

Theoretical aspects of electroweak symmetry breaking in SUSY models

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The electroweak (EW) symmetry breaking in the simplest supersymmetric (SUSY) extensions of the standard model (SM), i.e. minimal and next-to-minimal supersymmetric standard models (MSSM and NMSSM), is considered. The spectrum of the Higgs particles, upper bound on the mass of the lightest Higgs boson and little hierarchy problem are discussed. The breakdown of gauge symmetry and Higgs phenomenology within the E_6 inspired SUSY models with extra $U(1)'$ factor are briefly reviewed.

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1. Introduction

Nowadays, there are serious reasons to believe that new particles and interactions should exist at the TeV scale. We expect that LHC might be able to shed light on the physics beyond the Standard Model (SM), origin of dark matter and the mechanism of electroweak (EW) symmetry breaking in the near future. At the moment the supersymmetric (SUSY) models are the best motivated extensions of the SM. SUSY models are defined by the field content, structure of gauge interactions and superpotential. In order to give both up- and down-type fermions a mass and to ensure anomaly cancellation the Higgs sector of the minimal supersymmetric standard model (MSSM) includes two Higgs doublets (H_1 and H_2) with the opposite hypercharge. Also in SUSY models R-parity is normally imposed ($R = (-1)^{3(B-L)+2S}$). It forbids the baryon and lepton number violating operators that lead to rapid proton decay. R-parity ensures that the lightest supersymmetric particle (LSP) is absolutely stable and can play the role of dark matter. The most general renormalizable superpotential of the MSSM can be presented in the following form

$$W_{MSSM} = \varepsilon_{ij}(y_{ab}^U Q_a^j u_b^c H_2^i + y_{ab}^D Q_a^j d_b^c H_1^i + y_{ab}^L L_a^j e_b^c H_1^i + \mu H_1^i H_2^j), \quad (1.1)$$

where $i, j = 1, 2, 3$ are the $SU(2)$ and $a, b = 1, 2, 3$ are the generation indices. In Eq. (1.1) colour indices are suppressed.

Since SUSY particles are considerably heavier than observed quarks and leptons supersymmetry must be broken. Thus the Lagrangian of SUSY models can be written as

$$\mathcal{L} = \mathcal{L}_{SUSY} + \mathcal{L}_{soft}. \quad (1.2)$$

where \mathcal{L}_{soft} contains soft SUSY breaking terms that break supersymmetry but do not induce quadratic divergences that get cancelled automatically within the models based on the exact global supersymmetry. In the MSSM the set of the soft SUSY breaking terms includes

$$-\mathcal{L}_{soft} = \sum_i m_i^2 |\varphi_i|^2 + (\frac{1}{2} \sum_\alpha M_\alpha \tilde{\lambda}_\alpha \tilde{\lambda}_\alpha + \sum_{a,b} [A_{ab}^U y_{ab}^U \tilde{Q}_a \tilde{u}_b^c H_2 + A_{ab}^D y_{ab}^D \tilde{Q}_a \tilde{d}_b^c H_1 + A_{ab}^L y_{ab}^L \tilde{L}_a \tilde{e}_b^c H_1] + B\mu H_1 H_2 + h.c.), \quad (1.3)$$

where φ_i are scalar components of chiral superfields while $\tilde{\lambda}_\alpha$ are gaugino fields. To avoid fine-tuning SUSY breaking mass parameters are expected to be in the TeV range.

2. EW symmetry breaking in the MSSM

Including soft SUSY breaking terms and radiative corrections, the Higgs effective potential in the MSSM can be written as

$$V = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_3^2 (H_1 H_2 + h.c.) + \frac{g_2^2}{8} (H_1^+ \sigma_a H_1 + H_2^+ \sigma_a H_2)^2 + \frac{g'^2}{8} (|H_1|^2 - |H_2|^2)^2 + \Delta V, \quad (2.1)$$

where $g' = \sqrt{3/5}g_1$, g_2 and g_1 are the low energy (GUT normalised) $SU(2)_W$ and $U(1)_Y$ gauge couplings, $m_1^2 = m_{H_1}^2 + \mu^2$, $m_2^2 = m_{H_2}^2 + \mu^2$ and $m_3^2 = -B\mu$. In Eq. (2.1) ΔV represents the contribution of loop corrections to the Higgs effective potential.

At the physical minimum of the scalar potential (2.1) the Higgs fields develop vacuum expectation values (VEVs)

$$\langle H_1 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \quad \langle H_2 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_2 \end{pmatrix}. \quad (2.2)$$

breaking the $SU(2)_W \times U(1)_Y$ gauge symmetry to $U(1)_{em}$ associated with electromagnetism and generating the masses of all bosons and fermions. The value of $v = \sqrt{v_1^2 + v_2^2} \simeq 246 \text{ GeV}$ is fixed by the Fermi scale. At the same time the ratio of the Higgs VEVs remains arbitrary. Hence it is convenient to introduce $\tan \beta = v_2/v_1$.

The vacuum configuration (2.2) is not the most general one. Because of the $SU(2)$ invariance of the Higgs potential (2.1) one can always make $\langle H_2^+ \rangle = 0$ by virtue of a suitable gauge rotation. Then the requirement $\langle H_1^- \rangle = 0$, which is a necessary condition to preserve $U(1)_{em}$ in the physical vacuum, is equivalent to requiring the squared mass of the physical charged scalar to be positive. It imposes additional constraints on the parameter space.

At tree-level $\Delta V = 0$ and the Higgs potential is described by the sum of the first five terms in Eq. (2.1). Notice that the Higgs self-interaction couplings in Eq. (2.1) are fixed and determined by the gauge couplings in contrast with the SM. At the tree-level the MSSM Higgs potential contains only three independent parameters: m_1^2, m_2^2, m_3^2 . The stable vacuum of the scalar potential (2.1) exists only if

$$m_1^2 + m_2^2 > 2|m_3|^2. \quad (2.3)$$

This can be easily understood if one notice that in the limit $v_1^2 = v_2^2$ the quartic terms in the Higgs potential vanish. In the considered case the scalar potential (2.1) remains positive definite only if the condition (2.3) is satisfied. Otherwise physical vacuum becomes unstable. On the other hand Higgs doublets acquire non-zero VEVs only when

$$m_1^2 m_2^2 < |m_3|^4. \quad (2.4)$$

Indeed, if $m_1^2 m_2^2 > |m_3|^4$ then all Higgs fields have positive masses for $v_1 = v_2 = 0$ and the breakdown of EW symmetry does not take place. The conditions (2.3) and (2.4) also follow from the equations for the extrema of the Higgs boson potential. At tree-level the minimization conditions in the directions (2.2) in field space read:

$$\begin{aligned} \frac{\partial V}{\partial v_1} &= \left(m_1^2 + \frac{\bar{g}^2}{8}(v_1^2 - v_2^2) \right) v_1 - m_3^2 v_2 = 0, \\ \frac{\partial V}{\partial v_2} &= \left(m_2^2 + \frac{\bar{g}^2}{8}(v_2^2 - v_1^2) \right) v_2 - m_3^2 v_1 = 0, \end{aligned} \quad (2.5)$$

where $\bar{g} = \sqrt{g_2^2 + g'^2}$. Solution of the minimization conditions (2.5) can be written in the following form

$$\sin 2\beta = \frac{2m_3^2}{m_1^2 + m_2^2}, \quad \frac{\bar{g}^2}{4} v^2 = \frac{2(m_1^2 - m_2^2 \tan^2 \beta)}{\tan^2 \beta - 1}. \quad (2.6)$$

Requiring that $v^2 > 0$ and $|\sin 2\beta| < 1$ one can reproduce conditions (2.3) and (2.4).

