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LHC signatures and cosmological implications of the E_6 inspired SUSY models

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The phenomenological implications of the E_6 inspired supersymmetric models based on the Standard Model gauge group together with extra $U(1)_N$ gauge symmetry under which right–handed neutrinos have zero charge are examined. In these models single discrete $\tilde{Z}_2^H$ symmetry forbids the tree-level flavour changing processes and the most dangerous operators that violate baryon and lepton numbers. The two–loop renormalisation group flow of the gauge and Yukawa couplings is explored and the qualitative pattern of the Higgs spectrum in the case of the quasi-fixed point scenario is discussed. These E_6 inspired models contain two dark-matter candidates. The presence of exotic states in these models gives rise to the nonstandard decays of the lightest Higgs boson which are also considered.

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1. E_6 inspired SUSY models with extra $U(1)_N$ gauge symmetry

It is well known that Supersymmetric (SUSY) Grand Unified Theories (GUTs) can lead to the $U(1)$ extensions of the Minimal Supersymmetric Standard Model (MSSM), i.e. SUSY models based on the Standard Model (SM) gauge group together with extra $U(1)$ gauge symmetry. In particular, near the GUT scale E_6 can be broken to $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N \times Z_2^M$ group [1] where $Z_2^M = (-1)^{3(B-L)}$ is a matter parity and

$$U(1)_N = \frac{1}{4}U(1)_\chi + \frac{\sqrt{15}}{4}U(1)_\psi\,. \tag{1.1}$$

Two $U(1)_\psi$ and $U(1)_\chi$ symmetries can originate from breakings $E_6 \to SO(10)\times U(1)_\psi$, $SO(10) \to SU(5)\times U(1)_\chi$ [2]. To ensure anomaly cancellation the low energy matter content of the E_6 inspired SUSY models with extra $U(1)_N$ gauge symmetry is extended to fill out three complete 27 representations of E_6 [2]–[3]. Each 27_i multiplet contains SM family of quarks and leptons, right–handed neutrino N_i^c, SM singlet field S_i which carry non–zero $U(1)_N$ charge, a pair of $SU(2)_W$–doublets H_i^d and H_i^u, which have the quantum numbers of Higgs doublets, and charged $\pm 1/3$ coloured triplets of exotic quarks D_i and $\bar{D}_i$. In addition to the complete 27_i multiplets the low energy particle spectrum is supplemented by $SU(2)_W$ doublets L_4 and $\overline{L}_4$ from extra $27'$ and $\overline{27'}$ to preserve gauge coupling unification [4]. Since in these models N_i^c do not participate in the gauge interactions these states are expected to gain masses at some intermediate scale, shedding light on the origin of the mass hierarchy in the lepton sector and providing a mechanism for the generation of the baryon asymmetry in the Universe via leptogenesis [5]. The remaining matter supermultiplets survive down to the TeV scale. Different phenomenological implications of the E_6 inspired SUSY models with extra $U(1)_N$ gauge symmetry were considered in [1]–[12]. Recently the particle spectrum and the corresponding collider signatures were analysed within the constrained version of this $U(1)_N$ SUSY extension of the SM [13]–[16].

The presence of exotic matter in the E_6 inspired SUSY models lead to non–diagonal flavour transitions and rapid proton decay. In order to suppress flavour changing processes as well as baryon and lepton number violating operators one can impose a $\tilde{Z}_2^H$ symmetry. Under this symmetry all superfields except L_4, $\overline{L}_4$, one pair of H_i^u and H_i^d (i.e. H_u and H_d) and one of the SM-type singlet superfields S_i (i.e. S) are odd. The $\tilde{Z}_2^H$ symmetry reduces the structure of the Yukawa interactions to

$$\begin{aligned} W = {} & \lambda S(H_u H_d) + \lambda_{\alpha\beta} S(H_\alpha^d H_\beta^u) + \kappa_{ij} S(D_i \overline{D}_j) + \tilde{f}_{\alpha\beta} S_\alpha (H_\beta^d H_u) + f_{\alpha\beta} S_\alpha (H_d H_\beta^u) \\ & + g_{ij}(Q_i L_4)\overline{D}_j + h_{i\alpha} e_i^c (H_\alpha^d L_4) + \mu_L L_4 \overline{L}_4 + W_{MSSM}(\mu = 0)\,, \end{aligned} \tag{1.2}$$

