Indirect heavy SUSY signals in Higgs and top decays

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We summarize the recent results on the supersymmetric QCD radiative corrections, at the one-loop level, in Higgs and top quark decays, in the context of the Minimal Supersymmetric Standard Model and in the decoupling limit of very heavy SUSY particles. Special attention is devoted to the particular decays $h^0 \rightarrow b\bar{b}$, $H^+ \rightarrow t\bar{b}$, $H^0 \rightarrow b\bar{b}$, $A^0 \rightarrow b\bar{b}$ and $t \rightarrow H^+\bar{b}$ where the radiative corrections from heavy squarks and heavy gluinos do not decouple and are enhanced at large $\tan \beta$. Some interesting phenomenological consequences are also briefly summarized.

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1 Introduction

One of the most challenging goals of the next generation colliders is the discovery of supersymmetric (SUSY) particles and the study of their rich associated phenomenology. In the context of the Minimal Supersymmetric Standard Model (MSSM) [1], the SUSY physical spectrum consists of squarks, $\tilde{q}_1, \tilde{q}_2$ ($q = u, d, s, c, b, t$), sleptons, $\tilde{l}_1, \tilde{l}_2$ ($l = e, \mu, \tau$), sneutrinos, $\tilde{\nu}_l$ ($l = e, \mu, \tau$), gluinos, $\tilde{g}_a$ ($a = 1, \ldots, 8$), charginos, $\tilde{\chi}^\pm_i$ ($i = 1, 2$) and neutralinos, $\tilde{\chi}^0_j$ ($j = 1, \ldots, 4$). In addition, the MSSM incorporates an extended Higgs sector with two Higgs doublets that lead to five physical Higgs boson particles: two CP-even scalars, $h^0$ and $H^0$, one CP-odd scalar, $A^0$, and two charged scalars, $H^\pm$ [2]. None of these non-standard particles have been discovered yet, and the present experiments have placed some lower mass limits. From the last published Review of Particle Physics [3] the following 95\%CL limits have been extracted: $M_{\tilde{q}} > 250\, GeV$ ($q \neq t, b$), $M_{\tilde{t}} > 86.4\, GeV$, $M_{\tilde{b}} > 75\, GeV$, $M_{\tilde{\chi}^\pm} > 87.1\, GeV$, $M_{\tilde{\chi}^0} > 82.3\, GeV$, $M_{\tilde{\chi}_1} > 81.0\, GeV$, $M_{\tilde{\chi}_2} > 43.1\, GeV$, $M_{\tilde{\chi}_3} > 190\, GeV$, $M_{\tilde{\chi}_4} > 67.7\, GeV$, $M_{\tilde{\chi}_5} > 32.5\, GeV$, $M_{\tilde{\chi}_6} > 82.6\, GeV$, $M_{\tilde{\chi}_7} > 84.1\, GeV$, $M_{H^\pm} > 69.0\, GeV$. Most of these limits are dependent on particular assumptions on the MSSM parameters and have been improved after that Review of Particle Physics [3].

The absence of SUSY particles in the present colliders still leaves the possibility that these particles manifest at energies larger than the present explored energies, being these latter typically of the order of the electroweak scale, $M_{EW} \sim O(245\, GeV)$. In this work we will assume that the whole spectrum of genuine SUSY particles is heavy, such that they can not be produced directly at the present or next generation colliders, and different strategies based on indirect searches must be performed.

The study of radiative corrections from SUSY particles to Standard Model (SM) couplings and SM observables may provide crucial clues in this search of indirect SUSY signals if the SUSY masses are beyond the reach of present and planned accelerators [4]. In particular, if a light Higgs boson, $h^0$, were discovered in the mass range predicted by the MSSM, $M_{h^0} \leq 135\, GeV$ [3], but the SUSY particles were not found, a precise measurement of Higgs couplings to SM particles, which are sensitive to radiative corrections, could provide indirect information about the existence of SUSY in Nature and even serve to infer the SUSY particle masses. For example, one could conclude (in the context of the MSSM) whether the data favor a SUSY spectrum below the 1 TeV energy scale. Similar studies can be performed by considering some relevant observables as, for instance, the partial widths of the top quark and Higgs boson $h^0$ decays into SM particles, and by comparing their predictions in the MSSM and the SM. Furthermore, in case the extra Higgs boson particles, $H^0, A^0$ and $H^\pm$, be