From Eqs. (2.3)–(2.6) it is easy to see that the appropriate breakdown of the EW symmetry can not be always achieved. In particular the breakdown of $SU(2)_W \times U(1)_Y \rightarrow U(1)_{em}$ does not take place when $m_2^2 = m_1^2$. This is exactly what happens in the constrained version of the MSSM (cMSSM) which implies that $m_i^2(M_X) = m_0^2$, $A_{ab}^k(M_X) = A$, $M_\alpha(M_X) = M_{1/2}$. Thus in the cMSSM $m_2^2(M_X) = m_1^2(M_X)$. Nevertheless the correct pattern of the EW symmetry breaking can be achieved within this SUSY model. Since the top–quark Yukawa coupling is large it affects the renormalisation group (RG) flow of $m_2^2(Q)$ rather strongly resulting in small or even negative values of $m_2^2(Q)$ at low energies that triggers the breakdown of the EW symmetry. This is the so-called radiative mechanism of the EW symmetry breaking (EWSB).

The radiative EWSB demonstrates the importance of loop effects in the considered process. In addition to the RG flow of all couplings one has to take into account loop corrections to the Higgs effective potential which are associated with the last term ΔV in Eq. (2.1). In the simplest SUSY extensions of the SM the dominant contribution to ΔV comes from the loops involving the top–quark and its superpartners because of their large Yukawa coupling h_t . In the one–loop approximation the contribution of the top–quark and its superpartners to ΔV is determined by the masses of the corresponding bosonic and fermionic states, i.e.

$$\Delta V = \frac{3}{32\pi^2} \left[m_{\tilde{t}_1}^4 \left(\ln \frac{m_{\tilde{t}_1}^2}{Q^2} - \frac{3}{2} \right) + m_{\tilde{t}_2}^4 \left(\ln \frac{m_{\tilde{t}_2}^2}{Q^2} - \frac{3}{2} \right) - 2m_t^4 \left(\ln \frac{m_t^2}{Q^2} - \frac{3}{2} \right) \right], \quad (2.7)$$

$$m_{\tilde{t}_{1,2}} = \frac{1}{2} \left(m_Q^2 + m_U^2 + 2m_t^2 \pm \sqrt{(m_Q^2 - m_U^2)^2 + 4m_t^2 X_t^2} \right),$$

where $X_t = A_t - \mu/\tan\beta$ is a stop mixing parameter, A_t is a trilinear scalar coupling associated with the top quark Yukawa coupling and m_t is the running top quark mass

$$m_t(M_t) = \frac{h_t(M_t)}{\sqrt{2}} v \sin\beta.$$

Initially the sector of EWSB involves eight degrees of freedom. However three of them are massless Goldstone modes which are swallowed by the W^\pm and Z gauge bosons. The W^\pm and Z bosons gain masses via the interaction with the neutral components of the Higgs doublets so that

$$M_W = \frac{g^2}{2} v, \quad M_Z = \frac{\bar{g}}{2} v.$$

When CP in the MSSM Higgs sector is conserved the remaining five physical degrees of freedom form two charged, one CP–odd and two CP–even Higgs states. The masses of the charged and CP–odd Higgs bosons are

$$m_A^2 = m_1^2 + m_2^2 + \Delta_A, \quad M_{H^\pm}^2 = m_A^2 + M_W^2 + \Delta_\pm, \quad (2.8)$$

where Δ_\pm and Δ_A are the loop corrections. The CP–even states are mixed and form a 2×2 mass matrix. It is convenient to introduce a new field space basis (h, H) rotated by the angle β with respect to the initial one:

$$\begin{aligned} \text{Re} H_1^0 &= (h \cos\beta - H \sin\beta + v_1)/\sqrt{2}, \\ \text{Re} H_2^0 &= (h \sin\beta + H \cos\beta + v_2)/\sqrt{2}. \end{aligned} \quad (2.9)$$

In this new basis the mass matrix of the Higgs scalars takes the form [1]

$$M^2 = \begin{pmatrix} M_{11}^2 & M_{12}^2 \\ M_{21}^2 & M_{22}^2 \end{pmatrix} = \begin{pmatrix} \frac{\partial^2 V}{\partial v^2} & \frac{1}{v} \frac{\partial^2 V}{\partial v \partial \beta} \\ \frac{1}{v} \frac{\partial^2 V}{\partial v \partial \beta} & \frac{1}{v^2} \frac{\partial^2 V}{\partial \beta^2} \end{pmatrix}, \quad (2.10)$$

$$M_{22}^2 = m_A^2 + M_Z^2 \sin^2 2\beta + \Delta_{22}, \quad M_{12}^2 = M_{21}^2 = -\frac{1}{2} M_Z^2 \sin 4\beta + \Delta_{12}, \quad (2.11)$$

$$M_{11}^2 = M_Z^2 \cos^2 2\beta + \Delta_{11},$$

where Δ_{ij} represent the contributions from loop corrections. In Eqs. (2.11) the equations for the extrema of the Higgs boson effective potential are used to eliminate m_1^2 and m_2^2 .

From Eqs. (2.8) and (2.11) one can see that at tree-level the masses and couplings of the Higgs bosons in the MSSM can be parametrised in terms m_A and $\tan \beta$ only. The masses of the two CP-even eigenstates obtained by diagonalizing the matrix (2.10)–(2.11) are given by

$$m_{h_1, h_2}^2 = \frac{1}{2} \left(M_{11}^2 + M_{22}^2 \mp \sqrt{(M_{22}^2 - M_{11}^2)^2 + 4M_{12}^4} \right). \quad (2.12)$$

The qualitative pattern of the Higgs spectrum depends very strongly on the mass m_A of the pseudoscalar Higgs boson. With increasing m_A the masses of all the Higgs particles grow. At very large values of m_A ($m_A^2 \gg v^2$), the lightest Higgs boson mass approaches its theoretical upper limit $\sqrt{M_{11}^2}$. Thus the top-left entry of the CP-even mass matrix (2.10)–(2.11) represents an upper bound on the lightest Higgs boson mass-squared. In the leading two-loop approximation the mass of the lightest Higgs boson in the MSSM does not exceed 130 – 135 GeV. This is one of the most important predictions of the minimal SUSY model that can be tested at the LHC in the near future.

