Exceptional Supersymmetric Standard Model

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**Abstract**

We discuss some phenomenological aspects of an $E_6$ inspired supersymmetric standard model with an extra $U(1)_N$ gauge symmetry under which right-handed neutrinos have zero charge, allowing a conventional see-saw mechanism. The $\mu$ problem is solved in a similar way to the NMSSM, but without the accompanying problems of singlet tadpoles or domain walls. The above exceptional supersymmetric standard model (ESSM) involves the low energy matter content of three 27 representations of $E_6$, which is broken at the GUT scale, and allows gauge coupling unification due to an additional pair of Higgs-like doublets. The ESSM predicts a $Z'$ boson and exotic quarks which, if light enough, will provide spectacular new physics signals at the LHC. We study the LHC phenomenology of the $Z'$ and extra quarks, including their production and decay signatures particular to the ESSM. We also discuss the two-loop upper bound on the mass of the lightest CP-even Higgs boson, and show that it can be significantly heavier than in either the MSSM or the NMSSM.

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1. Introduction

The minimal supersymmetric standard model (MSSM) provides a very attractive supersymmetric extension of the standard model (SM) in which the superpotential contains the bilinear term $\mu H_d H_u$, where $H_d, H_u$ are the two Higgs doublets which develop vacuum expectation values (VEVs) at the weak scale and $\mu$ is the supersymmetric Higgs mass parameter which can be present before SUSY is broken. However, despite its attractiveness, the MSSM suffers from the $\mu$ problem: one would naturally expect $\mu$ to be either zero or of the order of the Planck scale, while, in order to get the correct pattern of electroweak symmetry breaking (EWSB), $\mu$ is required to be in the TeV range. The next-to-minimal supersymmetric standard model (NMSSM) is an attempt to solve the $\mu$ problem of the MSSM by generating the aforementioned term dynamically as the low energy VEV of a singlet field $S$ via the interaction $\lambda SH_d H_u$. In order to avoid a low energy global $U(1)$ symmetry, the superpotential is also supplemented by a trilinear term $S^3$. However the superpotential of the NMSSM remains invariant under a discrete $Z_3$ symmetry which, when broken at the weak scale, leads to the formation of domain walls in the early universe, which are inconsistent with modern cosmology. In an attempt to break the $Z_3$ symmetry, operators suppressed by powers of the Planck scale could be introduced. But these give rise to quadratically divergent tadpole contributions which would destabilize the mass hierarchy. (For a review of the MSSM and NMSSM see e.g. [1].)

An elegant solution to the $\mu$ problem can emerge in the framework of ten dimensional heterotic superstring theory based on $E_8 \times E_8$ [2]. Compactification of the extra dimensions results in the breakdown of $E_8$ down to $E_6$ or one of its subgroups in the observable sector [3]. At the string scale, $E_6$ can be broken directly to the rank-6 subgroup $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\psi \times U(1)_\chi$ via the Hosotani mechanism [4]. Two anomaly-free $U(1)_\psi$ and $U(1)_\chi$ symmetries of the rank-6 model are defined by [5]:

$$E_6 \to SO(10) \times U(1)_\psi, \quad SO(10) \to SU(5) \times U(1)_\chi.$$  

In this article we explore a particular $E_6$ inspired supersymmetric model with one extra $U(1)_N$ gauge symmetry defined by:

$$U(1)_N = \frac{1}{4} U(1)_\chi + \frac{\sqrt{15}}{4} U(1)_\psi,$$

under which right-handed neutrinos have no charge and thus may gain large Majorana masses in accordance with the see-saw mechanism (for a review see e.g. [6]). The extra $U(1)_N$ gauge symmetry survives to low energies and serves to forbid an elementary $\mu$ term as well as terms like $S^a$ in the superpotential but allows the interaction $\lambda SH_d H_u$. After EWSB the scalar component of the singlet superfield acquires a non-zero VEV, $\langle S \rangle = s/\sqrt{2}$, breaking $U(1)_N$ and an effective $\mu = \lambda s/\sqrt{2}$ term is automatically generated. Clearly there are no domain wall problems in such a model since there is no discrete $Z_3$ symmetry, and instead of a global symmetry there is a gauged $U(1)_N$. Anomalies are
cancelled by complete 27 representations of $E_6$ which survive to low energies, even though $E_6$ is broken at the GUT scale.