where $\alpha, \beta = 1,2$ and $i, j = 1,2,3$. At low energies the superfields H_u, H_d and S play the role of Higgs fields. The vacuum expectation value (VEV) of $\langle S \rangle = s/\sqrt{2}$ breaks the extra $U(1)_N$ symmetry providing an effective μ term as well as the masses of the exotic fermions and the Z' boson ($M_{Z'}$). The VEVs of the $SU(2)_W$ doublets $\langle H_d \rangle = v_1/\sqrt{2}$ and $\langle H_u \rangle = v_2/\sqrt{2}$ result in the electroweak (EW) symmetry breaking (EWSB), inducing the masses of quarks and leptons.

2. Quasi-fixed points and spectrum of Higgs bosons

The superpotential (1.2) involves a lot of new Yukawa couplings. To simplify our analysis of the renormalisation group (RG) flow of the gauge and Yukawa couplings we assume that all

Yukawa couplings except λ and the top–quark Yukawa coupling h_t are sufficiently small and can be neglected in the leading approximation. Then the superpotential (1.2) reduces to

$$W \approx \lambda S(H_d H_u) + h_t(H_u Q_3)u_3^c\,. \tag{2.1}$$

For the purposes of RG analysis, it is convenient to introduce $\rho_t = h_t^2/g_3^2$ and $\rho_\lambda = \lambda^2/g_3^2$, where g_3 is a strong gauge coupling.