It is important to study the spectrum of the Higgs bosons together with their couplings to the gauge bosons because LEP set stringent limits on the masses and couplings of the Higgs states. In the rotated field basis (h, H) the trilinear part of the Lagrangian, which determines the interactions of the neutral Higgs states with the Z -boson, is simplified:

$$L_{AZH} = \frac{\bar{g}}{2} M_Z Z_\mu Z_{1\mu} h + \frac{\bar{g}}{2} Z_\mu \left[H(\partial_\mu A) - (\partial_\mu H)A \right]. \quad (2.13)$$

In this basis the SM-like CP-even component h couples to a pair of Z bosons, while the other one H interacts with the pseudoscalar A and Z . The coupling of h to the Z pair is exactly the same as in the SM. The couplings of the Higgs mass eigenstates to a Z pair (g_{ZZh_i} , $i = 1, 2, 3$) and to the Higgs pseudoscalar and Z boson ($g_{ZA h_i}$) appear because of the mixing between h and H .

Following the traditional notations, one can define the normalised R -couplings as: $g_{VVh_i} = R_{VVh_i} \times \text{SM coupling}$ ($V = Z, W^\pm$); $g_{ZA h_i} = \frac{\bar{g}}{2} R_{ZA h_i}$. The absolute values of all these R -couplings vary from zero to unity. The relative couplings R_{ZZh_i} and $R_{ZA h_i}$ are given in terms of the angles α and β :

$$R_{VVh_1} = -R_{ZA h_2} = \sin(\beta - \alpha), \quad R_{VVh_2} = R_{ZA h_1} = \cos(\beta - \alpha), \quad (2.14)$$

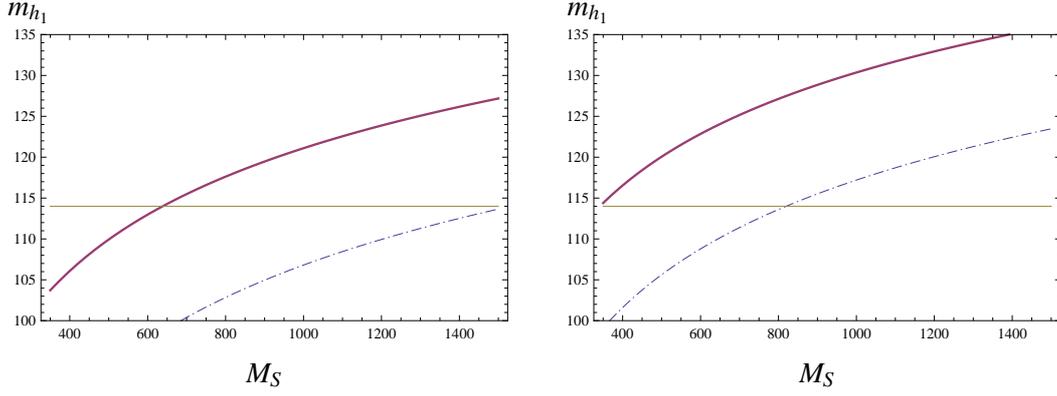


Figure 1: *Left:* The dependence of the one-loop lightest Higgs boson mass on M_S for $\tan\beta = 2$. *Right:* The one-loop mass of the lightest CP-even Higgs state versus M_S for $\tan\beta = 3$. The solid and dashed-dotted lines correspond to $X_t = 2M_S$ and $X_t = M_S$ respectively. The horizontal line represents the current LEP limit.

where the angle α is defined as follows:

$$\begin{aligned} h_1 &= -(\sqrt{2}\operatorname{Re}H_1^0 - v_1)\sin\alpha + (\sqrt{2}\operatorname{Re}H_2^0 - v_2)\cos\alpha, \\ h_2 &= (\sqrt{2}\operatorname{Re}H_1^0 - v_1)\cos\alpha + (\sqrt{2}\operatorname{Re}H_2^0 - v_2)\sin\alpha. \end{aligned} \quad (2.15)$$

From Eqs. (2.14) it becomes clear that in the MSSM the couplings of the lightest Higgs boson to the Z pair can be substantially smaller than in the SM. Therefore the experimental lower bound on the lightest Higgs mass in the MSSM is weaker than in the SM. On the other hand, when $m_A^2 \gg v^2$ the lightest CP-even Higgs state is predominantly the SM-like superposition of the neutral components of Higgs doublets h so that $R_{ZZh_1} \simeq 1$. In this case the lightest Higgs scalar has to satisfy LEP constraint on the mass of the SM-like Higgs boson, i.e. it should be heavier than 114.4 GeV.

Recent studies indicate that in the MSSM the scenarios with the light Higgs pseudoscalar ($m_A \sim 100$ GeV) are almost ruled out by LEP. Since m_A tend to be large the SM-like Higgs boson must be relatively heavy. At the same time the lightest Higgs boson mass in the MSSM does not exceed $M_Z \simeq 91$ GeV at the tree-level. As a consequence in order to satisfy LEP constraints large contribution of loop corrections to the mass of the lightest CP-even Higgs state is required. When SUSY breaking scale M_S is considerably larger than M_Z and $m_Q^2 \simeq m_U^2 \simeq M_S^2$ the contribution of the one-loop corrections to $m_{h_1}^2$ in the leading approximation can be written as

$$\Delta_{11}^{(1)} \simeq \frac{3M_t^4}{2\pi^2 v^2} \left[\frac{X_t^2}{M_S^2} \left(1 - \frac{1}{12} \frac{X_t^2}{M_S^2} \right) + \ln \left(\frac{M_S^2}{m_t^2} \right) \right]. \quad (2.16)$$

The large values of $\Delta_{11}^{(1)} \sim M_Z^2$ can be obtained only if $M_S \gg m_t$ and the ratio $|X_t/M_S|$ is also large. The contribution of the one-loop corrections (2.16) attains its maximal value for $X_t^2 = 6M_S^2$. This is the so-called maximal mixing scenario. In Figs. 1a and 1b the dependence of the one-loop lightest Higgs boson mass on the SUSY breaking scale M_S for $\tan\beta = 2$ and $\tan\beta = 3$ is examined. Two different cases $X_t = M_S$ and $X_t = 2M_S$ are considered. From Figs. 1a and 1b one can see that in order to satisfy LEP constraints M_S should be larger than 400 – 600 GeV. Leading two-loop corrections reduce the SM-like Higgs mass even further. As a result larger values of the SUSY breaking scale are required to overcome LEP limit.

