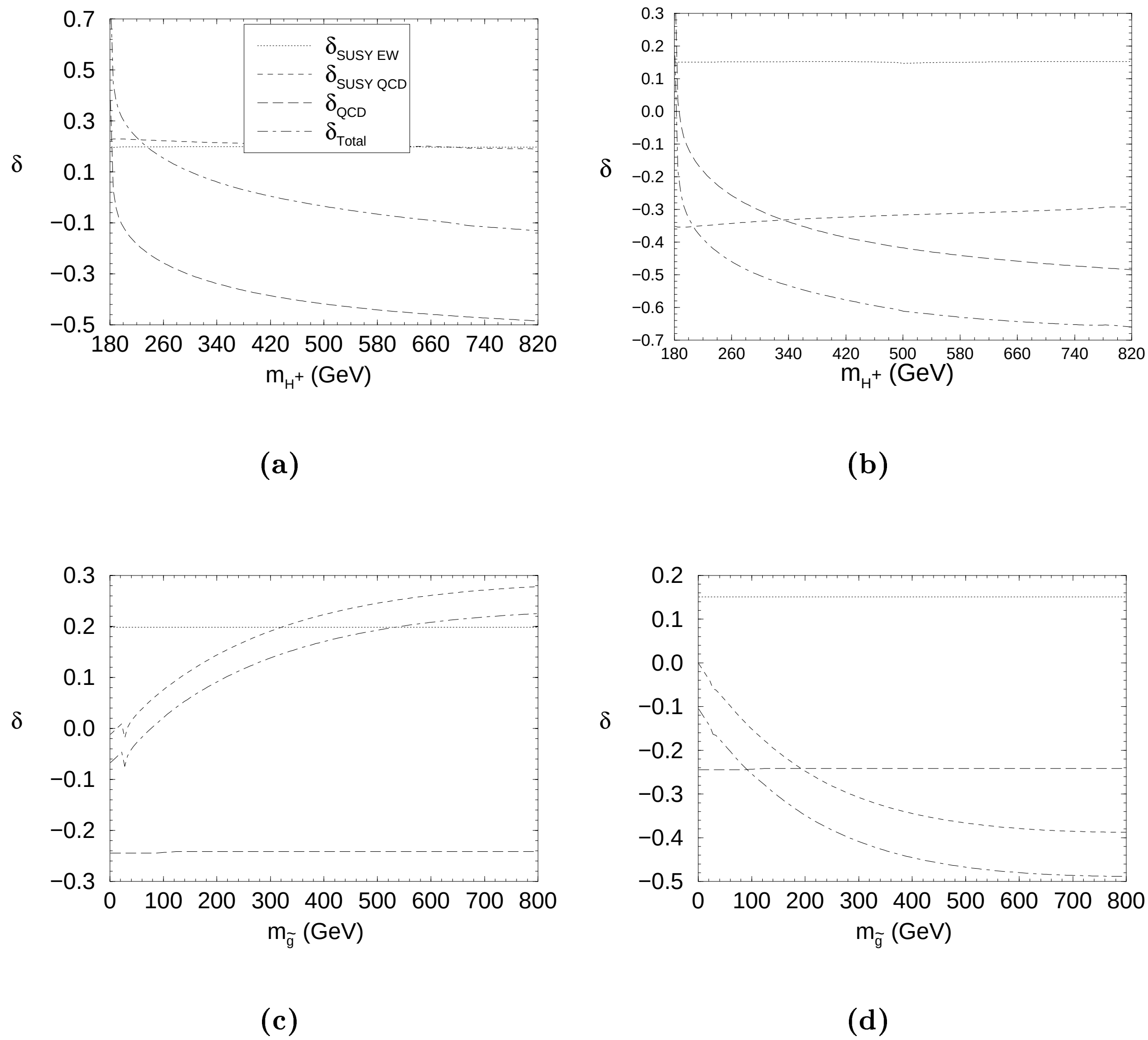


Figure 4.24: **(a)** *The SUSY-EW, SUSY-QCD, standard QCD and full MSSM contributions to* δ*, eq.(4.67), as a function of* M_{H^+}*. Inputs as in Fig.4.23.* **(b)** *As in (a), but for the positive* μ *case. Inputs as in Fig.4.25.* **(c)** *As in (a), but as a function of* $m_{\tilde{g}}$ *Remaining.* **(d)** *As in (b), but as a function of* $m_{\tilde{g}}$ *Remaining.*

important, the direct SUSY decay $H^+ \to \tilde{t}_i\,\bar{\tilde{b}}_j$ mentioned above is blocked up kinematically and plays no role in our analysis. On the other hand, the SUSY-EW output is basically controlled by the lightest stop mass, as it is plain in Fig.4.23d, where we vary it in a range past the LEP 200 threshold.

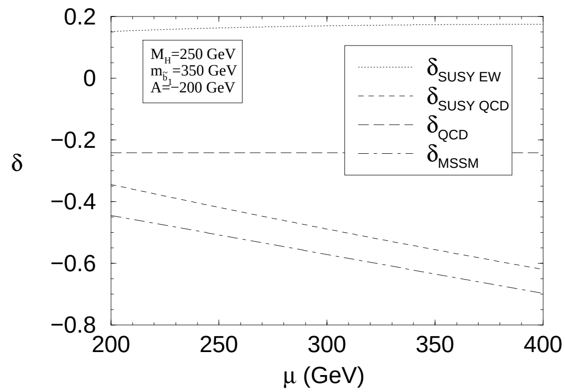
We have also checked that in the alternative $\mu > 0$, $A_t < 0$ scenario (also admissible according to Figs. 4.5-4.6), the SUSY-QCD correction is negative but it is largely cancelled by the SUSY-EW part, which stays positive, so that the total δ_{MSSM} is negative and larger (in absolute value) than the standard QCD correction. The results for this case are shown in Fig. 4.25.

Finally, coming back to Fig.4.22 we remark that if we take the standard QCD-corrected branching ratio (central curve in that figure) as a fiducial quantity, rather than the corresponding tree-level result, then $BR(H^+ \to \tau^+\,\nu_\tau)$ undergoes an effective MSSM correction of order $\pm(40-50)\%$. The sign of this effect is given by the sign of μ. In practice, $BR(H^+ \to \tau^+\,\nu_\tau)$ should be directly measurable from the cross-section for τ-production [56–58].