\footnote{The 2001 summer conferences (see for instance ref. [2]) have announced updated 95\%CL limits: $M_{\tilde{q}} > 300\, GeV$ ($q \neq t, b$), $M_{\tilde{t}} > 95\, GeV$, $M_{\tilde{b}} > 92\, GeV$, $M_{\tilde{\chi}^\pm} > 99\, GeV$, $M_{\tilde{\chi}^0} > 95\, GeV$, $M_{\tilde{\chi}_1} > 80.0\, GeV$, $M_{\tilde{\chi}_2} > 43.1\, GeV$, $M_{\tilde{\chi}_3} > 195\, GeV$, $M_{\tilde{\chi}_4} > 101\, GeV$, $M_{\tilde{\chi}_5} > 45.6\, GeV$, $M_{\tilde{\chi}_6} > 91.0\, GeV$, $M_{A^0} > 91.9\, GeV$, $M_{H^\pm} > 78.6\, GeV$.}
accessible as well to the present or future colliders, one can also study the sensitivity
to a heavy SUSY spectrum via the SUSY radiative corrections to their relevant partial
decay widths and compare their predictions with those of a more general two-Higgs-
doublets model (2HDM). There is an extensive literature on SUSY radiative corrections
to the decays of Higgs particles and the top quark, in the context of the MSSM.
We refer the reader to refs. [7,8,9,10,11,12,13,14,15,16,17,18,19,20,21,22,23,24,25] for
a selection of works closely related with the subject of this contribution.

In this review we will consider the dominant SUSY radiative corrections to the
relevant Higgs boson particles and top quark decays, at the one-loop level, that come
from the SUSY-QCD (SQCD) sector, that is from squarks and gluinos, and study
their behavior in the limit where the SUSY particles are very heavy as compared to the
electroweak scale, $M_{\text{SUSY}} \gg M_{\text{EW}}$, where $M_{\text{SUSY}}$ represents generically the masses
of the SUSY particles. This situation corresponds to the decoupling of SUSY particles
from the rest of the MSSM spectrum, namely, the SM particles and the MSSM Higgs
sector. Special attention will be devoted to the particular decays $h^0 \to b\bar{b}$, $H^+ \to t\bar{b}$,
$H^0 \to b\bar{b}$, $A^0 \to b\bar{b}$ and $t \to H^+ b$ where the SUSY-QCD radiative corrections from
heavy squarks and heavy gluinos turn out to be non-decoupling from the low energy
physics and, furthermore, they are enhanced at large $\tan \beta$. This non-decoupling
behavior is genuine of Higgs sector physics and has not been found yet in other
MSSM sectors, as for instance in electroweak gauge bosons physics [26]. It has not
been seen in the $h^0$ self-couplings either [27]. The interest of these non-decoupling
effects is that they can produce sizeable contributions to the Higgs particles and top
quark partial decay widths for large enough $\tan \beta$ values and, therefore, can provide
indirect SUSY signals, even for very heavy squarks and gluinos. We will briefly
comment on some recently proposed observables that are defined as ratios of Higgs
branching ratios into quarks divided by the corresponding Higgs branching ratios into
leptons, and that turn out to be the most sensitive to these SUSY non-decoupling
contributions. We will devote as well some comments to the alternative plausible
possibility where, not just the SUSY particles but also the extra Higgses are heavy.
This limiting situation occurs when both $M_{\text{SUSY}}$ and the CP-odd Higgs boson mass
$M_A$ are larger than $M_{\text{EW}}$, and the decoupling of all non-standard particles from the
SM physics is expected. In this limit, the lightest Higgs boson, $h^0$, resembles the
SM Higgs particle, and a distinction between the MSSM and the SM will be very
difficult [28].

We present here just a summary of the main results and refer the reader to
refs. [19,20,21,22,23,24,25,26] for more details.
2 Decoupling limit in the SUSY-QCD sector

In this section we shortly review the SUSY-QCD sector of the MSSM, consisting of squarks and gluinos, and define the decoupling limit, where these SUSY particles are much heavier than the electroweak scale. We will concentrate on the third-generation squarks since they provide the dominant radiative corrections to the Higgs particles and top quark decays.