Large values of the soft scalar masses of the superpartners of the top quark ($m_Q^2, m_U^2 \gg M_Z^2$) tend to induce large mass parameter m_2^2 at low energies due to the RG flow. This leads to the fine tuning because m_1^2 and m_2^2 determine the EW scale (see Eqs. (2.6)). Generically the fine tuning which is required to overcome LEP constraints in the MSSM is of the order of 1% (little hierarchy problem). This fine tuning should be compared with the fine tuning in other theories which are used in particle physics. In particular, it is well known that QCD is a highly fine-tuned theory. Indeed, QCD Lagrangian should contain “ θ -term”

$$\mathcal{L}_\theta = \theta_{\text{eff}} \frac{\alpha_s}{8\pi} F^{\mu\nu a} \tilde{F}_{\mu\nu}^a, \quad \theta_{\text{eff}} = \theta + \arg \det M_q, \quad (2.17)$$

where $F^{\mu\nu a}$ is the gluon field strength and $\tilde{F}_{\mu\nu}^a \equiv \frac{1}{2} \epsilon_{\mu\nu\rho\sigma} F^{\rho\sigma a}$ is its dual. This term is not forbidden by the gauge invariance. On the other hand the parameter θ must be extremely small, i.e. $|\theta_{\text{eff}}| \lesssim 10^{-9}$ (strong CP problem). Otherwise it results in too large value of the neutron electric dipole moment. Eqs. (2.17) demonstrate that so small value of θ_{eff} implies enormous fine tuning which is much higher than in the MSSM.

The little hierarchy problem can be solved within SUSY models that allow to get relatively large mass of the SM-like Higgs boson ($m_{h_1} \gtrsim 100 - 110 \text{ GeV}$) at the tree-level. Alternatively, one can try to avoid stringent LEP constraints by allowing exotic decays of the lightest Higgs particle. If usual branching ratios of the lightest Higgs state are dramatically reduced then the lower LEP bound on the SM-like Higgs mass may become inapplicable. In this case the lightest Higgs boson can be still relatively light so that large contribution of loop corrections is not required. Both possibilities mentioned above imply the presence of new particles and interactions. These new particles and interactions can be also used to solve the so-called μ problem. This problem arises when MSSM gets incorporated into supergravity and/or GUT models. Within these models the parameter μ is expected to be either zero or of the order of Planck/GUT scale. At the same time in order to provide the correct pattern of the EWSB μ is required to be of the order of the EW scale.

3. Higgs sector of the NMSSM

In the simplest extension of the MSSM, the Next-to-Minimal Supersymmetric Standard Model (NMSSM), the superpotential is invariant with respect to the discrete transformations $\phi'_\alpha \rightarrow e^{2\pi i/3} \phi_\alpha$ of the Z_3 group (for recent review see [2]). The term $\mu(H_1 H_2)$ does not meet this requirement. Therefore it is replaced in the superpotential by

$$W_H = \lambda S(H_1 H_2) + \frac{1}{3} \kappa S^3, \quad (3.1)$$

where S is an additional superfield which is a singlet with respect to $SU(2)_W$ and $U(1)_Y$ gauge transformations. A spontaneous breakdown of the EW symmetry leads to the emergence of the VEV of extra singlet field $\langle S \rangle = s/\sqrt{2}$ and an effective μ parameter is generated ($\mu = \lambda s/\sqrt{2}$).

The potential energy of the Higgs field interaction can be written as a sum

$$V = V_F + V_D + V_{\text{soft}} + \Delta V, \quad (3.2)$$

$$V_F = \lambda^2 |S|^2 (|H_1|^2 + |H_2|^2) + \lambda^2 |(H_1 H_2)|^2 + \lambda \kappa [S^{*2} (H_1 H_2) + h.c.] + \kappa^2 |S|^4, \quad (3.3)$$

$$V_D = \frac{g_2^2}{8} (H_1^+ \sigma_a H_1 + H_2^+ \sigma_a H_2)^2 + \frac{g_1^2}{8} (|H_1|^2 - |H_2|^2)^2, \quad (3.4)$$

$$V_{soft} = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_S^2 |S|^2 + \left[\lambda A_\lambda S(H_1 H_2) + \frac{\kappa}{3} A_\kappa S^3 + h.c. \right]. \quad (3.5)$$

At the tree level the Higgs potential (3.2) is described by the sum of the first three terms. V_F and V_D are the F and D terms. Their structure is fixed by the superpotential (3.1) and the EW gauge interactions in the common manner. The soft SUSY breaking terms are collected in V_{soft} . The set of soft SUSY breaking parameters involves soft masses m_1^2, m_2^2, m_S^2 and trilinear couplings A_κ, A_λ . The last term in Eq. (3.2), ΔV , corresponds to the contribution of loop corrections. In the leading one-loop approximation ΔV in the NMSSM is given by Eqs. (2.8) in which μ has to be replaced by $\lambda s/\sqrt{2}$. Further we assume that λ, κ and all soft SUSY breaking parameters are real so that CP is conserved.

At the physical vacuum of the Higgs potential

$$\langle H_1 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \quad \langle H_2 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_2 \end{pmatrix}, \quad \langle S \rangle = \frac{s}{\sqrt{2}}. \quad (3.6)$$

The equations for the extrema of the full Higgs boson effective potential in the directions (3.6) in the field space are given by

$$\frac{\partial V}{\partial s} = \left(m_S^2 + \frac{\lambda^2}{2} (v_1^2 + v_2^2) - \lambda \kappa v_1 v_2 \right) s - \frac{\lambda A_\lambda}{\sqrt{2}} v_1 v_2 + \frac{\kappa A_\kappa}{\sqrt{2}} s^2 + \kappa^2 s^3 + \frac{\partial \Delta V}{\partial s} = 0, \quad (3.7)$$

$$\frac{\partial V}{\partial v_1} = \left(m_1^2 + \frac{\lambda^2}{2} (v_2^2 + s^2) + \frac{g_2^2}{8} (v_1^2 - v_2^2) \right) v_1 - \left(\frac{\lambda \kappa}{2} s^2 + \frac{\lambda A_\lambda}{\sqrt{2}} s \right) v_2 + \frac{\partial \Delta V}{\partial v_1} = 0, \quad (3.8)$$

$$\frac{\partial V}{\partial v_2} = \left(m_2^2 + \frac{\lambda^2}{2} (v_1^2 + s^2) + \frac{g_2^2}{8} (v_2^2 - v_1^2) \right) v_2 - \left(\frac{\lambda \kappa}{2} s^2 + \frac{\lambda A_\lambda}{\sqrt{2}} s \right) v_1 + \frac{\partial \Delta V}{\partial v_2} = 0. \quad (3.9)$$

As in the MSSM upon the breakdown of the EW symmetry three goldstone modes (G^\pm and G^0) emerge, and are absorbed by the W^\pm and Z bosons. In the field space basis rotated by an angle β with respect to the initial direction, i.e.