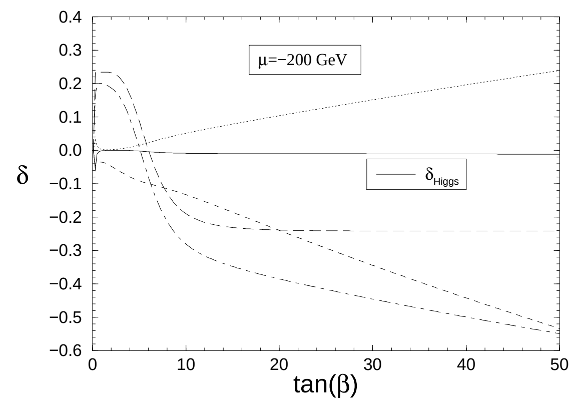
4.6 Conclusions

To summarize, we have presented a fairly complete treatment of the supersymmetric quantum effects ($\widetilde{\text{QCD}}$ and $\widetilde{\text{EW}}$) on the decay width of $H^+ \to t\,\bar{b}$ and have put forward plenty of evidence that they could be sizable enough to seriously compete with the ordinary QCD corrections. Consequently, they can either reinforce the conventional QCD corrections or counterbalance them, and even reverse their sign; the QCD corrections would then be found much "larger", "missing" or with the "wrong" sign, respectively. This should be helpful to differentiate H^+ from alternative charged pseudoscalar decays leading to the same final states.

Furthermore, our computation shows that these effects are compatible with CLEO data from low-energy B-meson phenomenology. The present study completes preliminary supersymmetric treatments where only the SUSY-QCD corrections were calculated [40,41,47,153]



(a)



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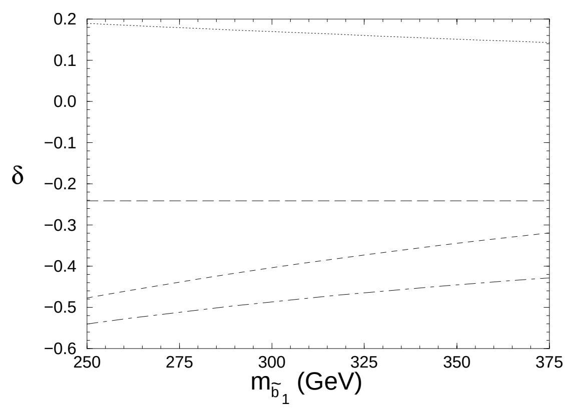
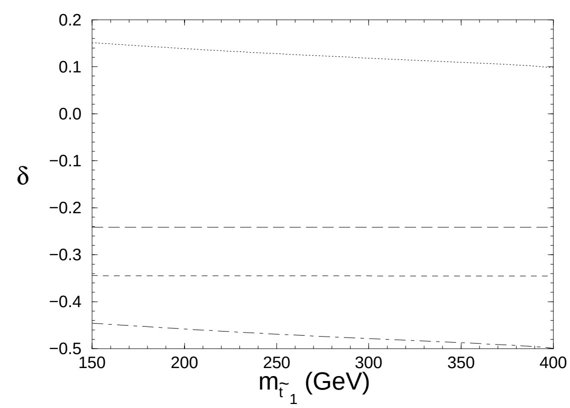


(c)



(d)

Figure 4.25: *The SUSY-EW, SUSY-QCD, standard QCD and full MSSM contributions to δ, eq.(5.6), as a function of μ;* **(b)** *As in (a), but as a function of* $\tan\beta$*. Also shown in (b) is the Higgs contribution, δ_{Higgs};* **(c)** *As in (a), but as a function of $m_{\tilde{b}_1}$;* **(d)** *As a function of $m_{\tilde{t}_1}$. Remaining inputs as in Fig.4.22.*

within the $(b \to s\,\gamma)$-unconstrained MSSM parameter space. Here we have evaluated the leading SUSY-EW effects and combined them with the SUSY-QCD ones both within the domain of compatibility with $b \to s\,\gamma$. As a result, we confirm that also in the constrained case the SUSY-QCD effects are generally very important, that is, $\widetilde{\rm QCD}$ corrections can be rather large (typically between $10\% - 50\%$), slowly decoupling and of both signs [40]. However, we have exemplified an scenario with sparticle masses above the LEP 200 discovery range where the SUSY electroweak corrections triggered by large Yukawa couplings can be comparable to the SUSY-QCD effects. In this context the total SUSY correction remains fairly large –around $+(30-50)\%$– with a $\sim 50\%$ component from electroweak supersymmetric origin. This situation occurs for

- large $\tan\beta$ (> 20),
- huge sbottom masses ($> 300\,GeV$) and
- relatively light stop and charginos ($100 - 200\,GeV$).

If the charged Higgs mass lies in the intermediate window (4.9), a chance is still left for Tevatron to produce a charged Higgs heavier than the top quark by means of "charged Higgsstrahlung" off top and bottom quarks. Should, however, a heavier $H^\pm$ exist outside the window (4.9), the LHC could continue the searching task mainly from gluon-gluon fusion where again $H^\pm$ is produced in association with the top quark.

The upshot is that the whole range of charged Higgs masses up to about $1\,TeV$ could be probed and, within the present renormalization framework, its potential supersymmetric nature be unravelled through a measurement of $\Gamma(H^+ \to t\,\bar{b})$ with a modest precision of $\sim 20\%$. Alternatively, one could look for indirect SUSY quantum effects on the branching ratio of $H^+ \to \tau^+\,\nu_\tau$ by measuring this observable to within a similar degree of precision.

Summarising, we have shown that the mechanisms capable of producing a charged Higgs scalar in a hadron collider, e.g. $t\,\bar{b}$ and $b\,\bar{t}$ fusion, which can be greatly enhanced at large $\tan\beta$, are very sensitive to potentially large SUSY quantum effects. If these effects are eventually found, they could be the smoking gun needed to recognize that the produced $H^\pm$ in a hadron collider is, truly, a SUSY Higgs.

This study is complementary to a previous one [48] in which $H^{\pm}$ was light enough for $t \to H^{+}b$ to be allowed.

Chapter 5

Strong effects on the hadronic widths of the neutral Higgs bosons in the MSSM

We analyze the correlation of QCD one-loop effects on the partial widths of the three neutral Higgs bosons of the MSSM decaying into quark-antiquark pairs. The SUSY-QCD effects turn out to be comparable or even larger than the standard QCD effects and are slowly decoupling in a wide window of the parameter space. Our study is aimed at elucidating the possible supersymmetric nature of the neutral Higgs bosons that might be discovered in the near future at the Tevatron and/or at the LHC. In particular, we point out the presence of potentially large SUSY corrections to the various neutral Higgs production cross-sections.

5.1 Motivation

To glimpse into the relevance of addressing the issue of the width of a Higgs boson, notice that if a heavy neutral Higgs is discovered and is found to have a narrow width, it would certainly not be the SM Higgs, whilst it could be a SUSY Higgs. In fact, a heavy enough SM Higgs boson is expected to rapidly develop a broad width through decays into gauge boson pairs whereas the SUSY Higgs bosons cannot in general be that broad since their couplings to gauge bosons are well-known to be suppressed [67]. In compensation, their couplings to fermions (especially to heavy quarks) can be considerably augmented. Thus the width of a SUSY Higgs should to a great extent be given by its hadronic width; even so a heavy H^0 and A^0 is in general narrower than a SM Higgs of the same mass. Alternatively, if the discovered neutral Higgs is sufficiently light that it cannot decay into gauge boson pairs, its decay width into relatively heavy fermion pairs such as $\tau^+\tau^-$, and especially into $b\bar{b}$, could be much larger than that of the SM Higgs, because of $\tan\beta$-enhancement of the fermion couplings [67]. Hence, it becomes clear that the hadronic width may play a very important role in the study of the MSSM higgses, already at the tree-level.