The sbottom and stop mass matrices are given in terms of the MSSM parameters respectively by:

$$
\hat{M}_b^2 = \left( \begin{array}{cc} M_Q^2 + m_b^2 - M_Z^2 (\frac{1}{2} + Q_b s_w^2) \cos 2\beta & m_b (A_b - \mu \tan \beta) \\ m_b (A_b - \mu \tan \beta) & M_D^2 + m_b^2 + M_Z^2 Q_b s_w^2 \cos 2\beta \end{array} \right),
$$

and

$$
\hat{M}_t^2 = \left( \begin{array}{cc} M_Q^2 + m_t^2 + M_Z^2 (\frac{1}{2} - Q_t s_w^2) \cos 2\beta & m_t (A_t - \mu \cot \beta) \\ m_t (A_t - \mu \cot \beta) & M_U^2 + m_t^2 + M_Z^2 Q_t s_w^2 \cos 2\beta \end{array} \right),
$$

where $M_Q$, $M_D$, $M_U$, are the soft-SUSY-breaking mass parameters for the squark doublet $\tilde{q}_L$ and the squark singlets $\tilde{t}_R$ and $\tilde{b}_R$, respectively. $A_{t,b}$ are the stop and sbottom soft-SUSY-breaking trilinear couplings, respectively. The $\mu$-parameter is the SUSY-preserving bilinear Higgs coupling. The ratio of the two Higgs vacuum expectation values is given by $\tan \beta = v_2/v_1$. $M_Z$, $m_t$ and $m_b$ are the standard Z boson, top quark and bottom quark masses. $Q_t = 2/3$, $Q_b = -1/3$ and $s_w = \sin \theta_W$.

The physical squared masses of the sbottoms, $\hat{M}_{b1,2}^2$, and stops, $\hat{M}_{t1,2}^2$, are the eigenvalues of the previous mass matrices, $\hat{M}_b^2$ and $\hat{M}_t^2$, respectively. The mass of the gluinos, $\tilde{g}_a$ ($a = 1, \ldots, 8$), is given by the soft-SUSY-breaking Majorana mass $M_{\tilde{g}}$.

In order to get heavy squarks and heavy gluinos, we need to choose properly the soft-SUSY-breaking parameters and the $\mu$-parameter. Since here we are interested in the limiting situation where the whole SUSY spectrum is heavier than the electroweak scale, we have made the simplest assumption for the soft breaking squark mass parameters, trilinear terms, $\mu$-parameter and gluino mass (see refs. [19,20] for more details),

$$
M_{SUSY} \sim M_Q \sim M_D \sim M_U \sim M_{\tilde{g}} \sim \mu \sim A_b \sim A_t \gg M_{EW},
$$

where $M_{SUSY}$ represents generically a common SUSY large mass scale and the symbol ‘$\sim$’ means quantities of the same order of magnitude but not necessarily equal. In the following of this work, this general condition will be referred to as ‘large SUSY mass limit’. 

3
We have considered two extreme cases, maximal and minimal squark mixing, which, for the previous 'large SUSY mass limit', imply certain constraints on the squark mass differences. Thus, given the generic mass matrix,

\[ \hat{M}_q^2 \equiv \begin{pmatrix} M_{q,L}^2 & m_q X_q \\ m_q X_q & M_{q,R}^2 \end{pmatrix}, \]

the two limiting cases are reached by choosing the relative size of \( M_{q,L,R} \) and \( X_q \) as follows,

A.- Close to maximal mixing: \( \theta_{\tilde{q}} \sim \pm 45^\circ \)

\[ |M_{q,L}^2 - M_{q,R}^2| \ll m_q X_q \Rightarrow |\tilde{q}_1^2 - \tilde{q}_2^2| \ll |\tilde{q}_1^2 + \tilde{q}_2^2| \]

B.- Close to minimal mixing: \( \theta_{\tilde{q}} \sim 0^\circ \)

\[ |M_{q,L}^2 - M_{q,R}^2| \gg m_q X_q \Rightarrow |\tilde{q}_1^2 - \tilde{q}_2^2| \sim \mathcal{O}|\tilde{q}_1^2 + \tilde{q}_2^2| \]

where we have included the corresponding implications for the physical squark mass differences.

