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Exploring the SUSY Higgs Sector at e^+e^- Linear Colliders: A Synopsis

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ABSTRACT

We discuss the Higgs scenario in the minimal supersymmetric extension of the Standard Model ($\mathcal{MSSM}$) at e^+e^- linear colliders operating in the c.m. energy range between 300 and 500 GeV. Besides decays of the Higgs particles into ordinary fermions and cascade decays, we analyze also decays into gaugino/Higgsinos and in particular, neutral Higgs decays into the lightest supersymmetric particles which are invisible if R-parity is conserved. The cross sections for the various production channels of $\mathcal{SUSY}$ Higgs particles in e^+e^- collisions are discussed in detail. The lightest Higgs boson cannot escape detection, and in major parts of the $\mathcal{MSSM}$ parameter space all five Higgs particles can be observed.

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1 Introduction

Supersymmetric theories are very attractive extensions of the Standard Model ($\mathcal{SM}$), incorporating the most general symmetry of the $\mathcal{S}$ matrix in field theory. At low energies they provide a theoretical framework in which the problem of naturalness and hierarchy in the Higgs sector is solved while retaining Higgs bosons with moderate masses as elementary particles in the context of high mass scales demanded by grand unification. The Minimal Supersymmetric extension of the Standard Model ($\mathcal{MSSM}$) [1, 2] may serve as a useful guideline into this domain. This point is underlined by the fact that the model led to a prediction of the electroweak mixing angle [3] that is in very nice agreement with present high-precision measurements of $\sin^2\theta_W$. Though some of the phenomena will be specific to this minimal version, the general pattern will nevertheless be characteristic to more general extensions [2, 4, 5] so that the analyses can be considered as representative for a wide class of $\mathcal{SUSY}$ models.

The $\mathcal{MSSM}$ requires the existence of two isodoublets of scalar Higgs fields giving mass separately to up and down-type fermions. Three neutral [h/H with $\mathcal{CP} = +$ and A with $\mathcal{CP} = -$] and a pair of charged scalar particles [$H^\pm$] are introduced by this extension of the Higgs sector. At tree level, the mass of the lightest Higgs boson h is smaller than the Z mass. Radiative corrections [6] however, which grow as the fourth power of the top quark mass, can shift the upper limit from M_Z to ~ 140 GeV. Beyond ~ 90 GeV, Higgs particles are not accessible anymore at LEP200 and higher energies are required to search for these particles. The masses of the heavy neutral and of the charged Higgs particles are expected in the range of the scale of electroweak symmetry breaking.

Apart from cascade decays in some corners of the parameter space, the main decay modes of the neutral Higgs particles are in general $b\bar{b}$ decays [$\sim 90\%$] and $\tau^+\tau^-$ decays [$\sim 10\%$] which are easy to detect experimentally in e^+e^- colliders [7]. The gold-plated ZZ decays of the $\mathcal{SM}$ Higgs above 140 GeV play only a minor role in the $\mathcal{SUSY}$ Higgs sector — and in large parts of the parameter space their role is even negligible. Charged Higgs particles decay predominantly into $\tau\nu_\tau$ and $t\bar{b}$ pairs [primarily into the latter channel, if kinematically allowed]. The total width of the states remains small, $\mathcal{O}$ (1 GeV) anywhere in the intermediate mass range. In addition to these conventional decays, the Higgs particles may also decay into chargino and neutralino pairs. Depending on the details of the $\mathcal{SUSY}$ parameters, the branching ratios for decays into these channels can add up to a few tens of percent. When kinematically allowed, the Higgs particles also decay into squarks and sleptons.

The neutral Higgs particles will be searched for mainly in the decay channels $\tau^+\tau^-$ and $\gamma\gamma$ at the hadron colliders LHC/SSC [8] since the large QCD jet background does not allow the analysis of the dominating $b\bar{b}$ final states. Charged Higgs particles appear to be accessible in these colliders only through top quark decays.

At e^+e^- colliders, $\mathcal{SUSY}$ Higgs particles can be created through a variety of mechanisms, covering a large part of the mass parameter space at a c.m. energy of 500 GeV. Besides the classical bremsstrahlung and fusion processes, $e^+e^- \to Z+h/H$ and $e^+e^- \to \bar{\nu}\nu/e^+e^-+h/H$, neutral and charged Higgs particles can also be produced pairwise, $e^+e^- \to A + h/H$ and

$e^+e^- \to H^+ + H^-$. The cross sections for the light and the heavy $\mathcal{CP}$–even neutral Higgs bosons h/H add up to the cross section for the $\mathcal{SM}$ Higgs particle, while their relative weights can vary between wide margins [9, 10]. Nevertheless, combining all channels at least the lightest neutral Higgs particle must be found at e^+e^- colliders. If not, the minimal supersymmetric extension of the $\mathcal{SM}$ would be ruled out. Furthermore, in a large part of the parameter space *all* three neutral Higgs particles and the charged Higgs particle can be found. e^+e^- colliders are therefore ideal machines to analyze the $\mathcal{SUSY}$ Higgs sector.

The paper is organized as follows. After a brief summary of the mass parameters and couplings, including radiative corrections, we shall discuss the decay modes of the Higgs particles into $\mathcal{SM}$ fermions and gauge bosons as well as cascade decays. We will also present a comprehensive analysis of Higgs decays into charginos and neutralinos. Finally we will discuss the production cross sections of neutral and charged Higgs particles at e^+e^- linear colliders with c.m. energies up to 500 GeV. For the sake of completeness, the new material worked out in this paper will be mixed with updates on material discussed already in the past, and it will be embedded into a comprehensive description of $\mathcal{SUSY}$ Higgs physics at e^+e^- colliders.

2 Mass Parameters and Couplings

Supersymmetry requires the existence of at least two isodoublet scalar fields Φ_1 and Φ_2, thus extending the physical spectrum of scalar particles to five. The $\mathcal{MSSM}$ is restricted to this minimal extension. The field Φ_2 [with vacuum expectation value v_2] couples only to up–type quarks while Φ_1 [with vacuum expectation value v_1] couples to down–type quarks and charged leptons. The physical Higgs bosons introduced by this extension are of the following type: two $\mathcal{CP}$–even neutral bosons h and H [where h will be the lightest particle], a $\mathcal{CP}$–odd neutral boson A [usually called pseudoscalar] and two charged Higgs bosons $H^\pm$. Besides the four masses M_h, M_H, M_A and $M_{H^\pm}$, two additional parameters define the properties of the scalar particles and their interactions with gauge bosons and fermions: the ratio of the two vacuum expectation values $\mathrm{tg}\beta = v_2/v_1$ and a mixing angle α in the neutral $\mathcal{CP}$–even sector. Supersymmetry leads to several relations among these parameters and, in fact, only two of them are independent. These relations impose a strong hierarchical structure on the mass spectrum $[M_h < M_Z, M_A < M_H$ and $M_W < M_{H^\pm}]$ which however is broken by radiative corrections if the top quark mass is large [6]. The parameter $\mathrm{tg}\beta$ will in general be assumed in the range $1 < \mathrm{tg}\beta < \mathrm{m_t/m_b}$ $[\pi/4 < \beta < \pi/2]$, consistent with restrictions [11] that follow from interpreting the $\mathcal{MSSM}$ as low energy limit of a supergravity model.

Since the lightest $\mathcal{CP}$–even Higgs boson h is likely to be the particle which will be discovered first, an attractive choice of the two input parameters is the set $[M_h, \mathrm{tg}\beta]$, with $\mathrm{tg}\beta$ to be determined by the production cross sections. Once these two parameters [as well as the top quark mass and the associated squark masses which enter through radiative corrections] are specified, all other masses and the mixing angle α can be predicted. To incorporate radiative corrections we first shall neglect, for the sake of simplicity, non–leading effects due to non–zero values of the supersymmetric Higgs mass parameter μ and of the parameters A_t and A_b in the soft symmetry breaking interaction. The radiative corrections are then

determined by the parameter ϵ which grows as the fourth power of the top quark mass m_t and logarithmically with the squark mass M_S

$$\epsilon = \frac{3\alpha}{2\pi}\frac{1}{s_W^2 c_W^2}\frac{1}{\sin^2\beta}\frac{m_t^4}{M_Z^2}\log\left(1+\frac{M_S^2}{m_t^2}\right) \tag{1}$$

$[s_W^2 = 1 - c_W^2 \equiv \sin^2\theta_W]$. These corrections are positive and they shift the mass of the light neutral Higgs boson h upward with increasing top mass. The variation of the upper limit on M_h with the top quark mass is shown in Fig. 1a for $M_S = 1$ TeV and two representative values of $\mathrm{tg}\beta = 2.5$ and 20. While the dashed lines correspond to the leading radiative corrections in eq.(1) $[\mu = A_t = A_b = 0]$, the solid lines correspond to $\mu = -200, 0, +200$ GeV and $A_t = A_b = 1$ TeV. The upper bound on M_h is shifted from the tree level value M_Z up to ~ 140 GeV for $m_t = 180$ GeV.

Taking M_h and $\mathrm{tg}\beta$ as the base parameters, the mass of the pseudoscalar state A is given to leading order by

$$M_A^2 = \frac{M_h^2(M_Z^2 - M_h^2 + \epsilon) - \epsilon M_Z^2\cos^2\beta}{M_Z^2\cos^2 2\beta - M_h^2 + \epsilon\sin^2\beta} \tag{2}$$

The masses of the heavy neutral and charged Higgs bosons follow from the sum rules

$$M_H^2 = M_A^2 + M_Z^2 - M_h^2 + \epsilon \tag{3}$$

$$M_{H^\pm}^2 = M_A^2 + M_W^2 \tag{4}$$

In the subsequent discussion, we will assume for definiteness that $m_t = 140$ GeV and $M_S = 1$ TeV. For the two representative values of $\mathrm{tg}\beta$ introduced above, the masses M_A, M_H and $M_{H^\pm}$ are displayed in Fig. 1b–1d as a function of the light neutral Higgs mass M_h. Apart from the range near the upper limit of M_h for a given value of $\mathrm{tg}\beta$, the masses cluster in characteristic bands of 100 to 150 GeV for M_H and $M_{H^\pm}$, and up to ~ 100 GeV for M_A [similarly to M_h]. On general grounds, the masses of the heavy neutral and charged Higgs bosons are expected to be of the order of the electroweak symmetry breaking scale.

The mixing parameter α is determined by $\mathrm{tg}\beta$ and the Higgs masses,

$$\mathrm{tg}2\alpha = \mathrm{tg}2\beta\frac{M_A^2 + M_Z^2}{M_A^2 - M_Z^2 + \epsilon/\cos 2\beta} \qquad \left[-\frac{\pi}{2} \le \alpha \le 0\right] \tag{5}$$

The couplings of the various neutral Higgs bosons to fermions and gauge bosons will in general depend on the angles α and β. Normalized to the $\mathcal{SM}$ Higgs couplings, they are summarized in Table 1. The pseudoscalar particle A has no tree level couplings to gauge bosons, and its couplings to down (up) type fermions are (inversely) proportional to $\mathrm{tg}\beta$.

Φ	$g_{\Phi\bar{u}u}$	$g_{\Phi\bar{d}d}$	$g_{\Phi VV}$
H_{SM}	1	1	1
h	$\cos\alpha/\sin\beta$	$-\sin\alpha/\cos\beta$	$\sin(\beta-\alpha)$
H	$\sin\alpha/\sin\beta$	$\cos\alpha/\cos\beta$	$\cos(\beta-\alpha)$
A	$1/\mathrm{tg}\beta$	$\mathrm{tg}\beta$	0

Tab.1: Higgs couplings in the $\mathcal{MSSM}$ to fermions and gauge bosons relative to $\mathcal{SM}$ couplings.

Typical numerical values of these couplings are shown in Fig. 2 as a function of the lightest neutral Higgs boson mass and for two values of tgβ; the same set of μ and A_t, A_b values as in Fig. 1 has been chosen. The figure demonstrates that the dependence on the parameters μ and A_t, A_b is very weak and the leading radiative corrections provide a very good approximation. [This is also the case for the radiative corrections to the Higgs boson masses, Fig. 1, although the dependence on A_t, A_t is slightly stronger in this case, leading to a shift of a few GeV]. There is in general a strong dependence on the input parameters tgβ and M_h. The couplings to down (up) type fermions are enhanced (suppressed) compared to the $\mathcal{SM}$ Higgs couplings. If M_h is very close to its upper limit for a given value of tgβ, the couplings to fermions and gauge bosons are $\mathcal{SM}$ like. It is therefore very difficult to distinguish the Higgs sector of the $\mathcal{MSSM}$ from the $\mathcal{SM}$, if all other Higgs bosons are very heavy.

Charginos and neutralinos are expected to be the lightest supersymmetric particles. In general, they are mixtures of the [non-colored] gauginos and Higgsinos, the spin 1/2 supersymmetric partners of the gauge bosons and Higgs bosons. There are two chargino $\tilde{\chi}_i^\pm$ [$i = 1,2$] and four neutralino $\tilde{\chi}_i^0$ [$i = 1,..,4$] states, where $\tilde{\chi}_1^0$ will be assumed to be the stable lightest supersymmetric particle. The masses and the couplings of these particles are obtained by diagonalizing the charged and neutral mass matrices; the unitary matrices which diagonalize the mass matrices can be found in [12]. Interpreting the $\mathcal{MSSM}$ as the low-energy limit of a supergravity model, the matrix elements will depend on four parameters, one Higgs mass [M_h or M_A], tgβ, μ and the SU(2) gaugino mass M which, without loss of generality, can be taken to be positive. [The fifth $\mathcal{SUSY}$ parameter of the model is the universal mass parameter m_0 of the scalar particles at the unification scale.]

