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Radiative corrections to Higgs boson masses in supersymmetric models

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Abstract. The effect of radiative corrections on the Higgs masses and couplings in supersymmetric models is summarized. Radiative corrections in both the minimal and nonminimal models are discussed. It is pointed out that large singlet Higgs vacuum expectation values are screened out from the radiative corrections to the lightest Higgs mass in nonminimal models. In supersymmetric models the crucial mass limit for the Higgs search may be around 150 GeV.

1. Introduction

It is now widely accepted that the only way to have an elementary Higgs scalar at the weak scale, as required by the consistency of the Standard Model, is to have approximate low energy supersymmetry [1]. The main motivation for this, i.e. supersymmetry effectively broken around the electroweak scale, comes from the naturalness or hierarchy problem [2] of the Standard Model. This problem is related to the fact that the scalar sector of the Standard Model receives large quadratically divergent radiative corrections of $O(\alpha/\pi)\Lambda^2$ where Λ is a cutoff which can be either the Planck mass or a mass scale associated with a grand unified theory (GUT). Thus, even if one could find an explanation as to why the Higgs mass was small at the tree level, one would have to explain why the radiative corrections were not much larger than the physical value of the Higgs mass. A light Higgs scalar is natural in a softly broken supersymmetric model since the quadratically divergent loop contributions mentioned above cancel between bosons and fermions leaving a finite supersymmetry breaking contribution to the Higgs mass. In the Minimal Supersymmetric version of the Standard Model (MSSM), based on the gauge group $G = SU(3)_C \times SU(2)_L \times U(1)_Y$, the spin-1 gauge bosons are accompanied by their spin- $\frac{1}{2}$ superpartners, the gauginos. Similarly, the three generations of quarks and leptons of the Standard Model are accompanied by their spin-0 superpartners, the squarks and sleptons. In addition, to give masses to all quarks and leptons, and to cancel gauge anomalies, at least two Higgs doublets H_1, H_2, with hypercharge Y = -1, +1, respectively, are required in the MSSM. To enforce baryon and lepton number conservation in renormalizable interactions, one imposes a discrete, multiplicative symmetry, called R-parity defined as R = +1 for ordinary particles, R = -1 for the superpartners, in the MSSM. The most general superpotential compatible with gauge invariance, renormalizibility and R- parity is

$$W = h^U QU^c H_2 + h^D QD^c H_1 + h^E LE^c H_1 + \mu H_1 H_2, \tag{1}$$

where Q, U^c, D^c, L, E^c are chiral superfields containing the left-handed components of ordinary quarks and leptons, H_1 and H_2 are the two Higgs chiral superfields, and family and group indices have been suppressed for notational simplicity. The first three terms in (1) represent the supersymmetric generalization of the Standard Model Yukawa couplings, whereas the last term is a globally supersymmetric Higgs mass term which is necessary to break $SU(2)\times U(1)$ without producing an unwanted axion.

The effective potential at tree level which follows from (1) can be written as

$$V = V_F + V_D, \tag{2}$$

where

$$V_F = \sum_i | F_i |^2 \equiv \sum_i | \frac{\partial W}{\partial \phi_i} |^2, \tag{3}$$

and

$$V_D = \frac{1}{2}\sum_a [D_a]^2 \equiv \frac{1}{2}\sum_a (-g_a \phi_i^* T_{ij}^a \phi_j)^2, \tag{4}$$

are, respectively the F and D term contributions to the potential, and where ϕ_i denotes a generic spin-0 field. Here g_a stands for the gauge coupling constants g_3, g and g' corresponding to the different factors of the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$, respectively. In addition, one must add explicit but soft supersymmetry-breaking terms to the potential (2) to take into account the fact that supersymmetry must be broken in nature. These terms, which must preserve the good ultraviolet properties of supersymmetric theories in order to solve the naturalness or hierarchy problem, include [3] mass terms for scalar fields and gauginos, as well as a restricted set of scalar interaction terms proportional to the corresponding superpotential couplings:

$$\begin{aligned} V_{soft} &= \sum_i m_i^2 | \phi_i |^2 + \frac{1}{2}\sum_a M_a \bar{\lambda}_a \lambda_a + (h^U A^U QU^c H_2 \\ &\quad + h^D A^D QD^c H_1 + h^E A^E LE^c H_1 + m_3^2 H_1 H_2 + h.c.), \end{aligned} \tag{5}$$

where $\lambda_a (a = 1, 2, 3)$ denotes the generic gaugino field. We note that A^U, A^D, A^E are matrices in generation space. The properties of Higgs bosons in the MSSM, including masses, mixing angles, and couplings are completely determined by Eqs. (2) – (5).