3 Decoupling limit in the Higgs sector

The decoupling limit in the Higgs sector of the MSSM was first studied in ref. \[28\]. In short, it is defined by considering the CP-odd Higgs mass much larger than the electroweak mass, \( M_A \gg M_Z \), and leads to a particular spectrum in the Higgs sector with very heavy \( H^0 \), \( H^\pm \) and \( A^0 \) bosons, and a light \( h^0 \) boson. For a review of the MSSM Higgs sector, see ref. \[3\].

At tree level, if \( M_A \gg M_Z \), the Higgs masses are,

\[ M_{h^0} \simeq M_{H^\pm} \simeq M_A \gg M_Z, \quad M_{h^0} \simeq M_Z |\cos 2\beta| . \]

That is, at tree-level there exists a CP-even Higgs, \( h^0 \), lighter than the \( Z \) boson.

| \( \phi \) | \( g_{\phi \bar{t} t} \) | \( g_{\phi \bar{b} b} \) | \( g_{\phi V V} \) |
|---|---|---|---|
| SM \( H \) | 1 | 1 | 1 |
| MSSM \( h^0 \) | \( \cos \alpha / \sin \beta \) | \( - \sin \alpha / \cos \beta \) | \( \sin (\beta - \alpha) \) |
| \( H^0 \) | \( \sin \alpha / \sin \beta \) | \( \cos \alpha / \cos \beta \) | \( \cos (\beta - \alpha) \) |
| \( A^0 \) | 1 / \( \tan \beta \) | \( \tan \beta \) | 0 |

Table 1: Higgs couplings in the MSSM normalized to SM couplings, in terms of \( \beta \) and the mixing angle of the neutral CP-even Higgs sector \( \alpha \). Here \( V = Z,W \).
Concerning the neutral Higgs couplings, their tree-level values in the MSSM, normalized to SM couplings and for arbitrary $M_A$, are given in table 1.

Notice that by expanding in inverse powers of $M_A$, we get:

$$
\frac{\cos \alpha}{\sin \beta} \sim 1 + \mathcal{O}(M_Z^2/M_A^2), \quad -\frac{\sin \alpha}{\cos \beta} \sim 1 + \mathcal{O}(M_Z^2/M_A^2)
$$

$$
\sin(\beta - \alpha) \sim 1 + \mathcal{O}(M_A^4/M_Z^4).
$$

Therefore, the $h^0$ tree-level couplings in the decoupling limit, $M_A \gg M_Z$, tend to their SM values, as expected.

Beyond tree level, it has been shown [28] that, in this same decoupling limit, the Higgs masses keep a similar pattern as at tree level, that is, very heavy $H^0$, $H^\pm$ and $A^0$ bosons, and a light $h^0$ boson. The particular values of their masses depend of course on the MSSM parameters, but for $M_A \gg M_Z$,

$$
M_{h^0} \simeq M_{H^\pm} \simeq M_A \gg M_Z, \quad \text{and} \quad M_{h^0} \leq 135 \text{GeV}.
$$

In this work we will go beyond tree level and study the decoupling behavior of heavy SUSY particles and heavy Higgses, at one-loop level.

4 Decoupling in low-energy electroweak gauge boson physics

It has been shown that all one-loop corrections to low-energy electroweak gauge boson physics involving SUSY particles and extra Higgs bosons decouple in the limit of large sparticle masses and large $M_A$ [26]. The formal proof of this decoupling involves the computation of the effective action for electroweak gauge bosons $\Gamma_{\text{eff}}[A, Z, W^\pm]$ by explicit integration, in the path integral, of all the heavy sparticles and heavy extra Higgs bosons, $\tilde{f}(\tilde{q}, \tilde{\bar{q}}, \tilde{\nu}, \tilde{\nu}^c), \tilde{\chi}^\pm_i, \tilde{\chi}_i^0, H(H = H^\pm, H^0, A^0)$ at the one-loop level,

$$
e^{i\Gamma_{\text{eff}}[V]} = \int [d\tilde{f}] [d\tilde{f}^*] [d\tilde{\chi}^+] [d\tilde{\chi}^-] [d\tilde{\chi}_i^0] [dH] e^{i\Gamma_{\text{MSSM}}[V, \tilde{f}, \tilde{\chi}^+, \tilde{\chi}_i^0, H]}; \quad V = A, Z, W^\pm.
$$

The proof involves, in addition, a large sparticle masses and large $M_A$ expansion of $\Gamma_{\text{eff}}[A, Z, W^\pm]$ that is valid in the heavy SUSY masses limit, $M_{\text{SUSY}} \sim m_{\tilde{f}}, m_{\tilde{\chi}_i^\pm}, m_{\tilde{\chi}_i^0}, M_A \gg M_{\text{EW}}$, with $|\tilde{m}_i^2 - \tilde{m}_j^2| \ll |\tilde{m}_i^2 + \tilde{m}_j^2|$ for $i \neq j$. The result of this expansion is shown schematically in Fig. 1.