The aim of this work is to complete the analysis of the strong SUSY corrections to the hadronic decay widths of the Higgs bosons of the MSSM that was initiated in [40, 47]. In the latter reference two interesting domains of the general MSSM parameter were defined, the so-called Regions I and II, where the physics of the supersymmetric Higgs bosons can be especially relevant. Region I is characterized by a large value of $\tan\beta(> 30)$ in conjunction with a moderate value of both the CP-odd Higgs and a light chargino (a light stop, $m_{\tilde{t}_1}$, is not strictly needed in this region); and Region II by a large value of the CP-odd Higgs mass, m_{A^0}, of a few hundred GeV, by a relative light chargino and stop and by a moderate value of $\tan\beta$, $2 \lesssim \tan\beta \lesssim 20$. In our case, we will not limit to this regions exclusively.

Depending on the region of parameter space considered, not all the Higgs particles of the MSSM are allowed to decay hadronically in a significant manner. On the one hand, the process $H^+ \to t\,\bar{b}$, which requires $M_{H^+} > m_t + m_b \sim 180\,GeV$, is permitted in Region II and the SUSY-QCD corrections can be relevant in that region [40]. When $H^+ \to t\,\bar{b}$ is allowed, such a decay is by far the main hadronic decay mode of a SUSY charged Higgs boson. If, however, the condition $M_{H^+} > m_t + m_b$ is not satisfied, the remaining hadronic decays available to H^+ are not so appealing since the corresponding branching fractions lie below the leptonic τ-mode $H^+ \to \tau^+\nu_\tau$ as seen in Chapter 4.

On the other hand, we may turn our attention to the various hadronic neutral Higgs decays $\Phi^i \to q\bar{q}$ ($\Phi^1 \equiv A^0$, $\Phi^2 \equiv h^0$, $\Phi^3 \equiv H^0$). Of these, we will neglect the decays leading to light $q\,\bar{q}$ final states since their branching ratio is very small. Thus, for the lightest neutral

Higgs, h^0, we will concentrate on just the decay $h^0 \to b\,\bar{b}$, whereas for A^0, H^0 (which can be arbitrarily heavy) we shall consider the channels $A^0, H^0 \to b\,\bar{b}$ and $A^0, H^0 \to t\,\bar{t}$.

Moreover, should the physical domain of the MSSM parameter space turn out to lie in Regions I or II, then we shall see that the hadronic widths of the MSSM Higgs bosons must incorporate important virtual SUSY signatures. The latters could be extracted from measured quantities by subtracting the corresponding conventional QCD effects, which can be easily computed by adapting the results of Refs. [154, 155]. Knowledge of the SUSY corrections could be determinant to trace the nature of those scalars and establish whether they are truly supersymmetric Higgs particles.

Although we elaborate here mostly on the Higgs strategies at hadron colliders, such as the Tevatron and especially the LHC, it should be clear that the kind of effects that we wish to study have an impact on Higgs physics in $e^+\,e^-$ machines as well, where the neutral Higgs states can be produced through e.g. $e^+\,e^- \to Z\,h^0(H^0)$ and $e^+\,e^- \to A\,h^0(H^0)$. The observed cross-sections for these processes are equal to the production cross-sections times the Higgs branching ratios. Thus, in an $e^+\,e^-$ environment one aims more at a measurement of the various branching ratios (or, more precisely: ratios of branching ratios) of the fermionic Higgs decay modes rather than of the partial widths themselves. For instance, in $e^+\,e^-$ one would naturally address the measurement of $BR(\Phi^i \to b\,\bar{b})/BR(\Phi^i \to \tau^+\,\tau^-)$; in fact, this observable should receive large SUSY-QCD corrections if $\Phi^i \to b\,\bar{b}$ proves to be, as we shall see, very sensitive to the strong supersymmetric effects. This could have an impact for future LC-physics [156]

In hadron machines an actual measurement of the hadronic partial widths and in general of the effective hadronic vertices $\Phi^i\,q\,\bar{q}\,(q = t, b)$ should be feasible. Let us briefly remind of the five basic mechanisms for neutral Higgs production in a hadron collider [157]. They have been primarily described for the SM Higgs, H^0_{SM}, but can be straightforwardly extended to the three neutral higgses, Φ^i, of any two-doublet Higgs sector (see Fig. 5.1a. for a sketch of some of these mechanisms):

- **(i)** Gluon-gluon fusion: $g\,g \to \Phi^i$;
- **(ii)** $WW(ZZ)$ fusion: $q\,q \to q\,q\,\Phi^i$;

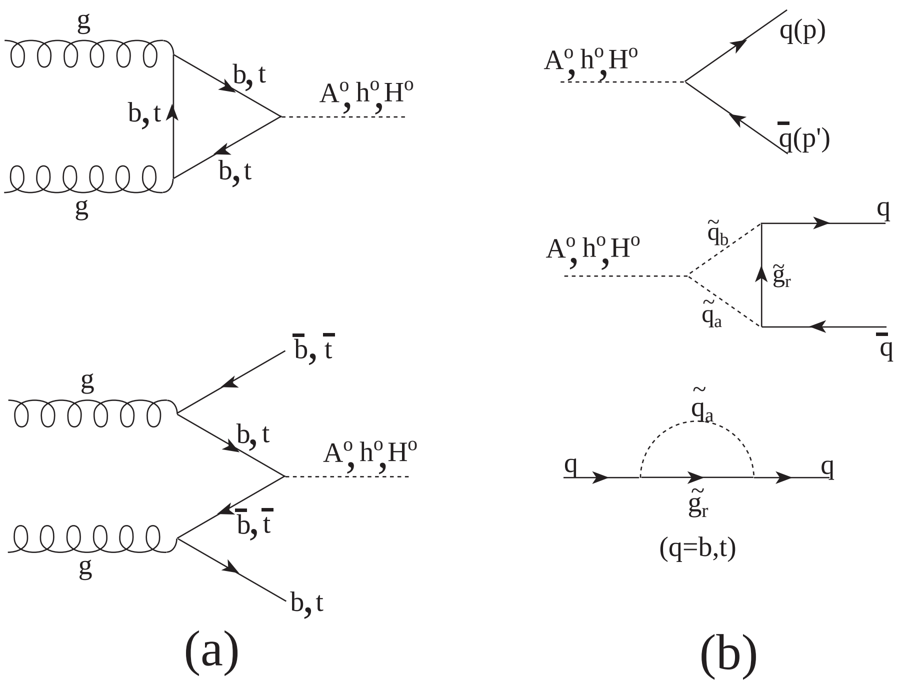


Figure 5.1: **(a)** *Typical mechanisms for neutral Higgs production at hadron colliders;* **(b)** *SUSY-QCD Feynman diagrams, up to one-loop level, correcting the partial widths of* $A^0, h^0, H^0 \to q\,\bar{q}$*. Our proposal to compute this latter process is now realised in the literature [50]*

- **(iii)** Associated $W(Z)$ production: $q\,\bar{q} \to W(Z)\,\Phi^i$;
- **(iv)** $t\,\bar{t}$ fusion: $g\,g \to t\,\bar{t}\,\Phi^i$, and
- **(v)** $b\bar{b}$ fusion: $g\,g \to b\,\bar{b}\,\Phi^i$.