We have noted above that a bilinear term $\mu H_1 H_2$ is necessary in the superpotential (1) in order to avoid an unacceptable electroweak axion in the MSSM. Softly broken supersymmetry ensures that radiative corrections to the mass parameter μ are under control so that μ is naturally of the order of electroweak scale. However, supersymmetry does not provide a dynamical reason why μ should be so small in the first place. The simplest mechanism that provides a dynamical source for a such a bilinear term is the inclusion of an additional $SU(2) \times U(1)$ singlet chiral superfield N [4, 5]. Then, if the superpotential contains a trilinear term $\lambda N H_1 H_2$, and if N develops a vacuum expectation value $< N > \equiv x$, a bilinear term $\mu H_1 H_2$

with $\mu = \lambda x$ is generated. In the presence of soft supersymmetry breaking, we expect $x = O(m_W)$, and hence $\mu = O(m_W)$. Such a mechanism does operate in many superstring models based on E_6 [6] and $SU(5) \times U(1)$ grand unified groups [7] in which the renormalizable superpotential is purely trilinear, and automatically contains a trilinear term coupling the two Higgs doublets to a singlet superfield N. The superpotential for such a minimal extension of the Minimal Supersymmetric Standard Model (NMSSM) containing two Higgs doublets and a singlet chiral superfield can then be written as

$$W = h^U QU^c H_2 + h^D QD^c H_1 + h^E LE^c H_1 + \lambda H_1 H_2 N - \frac{k}{3} N^3, \tag{6}$$

where the last term is needed to avoid an unacceptable axion. The effective potential at the tree level for the Nonminimal Supersymmetric Standard Model (NMSSM) (6) can then be obtained in analogy with the MSSM by using Eqs. (2) – (5), where only trilinear soft supersymmetry breaking terms corresponding to the superpotential (6) appear in Eq. (5).

2. Higgs Boson Masses and Couplings

2.1 The Minimal Supersymmetric Standard Model

Assuming that squark and slepton fields have vanishing vacuum expectation values, the tree level scalar potential in the MSSM involving only the Higgs fields $H_1 \equiv (H_1^0, H^-)$ and $H_2 \equiv (H^+, H_2^0)$ can be written as [8]

$$\begin{aligned} V_0 &= m_1^2 \mid H_1 \mid^2 + m_2^2 \mid H_2 \mid^2 + m_3^2 (H_1 H_2 + h.c.) \\ &\quad + \frac{1}{8} g^2 \left(H_2^+ \vec{\sigma} H_2 + H_1^+ \vec{\sigma} H_1 \right)^2 + \frac{1}{8} g'^2 \left(\mid H_2 \mid^2 - \mid H_1 \mid^2 \right)^2, \end{aligned} \tag{7}$$

where

$$m_1^2 \equiv m_{H_1}^2 + \mu^2, \qquad m_2^2 \equiv m_{H_2}^2 + \mu^2, \tag{8}$$

and it is possible to choose a basis such that $m_3^2 \leq 0$. After the neutral Higgs bosons H_1^0 and H_2^0 develop vacuum expectation values, $< H_1^0 > \equiv v_1, < H_2^0 > \equiv v_2$, which can be taken to be real and positive without loss of generality, we are left with five physical degrees of freedom. These consist of two CP-even $\left(h^0, H^0; m_{h^0} < m_{H^0}\right)$, one CP-odd (A°), and two charged $(H^\pm)$ Higgs bosons. Taking into account the minimization conditions on the potential together with the physical constraint $m_W^2 = \frac{1}{2} g^2 (v_1^2 + v_2^2)$, the Higgs sector of the MSSM contains only two parameters for which a convenient choice usually adopted is m_A, the physical mass of the CP-odd neutral Higgs boson, and $\tan\beta \equiv v_2/v_1$. The parameter m_A is essentially unconstrained, whereas $1 \leq \tan\beta \lesssim m_t/m_b$ in the Minimal Supersymmetric Standard Model with radiative symmetry breaking [9]. In terms of these parameters the mass squared matrix for the neutral CP-even Higgs bosons can be written as

$$(\mathcal{M}_S^0)^2 = \begin{pmatrix} m_A^2 \sin^2\beta + m_Z^2 \cos^2\beta & -(m_A^2 + m_Z^2) \sin\beta\cos\beta \\ -(m_A^2 + m_Z^2) \sin\beta\cos\beta & m_A^2 \cos^2\beta + m_Z^2 \sin^2\beta \end{pmatrix} \tag{9}$$

from which we obtain the tree level squares masses

$$m^2_{h,H} = \frac{1}{2}\left[m^2_A + m^2_Z \mp \sqrt{(m^2_A + m^2_Z)^2 - 4m^2_A m^2_Z \cos^2 2\beta}\right] \tag{10}$$

through diagonalization with a mixing angle α given by

$$\cos 2\alpha = -\cos 2\beta \frac{(m^2_A - m^2_Z)}{(m^2_H - m^2_h)}. \tag{11}$$