$$\begin{aligned} H_1^- &= G^- \cos \beta + H^- \sin \beta, & H_2^+ &= H^+ \cos \beta - G^+ \sin \beta, \\ \text{Im} H_1^0 &= (P \sin \beta + G^0 \cos \beta)/\sqrt{2}, & \text{Re} H_1^0 &= (h \cos \beta - H \sin \beta + v_1)/\sqrt{2}, \\ \text{Im} H_2^0 &= (P \cos \beta - G^0 \sin \beta)/\sqrt{2}, & \text{Re} H_2^0 &= (h \sin \beta + H \cos \beta + v_2)/\sqrt{2}, \\ \text{Im} S &= P_S/\sqrt{2}, & \text{Re} S &= (s + N)/\sqrt{2}, \end{aligned} \quad (3.10)$$

these unphysical degrees of freedom decouple and the mass terms in the Higgs boson potential can be written as follows

$$V_{mass} = M_{H^\pm}^2 H^+ H^- + \frac{1}{2} (P P_S) \tilde{M}^2 \begin{pmatrix} P \\ P_S \end{pmatrix} + \frac{1}{2} (H h S) M^2 \begin{pmatrix} H \\ h \\ S \end{pmatrix}. \quad (3.11)$$

From the conditions for the extrema (3.7)–(3.9) one can express m_S^2, m_1^2, m_2^2 via other fundamental parameters, $\tan \beta$ and s . Substituting the obtained relations for the soft masses in the 2×2 CP-odd mass matrix \tilde{M}_{ij}^2 we get:

$$\tilde{M}_{11}^2 = m_A^2 = \frac{4\mu^2}{\sin^2 2\beta} \left(x - \frac{\kappa}{2\lambda} \sin 2\beta \right) + \tilde{\Delta}_{11}, \quad \tilde{M}_{22}^2 = \frac{\lambda^2 v^2}{2} x + \frac{\lambda \kappa}{2} v^2 \sin 2\beta - 3 \frac{\kappa}{\lambda} A_\kappa \mu + \tilde{\Delta}_{22},$$

$$\tilde{M}_{12}^2 = \tilde{M}_{21}^2 = \sqrt{2}\lambda v\mu \left(\frac{x}{\sin 2\beta} - 2\frac{\kappa}{\lambda} \right) + \tilde{\Delta}_{12}, \quad (3.12)$$

where $x = \frac{1}{2\mu} \left(A_\lambda + 2\frac{\kappa}{\lambda}\mu \right) \sin 2\beta$, $\mu = \frac{\lambda s}{\sqrt{2}}$ and $\tilde{\Delta}_{ij}$ are contributions of the loop corrections to the mass matrix elements. The mass matrix (3.12) can be easily diagonalized. The corresponding eigenvalues are given by

$$m_{A_2, A_1}^2 = \frac{1}{2} \left(\tilde{M}_{11}^2 + \tilde{M}_{22}^2 \pm \sqrt{(\tilde{M}_{11}^2 - \tilde{M}_{22}^2)^2 + 4\tilde{M}_{12}^4} \right), \quad (3.13)$$

Because the charged components of the Higgs doublets are not mixed with the neutral Higgs states the charged Higgs fields H^\pm are already physical mass eigenstates with

$$M_{H^\pm}^2 = m_A^2 - \frac{\lambda^2 v^2}{2} + M_W^2 + \Delta_\pm. \quad (3.14)$$

Here Δ_\pm includes loop corrections to the charged Higgs mass.

In the rotated basis H, h, N the mass matrix of the CP–even Higgs sector takes the form [1],[3]:

$$M^2 = \begin{pmatrix} M_{11}^2 & M_{12}^2 & M_{13}^2 \\ M_{21}^2 & M_{22}^2 & M_{23}^2 \\ M_{31}^2 & M_{32}^2 & M_{33}^2 \end{pmatrix} = \begin{pmatrix} \frac{1}{v^2} \frac{\partial^2 V}{\partial \beta^2} & \frac{1}{v} \frac{\partial^2 V}{\partial v \partial \beta} & \frac{1}{v} \frac{\partial^2 V}{\partial s \partial \beta} \\ \frac{1}{v} \frac{\partial^2 V}{\partial v \partial \beta} & \frac{\partial^2 V}{\partial v^2} & \frac{\partial^2 V}{\partial v \partial s} \\ \frac{1}{v} \frac{\partial^2 V}{\partial s \partial \beta} & \frac{\partial^2 V}{\partial v \partial s} & \frac{\partial^2 V}{\partial s^2} \end{pmatrix}, \quad (3.15)$$

$$\begin{aligned} M_{11}^2 &= m_A^2 + \left(\frac{\bar{g}^2}{4} - \frac{\lambda^2}{2} \right) v^2 \sin^2 2\beta + \Delta_{11}, \\ M_{22}^2 &= M_Z^2 \cos^2 2\beta + \frac{\lambda^2}{2} v^2 \sin^2 2\beta + \Delta_{22}, \\ M_{33}^2 &= 4\frac{\kappa^2}{\lambda^2} \mu^2 + \frac{\kappa}{\lambda} A_\kappa \mu + \frac{\lambda^2 v^2}{2} x - \frac{\kappa \lambda}{2} v^2 \sin 2\beta + \Delta_{33}, \\ M_{12}^2 &= M_{21}^2 = \left(\frac{\lambda^2}{4} - \frac{\bar{g}^2}{8} \right) v^2 \sin 4\beta + \Delta_{12}, \\ M_{13}^2 &= M_{31}^2 = -\sqrt{2}\lambda v\mu x \operatorname{ctg} 2\beta + \Delta_{13}, \\ M_{23}^2 &= M_{32}^2 = \sqrt{2}\lambda v\mu (1-x) + \Delta_{23}, \end{aligned} \quad (3.16)$$

where Δ_{ij} can be calculated by differentiating ΔV .

At least one Higgs state in the CP–even sector is always light. Since the minimal eigenvalue of a Hermitian matrix does not exceed its smallest diagonal element the lightest CP–even Higgs boson squared mass $m_{h_1}^2$ remains smaller than $M_{22}^2 \sim M_Z^2$ even when the supersymmetry breaking scale is much larger than the EW scale¹. The upper bound on the lightest Higgs mass in the NMSSM differs from the theoretical bound in the MSSM only for moderate values of $\tan \beta$. In the leading two–loop approximation the lightest Higgs boson mass in the NMSSM does not exceed 135 GeV. As follows from the explicit form of the mass matrices (3.12) and (3.16) at the tree–level, the spectrum of the Higgs bosons and their couplings depend on the six parameters: $\lambda, \kappa, \mu, \tan \beta, A_\kappa$ and m_A (or x).

¹The same theorem may lead to the upper bound on the mass of the lightest neutralino [4].