It has been known for a long time [158] that in the SM, where only one neutral Higgs H^0_{SM} is present, mechanism (i) provides the dominant contribution over most of the accessible range, but as a drawback it has too big backgrounds from QCD. For very large (obese) SM Higgs mass, however, mechanism (ii) eventually takes over; the rest of the mechanisms are subleading, and in particular $b\,\bar{b}$ fusion is negligible in the SM.

Remarkably enough, this situation could drastically change in the MSSM. As noted above for the Higgs decays, also the production mechanisms of the MSSM Higgs scalars can be very different from the SM [157]. For instance, whereas one-loop $g\,g$-fusion in the SM is dominated by a top quark in the loop, this is not always so in the MSSM where the new couplings turn out to enhance, at high $\tan\beta$, the b-quark loops and make them fully competitive with the top quark loops (Fig. 5.1a). Furthermore, mechanism (ii) becomes suppressed by the SUSY couplings; e.g. in Region I the lightest neutral Higgs, h^0, has a very small coupling to the weak gauge bosons as compared to H^0_{SM}. In this respect the situation with the CP-odd scalar, A^0, is even worse, for it can never be produced by mechanism (ii) at the tree-level. In contrast, $b\,\bar{b}$ fusion (Fig. 5.1a), which was negligible in the SM, can be very important in the MSSM at large $\tan\beta$, especially in Region I but also in Region II. As a matter of fact, for large enough $\tan\beta$, the $b\,\bar{b}$-fusion cross-section can be larger than that for any mechanism for producing a SM Higgs boson of similar mass [157].

Our interest in the production mechanisms mentioned above stems from the fact that the radiative effects could play a crucial role. This is true already in the SM. For example, the conventional QCD corrections to $g\,g \to H^0_{SM}$, which is the dominant process for the production of a light and an intermediately heavy Higgs boson, are known to be large [159, 160]. A similar conclusion holds for an obese SM Higgs boson produced at very high energies by means of the $WW(ZZ)$-fusion mechanisms; here, again, non-negligible radiative effects do appear [161, 162]. Therefore, the production cross-section for H^0_{SM} is expected to acquire valuable quantum corrections both for light and for heavy Higgs masses. This is not so, however, for the corresponding width. In fact, only for a heavy SM Higgs, namely, with a mass above the vector boson thresholds, the corrections to its decay width can be of interest; for a light SM Higgs, instead, light enough that it cannot decay into gauge boson pairs, the decay width is very small and thus the corresponding quantum effects are of no practical interest.

Now, in contradistinction to the SM case, the hadronic vertices $\Phi^i\, q\, \bar{q}\, (q = t, b)$ could be the most significant interactions for MSSM higgses, irrespective of the value of the Higgs masses. In fact, these vertices can be greatly enhanced. Therefore, if large radiative corrections may modify the effective structure of these interactions, it is clear that they should be taken into account and could be of much practical interest. In what follows we shall substantiate that in the MSSM the $\Phi^i\, b\bar{b}$ and $\Phi^i\, t\bar{t}$ vertices involved in mechanisms (i), (iv) and (v) above could receive very large SUSY-QCD corrections (in some cases above 50%) and so, if these effects are there, they will be reflected in the Higgs boson partial widths and to a large extent also in the production cross-sections. In this respect the aforementioned $b\,\bar{b}$-fusion mechanism, which is highly operative at large $\tan\beta$, can be very sensitive to these SUSY-QCD corrections. To our knowledge, these matters have not been discussed in the literature and could play a momentous role to decide whether a neutral Higgs hypothetically produced in a hadron collider is supersymmetric or not.

While it goes beyond the scope of this note to compute the SUSY-QCD corrections to the production processes themselves (a recent estimation can be found in [50]), we have performed a detailed analysis of the partial decay widths, which are the canonical observables that should be first addressed to probe the new quantum corrections to the basic interaction vertices. In this way, a definite prediction is made on the properties of a physical observable, and moreover this should suffice both to exhibit the relevance of the SUSY quantum effects and to demonstrate the necessity to incorporate these corrections in a future, truly comprehensive, analysis of the cross-sections, namely, an analysis where one would include the quantum effects on all the relevant production mechanisms within the framework of the MSSM. This has been recently recognized at the SUSY/Higgs workshop at Fermilab (November, 1998) where the process $q\bar{q} + gg \to hb\bar{b} \to b\bar{b}b\bar{b}$ is expected to be the leading mode for SUSY-Higgs production at high $\tan\beta$ at the Tevatron.

5.2 One loop corrected hadronic neutral Higgs bosons decay width

Let us now concentrate on the diagrams depicted in Fig. 5.1b. The interaction Lagrangian describing the $\Phi^i\, q\, \bar{q}$-vertex in the MSSM is:

$$\mathcal{L}_{\Phi qq} = \frac{g m_q}{2M_W} \Phi^i \bar{q} \left[a^i_L(q) P_L + a^i_R(q) P_R \right] q \,. \tag{5.1}$$

We shall focus on top and bottom quarks ($q = t, b$). In a condensed and self-explaining notation we have defined

$$\begin{aligned} a^1_R(t,b) &= -a^1_L(t,b) = (i\cot\beta, i\tan\beta)\,, \\ a^2_R(t,b) &= a^2_L(t,b) = (-c_\alpha/s_\beta, s_\alpha/c_\beta)\,, \\ a^3_R(t,b) &= a^3_L(t,b) = (-s_\alpha/s_\beta, -c_\alpha/c_\beta)\,, \end{aligned} \tag{5.2}$$

with $c_\alpha \equiv \cos\alpha, s_\beta \equiv \sin\beta$ etc. (Angles α and β are related in the usual manner prescribed by the MSSM [67].) Apart from the SUSY-QCD interactions involving gluinos and squarks, a very relevant piece of our calculation is the interaction Lagrangian between neutral higgses and squarks (eq. 2.35).