The mass of the charged Higgs boson is given by

$$m^2_{H^\pm} = m^2_W + m^2_A. \tag{12}$$

Eqs. (10) and (12) lead to the tree level inequalities among the Higgs boson masses

$$m_h < m_Z < m_H, \quad m_h < m_A < m_H, \quad m_W, m_A < m_{H^\pm}. \tag{13}$$

The first inequality, which raised the hope that nonobservation of a Higgs with mass less than the Z boson mass could rule out the MSSM, is phenomenologically the most important relation. This relation, as we will see, is no longer true when we take into account the radiative corrections.

The couplings of the Higgs bosons to known particles are controlled by the parameters α and β [10] through the following factors :

	h	H	A
$t\bar{t}$	$\cos\alpha/\sin\beta$	$\sin\alpha/\sin\beta$	$\gamma_5\cos\beta$
$b\bar{b},\ \tau\bar{\tau}$	$-\sin\alpha/\cos\beta$	$\cos\alpha/\cos\beta$	$\gamma_5\tan\beta$
$WW,\ ZZ$	$\sin(\beta-\alpha)$	$\cos(\beta-\alpha)$	0
ZA	$\cos(\beta-\alpha)$	$\sin(\beta-\alpha)$	0

(14)

	H^+
$t\bar{b}$	$m_t\cos\beta(1-\gamma_5) + m_b\tan\beta(1+\gamma_5)$
$\nu\bar{\tau}$	$m_\tau\tan\beta(1+\gamma_5)$
ZH^+	$1-2\sin^2\theta_W$

(15)

We note that there is no tree-level WZH vertex. It is conventional to choose $0 \le \beta \le \pi/2$ and $-\pi/2 \le \alpha \le 0$. In the limit $m_A \to \infty, \alpha \to \beta - \pi/2$ for a fixed β, whereas A, H and $H^\pm$ become approximately degenerate and physically irrelevant. In this limit the couplings of h approach those of Standard Model Higgs boson. The Higgs boson h of MSSM would be almost indistinguishable from the Standard Model Higgs boson for $m_A \gtrsim 200$ GeV and $\tan\beta > 1$ unless some of the other Higgs bosons $H, H^\pm$ or A are discovered [11].

2.2 The Nonminimal Supersymmetric Standard Model

For the case of the nonminimal model (6), the tree level Higgs potential can be written as

$$V_0 = V_F + V_D + V_{soft}, \tag{16}$$

$$\begin{aligned} V_F &= \lambda^2\left[(|H_1|^2 + |H_2|^2)|N|^2 + |H_1H_2|^2\right] + k^2|N|^4 \\ &\quad -(\lambda k H_1 H_2 N^{*2} + h.c.) - \lambda^2(H^0_1 H^0_2 H^{+*} H^{-*} + h.c.), \end{aligned} \tag{17}$$

$$V_D = \frac{g^2}{8}\left(H_2^+\vec{\sigma}H_2 + H_1^+\vec{\sigma}H_1\right)^2 + \frac{g'^2}{8}\left(\mid H_2\mid^2 - \mid H_1\mid^2\right)^2, \tag{18}$$

$$\begin{aligned} V_{soft} &= m_{H_1}^2 \mid H_1\mid^2 + m_{H_2}^2 \mid H_2\mid^2 + m_N^2|N|^2 \\ &\quad -(\lambda A_\lambda H_1 H_2 N + h.c.) - (\frac{1}{3}kA_k N^3 + h.c.), \end{aligned} \tag{19}$$

where we have chosen $\lambda, k, A_\lambda, A_k$ to be real so that there is no explicit CP violation [5]. After spontaneous symmetry breaking, with $< H_1^0 > \equiv v_1$, $< H_2^0 > \equiv v_2$ and $< N > \equiv x$ chosen to be real and positive in order for the vacuum to conserve CP [5], the physical Higgs boson spectrum consists of three scalars, two pseudoscalars and two charged particles. The minimization conditions on the potential together with the constraint $m_W^2 = \frac{1}{2}g^2(v_1^2+v_2^2)$ imply that all Higgs boson masses can be expressed in terms of six parameters $\lambda, k, A_\lambda, A_k, x$ and $\tan\beta$. In this model there is a bound on the mass of the lightest neutral CP-even Higgs boson [12] given by

$$m_h^2 \leq m_Z^2\left[\cos^2 2\beta + \frac{2\lambda^2}{g^2}\cos^2\theta_W \sin^2 2\beta\right], \tag{20}$$

which follows from the fact that the smallest eigenvalue of a real, symmetric $n\times n$ matrix is smaller that the smallest eigenvalue of the upper left 2×2 submatrix. The charged Higgs boson mass is given by

$$m_{H^\pm}^2 = m_W^2 - \lambda^2(v_1^2+v_2^2) + \lambda(A_\lambda + kx)\frac{2x}{\sin 2\beta}, \tag{21}$$

which could be less or greater than m_W^2, depending upon the relative magnitudes of the last two terms, in contrast to the MSSM inequality (13).

