E\textsubscript{6} inspired SUSY models with Custodial Symmetry

R. Nevzorov\textsuperscript{*}

Alikhanov Institute for Theoretical and Experimental Physics, Moscow, 117218, Russia
*E-mail: nevzorov@itep.ru

The breakdown of E\textsubscript{6} within the supersymmetric (SUSY) Grand Unified Theories (GUTs) can result in SUSY extensions of the standard model (SM) based on the SM gauge group together with extra U(1) gauge symmetry under which right–handed neutrinos have zero charge. In these U(1)\textsubscript{N} extensions of the minimal supersymmetric standard model (MSSM) a single discrete 2\textsuperscript{N} Z\textsuperscript{'} symmetry may be used to suppress the most dangerous operators, that give rise to proton decay as well as non–diagonal flavour transitions at low energies. The SUSY models under consideration involves Z\textsuperscript{'} and extra exotic matter beyond the MSSM. We discuss leptogenesis within this SUSY model and argue that the extra exotic states may lead to the non–standard Higgs decays.

Keywords: Grand Unified Theories; Supersymmetry; Leptogenesis; Higgs boson.

1. Introduction

Supersymmetric (SUSY) extensions of the standard model (SM) allows one to embed SM into Grand Unified Theories (GUTs) based on simple gauge groups such as SU(5), SO(10) or E\textsubscript{6}. Indeed, it was found that the electroweak (EW) and strong gauge couplings extrapolated to high energies using the renormalisation group equation (RGE) evolution converge to a common value at some high energy scale in the framework of the minimal SUSY standard model (MSSM)\textsuperscript{1–4}. The incorporation of the SM gauge interactions within GUTs permits, in particular, to explain the peculiar assignment of U(1)\textsubscript{Y} charges postulated in the SM.

It is well known that each family of quarks and leptons fills in complete 16 dimensional spinor representation of SO(10) that also predicts the existence of right–handed neutrino, allowing it to be used for both the see–saw mechanism and leptogenesis. In N = 1 SUSY GUT based on E\textsubscript{6} the complete fundamental 27 representation, that decomposes under SO(10) \times U(1)\textsubscript{\phi} subgroup as

\[27 \rightarrow \begin{pmatrix} 16, \frac{1}{\sqrt{24}} \\ 10, -\frac{2}{\sqrt{24}} \\ 1, \frac{4}{\sqrt{24}} \end{pmatrix},\]

contains Higgs doublet. It is assigned to \(\begin{pmatrix} 10, -\frac{2}{\sqrt{24}} \end{pmatrix}\). The SM gauge bosons belong to the adjoint representation of E\textsubscript{6}, i.e. 78–plet. In N = 2 SUSY GUT based on the E\textsubscript{8} gauge group all SM particles belong to 248 dimensional representation of E\textsubscript{8} that decomposes under its E\textsubscript{6} subgroup as follows

\[248 \rightarrow 78 \oplus 3 \times 27 \oplus 3 \times 27 \oplus 8 \times 1.\]
renormalisable theory and has to be considered as an effective low energy limit of some renormalisable or even finite theory. Currently, the best candidate for such an underlying theory, i.e. hypothetical single framework that explains and links together all physical aspects of the universe, is ten–dimensional heterotic superstring theory based on $E_8 \times E_8'$. Compactification of extra dimensions leads to an effective supergravity and results in the breakdown of $E_8$ to $E_6$ or its subgroups in the observable sector. The remaining $E_8'$ plays the role of a hidden sector which gives rise to spontaneous breakdown of SUGRA.

2. The $U(1)_N$ extensions of the MSSM

In orbifold SUSY GUTs the $E_6$ gauge group can be broken down to $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_\chi \times U(1)_\psi$, where the $U(1)_\psi$ and $U(1)_\chi$ symmetries are defined by: $E_6 \rightarrow SO(10) \times U(1)_\psi$, $SO(10) \rightarrow SU(5) \times U(1)_\chi$. In order to ensure anomaly cancellation in this case the particle content below the GUT scale $M_X$ should be extended to include three 27–plets. Each 27–plet, referred to as $27_i$ with $i = 1, 2, 3$, includes one generation of ordinary matter, a SM singlet field $S_i$ (see last term in Eq. (1)), that carries non–zero $U(1)_\psi$ charge, as well as Higgs–like doublets ($H_u^i$ and $H_d^i$) and charged $\pm 1/3$ exotic quarks ($D_i$ and $\bar{D}_i$) which are associated with $\left(10, -\frac{2}{\sqrt{24}}\right)$ in Eq. (1). In addition the splitting of bulk 27–plets can give rise to a set of $M_l$ and $\overline{M}_l$ supermultiplets with opposite quantum numbers.