The one-loop renormalized vertices for any of the relevant hadronic decays $\Phi^i \to q\,\bar{q}$ are derived from the on-shell renormalized Lagrangian and can be parametrized in terms of two bare form factors $K^i_L(q)$, $K^i_R(q)$ and the corresponding mass and wave-function renormalization counterterms δm_q and $\delta Z^q_{L,R}$ associated to the external quark lines:

$$O^i(q) = \frac{g\, m_q}{2M_W} \left[a^i_L(q) \left(1 + O^i_L(q)\right) P_L + a^i_R(q) \left(1 + O^i_R(q)\right) P_R \right] , \tag{5.3}$$

the renormalized form factors being

$$\begin{aligned} O^i_L(q) &= K^i_L(q) + \frac{\delta m_q}{m_q} + \frac{1}{2}\delta Z^q_L + \frac{1}{2}\delta Z^q_R\,, \\ O^i_R(q) &= K^i_R(q) + \frac{\delta m_q}{m_q} + \frac{1}{2}\delta Z^q_L + \frac{1}{2}\delta Z^q_R\,. \end{aligned} \tag{5.4}$$

For each $\Phi^i = A^0,\, h^0, H^0$ decaying into $q\,\bar{q}$ a straightforward calculation of the diagrams in Fig. 5.1b yields a generic contribution of the form (summation is understood over a, b)

$$K^i_L(q) = 8\pi\alpha_s i C_F \frac{G^{(q)}_{i\ ab}}{a^i_L(q)}\left[R^{(q)}_{1a}R^{(q)*}_{1b}\,(C_{11}-C_{12}) + R^{(q)}_{2a}R^{(q)*}_{2b}C_{12} + \frac{m_{\tilde{g}}}{m_q}R^{(q)}_{2a}R^{(q)*}_{1b}C_0\right],$$
$$K^i_R(q) = 8\pi\alpha_s i C_F \frac{G^{(q)}_{i\ ab}}{a^i_R(q)}\left[R^{(q)}_{2a}R^{(q)*}_{2b}\,(C_{11}-C_{12}) + R^{(q)}_{1a}R^{(q)*}_{1b}C_{12} + \frac{m_{\tilde{g}}}{m_q}R^{(q)}_{1a}R^{(q)*}_{2b}C_0\right]. \quad (5.5)$$

The explicit expressions for the mass and wave-function renormalization counterterms are borrowed from Sec. 4.4.3 and will not be repeated here, and the various 3-point functions in eq.(5.5) have the arguments $C_{...} = C_{...}(p, p', m_{\tilde{g}}, m_{\tilde{q}_a}, m_{\tilde{q}_b})$ [109]; they carry indices a, b summed over. Finally, $C_F = (N_C^2 - 1)/2N_C = 4/3$ follows from summation over color indices.

At the end of the day, the relative SUSY-QCD correction to each decay width of $\Phi^i \to q\,\bar{q}$ with respect to the corresponding tree-level width reads as follows,

$$\delta^i_{\tilde{g}}(q) = \frac{\Gamma^i(q) - \Gamma^i_0(q)}{\Gamma^i_0(q)} = Re[O^i_L(q) + O^i_R(q)]\,, \quad (5.6)$$

where $\Gamma^i(q) \equiv \Gamma(\Phi^i \to q\,\bar{q})$ is the corrected width, and

$$\Gamma^i_0(q) = \left(\frac{N_C G_F}{4\pi\tilde{q}rt2}\right)\left|a^i_R(q)\right|^2 M_{\Phi^i}\, m_q^2\,\lambda^{(1/2+j)}\,(1, \frac{m_q^2}{M^2_{\Phi^i}}, \frac{m_q^2}{M^2_{\Phi^i}})\,,$$
$$(j = 0 \ \text{ for } \ i = 1 \ \text{ and } \ j = 1 \ \text{ for } \ i = 2, 3)\,, \quad (5.7)$$

are the corresponding tree-level widths.

5.3 Numerical Analysis

The upshot of our exhaustive numerical analysis is synthesized in Figs. 5.2-5.5.

In Fig. 5.2a, where we fix $M_{A^0} = 60\,GeV$, we study the dependence of the SUSY-QCD correction (5.6) on the Higgs mixing mass parameter μ for the three decays $\Phi^i \to b\,\bar{b}$. We immediately gather that the sign of the correction is opposite to that of μ. For A^0 and h^0 the correction is basically the same and can reach large values, e.g. $|\delta_{\tilde{g}}| \simeq 50\%$ at $\tan\beta = 30$ and $|\mu| \simeq 100\,GeV$. As stressed in Ref. [40] and made obvious in Sec. 4.4.4, the origin of the