We note that the upper bound (20) is controlled by m_Z, and is independent of the supersymmetric particle masses and the scale of supersymmetry breaking. It is also independent of the vacuum expectation value, x, of the singlet Higgs field N. Furthermore, if one assumes that the underlying supersymmetric theory remains perturbative upto some scale M_{GUT}, in the energy range where the theory holds, then one obtains an upper bound on λ as a function of top quark mass (m_t) which, through Eq.(20), translates into a bound on $m_h \lesssim 140$ GeV [13]. The couplings of the Higgs bosons in the nonminimal model can, in general, be obtained numerically and are discussed in Ref. [5].

3. Radiative Corrections

3.1. Radiative Correction in the Minimal Supersymmetric Standard Model

We have seen above that the lightest neutral Higgs mass is less than the Z° mass in MSSM, and is controlled by m_Z in the NMSSM. However, it has recently been pointed out [14-20] that the Higgs masses are subject to large radiative corrections associated with the top quark and its SUSY partner, the stop. These corrections have been computed by using different methods: standard diagrammatic techniques [20], the one-loop effective potential [18], and the renormalization group equation approach [17,19]. In the following we shall follow [18] the method of the one-loop effective potential [21] to discuss the radiative corrections to the masses and couplings of the Higgs bosons in supersymmetric models. It has been shown that this method gives a fairly accurate estimate of the radiative corrections when compared with the explicit diagrammatic calculations [22].

When one-loop radiative corrections are included, the Higgs boson masses depend, besides the one-loop corrected m_A and $\tan\beta$ as independent parameters, on a number of additional parameters which enter through the loop effects. These include the top-quark mass (m_t), the squark (top and bottom) mass parameters which include the soft SUSY-breaking masses m_Q, m_U, m_D, the coefficients of trilinear terms, A_t and A_b (see Eq. (5)), and the parameter μ of the Higgs mixing term in the superpotential (1). The essential features of the one-loop corrections to the neutral CP-even Higgs boson masses can be studied as a function of m_t and a common SUSY-breaking mass $m_Q = m_U = m_D = \tilde{m}$, neglecting the small effects due to A_t, A_b, μ and the b-quark mass m_b. In this approximation the mass-squared matrix for the neutral CP-even Higgs bosons can be written as

$$(\mathcal{M}_S^R)^2 = \begin{pmatrix} m_A^2 \sin^2\beta + m_Z^2 \cos^2\beta & -(m_A^2 + m_Z^2)\sin\beta\cos\beta \\ -(m_A^2 + m_Z^2)\sin\beta\cos\beta & m_A^2\cos^2\beta + m_Z^2\sin^2\beta + \frac{\varepsilon}{\sin^2\beta} \end{pmatrix}, \quad (22)$$

where the effect of radiative corrections is summarized by a single parameter

$$\varepsilon = \frac{3g^2 m_t^4}{8\pi^2 m_W^2} \ell n\left(1 + \frac{\tilde{m}^2}{m_t^2}\right). \quad (23)$$

This leads to the masses

$$\begin{aligned} m_{h,H}^2 &= \frac{1}{2}\left[m_A^2 + m_Z^2 + \varepsilon/\sin^2\beta\right] \\ &\mp \frac{1}{2}\{[(m_A^2 - m_Z^2)\cos 2\beta + \varepsilon/\sin^2\beta]^2 + (m_A^2 + m_Z^2)^2\sin^2 2\beta\}^{\frac{1}{2}}. \end{aligned} \quad (24)$$

For $\tan\beta \geq 1$, these mass eigenvalues increase monotonically with m_A. The upper bound on the radiatively corrected mass of the lightest CP-even Higgs boson h is

$$m_h^2 \lesssim m_Z^2\cos^2 2\beta + \frac{3G_F}{\pi^2\sqrt{2}}\; m_t^4\;\; \ell n\left(1 + \frac{\tilde{m}^2}{m_t^2}\right), \quad (25)$$

For $m_t = 150$ GeV and $\tilde{m} = 1$ TeV, this gives an upper bound of $m_h < 115$ GeV. There is a corresponding lower bound on the heavier Higgs mass m_H,

$$m_H^2 > \frac{1}{2}\left[m_Z^2 + \varepsilon/\sin^2\beta\right] + \frac{1}{2}\left[(m_Z^2 - \varepsilon/\sin^2\beta)^2 + 4m_Z^2\varepsilon\right]^{\frac{1}{2}}. \quad (26)$$

We note that these bounds can be made more precise once $\tan\beta$ is specified. From (24) and (25) it is clear that the tree level predictions for the mass of lightest Higgs can be badly violated by the radiative corrections which are proportional to m_t^4/m_W^2.