The masses of the two lightest neutralinos and of the lightest chargino are shown in Fig. 3 in the [μ, M] plane for the two values of tgβ = 2.5 and 20. The contours are given for $M_{\tilde{\chi}_1^0} = 30, 60$ and 90 GeV and $M_{\tilde{\chi}_2^0}, M_{\tilde{\chi}_1^+} = 50, 100$ and 150 GeV. These particles can be produced at 500 GeV e^+e^- colliders in large areas of the [μ, M] plane. Also shown in this figure are the contours of the masses which are excluded by the negative search of supersymmetric particles in Z decays as well as the mass bounds which can be probed at LEP200. From Z decay data [21], the lightest neutralino [$\tilde{\chi}_1^0$] mass is restricted to be larger than 20 GeV for tgβ = 2.5 and larger than 22 GeV for tg$\beta > 4$; the second lightest neutralino [$\tilde{\chi}_2^0$] and the charginos are excluded if their masses are less than $\sim M_Z/2$. If the search at LEP200 with a c.m. energy of 180 GeV is negative, charginos with masses $m_{\tilde{\chi}_1^+} < 90$ GeV will also be excluded.

The present LEP data also exclude sleptons with masses below $\sim M_Z/2$; if sleptons will not be observed at LEP200 these limits can be improved by roughly a factor of two. On the other hand, CDF data restrict the squarks masses to be larger than $\sim$ 150 GeV if cascade decays are suppressed. Since the sfermion masses depend on the universal mass parameter m_0, they are not fixed by the mass scales which appear in the chargino/neutralino mass matrices. We shall assume in our analysis that squarks and sleptons are heavier than $\sim$ 150 GeV so that they will not affect Higgs boson decays in the mass range studied here.

3 Decay Modes

3.1 Decays to Standard Particles

Since its mass is smaller than $\sim$ 140 GeV, the lightest <u>neutral</u> Higgs boson will decay mainly into fermion pairs. This is also the dominant decay mode of the pseudoscalar boson A which has no tree-level couplings to gauge bosons. The partial decay width of a neutral Higgs boson Φ into fermion pairs is given by

$$\Gamma(\Phi \to \bar{f}f) \;=\; N_c \frac{G_F m_f^2}{4\sqrt{2}\pi}\, g_{\Phi ff}^2\, M_\Phi\, \beta^p \tag{6}$$

where $\beta = (1 - 4m_f^2/M_\Phi^2)^{1/2}$ and $p = 3(1)$ for the $\mathcal{CP}$-even (odd) Higgs boson; the couplings $g_{\Phi ff}$ are listed in Tab.1. For quarks, one has to use the running masses which take into account the bulk of the QCD corrections [14]. [In the case of bottom quarks, the b mass $m_b(m_b^2) = 4.2$ GeV drops to the effective value $m_b(M_\Phi^2) = 3$ GeV for $M_\Phi = 120$ GeV so that the partial decay width $\Gamma(\Phi \to b\bar{b})$ decreases by roughly a factor of two]. For values of tgβ larger than unity and for masses less than $\sim$ 130 GeV, the main decay modes of the neutral Higgs bosons will be decays into $b\bar{b}$ and $\tau^+\tau^-$ pairs. These branching ratios are larger than $\sim$ 90% and 5%, respectively. The decays into $c\bar{c}$ pairs are in general strongly suppressed, especially for large values of tgβ. For large masses, the top decay channels $H, A \to t\bar{t}$ open up; yet this mode remains suppressed for large tgβ.

If the mass is high enough, the heavy $\mathcal{CP}$-even Higgs boson can in principle decay into weak gauge bosons [15] $H \to VV$, $V = W$ or Z, with a partial width

$$\Gamma(H \to VV) = \delta_V \,\frac{G_F}{16\sqrt{2}\pi} \cos^2(\beta - \alpha)\, M_H^3 \left(1 - 4\frac{M_V^2}{M_H^2} + 12\frac{M_V^4}{M_H^4}\right)\left(1 - \frac{4M_V^2}{M_H^2}\right)^{1/2} \tag{7}$$

where $\delta_V = 1(2)$ for $V = Z(W)$. Below the threshold for two real bosons, the $\mathcal{CP}$-even neutral Higgs bosons H and h can decay into VV^* pairs, one of the vector bosons being virtual. The partial decay width is given by

$$\Gamma(\Phi \to VV^*) = \frac{3G_F^2 M_V^4}{16\pi^3}\, M_\Phi\, R(M_V^2/M_\Phi^2)\, \delta_V'\, g_{\Phi VV}^2 \tag{8}$$

with $\Phi = h$ or H, $\delta_W' = 1, \delta_Z' = 7/12 - 10s_W^2/9 + 40s_W^4/27$ and

$$R(x) = \frac{3(1 - 8x + 20x^2)}{(4x - 1)^{1/2}} \arccos\left(\frac{3x - 1}{2x^{3/2}}\right) - \frac{1 - x}{2x}(2 - 13x + 47x^2) - \frac{3}{2}(1 - 6x + 4x^2)\log x$$

Since h is light and the H partial width is proportional to $\cos^2(\beta - \alpha)$, they are strongly suppressed and the gold-plated ZZ signal of the heavy $\mathcal{SM}$ Higgs boson is lost in the supersymmetric extension. [If M_H is large enough for these decay modes to be kinematically allowed, M_h is very close to its maximum so that $\cos^2(\beta - \alpha) \to 0$.] For the same reason, the cascade decay of the $\mathcal{CP}$-odd Higgs boson, $A \to Zh$, is suppressed in general [2]

$$\Gamma(A \to Zh) = \frac{G_F}{8\sqrt{2}\pi} \cos^2(\beta - \alpha)\, \frac{M_Z^4}{M_A} \lambda^{1/2}(M_Z^2, M_h^2; M_A^2)\lambda(M_A^2, M_h^2; M_Z^2) \tag{9}$$

with $\lambda(x,y;z) = (1 - x/z - y/z)^2 - 4xy/z^2$ being the usual two–body phase space function.

The heavy neutral Higgs boson H can also decay into two lighter Higgs bosons [2],

$$\Gamma(H \to hh) = \frac{G_F}{16\sqrt{2}\pi}\frac{M_Z^4}{M_H}\left(1 - 4\frac{M_h^2}{M_H^2}\right)^{1/2}[\cos 2\alpha\cos(\beta+\alpha) - 2\sin 2\alpha\sin(\beta+\alpha)]^2$$

$$\Gamma(H \to AA) = \frac{G_F}{16\sqrt{2}\pi}\frac{M_Z^4}{M_H}\left(1 - 4\frac{M_A^2}{M_H^2}\right)^{1/2}[\cos 2\beta\cos(\beta+\alpha)]^2 \qquad (10)$$

These modes, however, are restricted to very small domains in the parameter space.

Gluonic Higgs decays $\Phi \to gg$ are mediated by top and bottom quark loops [the squarks decouple from the effective Φgg vertex for high masses]. For the light Higgs particle this decay mode is significant only for h masses close to the maximal value where h has $\mathcal{SM}$ like couplings, and for H masses only below 140 GeV and small values of tgβ where the coupling to top quarks is sufficiently large. Therefore, one can neglect the b loop contribution, and the decay width $\Gamma(\Phi \to gg)$ with $\Phi = h$ or H can be cast into the approximate form [16]

$$\Gamma(\Phi \to gg) = \frac{G_F\alpha_s^2(M_\Phi^2)}{36\sqrt{2}\pi^3}\, g_{\Phi u}^2\, M_\Phi^3\left[1 + \frac{215}{12}\frac{\alpha_s(M_\Phi^2)}{\pi}\right] \qquad (11)$$

The QCD radiative corrections which include ggg and $gq\bar{q}$ final states, are very important; they increase the partial width by $\sim 65\%$. For the pseudoscalar Higgs particle, the gluonic decay mode is marginal. Decays into $\gamma\gamma$ and $Z\gamma$ final states are very rare and they do not play a significant role at e^+e^- colliders; they will therefore not be discussed here.

The couplings of the charged Higgs particle to fermions is a $\mathcal{P}$ violating mixture of scalar and pseudoscalar currents

$$g_{H^+ u d} = \left(\frac{G_F}{\sqrt{2}}\right)^{1/2}\left[(1-\gamma_5)\frac{m_u}{\mathrm{tg}\beta} + (1+\gamma_5)m_d\mathrm{tg}\beta\right] \qquad (12)$$

The charged Higgs particles decay into fermions but also, if allowed kinematically, into the lightest neutral Higgs plus a W boson,

$$\Gamma(H^+ \to u\bar{d}) = N_c\frac{G_F}{4\sqrt{2}\pi}\frac{\lambda^{1/2}(m_u^2, m_d^2, M_{H^\pm}^2)}{M_{H^\pm}}\left[(M_{H^\pm}^2 - m_u^2 - m_d^2)(m_d^2\mathrm{tg}^2\beta + \frac{m_u^2}{\mathrm{tg}^2\beta}) - 4m_u^2m_d^2\right]$$

$$\Gamma(H^+ \to Wh) = \frac{G_F}{8\sqrt{2}\pi c_W^2}\cos^2(\beta-\alpha)\frac{M_W^4}{M_{H^\pm}}\lambda^{1/2}(M_W^2, M_H^2; M_{H^\pm}^2)\lambda(M_{H^\pm}^2, M_h^2; M_W^2) \qquad (13)$$

Below the tb and Wh thresholds, the charged Higgs particles will decay mostly into $\tau\nu_\tau$ and $c\bar{s}$ pairs, the former being dominant for tg$\beta > 1$. For large $M_{H^\pm}$ and large tgβ values, the top–bottom decay $H^+ \to t\bar{b}$ becomes dominant.

Adding up the various decay modes, the width of all five Higgs bosons remains very narrow, even for large masses. This is shown for the two representative values tg$\beta = 2.5$ and 20 in Fig. 4. Apart from the $\mathcal{CP}$–even heavy neutral Higgs boson H and small tgβ, the pattern of

branching ratios is in general quite simple. The neutral Higgs bosons decay preferentially to $b\bar{b}$, and to a lesser extent to $\tau^+\tau^-$ pairs; the charged Higgs bosons to $\tau\nu_\tau$ and, preferentially, $t\bar{b}$ pairs above this threshold [see Fig. 5]. The decay patterns lead to clear signatures for the detection of these particles at e^+e^- colliders.

3.2 Decays to Charginos and Neutralinos

In most studies of supersymmetric Higgs particles at future colliders it is assumed that they do not decay into supersymmetric particles. However, while sfermions are probably too heavy to affect Higgs decays, the Higgs boson decays into charginos and neutralinos could eventually play a significant role since some of these particles are expected to be lighter or of the same order as M_Z. These new channels could open up at least for the heavy Higgs bosons H, A and $H^\pm$. They could be so large that they reduce the branching fractions for the $\mathcal{SM}$ decays in a sizable way, therefore altering the signals, and as a result, the search strategies for these particles.

The decay widths of the neutral Higgs bosons H_k, with $H_1 = H$, $H_2 = h$ and $H_3 = A$, into chargino or neutralino pairs are given by [13]

$$\Gamma(H_k \to \tilde{\chi}_i\tilde{\chi}_j) = \frac{G_F}{2\sqrt{2}\pi}\lambda^{\frac{1}{2}}\frac{M_W^2 M_H}{1+\delta(i,j)}\left[(F_{ijk}^2 + F_{jik}^2)\left(1 - \frac{M_i^2}{M_H^2} - \frac{M_j^2}{M_H^2}\right) - 4F_{ijk}F_{jik}\eta_k\frac{|M_iM_j|}{M_H^2}\right] \qquad (14)$$

where $\eta_{1,2} = +1$, $\eta_3 = -1$ and $\delta(i,j)$ is zero unless the final state consists of two identical Majorana neutralinos in which case $\delta(i,j) = 1$. The coefficients F_{ijk} can be expressed in terms of the elements of the matrices U, V which diagonalize the chargino mass matrices, and of the neutralino matrix Z, given in Ref. [13]. For completeness we list these coefficients,

$$H_k \to \tilde{\chi}_i^+\tilde{\chi}_j^- \;:\; F_{ijk} = \frac{1}{\sqrt{2}}[e_kV_{i1}U_{j2} - d_kV_{i2}U_{j1}]$$

$$H_k \to \tilde{\chi}_i^0\tilde{\chi}_j^0 \;:\; F_{ijk} = \frac{e_k}{2}\left[Z_{i3}Z_{j2} + Z_{j3}Z_{i2} - \frac{s_W}{c_W}(Z_{i3}Z_{j1} + Z_{j3}Z_{i1})\right] + \frac{d_k}{2}\left[Z_{i4}Z_{j2} + Z_{j4}Z_{i2} - \frac{s_W}{c_W}(Z_{i4}Z_{j1} + Z_{j4}Z_{i1})\right] \qquad (15)$$

where the constants e_k and d_k are given by

$$e_1/d_1 = \cos\alpha/-\sin\alpha\,, \quad e_2/d_2 = \sin\alpha/\cos\alpha\,, \quad e_3/d_3 = -\sin\beta/\cos\beta \qquad (16)$$

For the charged Higgs boson decays into neutralino/chargino pairs, the partial widths read

$$\Gamma(H^\pm \to \tilde{\chi}_i^\pm\tilde{\chi}_j^0) = \frac{G_F}{2\sqrt{2}\pi}M_W^2 M_{H^\pm}\lambda^{\frac{1}{2}}\left[(F_L^2 + F_R^2)\left(1 - \frac{M_i^2}{M_{H^\pm}^2} - \frac{M_j^2}{M_{H^\pm}^2}\right) - 4F_LF_R\frac{|M_iM_j|}{M_{H^\pm}^2}\right] \qquad (17)$$

with

$$F_L = \cos\beta\left[Z_{j4}V_{i1} + \frac{1}{\sqrt{2}}\left(Z_{j2} + \frac{s_W}{c_W}Z_{j1}\right)V_{i2}\right]$$

$$F_R = \sin\beta\left[Z_{j3}U_{i1} - \frac{1}{\sqrt{2}}\left(Z_{j2} + \frac{s_W}{c_W}Z_{j1}\right)U_{i2}\right] \qquad (18)$$

We have updated the branching ratios of the Higgs boson decays into chargino and neutralino pairs [17], taking into account the new LEP100 constraints [on both the supersymmetric particles and the Higgs boson masses] and including the large radiative corrections to the Higgs masses and couplings. The results of the analysis are summarized in Fig. 6–8.