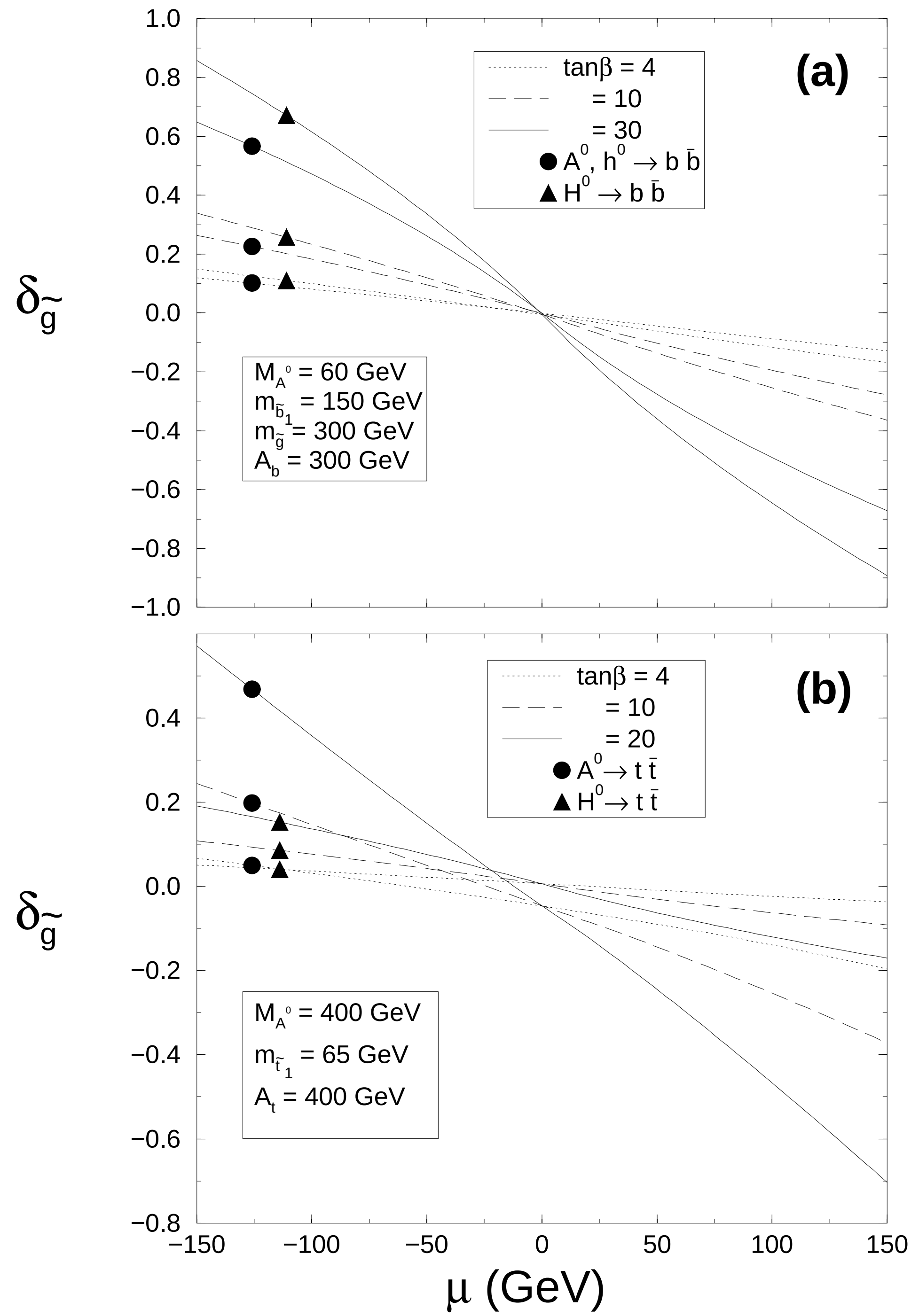


Figure 5.2: **(a)** *Dependence of the relative SUSY-QCD corrections* $\delta^i_{\tilde{g}}(q)$ *– eq.(5.6) – for* $i = 1, 2, 3$ *and* $q = b$*, upon the supersymmetric Higgs mass mixing term,* μ*, for light* $M_{A^0} = 60\,GeV$ *and given values of the other parameters (the scale of the abscissa is common to (b) below);* **(b)** *As before but for* $\delta^{1,3}_{\tilde{g}}(t)$ *and heavy* $M_{A^0} = 400\,GeV$*, for fixed* $m_{\tilde{t}_1} = 65\,GeV$*.*

large SUSY-QCD contributions obtained in the presence of final states involving the b-quark can be ascribed to their self-energy renormalization effects [133,134], which in our case go to the counterterm $\delta m_b/m_b$ on eq.(5.4). We remark that the corrections affecting the $b\,\bar{b}$ final states are larger for H^0 than for A^0 and h^0. The drawback, however, is that the huge effects obtained for $H^0 \rightarrow b\,\bar{b}$ at the highest values of $\tan\beta$ correspond to the smallest tree-level decay widths. In contrast, the less ambitious but still quite respectable quantum effects on $A^0 \rightarrow b\,\bar{b}$ are larger the larger is its decay width.

In Fig. 5.2b we study the alternative decays of A^0 and H^0 into $t\bar{t}$ final states for parameter values in Region II. Here, the minimum value of the lightest stop mass has to be specified and we take $m_{\tilde{t}_1} = 65\,GeV$. Even though $A^0 \rightarrow b\,\bar{b}$ is also dominant in Region II for the largest allowed values of $\tan\beta$ in this region, it has large QCD backgrounds. The heavy $t\,\bar{t}$ final states, however, are projected in the direction of the beam and can be identified through high-p_T leptons from semileptonic t-quark decays. Thus the heavy Higgs decays into $t\,\bar{t}$ final states, though they have a branching fraction smaller than that of the $b\,\bar{b}$ final states for $\tan\beta \gtrsim 6$, may in compensation be more manageable from the experimental point of view. For these decays we generally select more moderate values for $\tan\beta$ in order to make them sufficiently operative. From Fig. 5.2b we realize that of the two decays into $t\,\bar{t}$, the most sensitive to SUSY-QCD radiative corrections is that of the CP-odd Higgs boson. Here, in contradistinction to the b-quark final states, the main source of the corrections lies in the structure of the form factors $K_{L,R}$ on eqs.(5.4)-(5.5) – the top quark self-energies being negligible.

Of course, we expect –and we have numerically verified– that the SUSY-QCD contributions drop off upon freely increasing the squark masses. However, in practice the asymptotic regime begins for fairly large values of these masses. For example, in Fig. 5.3 we study the Higgs decays into $b\,\bar{b}$ as a function of $m_{\tilde{b}_1}$, for various $\tan\beta$. We see that, for $\tan\beta \gtrsim 10$, the corrections can reach several 10% even for $m_{\tilde{b}_1}$ in the few hundred GeV range.

Worth noticing is the asymptotic behaviour of the correction (5.6) versus the gluino mass for the various Higgs decays. As shown in Figs. 5.4a-5.4d, it takes a long time, so to speak, for the gluino to decouple. Corrections of $\sim 50-60\%$ for $A^0, h^0, H^0 \rightarrow b\,\bar{b}$ are obtained at

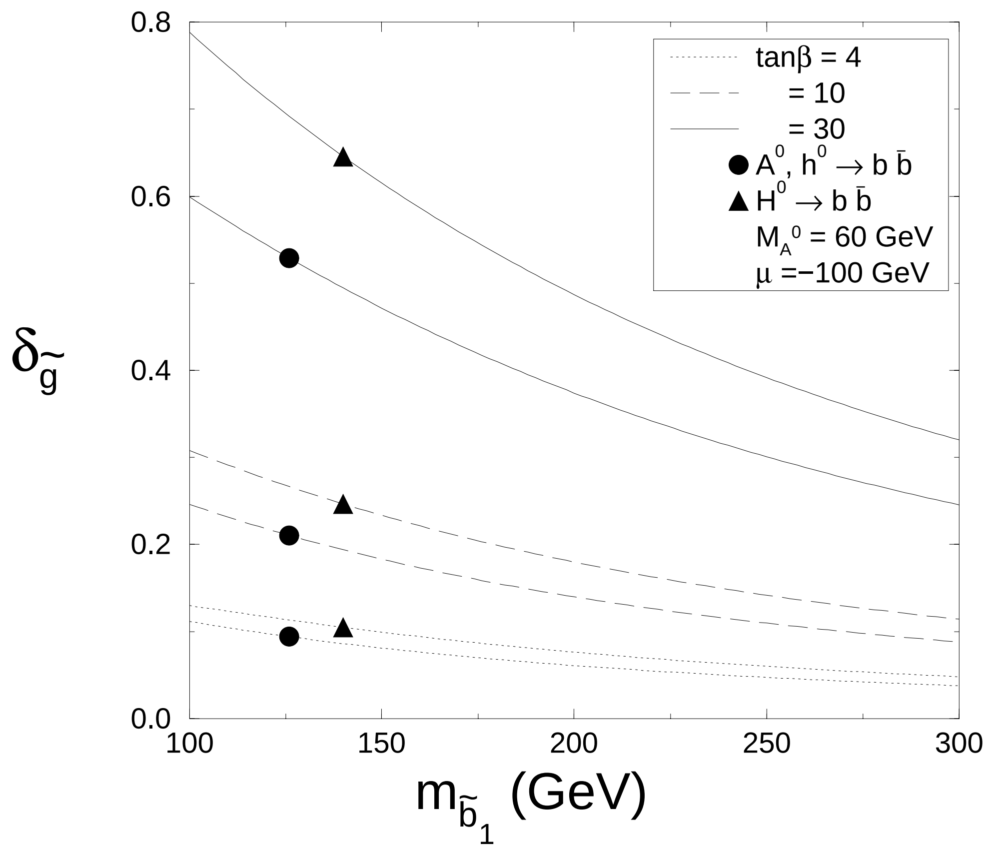


Figure 5.3: $\delta_{\tilde{g}}^{1,2,3}(b)$ *as a function of the lightest sbottom mass* $m_{\tilde{b}_1}$ *for various* $\tan\beta$ *and fixed* μ *and* M_{A^0}. *The other fixed parameters are as in Fig. 5.2.*