The effective potential method [23] has also been used to calculate the one loop radiative corrections to the charged Higgs mass in the MSSM [24]. The leading corrections in this case are proportional to m_t^2/m_W^2 and are considerably smaller. In Figure 1 we show the impact of these radiative corrections by displaying [25] the contours of the maximum allowed values of m_h in the $(m_t, \tan\beta)$ plane. These maximum values are reached for $m_A \to \infty$. A value of $\tilde{m} = 1$ TeV has been chosen to obtain these results. To plot the actual values of the masses in the $(m_A, \tan\beta)$ plane, we fix the value of $m_t = 140$ GeV. Figure 2 shows the contours of constant m_h and m_H in the $(m_A, \tan\beta)$ plane for $\tilde{m} = 1$ TeV.

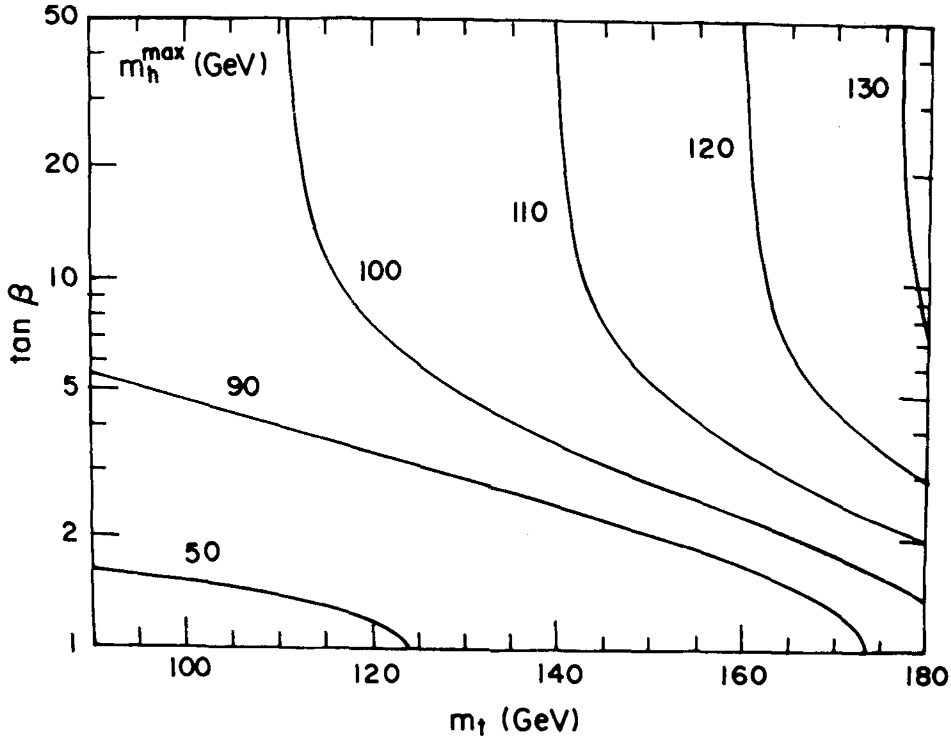

Figure 1. The contours of the maximum value of m_h, denoted by m_h^{max}, in the $(m_t, \tan\beta)$ plane for $\tilde{m} = 1$ TeV. The maximum is reached for $m_A \to \infty$ [25].

The one-loop radiatively corrected mixing angle α in the CP-even sector, which can be determined from the diagonalization of the mass squared matrix (22), is now given by

$$\tan 2\alpha = \frac{(m_A^2 + m_Z^2)}{(m_A^2 - m_Z^2)\cos 2\beta + \varepsilon/\sin^2\beta}. \tag{27}$$

The one-loop corrections to the mixing angle α depend on ε and can be large if $\tilde{m}$ is well above m_t and m_t is well above m_W. The couplings of the Higgs bosons in MSSM depend on α and β through the factors which are shown in Eq. (14,15), now with the radiatively corrected α, Eq. (27), substituted therein. Figure 3 shows the regions [11] where the W, t, b couplings of h are within 10% of the corresponding Standard Model Higgs couplings. The relevant processes for MSSM Higgs boson search at LEP I are $Z \to hZ^*$ and $Z \to hA$, which play a complementary role since their rates are proportional to $\sin^2(\beta - \alpha)$ and $\cos^2(\beta - \alpha)$, respectively;

thus any suppression in one coupling is accompanied by an enhancement in the complementary coupling. Similarly, the decays of h (or H) to WW, ZZ and ZA are complementary. Fig. 3 shows the contour plots of various coupling factors in the $(m_A, \tan\beta)$ plane which govern the couplings of Higgs bosons to fermions and vector bosons in the MSSM. We see that these couplings vary strongly across the parameter space.