The presence of exotic matter in the $E_6$ inspired SUSY models generically leads to rapid proton decay and non–diagonal flavour transitions at low energies. To suppress flavour changing processes as well as the most dangerous baryon and lepton number violating operators one can impose a single discrete $\tilde{Z}_H^H$ symmetry. All states from complete $27_i$–plets are odd whereas all supermultiplets $M_l$ are even under this $\tilde{Z}_H^H$ symmetry. Because $M_l$ can be used for the breakdown of gauge symmetry this set of supermultiplets should contain $H_u$, $H_d$, $S$ and $N_{H}^i$. Since superfield $N_{H}^i$ has the same $U(1)_\psi$ and $U(1)_\chi$ charges as the right–handed neutrino the large vacuum expectation values (VEVs) of $N_{H}^i$ and $\overline{N}_{H}$ break $U(1)_\psi \times U(1)_\chi$ down to $U(1)_N$ generating large Majorana masses for the right–handed neutrinos. Only in this $E_6$ inspired $U(1)$ extension of the MSSM, i.e. in the so–called Exceptional Supersymmetric Standard Model (E$_6$SSM)$^{7,8}$, the right–handed neutrinos may be superheavy, shedding light on the origin of the mass hierarchy in the lepton sector and providing a mechanism for the generation of lepton and baryon asymmetry of the universe$^9$. Different phenomenological implications of the several variants of the $E_6$SSM were considered in Refs. $^{6–21}$. In particular, the renormalisation group (RG) flow of the gauge and Yukawa couplings as well as the theoretical upper bound on the lightest Higgs boson mass were examined in the vicinity of the quasi–fixed point$^{21}$ that appears as a result of the intersection of the invariant and quasi–fixed lines$^{22,23}$. Within the constrained version of the $E_6$SSM and its modifications the particle spectrum, the corresponding collider signatures and the implications for dark matter were analysed in Refs. $^{24–29}$. Here we assume that $U(1)_\psi \times U(1)_\chi$
symmetry is broken down to $U(1)_N \times Z_2^M$, where $Z_2^M = (-1)^{3(B-L)}$ is a matter parity. This can occur because $Z_2^M$ is a discrete subgroup of $U(1)_\psi$ and $U(1)_N$.

In the simplest case the set of the $Z_2^M$–even supermultiplets $M_i$ should also include a lepton $SU(2)_W$ doublet $L_4$ to allow the lightest exotic quarks to decay. The supermultiplets $\overline{M}_i$ can be either even or odd under the $Z_2^M$ symmetry. The simplest scenario imply that $\overline{S}$, $\overline{U}_u$ and $\overline{D}_d$ are odd whereas $\overline{L}_4$ is even under $Z_2^M$. It is expected that the $\overline{Z}_M$–odd supermultiplets $\overline{S}$, $\overline{U}_u$ and $\overline{D}_d$ get combined with the superposition of the appropriate components from $27_i$ forming vectorlike states with masses of order of $M_X$. At the same time the supermultiplets $L_4$ and $\overline{L}_4$ should form TeV scale vectorlike states to render the lightest exotic quarks unstable. The most general renormalisable superpotential which is allowed by the $Z_2^H$, $Z_2^M$ and $SU(3) \times SU(2)_W \times U(1)_Y \times U(1)_N$ symmetries can be written as

$$W = \lambda S(H_u H_d) + \lambda_{\alpha \beta} S(H_u^d H_d^\alpha) + \kappa_{ij} S(D_i \overline{D}_j) + \tilde{f}_{\alpha \beta} S_0 (H_d^\alpha H_u^\beta) + f_{\alpha \beta} S_0 (H_d^\alpha H_u^\beta) + \mu L_4 \overline{L}_4 + W_N$$

(3)

where $\alpha, \beta = 1, 2$ and $i, j = 1, 2, 3$ while $W_{\text{MSSM}}(\mu = 0)$ is the MSSM superpotential with the bilinear mass parameter $\mu$ set to zero and

$$W_N = \frac{1}{2} M_{ij} N_i^c N_j^c + \tilde{h}_{ij} N_i^c (H_u L_j) + h_{i0} N_i^c (H_u^0 L_4).$$

(4)

In Eqs. (3) and (4) $e_i^c$ and $N_i^c$ are the right-handed charged leptons and neutrinos whereas $Q_i$ and $L_j$ are the left-handed quark and lepton doublets respectively. The second last term in Eq. (3) ensures that the lightest exotic quarks decay within a reasonable time when the couplings $g_{ij}^D$ are sufficiently large and the components of the supermultiplets $L_4$ and $\overline{L}_4$ have masses of the order of a few TeV. Since in this case extra matter beyond the MSSM fill in complete $SU(5)$ representations the gauge coupling unification in the SUSY model under consideration can be achieved for any phenomenologically acceptable value of $\alpha_3(M_Z)$, consistent with its central measured low energy value $\alpha_3 \simeq 0.01$.

