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On radiative corrections to supersymmetric Higgs boson masses and their implications for LEP searches

John Ellis

Theory Division, CERN, Geneva, Switzerland

Giovanni Ridolfi

INFN, Sezione di Genova, Italy

and

Fabio Zwirner¹

Theory Division, CERN, Geneva, Switzerland

Abstract

We present calculations of the one-loop radiative corrections to the mass of the neutral CP-odd Higgs boson (A) in the minimal supersymmetric extension of the Standard Model, as well as to the neutral CP-even Higgs (h, H) masses and mixing angles. We use these results to recalculate the cross-sections for Higgs production at LEP in the $h + (Z^* + f\bar{f})$, $h(H)Z^0$ and $h(H)A$ final states. We delineate the domains of parameter space accessible at LEP I at the Z^0 peak and at LEP II with a center-of-mass energy of 180, 190 or 200 GeV.

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¹On leave from Istituto Nazionale di Fisica Nucleare, Sezione di Padova, Italy.

The search for a Higgs boson is one of the highest priorities in the LEP experimental programme [1]. Quite stringent limits have already been obtained from unsuccessful searches at LEP I [2], and further progress is expected in the rest of the LEP I phase. At LEP II it should be possible to detect any neutral Higgs boson weighing less than or about the Z^0 [3], provided that its couplings to the Z^0 are not significantly suppressed with respect to the ones predicted by the Standard Model. The minimal supersymmetric extension of the Standard Model [4], with just two complex Higgs doublets $H_{1,2}$, has provided encouragement for this search, because its tree-level relations indicate that there always exists at least one neutral CP-even Higgs boson, weighing less than or about m_Z and with approximately standard couplings to the Z^0 [5]. Moreover, the associated production of the CP-even (h) and the CP-odd (A) Higgses offers another possibility for detection.

However, it has recently been pointed out [6-11] that there are large radiative corrections to the masses of the neutral CP-even Higgs bosons, h and H , due to loops involving the t quarks and $\bar{t}$ squarks that couple to the H_2 doublet. Numerically, for m_t in the range 90 to 180 GeV, indicated by direct searches [12] and by analyses of precision electroweak measurements at LEP and elsewhere [13], the radiative corrections to m_h and m_H can be very significant: they can increase m_h (and/or m_H) by several dozen GeV, implying that m_h could exceed 50 GeV even if $\tan\beta \equiv v_2/v_1 \sim 1$, and possibly pushing m_h beyond m_Z .

These results clearly have dramatic implications for Higgs searches at LEP [11]. The Higgs production processes of most interest at LEP are: $e^+e^- \rightarrow Z^* \rightarrow h + (Z^* \rightarrow f\bar{f})$, $H + (Z^* \rightarrow f\bar{f})$, $h + A$, $H + A$. The first two processes involve the trilinear ZZh (or ZZH) coupling, whereas the last the cross-sections for these processes requires, in addition to the radiatively-corrected values of m_h and m_H , calculating the radiative corrections to m_A and reevaluating the $ZZh(H)$ and $Zh(H)A$ couplings, which we do in this paper.

We first calculate the one-loop radiatively-corrected squared-mass matrix in the Higgs CP-odd sector, which includes terms depending on the amount of left-right squark mixing. We then give an analytic expression for the Higgs squared-mass matrix in the CP-even sector in terms of the one-loop radiatively-corrected mass m_A . We also discuss the renormalization-scale-independence of our one-loop expressions for $m_{h,H}$ and m_A . We fi-

nally present contours in the $(m_A, \tan\beta)$ plane for the branching ratios $B(Z^0 \rightarrow h f f)$, $B(Z^0 \rightarrow h A)$, and for the cross sections $\sigma(e^+e^- \rightarrow h Z^0)$, $\sigma(e^+e^- \rightarrow h A)$ at LEP II. In the latter case, we indicate explicitly the increased coverage of parameter space that would result from increasing the center-of-mass energy from 180 to 190 or 200 GeV.