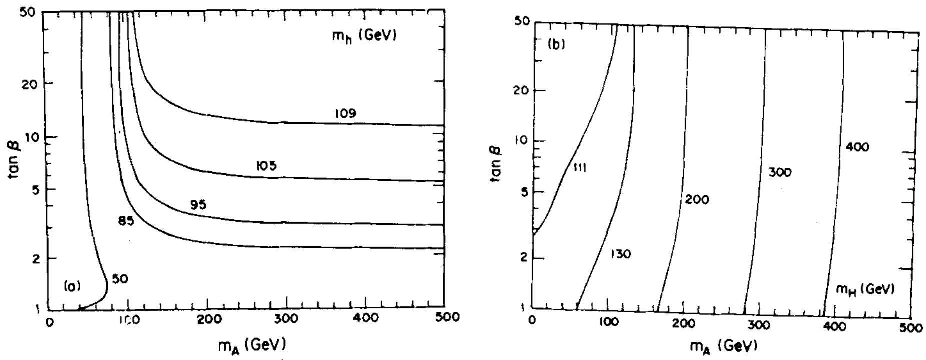

Figure 2. The contours of (a) the lightest Higgs mass m_h and (b) the heavier Higgs mass m_H, in the $(m_A, \tan\beta)$ plane. The values chosen for $\tilde{m}$ and m_t are 1 TeV and 140 GeV, respectively [25].

The method of 1-loop effective potential can also be used to calculate the leading corrections to the trilinear and quartic Higgs self-couplings. As an example, the leading radiative correction to the trilinear hAA coupling, which is important in the determination of the h branching ration, is [23]

$$\lambda_{hAA} = \lambda^{\circ}_{hAA} + \Delta\lambda_{hAA}, \tag{28}$$

where

$$\lambda^{\circ}_{hAA} = \frac{-igm_Z}{2\cos\theta_W}\cos 2\beta \sin(\beta+\alpha), \tag{29}$$

and

$$\Delta\lambda_{hAA} = \frac{-igm_Z}{2\cos\theta_W}\frac{3g^2\cos^2\theta_W}{8\pi^2}\frac{\cos\alpha\cos^2\beta}{\sin^3\beta}\frac{m_t^4}{m_W^4}\ell n\left(1+\frac{\tilde{m}^2}{m_t^2}\right). \tag{30}$$

Here we have neglected the D-term contributions, and the contribution from the b-quarks. Other Higgs couplings receive similar m_t^4/m_W^4 corrections. In contrast to the tree level results, the decay $h \to AA$ is now possible, due to radiative corrections, for some range of parameters.

Experimental limits which take into account the radiative corrections have now been obtained by the four LEP collaborations [26], using different methods to present and analyze the data, and different ranges of parameters in the evaluation of radiative corrections. The excluded regions in the (m_h, m_A) plane are shown