The $Z_2^H$–even supermultiplets $H_u$, $H_d$ and $S$ gain non–zero VEVs, i.e. $\langle H_u \rangle = v_1 / \sqrt{2}$, $\langle H_d \rangle = v_2 / \sqrt{2}$ and $\langle S \rangle = s / \sqrt{2}$, which are much smaller than the VEVs of $N_i^c$ and $\overline{N}_i$. In phenomenologically viable scenarios the SM singlet superfield $S$ has to acquire VEV which is much larger than 1 TeV breaking $U(1)_N$ gauge symmetry and inducing sufficiently large masses of $Z'$ boson and exotic fermion states. The neutral components of $H_u$ and $H_d$ develop VEVs, so that $v = \sqrt{v_1^2 + v_2^2} \simeq 246$ GeV. These VEVs trigger the breakdown of the $SU(2)_W \times U(1)_Y$ symmetry down to $U(1)_{\text{em}}$ associated with electromagnetism and give rise to the masses of ordinary quarks and leptons.

In the framework of the $E_6$SSM the Higgs sector was explored in Ref. [3]. When CP-invariance is preserved the Higgs spectrum contains three CP-even, one CP-odd and two charged states. The SM singlet dominated CP-even state and the $Z'$ boson are almost degenerate. If $\lambda < g_4'$, where $g_4'$ is the $U(1)_N$ gauge coupling,
the SM singlet dominated Higgs boson is the heaviest CP-even state. In this case the rest of the Higgs spectrum is basically indistinguishable from the one in the MSSM. When \( \lambda \gtrsim g' \), the Higgs spectrum has a very hierarchical structure, which is similar to the one in the NMSSM with the approximate PQ symmetry. As a consequence the mass matrix of the CP–even Higgs sector can be diagonalised using the perturbation theory. If \( \lambda \gtrsim g' \), the MSSM–like CP-even, CP-odd and charged states have almost the same masses and lie beyond the TeV range.

For the analysis of the phenomenological implications of the SUSY models discussed above it is convenient to introduce the \( Z^E_2 \) symmetry, which can be defined such that \( \tilde{Z}^H_2 = Z^M_2 \times Z^E_2 \). The supermultiplets \( S_\alpha, H^u_\alpha, H^d_\alpha, D_i, \bar{D}_i, L^4, \bar{L}^4 \) are odd under the \( Z^E_2 \) symmetry. The components of all other supermultiplets are \( Z^E_2 \) even. Because the Lagrangian is invariant under both \( Z^M_2 \) and \( \tilde{Z}^H_2 \) symmetries, the \( Z^E_2 \) symmetry is also conserved. This implies that in collider experiments the exotic particles, which are odd under the \( Z^E_2 \) symmetry, can only be created in pairs and the lightest exotic state has to be absolutely stable. Using the method proposed in Ref. 34 it was argued that the masses of the lightest exotic fermions, which are predominantly linear superpositions of the fermion components of the superfields \( S_\alpha \), do not exceed \( 60 - 65 \text{ GeV} \). Thus these states tend to be the lightest exotic particles in the spectrum. Moreover the lightest exotic fermion is also the lightest SUSY particle (LSP). Although the couplings of the corresponding states to the SM gauge bosons and fermions are quite small the lightest exotic state could account for all or some of the observed cold dark matter density if it had a mass close to half the \( Z^E_2 \) mass. However in this case the SM–like Higgs boson would decay almost 100% of the time into the fermion components of \( S_\alpha \). All other branching ratios would be strongly suppressed. Basically such scenario has been already ruled out by the LHC experiments. On the other hand if the lightest exotic fermions are substantially lighter than \( M_Z \) the annihilation cross section for LSP + LSP \( \rightarrow \) SM particles becomes too small leading to a relic density that is much larger than its measured value.

The simplest phenomenologically viable scenarios imply that the fermion components of \( S_\alpha \) are significantly lighter than 1 eV. In this scenario the lightest SUSY particles form hot dark matter in the Universe. When the masses of the fermion components of \( S_\alpha \) are considerably smaller than 1 eV these states give only a very minor contribution to the dark matter density. At the same time the invariance of the Lagrangian under the \( Z^M_2 \) symmetry ensures that the \( R \)-parity is also conserved and the lightest ordinary neutralino is stable. In this case the lightest ordinary neutralino may account for all or some of the observed cold dark matter density.