All the effects from heavy SUSY particles and extra Higgses in the effective action for electroweak gauge bosons, or equivalently in the n-point functions, are either absorbed into redefinitions of the electroweak parameters (represented generically by $\hat{g}$ in Fig. 1) and external wave functions, $(\hat{V}_1, \hat{V}_n)$, or else they are suppressed by inverse powers of the large masses. Therefore, in the asymptotic limit of an infinitely heavy SUSY spectrum, these effects decouple in the physical observables involving external electroweak gauge bosons. This decoupling is referred to as decoupling a la Appelquist Carazzone [29].
Figure 1: Decoupling of heavy sparticles and heavy Higgses in the n-point functions of the electroweak gauge bosons $V = A, Z, W^\pm$.

5 SUSY-QCD corrections to $h^o \rightarrow \bar{b}b$ in the decoupling limit

In this section we study the SUSY-QCD corrections to the partial decay width $\Gamma(h^o \rightarrow \bar{b}b)$ at the one-loop level and to leading order in perturbative QCD, that is $\mathcal{O}(\alpha_S)$. We will then explore the decoupling behaviour of these corrections for large SUSY masses, $M_{\text{SUSY}}$, and/or large $M_A$. Both numerical and analytical results will be presented [19].

For the $h^0$ mass range predicted by the MSSM, the decay channel $h^o \rightarrow \bar{b}b$ is by far the dominant one and the precise value of its branching ratio will be crucial for the $h^0$ final experimental reach at the Tevatron.

Among the various contributions to this decay width, the QCD corrections are known to be the dominant ones. At the one-loop level and to order $\alpha_S$ these can be written as,

$$\Gamma_1(h^o \rightarrow \bar{b}b) \equiv \Gamma_0(h^o \rightarrow \bar{b}b)(1 + 2\Delta_{QCD} + 2\Delta_{SQCD}),$$

where, $\Gamma_0(h^o \rightarrow \bar{b}b)$ is the tree level width, $\Delta_{QCD}$ is the one-loop contribution from standard QCD and $\Delta_{SQCD}$ is the one-loop contribution from the SUSY-QCD sector of the MSSM. The factor 2 is just a convention. The QCD correction, $\Delta_{QCD}$, gives a $\sim 50\%$ reduction in the $\Gamma(h^o \rightarrow \bar{b}b)$ decay rate for $M_{h^0}$ in its MSSM range [30]. This correction has the same form in the MSSM as in the SM, so that it gives no information in distinguishing the MSSM from the SM. The SQCD correction, $\Delta_{SQCD}$, was first computed in the on-shell scheme by using a diagrammatic approach in ref. [7] and later studied in more detail in [8]. The SQCD corrections to the $h^0\bar{b}b$ coupling were also computed in an effective Lagrangian approach in ref. [14], using the SUSY contributions to the $b$-quark self energy [31] and neglecting terms suppressed by inverse powers of SUSY masses. Radiatively induced Higgs-fermion-fermion couplings in supersymmetric theories were also studied in ref [15]. The size of the SQCD correction, $\Delta_{SQCD}$, and the QCD correction, $\Delta_{QCD}$, are comparable for a wide window
of the MSSM parameter space. In some regions of the MSSM parameter space, the SQCD corrections become so large that it is important to take into account higher-order corrections. The two-loop SQCD corrections have been studied in a diagrammatic approach in ref. \[33\]. A higher-order analysis has also been carried out in refs. \[17,18\] by performing a resummation of the leading tan $\beta$ contributions to all orders of perturbation theory and by using an effective Lagrangian approach. However, this resummation is not important in our present work because we are interested in the decoupling limit, in which the one-loop corrections to the $h^0 b\bar{b}$ coupling are small enough. Thus, for the present analysis we will just keep the one-loop corrections.