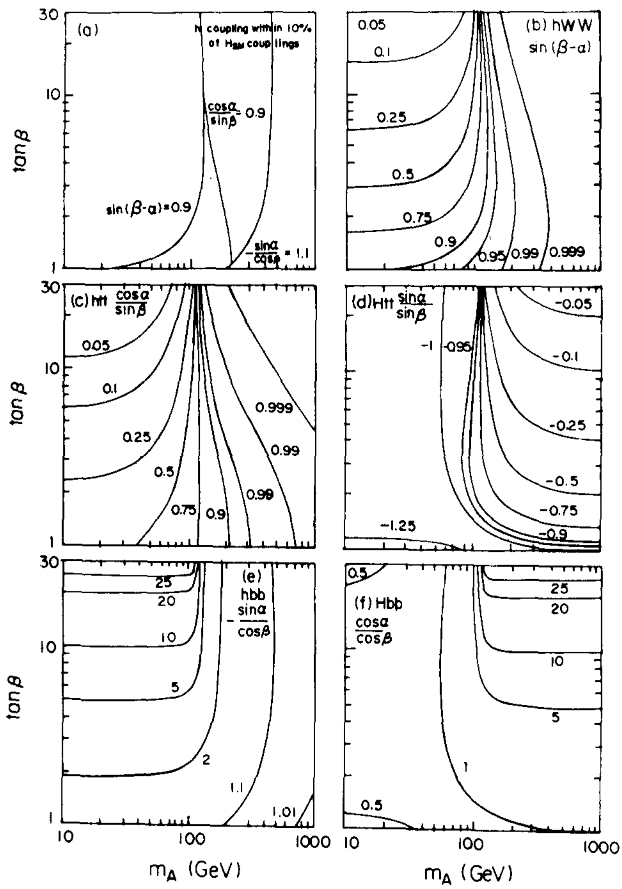

Figure 3. Contour plots in the $(m_A, \tan\beta)$ plane of the α-dependent factors that control the h and H couplings to Standard Model fermions and vector bosons. In (a) we show the boundaries, to the right of which the h couplings are within 10% of the Standard Model couplings. In (b) - (f) we show the contour plots of the individual coupling factors [11].

in Fig. 4. The combination of these regative searches for the neutral Higgs boson excludes almost the entire (m_h, m_A) mass region which is kinematically accessible at LEP I. The impact of radiative corrections on the search for the MSSM Higgs boson is much more dramatic at LEP II [25]. We note that because of kinematical reasons, there is very little hope of finding the charged Higgs boson of the MSSM at LEP II.

3.2. Radiative Corrections in the Nonminimal Supersymmetric Standard Model

The calculation of radiative corrections is much more involved in the case of NMSSM, because the Higgs boson spectrum is much richer. The corrections are also much more model dependent because of the presence of unknown trilinear couplings λ and k in the superpotential (6). However, the strategy for computing the one-loop radiative corrections in the effective potential approximation remains the same. As

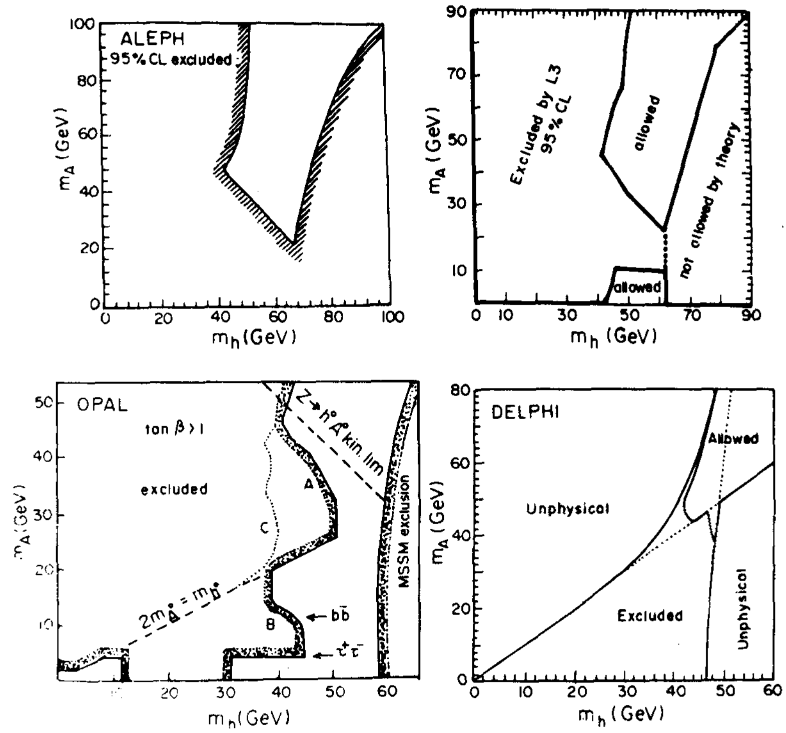

Figure 4. In the (m_h, m_A) plane of the MSSM with $\tan\beta > 1$, the domain excluded by the various LEP experiments at 95% CL. The ALEPH, L3 and OPAL excluded regions are for any value of $\tilde{m}$ and m_t. The excluded region by DELPHI is given for $m_t = 135$ GeV and $m_{stop} = 1$ TeV [26].

in the case of tree level result (20), we will be interested in the 2×2 radiatively corrected CP-even neutral Higgs mass matrix in order to obtain an upper bound on the lightest Higgs mass. The result [27] for the upper bound obtained in this manner is shown in Fig. 5, where we have chosen the fixed point value for $\lambda \simeq 0.87$ [5]. The corresponding result for the MSSM is also shown in the same diagram. We see that the upper bound on the mass of lightest Higgs lies substantially higher than the corresponding value in the MSSM for all values of m_t. We further note that the upper bound in NMSSM is independent of k. More important is the fact that in the limit of a large value of the singlet expectation value x, the upper bound depends only logarithmically on x. Thus, the lightest Higgs mass decomples from the large vacuum expectation value effects [27,28]. This is similar to the screening theorem for electroweak radiative corrections in the Standard Model [29], wherein the one-loop radiative corrections have only a logarithmic dependence on the Higgs mass. From Fig. 5 one obtains $m_h \lesssim 160$ GeV for $m_t < 150$ GeV. A somewhat