The scenarios discussed above are realised if \( \tilde{f}_{\alpha\beta} \sim f_{\alpha\beta} \lesssim 10^{-6} \). When the Yukawa couplings of the superfields \( S_\alpha \) are very small the terms \( \tilde{f}_{\alpha\beta} S_\alpha (H^d_\beta H_u) \) and \( f_{\alpha\beta} S_\alpha (H^d_\beta H_u) \) in the superpotential can be ignored. In this limit the low–energy

---

*a The presence of very light neutral fermions in the particle spectrum might have interesting implications for neutrino physics.*
effective Lagrangian possesses an approximate global $\mathbb{U}(1)_{E}$ symmetry below the scale $M_1$ where $M_1$ is the mass of the lightest right–handed neutrinos. The $\mathbb{U}(1)_{E}$ charges of the exotic matter fields are summarised in Table 1. Both $\mathbb{U}(1)_{B-L}$ and $\mathbb{U}(1)_{E}$ symmetries are explicitly broken because of the interactions of matter supermultiplets with $N_1^\alpha$ in $W_{N}$. As a consequence the decays of the lightest right–handed neutrino/sneutrino induce simultaneously $\mathbb{U}(1)_{E}$ and $\mathbb{U}(1)_{B-L}$ asymmetries. These asymmetries would not be washed out in the limit $f_{\alpha\beta}$, $f_{\alpha\beta} \to 0$. Moreover the sufficiently small values of the $\mathbb{U}(1)_{E}$ violating Yukawa couplings, i.e $f_{\alpha\beta}$, $f_{\alpha\beta} \lesssim 10^{-7}$, should not erase the induced $\mathbb{U}(1)_{E}$ asymmetry. The non-zero values of $f_{\alpha\beta}$ and $f_{\alpha\beta}$ break the $\mathbb{U}(1)_{E}$ symmetry and the lightest exotic state that carries the $\mathbb{U}(1)_{E}$ charge becomes unstable. It decays into the fermion components of $S_\alpha$ so that the generated $\mathbb{U}(1)_{E}$ asymmetry gets converted into the hot dark matter density.

Table 1. The $\mathbb{U}(1)_{E}$ charges of exotic matter supermultiplets.

| $H_\alpha^u$ | $H_\alpha^d$ | $D_4$ | $\overline{T}_4$ | $L_4$ | $\overline{T}_4$ |
|-------------|-------------|-------|----------------|-------|----------------|
| +1          | -1          | +1    | -1             | +1    | -1             |

3. Generation of baryon asymmetry

A potential drawback of supersymmetric thermal leptogenesis is the lower bound on $M_1$. Indeed, it was shown that the appropriate amount of the baryon asymmetry in the SM and MSSM can be induced only if $M_1$ is larger than $10^9$ GeV. In the framework of supergravity this lower bound on $M_1$ leads to the gravitino problem. After inflation the universe thermalizes with a reheat temperature $T_R$. If $T_R > M_1$, the right-handed neutrinos are produced by thermal scattering and thermal leptogenesis could take place. At the same time when $T_R \gtrsim 10^9$ GeV such a high reheating temperature results in an overproduction of gravitinos which tend to decay during or after Big Bang Nucleosynthesis (BBN) destroying the agreement between the predicted and observed light element abundances. It was argued that the gravitino density becomes low enough if $T_R \lesssim 10^{6-7}$ GeV. In order to avoid the gravitino problem we fix $M_1 \simeq 10^6$ GeV. We also assume that two other right-handed neutrino states have masses $M_{2,3} \sim 10^{6-7}$ GeV. For so low $M_i$ the absolute values of the Yukawa couplings $|\tilde{h}_{\alpha ij}|$ should be rather small to reproduce the left–handed neutrino mass scale $m_\nu \lesssim 0.1$ eV, i.e. $|\tilde{h}_{\alpha ij}|^2 \ll 10^{-8}$. So small Yukawa couplings can be ignored in the leading approximation. Then only the new channels of the decays of the lightest right–handed neutrino $N_1$ and its superpartner $\tilde{N}_1$, i.e.

$$N_1 \to L_4 + H_\alpha^u, \quad N_1 \to \tilde{L}_4 + \tilde{H}_\alpha^u, \quad \tilde{N}_1 \to L_4 + \tilde{H}_\alpha^u, \quad \tilde{N}_1 \to \tilde{L}_4 + H_\alpha^u.$$  (5)
give rise to the generation of lepton asymmetry. This process is controlled by the CP (decay) asymmetries associated with the decays of $N_1$, i.e.

$$
\varepsilon_{1,4}^\alpha = \frac{\Gamma^\alpha_{N_1 \ell_4} - \Gamma^\alpha_{\tilde{N}_1 \ell_4}}{\sum_{\beta} \left( \Gamma^\beta_{N_1 \ell_4} + \Gamma^\beta_{\tilde{N}_1 \ell_4} \right)}, \quad \varepsilon_{\bar{1},\bar{4}}^\alpha = \frac{\Gamma^\alpha_{N_1 \bar{\ell}_4} - \Gamma^\alpha_{\tilde{N}_1 \bar{\ell}_4}}{\sum_{\beta} \left( \Gamma^\beta_{N_1 \bar{\ell}_4} + \Gamma^\beta_{\tilde{N}_1 \bar{\ell}_4} \right)},
$$

(6)

and $\tilde{N}_1$, i.e.