We refer to the model described above as the exceptional supersymmetric standard model (ESSM). The ESSM thus represents a low energy alternative to the MSSM or NMSSM, and provides a solution to the $\mu$ problem without domain wall problems. The ESSM contains a rich phenomenology accessible to the LHC in the form of a $Z'$ plus three families of exotic quarks and non-Higgs doublets. In a companion paper we have made a comprehensive study of the theory and phenomenology of the ESSM \cite{7}. The purpose of this accompanying letter is to summarize the phenomenological highlights of our study, including the two loop upper bound on the lightest CP-even Higgs mass, and the LHC phenomenology of the $Z'$ and exotic quarks (including some new phenomenological results) in a form that will be more easily accessible to our phenomenological and experimental colleagues. For more details we refer the interested reader to the accompanying full length paper \cite{7}. Previously, the implications of SUSY models with an additional $U(1)_N$ gauge symmetry had been studied in the context of leptogenesis \cite{8}, EW baryogenesis \cite{9} and neutrino physics \cite{10}. Supersymmetric models with a $U(1)_N$ gauge symmetry under which right-handed neutrinos are neutral have been specifically considered in \cite{11} from the point of view of $Z - Z'$ mixing and the neutralino sector, in \cite{12} where a renormalization group (RG) analysis was performed, and in \cite{13} where a one-loop Higgs mass upper bound was presented.

In the next section we briefly review the ESSM. In sect. 3 we analyse the upper bound on the lightest CP-even Higgs boson mass including leading two-loop corrections. Then in sect. 4 we discuss the phenomenology of some of the extra particles predicted by the ESSM and analyze their production cross-sections and signatures at the LHC. Our results are summarized in sect. 5.

2. The ESSM

One of the most important issues in models with additional Abelian gauge symmetries is the cancellation of anomalies. In $E_6$ theories the anomalies are cancelled automatically. Therefore any model based on $E_6$ subgroups which contains complete representations should be anomaly-free. Thus in order to ensure anomaly cancellation the particle content of the ESSM should include complete fundamental 27 representations of $E_6$. These multiplets decompose under the $SU(5) \times U(1)_N$ subgroup of $E_6$ \cite{12} as follows:

$$27_i \rightarrow (10, 1)_i + (5^*, 2)_i + (5^*, -3)_i + (5, -2)_i + (1, 5)_i + (1, 0)_i.$$  \hspace{1cm} (2)

The first and second quantities in the brackets are the $SU(5)$ representation and extra $U(1)_N$ charge while $i$ is a family index that runs from 1 to 3. An ordinary SM family
which contains 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, is assigned to $(10, 1)_i + (5^*, 2)_i$. Right-handed neutrinos $N_i^c$ should be associated with the last term in Eq. (2) $(1, 0)_i$. The next-to-last term in Eq. (2) $(1, 5)_i$ represents SM-type singlet fields $S_i$ which carry non-zero $U(1)_N$ charges and therefore survive down to the EW scale. The pair of $SU(2)$-doublets ($H_{1i}$ and $H_{2i}$) that are contained in $(5^*, -3)_i$ and $(5, -2)_i$ have the quantum numbers of Higgs doublets. Other components of these $SU(5)$ multiplets form color triplet of exotic quarks $D_i$ and $\overline{D}_i$ with electric charges $-1/3$ and $+1/3$ respectively.

The matter content and correctly normalized Abelian charge assignment are in Tab. 1.

|   | $Q$ | $u^c$ | $d^c$ | $L$ | $e^c$ | $N^c$ | $S$ | $H_2$ | $H_1$ | $D$ | $\overline{D}$ | $H'$ | $H''$ |
|---|-----|------|------|-----|------|------|-----|-------|-------|----|--------|-----|------|
| $\sqrt{\frac{5}{6}}Q_i^N$ | $\frac{1}{6}$ | $-\frac{2}{3}$ | $\frac{1}{3}$ | $-\frac{1}{2}$ | 1 | 0 | 0 | $\frac{1}{2}$ | $-\frac{1}{2}$ | $\frac{1}{3}$ | $-\frac{1}{2}$ | $\frac{1}{2}$ |
| $\sqrt{40}Q_i^Y$ | 1 | 1 | 2 | 2 | 1 | 0 | 5 | $-2$ | $-3$ | $-3$ | 2 | $-2$ |

Table 1: The $U(1)_Y$ and $U(1)_N$ charges of matter fields in the ESSM, where $Q_i^N$ and $Q_i^Y$ are here defined with the correct $E_6$ normalization factor required for the RG analysis.