In Fig. 6, we display the decay branching ratios of the lightest $\mathcal{CP}$–even neutral Higgs boson h into the lightest chargino pair and the lightest and next–to–lightest neutralino pair. The contours are shown in the $[\mu, M]$ plane where the sum of these branching ratios exceeds 5% [dashed lines] and 50% [solid lines]; the dotted lines are the countours which are excluded by LEP100 data and which can be probed at LEP200. These decays can be sizable in the area of the $[\mu, M]$ plane between the two LEP contours. They are particularly important for h masses close to the maximum allowed values [see Fig. 1]; in this case the lightest Higgs boson has $\mathcal{SM}$ couplings and the dominant $b\bar{b}$ decay mode is not enhanced anymore for large tgβ values, so that other decay modes can become significant. For these masses and for large tgβ values, the branching ratios for neutralino/chargino decays are sizable even outside the regions which can be probed at LEP200.

In Figs. 7a–c, the branching ratios of the heavy $\mathcal{CP}$–even, the $\mathcal{CP}$–odd and the charged Higgs boson decays into chargino and neutralino pairs are displayed for the usual values of tgβ and for a set of μ and M values. These branching fractions can be very large and they can even be dominant in some areas of the $\mathcal{MSSM}$ parameter space, exeeding 50% for positive values of μ and/or M values below 200 GeV. This is mainly due to the fact that the couplings of the Higgs bosons to charginos and neutralinos are gauge couplings which can be larger than the Yukawa couplings to standard fermions and the couplings to the gauge bosons [the latter being zero at tree level for the $\mathcal{CP}$–odd and being suppressed for the heavy $\mathcal{CP}$–even Higgs particles].

Finally, the branching fractions of the invisible neutral Higgs decays are shown in Fig. 8 for the same values of μ, M and tgβ as in the previous figures. These decays can be important and they can jeopardize the search for the Higgs particles at hadron colliders. At e^+e^- colliders however, missing mass techniques allow us to tag these events in the case of the $\mathcal{CP}$–even Higgs bosons which can be produced in association with the Z boson, or in mixed visible and invisible decay modes of $\mathcal{CP}$–odd and $\mathcal{CP}$–even Ah and AH pairs.

4 Production Mechanisms

4.1 Neutral Higgs Bosons

The search for neutral $\mathcal{SUSY}$ Higgs bosons at 500 GeV e^+e^- colliders will be a direct extension of the search to be performed at LEP200, which is expected to cover the mass range up to ~ 85 GeV. Higher energies, $\sqrt{s} > 250$ GeV, are required to sweep the entire parameter space of the $\mathcal{MSSM}$.

The main production mechanisms of the $\mathcal{CP}$–even neutral Higgs bosons at e^+e^- colliders are the bremsstrahlung [18] and the fusion processes [19], which are well–known from the

Standard Model,

$$
\begin{array}{llll}
(a) & \text{bremsstrahlung} & e^+e^- \rightarrow & (Z) \rightarrow Z + h/H \\
(b) & \text{fusion processes} & e^+e^- \rightarrow & \nu\,\bar{\nu}\,(WW) \rightarrow \nu\,\bar{\nu}\; + h/H \\
 & & e^+e^- \rightarrow & e^+e^-(ZZ) \rightarrow e^+e^- + h/H
\end{array}
$$

and pair production of the $\mathcal{CP}$–odd higgs boson A together with one of the $\mathcal{CP}$–even Higgs bosons,

$$(c)\ \ \text{pair production} \qquad e^+e^- \rightarrow (Z) \rightarrow A + h/H$$

Since the $\mathcal{CP}$–odd Higgs boson A does not couple directly to two gauge bosons, pair production is the only source for this particle.

The cross sections for the four bremsstrahlung and pair production processes can be expressed as [9]

$$
\begin{aligned}
\sigma(e^+e^- \rightarrow Zh) &= \sin^2(\beta-\alpha)\,\sigma_{SM} \\
\sigma(e^+e^- \rightarrow ZH) &= \cos^2(\beta-\alpha)\,\sigma_{SM} \\
\sigma(e^+e^- \rightarrow Ah) &= \cos^2(\beta-\alpha)\,\sigma_{SM}\,\bar{\lambda} \\
\sigma(e^+e^- \rightarrow AH) &= \sin^2(\beta-\alpha)\,\sigma_{SM}\,\bar{\lambda}
\end{aligned}
\tag{19}
$$

where

$$\sigma_{SM} = \frac{G_F^2 M_Z^4}{92\pi s}\,(v_e^2 + a_e^2)\,\lambda^{1/2}(M_{h,H}^2, M_Z^2; s)\,\frac{\lambda(M_{h,H}^2, M_Z^2; s) + 12M_Z^2/s}{(1 - M_Z^2/s)^2} \tag{20}$$

is the $\mathcal{SM}$ cross section for Higgs bremsstrahlung and the factor $\bar{\lambda}$ accounts for the correct suppression of the P–wave cross sections near the threshold

$$\bar{\lambda} = \frac{\lambda^{3/2}(M_j^2, M_A^2; s)}{\lambda^{1/2}(M_j^2, M_Z^2; s)\,[\lambda(M_j^2, M_A^2; s) + 12M_Z^2/s]} \qquad [\, j = h, H\,] \tag{21}$$

with the velocity factor $\lambda(x, y; z) = (1 - x/z - y/z)^2 - 4xy/z^2$.

The cross sections for bremsstrahlung and pair production as well as the cross sections for the production of the light and the heavy neutral Higgs bosons h and H are mutually complementary to each other, coming either with a coefficient $\sin^2(\beta-\alpha)$ or $\cos^2(\beta-\alpha)$. Since σ_{SM} is large, at least the lightest of the $\mathcal{CP}$–even Higgs bosons should be detected in these reactions. From the mass and the $\sin^2/\cos^2(\beta-\alpha)$ plots we conclude that, depending on the values of M_h and tgβ, the following final states will be observed [Fig. 9]:

M_h "small" ,	tgβ small :	hZ ,	HZ ,	hA ,	HA
	tgβ large :		HZ ,	hA ,	
M_h "large" ,	tgβ small :	hZ ,			$[HA]$
	tgβ large :	hZ ,			$[HA]$

where "M_h small" and "large" are synonymous for "considerably below" and "close to the upper limit of the light $\mathcal{CP}$–even Higgs boson h" for a given value of tgβ. If M_h is "large" the H, A masses can exceed the kinematical limit for HA pair production.

The cross sections for the production of the $\mathcal{CP}$–even light and heavy Higgs bosons h and H via bremsstrahlung are shown as functions of the masses in Fig. 10a. The cross section for h is large for small values of tgβ and/or large values of M_h where $\sin^2(\beta-\alpha)$ approaches its maximal value. In these two cases the cross section is of the order of ~ 50 fb, which for an integrated luminosity of 10 fb^{-1} corresponds to ~ 500 events. By contrast, the cross section for H is large for large tgβ and light h [implying small M_H]. In the case of h [and also for H in most parts of the parameter space] the signal consists of a Z boson accompanied by a $b\bar{b}$ or $\tau^+\tau^-$ pair. The signal is easy to separate from the background which comes mainly from ZZ production if the Higgs mass is close to M_Z.

The cross sections for the associated production channels $e^+e^- \to Ah$ and AH are displayed in Fig. 10b. As anticipated, the situation is opposite to the previous case: the cross section for Ah is large for light h and/or large values of tgβ whereas AH production is preferred in the complementary region. The sum of the two cross sections decreases from ~ 50 to 10 fb if M_A increases from ~ 50 to 200 GeV. In major parts of the parameter space, the signal consists of four b quarks in the final state, requiring instruments for efficient b quark tagging. Mass constraints will help to eliminate the backgrounds from QCD jets as well as ZZ final states.

<u>WW and ZZ fusion</u> provide additional mechanisms for the production of the $\mathcal{CP}$–even neutral Higgs bosons. They lead to Higgs bosons in association with $\nu\bar{\nu}$ or e^+e^- pairs in the final state. The cross sections can again be expressed in terms of the corresponding $\mathcal{SM}$ cross sections,

$$\begin{aligned} \sigma(e^+e^- \to (VV) \to h) &= \sin^2(\beta-\alpha)\sigma_{SM}^{VV} \\ \sigma(e^+e^- \to (VV) \to H) &= \cos^2(\beta-\alpha)\sigma_{SM}^{VV} \end{aligned} \qquad (22)$$

for $V = W, Z$. Even though the analytic expressions for σ_{SM}^{VV} are rather involved, they can nevertheless be cast into a compact integral form [20],

$$\sigma_{SM}^{VV} = \frac{G_F^3 M_V^4}{64\sqrt{2}\pi^3} \int_{\kappa_H}^1 \mathrm{d}x \int_x^1 \frac{\mathrm{d}y}{[1+(y-x)/\kappa_V]^2} \left[(v^2+a^2)^2 f(x,y) + 4v^2a^2 g(x,y)\right] \qquad (23)$$

$$\begin{aligned} f(x,y) &= \left(\frac{2x}{y^3} - \frac{1+2x}{y^2} + \frac{2+x}{2y} - \frac{1}{2}\right)\left[\frac{z}{1+z} - \log(1+z)\right] + \frac{x}{y^3}\frac{z^2(1-y)}{1+z} \\ g(x,y) &= \left(-\frac{x}{y^2} + \frac{2+x}{2y} - \frac{1}{2}\right)\left[\frac{z}{1+z} - \log(1+z)\right] \end{aligned}$$

with $\kappa_H = M_H^2/s$, $\kappa_V = M_V^2/s$, $z = y(x-h)/(\kappa_V x)$ and v, a denoting the electron couplings to the gauge bosons [$v = -1+4s_W^2, a = -1$ for the Z and $v = a = \sqrt{2}$ for the W boson].

[For $\sqrt{s}, M_H \gg M_W$ the cross section simplifies, in the equivalent W approximation, to

$$\sigma = \frac{G_F^3 M_W^4}{4\sqrt{2}\pi^3}\left[\left(1+\frac{M_H^2}{s}\right)\log\frac{s}{M_H^2} - 2\left(1-\frac{M_H^2}{s}\right)\right] \qquad (24)$$

At $\sqrt{s} = 500$ GeV, however, the equivalent W approximation gives a result still twice as large as the exact cross section.]

The cross sections for the WW fusion mechanism are larger than those for the bremsstrahlung process if the Higgs mass is moderately small – less than 160 GeV at $\sqrt{s} = 500$ GeV. However, since the final state cannot be fully reconstructed, the signal is more difficult to extract. As in the case of the bremsstrahlung process, the production of light h and heavy H Higgs bosons are complementary. The cross sections for the ZZ fusion mechanism are about an order of magnitude smaller than those for the WW fusion process. ZZ fusion will nevertheless be useful as the final state can be fully reconstructed. The cross sections are displayed for representative values of tgβ in Fig. 11.

The preceding discussion on the neutral $\mathcal{MSSM}$ Higgs sector at e^+e^- linear colliders can be summarized in the following points [see also Ref. [7]]:

(i) The lightest $\mathcal{CP}$–even Higgs particle h can be detected in the entire range of the $\mathcal{MSSM}$ parameter space, either through the bremsstrahlung process $e^+e^- \to hZ$ or through pair production $e^+e^- \to hA$. In fact, this conclusion holds true even at a c.m. energy of 300 GeV, independently of the top and squark mass values, and also if invisible neutralino decays are allowed for.

(ii) There is a subtantial area of the $[M_h, \mathrm{tg}\beta]$ parameter space where *all* neutral $\mathcal{SUSY}$ Higgs bosons can be discovered at a 500 GeV collider. This is possible if the masses of the heavy scalar H and the pseudoscalar A boson are less than ~ 230 GeV.