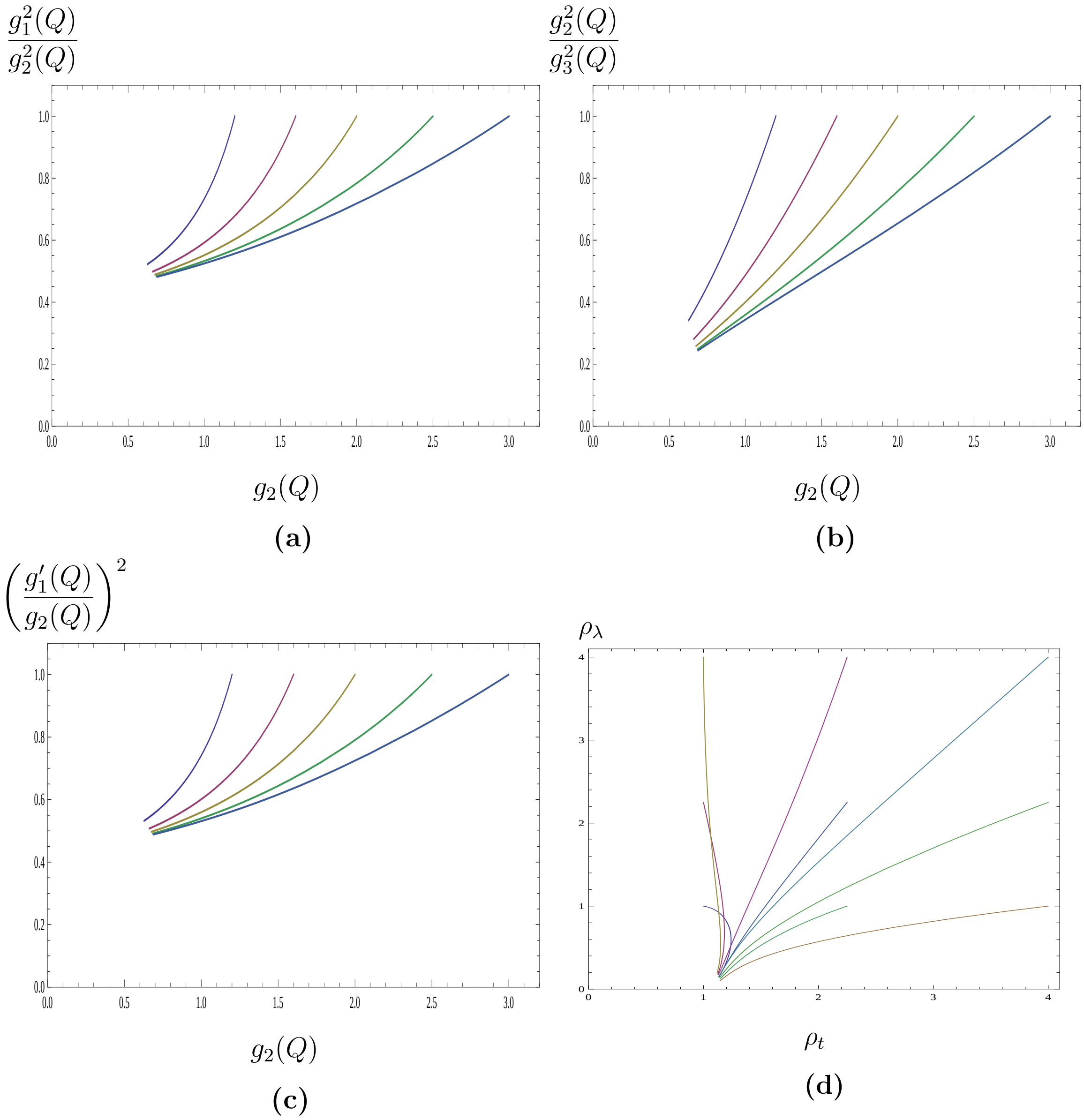

Figure 1: Two–loop RG flow of gauge and Yukawa couplings from $Q = M_X$ to EW scale: *(a)-(c)* evolution of gauge couplings for $g_1(M_X) = g_1'(M_X) = g_2(M_X) = g_3(M_X) = h_t(M_X) = \lambda(M_X) = g_0$, and different values of g_0; *(d)* running of Yukawa couplings in the $\rho_\lambda - \rho_t$ plane for $g_0 = 1.5$.

In principle, the RG flow of the gauge couplings is affected by the kinetic term mixing. However if the corresponding mixing is small at the GUT scale M_X then it remains small at any scale below M_X. Thus in the leading approximation the kinetic term mixing can be ignored. Our analysis revealed that the solutions of the two–loop RG equations (RGEs) for the $SU(2)_W$, $U(1)_Y$ and $U(1)_N$ gauge couplings (g_2, g_1 and g'_1) are focused in the infrared region near the quasi-fixed points (see Figs. 1a and 1c) which are rather close to the measured values of these couplings at the EW scale. On the other hand from Fig. 1b it follows that the convergence of the solutions for the strong gauge coupling $g_3(Q)$ to the fixed point is quite weak because the corresponding one–loop beta function vanishes. One can show that the values of $g_i(M_X) = g_0 \simeq 1.5$ lead to $g_i(M_Z)$ which are rather close to the measured central values of these couplings at the EW scale.

As $h_t(M_X)$ and $\lambda(M_X)$ grow, the region at the EW scale in which the solutions of the RGEs for ρ_t and ρ_λ are concentrated shrinks drastically. Thus the corresponding solutions are focused near the quasi–fixed point. This follows from Fig. 1d. In this case the quasi–fixed point is an intersection point of the invariant and quasi–fixed lines [17]–[18]. In the two–loop approximation the coordinates of the quasi–fixed point are $\rho_t = 1.16$ and $\rho_\lambda = 0.14$. The quasi–fixed point solution corresponds to $\tan\beta = \frac{v_2}{v_1} \simeq 1$. Our estimations show that for $1.5 \lesssim h_t(M_X), \lambda(M_X) \lesssim 3$ two–loop upper bound on the lightest Higgs boson mass varies between $120-127\,\mathrm{GeV}$ [9].

It is worth noting that for $\tan\beta \simeq 1$ the solutions with $125-126\,\mathrm{GeV}$ SM–like Higgs boson can be obtained only if $\lambda \gtrsim g'_1$ at the EW scale which leads to extremely hierarchical structure of the Higgs spectrum [2]. In this case the qualitative pattern of the Higgs spectrum is rather similar to the one that arises in the Peccei-Quinn (PQ) symmetric NMSSM in which the heaviest CP–even, CP–odd and charged states are almost degenerate and much heavier than the lightest and second lightest CP-even Higgs bosons [19]. Since the second lightest CP–even Higgs state and Z' boson have approximately the same masses and $M_{Z'} \gtrsim 2.5\,\mathrm{TeV}$ the heaviest Higgs boson masses lie beyond the multi TeV range and the mass matrix of the CP–even Higgs sector can be diagonalized using the perturbation theory [19]–[22]. In this limit the lightest CP–even Higgs boson is the analogue of the SM Higgs field. All other Higgs states are so heavy that they can not be discovered at the LHC. Extremely hierarchical structure of the Higgs spectrum also implies that all phenomenologically viable scenarios associated with the quasi–fixed point are very fine–tuned.

