well as right-handed charged leptons (e_i^c), corresponds 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 Equation (2), $\left(1, \frac{1}{\sqrt{24}}, -\frac{5}{\sqrt{40}}\right)_i$, represents the right-handed neutrinos N_i^c. The next-to-last term, $\left(1, \frac{4}{\sqrt{24}}, 0\right)_i$, is associated with new SM-singlet fields S_i, that carry non-zero $U(1)_\psi$ charges. The $SU(2)_W$-doublets (H_i^d and H_i^u) from $\left(5^*, -\frac{2}{\sqrt{24}}, -\frac{2}{\sqrt{40}}\right)_i$ and $\left(5, -\frac{2}{\sqrt{24}}, \frac{2}{\sqrt{40}}\right)_i$ have the quantum numbers of the MSSM Higgs doublets. The colour triplets from these $SU(5)$ supermultiplets are associated with exotic quarks $\overline{D}_i$ and D_i with electric charges $+1/3$ and $-1/3$ respectively. They carry a $B - L$ charge $\left(\pm\frac{2}{3}\right)$ which is twice as large as the $B - L$ charges of ordinary quarks.

Among the E_6-inspired $U(1)$ extensions of the MSSM there is a unique combination of $U(1)_\psi$ and $U(1)_\chi$ corresponding to $\theta_{E_6} = \arctan\sqrt{15}$ for which N_i^c do not participate in the gauge interactions. Only in this SUSY model with extra $U(1)_N$ gauge symmetry, i.e., the so-called Exceptional Supersymmetric Standard Model (E$_6$SSM) [32,33] (for recent review see [34]), the right-handed neutrinos can be rather heavy inducing the mass hierarchy in the lepton sector. The heavy Majorana right-handed neutrinos are allowed to decay into final states with lepton number $L = \pm 1$, resulting in lepton and baryon asymmetries in the early Universe [35–37]. In the $U(1)_N$ extensions of the MSSM the extra exotic states with the TeV scale masses can give rise to rapid proton decay and flavor-changing transitions. The corresponding operators can be suppressed in the E$_6$SSM using a set of discrete symmetries [32,33].

In this article we focus on the variant of the E$_6$SSM (SE$_6$SSM) in which a single $\tilde{Z}_2^H$ symmetry forbids non-diagonal flavor transitions and most dangerous operators that violate baryon and lepton numbers. In the next section the SE$_6$SSM is specified. In Section 3 the thermal leptogenesis within this $U(1)_N$ extension of the MSSM is considered. The interactions of the dark matter states with the nucleons is explored in Section 4. In Section 5 we summarize the results of our studies and discuss the implications of the SUSY model under consideration for collider phenomenology.

2. The $U(1)_N$ Extension of the MSSM with Exact Custodial $\tilde{Z}_2^H$ Symmetry

At very high energies the E_6 orbifold SUSY GUTs can be reduced to an effective rank-6 SUSY model based on the $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_\chi \times U(1)_\psi$ gauge symmetry [38]. If the particle content of this rank-6 model involves just three 27-plets at low energies then the most general renormalisable superpotential comes from the $27 \times 27 \times 27$ decomposition of E_6 and can be written as

$$\begin{aligned} W_{E_6} &= W_0 + W_1 + W_2\,, \\ W_0 &= \lambda_{ijk} S_i (H_j^d H_k^u) + \kappa_{ijk} S_i (D_j \overline{D}_k) + h_{ijk}^N N_i^c (H_j^u L_k) + h_{ijk}^U u_i^c (H_j^u Q_k) + \\ &+ h_{ijk}^D d_i^c (H_j^d Q_k) + h_{ijk}^E e_i^c (H_j^d L_k)\,, \\ W_1 &= g_{ijk}^Q D_i (Q_j Q_k) + g_{ijk}^q \overline{D}_i d_j^c u_k^c\,, \\ W_2 &= g_{ijk}^N N_i^c D_j d_k^c + g_{ijk}^E e_i^c D_j u_k^c + g_{ijk}^D (Q_i L_j) \overline{D}_k\,, \end{aligned} \tag{3}$$

where the summation over repeated indexes is implied and $i, j, k = 1, 2, 3$.

From Equation (3) it follows that if all Yukawa couplings in W_1 and W_2 have non-zero values then one cannot define the baryon (B) and lepton (L) numbers so that the Lagrangian of this model is invariant under the corresponding $U(1)_B$ and $U(1)_L$ global symmetries. Therefore, as in the simplest SUSY extensions of the SM, the gauge symmetry in the E_6-inspired SUSY models does not forbid the operators which violate lepton and baryon numbers. This means that in general these models lead to rapid proton decay. Moreover, since three pairs of H_i^u and H_i^d couple to charged leptons and ordinary quarks the corresponding Yukawa interactions may give rise to unacceptably large flavor changing processes at the tree level. In particular, these interactions can induce new channels of muon decay such as $\mu \to e^- e^+ e^-$ and contribute to the amplitude of K^0–$\overline{K}^0$ oscillations.

Although $U(1)_B$ and $U(1)_L$ symmetries are not conserved the superpotential (3) possesses $U(1)_{B-L}$ symmetry associated with $B-L$ number conservation if the exotic quark supermultiplets $\overline{D}_i$ (D_i) carry $B-L$ numbers $B-L = \dfrac{2}{3}(-\dfrac{2}{3})$. As a consequence $U(1)_\chi \times U(1)_\psi$ gauge symmetry can be broken down to matter parity $Z_2^M = (-1)^{3(B-L)}$, which is a discrete subgroup of $U(1)_{B-L}$. In the case of the E$_6$SSM $U(1)_\chi \times U(1)_\psi$ symmetry is expected to be broken to $U(1)_N \times Z_2^M$ near the GUT scale M_X [32,33]. Such breakdown can be attained if N_H^c and $\overline{N}_H^c$ components of some extra 27_H and $\overline{27}_H$ representations develop vacuum expectation values (VEVs) along the D-flat direction $< N_H^c > = < \overline{N}_H^c >$ [38]. These VEVs may also induce the Majorana mass terms of the right-handed neutrinos (i.e., $\dfrac{1}{2} M_{ij} N_i^c N_j^c$) in the superpotential with the intermediate scale mass parameters M_{ij} through the non-renormalizable operators of the form

$$\delta W = \frac{\varkappa_{ij}}{M_{Pl}} (\overline{27}_H\, 27_i)(\overline{27}_H\, 27_j) \qquad \Longrightarrow \qquad M_{ij} = \frac{2\varkappa_{ij}}{M_{Pl}} < \overline{N}_H^c >^2\,, \tag{4}$$

where $M_{Pl} = (8\pi G_N)^{-1/2} \simeq 2.4 \cdot 10^{18}$ GeV is the reduced Planck mass. When $< N_H^c > \simeq < \overline{N}_H^c > \simeq M_X \simeq 2 - 3 \cdot 10^{16}$ GeV, the observed pattern of masses and mixing angles of the left-handed neutrinos can be obtained.

Over the last fifteen years, several modifications of the E$_6$SSM, in which the operators leading to rapid proton decay and flavor changing processes are suppressed, have been proposed [32,33,38–49]. The implications of the $U(1)_N$ extensions of the MSSM were explored for Z' physics [50], neutralino sector [50–52], electroweak (EW) symmetry breaking (EWSB) [51,53,54], the renormalization group (RG) flow of couplings [51,55], the renormalization of VEVs [56,57], non-standard neutrino models [58] and dark matter [48,49,59–61]. Within the E$_6$SSM the upper bound on the lightest Higgs mass near the quasi-fixed point was studied in [62]. This quasi-fixed point is an intersection of the invariant and quasi-fixed lines [63,64]. The particle spectrum and corresponding phenomenological implications in the constrained E_6SSM (cE_6SSM) and its modifications were analyzed in [65–68]. The degree of fine tuning and threshold corrections were examined in [69,70] and [71] respectively. In the E_6SSM extra exotic matter may lead to distinctive LHC signatures [32,33,41,44,72–75] and can give rise to non-standard decays of the lightest Higgs boson [39,60,76–79].