$$
\varepsilon_{1,4}^\alpha = \frac{\Gamma^\alpha_{\tilde{N}_1 \ell_4} - \Gamma^\alpha_{\tilde{N}_1 \bar{\ell}_4}}{\sum_{\beta} \left( \Gamma^\beta_{\tilde{N}_1 \ell_4} + \Gamma^\beta_{\tilde{N}_1 \bar{\ell}_4} \right)}, \quad \varepsilon_{\bar{1},\bar{4}}^\alpha = \frac{\Gamma^\alpha_{\tilde{N}_1 \bar{\ell}_4} - \Gamma^\alpha_{\tilde{N}_1 \ell_4}}{\sum_{\beta} \left( \Gamma^\beta_{\tilde{N}_1 \bar{\ell}_4} + \Gamma^\beta_{\tilde{N}_1 \ell_4} \right)}.
$$

(7)

In Eqs. (6) and (7) the superscripts $\alpha$ and $\beta$ represent the components of the supermultiplets $H^u_\alpha$ and $H^d_\beta$ in the final state. At the tree level the partial decay widths associated with the new channels (5) are given by

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

(8)

and all decay asymmetries (6) and (7) vanish.

The non-zero values of the CP asymmetries arise after the inclusion of one–loop vertex and self–energy corrections to the decay amplitudes of $N_1$ and $\tilde{N}_1$. In this context it is worth noting that the supermultiplets $H^u_\alpha$ can be redefined so that only one doublet $H^u_1$ interacts with $L_4$ and $\tilde{N}_1$. Therefore without loss of generality $h_{12}$ in $W_N$ may be set to zero. In this limit $\varepsilon_{1,4}^2 = \varepsilon_{1,\bar{4}}^2 = \varepsilon_{\bar{1},\bar{4}}^2 = 0$. When SUSY breaking scale is negligibly small as compared with $M_1$, $h_{j1} = |h_{j1}| e^{i\varphi_{j1}}$ and $M_j$ are real the non–zero asymmetries are given by

$$
\varepsilon_{1,\bar{4}}^1 = \varepsilon_{1,\bar{4}}^\varphi = \varepsilon_{1,\bar{4}}^{\Delta \varphi} = \frac{1}{8\pi} \left[ \sum_{j=1,2} |h_{j1}|^2 f \left( \frac{M_j^2}{M_1^2} \right) \sin 2\Delta \varphi_{j1} \right],
$$

(9)

where $\Delta \varphi_{j1} = \varphi_{j1} - \varphi_{11}$ and

$$
f(z) = f^V(z) + f^S(z), \quad f^S(z) = \frac{2\sqrt{z}}{1-z}, \quad f^V(z) = -\sqrt{z} \ln \left( \frac{1+z}{z} \right).
$$

(10)

Because the Yukawa couplings of the superfields $N^c_1$ to the supermultiplets $H^u_\alpha$ and $L_4$ violate both $U(1)_E$ and $U(1)_{B-L}$ the decay channels of the lightest right–handed neutrino/sneutrino [5] induce simultaneously $U(1)_{B-L}$ and $U(1)_E$ asymmetries. These asymmetries are determined by the same set of the CP asymmetries [8].

The evolution of the $U(1)_{B-L}$ and $U(1)_E$ asymmetries are described by the system of Boltzmann equations. The generated baryon asymmetry can be estimated as follows

$$
Y_{\Delta B} \sim 10^{-3} \varepsilon_{1,\bar{4}}^1 \eta,
$$

(11)

where $Y_{\Delta B}$ is the baryon asymmetry relative to the entropy density, i.e.

$$
Y_{\Delta B} = \frac{n_B - \bar{n_B}}{s} \bigg|_0 = (8.75 \pm 0.23) \times 10^{-11}.
$$

(12)
Fig. 1. Logarithm (base 10) of the absolute value of $\omega = \epsilon_1 \ell_4 \eta$ as a function of logarithm (base 10) of $|h_{11}|$ for $h_{21} = h_{31} = 0$, $\Delta \varphi_{21} = \pi/4$ and $M_2 = 10 \cdot M_1$. The thick, solid and dashed lines correspond to $|h_{21}| = 0.3$, $|h_{21}| = 0.1$ and $|h_{21}| = 0.03$ respectively.