(iii) In some part of the $\mathcal{MSSM}$ parameter space, the lightest Higgs particle h can be detected, but it cannot be distinguished from the $\mathcal{SM}$ Higgs boson [if $\mathcal{SUSY}$ decays are not allowed]. This happens if, for a given value of tgβ, M_h is very close to its maximum value and H and A are too heavy to be produced in association. In this case, the couplings of h to gauge bosons and fermions are $\mathcal{SM}$ like. As long as this ambiguity cannot be resolved by proceeding to higher collider energies, the only way to distinguish h from the $\mathcal{SM}$ Higgs particle is provided by Higgs production in $\gamma\gamma$ fusion. While in the Standard Model this process is built-up by W and top quark loops, additional contributions in $\mathcal{SUSY}$ models are provided by supersymmetric particle loops, such as chargino loops, which alter the $\mathcal{SM}$ production rates.

(iv) Higgs boson decays into charginos and neutralinos can be very important in some areas of the MSSM parameter space; in particular, invisible decays of the neutral Higgs bosons can be larger than the decays into standard particles. At e^+e^- colliders, missing mass techniques allow to isolate these events in the bremsstrahlung process for the $\mathcal{CP}$–even Higgs bosons or in a mixture of visible and invisible decay modes of Ah and AH in the pair production processes.

Some of the features are not specific to the minimal extension but they are expected to be realized also in more general $\mathcal{SUSY}$ models. For example, the existence of a light Higgs boson with a mass in the intermediate range is quite generally predicted in supersymmetric theories[1] [2, 4, 5].

[1] For instance, in low energy supersymmetric models derived from E_6–based grand unified theories or

4.2 Charged Higgs Bosons

An unambiguous signal for an extended Higgs sector would be the discovery of a charged Higgs boson. In a general two–Higgs doublet model, charged Higgs bosons can be as light as ~ 45 GeV, the lower limit derived from the negative search at LEP100 [21]. In the $\mathcal{MSSM}$ however, $H^\pm$ is constrained to be heavier than the W boson. More precisely, the lower limit $M_h > 40$ GeV obtained at LEP100 [21] implies $M_{H^\pm} > 90$ GeV.

In e^+e^- collisions, the production of a pair of charged Higgs bosons [22] proceeds through virtual photon and Z boson exchange. The cross section depends only on the charged Higgs mass [and does not depend on any extra parameter],

$$\sigma(e^+e^- \longrightarrow H^+H^-) = \frac{\pi\alpha^2}{3s}\left[1 - \frac{2\hat{v}_e\hat{v}_H}{1-M_Z^2/s} + \frac{(\hat{a}_e^2+\hat{v}_e^2)\hat{v}_H^2}{(1-M_Z^2/s)^2}\right]\beta^3 \tag{25}$$

with the standard Z charges $\hat{v}_e = (-1+4s_W^2)/4c_W s_W$, $\hat{a}_e = -1/4c_W s_W$ and $\hat{v}_H = (-1+2s_W^2)/2c_W s_W$, and $\beta = (1-4M_{H^\pm}^2/s)^{1/2}$ the velocity of the Higgs particles. The cross section is shown in Fig.12a as a function of the charged Higgs mass for the c.m. energy $\sqrt{s} = 500$ GeV. For small Higgs masses the cross section is of order 100 fb, but it drops very quickly due to the P–wave suppression factor β^3 near the threshold. For $M_{H^\pm} = 220$ GeV, the cross section has fallen to a level of $\simeq 5$ fb, which for an integrated luminosity of 10 fb^{-1} corresponds to 50 events. The angular distribution of the charged Higgs bosons follows the $\sin^2\theta$ law,

$$\frac{1}{\sigma}\frac{\mathrm{d}\sigma}{\mathrm{d}\cos\theta} = \frac{3}{4}\sin^2\theta \tag{26}$$

typical for spin–zero particle production.

Charged Higgs particles can also be created in $\gamma\gamma$ collisions. Generating the γ beams through back–scattering of laser light, the total energy of the $\gamma\gamma$ collider can go up to $\sim 80\%$ of the original e^+e^- energy, which corresponds to $\sqrt{s_{\gamma\gamma}} \simeq 400$ GeV for a 500 GeV e^+e^- collider. The $\gamma\gamma$ luminosity is expected to be of the same magnitude as the original e^+e^- luminosity. The cross section is given by

$$\sigma(\gamma\gamma \to H^+H^-) = \frac{2\pi\alpha^2}{s}\beta\left[2 - \beta^2 - \frac{1-\beta^4}{2\beta}\log\frac{1+\beta}{1-\beta}\right] \tag{27}$$

where β is the velocity of the Higgs particle. The numerical result is displayed in Fig. 12a for the $\gamma\gamma$ luminosity without beam polarization [23]. Due to the reduced energy, the maximum Higgs mass which can be probed in $\gamma\gamma$ collisions is smaller than in the original e^+e^- collisions; the cross section however is enhanced by a factor ~ 3 in the low mass range.

The charged Higgs boson, if lighter than the top quark, can also be produced in top decays [24]. The ratio of the decay width $t \to bH^+$ to the standard mode $t \to bW^+$ is given by

$$\frac{\Gamma(t\to bH^+)}{\Gamma(t\to bW^+)} = \frac{\lambda(M_{H^\pm}^2, m_b^2; m_t^2)^{1/2}}{\lambda(M_W^2, m_b^2; m_t^2)^{1/2}}\,\frac{(m_t^2+m_b^2-M_{H^\pm}^2)(m_t^2\mathrm{ctg}^2\beta + m_b^2\mathrm{tg}^2\beta) + 4m_t^2m_b^2}{M_W^2(m_t^2+m_b^2-2M_W^2) + (m_t^2-m_b^2)^2} \tag{28}$$

superstring–inspired theories, which require the existence of two–doublets and one or two singlets of scalar fields, one of the $\mathcal{CP}$–even Higgs bosons is constrained to be lighter than ~ 120 GeV [5]. The couplings of this particle are similar to those of the lightest Higgs boson of the $\mathcal{MSSM}$.

The branching ratio BR($t \to bH^+$) is shown in Fig. 12b as a function of tgβ for $m_t = 140$ GeV and $M_{H^\pm} = 100, 120$ and 135 GeV. In the range $1 < \mathrm{tg}\beta < m_t/m_b$ favored by $\mathcal{SUSY}$ models and for $M_{H^\pm} = 100$ GeV, the branching ratio varies between $\sim 2\%$ and 20%. Since the cross section for top pair production is of order of 0.5 pb at $\sqrt{s} = 500$ GeV, this corresponds to 10,000 t and $\bar{t}$ quarks at a luminosity $\int\mathcal{L} = 10$ fb^{-1} and between 200 and 2000 charged Higgs bosons.

The signature for $H^\pm$ production can be read off the graphs displaying the branching ratios in the previous section. If $M_{H^\pm} < m_t + m_b$, the charged Higgs boson will decay mainly into $\tau\nu_\tau$ and $c\bar{s}$ pairs, the $\tau\nu_\tau$ mode dominating for tgβ larger than unity. This results in a surplus of τ final states over e,μ final states, an apparent breaking of τ *vs.* e,μ universality. For large Higgs masses the dominant decay mode is the top decay $H^+ \to t\bar{b}$. In some part of the parameter space also the decay $H^+ \to W^+h$ is allowed, leading to cascades with heavy τ and b particles in the final state. In addition, cascades from charginos and neutralinos could be observed in Higgs decays.

5 Conclusions

We have presented a comprehensive synopsis of the Higgs scenario in the minimal supersymmetric extension of the Standard Model at e^+e^- linear colliders operating at c.m. energies between 300 to 500 GeV [which may later be expanded to 1 TeV]. The characteristic pattern of the $\mathcal{MSSM}$ Higgs scenario can be taken as representative for a wide class of supersymmetric models, in particular the existence of a light neutral Higgs boson with a mass of order M_Z. Also the additional neutral and charged Higgs bosons could have masses near the scale of electroweak symmetry breaking.

In general several independent production channels can be exploited to search for the Higgs particles at e^+e^- colliders. At least the lightest neutral Higgs particle of the $\mathcal{MSSM}$ must be detected at these colliders; in major parts of the $\mathcal{SUSY}$ parameter space all five neutral and charged Higgs bosons can be observed. This holds true for all Higgs decay modes. In addition to the decays of the neutral Higgs particles to $\tau^+\tau^-$ and, predominantly, to $b\bar{b}$, cascade Higgs decays could eventually occur. We have also analyzed, in some detail, Higgs decays to supersymmetric particles. For the mass range of a 500 GeV collider, decays to sfermions cannot play a major role, decays to charginos and neutralinos however are potentially interesting channels.

In summary, comprehensive analyses of the Higgs scenarios in supersymmetric extensions of the Standard Model can be performed at high energy e^+e^- colliders.

Note added. After completing this work, we received a preprint by H. Baer et al. [25], where SUSY Higgs bosons decays into charginos and neutralinos have been discussed in the context of hadron colliders.

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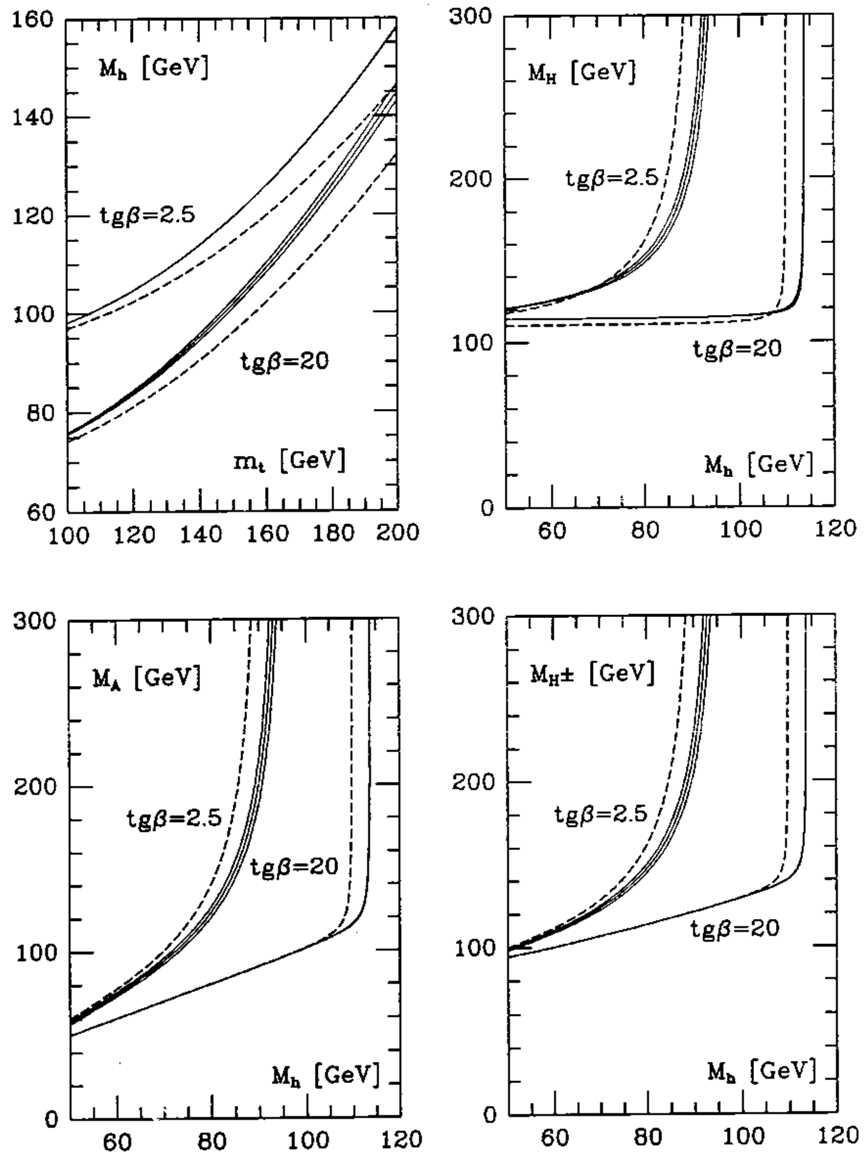

Fig. 1 The masses of the Higgs particles in the $\mathcal{MSSM}$ including radiative corrections; the squark masses are fixed to 1 TeV. The dashed curve shows the leading correction [$A_t = A_b \doteq \mu = 0$] while the solid curves include the full corrections [$A_t = A_b = 1$ TeV and $\mu = -200, 0, 200$ GeV]. (a) Upper limit on M_h as a function of m_t; (b)-(d) masses of the H, A and $H^\pm$ Higgs bosons as functions of M_h. The top mass is fixed to 140 GeV.