In addition to three complete 27-plets the splitting of bulk $27'$ supermultiplets in the E_6 orbifold SUSY GUTs can result in a set of M_l and $\overline{M}_l$ supermultiplets from extra $27'_l$ and $\overline{27}'_l$ representations [38]. Since M_l and $\overline{M}_l$ have opposite quantum numbers all gauge anomalies still cancel. In the case of SE$_6$SSM the set of M_l and $\overline{M}_l$ includes a pair of superfields S and $\overline{S}$ as well as three pairs of $SU(2)_W$ doublets, i.e., H_d and $\overline{H}_d$, H_u and $\overline{H}_u$, L_4 and $\overline{L}_4$. Only supermultiplets L_4, $\overline{L}_4$, S, $\overline{S}$, H_d and H_u are required to be even under the $\tilde{Z}_2^H$ symmetry that forbids the tree-level flavor-changing transitions, as well as the most dangerous baryon and lepton number violating operators. All other supermultiplets are expected to be odd under this discrete symmetry [39].

The $\tilde{Z}_2^H$ symmetry forbids all terms in the SE$_6$SSM superpotential that come from $27_i \times 27_j \times 27_k$ where $i, j, k = 1, 2, 3$ are family indexes. Nevertheless it allows the interactions which originate from $27'_l \times 27'_m \times 27'_n$ and $27'_l \times 27_i \times 27_k$. Here indexes l, m, n are associated with the supermultiplets M_l. As a consequence $\tilde{Z}_2^H$ symmetry forbids all Yukawa interactions in W_1 eliminating the most dangerous operators leading to rapid proton decay. On the other hand this symmetry allows the terms $(Q_i L_4)\overline{D}_k$ in the superpotential that permits the lightest exotic colored state (quark or squark) to decay. In the SE$_6$SSM all charged leptons and the down-type quarks couple to just H_d while the up-type quarks interact with H_u only. Thus at tree-level non-diagonal flavor transitions are suppressed.

Using the results of the analysis presented in [80–82], it was shown that within the E_6SSM and its simplest modifications the lightest SUSY particles (LSPs) are linear superpositions of the fermion components of the superfields S_i [60,77]. In the simplest scenarios these states are either massless or have masses which are much smaller than 1 eV forming hot dark matter in our Universe. The presence of very light neutral fermions may have interesting implications for neutrino physics (see, for example [83]).

To avoid the appearance of exotic fermions with tiny masses it is assumed that the low-energy matter content of the SE$_6$SSM involves at least four E_6 singlet superfields. One of these superfields ϕ is even under the $\tilde{Z}_2^H$ symmetry whereas three others ϕ_i are odd. The SE$_6$SSM implies that $\overline{H}_u$ and $\overline{H}_d$ are combined with the superposition of the appropriate components from the 27_i, composing vectorlike states with masses of order M_X. At the same time the components of the supermultiplets S and $\overline{S}$ as well as L_4 and $\overline{L}_4$ gain the TeV scale masses. The presence of L_4 and $\overline{L}_4$ at low energies facilitates the gauge coupling unification [55] and permits the lightest exotic colored state (quark or squark) to decay within a reasonable time. As a result the components of the supermultiplets

$$\begin{gathered}(Q_i,\, u_i^c,\, d_i^c,\, L_i,\, e_i^c,\, N_i^c) + (D_i,\, \bar{D}_i) + S_i + \phi_i + (H_\alpha^u,\, H_\alpha^d) \\ + L_4 + \overline{L}_4 + S + \overline{S} + H_u + H_d + \phi\,,\end{gathered} \tag{5}$$

survive to low energies, i.e., they have masses which are many orders of magnitude smaller than M_X. Here $i = 1, 2, 3$ and $\alpha = 1, 2$. The $U(1)_N$ and $U(1)_Y$ charges of the supermultiplets listed in Equation (5) are given in Table 1. It is worth noting that the superfields ϕ_i, N_i^c and ϕ do not participate in the gauge interactions. Therefore these superfields are not included in Table 1.

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

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

The most general renormalisable superpotential of the SE$_6$SSM, which is allowed by the $\tilde{Z}_2^H$, Z_2^M and $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$ symmetries, is given by

$$\begin{aligned} W_{\mathrm{SE_6SSM}} = {} & \lambda S(H_u H_d) - \sigma\phi S\overline{S} + \frac{\kappa}{3}\phi^3 + \frac{\mu}{2}\phi^2 + \Lambda\phi + \mu_L L_4\overline{L}_4 + \tilde{\sigma}\phi L_4\overline{L}_4 + W_{IH} \\ & + \kappa_{ij} S(D_i\overline{D}_j) + g^D_{ij}(Q_i L_4)\overline{D}_j + h^E_{i\alpha} e^c_i (H^d_\alpha L_4) + g_{ij}\phi_i \overline{L}_4 L_j + W_N + W_{\mathrm{MSSM}}(\mu = 0)\,. \end{aligned} \tag{6}$$

where

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

$$W_N = \frac{1}{2} M_{ij} N^c_i N^c_j + \tilde{h}_{ij} N^c_i (H_u L_j) + h_{i\alpha} N^c_i (H^u_\alpha L_4)\,. \tag{8}$$

In Equations (6)–(8) $\alpha, \beta = 1, 2$ and $i, j = 1, 2, 3$ as before. In the superpotential of the SE$_6$SSM the $U(1)_N$ symmetry forbids the term $\mu H_d H_u$. However all other terms, which are present in the MSSM superpotential, are allowed. In Equation (6) the sum of these terms is denoted as $W_{\mathrm{MSSM}}(\mu = 0)$. The sector responsible for the breakdown of the $SU(2)_W \times U(1)_Y \times U(1)_N$ symmetry involves the scalar components of ϕ, S, $\overline{S}$, H_u and H_d. If the superfields S and $\overline{S}$ develop VEVs along the D-flat direction, i.e., $\langle S \rangle \simeq \langle \overline{S} \rangle \simeq S_0$, then the value of S_0 can be much larger than the sparticle mass scale M_S resulting in an extremely heavy Z' boson. All extra exotic states may be also very heavy in this case. The neutral components of H_u and H_d have to gain non-zero VEVs, i.e., $\langle H_d \rangle = v_1/\sqrt{2}$ and $\langle H_u \rangle = v_2/\sqrt{2}$, so that $v = \sqrt{v_1^2 + v_2^2} \simeq 246\,\mathrm{GeV}$. These VEVs generate the masses of all SM particles triggering the breakdown of the $SU(2)_W \times U(1)_Y$ symmetry down to $U(1)_{em}$ associated with electromagnetism. Since we further focus on the scenarios with most sparticles at the multi-TeV scale a substantial degree of tuning is needed to stabilize the EW scale.

For the analysis of the phenomenological implications of the SE$_6$SSM it is worth introducing the Z_2^E symmetry, which is defined such that $\tilde{Z}_2^H = Z_2^M \times Z_2^E$ [38]. The supermultiplets $\overline{D}_i$, D_i, H^d_α, H^u_α, S_i, ϕ_i L_4 and $\overline{L}_4$ are odd under the Z_2^E symmetry. The components of all other supermultiplets are Z_2^E even. Because the Lagrangian of the SE$_6$SSM is invariant under both $\tilde{Z}_2^H$ and Z_2^M symmetries, the Z_2^E symmetry and R-parity are also conserved. This means that the exotic states, which are odd under the Z_2^E symmetry, can only be created in pairs in collider experiments and the lightest exotic particle as well as the lightest R-parity odd state have to be stable and may contribute to the density of dark matter. Here we focus on the scenarios in which the gravitino is the lightest R-parity odd state. Recently the cosmological implications of a gravitino with mass $m_{3/2} \sim$ KeV were discussed [84]. It is also assumed that the lightest stable exotic state is predominantly formed by the fermion components of H^d_α and H^u_α.