In Eq. 11 $\eta$ is an efficiency factor. It varies from 0 to 1. In the strong washout scenario $\eta$ is given by

$$\eta \simeq H(T = M_1)/\Gamma_1,$$

(13)

where $H$ is the Hubble expansion rate

$$H = 1.66g_s^{1/2} T^2/M_{Pl},$$

(14)

$g_s = n_b + \frac{7}{8} n_f$ is the number of relativistic degrees of freedom and

$$\Gamma_1 = \Gamma_{N_1 \ell_4}^{\dagger} + \Gamma_{N_1 \ell_4} = \frac{|h_{11}|^2}{8\pi} M_1.$$

(15)

As follows from Eq. 11 the values of the CP asymmetries are determined by the CP–violating phases $\Delta \varphi_{j1}$ and the absolute values of the Yukawa couplings $|h_{21}|$ and $|h_{31}|$ but do not depend on $|h_{11}|$. To simplify our analysis we fix $|h_{31}| = 0$ and $(M_2/M_1) = 10$. At the same time the efficiency factor $\eta$ is set by the lightest right–handed neutrino mass $M_1$ and $|h_{11}|$. We restrict our consideration here by the values of $|h_{11}|^2 \gg |\tilde{h}_{ik}|^2$, i.e. $|h_{11}|^2 \gtrsim 10^{-8}$. For $\Delta \varphi_{21} = \pi/4$ we find

$$\log |\eta| \simeq -2 \log |h_{11}| - 10.2, \quad \log |\omega| \simeq -2 \log |h_{11}| + 2 \log |h_{21}| - 12.1,$$

(16)

where $\omega = \epsilon_1 \ell_4 \eta$. Eq. 16 indicates that $\eta$ varies from $10^{-2}$ to $10^{-4}$ when $|h_{11}|$ increases from $10^{-4}$ to $10^{-3}$. The dependence of $|\omega|$, that determines the generated baryon asymmetry (11), on $|h_{21}|$ and $|h_{11}|$ is explored in Fig. 1. This figure illustrates that for $\Delta \varphi_{21} = \pi/4$ and $|h_{21}| \sim 0.1$ the phenomenologically acceptable baryon density, corresponding to $\omega \sim 10^{-7} - 10^{-6}$, can be obtained if $|h_{11}|$ varies between $10^{-4}$ and $10^{-3}$. If $f_{\alpha \beta}, f_{\alpha \beta} \lesssim 10^{-7}$ the induced dark matter and baryon number densities should be of the same order of magnitude.
4. Exotic Higgs decays

As it was mentioned before the lightest and second lightest exotic states ($\chi_1^0$ and $\chi_2^0$) are mostly linear superpositions of the fermion components of the superfields $S_i$. In the simplest phenomenologically viable scenarios $\chi_1^0$ should have mass $m_{\chi_1} \ll 1$ eV. At the same time $\chi_2^0$ can be considerably heavier if some of the Yukawa couplings $\tilde{f}_{\alpha\beta}$ and $f_{\alpha\beta}$ are much larger than $10^{-6} - 10^{-5}$. Although $\chi_1^0$ and $\chi_2^0$ tend to be rather light their couplings to the $Z$-boson and other SM particles can be negligibly small because these states are predominantly the fermion components of the SM singlet superfields $S_i$. As a result any possible signal, which $\chi_1^0$ and $\chi_2^0$ could give rise to at former and present collider experiments, would be extremely suppressed and such states could escape their experimental detection.

The couplings of the lightest Higgs boson $h_1$ to $\chi_1^0$ and $\chi_2^0$ are determined by their masses. Since $\chi_1^0$ is extremely light it does not affect Higgs phenomenology. The coupling of the SM-like Higgs state $h_1$ to the second lightest exotic particle $X_{22}^h \simeq |m_{\chi_2}|/v^{13}$. This coupling gives rise to the decays of $h_1$ into $\chi_2^0$ pairs with partial width given by

$$\Gamma(h_1 \to \chi_2^0 \chi_2^0) = \frac{(\chi_{22}^h)^2 m_{h_1}}{4\pi} \left(1 - 4 |m_{\chi_2}|^2/m_{h_1}^2 \right)^{3/2},$$

(17)

where $m_{h_1}$ is the lightest Higgs boson mass. From Eq. (17) it follows that the partial decay width of the non-standard Higgs decays depend rather strongly on $m_{\chi_2}$. The branching ratio of $h_1 \to \chi_2^0 \chi_2^0$, can be substantial if the second lightest exotic fermion has a mass of order of the $b$-quark mass $m_b$. To avoid the suppression of the branching ratios for Higgs decays into SM particles we restrict our consideration to the GeV scale masses of the second lightest exotic particle.

After being produced $\chi_2^0$ sequentially decay into $\chi_1^0$ and fermion–antifermion pair via virtual $Z$. Thus the exotic decays of the SM-like Higgs discussed above results in two fermion–antifermion pairs and missing energy in the final state. Nevertheless due to the small coupling of the lightest and second lightest exotic fermions to the $Z$-boson $\chi_2^0$ tends to live longer than $10^{-8}$ sec. Therefore it typically decays outside the detectors and can not be observed at the LHC directly. As a consequence the decay channel $h_1 \to \chi_2^0 \chi_2^0$ normally gives rise to an invisible branching ratio of $h_1$. If the second lightest exotic fermion is very long-lived then $\chi_2^0$ may decay during or after Big Bang Nucleosynthesis (BBN) destroying the agreement between the predicted and observed light element abundances. To preserve the success of the BBN, the lifetime $\tau_{\chi_2}$ of $\chi_2^0$ should not be longer than 1 sec. Because $\tau_{\chi_2} \sim 1/(m_{\chi_2}^2$ this requirement basically rules out too light $\chi_2^0$. Indeed, it is somewhat problematic to satisfy this restriction for $m_{\chi_2} \lesssim 100$ MeV.