First let us consider the MSSM limit of the NMSSM. Because the strength of the interaction of the extra SM singlet superfield S with H_1 and H_2 is determined by the size of the coupling λ in the superpotential (3.1) the MSSM expressions for the Higgs masses and couplings are reproduced when λ tends to be zero. On the other hand the equations (3.8)–(3.9) imply that s should grow with decreasing λ as M_Z/λ to ensure the correct breakdown of the EW symmetry. In the limit $\lambda \rightarrow 0$ all terms, which are proportional to λv_i , in the minimization conditions (3.7) can be neglected and the corresponding equation takes the form:

$$s \left(m_S^2 + \frac{\kappa A_\kappa}{\sqrt{2}} s + \kappa^2 s^2 \right) \simeq 0 \quad (3.17)$$

The Eq. (3.17) has always at least one solution $s_0 = 0$. In addition two non-trivial roots arise if $A_\kappa^2 > 8m_S^2$. They are given by

$$s_{1,2} \simeq \frac{-A_\kappa \pm \sqrt{A_\kappa^2 - 8m_S^2}}{2\sqrt{2}\kappa}. \quad (3.18)$$

When $m_S^2 > 0$ the root $s_0 = 0$ corresponds to the local minimum of the Higgs potential (3.2)–(3.5) that does not lead to the acceptable solution of the μ -problem. The second non-trivial vacuum, that appears if $A_\kappa^2 > 8m_S^2$, remains unstable for $A_\kappa^2 < 9m_S^2$. Larger absolute values of A_κ ($A_\kappa^2 > 9m_S^2$) stabilizes the second minimum which is attained at $s = s_1(s_2)$ for negative (positive) A_κ . From Eq. (3.18) it becomes clear that the increasing of s can be achieved either by decreasing κ or by raising m_S^2 and A_κ . Since there is no natural reason why m_S^2 and A_κ should be very large while all other soft SUSY breaking terms are left in the TeV range, the values of λ and κ are obliged to go to zero simultaneously so that their ratio remains unchanged.

Since in the MSSM limit of the NMSSM mixing between singlet states and neutral components of the Higgs doublets is small the mass matrices (3.12) and (3.15)–(3.16) can be diagonalised using the perturbation theory [1],[3],[5]–[7]. At the tree-level the masses of two Higgs pseudoscalars are given by

$$m_{A_2}^2 \simeq m_A^2 = \frac{4\mu^2}{\sin^2 2\beta} \left(x - \frac{\kappa}{2\lambda} \sin 2\beta \right), \quad m_{A_1}^2 \simeq -3 \frac{\kappa}{\lambda} A_\kappa \mu. \quad (3.19)$$

The masses of two CP-even Higgs bosons are the same as in the MSSM (see Eq. (2.12)) while the mass of the extra CP-even Higgs state, which is predominantly a SM singlet field, is set by $\frac{\kappa}{\lambda} \mu$

$$m_{h_3}^2 \approx 4 \frac{\kappa^2}{\lambda^2} \mu^2 + \frac{\kappa}{\lambda} A_\kappa \mu + \frac{\lambda^2 v^2}{2} x \sin^2 2\beta - \frac{2\lambda^2 v^2 \mu^2 (1-x)^2}{M_Z^2 \cos^2 2\beta}. \quad (3.20)$$

The parameter A_κ occurs in the masses of extra scalar m_{h_3} and pseudoscalar m_{A_1} with opposite sign and is therefore responsible for their splitting. To ensure that the physical vacuum is a global minimum of the Higgs potential (3.2)–(3.5) and the masses-squared of all Higgs states are positive in this vacuum the parameter A_κ must satisfy the following constraints

$$-3 \left(\frac{\kappa}{\lambda} \mu \right)^2 \lesssim A_\kappa \left(\frac{\kappa}{\lambda} \mu \right) \lesssim 0. \quad (3.21)$$

The experimental constraints on the SUSY parameters obtained in the MSSM remain valid in the NMSSM with small λ and κ . For example, non-observation of any neutral Higgs particle and chargino at the LEP II imply that $\tan \beta \gtrsim 2.5$ and $|\mu| \gtrsim 90 - 100 \text{ GeV}$.

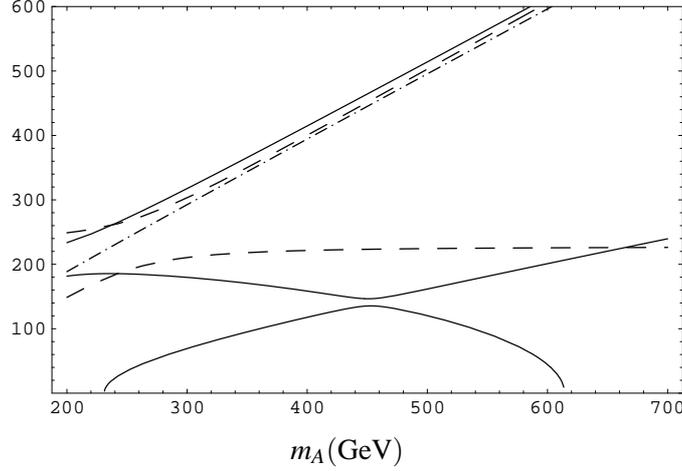


Figure 2: One-loop masses of the CP-even Higgs bosons versus m_A for $\lambda = 0.6$, $\kappa = 0.36$, $\tan\beta = 3$, $\mu = 150\text{ GeV}$, $A_\kappa = 135\text{ GeV}$, $m_Q^2 = m_U^2 = M_S^2$, $X_t = \sqrt{6}M_S$ and $M_S = 700\text{ GeV}$. Solid, dashed and dashed-dotted lines correspond to the masses of the CP-even, CP-odd and charged Higgs bosons respectively.

Decreasing κ reduces the masses of extra scalar and pseudoscalar states so that for $\kappa \ll \lambda$ they can be the lightest particles in the Higgs boson spectrum. In the limit $\kappa \rightarrow 0$ the mass of the lightest pseudoscalar state vanishes. In the considered limit the Lagrangian of the NMSSM is invariant under transformations of $SU(2) \times [U(1)]^2$ global symmetry. Extra $U(1)$ global symmetry gets spontaneously broken by the VEV of the singlet field S , giving rise to a massless Goldstone boson, the Peccei–Quinn (PQ) axion. In the PQ-symmetric NMSSM astrophysical observations exclude any choice of the parameters unless one allows s to be enormously large ($> 10^9 - 10^{11}\text{ GeV}$). These huge vacuum expectation values of the singlet field can be consistent with the EWSB only if $\lambda \sim 10^{-6} - 10^{-9}$ [8]–[9]. Therefore here we restrict our consideration to small but non-zero values of $\kappa \lesssim \lambda^2$ when the PQ-symmetry is only slightly broken.