The most general renormalizable superpotential which is allowed by the $E_6$ symmetry can be written in the following form:

$$W_{E_6} = W_0 + W_1 + W_2,$$

$$W_0 = \lambda_{ijk} S_i (H_{1j} H_{2k}) + \kappa_{ijk} S_i (D_j \overline{D}_k) + h_{ijk}^N N_i^c (H_{2j} L_k) + h_{ijk}^U u_i^c (H_{2j} Q_k) + h_{ijk}^D d_i^c (H_{1j} Q_k) + h_{ijk}^E e_i^c (H_{1j} L_k),$$

$$W_1 = g_{ijk}^Q D_i (Q_j Q_k) + g_{ijk}^D \overline{D}_i d_j^c u_k^c,$$

$$W_2 = g_{ijk}^N N_i^c D_j d_k^c + g_{ijk}^E e_i^c D_j e_k^c + g_{ijk}^D (Q_i L_j) \overline{D}_k.$$

Although $B - L$ is conserved automatically, some Yukawa interactions in Eq. (3) violate baryon number conservation resulting in rapid proton decay. The baryon and lepton number violating operators can be suppressed by imposing an appropriate $Z_2$ symmetry which is usually called $R$-parity. But the straightforward generalization of the definition of $R$-parity, assuming $B_D = 1/3$ and $B_{\overline{D}} = -1/3$, implies that $W_1$ and $W_2$ are forbidden by this symmetry and the lightest exotic quark is stable. Models with stable charged exotic particles are ruled out by different experiments [14].

To prevent rapid proton decay in $E_6$ supersymmetric models a generalized definition of $R$-parity should be used. There are two ways to do that. If $H_{1i}$, $H_{2i}$, $S_i$, $D_i$, and the quark superfields $(Q_i$, $u_i^c$, $d_i^c$) are even under a discrete $Z_2'$ symmetry while the lepton superfields $(L_i$, $e_i^c$, $N_i^c$) are odd all terms in $W_2$ are forbidden (Model I). Then the remaining superpotential is invariant with respect to a $U(1)_B$ global symmetry if the exotic quarks $\overline{D}_i$ and $D_i$ are diquark and anti-diquark, i.e. $B_D = -2/3$ and $B_{\overline{D}} = 2/3$. 

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An alternative possibility is to assume that the exotic quarks $D_i$ and $\bar{D}_i$ as well as lepton superfields are all odd under $Z_2^B$ whereas the others remain even. Then we get Model II in which all Yukawa interactions in $W_1$ are forbidden by the discrete $Z_2^B$ symmetry. Here, exotic quarks are leptoquarks. The two possible models are summarized as:

$$W_{ESSM1} = W_0 + W_1, \quad W_{ESSMII} = W_0 + W_2. \quad (4)$$

In addition to the complete $27_i$ representations some components of the extra $27'$ and $\overline{27'}$ representations must survive to low energies in order to preserve gauge coupling unification. We assume that an additional $SU(2)$ doublet components $H'$ of $(5^*, 2)$ from a $27'$ and corresponding anti-doublet $\overline{H'}$ from $\overline{27'}$ survive to low energies.

In either model the superpotential involves a lot of new Yukawa couplings in comparison to the SM. In general these new interactions induce non-diagonal flavor transitions. To suppress flavor changing processes one can postulate a $Z_2^H$ symmetry under which all superfields except one pair of $H_{1i}$ and $H_{2i}$ (say $H_d \equiv H_{13}$ and $H_u \equiv H_{23}$) and one SM-type singlet field $S \equiv S_3$ are odd. Then only one Higgs doublet $H_d$ interacts with the down-type quarks and charged leptons and only one Higgs doublet $H_u$ couples to up-type quarks while the couplings of all other exotic particles to ordinary quarks and leptons are forbidden. This eliminates any problem related with non-diagonal flavor transitions. The $SU(2)$ doublets $H_u$ and $H_d$ play the role of Higgs fields generating the masses of quarks and leptons after EWSB. Thus it is natural to assume that only $S$, $H_u$ and $H_d$ acquire non-zero VEVs.