In order to find a viable scenario with stable gravitinos one needs to ensure that the lightest unstable R-parity odd (or exotic) state Y decays before BBN, i.e., its lifetime $\tau_Y \lesssim 1\,\mathrm{s}$. Otherwise the decay products of Y can alter the abundances of light elements which are induced by the BBN. The lifetime of the particle Y decaying into gravitino and its SM partner (or the lightest Z_2^E odd state) can be estimated as [85]

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

where m_Y is its mass. For $m_Y \simeq 1\,\mathrm{TeV}$ one can get $\tau_Y \lesssim 1\,\mathrm{sec}$ if $m_{3/2} \lesssim 1\,\mathrm{GeV}$. When gravitinos originate from scattering of particles in the thermal bath their contribution to the dark matter density is proportional to the reheating temperature T_R [86,87]

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

In Equation (10) $M_{\tilde{g}}$ is a gluino mass. Since $\Omega_{3/2}h^2 \leq 0.12$ [88], for $m_{3/2} \simeq 1\,\text{GeV}$ and $M_{\tilde{g}} \gtrsim 3\,\text{TeV}$ one finds an upper bound $T_R \lesssim 10^{6-7}\,\text{GeV}$ [89].

3. Generation of Lepton and Baryon Asymmetries

Even for so low reheating temperatures the appropriate amount of the lepton asymmetry can be induced within the SE$_6$SSM via the out-of-equilibrium decays of the lightest right-handed neutrino/sneutrino. Due to $(B+L)$-violating sphaleron interactions the generated lepton asymmetry is converted into the baryon asymmetry.

In the SM the process of the generation of lepton asymmetry is controlled by the three flavor CP (decay) asymmetries ε_{1,ℓ_k} which are associated with three lepton flavors. These decay asymmetries appear on the right-hand side of the Boltzmann equations. They are defined as

$$\varepsilon_{1,\ell_k} = \frac{\Gamma_{N_1\ell_k} - \Gamma_{N_1\bar{\ell}_k}}{\sum_m \left(\Gamma_{N_1\ell_m} + \Gamma_{N_1\bar{\ell}_m}\right)}\,. \tag{11}$$

Here $\Gamma_{N_1\ell_k}$ and $\Gamma_{N_1\bar{\ell}_k}$ are partial widths of the lightest right-handed neutrino decays $N_1 \to L_k + H_u$ and $N_1 \to \overline{L}_k + H_u^*$ with $k,m = 1,2,3$. At the tree level $\Gamma_{N_1\ell_k} = \Gamma_{N_1\bar{\ell}_k}$ and CP asymmetries (11) vanish. The non-zero contributions to the decay asymmetries come from the interference between the tree-level amplitudes of the decays of N_1 and one-loop corrections to them if CP invariance is violated in the lepton sector.

In the MSSM the decays of the lightest right-handed neutrino into Higgsino $\widetilde{H}_u$ and sleptons $\widetilde{L}_k$ also contribute to the lepton asymmetry generation. The corresponding flavor decay asymmetries are given by

$$\varepsilon_{1,\tilde{\ell}_k} = \frac{\Gamma_{N_1\tilde{\ell}_k} - \Gamma_{N_1\tilde{\ell}^*_k}}{\sum_m \left(\Gamma_{N_1\tilde{\ell}_m} + \Gamma_{N_1\tilde{\ell}^*_m}\right)}\,. \tag{12}$$

Moreover, supersymmetry predicts the existence of the lightest right-handed sneutrino $\widetilde{N}_1$ which is a scalar partner of N_1. The decays of $\widetilde{N}_1$ into slepton and Higgs as well as into lepton and Higgsino provide another possible origin of lepton asymmetry. The corresponding decay asymmetries can be determined similarly to the neutrino ones

$$\varepsilon_{\tilde{1},\ell_k} = \frac{\Gamma_{\widetilde{N}^*_1\ell_k} - \Gamma_{\widetilde{N}_1\bar{\ell}_k}}{\sum_m \left(\Gamma_{\widetilde{N}^*_1\ell_m} + \Gamma_{\widetilde{N}_1\bar{\ell}_m}\right)}\,, \qquad \varepsilon_{\tilde{1},\tilde{\ell}_k} = \frac{\Gamma_{\widetilde{N}_1\tilde{\ell}_k} - \Gamma_{\widetilde{N}^*_1\tilde{\ell}^*_k}}{\sum_m \left(\Gamma_{\widetilde{N}_1\tilde{\ell}_m} + \Gamma_{\widetilde{N}^*_1\tilde{\ell}^*_m}\right)}\,. \tag{13}$$

When the sparticle mass scale M_S is considerably smaller than M_1

$$\varepsilon_{1,\ell_k} = \varepsilon_{1,\tilde{\ell}_k} = \varepsilon_{\tilde{1},\ell_k} = \varepsilon_{\tilde{1},\tilde{\ell}_k}\,. \tag{14}$$

Assuming the type I seesaw models of neutrino mass generation the decay asymmetries mentioned above were initially computed within the SM [90–93] and MSSM [94–96]. Flavor effects were ignored in the early studies of leptogenesis (see for example [97]). The importance of these effects was emphasised in [98–104].

The non-minimal SUSY models (like the SE$_6$SSM) may include additional $SU(2)_W$ doublets with quantum numbers of Higgs fields (H^d_α and H^u_α) and extra lepton multiplets (L_4 and $\overline{L}_4$) at low energies. It is convenient to denote all Higgs-like multiplets and $SU(2)_W$ lepton doublets that interact with the right-handed neutrino superfields as H^u_k and L_x respectively. In the case of the SE$_6$SSM $H^u_3 \equiv H_u$, $k = 1,2,3$ and $x = 1,2,3,4$. If the components of additional Higgs-like and lepton supermultiplets are lighter than N_1 and $\widetilde{N}_1$ they can give rise to new decay modes of the lightest right-handed neutrino and its superpartner. Each new channel of the decays of N_1 and $\widetilde{N}_1$ should lead to extra CP asymmetry that contributes to the lepton asymmetry generation. In this case the definitions of the decay asymmetries (11)–(13) need to be generalised. In particular the definitions (11) and (12) can be modified in the following way [36]

$$\varepsilon^k_{1,f} = \frac{\Gamma^k_{N_1 f} - \Gamma^k_{N_1 \bar{f}}}{\sum_{m,f'}\left(\Gamma^m_{N_1 f'} + \Gamma^m_{N_1 \bar{f}'}\right)}, \tag{15}$$

where f and f' may be either ℓ_x or $\widetilde{\ell}_x$ while $\bar{f}$ and $\bar{f}'$ should be associated with either $\bar{\ell}_x$ or $\widetilde{\ell}^*_x$. The superscripts k and m correspond to the components of the supermultiplets H^u_k and H^u_m in the final state. The denominator of Equation (15) contains a sum of partial widths of the decays of N_1. For ε^k_{1,ℓ_x} this sum involves all possible partial decay widths of the lightest right-handed neutrino whose final state includes fermion components of the supermultiplets L_x. The expressions for $\varepsilon^k_{1,\widetilde{\ell}_x}$ involve in the denominator a sum of partial widths of the decays of the lightest right-handed neutrino over all possible decay modes which have scalar components of the supermultiplets L_x in the final state.

The CP asymmetries associated with the decays of the lightest right-handed sneutrino $\varepsilon^k_{\widetilde{1},f}$ can be defined similarly to $\varepsilon^k_{1,f}$. In order to obtain the appropriate expressions for $\varepsilon^k_{\widetilde{1},f}$ the field of the lightest right-handed neutrino in Equation (15) ought to be replaced by either $\widetilde{N}_1$ or $\widetilde{N}^*_1$. In the limit, when the sparticle mass scale M_S is negligibly small as compared with M_1, all soft SUSY breaking terms can be safely ignored and the relation between different decay asymmetries (14) remains intact, i.e., $\varepsilon^k_{1,f} = \varepsilon^k_{\widetilde{1},f}$.