We assume the minimal supersymmetric extension of the Standard Model [4], which contains two complex Higgs doublets with the following $SU(3) \times SU(2) \times U(1)$ quantum numbers:

$$H_1 \equiv \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix} \sim (1, 2, -1/2), \quad H_2 \equiv \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} \sim (1, 2, +1/2). \quad (1)$$

The starting point for a discussion of gauge symmetry breaking and Higgs boson masses is the tree-level Higgs potential

$$V_0 = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_3^2 (H_1 H_2 + \text{h.c.}) + \frac{1}{8} g^2 (H_1^\dagger \bar{\sigma} H_2 + H_1^\dagger \sigma H_1)^2 + \frac{1}{8} g'^2 (|H_2|^2 - |H_1|^2)^2, \quad (2)$$

where for simplicity we have taken m_3^2 to be real and we have chosen a field basis such that $m_3^2 \leq 0$. In this case it is not restrictive to assume that the vacuum expectation values of H_1^0 and H_2^0 , denoted by v_1^0 and v_2^0 , are both real and non-negative. After introducing the auxiliary variables

$$v^0 \equiv v_1^0 + v_2^0, \quad \tan\beta^0 \equiv \frac{v_2^0}{v_1^0}, \quad (3)$$

the tree-level minimization conditions

$$\sin 2\beta^0 = -\frac{2m_3^2}{m_1^2 + m_2^2}, \quad v^0 = \frac{4}{g^2 + g'^2} \frac{m_1^2 - m_2^2 \tan^2 \beta^0}{\tan^2 \beta^0 - 1}, \quad (4)$$

combined with the tree-level expression for the Z mass,

$$m_Z^2 = \frac{1}{2}(g^2 + g'^2)v^0, \quad (5)$$

give for the masses of the CP-even neutral Higgses

$$m_{h,H}^0 = \frac{1}{2} \left[m_A^0 + m_Z^0 \mp \sqrt{(m_A^0 + m_Z^0)^2 - 4m_A^0 m_Z^0 \cos^2 2\beta^0} \right], \quad (6)$$

where

$$m_A^{0,2} = -\frac{2m_3^2}{\sin 2\beta^0} \quad (7)$$

is the tree-level expression for the CP-odd neutral Higgs mass. Equation (6) immediately implies the inequalities

$$m_h^0 < m_Z |\cos 2\beta^0| < m_Z < m_H^0 \quad (8)$$

and

$$m_h^0 < m_A^0 |\cos 2\beta^0| < m_A^0 < m_H^0. \quad (9)$$

Moreover, one can express $\tan \beta^0$ as a function of the Higgs boson masses m_h^0, m_A^0 with a two-fold ambiguity, which is removed if one requires $\tan \beta^0 \geq 1$, as suggested by model building. It has already been observed in refs. [6-8] that the upper bounds (8) on m_h^0 do not apply to the one-loop radiatively-corrected mass m_h : we will now see that the same can be said about the lower bounds (9) on m_A^0 , and that more than one value of $\tan \beta \geq 1$ can correspond to a given (m_h, m_A) pair.

Our approach to the radiative corrections to supersymmetric Higgs boson masses proceeds from the full one-loop effective potential:

$$V_1(Q) = V_0(Q) + \Delta V_1(Q) \quad (10)$$

$$\Delta V_1(Q) = \frac{1}{64\pi^2} \text{Str } \mathcal{M}^4 \left(\log \frac{\mathcal{M}^2}{Q^2} - \frac{3}{2} \right), \quad (11)$$

where $V_0(Q)$ is the tree-level potential evaluated with couplings renormalized at some scale Q , Str denotes the conventional supertrace and field-independent terms have been neglected. As discussed in ref. [7], the radiatively-corrected Higgs mass-squared matrix is given, to a good approximation, by the matrix of the second derivatives of V_1 with respect to the Higgs fields. In the neutral Higgs sector, one has a 4×4 matrix of second derivatives of V_1 , that breaks up into a pair of 2×2 matrices for the CP-even and CP-odd sectors, respectively, when CP is conserved. It has also been pointed out in refs. [6-8] that the dominant contributions to the supertrace in eq. (11) is that due to the top-stop system. However, for very large values of $\tan \beta$, also the bottom-bottom contributions can be non-negligible. Moreover, if D-terms are not neglected, $\text{Str } \mathcal{M}^2$ is field-independent only when all the

above contributions are included. We will then include both the top-stop and the bottom-sbottom contributions in the rest of the paper.