The $Z_2^H$ symmetry reduces the structure of the Yukawa interactions in (4):

$$W_{ESSM1,II} \rightarrow \lambda_i S(H_{1i}H_{2i}) + \kappa_i S(D_i\bar{D}_i) +$$

$$+ f_{\alpha\beta} S_\alpha(H_dH_{23}) + \tilde{f}_{\alpha\beta} S_\alpha(H_{13}H_u) + W_{MSSM}(\mu = 0), \quad (5)$$

where $\alpha, \beta = 1, 2$ and $i = 1, 2, 3$. In Eq. (5) we choose the basis of $H_{1\alpha}$, $H_{2\alpha}$, $D_i$ and $\bar{D}_i$ so that the Yukawa couplings of the singlet field $S$ have flavor diagonal structure. Here we define $\lambda \equiv \lambda_3$ and $\kappa \equiv \kappa_3$. If $\lambda$ or $\kappa_i$ are large at the Grand Unification scale $M_X$ they affect the evolution of the soft scalar mass $m_S^2$ of the singlet field $S$ rather strongly resulting in negative values of $m_S^2$ at low energies that triggers the breakdown of the $U(1)_X$ symmetry. To guarantee that only $H_u$, $H_d$ and $S$ acquire a VEV we impose a certain hierarchy between the couplings $H_{1i}$ and $H_{2i}$ to the SM-type singlet superfields $S_i$: $\lambda \gg \lambda_{1,2}$, $f_{\alpha\beta}$ and $\tilde{f}_{\alpha\beta}$. Although $\lambda_{1,2}$, $f_{\alpha\beta}$ and $\tilde{f}_{\alpha\beta}$ are expected to be considerably smaller than $\lambda$ they must be large enough to generate sufficiently large masses for the exotic particles to avoid conflict with direct particle searches at present and former accelerators. Keeping only Yukawa interactions whose couplings are allowed to be of order unity gives approximately:

$$W_{ESSM1,II} \approx \lambda S(H_dH_u) + \kappa_i S(D_i\bar{D}_i) + h_t(H_uQ)t^c + h_b(H_dQ)b^c + h_\tau(H_dL)\tau^c. \quad (6)$$
The $Z^H_2$ symmetry discussed above forbids all terms in $W_1$ or $W_2$ that would allow the exotic quarks to decay. Therefore the discrete $Z^H_2$ symmetry can only be approximate although $Z^D_2$ or $Z^L_2$ must be exact to prevent proton decay. In our model we allow only the third family $SU(2)$ doublets $H_d$ and $H_u$ to have Yukawa couplings to the ordinary quarks and leptons of order unity. This is a self-consistent assumption since the large Yukawa couplings of the third generation (in particular, the top-quark Yukawa coupling) provide a radiative mechanism for generating the Higgs VEVs [15]. The Yukawa couplings of two other pairs of $SU(2)$ doublets $H_{1i}$ and $H_{2i}$ as well as $H'$ and exotic quarks to the quarks and leptons of the third generation are supposed to be significantly smaller ($\lesssim 0.1$) so that none of the other exotic bosons gain VEVs. These couplings break the $Z^H_2$ symmetry explicitly resulting in flavor changing neutral currents (FCNCs). In order to suppress the contribution of new particles and interactions to $K^0 - \bar{K}^0$ oscillations and to the muon decay channel $\mu \rightarrow e^-e^+e^-$ in accordance with experimental limits, it is necessary to assume that the Yukawa couplings of new exotic particles to the quarks and leptons of the first and second generations are less than or of order $10^{-4}$.

3. Upper bound on the lightest CP-even Higgs mass

As in the NMSSM, the ESSM Higgs sector includes two doublets $H_u$ and $H_d$ as well as the SM-type singlet field $S$. The interactions between them are defined by the structure of the gauge interactions and by Eq.(6). At the physical vacuum Higgs fields develop the VEVs $\langle H_d \rangle = \frac{v_d}{\sqrt{2}}$, $\langle H_u \rangle = \frac{v_u}{\sqrt{2}}$ and $\langle S \rangle = \frac{s}{\sqrt{2}}$, thus breaking the $SU(2)_L \times U(1)_Y \times U(1)_N$ symmetry to $U(1)_{EM}$, associated with electromagnetism. Instead of $v_d$ and $v_u$ it is more convenient to use $\tan \beta = \frac{v_u}{v_d}$ and $v = \sqrt{v_d^2 + v_u^2}$, where $v = 246$ GeV. After the breakdown of the gauge symmetry two CP-odd and two charged Goldstone modes in the Higgs sector are absorbed by the $Z$, $Z'$ and $W^\pm$ gauge bosons so that only six physical degrees of freedom are left. They represent three CP-even (as in the NMSSM), and one CP-odd and two charged Higgs states (as in the MSSM). However, unlike the MSSM or NMSSM, there is an additional $Z'$, as well as exotic matter, as discussed in the next section.