As evident from Eq. (3.20) at small values of κ the mass-squared of the lightest Higgs scalar tends to be negative if $|\mu|$ is large and/or the auxiliary variable x differs too much from unity. Due to the vacuum stability requirement, which implies the positivity of the mass-squared of all Higgs particles, x has to be localized near unity, i.e.

$$1 - \left| \frac{\sqrt{2}\kappa M_Z}{\lambda^2 v} \right| < x < 1 + \left| \frac{\sqrt{2}\kappa M_Z}{\lambda^2 v} \right|. \quad (3.22)$$

This leads to the hierarchical structure of the Higgs spectrum. Indeed, combining LEP limits on $\tan\beta$ and μ one gets that $m_A^2 \gtrsim 9M_Z^2 x$. Because of this the heaviest CP-odd, CP-even and charged Higgs bosons are almost degenerate with masses around $m_A \simeq \mu \tan\beta$ while the SM-like Higgs state has a mass of the order of M_Z .

The main features of the NMSSM Higgs spectrum discussed above are retained when the couplings λ and κ increase. For the appreciable values of κ and λ the slight breaking of the PQ-symmetry can be caused by the RG flow of these couplings from the GUT scale M_X to M_Z . In the infrared region the solutions of the NMSSM RG equations are focused near the intersection of the Hill-type effective surface and invariant line [10]–[12]. As a result at the EW scale κ/λ tend to be

less than unity even when $\kappa(M_X) > \lambda(M_X)$ initially. In Fig. 2 the dependence of the masses and couplings of the Higgs bosons on m_A is examined. As a representative example we fix the Yukawa couplings so that $\lambda(M_X) = \kappa(M_X) = 2h_t(M_X) = 1.6$, that corresponds to $\tan\beta \gtrsim 3$, $\lambda(M_t) = 0.6$ and $\kappa(M_t) = 0.36$. In order to obtain a realistic spectrum, we include the leading one-loop corrections from the top and stop loops. From Fig. 2 it becomes clear that the requirement of stability of the physical vacuum limits the range of variations of m_A from below and above maintaining the mass hierarchy in the Higgs spectrum. Relying on this mass hierarchy the approximate solutions for the Higgs masses and couplings can be obtained [6],[7]. The numerical results in Fig. 2 reveal that the masses of the heaviest CP-even, CP-odd and charged Higgs states are approximately degenerate while the other three neutral states are considerably lighter. The hierarchical structure of the Higgs spectrum ensures that the heaviest CP-even and CP-odd Higgs bosons are predominantly composed of H and P . As before the lightest Higgs scalar and pseudoscalar are singlet dominated, making their observation quite problematic. The second lightest CP-even Higgs boson has a mass around 130 GeV, mimicking the lightest Higgs scalar in the MSSM. Observing two light scalars and one pseudoscalar Higgs particles but no charged Higgs boson at future colliders would yield an opportunity to differentiate the NMSSM with a slightly broken PQ-symmetry from the MSSM even if the heavy Higgs states are inaccessible.

The presence of light singlet scalar and pseudoscalar permits to weaken the LEP lower bound on the lightest Higgs boson mass. These states have reduced couplings to Z -boson that could allow them to escape the detection at LEP. On the other hand singlet scalar can mix with the SM-like superposition h of the neutral components of Higgs doublets resulting in the reduction of the couplings of the second lightest Higgs scalar to Z -boson. This relaxes LEP constraints so that the SM-like Higgs state does not need to be considerably heavier than 100 GeV. Therefore large contribution of loop corrections to the mass of the SM-like Higgs boson is not required. Another possibility to overcome the little hierarchy problem is to allow the SM-like Higgs state to decay predominantly into two light singlet pseudoscalars A_1 (for recent review see [13]). This can be achieved because the coupling of the SM-like Higgs boson to the b -quark is rather small. If this coupling is substantially smaller than the coupling of the SM-like Higgs state to A_1 then the decay mode $h \rightarrow A_1 + A_1$ dominates. The singlet pseudoscalar can sequentially decay into either $b\bar{b}$ or $\tau\bar{\tau}$ leading to four fermion decays of the SM-like Higgs boson. In this case, again, the corresponding Higgs eigenstate might be relatively light that permits to avoid little hierarchy problem.

However even when the couplings of the lightest CP-even Higgs state are almost the same as in the SM it is substantially easier to overcome LEP constraint on the mass of the SM-like Higgs boson in the NMSSM than in the MSSM. Indeed, in the NMSSM the theoretical upper bound on $m_{h_1}^2$, which is given by M_{22}^2 in Eq. (3.16), contains an extra term $\frac{\lambda^2}{2}v^2\sin^2 2\beta$ which is not present in the MSSM. Due to this term the maximum possible value of the mass of the lightest Higgs scalar in the NMSSM can be considerably larger as compared with the MSSM at moderate values of $\tan\beta$. In our analysis we require the validity of perturbation theory up to the scale M_X . This sets stringent upper limit on $\lambda(M_t)$ at low energies for each particular choice of $\tan\beta$. Using theoretical restrictions on $\lambda(M_t)$ one can compute the the maximum possible value of $m_{h_1}^2$ for each given value of $\tan\beta$. Fig. 3 shows the dependence of the upper bound on the lightest Higgs boson mass as a function of $\tan\beta$ in the MSSM and NMSSM. From Fig. 3 one can see that at the tree-level the lightest CP-even Higgs state in the NMSSM can be considerably heavier than in the

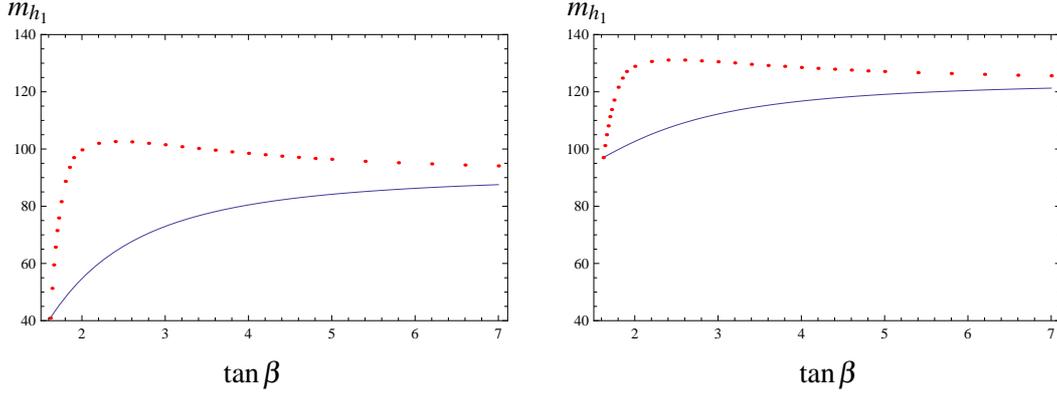


Figure 3: *Left:* Tree-level upper bound on the lightest Higgs boson mass in the MSSM and NMSSM as a function of $\tan\beta$. *Right:* The dependence of the two-loop upper bound on the lightest Higgs boson mass on $\tan\beta$ for $m_t(M_t) = 165\text{ GeV}$, $m_Q^2 = m_U^2 = M_S^2$, $X_t = \sqrt{6}M_S$ and $M_S = 700\text{ GeV}$. The solid and dotted lines represent the theoretical restrictions on m_{h_1} in the MSSM and NMSSM respectively.