The numerical analysis indicates that the branching ratio associated with the decays $h_1 \to \chi_2^0 \chi_2^0$ can vary from 0.2% to 20% when $m_{\chi_2}$ changes from 0.3 GeV to 2.7 GeV. When $\chi_2^0$ is lighter than 0.5 GeV the corresponding branching ratio can be as small as $10^{-3} - 10^{-4}$.
References

1. J. R. Ellis, S. Kelley and D. V. Nanopoulos, Probing the desert using gauge coupling unification, *Phys. Lett. B* **260**, 131 (1991).
2. P. Langacker and M. Luo, Implications of precision electroweak experiments for $M_t$, $\rho$, $\sin^2\theta_W$ and grand unification, *Phys. Rev. D* **44**, 817 (1991).
3. U. Amaldi, W. de Boer and H. Furstenau, Comparison of grand unified theories with electroweak and strong coupling constants measured at LEP, *Phys. Lett. B* **260**, 447 (1991).
4. F. Anselmo, L. Cifarelli, A. Peterman and A. Zichichi, The Effective experimental constraints on $M_{\text{susy}}$ and $M_{\text{gut}}$, *Nuovo Cim. A* **104**, 1817 (1991).
5. M. B. Green, J. H. Schwarz, E. Witten, *Superstring Theory* (Cambridge Univ. Press, Cambridge, 1987).
6. R. Nevzorov, $E_6$ inspired supersymmetric models with exact custodial symmetry, *Phys. Rev. D* **87**, 015029 (2013).
7. S. F. King, S. Moretti and R. Nevzorov, Theory and phenomenology of an exceptional supersymmetric standard model, *Phys. Rev. D* **73**, 035009 (2006).
8. S. F. King, S. Moretti and R. Nevzorov, Exceptional supersymmetric standard model, *Phys. Lett. B* **634**, 278 (2006).
9. S. F. King, R. Luo, D. J. Miller and R. Nevzorov, Leptogenesis in the Exceptional Supersymmetric Standard Model: Flavour dependent lepton asymmetries, *JHEP* **0812**, 042 (2008).
10. S. F. King, S. Moretti and R. Nevzorov, $E_6$SSM, *AIP Conf. Proc.* **881**, 138 (2007).
11. S. F. King, S. Moretti and R. Nevzorov, Gauge coupling unification in the exceptional supersymmetric standard model, *Phys. Lett. B* **650**, 57 (2007).
12. P. Athron, J. P. Hall, R. Howl, S. F. King, D. J. Miller, S. Moretti and R. Nevzorov, Aspects of the exceptional supersymmetric standard model, *Nucl. Phys. Proc. Suppl.* **200-202**, 120 (2010).
13. J. P. Hall, S. F. King, R. Nevzorov, S. Pakvasa and M. Sher, Novel Higgs Decays and Dark Matter in the $E_6$SSM, *Phys. Rev. D* **83**, 075013 (2011).
14. J. P. Hall, S. F. King, R. Nevzorov, S. Pakvasa and M. Sher, Nonstandard Higgs decays in the $E_6$SSM, *PoS QFTHEP* **2010**, 069 (2010).
15. R. Nevzorov and S. Pakvasa, Exotic Higgs decays in the $E_6$ inspired SUSY models, *Phys. Lett. B* **728**, 210 (2014).
16. P. Athron, M. Mihleitein, R. Nevzorov and A. G. Williams, Non-Standard Higgs Decays in U(1) Extensions of the MSSM, *JHEP* **1501**, 153 (2015).
17. R. Nevzorov and S. Pakvasa, Nonstandard Higgs decays in the $E_6$ inspired SUSY models, *Nucl. Part. Phys. Proc.* **273-275**, 690 (2016).
18. R. Nevzorov, LHC Signatures and Cosmological Implications of the $E_6$ Inspired SUSY Models, *PoS EPS -HEP2015*, 381 (2015).
19. S. F. King and R. Nevzorov, 750 GeV Diboson Resonance from Singlets in an Exceptional Supersymmetric Standard Model, *JHEP* **1603**, 139 (2016).
20. R. Nevzorov, Leptogenesis as an origin of hot dark matter and baryon asymmetry in the $E_6$ inspired SUSY models, *Phys. Lett. B* **779**, 223 (2018).
21. R. Nevzorov, Quasifixed point scenarios and the Higgs mass in the $E_6$ inspired supersymmetric models, *Phys. Rev. D* **89**, 055010 (2014).
22. R. B. Nevzorov and M. A. Trusov, Infrared quasifixed solutions in the NMSSM, *Phys. Atom. Nucl. Nucl. 64*, 1299 (2001).
23. R. B. Nevzorov and M. A. Trusov, Quasifixed point scenario in the modified NMSSM, *Phys. Atom. Nucl. 65*, 335 (2002).
24. P. Athron, S. F. King, D. J. Miller, S. Moretti and R. Nevzorov, Predictions of the
Constrained Exceptional Supersymmetric Standard Model, *Phys. Lett. B* **681**, 448 (2009).