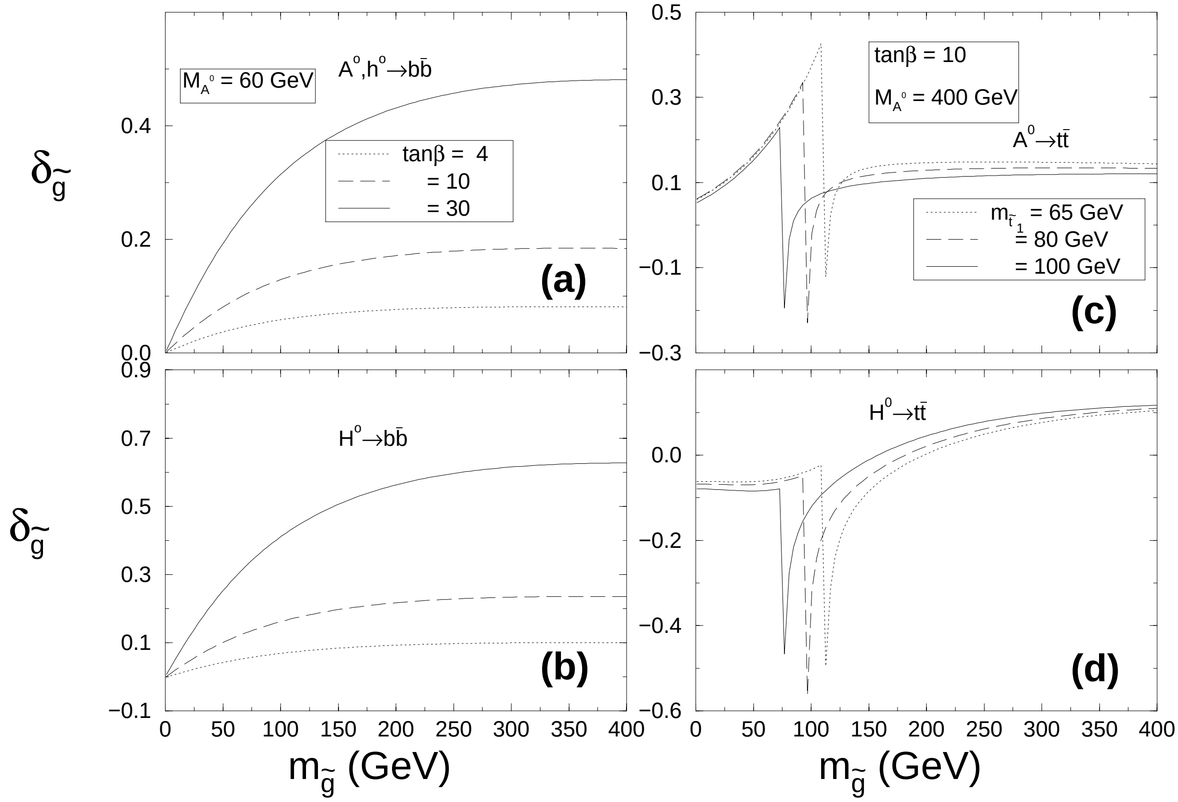


Figure 5.4: **(a)** *Evolution of* $\delta^1_{\tilde{g}}(b) \simeq \delta^2_{\tilde{g}}(b)$ *(almost indistinguishable) in terms of the gluino mass for various* $\tan\beta$*.* M_{A^0} *is chosen light and the remaining inputs as in Fig. 5.2 (the scale of the abscissa is common to (b) below);* **(b)** *As before but for* $\delta^3_{\tilde{g}}(b)$*;* **(c)** $\delta^1_{\tilde{g}}(t)$ *versus the gluino mass at fixed* $\tan\beta$ *and for various* $m_{\tilde{t}_1}$ *(the scale of the abscissa is common to (d) below);* **(d)** *As before but for* $\delta^3_{\tilde{g}}(t)$*.*

high $\tan\beta$ from a mass value $m_{\tilde{g}} \simeq 150\,GeV$ all the way out to $1\,TeV$ – hence far beyond the present phenomenological bounds. In Figs. 5.4c-5.4d we can also assess the dependence of $A^0, H^0 \to t\,\bar{t}$ on $m_{\tilde{t}_1}$, for fixed $m_{\tilde{b}_1} = 150\,GeV$ and a moderate value of $\tan\beta$; and we see that even for stop masses as heavy as $100\,GeV$ the corrections are longly sustained (above 10%) for practically any value of $m_{\tilde{g}}$ beyond the threshold singularities associated to points satisfying $m_{\tilde{g}} + m_{\tilde{t}_1} \simeq m_t$. For gluino masses below these points, the corrections to $A^0 \to t\,\bar{t}$ can be much larger.

From Figs. 5.5a and 5.5b we read off the dependence of the SUSY-QCD corrections on M_{A^0} for different values of $\tan\beta$, and they are compared with the ordinary QCD corrections. We remind the reader that the QCD corrections to $\Phi^i \to q\,\bar{q}$ can be very large for light quarks [154,155]. As in the decay of the charged Higgs [40], this is due to the appearance of a logarithmic term carrying a quark mass singularity, $\sim \log{(M_{\Phi^i}/m_q)}$, which stems from the anomalous dimension of the $\bar{q}\,q$ and $\bar{q}\,\gamma_5\,q$ operators. The complete one-loop (and renormalization group improved) formulae read as follows[1] ($b = \frac{33-2n_f}{6\pi}$):

$$\Gamma^i(q) = \Gamma^i_0(q)\left[1 - b\alpha_s(M_{\Phi^i})\log\frac{M_{\Phi^i}}{2m_q}\right]^{4/b\,\pi}\left\{1 + \frac{C_F\alpha_s(M_{\Phi^i})}{\pi}(\Delta_{\Phi^i} + 3\log\frac{M_{\Phi^i}}{2m_q})\right\}\,, \quad (5.8)$$

where the complicated functions Δ_{Φ^i} are given by eqs.(3.7) and (2.26) of Ref. [155] for $i = 1$ and $i = 2, 3$ respectively[2]. From eqs.(5.8) and (5.7) the standard QCD corrections $\delta^i_g = (\Gamma^i_{QCD}(q) - \Gamma^i_0)/\Gamma^i_0$ to the various MSSM neutral Higgs decays can be computed and are included in Figs. 5.5a and 5.5b, where they can be compared with the SUSY-QCD effects ($\delta^i_{\tilde{g}}$).