It is easy to see that the pseudoscalar mass m_A is related to the dependence of V_1 on the relative phases of the two neutral Higgs components, H_1^0 and H_2^0 . There can be a one-loop contribution to m_A only if there is mixing between the scalar partners of the left- and right-handed top and bottom quarks, $(\tilde{t}_L, \tilde{t}_R)$ and $(\tilde{b}_L, \tilde{b}_R)$. The 2×2 $\tilde{t}_L - \tilde{t}_R$ and $\tilde{b}_L - \tilde{b}_R$ mass-squared matrices are (keeping only the dependence on H_1^0 and H_2^0):

$$\begin{bmatrix} m_Q^2 + h_t^2 |H_2^0|^2 + \left(g^2 - \frac{1}{3}g'^2\right) \frac{|H_1^0|^2 - |H_2^0|^2}{4} & h_t(A_t H_2^0 + \mu H_1^{0*}) \\ h_t(A_t H_2^{0*} + \mu H_1^0) & m_Q^2 + h_t^2 |H_2^0|^2 + g'^2 \frac{|H_1^0|^2 - |H_2^0|^2}{3} \end{bmatrix}, \quad (12)$$

and

$$\begin{bmatrix} m_D^2 + h_b^2 |H_1^0|^2 - \left(g^2 + \frac{1}{3}g'^2\right) \frac{|H_1^0|^2 - |H_2^0|^2}{4} & h_b(A_b H_1^0 + \mu H_2^{0*}) \\ h_b(A_b H_1^{0*} + \mu H_2^0) & m_D^2 + h_b^2 |H_1^0|^2 - g'^2 \frac{|H_1^0|^2 - |H_2^0|^2}{6} \end{bmatrix}, \quad (13)$$

respectively, where A_t and A_b are the conventional trilinear soft supersymmetry-breaking parameters

$$V_{\phi\phi t} = h_t A_t \tilde{t}_L \tilde{t}_R^c H_2^0 + h_b A_b \tilde{b}_L \tilde{b}_R^c H_1^0 + \text{h.c.} + \dots, \quad (14)$$

μ is the coefficient of the $H_1 H_2$ mixing term in the superpotential, and the top and bottom quark masses squared are given by $h_t^2 |H_2^0|^2$ and $h_b^2 |H_1^0|^2$, respectively. We assume that A_t, A_b and μ are real, and the symbols m_Q, m_U and m_D stand for field-independent soft supersymmetry-breaking squark masses. Using eqs. (12) and (13) in $\Delta V_1(Q)$ as given by eq. (11), it is easy to calculate the radiatively-corrected expression for the squared masses in the CP-odd neutral sector. After imposing the appropriate one-loop minimization conditions, we find the following form for the 2×2 matrix of second derivatives of $V_1(Q)$ with respect to $\varphi_i \equiv \text{Im } H_i^0$:

$$\left[\frac{\partial^2 V_1}{\partial \varphi_i \partial \varphi_j} \right]_{v_1, v_2} = \begin{bmatrix} \tan \beta & 1 \\ 1 & \cot \beta \end{bmatrix} \Delta, \quad (15)$$

where

$$\Delta = -2m_3^2 - \frac{3g^2}{32\pi^2 \sin^2 \beta} \frac{m_t^2}{m_W^2} \frac{A_t \mu}{m_t^2 - m_{\tilde{t}_1}^2} \left[f(m_{\tilde{t}_1}^2) - f(m_{\tilde{t}_2}^2) \right]$$

