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LHC signatures and cosmological implications of the E_6 inspired SUSY models

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SUSY models with extra $U(1)_N$ symmetry

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 where $Z_2^M = (-1)^{3(B-L)}$ is a matter parity and

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

Two $U(1)_\psi$ and $U(1)_\chi$ symmetries can originate from breakings $E_6 \to SO(10) \times U(1)_\psi$, $SO(10) \to SU(5) \times U(1)_\chi$ [1].

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. 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. Since in these models N_i^c do not participate in the gauge interactions they are expected to gain masses at some intermediate scale, while the remaining matter survives down to the TeV scale.

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:

$$\begin{aligned} W = {} & \lambda S(H_u H_d) + \lambda_{\alpha\beta} S(H_\alpha^d H_\beta^u) + \kappa_{ij} S(D_i \overline{D}_j) \\ & + \tilde{f}_{\alpha\beta} S_\alpha (H_\beta^d H_u) + f_{\alpha\beta} S_\alpha (H_d H_\beta^u) + g_{ij}(Q_i L_4)\overline{D}_j \\ & + h_{i\alpha} e_i^c (H_\alpha^d L_4) + \mu_L L_4 \overline{L}_4 + W_{MSSM}(\mu = 0)\,, \end{aligned} \qquad (2)$$

where $\alpha = 1, 2$ and $i = 1, 2, 3$. At low energies the superfields H_u, H_d and S play the role of Higgs fields. The gauge group and field content of these SUSY models can originate from the orbifold GUT models [2].

Quasi-fixed points and the Higgs mass

Our analysis revealed that the solutions of the two–loop renormalization group (RG) equations 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 electroweak (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 rather weak because the corresponding one–loop beta function vanishes. The values of $g_i(M_X) = g_0$ around 1.5 lead to $g_i(M_Z)$ which are quite close to the measured central values of these couplings at the EW scale.