Within the SE$_6$SSM ε^3_{1,ℓ_n} ($\varepsilon^3_{1,\widetilde{\ell}_n}$) with $n = 1,2,3$ are flavor decay asymmetries associated with the decays of N_1 into Higgs doublet H_u and ordinary leptons (Higgsino $\widetilde{H}_u$ and sleptons), whereas $\varepsilon^3_{\widetilde{1},\ell_n}$ ($\varepsilon^3_{\widetilde{1},\widetilde{\ell}_n}$) are CP asymmetries corresponding to the decays of $\widetilde{N}_1$ into leptons and Higgsino $\widetilde{H}_u$ (sleptons and Higgs doublet H_u). Additional decay asymmetries $\varepsilon^\alpha_{1,\ell_4}$, $\varepsilon^\alpha_{1,\widetilde{\ell}_4}$, $\varepsilon^\alpha_{\widetilde{1},\ell_4}$ and $\varepsilon^\alpha_{\widetilde{1},\widetilde{\ell}_4}$ in this SUSY model arise due to the extra decay channels of N_1 and $\widetilde{N}_1$

$$N_1 \to L_4 + H^u_\alpha, \qquad N_1 \to \widetilde{L}_4 + \widetilde{H}^u_\alpha, \qquad \widetilde{N}^*_1 \to L_4 + \widetilde{H}^u_\alpha, \qquad \widetilde{N}_1 \to \widetilde{L}_4 + H^u_\alpha, \tag{16}$$

The structure of the part of the SE$_6$SSM superpotential W_N (8), that describes the interactions of the right-handed neutrino superfields with other supermultiplets, indicates that all other $\varepsilon^k_{1,f}$ and $\varepsilon^k_{\widetilde{1},f}$ vanish.

After inflation the lightest right-handed neutrino/sneutrino with mass M_1 may be produced by thermal scattering if $T_R > M_1$. Since in the scenarios under consideration $T_R \lesssim 10^{6-7}$ GeV here we require that $M_1 \lesssim 10^6$ GeV to guarantee that thermal leptogenesis can take place. It is also assumed that two other right-handed neutrino/sneutrino states have masses $M_{2,3} \lesssim 10^6$ GeV so that $M_1 \lesssim M_2 \lesssim M_3$ while the sparticle mass scale M_S is lower than 10 TeV. In order to reproduce the left-handed neutrino mass scale $m_\nu \lesssim 0.1$ eV the absolute values of the couplings $|\tilde{h}_{ij}|$ in Equation (8) should be rather small for such low M_i, i.e., $|\tilde{h}_{ij}|^2 \ll 10^{-8}$. Such couplings $\tilde{h}_{ij}$ induce quite small decay asymmetries $\varepsilon^3_{1,f}$ and $\varepsilon^3_{\widetilde{1},f}$ so that $\tilde{h}_{ij}$ can be ignored.

Nevertheless the new decay modes of N_1 and $\widetilde{N}_1$ (16) may give rise to the sufficiently large CP asymmetries $\varepsilon^\alpha_{1,\ell_4}$, $\varepsilon^\alpha_{1,\widetilde{\ell}_4}$, $\varepsilon^\alpha_{\widetilde{1},\ell_4}$ and $\varepsilon^\alpha_{\widetilde{1},\widetilde{\ell}_4}$ that control the process of lepton asymmetry generation. At the tree level the partial widths corresponding to the new decay channels (16) are given by

$$\Gamma^\alpha_{N_1\ell_4} + \Gamma^\alpha_{N_1\bar{\ell}_4} = \Gamma^\alpha_{N_1\widetilde{\ell}_4} + \Gamma^\alpha_{N_1\widetilde{\ell}^*_4} = \Gamma^\alpha_{\widetilde{N}^*_1\ell_4} = \Gamma^\alpha_{\widetilde{N}_1\bar{\ell}_4} = \Gamma^\alpha_{\widetilde{N}_1\widetilde{\ell}_4} = \Gamma^\alpha_{\widetilde{N}^*_1\widetilde{\ell}^*_4} = \frac{|h_{1\alpha}|^2}{8\pi}M_1 \tag{17}$$

and CP asymmetries $\varepsilon^{\alpha}_{1,\ell_4}$, $\varepsilon^{\alpha}_{1,\widetilde{\ell}_4}$, $\varepsilon^{\alpha}_{\widetilde{1},\ell_4}$ as well as $\varepsilon^{\alpha}_{\widetilde{1},\widetilde{\ell}_4}$ vanish. As in the SM and MSSM the non-zero values of the decay asymmetries in the SE$_6$SSM arise after the inclusion of one-loop self-energy and vertex corrections to the decay amplitudes of the lightest right-handed neutrino/sneutrino. Neglecting the Yukawa couplings $\tilde{h}_{ik}$ one finds [36]

$$\begin{gathered}
\varepsilon^{\alpha}_{1,\ell_4} = \varepsilon^{\alpha}_{1,\widetilde{\ell}_4} = \varepsilon^{\alpha}_{\widetilde{1},\ell_4} = \varepsilon^{\alpha}_{\widetilde{1},\widetilde{\ell}_4} = \frac{1}{8\pi}\frac{\sum_{j=2,3}\mathrm{Im}\left[h^{*}_{1\alpha}B_j h_{j\alpha}\right]}{\sum_{\beta}|h_{1\beta}|^2}\,,\\
B_j = \sum_{\beta}\left\{h^{*}_{1\beta}h_{j\beta}\,f\left(\frac{M_j^2}{M_1^2}\right) + \frac{M_1}{M_j}h_{1\beta}h^{*}_{j\beta}f^{S}\left(\frac{M_j^2}{M_1^2}\right)\right\},\\
f(z) = f^{V}(z) + f^{S}(z)\,,\qquad f^{S}(z) = \frac{2\sqrt{z}}{1-z}\,,\qquad f^{V}(z) = -\sqrt{z}\,\ln\left(\frac{1+z}{z}\right).
\end{gathered}\tag{18}$$

The analytical expressions for the CP asymmetries (18) are simplified dramatically if $|h_{12}|$ goes to zero. In this case $\varepsilon^{2}_{1,\ell_4} = \varepsilon^{2}_{1,\widetilde{\ell}_4} = \varepsilon^{2}_{\widetilde{1},\ell_4} = \varepsilon^{2}_{\widetilde{1},\widetilde{\ell}_4} = 0$. If M_j are real and $h_{j1} = |h_{j1}|e^{i\varphi_{j1}}$ the expressions for other CP asymmetries reduce to [36,37]

$$\varepsilon^{1}_{1,\ell_4} = \varepsilon^{1}_{1,\widetilde{\ell}_4} = \varepsilon^{1}_{\widetilde{1},\ell_4} = \varepsilon^{1}_{\widetilde{1},\widetilde{\ell}_4} = \varepsilon = \frac{1}{8\pi}\left[\sum_{j=2,3}|h_{j1}|^2 f\left(\frac{M_j^2}{M_1^2}\right)\sin 2\Delta\varphi_{j1}\right],\tag{19}$$

where $\Delta\varphi_{j1} = \varphi_{j1} - \varphi_{11}$. From the part of the SE$_6$SSM superpotential (8) one can see the supermultiplets H^u_α may be redefined so that only H^u_1 interacts with N^c_1 and L_4. Thus $\tilde{h}_{12}$ in W_N may be set to zero without loss of generality. It is also worth noting that the scalar and fermion components of the supermultiplet L_4 being produced in the decays of N_1 and $\widetilde{N}_1$ sequentially decay into ordinary leptons inducing lepton number asymmetries.

The evolution of the $U(1)_{B-L}$ number densities is described by the Boltzmann equations. In the scenarios under consideration the results for the baryon and lepton asymmetries obtained within the SM and MSSM can be easily generalised. In particular, the generated total baryon asymmetry may be estimated using an approximate formula [105]:

$$Y_{\Delta B} \sim 10^{-3}\varepsilon\eta\,,\qquad Y_{\Delta B} = \left.\frac{n_B - n_{\bar{B}}}{s}\right|_0 = (8.75 \pm 0.23)\times 10^{-11}\,,\tag{20}$$

where $Y_{\Delta B}$ is the baryon asymmetry relative to the entropy density. In Equation (20) η is an efficiency factor that varies from 0 to 1. The efficiency factor in the so-called strong washout scenario can be estimated as follows

$$\begin{gathered}
\eta \simeq H(T = M_1)/\Gamma_1\,,\\
\Gamma_1 = \Gamma^{1}_{N_1\ell_4} + \Gamma^{1}_{N_1\widetilde{\ell}_4} = \frac{|h_{11}|^2}{8\pi}M_1\,,\qquad H = 1.66 g_*^{1/2}\frac{T^2}{M_P}\,.
\end{gathered}\tag{21}$$

Here $M_P = 1.22\cdot 10^{19}$ GeV, H is the Hubble expansion rate and $g_* = n_b + \frac{7}{8}n_f$ is the number of relativistic degrees of freedom in the thermal bath. In the SE$_6$SSM $g_* \approx 360$.