Figure 2: One-loop SUSY diagrams contributing to $\mathcal{O}(\alpha_S)$ to $h^0 \to b\bar{b}$ decay

To one-loop and $\mathcal{O}(\alpha_S)$ there are two type of diagrams, shown in Fig. 2, that contribute to

$$\Delta_{SQCD} = \Delta_{SQCD}^{\text{loops}} + \Delta_{SQCD}^{\text{CT}}.$$  

The triangle diagram, with exchange of sbottoms and gluinos, contributes to $\Delta_{SQCD}^{\text{loops}}$, whereas the bottom self-energy diagram contributes to the counter-terms part $\Delta_{SQCD}^{\text{CT}}$. The exact results in the on-shell scheme are summarized by,

$$\Delta_{SQCD}^{\text{loops}} = \frac{\alpha_S}{3\pi} \left\{ \frac{2M_Z^2}{m_b} \frac{m_b \sin(\alpha + \beta)}{\sin \alpha} (I_3^b \cos^2 \theta_b - Q_b s_W^2 \cos 2\theta_b) + 2m_b + Y_b \sin 2\theta_b \right\} \times \left[ m_b C_{11} + M_{\tilde{g}} \sin 2\theta_b C_0 \right] \left( m_{b_1}^2, M_{\tilde{b}_1}, m_{\tilde{g}}, M_{\tilde{b}_1}, M_{\tilde{b}_1} \right)$$

$$+ \left[ \frac{2M_Z^2}{m_b} \frac{m_b \sin(\alpha + \beta)}{\sin \alpha} (I_3^b \sin^2 \theta_b + Q_b s_W \cos 2\theta_b) + 2m_b - Y_b \sin 2\theta_b \right] \times \left[ m_b C_{11} - M_{\tilde{g}} \sin 2\theta_b C_0 \right] \left( m_{b_2}^2, M_{\tilde{b}_2}, m_{\tilde{g}}, M_{\tilde{b}_2}, M_{\tilde{b}_2} \right)$$

$$+ \left[ \frac{M_Z^2}{m_b} \frac{m_b \sin(\alpha + \beta)}{\sin \alpha} (I_3^b - 2Q_b s_W \sin 2\theta_b + Y_b \cos 2\theta_b) \right] \times \left[ 2M_{\tilde{g}} \cos 2\theta_b C_0 (m_{b_1}^2, M_{\tilde{g}}^2, m_{\tilde{g}}, M_{\tilde{b}_1}^2, M_{\tilde{b}_2}^2) \right] \right\} ,$$

$$\Delta_{SQCD}^{\text{CT}} = -\frac{\alpha_S}{3\pi} \left\{ \frac{m_b^2 \sin 2\theta_b B_0 (m_{\tilde{b}_1}^2, M_{\tilde{g}}^2, M_{\tilde{b}_1}^2) - B_0 (m_{\tilde{b}_2}^2, M_{\tilde{g}}^2, M_{\tilde{b}_2}^2)}{m_b} \right\}$$

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\[-2m_b^2 \left[ B'_1(m_b^2; M_g^2, M_b^2) + B'_1(m_b^2; M_g^2, M_{b_1}^2) \right] \]
\[-2m_b M_g \sin 2\theta_b \left[ B'_0(m_b^2; M_g^2, M_b^2) - B'_0(m_b^2; M_g^2, M_{b_1}^2) \right] \}

where we have used the standard notation for masses, couplings and mixing angles, and we have followed the definitions and conventions for the one-loop integrals \( B_0, B'_0, B'_1, C_0 \) and \( C_{11} \) of ref. [19]. Notice that

\[ Y_b \equiv A_b + \mu \cot \alpha \]

appears in the \( h^0 R \bar{b}_b \) coupling and, therefore, it will be responsible for sizeable contributions in the large \( A_b \) and/or \( \mu \) limit. Our results agree with those of refs. [7,8].