SUSY models predict that the mass of the lightest Higgs particle is limited from above. The ESSM is not an exception. When the soft masses of the superpartners of the top-quark are equal, i.e. $m_{Q}^2 = m_{U}^2 = M_2^2$, the two-loop upper bound on the lightest CP-even Higgs boson mass $m_h$ in the ESSM can be written in the following form:

$$m_h^2 \lesssim \left[ \frac{\lambda^2}{2} v^2 \sin^2 2\beta + M_Z^2 \cos^2 2\beta + g_1^2 v^2 \left( \tilde{Q}_1 \cos^2 \beta + \tilde{Q}_2 \sin^2 \beta \right) \right] \left( 1 - \frac{3h^2}{8\pi^2} l \right) + \frac{3h^4 v^2 \sin^4 \beta}{8\pi^2} \left\{ \frac{1}{2} U_l + l + \frac{1}{16\pi^2} \left( \frac{3}{2} h_t^2 - 8g_3^2 \right) (U_l + l) \right\},$$

(7)
where

\[ U_t = 2 \frac{X_t^2}{M_S^2} \left( 1 - \frac{1}{12} \frac{X_t^2}{M_S^2} \right), \quad l = \ln \left[ \frac{M^2}{m^2} \right], \]  

(8)

\( X_t \) is the usual stop mixing parameter, \( \bar{g} = \sqrt{g_2^2 + g'^2} \), \( g' = \sqrt{3/5}g_1 \), whereas \( \tilde{Q}_1 \) and \( \tilde{Q}_2 \) are the \( U(1)_N \) charges of \( H_d \) and \( H_u \) respectively. Here \( g_2 \), \( g_1 \) and \( g'_1 \) are the gauge couplings of the \( SU(2)_L \), \( U(1)_Y \) and \( U(1)_N \) interactions. The couplings \( g_2 \) and \( g' \) are known precisely. Thus, by assuming gauge coupling unification one can determine the value of the extra \( U(1)_N \) gauge coupling. It turns out that \( g'_1(Q) \simeq g_1(Q) \) for any renormalization scale \( Q \) smaller than the unification one: \( Q \lesssim M_X \).

Eq. (7) is a simple generalization of the approximate expressions for the two-loop theoretical restriction on the mass of the lightest Higgs particle obtained in the MSSM [16] and NMSSM [17]. At tree level the upper limit on the mass of the lightest Higgs particle is described by the sum of the three terms in the square brackets. One-loop corrections from the top-quark and its superpartners increase the bound on the lightest CP-even Higgs boson mass in SUSY models substantially. The inclusion of leading two-loop corrections reduces the upper limit on \( m_h \) significantly and nearly compensates the growth of the theoretical restriction on \( m_h \) with increasing SUSY breaking scale \( M_S \) which is caused by one-loop corrections. In order to enhance the contribution of loop effects we assume maximal mixing in the stop sector (i.e. \( X_t = \sqrt{6}M_S \)). We also adopt \( M_S = 700 \) GeV. Then the theoretical restriction on the lightest Higgs mass [17] depends on \( \lambda \) and \( \tan \beta \) only. The requirement of validity of perturbation theory up to the Grand Unification scale constrains the parameter space further setting a limit on the Yukawa coupling \( \lambda \) for each value of \( \tan \beta \). Relying on the results of the analysis of the RG flow in the ESSM presented in [17] one can obtain the maximum possible value of the lightest Higgs scalar for each particular choice of \( \tan \beta \).

The dependence of the two-loop upper bound on the mass of the lightest Higgs particle is examined in Fig. 1 where it is compared with the corresponding limits in the MSSM and NMSSM. At moderate values of \( \tan \beta \) (\( \tan \beta = 1.6 - 3.5 \)) the upper limit on the lightest Higgs boson mass in the ESSM is considerably higher than in the MSSM and NMSSM. It reaches the maximum value \( \sim 150 \) GeV at \( \tan \beta = 1.5 - 2 \). In the considered part of the parameter space the theoretical restriction on the mass of the lightest CP-even Higgs boson in the NMSSM exceeds the corresponding limit in the MSSM because of the extra contribution to \( m_h^2 \) induced by the additional \( F \)-term in the Higgs scalar potential of the NMSSM. The size of this contribution, which is described by the first term in the square brackets of Eq. (7), is determined by the Yukawa coupling \( \lambda \). The upper limit on the coupling \( \lambda \) caused by the validity of perturbation theory in the NMSSM is more stringent than in the ESSM due to the presence of exotic 5 + \( \overline{5} \)-plets of matter in the particle
spectrum of the ESSM that leads to the growth of \( g_i(Q) \) at high energies. Since the \( g_i(Q) \) increase prevents the appearance of the Landau pole in the evolution of the Yukawa couplings the maximum allowed \( \lambda(m_t) \) value for each \( \tan \beta \) is greater in the ESSM than in the NMSSM. The increase of \( \lambda(m_t) \) is thus accompanied by the growth of the theoretical restriction on the mass of the lightest CP-even Higgs particle. This is the reason why the upper limit of \( m_h \) in the NMSSM is considerably less than in the ESSM at moderate values of \( \tan \beta \).