In order to simplify our numerical analysis we set $|h_{12}| = |h_{31}| = 0$. Our results are summarised in Figure 1. The decay asymmetries (19) are determined by $|h_{21}|$, (M_2/M_1) and combination of the CP-violating phases $\Delta\varphi_{21}$, but do not depend on $|h_{11}|$. Here we fix $\Delta\varphi_{21}$ so that these asymmetries attain their maximum absolute values, i.e., $|\sin 2\Delta\varphi_{21}| = 1$. In this case $|\varepsilon|$ changes from $2.4\cdot 10^{-5}$ to $2.4\cdot 10^{-3}$ if $|h_{21}|$ increases from 0.01 to 0.1 for $M_2 = 1.2\cdot M_1$. As follows from Equation (21) the efficiency factor η is set by $|h_{11}|$ and M_1. We restrict our consideration here by the values of $|h_{11}|^2 \gg |\tilde{h}_{ik}|^2$, i.e., $|h_{11}| \gtrsim 10^{-4}$. For $M_1 = 100$ TeV the efficiency factor varies from $6.5\cdot 10^{-4}$ to $6.5\cdot 10^{-6}$ when $|h_{11}|$ increases from 10^{-4} to 10^{-3}. Thus for $|h_{21}| \simeq 0.1$, $M_1 = 100$ TeV and $M_2 = 1.2\cdot M_1$ the observed baryon asymmetry can be reproduced even if $|h_{11}| \gtrsim 10^{-4}$.

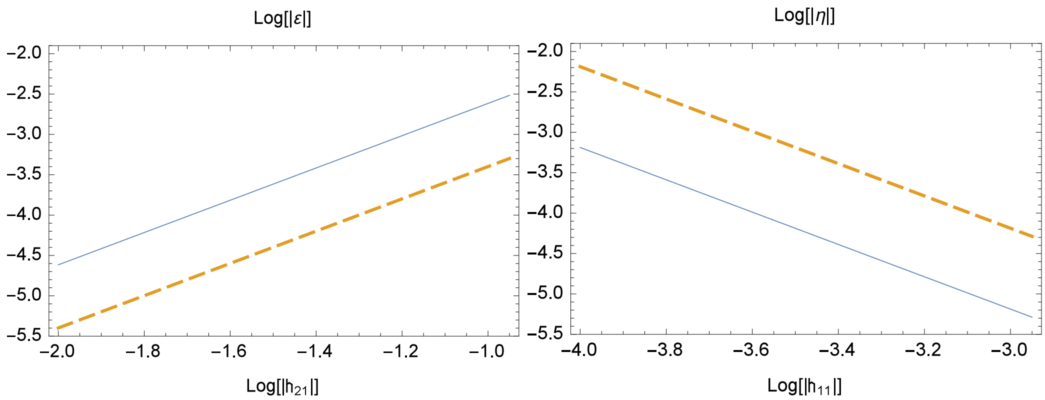

Figure 1. (**Left**) Logarithm (base 10) of the absolute value of the CP asymmetry $|\varepsilon|$ as a function of logarithm (base 10) of $|h_{21}|$ for $M_2 = 1.2 \cdot M_1$ (solid line) and $M_2 = 10 \cdot M_1$ (dashed line). (**Right**) Logarithm (base 10) of the absolute value of the efficiency factor $|\eta|$ as a function of logarithm (base 10) of $|h_{11}|$ for $M_1 = 100\,\text{TeV}$ (solid line) and $M_1 = 1000\,\text{TeV}$ (dashed line). Here we fix $|h_{12}| = |h_{31}| = 0$ and $\Delta\varphi_{21} = \pi/4$.

4. Dark Matter-Nucleon Scattering Cross Section

The scalar components of the supermultiplets ϕ_i, S_i, H^u_α and H^d_α do not acquire VEVs. Their fermion components form the exotic (inert) neutralino and chargino states. The signatures associated with the inert neutralino states were examined in [106,107]. When the components of ϕ_i are significantly heavier than the fermions and bosons from S_i, H^u_α and H^d_α, they can be integrated out so that W_{IH} reduces to

$$W_{IH} \to \widetilde{W}_{IH} \simeq -\widetilde{\mu}_{ij} S_i S_j + \lambda_{\alpha\beta} S (H^d_\alpha H^u_\beta) + \tilde{f}_{i\alpha} S_i (H^d_\alpha H_u) + f_{i\alpha} S_i (H_d H^u_\alpha) + \ldots . \qquad (22)$$

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

In this article we explore the scenarios in which the fermion components of the H^u_1 and H^d_1 compose the lightest exotic state with $Z^E_2 = -1$ while all other exotic states and all sparticles except gravitino have masses which are considerably larger than 1 TeV. We also assume that H^d_1 and H^u_1 mostly interact with S_1, H_u and H_d, whereas all other couplings of the supermultiplets H^u_1 and H^d_1 are very small. In this approximation the mass matrix, that determines the lightest exotic neutralino masses, takes the form [49]

$$M^{ab} = -\begin{pmatrix} 0 & \mu_{11} & \dfrac{\tilde{f}_{11}}{\sqrt{2}} v_2 \\ \mu_{11} & 0 & \dfrac{f_{11}}{\sqrt{2}} v_1 \\ \dfrac{\tilde{f}_{11}}{\sqrt{2}} v_2 & \dfrac{f_{11}}{\sqrt{2}} v_1 & \widetilde{\mu}_1 \end{pmatrix}, \qquad (23)$$

where $\mu_{11} \simeq \lambda_{11}\langle S\rangle$. Instead of the VEVs of H_d and H_u, i.e., v_1 and v_2, it is more convenient to introduce $\tan\beta = v_2/v_1$ and $v = \sqrt{v_1^2 + v_2^2} \approx 246\,\text{GeV}$. The charged fermion components of the supermultiplets H^u_1 and H^d_1 form the lightest exotic chargino. Its mass is determined by μ_{11}, i.e., $m_{\chi^\pm_1} = |\mu_{11}|$.

If $|\widetilde{\mu}_1|$ is considerably larger than $|\mu_{11}|$ and v the mass matrix (23) can be diagonalised. Using the perturbation theory method (see, for example, [108–111]), one finds [49]

$$\begin{gathered} m_{\chi_1} \simeq m_{\chi^\pm_1} - \Delta_1\,, \qquad m_{\chi_2} \simeq m_{\chi^\pm_1} + \Delta_2\,, \qquad m_{\chi_2} \simeq \widetilde{\mu}_1 + \Delta_1 + \Delta_2\,, \\ \Delta_1 \simeq \frac{(\tilde{f}_{11} v \sin\beta + f_{11} v \cos\beta)^2}{4(\widetilde{\mu}_1 - m_{\chi^\pm_1})}\,, \qquad \Delta_2 \simeq \frac{(\tilde{f}_{11} v \sin\beta - f_{11} v \cos\beta)^2}{4(\widetilde{\mu}_1 + m_{\chi^\pm_1})}\,. \end{gathered} \qquad (24)$$

As follows from Equation (24) the lightest exotic neutralino masses (m_{χ_2} and m_{χ_1}) are also set by μ_{11} in the leading approximation. We restrict our consideration here to the part of the parameter space of the SE$_6$SSM where $m_{\chi_2} - m_{\chi_1} > 200\,\mathrm{MeV}$. As a consequence the inelastic scattering processes $\chi_1 N \to \chi_2 N$, where N, χ_1 and χ_2 denote a nucleon as well as the lightest and second lightest exotic neutralino states, do not take place. In this part of the parameter space χ_2 decays before BBN, i.e., its lifetime is shorter than 1 s .

Since in the scenarios under consideration the lightest neutral exotic neutralino χ_1 is stable its contribution to the cold dark matter density may be estimated using formula

$$\Omega_{\tilde{H}} h^2 \simeq 0.1 \left(\frac{\mu_{11}}{1\,\mathrm{TeV}} \right)^2 . \tag{25}$$

The approximate formula (25) was derived within the MSSM [112,113]. On the other hand the Planck observations lead to $(\Omega h^2)_{\rm exp} = 0.1188 \pm 0.0010$ [88]. Therefore in the phenomenologically viable scenarios μ_{11} should be lower than 1.1 TeV. When $\mu_{11} < 1.1\,\mathrm{TeV}$, the gravitino can account for some or major part of the cold dark matter density.