MSSM at moderate values of $\tan\beta$. As a consequence in the leading two-loop approximation it is substantially easier to get $m_{h_1} \gtrsim 114.4\text{ GeV}$ in the NMSSM than in the MSSM for $\tan\beta = 2 - 4$.

4. Higgs spectrum in the E_6 inspired SUSY models with extra $U(1)'$ factor

Another solution to the μ problem arises within superstring inspired models based on the E_6 gauge group. At high energies E_6 can be broken $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)'$. An extra $U(1)'$ that appears at low energies is a linear superposition of $U(1)_\chi$ and $U(1)_\psi$:

$$U(1)' = U(1)_\chi \cos\theta + U(1)_\psi \sin\theta, \quad (4.1)$$

where two anomaly-free $U(1)_\psi$ and $U(1)_\chi$ symmetries are defined by: $E_6 \rightarrow SO(10) \times U(1)_\psi$, $SO(10) \rightarrow SU(5) \times U(1)_\chi$. If $\theta \neq 0$ or π the extra $U(1)'$ gauge symmetry forbids an elementary μ term but allows interaction $\lambda S(H_1 H_2)$ in the superpotential. After EWSB the scalar component of the SM singlet superfield S acquires a non-zero VEV breaking $U(1)'$ and an effective μ term of the required size is automatically generated.

The Higgs sector of the considered models includes two Higgs doublets as well as a SM-like singlet field S that carries $U(1)'$ charge. The Higgs effective potential can be written as

$$\begin{aligned} V &= V_F + V_D + V_{soft} + \Delta V, \\ V_F &= \lambda^2 |S|^2 (|H_d|^2 + |H_u|^2) + \lambda^2 |(H_d H_u)|^2, \\ V_D &= \frac{g_2^2}{8} \left(H_d^\dagger \sigma_a H_d + H_u^\dagger \sigma_a H_u \right)^2 + \frac{g_1^2}{8} (|H_d|^2 - |H_u|^2)^2 \\ &\quad + \frac{g_1'^2}{2} (\tilde{Q}_1 |H_d|^2 + \tilde{Q}_2 |H_u|^2 + \tilde{Q}_S |S|^2)^2, \\ V_{soft} &= m_S^2 |S|^2 + m_1^2 |H_d|^2 + m_2^2 |H_u|^2 + \left[\lambda A_\lambda S (H_u H_d) + h.c. \right], \end{aligned} \quad (4.2)$$

where g_1' is $U(1)'$ gauge coupling and \tilde{Q}_1 , \tilde{Q}_2 and \tilde{Q}_S are effective $U(1)'$ charges of H_1 , H_2 and S respectively. In Eq. (4.2) V_F and V_D are the F and D terms, V_{soft} contains a set of soft SUSY breaking terms while ΔV represents the contribution of loop corrections.

At the physical vacuum the Higgs fields acquire VEVs given by Eq. (3.6) thus breaking the $SU(2)_W \times U(1)_Y \times U(1)'$ symmetry to $U(1)_{em}$. As a result 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 form one CP-odd, three CP-even and two charged states. The masses of the CP-odd and charged Higgs bosons can be written as

$$m_A^2 = \frac{2\lambda^2 s^2 x}{\sin^2 2\beta} + O(M_Z^2), \quad m_{H^\pm}^2 = m_A^2 + O(M_Z^2), \quad (4.3)$$

where $x = \frac{A_\lambda}{\sqrt{2}\lambda_s} \sin 2\beta$. The masses of two heaviest CP-even states are set by $M_{Z'}$ and m_A , i.e.

$$m_{h_3}^2 = m_A^2 + O(M_Z^2), \quad m_{h_2}^2 = M_{Z'}^2 + O(M_Z^2), \quad (4.4)$$

where $M_{Z'} \simeq g_1' \tilde{Q}_{SS}$. The lightest CP-even Higgs boson has a mass which is less than

$$m_{h_1}^2 \lesssim \frac{\lambda^2}{2} v^2 \sin^2 2\beta + M_Z^2 \cos^2 2\beta + g_1'^2 v^2 \left(\tilde{Q}_1 \cos^2 \beta + \tilde{Q}_2 \sin^2 \beta \right)^2 + \Delta. \quad (4.5)$$

In Eq. (4.5) Δ represents the contribution of loop corrections. Since the mass of the Z' boson in the E_6 inspired models has to be heavier than 800 – 900 GeV at least one CP-even Higgs state, which is singlet dominated, is always heavy. If $m_A < M_{Z'}$ then we get MSSM-type Higgs spectrum. When $m_A > M_{Z'}$ the heaviest CP-even, CP-odd and charged states are almost degenerate with masses around m_A . In this case the lightest Higgs state is predominantly the SM-like superposition h of the neutral components of Higgs doublets.

Recently the detailed analysis of the Higgs sector was performed within a particular E_6 inspired SUSY model with an extra $U(1)_N$ gauge symmetry that corresponds to $\theta = \arctan \sqrt{15}$ [14]-[15]. The extra $U(1)_N$ gauge symmetry is defined such that right-handed neutrinos do not participate in the gauge interactions. Only in this Exceptional Supersymmetric Standard Model (E_6 SSM) right-handed may be superheavy, 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 [16]. To ensure anomaly cancellation the particle content of the E_6 SSM is extended to include three complete fundamental 27 representations of E_6 . In addition to the complete 27_i multiplets the low energy particle spectrum of the E_6 SSM is supplemented by $SU(2)_W$ doublet H' and anti-doublet \bar{H}' states from extra $27'$ and $\bar{27}'$ to preserve gauge coupling unification. The unification of gauge couplings in the considered model can be achieved for any phenomenologically acceptable value of $\alpha_3(M_Z)$ consistent with the measured low energy central value [17]. The Higgs spectrum within the E_6 SSM was studied in [14]-[15], [18]-[19]. It was argued that even at the tree level the lightest Higgs boson mass in this model can be larger than 120 GeV. Therefore nonobservation of the Higgs boson at LEP does not cause any trouble for the E_6 SSM, even at tree-level. In the leading two-loop approximation the mass of the lightest CP-even Higgs boson in the considered model does not exceed 150 – 155 GeV [14]. The presence of light exotic particles in the E_6 SSM spectrum lead to the nonstandard decays of the SM-like Higgs boson which were discussed in [20].

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