\[
\begin{array}{c}
\text{m}_h \\
\end{array}
\]

Figure 1: The dependence of the two-loop upper bound on the lightest Higgs boson mass on \( \tan \beta \) for \( m_t(m_t) = 165 \) GeV, \( m_Q^2 = m_U^2 = M_S^2 \), \( X_t = \sqrt{6}M_S \) and \( M_S = 700 \) GeV. The solid, lower and upper dotted lines represent the limit on \( m_h \) in the MSSM, NMSSM and ESSM respectively.

At large \( \tan \beta \gtrsim 10 \) the contribution of the \( F \)-term of the SM-type singlet field to \( m_h^2 \) vanishes. Therefore with increasing \( \tan \beta \) the upper bound on the lightest Higgs boson mass in the NMSSM approaches the corresponding limit in the MSSM. In the ESSM the theoretical restriction on the mass of the lightest Higgs scalar also diminishes when \( \tan \beta \) rises. But even at very large values of \( \tan \beta \) the upper limit on \( m_h \) in the ESSM is still 4 – 5 GeV larger than the ones in the MSSM and NMSSM because of the \( U(1)_N \) \( D \)-term contribution to \( m_h \) (last term in the square brackets of Eq. (7)).

Note that the quoted upper limits for the ESSM, as well as the MSSM and NMSSM, are sensitive to the value of the top-quark mass, the SUSY breaking scale and also depend on the precise form of the two-loop approximations used. Here we have used an analytic approximation of the two-loop effects which slightly underestimates the full two-loop corrections. The upper bounds quoted here may therefore be further increased by several
GeV. However the main point we wish to make is that the upper bound on the lightest CP-even Higgs scalar in the ESSM is always significantly larger than in the MSSM as well as the NMSSM.

4. \( Z' \) and exotica phenomenology at the LHC

The presence of a relatively light \( Z' \) or of exotic multiplets of matter permits to distinguish the ESSM from the MSSM or NMSSM. Collider experiments \[18\] and precision EW tests \[19\] imply that the \( Z' \) is typically heavier than \( 500 - 600 \) GeV while the mixing angle between \( Z \) and \( Z' \) is smaller than a few \( \times 10^{-3} \). The analysis performed in \[20\] revealed that a \( Z' \) boson in \( E_6 \) inspired models can be discovered at the LHC if its mass is less than \( 4 - 4.5 \) TeV. At the same time the determination of the \( Z' \) couplings should be possible up to \( M_{Z'} \sim 2 - 2.5 \) TeV using Drell-Yan (DY) production \[21\].

Fig. 2 shows the differential distribution in invariant mass of the lepton pair \( l^+l^- \) (for one species of lepton \( l = e, \mu \)) in DY production at the LHC with and without light exotic quarks with representative masses \( \mu_{D_i} = 250 \) GeV for all three generations and with \( M_{Z'} = 1.2 \) TeV. This distribution is promptly measurable at the CERN collider with a high resolution and would enable one to not only confirm the existence of a \( Z' \) state but also to establish the possible presence of additional exotic matter, by simply fitting to the data the width of the \( Z' \) resonance \[21\]. In order to perform such an exercise, the \( Z' \) couplings to ordinary matter ought to have been previously established elsewhere, as a modification of the latter may well lead to effects similar to those induced by our exotic matter.

The exotic quarks can also be pair produced directly and decay with novel signatures. In the ESSM the exotic quarks and squarks receive their masses from the large VEV of the singlet \( S \), according to the terms \( \kappa_i S(D_i \bar{D}_i) \) in Eq. (6). Their couplings to the quarks and leptons of the first and second generation should be rather small, as previously discussed in sect.2. The exotic quarks can be relatively light in the ESSM since their masses are set by the Yukawa couplings \( \kappa_i \) and \( \lambda_i \) that may be small. This happens, for example, when the Yukawa couplings of the exotic particles have hierarchical structure which is similar to the one observed in the ordinary quark and lepton sectors. Since the exotic squarks also receive soft masses from SUSY breaking, they are expected to be much heavier, and their production cross-sections will be considerably smaller.