From Fig. 5.5a we see that the decays $\Phi^i \to b\,\bar{b}$ receive negative standard QCD corrections around $30 - 45\%$ (Notice that we have plotted $-\delta^i_g$ in Fig. 5.5a.). For $M_{A^0} \gtrsim 100\,GeV$, the standard QCD correction to $h^0 \to b\,\bar{b}$ remains saturated at about -30% since the mass M_{h^0} also saturates at its maximum value, whereas the modes $A^0, H^0 \to b\,\bar{b}$ obtain slowly increasing negative corrections. In contrast, $A^0 \to t\,\bar{t}$ and $H^0 \to t\,\bar{t}$, receive positive standard QCD corrections rapidly varying with the Higgs mass (Cf. Fig. 5.5b) Comparison with our

[1] This equation is equivalent to eq.(4.5) of Ref. [155], except that we have corrected a missing factor of 2 in the last logarithm.

[2] We have also corrected a missing factor of 2 in the third term on the RHS of eq.(2.27) of Ref. [155].

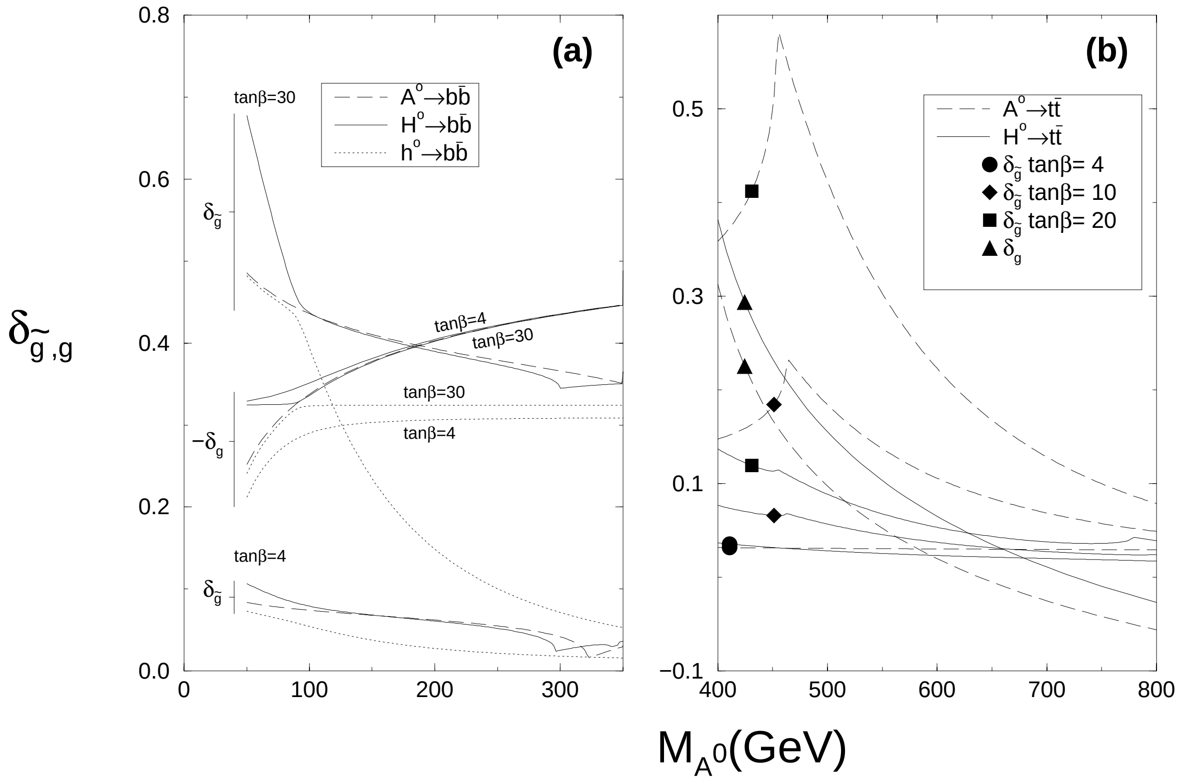


Figure 5.5: **(a)** $\delta_{\tilde{g}}^{1,2,3}(b)$ *for* $\tan\beta = 30$ *(upper-born curves) and* $\tan\beta = 4$ *(lower-born curves) as a function of* M_{A^0} *and the rest of the parameters as in Fig. 5.2. The middle-born curves stand for the corresponding standard QCD corrections,* δ_g^i*, to the three decay processes. The latters being negative, we plot* $-\delta_g^i$ *to ease comparison with the SUSY-QCD corrections at* $\mu < 0$*;* **(b)** *As before but for* $\delta_{\tilde{g}}^{1,3}(t)$ *and the range of* M_{A^0} *selected deep into Region II and three values of* $\tan\beta$*. The scale of the ordinate is common to (a).*

Figs. 5.2-5.5 clearly shows that in many cases the supersymmetric effects are important since they can be of the same order as the standard QCD corrections.

5.4 Conclusions

Overall, large SUSY-QCD quantum corrections are expected in the hadronic widths of the neutral Higgs bosons of the MSSM. They should be measurable, though with different techniques, both in $e^+ e^-$ and in hadron machines. These supersymmetric effects can not only be comparable to the ordinary QCD corrections, but even dominant in some cases. Since they can have either sign, the net QCD correction would be found either much “larger” than expected, perhaps “missing” or even with the “wrong” sign; in any case, it should be revealing to hint at the SUSY nature of these higgses. Furthermore, we have found that, contrary to the situation with SUSY corrections on gauge boson observables [94,111,163,164], these effects decouple very slowly, especially with the gluino mass. Therefore, if SUSY is there, these corrections should also be there, and cannot be missed for a wide range of sparticle masses. However, Region II is out of reach of LEP 200, and even though part of the Higgs spectrum characterizing Region I is within its discovery range a complete experimental account of the MSSM Higgs sector will not be possible at LEP 200. In this respect, we have put forward the convenience of trying to see the kind of effects studied here in the large hadron machines (Tev II, LHC), perhaps before an $e^+ e^-$ supercollider (LC) be at work. In fact, the potentially large size of these effects indicates that they ought to be included in any serious analysis of supersymmetric Higgs production processes in hadron colliders [44,45]. The combined information on branching ratios (from $e^+ e^-$) and on cross-sections (from the Tevatron and/or at the LHC) should be very useful to pin down the nature of the observed effects. A more complete study should include the electroweak SUSY effects (in particular, the effect from the one-loop Higgs mass relations), but as already mentioned in [40], these are expected not to alter qualitatively the SUSY-QCD picture presented here. Our general conclusion is that quantum corrections on Higgs physics may be the clue to “virtual” Supersymmetry.