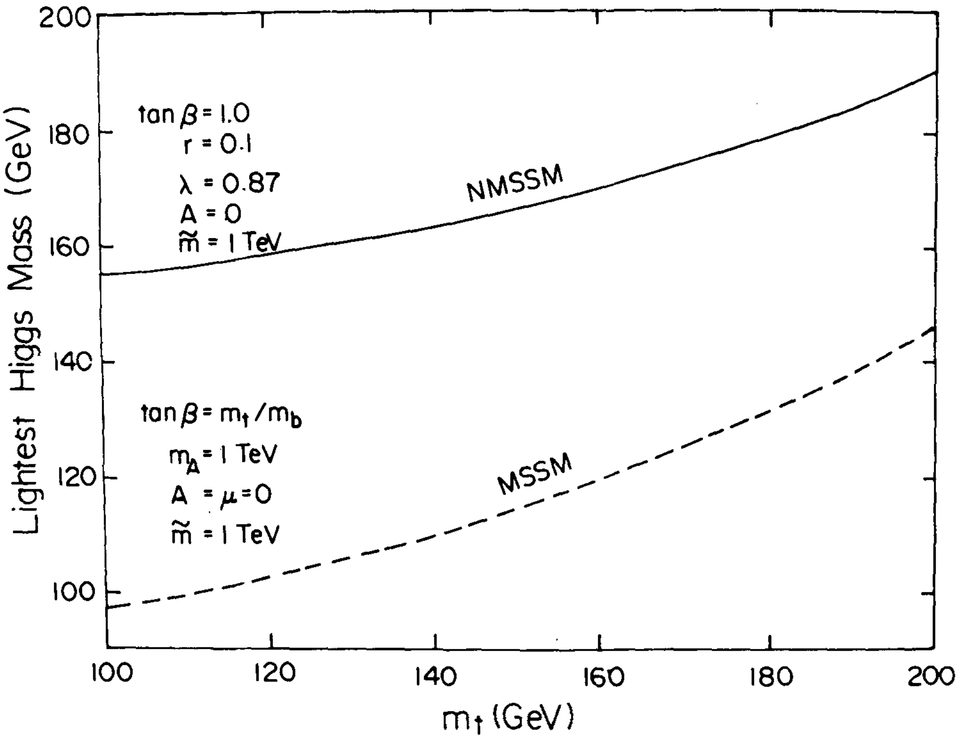

Figure 5. Upper bound on the mass of lightest Higgs in the Nonminimal Supersymmetric Standard Model with one Higgs singlet compared with the corresponding upper bound in the Minimal Supersymmetric Standard Model for the value of parameters shown in the figure [27].

smaller bound of $m_h \lesssim 140$ GeV for $m_t > 91$ GeV has been obtained in a model with an arbitrary number of gauge singlet Higgs superfields, within the framework of renormalization group equation approach, by assuming that the theory remains perturbature upto some large scale [30].

In conclusion, there are large radiative corrections to the lightest Higgs mass in both the minimal and nonminimal supersymmetric models. These corrections may push the lightest Higgs mass in the intermediate range 100 – 200 GeV. Thus, if the world is supersymmetric, the crucial mass limit for the Higgs search may be around 150 GeV rather than 800 GeV as in the case of Standard Model.

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Discussion

Probir Roy : Any comments on the pseduscalar mass?

P.N. Pandita : Assuming $\tan\beta > 1$ and including one-loop radiative corrections the mass of the pseudoscalar and the lightest scalar have to exceed 21 GeV/c^2 and 43 GeV/c^2 respectively.

Probir Roy : If m_t is measured, whatever be your parameters, there will always be some absolute upper bound on the lightest Higgs mass.

P.N. Pandita : Such an upper bound will depend on other parameters, namely A and μ etc., in the minimal supersymmetric model.

A. Khare : After radiative corrections is there a firm, testable prediction about the Higgs masses in minimal supersymmetric model?

P.N. Pandita : After the inclusion of radiative corrections, the lightest Higgs mass in the minimal supersymmetric model depends on a number of parameters. However, for a reasonable range of parameters, the lightest Higgs mass should be less than $110 - 120$ GeV/c^2 for m_t less than 150 GeV/c^2.