$$- \frac{3g^2}{32\pi^2 \cos^2 \beta} \frac{m_b^2}{m_W^2 m_{b_1}^2 - m_{b_2}^2} \left[f(m_{b_1}^2) - f(m_{b_2}^2) \right] \quad (16)$$

and

$$f(m^2) \equiv 2m^2 \left(\log \frac{m^2}{Q^2} - 1 \right). \quad (17)$$

The four eigenvalues $m_{i,2}^2, m_{b,2}^2$ of the $\tilde{t}, \tilde{b}$ mass-squared matrices are given by

$$\begin{aligned} m_{i,2}^2 &= m_t^2 + \frac{1}{2}(m_Q^2 + m_U^2) + \frac{1}{4}m_Z^2 \cos 2\beta \\ &\pm \sqrt{\left[\frac{1}{2}(m_Q^2 - m_U^2) + \frac{1}{12}(8m_W^2 - 5m_Z^2) \cos 2\beta \right]^2 + m_t^2 (A_t + \mu \cot \beta)^2}, \end{aligned} \quad (18)$$

$$\begin{aligned} m_{b,2}^2 &= m_b^2 + \frac{1}{2}(m_Q^2 + m_D^2) - \frac{1}{4}m_Z^2 \cos 2\beta \\ &\pm \sqrt{\left[\frac{1}{2}(m_Q^2 - m_D^2) - \frac{1}{12}(4m_W^2 - m_Z^2) \cos 2\beta \right]^2 + m_b^2 (A_b + \mu \tan \beta)^2}. \end{aligned} \quad (19)$$

The determinant of the matrix (15) vanishes, corresponding to the masslessness of the Goldstone boson eaten by the Z^0 in the Higgs mechanism. Furthermore, the null eigenvector has the same expression, as a function of $\tan \beta$, as the tree-level Goldstone mode has as a function of $\tan \beta^0$. The mass of the remaining non-null mode is given by

$$m_A^2 = \frac{\Delta}{\sin 2\beta}. \quad (20)$$

We note that the expression (20) for m_A^2 has the same form as the tree-level one if either $\mu = 0$ or $A_t = A_b = 0$, as one would expect from the structure of the mass-squared matrices (12,13). The explicit Q -dependence in the one-loop expression (20) for m_A^2 is actually cancelled by the implicit Q -dependence of m_t^2 and $\tan \beta$ to leading order, so that the physical value of m_A^2 can be renormalization-scale-independent, as it should be. Residual

Q -dependence in the expression (20) is a two-loop effect that would cancel in a full calculation of m_A using the two-loop effective potential.

To specify quantities at the one-loop level, one should use renormalized parameters as inputs. We choose to use $\tan \beta$ and m_A , since all dependent variables, including the CP-even Higgs boson masses and mixing angle, can be expressed unambiguously as functions of $(m_A, \tan \beta)$, and of the parameters of the top-stop and bottom-bottom sectors. As already mentioned, there would be an unavoidable ambiguity in $\tan \beta$ if one chose as independent variables m_A and m_A .

We now proceed to derive an analytic expression for the squared-mass matrix of the CP-even mass eigenstates in terms of the one-loop corrected mass m_A as given by eq. (20). Defining $\psi_i \equiv \text{Re} H_i^0$, we find

$$\begin{aligned} \left[\frac{\partial^2 V_1}{\partial \psi_i \partial \psi_j} \right]_{v_1, v_2} &= \begin{bmatrix} \cot \beta & -1 \\ -1 & \tan \beta \end{bmatrix} m_Z^2 \sin 2\beta + \begin{bmatrix} \tan \beta & -1 \\ -1 & \cot \beta \end{bmatrix} \Delta \\ &\quad + \frac{3g^2}{8\pi^2 m_W^2} \begin{bmatrix} \Delta_{11} & \Delta_{12} \\ \Delta_{12} & \Delta_{22} \end{bmatrix}, \end{aligned} \quad (21)$$