3. Dark matter and exotic Higgs decays

The fermionic components of the Higgs–like and SM singlet superfields, which are $\tilde{Z}_2^H$ odd, form a set of inert neutralino and chargino states. Using the method proposed in [23] one can show that there are theoretical upper bounds on the masses of the lightest and second lightest inert neutralino states ($\tilde{H}_1^0$ and $\tilde{H}_2^0$) [7]–[8]. Their masses do not exceed $60-65\,\mathrm{GeV}$ [7]. Therefore these states, which are predominantly inert singlinos, tend to be the lightest and next–to–lightest SUSY particles (LSP and NLSP). In the simplest phenomenologically viable scenarios LSP is considerably lighter than $1\,\mathrm{eV}$ and form hot dark matter in the Universe. The existence of very light neutral fermions in the particle spectrum may lead to some interesting implications for the neutrino physics (see, for example [24]). Because LSP is so light it gives only minor contribution to the dark matter density. At the same time the conservation of the Z_2^M and $\tilde{Z}_2^H$ symmetries ensures that the lightest ordinary neutralino is also stable and may account for all or some of the observed cold dark matter density.

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The NLSP with the GeV scale mass gives rise to the exotic decays of the SM–like Higgs boson $h_1 \to \tilde{H}_2^0 \tilde{H}_2^0$. The couplings of the lightest Higgs state to the LSP and NLSP are determined by their masses. Because $\tilde{H}_1^0$ is extremely light it does not affect Higgs phenomenology. On the other hand the branching ratio of the nonstandard decays of the SM–like Higgs boson into a pair of the NLSP states, i.e. $h_1 \to \tilde{H}_2^0 \tilde{H}_2^0$, can be substantial if NLSP has a mass $m_{\tilde{H}_2^0}$ of order of the b–quark mass m_b. Nonetheless the couplings of $\tilde{H}_1^0$ and $\tilde{H}_2^0$ to the Z–boson and other SM particles can be negligibly small because of the inert singlino admixture in these states [10]. As a consequence these states could escape detection at former and present experiments.

After being produced the NLSP sequentially decay into the LSP and fermion–antifermion pairs via virtual Z. Thus the exotic decays of the lightest CP–even Higgs state discussed above lead to two fermion–antifermion pairs and missing energy in the final state. However since $\tilde{H}_2^0$ tend to be longlived particle it decays outside the detectors resulting in the invisible decays of h_1. If $m_{\tilde{H}_2^0} \gg m_b(m_{h_1})$ the lightest Higgs boson decays mainly into $\tilde{H}_2^0 \tilde{H}_2^0$ resulting in the strong suppression of the branching ratios for the decays of h_1 into SM particles. To avoid such suppression we restrict our consideration to the GeV scale masses of $\tilde{H}_2^0$. In our analysis we require that the NLSP decays before BBN, i.e. its lifetime is shorter than 1 sec. This requirement rules out too light $\tilde{H}_2^0$ because $\tau_{\tilde{H}_2^0} \sim 1/(m^5_{\tilde{H}_2^0})$. One can easily find that it is rather problematic to satisfy this restriction for $m_{\tilde{H}_2^0} \lesssim 100\,\mathrm{MeV}$. The numerical analysis indicates that the branching ratio associated with the decays $h_1 \to \tilde{H}_2^0 \tilde{H}_2^0$ can be as large as 20-30% if $\tilde{H}_2^0$ is heavier than 2.5 GeV [10]. When $\tilde{H}_2^0$ is lighter than 0.5 GeV this branching ratio can be as small as $10^{-3} - 10^{-4}$ [10].

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