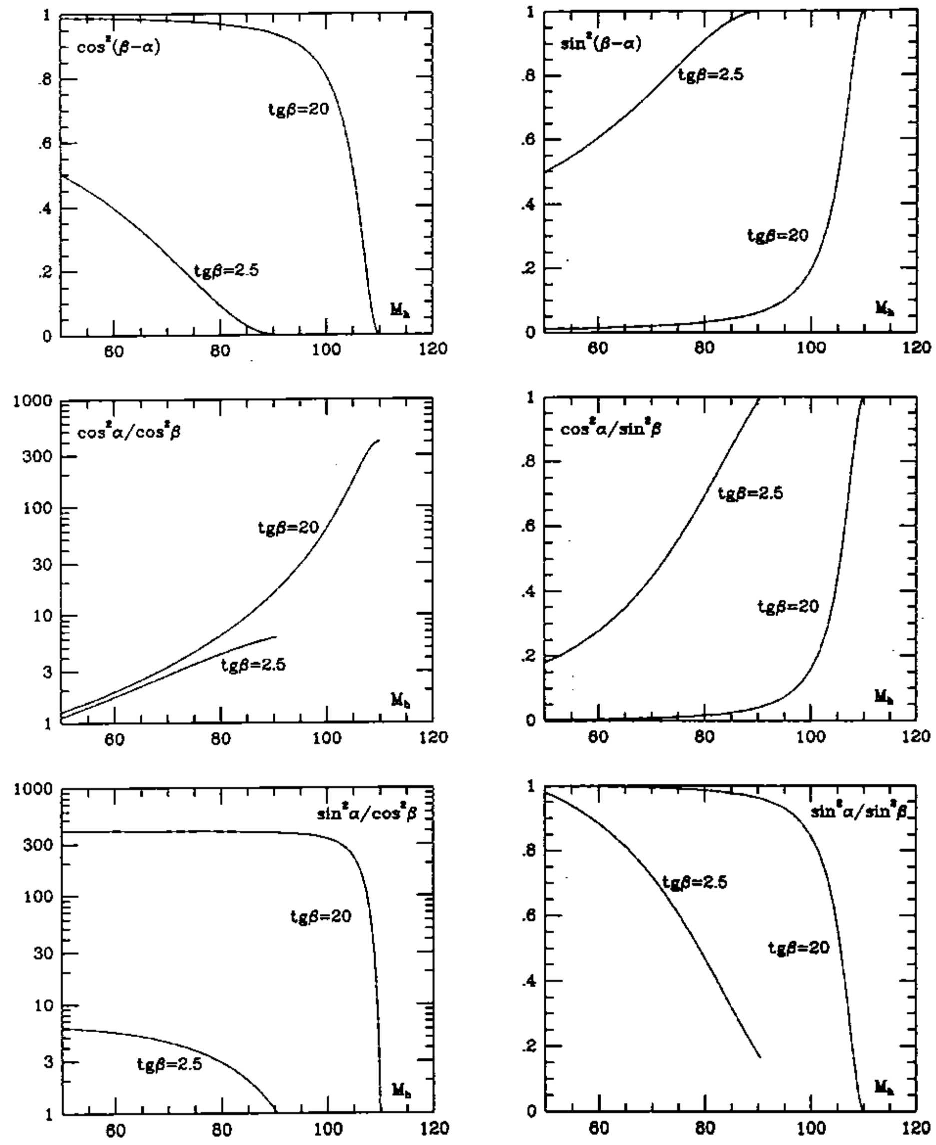

Fig. 2 Coupling parameters of $\mathcal{SUSY}$ Higgs bosons as functions of the lightest Higgs h mass. The couplings are normalized to the $\mathcal{SM}$ couplings as given in Tab. 1. The parameters are the same as in Fig. 1.

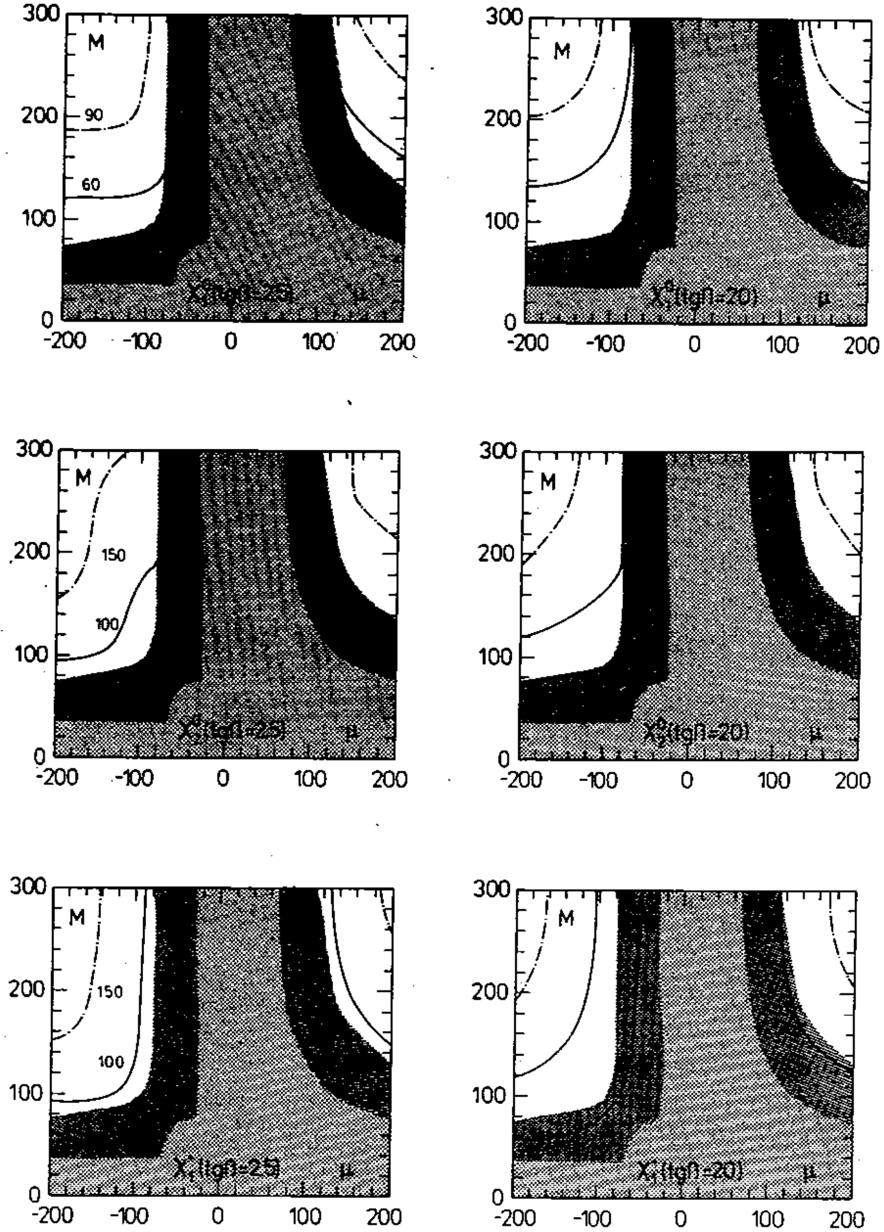

Fig. 3 Contour lines in the $[\mu, M]$ plane for the masses of the lightest and second-lightest neutralino and for the lightest chargino. The numbers at the contour lines denote the masses in GeV. The shaded areas are the regions which are (can be) excluded at LEP100 (LEP200).

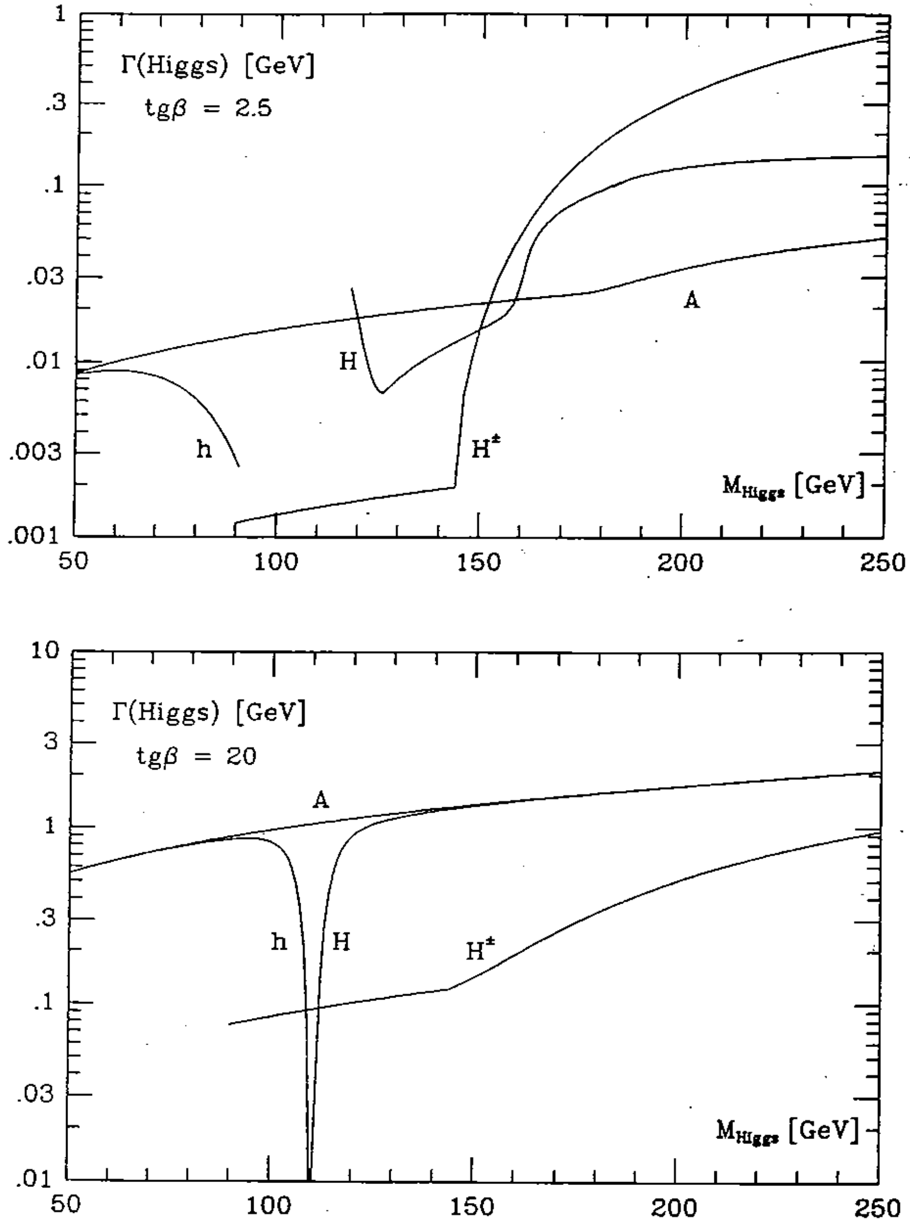

Fig. 4 Total decay widths of the $\mathcal{SUSY}$ Higgs bosons [without $\mathcal{SUSY}$ decays] as functions of their masses for (a) $tg\beta = 2.5$ and (b) $tg\beta = 20$. The top mass was chosen as $m_t = 140$ GeV.

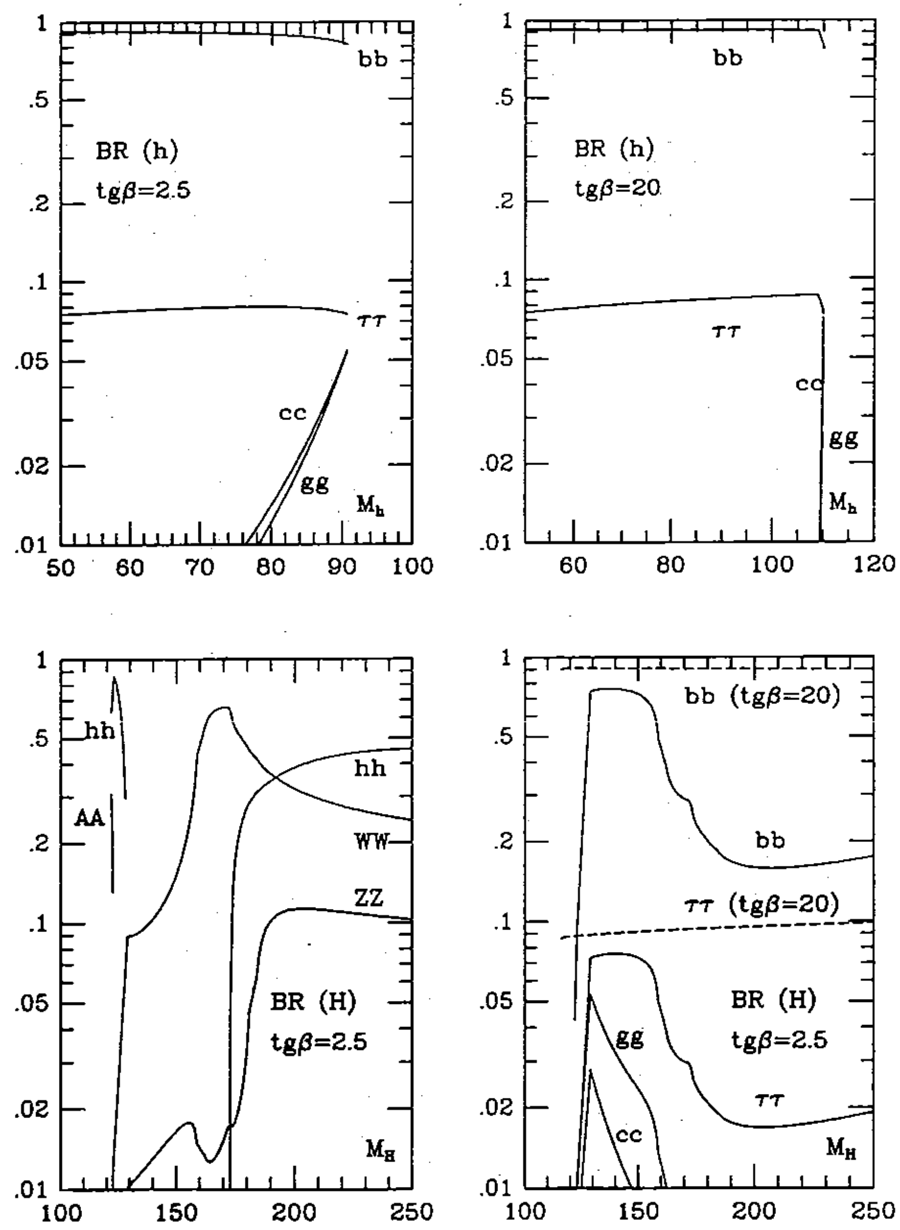

Fig. 5a–b Decay branching ratios of the $\mathcal{CP}$–even neutral Higgs bosons [without $\mathcal{SUSY}$ decays] as functions of their masses for two values of $tg\beta = 2.5$ and 20.