25. P. Athron, S. F. King, D. J. Miller, S. Moretti and R. Nevzorov, The Constrained Exceptional Supersymmetric Standard Model, *Phys. Rev. D* **80**, 035009 (2009).

26. P. Athron, S. F. King, D. J. Miller, S. Moretti and R. Nevzorov, LHC Signatures of the Constrained Exceptional Supersymmetric Standard Model, *Phys. Rev. D* **84**, 055006 (2011).

27. P. Athron, S. F. King, D. J. Miller, S. Moretti and R. Nevzorov, Constrained Exceptional Supersymmetric Standard Model with a Higgs Near 125 GeV, *Phys. Rev. D* **86**, 095003 (2012).

28. P. Athron, D. Harries, R. Nevzorov and A. G. Williams, $E_6$ Inspired SUSY benchmarks, dark matter relic density and a 125 GeV Higgs, *Phys. Lett. B* **760**, 19 (2016).

29. P. Athron, D. Harries, R. Nevzorov and A. G. Williams, Dark matter in a constrained $E_6$ inspired SUSY model, *JHEP* **1612**, 128 (2016).

30. D. J. Miller, R. Nevzorov and P. M. Zerwas, The Higgs sector of the next-to-minimal supersymmetric standard model, *Nucl. Phys. B* **818**, 3 (2016).

31. P. A. Kovalenko, R. Nevzorov and K. A. Ter-Martirosyan, Masses of Higgs bosons in supersymmetric theories, *Phys. Rev. D* **50**, 265 (1994).

32. R. Nevzorov and M. A. Trusov, Particle spectrum in the modified NMSSM in the strong Yukawa coupling limit, *J. Exp. Theor. Phys.** **91**, 1079 (2000).

33. R. Nevzorov, K. A. Ter-Martirosyan and M. A. Trusov, Higgs bosons in the simplest SUSY models, *Phys. Rev. Lett.** **85**, 265 (2000).

34. S. Hesselbach, D. J. Miller, G. Moortgat-Pick, R. Nevzorov and M. Trusov, Theoretical upper bound on the mass of the LSP in the MNSSM, *Phys. Lett. B* **662**, 199 (2008).

35. J. M. Frere, R. Nevzorov and M. I. Vysotsky, Stimulated neutrino conversion and bounds on neutrino magnetic moments, *Phys. Lett. B* **394**, 127 (1997).

36. B. A. Campbell, S. Davidson, J. R. Ellis and K. A. Olive, Cosmological baryon asymmetry constraints on extensions of the standard model, *Phys. Lett. B* **256**, 484 (1991).

37. S. Davidson and R. Hempfling, Protecting the baryon asymmetry in theories with R-parity violation, *Phys. Lett. B* **391**, 287 (1997).

38. R. N. Mohapatra, Supersymmetry and R-parity: an Overview, *Phys. Scripta** **90**, 088004 (2015).

39. S. Davidson and A. Ibarra, A Lower bound on the right-handed neutrino mass from leptogenesis, *Phys. Lett. B* **535**, 25 (2002).

40. K. Hamaguchi, H. Murayama and T. Yanagida, Leptogenesis from N dominated early universe, *Phys. Rev. D* **65**, 043512 (2002).

41. M. Y. Khlopov and A. D. Linde, Is It Easy to Save the Gravitino?, *Phys. Lett. B* **138**, 265 (1984).

42. J. R. Ellis, J. E. Kim and D. V. Nanopoulos, Cosmological Gravitino Regeneration and Decay, *Phys. Lett. B* **145**, 181 (1984).

43. M. Y. Khlopov, Y. L. Levitan, E. V. Sedelnikov and I. M. Sobol, Nonequilibrium cosmological nucleosynthesis of light elements: Calculations by the Monte Carlo method, *Phys. Atom. Nucl.* **57**, 1393 (1994).

44. M. Kawasaki, K. Kohri and T. Moroi, Big-Bang nucleosynthesis and hadronic decay of long-lived massive particles, *Phys. Rev. D* **71**, 083502 (2005).

45. K. Kohri, T. Moroi and A. Yotsuyanagi, Big-bang nucleosynthesis with unstable gravitino and upper bound on the reheating temperature, *Phys. Rev. D* **73**, 123511 (2006).