In order to compute \( \Delta_{SQCD} \) in the decoupling limit of very heavy sbottoms and gluinos, we have considered the previously announced assumption for the MSSM parameters, i.e, the 'large SUSY mass limit'. We have performed a systematic expansion of the one-loop integrals and the mixing angle \( \theta_b \) in inverse powers of the large SUSY mass parameters. The resulting formulas of these expansions can be found in ref. [19]. Thus, by defining

\[ \tilde{M}_S^2 \equiv \frac{1}{2} (M_{b_1}^2 + M_{b_2}^2), \quad R \equiv \frac{M_{\tilde{g}}}{M_S}, \quad X_b \equiv A_b - \mu \tan \beta \]

and including terms up to \( \mathcal{O}(M_{Z,h^0}^2/\tilde{M}_S^2) \) in the expansion, we get the following result for the maximal mixing case, \( \theta_b^2 \sim \pm 45^\circ \):

\[ \Delta_{SQCD} = \frac{3}{8\pi} \left\{ -\frac{\mu M_b}{M_S^2} (\tan \beta + \cot \alpha) f_1(R) - \frac{Y_b M_b M_S^2}{12 M_S^4} f_4(R) \right. \]
\[ + \frac{2 M_S^2}{3 \tilde{M}_S^2} \frac{\cos \beta \sin(\alpha + \beta)}{\sin \alpha} I_3^b \left( f_5(R) + \frac{M_b X_b}{M_S^2} f_2(R) \right) + \mathcal{O} \left( \frac{m_{h^0}^2}{\tilde{M}_S^2} \right) \]

where \( I_3^b = -1/2 \) and the functions \( f_i(R) \) are defined in ref. [19] and are normalized as \( f_i(1) = 1 \).

Notice that the first term is the dominant one in the limit of large \( M_{SU3Y} \) mass parameters and does not vanish in the asymptotic limit of infinitely large \( M_S, M_{\tilde{g}} \) and \( \mu \). The second and third terms are respectively of \( \mathcal{O}(M_{h^0}^2/M_{SU3Y}^2) \) and \( \mathcal{O}(M_{Z}^2/M_{SU3Y}^2) \) and vanish in the previous asymptotic limit. Therefore the first term gives the non-decoupling SUSY contribution to the \( \Gamma(h^0 \to b\bar{b}) \) partial width which can be of phenomenological interest. Moreover, since this term is enhanced at large \( \tan \beta \) it can provide important corrections to the branching ratio \( BR(h^0 \to b\bar{b}) \), even for a very heavy SUSY spectrum. The sign of these corrections depends crucially on the sign of \( \mu M_{\tilde{g}} \). The previous result, when expressed in terms of the \( h^0 \) effective coupling to \( b\bar{b} \), agrees with the result in ref. [14] based on the zero external momentum approximation or, equivalently, the effective Lagrangian approach.
From our previous result, we conclude that there is no decoupling of sbottoms and gluinos in the limit of large SUSY mass parameters for fixed $M_A$. How do we then recover decoupling of the heavy MSSM spectra from the SM low energy physics? The answer to this question relies in the fact that in order to converge to SM predictions we need to consider not just a heavy SUSY spectra but also a heavy Higgs sector. That is, besides large $M_{SUSY}$, the condition of large $M_A$ is also needed. Thus, if $M_A \gg M_Z$ the light Higgs $h^0$ behaves as the SM Higgs boson, and the extra heavy Higgses $A$, $H^\pm$ and $H^0$ decouple. The decoupling of SUSY particles and the extra Higgs bosons in $\Delta_{SQCD}$ is seen explicitly once the large $M_A$ limit of the mixing angle $\alpha$ is considered,

$$\cot \alpha = -\tan \beta - \frac{2M_Z^2}{M_A^2} \tan \beta \cos 2\beta + O\left(\frac{M_A^4}{M_Z^4}\right).$$

By substituting this into our previous result we see that the non-decoupling terms cancel out and we get finally,

$$\Delta_{SQCD} = \frac{\alpha}{3\pi} \left\{ \frac{2\mu M^2_3}{M_Z^2} f_1(R) \tan \beta \cos 2\beta \frac{M_Z^2}{M_A^2} - X_b^b \frac{M_Z^2}{M_S^2} f_4(R) \right\}$$

$$+ \frac{2M_Z^2}{3M_S^2} \cos 2\beta I_b^b \left( f_5(R) + \frac{M_Z^2}{M_S^2} f_2(R) \right) + O\left(\frac{m_b^2}{M_Z^2}\right)$$

which clearly vanishes in the asymptotic limit of $M_{SUSY}$ and $M_A \to \infty$.