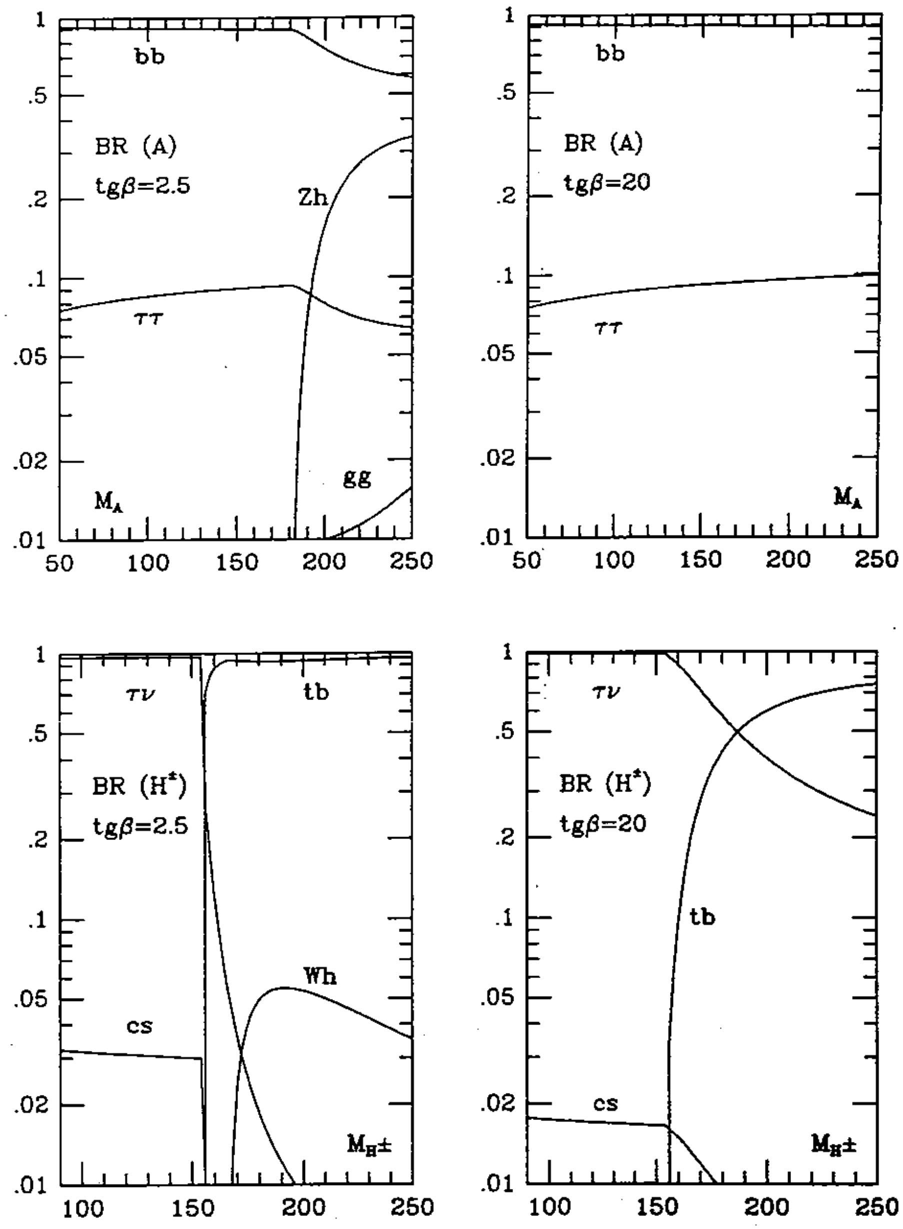

Fig. 5c–d Decay branching ratios of the pseudoscalar and charged Higgs bosons [without $\mathcal{SUSY}$ decays] as functions of their masses for two values of $tg\beta = 2.5$ and 20.

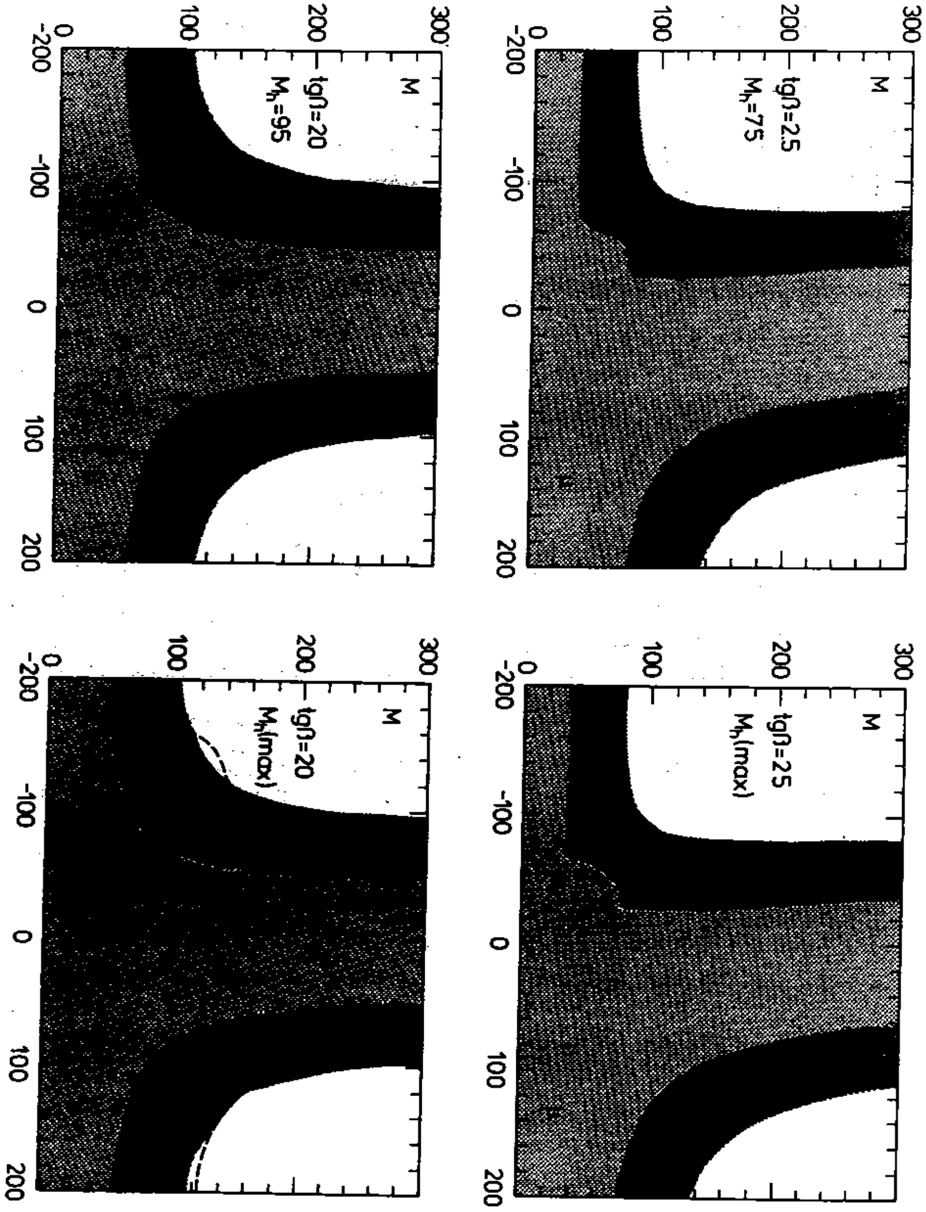

Fig. 6: Contour lines in the $[\mu, M]$ plane where the sum of the branching ratios of the lightest Higgs boson h into charginos and neutralinos exceeds 5% (dashed lines) and 50% (full lines) for two values of M_h and tgβ. The shaded areas are the regions which are (can be) excluded at LEP100 (LEP200).

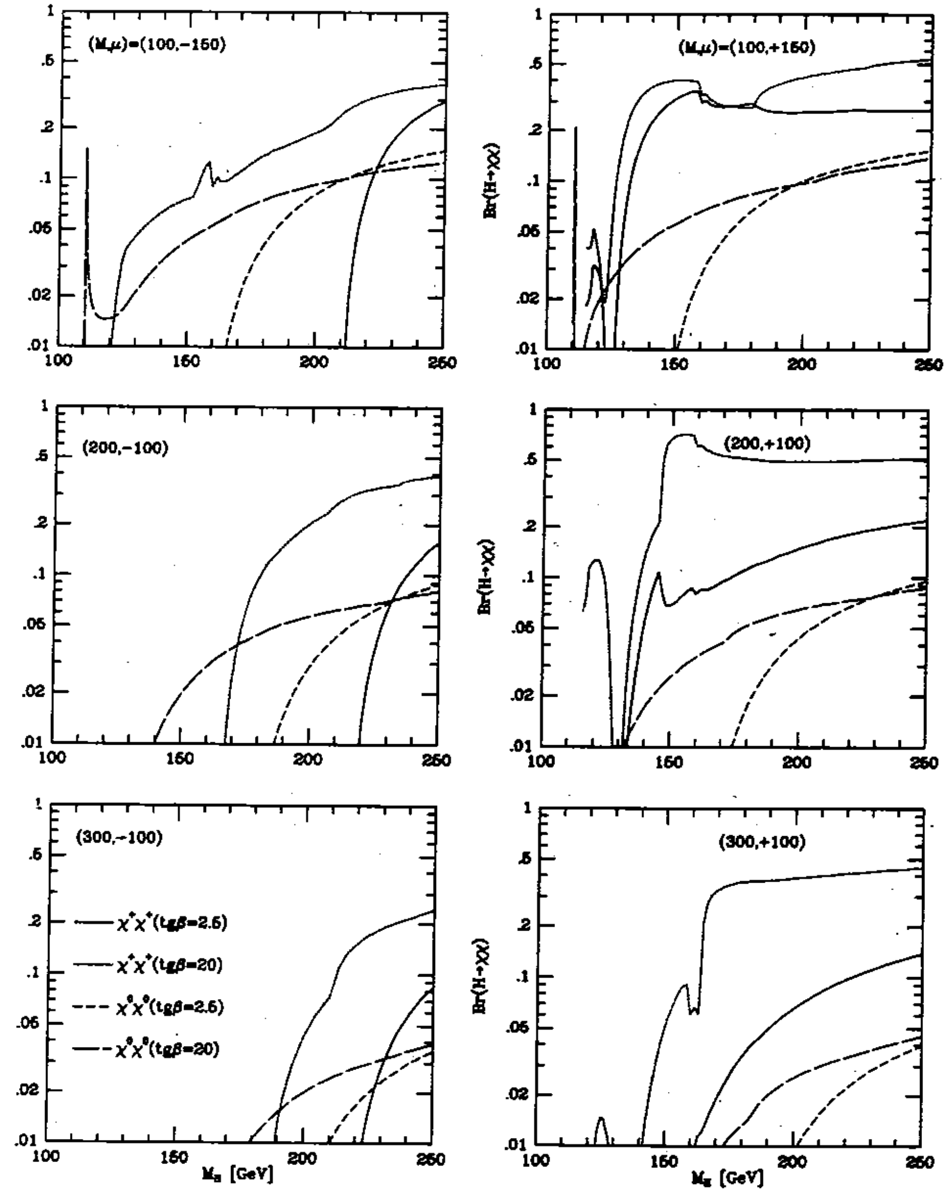

Fig. 7a Branching ratios of the heavy $\mathcal{CP}$–even Higgs boson decay into the sum of charginos and the sum of neutralinos as functions of the mass for tgβ = 2.5 and 20. The values of μ and M are given in GeV.