If exotic quarks of the nature described here do exist at low scales, they could possibly be accessed through direct pair hadroproduction at the LHC. Fig. 3 shows the production cross section of exotic quark pairs at the LHC as a function of the invariant mass of the final state. The lifetime and decay modes of these particles are determined by the
Figure 2: Differential cross section in the final state invariant mass, denoted by \( M_{l^+l^-} \), at the LHC for DY production \((l = e \) or \( \mu \) only) in presence of a \( Z' \) with and without the (separate) contribution of exotic \( D \)-quarks with \( \mu_{Di} = 250 \) GeV for \( M_{Z'} = 1.2 \) TeV.

Figure 3: Cross section at the LHC for pair production of exotic \( D \)-quarks as a function of the invariant mass of \( DD \) pair. Similar cross sections of \( tt \) and \( bb \) production are also included for comparison.

operators that break the \( Z_2^H \) symmetry. If \( Z_2^H \) is only slightly broken then exotic quarks may live for a long time. Then they will form bound states with ordinary quarks. It means that at the LHC it may be possible to study the spectroscopy of new composite
scalar leptons or baryons. When \( Z_2^H \) is broken significantly exotic quarks can also produce a remarkable signature. Since according to our initial assumptions the \( Z_2^H \) symmetry is mostly broken by the operators involving quarks and leptons of the third generation the exotic quarks decay either via \( \overline{D} \to t + \bar{b}, \overline{D} \to b + \bar{t} \), if the exotic quarks \( \overline{D} \) are diquarks or via \( D \to t + \bar{t}, D \to \tau + \bar{\nu}_\tau, D \to \nu_\tau + \bar{b} \) if they are leptoquarks. Because in general sfermions decay into corresponding fermion and neutralino one can expect that each diquark will decay further into \( t \)- and \( b \)-quarks while a leptoquark will produce a \( t \)-quark and a \( \tau \)-lepton in the final state with rather high probability.

As each \( t \)-quark decays into a \( b \)-quark whereas a \( \tau \)-lepton gives one charged lepton \( l \) in the final state with a probability of 35%, both these scenarios would generate an excess in the \( b \)-quark production cross section. In this respect SM data samples which should be altered by the presence of exotic \( D \)-quarks are those involving \( t \bar{t} \) production and decay as well as direct \( b \bar{b} \) production. For this reason, Fig. 3 shows the cross sections for these two genuine SM processes alongside those for the exotica. Detailed LHC analyses will be required to establish the feasibility of extracting the excess due to the light exotic particles predicted by our model. However, Fig. 3 should clearly make the point that - for the discussed parameter configuration - one is in a favorable position in this respect, as the decay BRs of the exotic objects are much larger than the expected four-body cross section involving heavy quarks and/or leptons, as discussed in [7]. Thus the presence of light exotic quarks in the particle spectrum could result in an appreciable enhancement of the cross section of either \( pp \to t\overline{t}b\overline{b} + X \) and \( pp \to b\overline{b}b\overline{b} + X \) if exotic quarks are diquarks or \( pp \to t\overline{t}l^+l^- + X \) and \( pp \to b\overline{b}l^+l^- + X \) if new quark states are leptoquarks.

5. Conclusions

We have considered the ESSM defined by a \( U(1)_N \) gauge group extension of the SM under which right-handed neutrinos are neutral, with anomalies cancelled by the low energy matter content of three 27 representations of an \( E_6 \) gauge group broken at the GUT scale. In the ESSM the \( \mu \) problem is solved in a similar way to the NMSSM, but without the accompanying problems of singlet tadpoles or domain walls. The ESSM allows the conventional see-saw mechanism, and gauge coupling unification due to an additional pair of Higgs-like doublets.

In this letter we have focussed on some interesting phenomenological aspects of the model, with the full details given in an accompanying longer paper [7]. If the \( Z' \) boson and exotic quarks are light enough they will be visible at the LHC, providing spectacular new physics signals. For example, the three extra families of exotic charge 1/3 quarks, which must be either diquarks or leptoquarks to ensure baryon number conservation, may be
directly produced in pairs with decay signatures determined by the hierarchical structure of the Yukawa interactions in the ESSM. The exotic quarks may decay into either a heavy quark-antiquark pair $Qar{Q}$ or a heavy quark and lepton pair $Q\tau(\nu_{\tau})$ where $Q$ is either $b$ or $t$ quark. This would result in the growth of the cross section of either $pp \rightarrow QQ\bar{Q} + X$ or $pp \rightarrow l^+l^-Q\bar{Q} + X$ at the LHC.