In the SE$_6$SSM the interactions of the cold dark matter with the SM particles are determined by the couplings of χ_1 because the corresponding gravitino couplings are negligibly small. The low-energy effective Lagrangian, that describes the interactions of the lightest exotic neutralino with quarks can be written as

$$\mathcal{L}_{\chi_1 q} = \sum_q \Big(a_q \bar{\chi}_1 \chi_1 \bar{q} q + d_q \bar{\chi}_1 \gamma^\mu \gamma_5 \chi_1 \bar{q} \gamma_\mu \gamma_5 q \Big) . \tag{26}$$

The first term in the brackets results in a spin-independent interaction whereas the second one gives rise to a spin-dependent interaction.

In the scenarios under consideration the dominant contribution to the parameters d_q in the Lagrangian (26) stems from t-channel Z boson exchange. Taking into account that in the field basis $(\tilde{H}_1^{d0},\, \tilde{H}_1^{u0},\, \tilde{S}_1)$

$$\chi_\alpha = N_\alpha^1 \tilde{H}_1^{d0} + N_\alpha^2 \tilde{H}_1^{u0} + N_\alpha^3 \tilde{S}_1 \,, \qquad\qquad \alpha = 1, 2 \,, \tag{27}$$

where N_i^a is the exotic neutralino mixing matrix defined by

$$N_i^a M^{ab} N_j^b = m_i \delta_{ij}, \qquad\qquad \text{no sum on } i \,, \tag{28}$$

the part of the Lagrangian, which describes the interactions of the lightest and second-lightest exotic neutralino states with Z, may be presented in the following form:

$$\mathcal{L}_{Z\chi\chi} = \sum_{\alpha,\beta} \frac{M_Z}{2v} Z_\mu \left(\chi_\alpha^T \gamma_\mu \gamma_5 \chi_\beta \right) R_{Z\alpha\beta} \,, \qquad\qquad R_{Z\alpha\beta} = N_\alpha^1 N_\beta^1 - N_\alpha^2 N_\beta^2 \,. \tag{29}$$

In Equation (28) M^{ab} is 3×3 mass matrix (23). Then the parameters d_q as well as the corresponding χ_1-proton and χ_1-neutron scattering cross sections (σ^p and σ^n) are given by

$$\sigma^{p,n} = \frac{12 m_r^2}{\pi} \left(\sum_{q=u,d,s} d_q \Delta_q^{p,n} \right)^2 , \qquad d_q = \frac{T_{3q}}{2v^2} R_{Z11} \,, \qquad m_r = \frac{m_{\chi_1} m_N}{m_{\chi_1} + m_N} \,. \tag{30}$$

Here m_N is a nucleon mass and T_{3q} is the third component of isospin. We set $\Delta_u^p = \Delta_d^n = 0.842$, $\Delta_d^p = \Delta_u^n = -0.427$ and $\Delta_s^p = \Delta_s^n = -0.085$ [113].

In the SE$_6$SSM χ_1 does not couple to squarks and quarks. As a consequence the only contributions that parameters a_q receive come from the t-channel exchange of Higgs scalars. Since in the scenarios under consideration all Higgs bosons except the lightest Higgs scalar h_1 are expected to be considerably heavier than 1 TeV, all contributions caused by the heavy

Higgs exchange can be neglected. The lightest Higgs boson with mass $m_{h_1} \approx 125\,\text{GeV}$ manifests itself in the interactions with the SM states as a SM-like Higgs in this case so that

$$\frac{a_q}{m_q} \simeq \frac{g_{h\chi\chi}}{v m_{h_1}^2}\,, \qquad\qquad g_{h\chi\chi} = -\frac{1}{\sqrt{2}}\left(f_{11} N_1^3 N_1^2 \cos\beta + \tilde{f}_{11} N_1^3 N_1^1 \sin\beta\right), \tag{31}$$

where m_q is a quark mass and $g_{h\chi\chi}$ is the coupling of the lightest exotic neutralino to h_1. The spin-independent part of χ_1-nucleon cross section takes the form [114,115]

$$\begin{aligned} \sigma_{SI} &= \frac{4 m_r^2 m_N^2}{\pi v^2 m_{h_1}^4} |g_{h\chi\chi} F^N|^2\,, & F^N &= \textstyle\sum_{q=u,d,s} f_{Tq}^N + \frac{2}{27}\sum_{Q=c,b,t} f_{TQ}^N\,, \\ m_N f_{Tq}^N &= \langle N | m_q \bar{q} q | N \rangle\,, & f_{TQ}^N &= 1 - \textstyle\sum_{q=u,d,s} f_{Tq}^N\,. \end{aligned} \tag{32}$$

The value of σ_{SI} depends quite strongly on f_{Tq}^N, i.e., hadronic matrix elements. We fix $f_{Ts}^N \simeq 0.0447$, $f_{Td}^N \simeq 0.0191$ and $f_{Tu}^N \simeq 0.0153$ which are the default values used in micrOMEGAs [116] (see also [117–120]). Using the perturbation theory method it is straightforward to obtain the approximate expressions for $g_{h\chi\chi}$ and R_{Z11}. If $\tilde{\mu}_1 \gg \mu_{11} > 0$ and μ_{11} is substantially larger than $f_{11} v \cos\beta$ and $\tilde{f}_{11} v \sin\beta$, one finds

$$|g_{h\chi\chi}| \simeq \frac{\Delta_1}{v}\,, \qquad\qquad R_{Z11} \simeq \frac{v^2 (f_{11}^2 \cos^2\beta - \tilde{f}_{11}^2 \sin^2\beta)}{4\mu_{11}(\tilde{\mu}_1 - \mu_{11})}\,. \tag{33}$$

In our analysis we restrict our considerations to moderate values of $\tan\beta$, i.e., $\tan\beta \simeq 2$. For $\tan\beta \leq 4$ one can get $m_{h_1} \approx 125\,\text{GeV}$ in the SE_6SSM only if $\lambda \gtrsim \sqrt{2}(M_Z/v) \simeq 0.5$. When coupling λ is so large all Higgs states except the SM-like Higgs boson have masses beyond the multi-TeV range [32,33,72,73]. Therefore they cannot be observed at the LHC. The realisation of such scenarios requires significant fine-tuning, $\sim$0.01%, of the parameters of the model under consideration [121]. LHC experiments ruled out the $U(1)_N$ gauge boson with masses $M_{Z'}$ below 4.5 TeV [122,123]. If $\langle S \rangle \simeq \langle \overline{S} \rangle$ the mass of the Z' boson in the SE_6SSM is given by

$$M_{Z'} \approx 2 g_1' Q_S \langle S \rangle\,, \tag{34}$$

where g_1' and Q_S are the $U(1)_N$ gauge coupling and the $U(1)_N$ charge of the superfield S. The low-energy value of g_1' can be calculated assuming the unification of gauge couplings [32]. Then $M_{Z'} \gtrsim 4.5\,\text{TeV}$ can be obtained when $\langle S \rangle \simeq \langle \overline{S} \rangle \gtrsim 6\,\text{TeV}$. To avoid the lower experimental bound on the lightest exotic chargino mass we assume that $\mu_{11} \gtrsim 200\,\text{GeV}$. To ensure that χ_1 leads to the phenomenologically acceptable density of the cold dark matter the interval of variations of μ_{11} is limited from above by 1 TeV so that $\lambda_{11} \lesssim 0.17$. In addition, the validity of perturbation theory up to the GUT scale is required that constrains the range of variations of f_{11} and $\tilde{f}_{11}$ at low energies. We also set $\tilde{\mu}_1 \simeq 2\,\text{TeV}$.