where Δ is defined in eq. (16),

$$\begin{aligned} \Delta_{11} &= \frac{m_b^4}{\cos^2 \beta} \left[\log \frac{m_b^2 m_{b_2}^2}{m_s^4} + \frac{2A_b(A_b + \mu \tan \beta)}{m_{b_1}^2 - m_{b_2}^2} \log \frac{m_{b_1}^2}{m_{b_2}^2} \right] \\ &\quad + \frac{m_b^4}{\cos^2 \beta} \left[\frac{A_b(A_b + \mu \tan \beta)}{m_{b_1}^2 - m_{b_2}^2} \right]^2 g(m_{b_1}^2, m_{b_2}^2) \\ &\quad + \frac{m_t^4}{\sin^2 \beta} \left[\frac{\mu(A_t + \mu \cot \beta)}{m_{t_1}^2 - m_{t_2}^2} \right]^2 g(m_{t_1}^2, m_{t_2}^2), \end{aligned} \quad (22)$$

$$\begin{aligned} \Delta_{22} &= \frac{m_t^4}{\sin^2 \beta} \left[\log \frac{m_{t_1}^2 m_{t_2}^2}{m_t^4} + \frac{2A_t(A_t + \mu \cot \beta)}{m_{t_1}^2 - m_{t_2}^2} \log \frac{m_{t_1}^2}{m_{t_2}^2} \right] \\ &\quad + \frac{m_t^4}{\sin^2 \beta} \left[\frac{A_t(A_t + \mu \cot \beta)}{m_{t_1}^2 - m_{t_2}^2} \right]^2 g(m_{t_1}^2, m_{t_2}^2) \\ &\quad + \frac{m_b^4}{\cos^2 \beta} \left[\frac{\mu(A_b + \mu \tan \beta)}{m_{b_1}^2 - m_{b_2}^2} \right]^2 g(m_{b_1}^2, m_{b_2}^2), \end{aligned} \quad (23)$$

$$\Delta_{12} = \frac{m_{t_1}^4}{\sin^2 \beta} \frac{\mu(A_t + \mu \cot \beta)}{m_{t_1}^2 - m_{t_2}^2} \left[\log \frac{m_{t_1}^2}{m_{t_2}^2} + \frac{A_t(A_t + \mu \cot \beta)}{m_{t_1}^2 - m_{t_2}^2} g(m_{t_1}^2, m_{t_2}^2) \right] \\ + \frac{m_{b_1}^4}{\cos^2 \beta} \frac{\mu(A_b + \mu \tan \beta)}{m_{b_1}^2 - m_{b_2}^2} \left[\log \frac{m_{b_1}^2}{m_{b_2}^2} + \frac{A_b(A_b + \mu \tan \beta)}{m_{b_1}^2 - m_{b_2}^2} g(m_{b_1}^2, m_{b_2}^2) \right], \quad (24)$$

and

$$g(m_1^2, m_2^2) = 2 - \frac{m_1^2 + m_2^2}{m_1^2 - m_2^2} \log \frac{m_1^2}{m_2^2}. \quad (25)$$

Using eqs. (21-24), it is easy to compute $m_{h,H}^2$ and the mixing angle α , which specifies the eigenvectors of the matrix (21), as functions of $m_A^2 = \Delta/\sin 2\beta$. The angle α (together with β) determines all the couplings of the Higgs bosons. Observe that, when $A_t = A_b = \mu = 0$ and the bottom contributions is neglected, eq. (21) reduces to the analytic expressions that can be found, for example, in refs. [7] and [11]. The terms proportional to A_t, A_b and μ in eqs. (22-24) are obtained neglecting the D-terms in the squark masses. Actually, the D-terms should be neglected everywhere in our formulae in order to gain approximate independence of the renormalization scale Q , since we are including only the quark-squark contributions to ΔV_1 . One can see from eqs. (18,19) that the D-term contributions to squark masses are indeed very small, especially when the soft supersymmetry-breaking mass parameters are large with respect to m_W and m_Z , and/or $\tan \beta$ is not too far from 1.