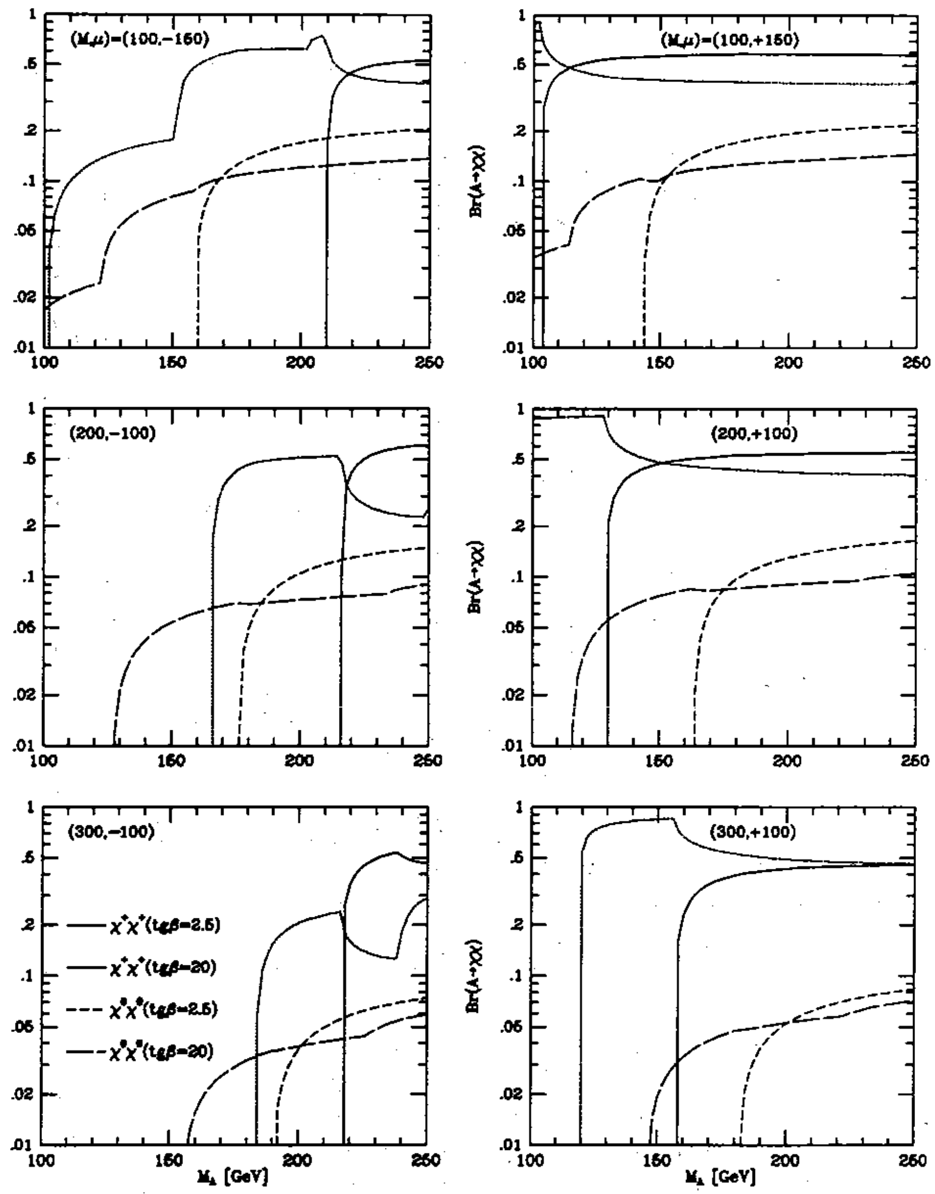

Fig. 7b Branching ratios of the $\mathcal{CP}$–odd Higgs boson decays into the sum of charginos and the sum of neutralinos as functions of the mass for tgβ = 2.5 and 20. The values of μ and M are given in GeV.

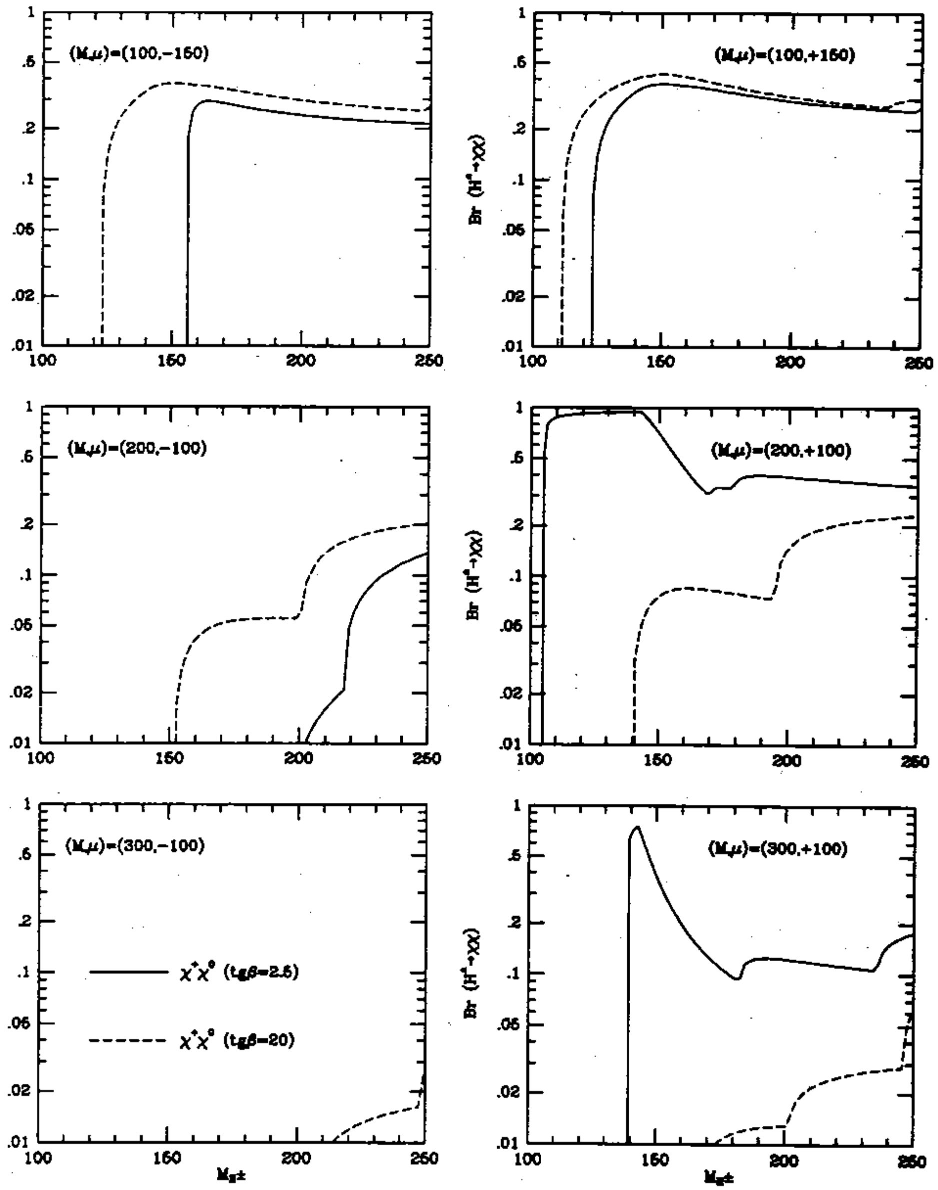

Fig. 7c Branching ratios of the charged Higgs boson decays into charginos and neutralinos as functions of the mass for tgβ = 2.5 and 20. The values of μ and M are given in GeV.

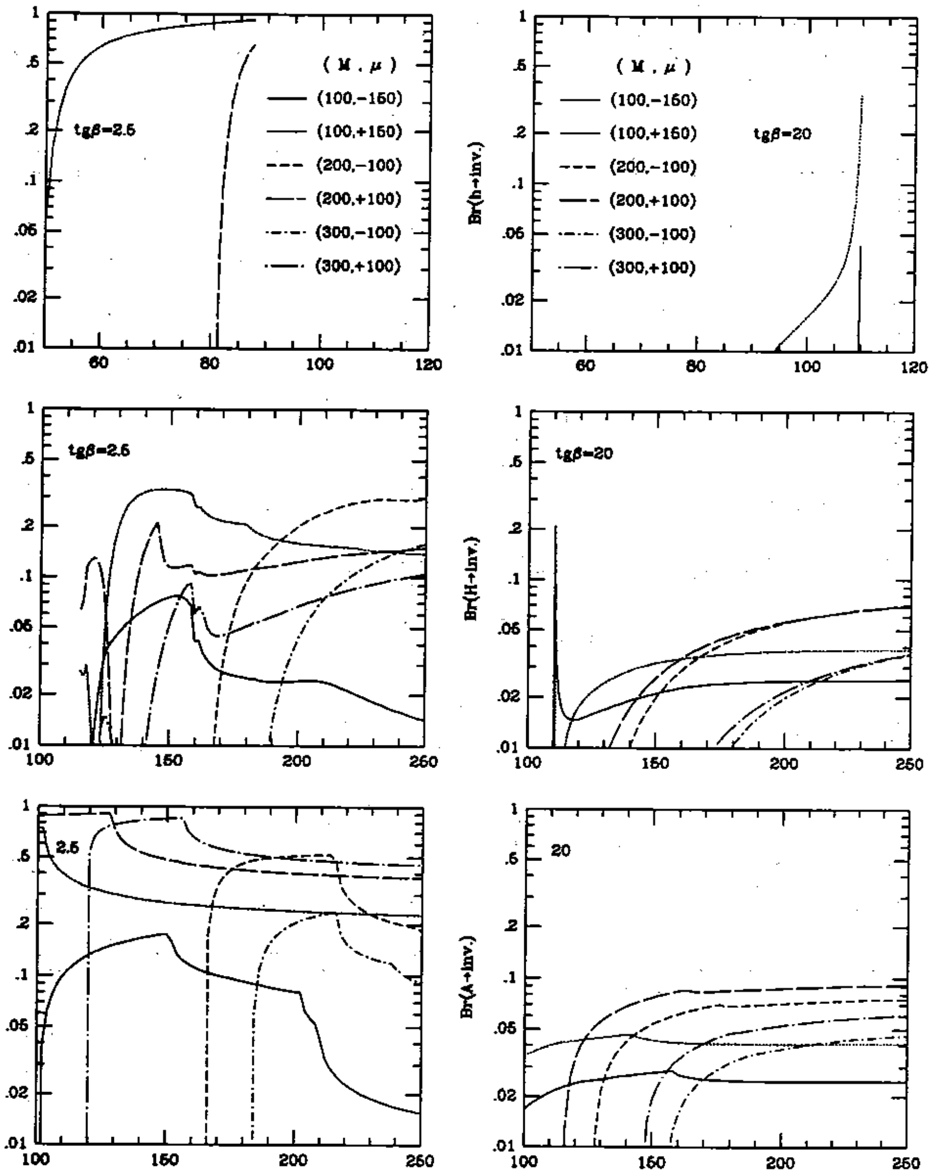

Fig. 8 Branching ratios of the decays of the three neutral Higgs bosons into the lightest neutralino pair [invisible decays] as a function of their masses for tgβ = 2.5 and 20. The values of μ and M are given in GeV.

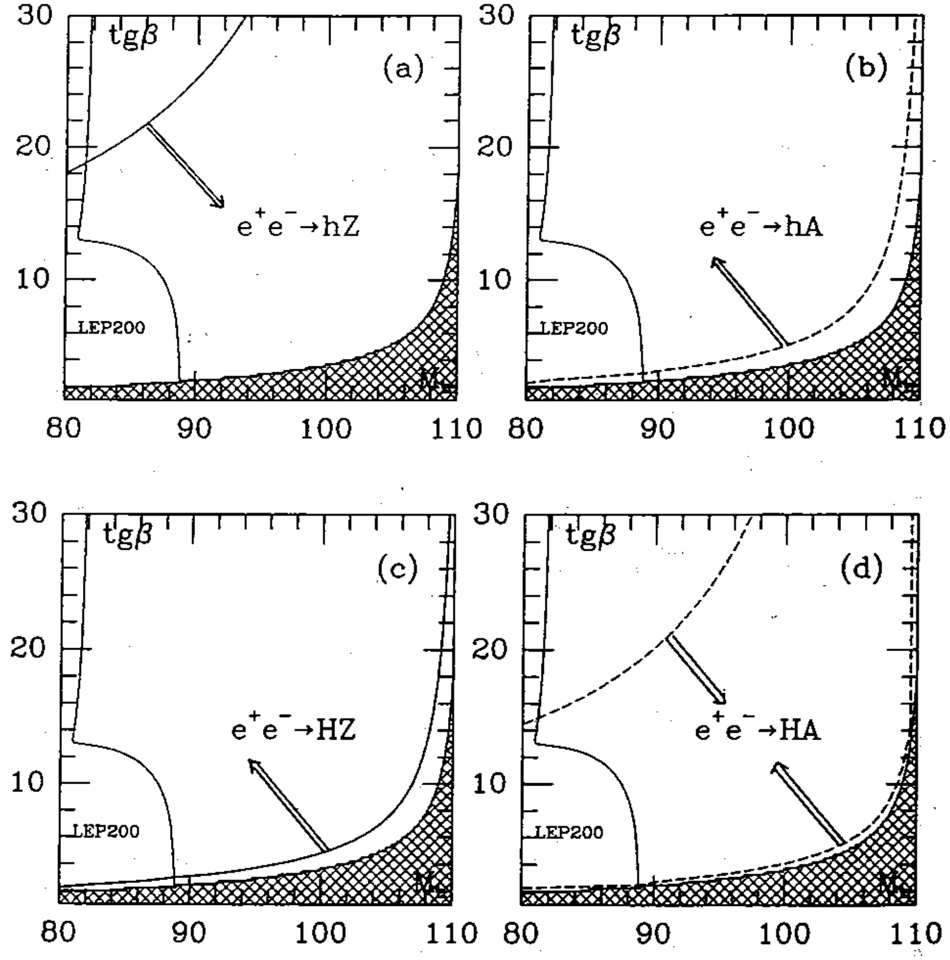

Fig. 9 Regions of the $[M_h, \mathrm{tg}\beta]$ plane where the four cross sections $e^+e^- \to hZ, hA, HZ$ and HA are observable. The contours are defined such that the cross sections are larger than 2.5 fb [corresponding to 25 events for an integrated luminosity of $\int \mathcal{L} = 10\ \mathrm{fb}^{-1}$]. The dashed area corresponds to the theoretically forbidden region $[M_h^2 \not> M_Z^2 \cos^2 2\beta + \epsilon \sin^2 \beta]$. The region which can be probed at LEP200 [defined such that for $\sqrt{s} = 180$ GeV and $\int \mathcal{L} = 500\ \mathrm{pb}^{-1}$, the number of events in one of the two processes $e^+e^- \to hZ$ or hA is larger than 25] is the area to the left of the thin line. The process $e^+e^- \to hZ$ is accessible in the entire area below the full line (a), hA in the entire area above the broken line (b) and HZ in the entire area above the full line (c); HA final states can be detected in the area between the two dashed lines (d).