The results of our numerical analysis are presented in Figures 2 and 3. From these Figures it follows that the approximate expressions (33) describe quite well the dependence of the couplings $|R_{Z11}|$ and $g_{h\chi\chi}$ on the SE_6SSM parameters if $|\tilde{\mu}_1| \gg |\mu_{11}|$. In particular, $|R_{Z11}|$ and $\sigma^{p,n}$ diminish whereas $g_{h\chi\chi}$ and σ_{SI} grow with increasing μ_{11} from 200 GeV to 1 TeV. The approximate expressions (33) indicate that the couplings $g_{h\chi\chi}$ and R_{Z11} go to zero when $f_{11} \approx -\tilde{f}_{11}\tan\beta$. This means that in the corresponding part of the parameter space the interactions of χ_1 with the baryons tend to extremely weak. The vanishing of $g_{h\chi\chi}$ and R_{Z11} can be attained only if the parameters f_{11}, $\tilde{f}_{11}$ and $\tan\beta$ are fine-tuned. However as follows from Figures 2 and 3 in order to achieve the desirable suppression of the σ_{SI} and $\sigma^{p,n}$ the fine-tuning is not needed. In our analysis f_{11} is chosen to be positive while $\tilde{f}_{11}$ is fixed to be negative. As a consequence $g_{h\chi\chi}$, $|R_{Z11}|$, the spin-independent and spin-dependent cross sections decrease when f_{11} grows and approaches $-\tilde{f}_{11}\tan\beta$. It is worth pointing out that in the part of the parameter space of the SE_6SSM where the couplings $|g_{h\chi\chi}|$ and $|R_{Z11}|$ become considerably smaller than 10^{-3} one cannot neglect the contributions to σ_{SI} and $\sigma^{p,n}$ which are induced by the heavy Higgs states and Z' boson. Moreover one should take into account the quantum corrections to the Lagrangian (26) that stem from the one-loop diagrams involving the electroweak gauge bosons [124,125]. The inclusion of these quantum corrections lead to $\sigma_{SI} \sim 0.1$ yb even if at the tree level $|g_{h\chi\chi}| \ll 10^{-3}$ [126].

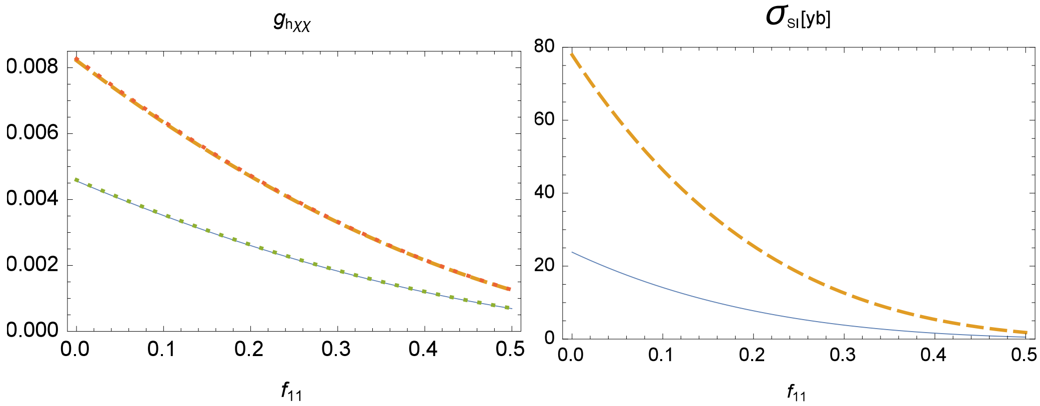

Figure 2. (**Left**) The coupling $g_{h\chi\chi}$ and (**Right**) the cross-section σ_{SI} as a function of f_{11} for $\tilde{f}_{11} = -0.41$, $\tan\beta = 2$, $\tilde{\mu}_1 = 2$ TeV, $\mu_{11} = 200$ GeV (solid lines) and $\mu_{11} = 1$ TeV (dashed lines). The dotted lines correspond to the approximate expression for $g_{h\chi\chi}$ (33).

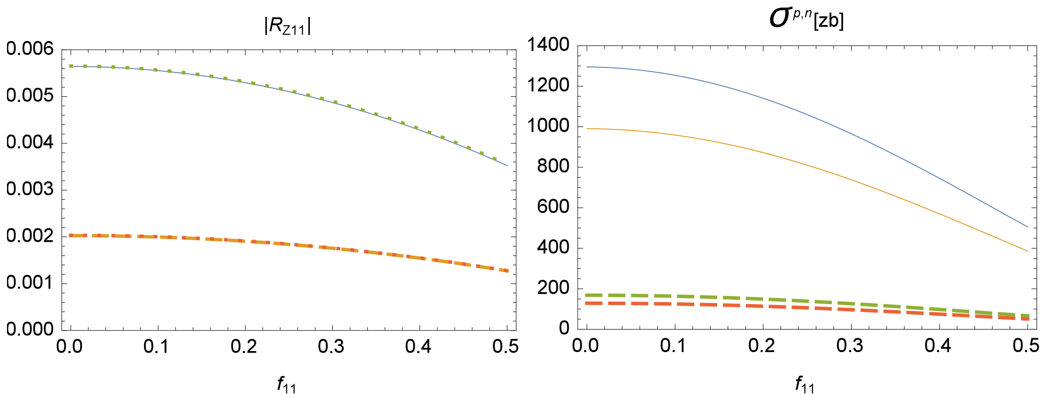

Figure 3. (**Left**) The coupling $|R_{Z11}|$ and (**Right**) the cross-section $\sigma^{p,n}$ as a function of f_{11} for $\tilde{f}_{11} = -0.41$, $\tan\beta = 2$, $\tilde{\mu}_1 = 2$ TeV, $\mu_{11} = 200$ GeV (solid lines) and $\mu_{11} = 1$ TeV (dashed lines). The upper solid and upper dashed lines represent σ^p while lower solid and lower dashed lines are associated with σ^n. The dotted lines correspond to the approximate expression for R_{Z11} (33).

Now let us compare the computed values of σ_{SI} and $\sigma^{p,n}$ with the corresponding experimental bounds $(\sigma_{SI})_{\text{exp}}$ and $(\sigma^{p,n})_{\text{exp}}$. The spin-independent χ_1-nucleon scattering cross sections presented in Figure 2 are always smaller than LUX-ZEPLIN (LZ) experimental limits, i.e., $(\sigma_{SI})_{\text{exp}} \approx 60\,\text{yb}(300\,\text{yb})$ for $m_{\chi_1} \approx 200\,\text{GeV}(1\,\text{TeV})$ [127]. The most stringent experimental bound on the spin-dependent WIMP-proton scattering cross section was obtained by the PICO-60 experiment, i.e., $(\sigma^p)_{\text{exp}} \approx 10^5\,\text{zb}\,(4\cdot 10^5\,\text{zb})$ for $m_{\chi_1} \approx 200\,\text{GeV}(1\,\text{TeV})$ [128]. The spin-dependent WIMP-neutron scattering cross section is more tightly constrained, i.e., $(\sigma^n)_{\text{exp}} \approx 0.9\cdot 10^4\,\text{zb}\,(5\cdot 10^4\,\text{zb})$ for $m_{\chi_1} \approx 200\,\text{GeV}(1\,\text{TeV})$ [127]. The values of $\sigma^{p,n}$ shown in Figure 3 are considerably smaller than $(\sigma^{p,n})_{\text{exp}}$.

In the SE$_6$SSM the maximal values of the spin-independent and spin-dependent χ_1-nucleon scattering cross sections are much larger than σ_{SI} and $\sigma^{p,n}$ presented in Figures 2 and 3. These cross sections attain their maximal possible values for $\mu_{11} \simeq \tilde{\mu}_1$ and $\tilde{f}_{11} \sim f_{11} \sim 1$. In this case σ_{SI} can reach $20 - 30$ zb which is considerably larger than the corresponding experimental limit[127]. Since for a given mass of the lightest exotic fermion σ_{SI} vanishes when $f_{11} = -\tilde{f}_{11}\tan\beta$, the spin-independent χ_1-nucleon scattering cross section varies from zero to its maximal value for each m_{χ_1}. In the scenarios under consideration the suppression of σ_{SI} and $\sigma^{p,n}$ is caused by the cancellations of different contributions to R_{Z11} and $g_{h\chi\chi}$ as well as by the large value of $\widetilde{\mu}_1$ that should be associated with the sparticle mass scale M_S. In the near future the experiments LZ [129], XENONnT [130], DARWIN [131] and DarkSide-20k [132] can set even more stringent limits on σ_{SI} and $\sigma^{p,n}$ constraining further the parameter space of the SE$_6$SSM.