Fig. 1 shows the (m_h, m_A) plane for various different choices of the top mass, the soft terms in the squark mass matrices, and the amount of left-right squark mixing. In each panel we show contour curves of $\tan \beta$, and indicate by a dotted line the boundary of the region $m_h \leq m_A, m_Z$ allowed at the tree level. For simplicity, we take m_Q, m_U and m_D all equal to a common mass parameter, which we call m_t . Figs. 1a and 1c show cases for which the radiative corrections are relatively small: $m_t = 120$ GeV is in the lower part of the range indicated by electroweak data. We see that there is no lower bound on either m_h or m_A , that $m_h > m_A$ is possible (although only if both are small), and that m_h never gets much larger than m_Z . Fig. 1b shows the effect of increasing m_t to 160 GeV. Again there is no lower bound on m_h or m_A , but $m_h > 53$ GeV for $\tan \beta = 1$, there is a

large region where $m_h > m_A$, and m_h can be as large as 120 GeV¹. Finally, fig. 1d shows the effect of including left-right mixing in the squark mass matrices, in the case $m_t = 160$ GeV and $m_{\tilde{q}} = 1000$ GeV. We have chosen $A_t = A_b = 1000$ GeV and $\mu = 200$ GeV. Comparing figs. 1b and 1d, we can see that the effect of left-right mixing is most important for small m_h , and relatively unimportant for $m_h \sim m_Z$. Notice that there are sometimes two values of $\tan \beta$ that yield the same values of m_h and m_A , which does not happen at the tree level. For this reason, we have chosen not to present our other results in the (m_h, m_A) plane favoured in some tree-level studies. Notice also the following general pattern of the shifts in m_h with respect to the tree-level prediction. For $\tan \beta \sim 1$ the shift is large and approximately independent of m_A , whilst for very large $\tan \beta$ the shift is negligible as long as $m_A < m_Z$ (in this case m_H receives the most important shift), and rapidly tends to a constant asymptotic value for increasing $m_A > m_Z$.

Knowing the one-loop values of the physical h and A masses and the mixing angle α , we can make revised estimates of the Higgs production cross-sections at LEP. Fig. 2 exhibits branching ratios for $Z^0 \rightarrow h + (Z^* \rightarrow f\bar{f})$ and $Z^0 \rightarrow hA$ for $m_t = 140$ GeV, $m_{\tilde{q}} = 1000$ GeV, $A_t = A_b = \mu = 0$ or $A_t = A_b = 1000$ GeV, $\mu = 200$ GeV. Searches for these Z^0 decay modes have been used in the past by the LEP experiments [2] to give excluded domains in the (m_h, m_A) plane, using the tree-level formulae for the masses and couplings. The large size of the one-loop radiative corrections to the Higgs masses requires revision of the experimental exclusion contours in the (m_h, m_A) plane, even if we do not expect substantial modifications to the absolute lower bounds on m_h and m_A . In the future, it would be helpful if the experimental results were plotted as contours of upper limits on the $Z^0 \rightarrow h + (Z^* \rightarrow f\bar{f})$ and $Z^0 \rightarrow hA$ branching ratios. In the mass region of present interest for LEP, radiative corrections should not substantially affect the branching ratios for the most important h, A decay modes ($b\bar{b}, \tau^+\tau^-$). It would then be possible, for any given values of the relevant parameters, to infer an exclusion domain in the $(m_h, \tan \beta)$ plane. We recall that the (m_h, m_A) plane is less useful at the one-loop level. If it is nevertheless preferred to present the experimental data as exclusion contours in the (m_h, m_A) plane, one should take the $\tan \beta$

¹One should notice, however, that the simultaneous existence of $\tan \beta \sim 1$ and a very large m_t (> 140 GeV) is disfavoured by supersymmetric grand unification, both because radiative symmetry breaking would require a fine-tuning of the parameters and because the top Yukawa coupling would develop a Landau pole below the grand-unification scale.

solution which corresponds to the more conservative values of the couplings, varying m_t and the soft supersymmetry-breaking parameters over plausible ranges. If one assumes that each LEP experiment will eventually have a sample of several million Z^0 decays, one could take the $B = 10^{-6}$ contours as useful guides for the eye: establishing a plausible limit of detectability would require an experimental discussion which goes beyond the aim of the present paper.