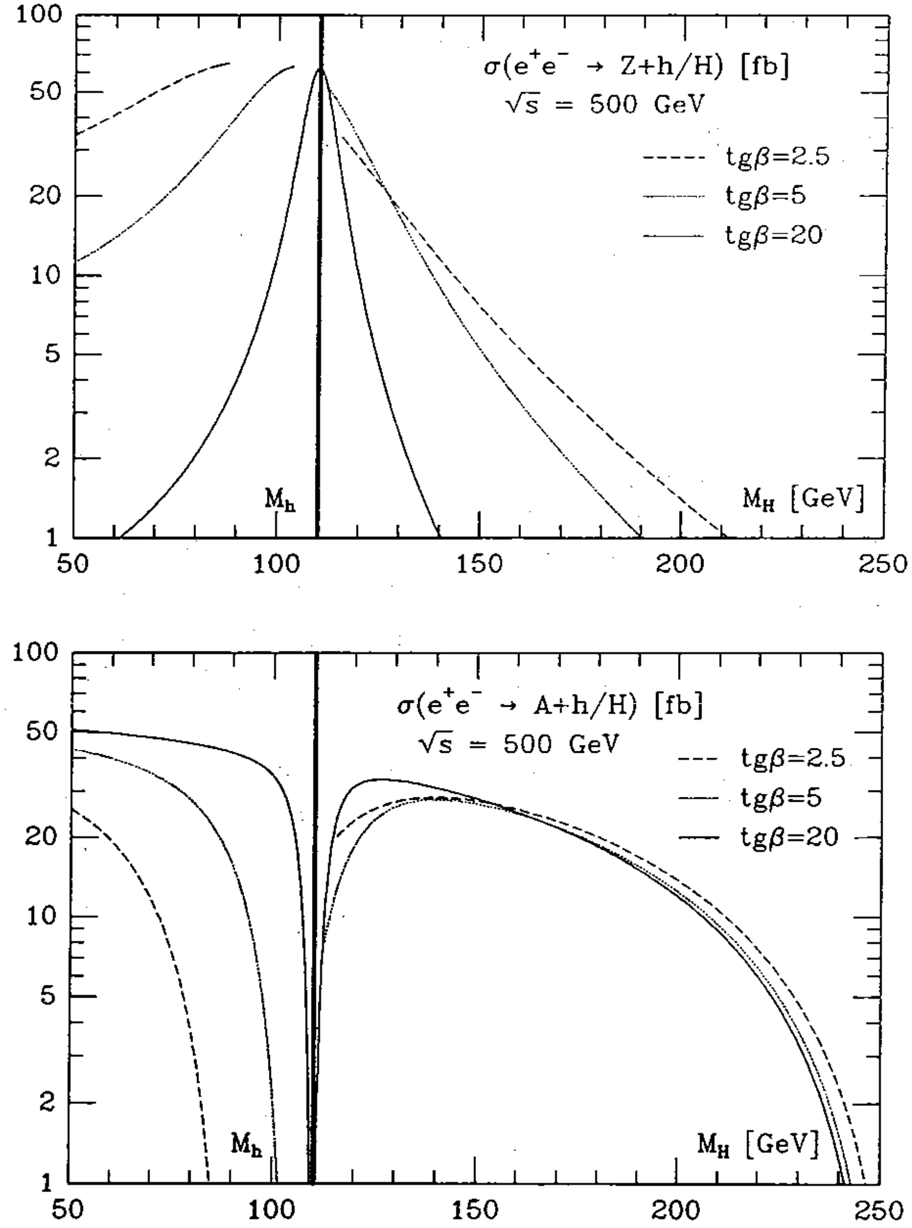

Fig. 10 Production cross sections of the $\mathcal{CP}$–even neutral Higgs bosons at $\sqrt{s}$ = 500 GeV as functions of their masses for three values of tgβ = 2.5, 5 and 20; (a) Bremsstrahl processes $e^+e^- \to Z + h/H$, and (b) in association with the pseudoscalar Higgs boson $e^+e^- \to A + h/H$.

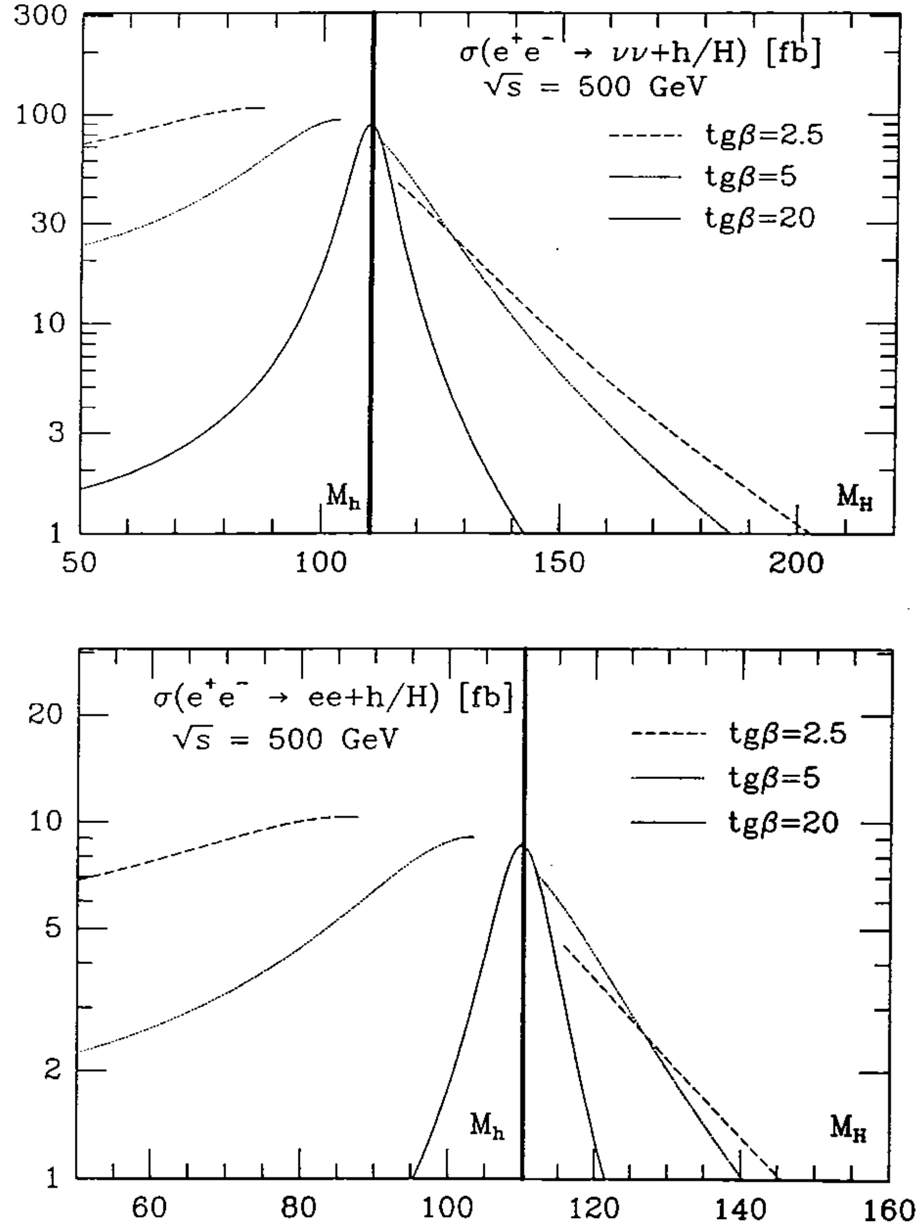

Fig. 11 Production cross sections of the $\mathcal{CP}$–even neutral Higgs bosons at $\sqrt{s}$ = 500 GeV as functions of their masses for three values of tgβ = 2.5, 5 and 20; (a) WW fusion process $e^+e^- \to \nu\bar{\nu} + h/H$, and (b) ZZ fusion process $e^+e^- \to e^+e^- + h/H$.

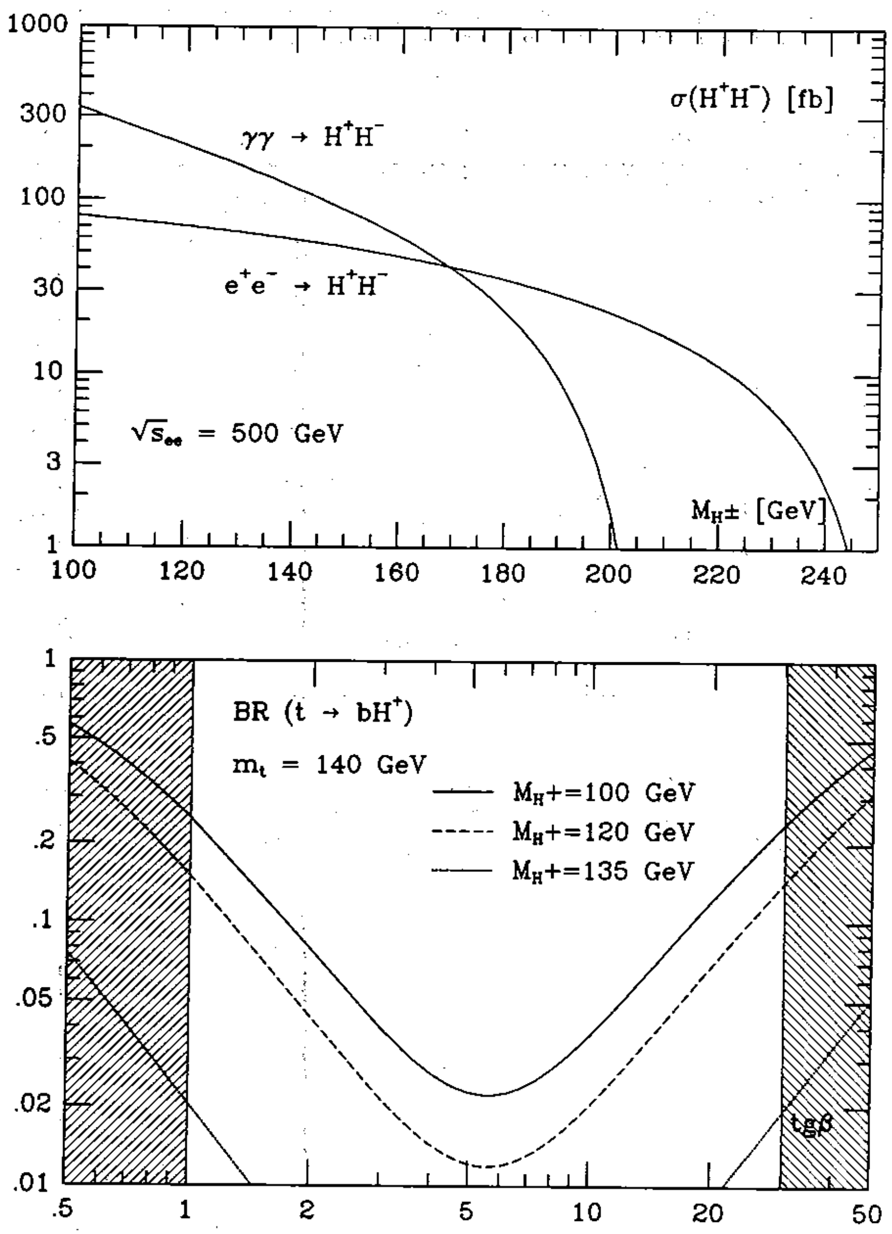

Fig. 12 (a) Production cross sections of the charged Higgs boson at e^+e^- colliders with $\sqrt{s} = 500$ GeV, and in $\gamma\gamma$ collisions generated by [unpolarized] Compton back-scattering; (b) Branching ratio $\mathrm{BR}(t \to bH^+)$ as a function of tgβ for $m_t = 140$ GeV and $M_{H^\pm} = 100, 120$ and 135 GeV. [The vertical lines delimit the regions of tgβ values favored in the $\mathcal{MSSM}$.]