5. Conclusions

In this article we considered leptogenesis and the interactions of dark matter with nucleons in the $U(1)_N$ extension of the MSSM, in which the single discrete $\tilde{Z}_2^H$ symmetry forbids flavor-changing transitions and the most dangerous baryon and lepton number violating operators. The low energy matter content of this SUSY model (SE$_6$SSM) includes three fundamental 27 representations of E_6, an additional 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, a pair of the SM singlet superfields S and $\overline{S}$ with opposite $U(1)_N$ charges as well as four E_6 singlet superfields. Thus the SE$_6$SSM contains extra exotic matter beyond the MSSM. The scalar components of the superfields S and $\overline{S}$ can develop VEVs along the D-flat direction, so that $\langle S \rangle \simeq \langle \bar{S} \rangle \gg 10\,\text{TeV}$, breaking the $U(1)_N$ symmetry and inducing TeV scale masses of all extra exotic fermions and Z' boson. The relatively light components of the supermultiplets L_4 and $\overline{L}_4$ allow the lightest exotic colored state to decay before BBN. They also facilitate the unification of gauge couplings.

In the SE$_6$SSM the cold dark matter density is formed by two stable neutral states. Here we focused on the scenarios in which one of these stable particles is gravitino. In this case all TeV scale states can decay before BBN only when gravitino mass $m_{3/2} \lesssim 1\,\text{GeV}$. Another stable state tends to be the lightest exotic neutralino χ_1 with mass $m_{\chi_1} \leq 1.1\,\text{TeV}$. Because it is a superposition of the neutral fermion components of the $SU(2)_W$ doublets, the lightest exotic chargino $\chi_1^{\pm}$, the second lightest exotic neutralino χ_2 and χ_1 are nearly degenerate around m_{χ_1}. Such scenarios result in the phenomenologically acceptable dark matter density if the reheating temperatures $T_R \lesssim 10^{6-7}\,\text{GeV}$. Even for so low T_R the decays of the lightest right-handed neutrino/sneutrino in the SE$_6$SSM can generate the appropriate lepton asymmetry due to the presence of L_4 and $\overline{L}_4$ in the particle spectrum. This lepton asymmetry is converted into the observed baryon asymmetry via sphaleron processes. In the scenarios under consideration there is a part of the SE$_6$SSM parameter space in which the dark matter–nucleon scattering cross section is substantially smaller than the present experimental limits.

The phenomenological viability of the scenarios under consideration requires $\chi_1^{\pm}$, χ_2 and χ_1 to be lighter than 1.1 TeV. Otherwise the annihilation cross section for $\chi_1 + \chi_1 \to$ SM particles becomes too small giving rise to the cold dark matter density which is considerably larger than its measured value. Relatively light charged and neutral fermions have been searched for in different experiments. If the mass of the lightest exotic chargino $m_{\chi_1^{\pm}}$ and the mass of the second lightest exotic neutralino m_{χ_2} are too close to m_{χ_1} the decay products of $\chi_1^{\pm}$ and χ_2 may escape detection. This happens, for example, within natural SUSY, where the mass splitting between the lightest and second lightest ordinary neutralino states as well as the mass splitting between the lightest ordinary chargino and the lightest ordinary neutralino are at least a few GeV [133–135]. The results of the searches for such degenerate states depend on $\Delta = m_{\chi_1^{\pm}} - m_{\chi_1}$ and $\Delta_1 = m_{\chi_2} - m_{\chi_1}$.

In the scenarios under consideration the SE$_6$SSM parameters are chosen so that $\Delta_1 > 200\,\text{MeV}$ while Δ is larger than 300 MeV [126,136]. Therefore $\chi_1^{\pm}$ and χ_2 cannot be long-lived. At the LHC the lightest exotic chargino and neutralino states can be produced in pairs via off-shell W and Z-bosons. Then $\chi_1^{\pm}$ and χ_2 subsequently decay into hadrons and χ_1. For $\Delta \simeq 4.7\,\text{GeV}(2\,\text{GeV})$ ATLAS ruled out $\chi_1^{\pm}$ with masses below $193\,\text{GeV}(140\,\text{GeV})$ [137]. For $\Delta = 1\,\text{GeV}$ CMS excluded $\chi_1^{\pm}$ with masses below 112 GeV [138]. The discovery prospects for such exotic chargino and neutralino states look more promising at future International Linear Collider (for a review see [139]).

The SE$_6$SSM also predicts the existence of other exotic neutralino and chargino states. Two exotic neutralino states and the second lightest exotic chargino are formed by the fermion components of the $SU(2)_W$ doublets. These fermions as well as their superpartners might be either light or heavy depending on the SE$_6$SSM parameters. Due to the Z_2^E symmetry conservation in the collider experiments all exotic particles can only be created in pairs. Since the exotic neutralino and chargino as well as their scalar partners do not couple to quarks/squarks directly at the LHC these states can be produced via the EW interactions. As a consequence their production cross section remains relatively small even if the corresponding states have masses around 1 TeV. The conservation of R-parity and Z_2^E symmetry implies that the final state in the decay of the exotic fermions involves at least one lightest exotic neutralino while the final state in the decay of its scalar partner should contain at least one lightest exotic neutralino and one gravitino. If both of the produced states decay into on-shell gauge bosons it is expected that they should result in some enhancements in the rates of

$$p\,p \to Z\,Z + E_T^{\rm miss} + X\,, \quad p\,p \to W\,Z + E_T^{\rm miss} + X\,, \quad p\,p \to W\,W + E_T^{\rm miss} + X\,. \tag{35}$$

where $E_T^{\rm miss}$ is associated with the lightest exotic fermion (and gravitino) and X should be identified with jets and/or extra charged leptons that may stem from the decays of intermediate states.

As mentioned before, the components of L_4 and $\overline{L}_4$ are expected to be relatively light. When all other exotic states and sparticles except $\chi_1^{\pm}$, χ_2, χ_1 and gravitino are rather heavy, the scalar ($\tilde{L}_4$) and fermionic (L_4) components of the supermultiplets L_4 and $\overline{L}_4$ can be produced in pairs via off-shell W and Z-bosons. Their decays always lead to either τ-lepton or electron/muon as well as missing energy in the final state. In the case of $\tilde{L}_4$ decays the missing energy in the final state can be associated with only one lightest exotic neutralino while the final state of the L_4 decays has to involve at least one lightest exotic neutralino and one gravitino to ensure the conservation of R-parity and Z_2^E symmetry. More efficiently L_4 and/or $\tilde{L}_4$ can be produced through the decays of the lightest exotic colored states if these states are relatively light and the corresponding decay channels are kinematically allowed.

Finally it is worth emphasising that the SE$_6$SSM predicts the existence of extra quarks and their scalar superpartners that carry lepton and baryon numbers simultaneously [38]. The LHC lower bounds on the scalar leptoquark masses [140–142] are not directly applicable in this case. Indeed, an ordinary scalar leptoquark with electric charge $-1/3$ decays either to the left-handed neutrino ν_j and down-type quark d_i or to charged lepton ℓ_j and up-type quark u_i. The lightest exotic colored state q_1 in the SE$_6$SSM, which is a superposition of either scalar or fermion components of the supermultiplets $\overline{D}_i$ and D_i, is odd under $\tilde{Z}_2^E$ symmetry. As a consequence its decays always lead to the missing energy $E_T^{\rm miss}$ in the final state

$$q_1 \to u_i(d_i) + \ell_j(\nu_j) + E_T^{\rm miss} + X\,. \tag{36}$$

The pair production of the lightest exotic colored states at the LHC may result in the enhancement of the cross sections of $pp \to jj + E_T^{\rm miss} + X$ and/or $pp \to jj\ell_k\overline{\ell}_m + E_T^{\rm miss} + X$. The LHC pair production cross section of the lightest exotic quarks changes from 10 fb to 1 fb if the mass of q_1 increases from 1.3 TeV to 1.7 TeV [143]. In the case of the lightest exotic squarks the production cross section is an order of magnitude smaller. The presence of Z' boson and exotic multiplets of matter in the particle spectrum is a very peculiar feature that should permit to distinguish the SE$_6$SSM from the MSSM and other extensions of the SM.

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