Turning now to LEP II, we plot in fig. 3 contours of $\sigma(e^+e^- \rightarrow hZ^0)$ and $\sigma(e^+e^- \rightarrow hA)$, for the same choices of parameters as in fig. 2. The solid lines are contours for cross-sections at a center-of-mass energy $\sqrt{s} = 180$ GeV. If one assumes an integrated LEP II luminosity of 500 pb^{-1} , one could take the $\sigma = 0.01 \text{ pb}$ contours as reference lines: again, and even more so than at LEP I, establishing a plausible limit of detectability would require a discussion of backgrounds, efficiencies, b-tagging and other experimental problems. Since there is considerable interest in the possibility of increasing the LEP II center-of-mass energy, we also show in fig. 3 the $\sigma = 0.01 \text{ pb}$ contours for $\sqrt{s} = 190$ GeV (dashed lines) and 200 GeV (dotted lines).

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Figure Captions

Figure 1

Contour curves of $\tan \beta$ in the (m_h, m_A) plane for (a) $m_t = 120$ GeV, $m_{\tilde{q}} = 1000$ GeV, $A_t = A_b = \mu = 0$; (b) $m_t = 160$ GeV, $m_{\tilde{q}} = 1000$ GeV, $A_t = A_b = \mu = 0$; (c) $m_t = 120$ GeV, $m_{\tilde{q}} = 150$ GeV, $A_t = A_b = \mu = 100$ GeV; (d) $m_t = 160$ GeV, $m_{\tilde{q}} = 1000$ GeV, $A_t = A_b = 1000$ GeV, $\mu = 200$ GeV. The boundary of the region $m_h < m_Z$, $m_A > m_h$ allowed at tree level is also shown (dotted lines).

Figure 2

(a) Branching ratio for the decay $Z^0 \rightarrow hZ^{0*}$ in the $(m_A, \tan \beta)$ plane, for $m_t = 140$ GeV, $m_{\tilde{q}} = 1$ TeV and $A_t = A_b = \mu = 0$. (b) Same as fig. 2a, with $A_t = A_b = 1$ TeV and $\mu = 200$ GeV. (c) Branching ratio for the decay $Z^0 \rightarrow hA$ in the $(m_A, \tan \beta)$ plane, for $m_t = 140$ GeV, $m_{\tilde{q}} = 1$ TeV and $A_t = A_b = \mu = 0$. (d) Same as fig. 2c, with $A_t = A_b = 1$ TeV and $\mu = 200$ GeV.

Figure 3

(a) Cross section (in pb) for $e^+e^- \rightarrow hZ^0$ in the $(m_A, \tan \beta)$ plane for $\sqrt{s} = 180$ GeV, $m_t = 140$ GeV, $m_{\tilde{q}} = 1$ TeV and $A_t = A_b = \mu = 0$. The dashed and dotted lines correspond to the $\sigma = 0.01$ pb at $\sqrt{s} = 190$ and 200 GeV, respectively. (b) Same as fig. 3a, with $A_t = A_b = 1$ TeV and $\mu = 200$ GeV. (c) Cross section (in pb) for $e^+e^- \rightarrow hA$ in the $(m_A, \tan \beta)$ plane for $\sqrt{s} = 180$ GeV, $m_t = 140$ GeV, $m_{\tilde{q}} = 1$ TeV and $A_t = A_b = \mu = 0$. The dashed and dotted lines correspond to the $\sigma = 0.01$ pb at $\sqrt{s} = 190$ and 200 GeV, respectively. (d) Same as fig. 3c, with $A_t = A_b = 1$ TeV and $\mu = 200$ GeV.

- Figure 1 -

- Figure 2 -

- Figure 3 -