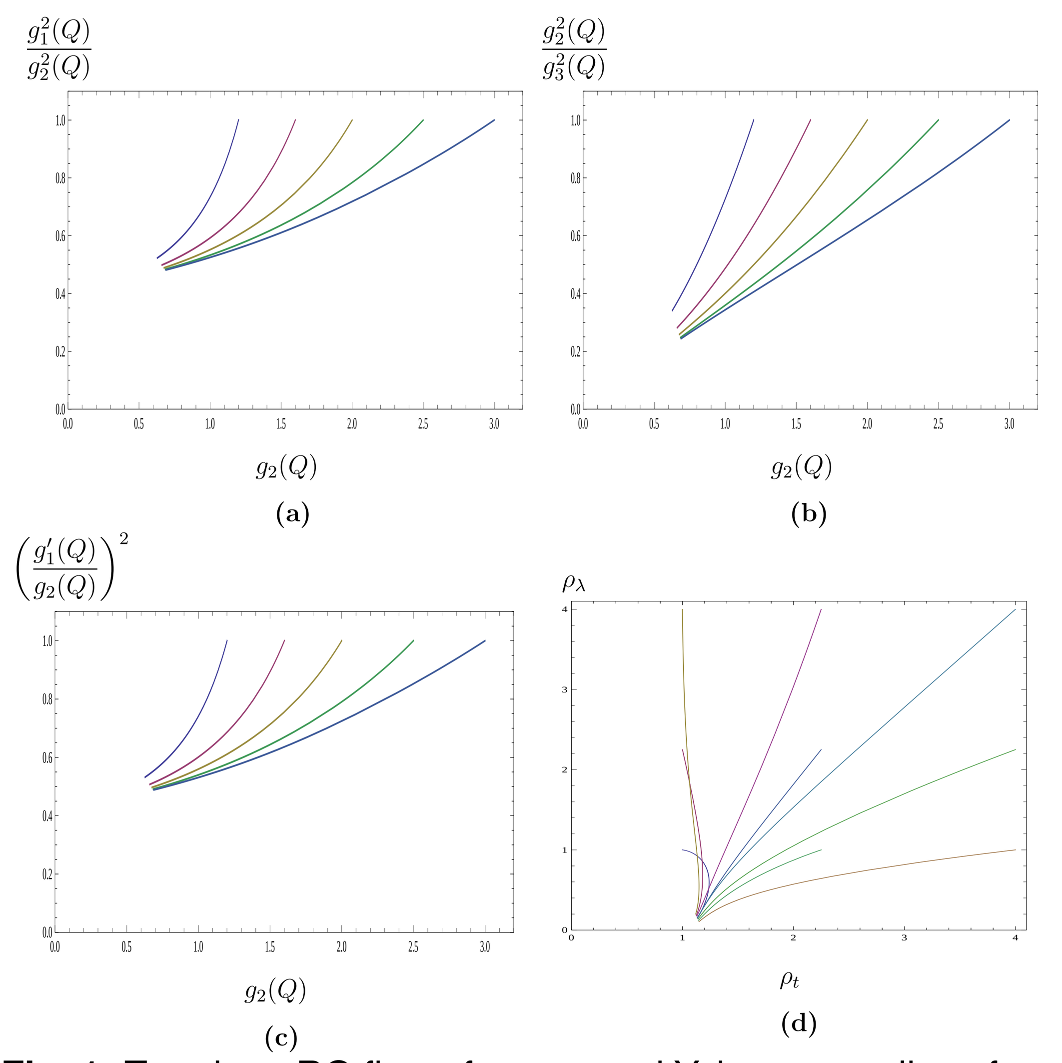

Fig. 1: Two–loop RG flow of gauge and Yukawa couplings from $Q = M_X$ to EW scale: *(a)-(c)* evolution of gauge couplings for $g_1(M_X) = g_1'(M_X) = g_2(M_X) = g_3(M_X) = h_t(M_X) = \lambda(M_X) = g_0$, $g_{11}(M_X) = 0$ and different values of g_0; *(d)* running of Yukawa couplings in the $\rho_\lambda - \rho_t$ plane for $g_0 = 1.5$.

To simplify the analysis of the RG flow of the Yukawa couplings we assumed that λ and the top–quark Yukawa coupling h_t are substantially larger than all other Yukawa couplings. For the purposes of RG studies, it is convenient to introduce $\rho_t = h_t^2/g_3^2$ and $\rho_\lambda = \lambda^2/g_3^2$. As one can see from Fig. 1d the solutions of the two–loop RG equations for the Yukawa couplings are concentrated near the quasi–fixed points when $h_t(M_X)$ and $\lambda(M_X)$ grow. For $1.5 \lesssim h_t(M_X)$, $\lambda(M_X) \lesssim 3$ two–loop upper bound on the lightest Higgs mass varies between $120 - 127$ GeV [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, compose a set of inert neutralino and chargino states. The lightest and second lightest inert neutralinos ($\tilde{H}_1^0$ and $\tilde{H}_2^0$), which are predominantly inert singlinos, tend to be LSP and NLSP. In the simplest phenomenologically viable scenarios LSP is expected to be substantially lighter than $1\,$eV and form hot dark matter in the Universe. Since LSP is so light it gives only minor contribution to the dark matter density. Because of the conservation of the Z_2^M and $\tilde{Z}_2^H$ symmetries the lightest ordinary neutralino can be also absolutely stable and may account for all or some of the observed cold dark matter density.

The NLSP with the GeV scale masses gives rise to the exotic decays of the lightest Higgs boson, i.e $h_1 \to \tilde{H}_2^0 \tilde{H}_2^0$. After being produced the NLSP sequentially decay into the LSP and fermion–antifermion pairs via virtual Z. Since $\tilde{H}_2^0$ tend to be longlived particle it decays outside the detectors resulting in the invisible decays of h_1. In our analysis we require that the NLSP decays before BBN, i.e. its lifetime is shorter than $1\,$sec. This requirement rules out $\tilde{H}_2^0$ with mass below $100\,$MeV. 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 [4]. When $\tilde{H}_2^0$ is lighter than 0.5 GeV this branching ratio can be as small as $10^{-3} - 10^{-4}$ [4].

References

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