We have also studied the $Z'$ which may be observed directly as a resonance in the di-lepton invariant mass of DY events, and shown that the exotic quarks can make their presence felt indirectly via a visible increase in the $Z'$ width. The extra matter also has the indirect effect of increasing the lightest CP-even Higgs boson mass. The point is that extra exotic matter predicted by the ESSM changes the RG flow of the gauge and Yukawa couplings relaxing the restrictions on the Yukawa couplings coming from the validity of perturbation theory up to the scale $M_X$ as compared with the MSSM and NMSSM. Larger values of Yukawa couplings at the EW scale and the extra $D$-term contribution in the Higgs potential both serve to increase the upper limit on the mass of the lightest CP-even Higgs boson. We have shown that, in the leading two-loop approximation, the lightest Higgs mass does not exceed about $150 - 155$ GeV which is much higher than in the MSSM or NMSSM.

The discovery of the $Z'$ and exotic quarks would provide a smoking gun signature of the ESSM, providing circumstantial evidence for an underlying $E_6$ gauge structure at the GUT scale, and a window into string theory.

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References

[1] D.J.H. Chung, L.L. Everett, G.L. Kane, S.F. King, J.D. Lykken, L.T. Wang, Phys. Rept. 407 (2005) 1.

[2] M.B.Green, J.H.Schwarz, E.Witten, ‘Superstring Theory’, Cambridge Univ. Press, 1987.

[3] F. del Aguila, G.A. Blair, M. Daniel, G.G. Ross, Nucl. Phys. B 272 (1986) 413.

[4] Y. Hosotani, Phys. Lett. B 129 (1983) 193.

[5] J.L. Hewett, T.G. Rizzo, Phys. Rept. 183 (1989) 193; M.Cvetič, P. Langacker, Phys. Rev. D 54 (1996) 3570; Mod. Phys. Lett. A 11 (1996) 1247; P. Langacker, J. Wang, Phys. Rev. D 58 (1998) 115010.

[6] S.F. King, Rept. Prog. Phys. 67 (2004) 107.

[7] S.F. King, S. Moretti, R. Nevzorov, arXiv:hep-ph/0510419
[8] T. Hambye, E. Ma, M. Raidal, U. Sarkar, Phys. Lett. B 512 (2001) 373.

[9] E. Ma, M. Raidal, J. Phys. G 28 (2002) 95.

[10] E. Ma, Phys. Lett. B 380 (1996) 286.

[11] E. Keith, E. Ma, Phys. Rev. D 54 (1996) 3587; D. Suematsu, Phys. Rev. D 57 (1998) 1738.

[12] E. Keith, E. Ma, Phys. Rev. D 56 (1997) 7155.

[13] Y. Daikoku, D. Suematsu, Phys. Rev. D 62 (2000) 095006.

[14] J. Rich, M. Spiro, J. Lloyd-Owen, Phys. Rept. 151 (1987) 239; P.F. Smith, Contemp. Phys. 29 (1988) 159; T.K. Hemmick et al., Phys. Rev. D 41 (1990) 2074.

[15] L.E. Ibáñez, G.G. Ross, Phys. Lett. B 110 (1982) 215; J. Ellis, D.V. Nanopoulos, K. Tamvakis, Phys. Lett. B 121 (1983) 123; J. Ellis, J. Hagelin, D.V. Nanopoulos, K. Tamvakis, Phys. Lett. B 125 (1983) 275; L. Alvarez-Gaumé, J. Polchinski, M. Wise, Nucl. Phys. B 221 (1983) 495.

[16] M. Carena, M. Quiros, C.E.M. Wagner, Nucl. Phys. B 461 (1996) 407.

[17] U. Ellwanger, C. Hugonie, Eur. Phys. J. C 25 (2002) 297.

[18] F. Abe et al., Phys. Rev. Lett. 79 (1997) 2192; P. Abreu et al., Phys. Lett. B 485 (2000) 45; R. Barate et al., Eur. Phys. J. C 12 (2000) 183.

[19] J. Erler, P. Langacker, Phys. Lett. B 456 (1999) 68.

[20] M. Cvetič, S. Godfrey, hep-ph/9504216; A. Leike, Phys. Rept. 317 (1999) 143; J. Kang, P. Langacker, Phys. Rev. D 71 (2005) 035014.

[21] M. Dittmar, A. Djouadi, A.-S. Nicollerat, Phys. Lett. B 583 (2004) 111.