In conclusion, we get decoupling of the SQCD sector in $h^0 \to b\bar{b}$ decays, if and only if, both $M_{SUSY}$ and $M_A$ are large. In this limit, the dominant terms go as,

$$\Delta_{SQCD} \sim C_1 \frac{M_Z^2}{M_A^2} + C_2 \frac{M_Z^2 h^0}{M_{SUSY}^2},$$

and, since both $C_1$ and $C_2$ are enhanced by $\tan \beta$, we expect this decoupling to be delayed for large $\tan \beta$ values. All these results are similar for the near zero mixing case, $\theta_b^b \sim 0^\circ$; for brevity we do not show these here (see ref. [19]).

Finally, in order to show this decoupling numerically, we have studied a simple example where there is just one relevant MSSM scale, $M_S$. More specifically, we have chosen,

$$M_S = M_Q = M_D = \mu = A_b = M_g = M_A,$$

which, in the limit $M_S \gg M_Z$, gives maximal mixing, $\theta_b^b \sim 45^\circ$. In Fig. 3 we show the numerical results for the exact one-loop SQCD corrections, as a function of this common MSSM mass scale $M_S$, and for several values of $\tan \beta$. We can see in this

\footnote{It should be noticed that, strictly speaking, the decoupling theorem [29] is not applicable to the MSSM case, since it is a theory that incorporates the SM chiral fermions and the SM electroweak spontaneous symmetry breaking. For a more detailed discussion on this, see ref. [20].}
Figure 3: Exact numerical results for $\Delta_{SQCD}$ in $h^0 \to b\bar{b}$ decay as a function of a common MSSM scale $M_S$ and for several values of $\tan \beta$.

figure clearly the decoupling of $\Delta_{SQCD}$ with $M_S$. This decoupling goes as $1/M_S^2$, in agreement with our analytical result, and is delayed for large $\tan \beta$ values. The typical size of this correction is $|\Delta_{SQCD}| \leq 10\%$ for $M_S \geq 250 \, \text{GeV}$. Notice that the sign of $\Delta_{SQCD}$ here is negative because of our choice of positive $\mu$ and $M_{\tilde{g}}$.

Finally we have studied the different behavior of decoupling of squarks and gluinos in the $\Gamma(h^0 \to b\bar{b})$ decay width. We have proved the independent decoupling of the gluinos and squarks whenever they are considered very heavy as compared to the electroweak scale while keeping a large gap among their masses, that is when $M_{\tilde{g}} \gg M_{\tilde{q}} \gg M_{EW}$ or $M_{\tilde{q}} \gg M_{\tilde{g}} \gg M_{EW}$ respectively. Furthermore, the decoupling of gluinos is much slower than the decoupling of squarks due to the logarithmic dependence on the gluino mass (see ref. [19] for more details).

6 SUSY-QCD corrections to $H^+ \to t\bar{b}$ in the decoupling limit

In this section we study the SUSY-QCD corrections to the partial decay width $\Gamma(H^+ \to t\bar{b})$ at the one-loop level and to $\mathcal{O}(\alpha_S)$. We will then analyze these corrections in the decoupling limit of large SUSY masses. We will present here just a short summary of the main numerical and analytical results, and refer the reader to ref. [20] for a more detailed study.

If all SUSY particles are heavy enough, $H^+$ decays dominantly into $t\bar{b}$ above the $t\bar{b}$ threshold. As in the case of $h^0 \to b\bar{b}$, the dominant radiative corrections to $H^+ \to t\bar{b}$
are the QCD corrections. At the one-loop level and to $\mathcal{O}(\alpha_S)$, the corresponding partial width can be written as,

$$\Gamma_1(H^+ \rightarrow t\bar{b}) \equiv \Gamma_0(H^+ \rightarrow t\bar{b})(1 + 2\Delta_{QCD} + 2\Delta_{SQCD}),$$

where $\Gamma_0(H^+ \rightarrow t\bar{b})$ is the tree-level width, $\Delta_{QCD}$ is the correction from standard QCD, and $\Delta_{SQCD}$ is the correction from SUSY-QCD. The standard QCD corrections were computed in ref. [35] and can be large (+10% to −50%). The SUSY-QCD corrections were computed by using the diagrammatic approach in refs. [9,10] and can be comparable or even larger than the standard QCD corrections in a large region of the SUSY parameter space.

At the one-loop level and to $\mathcal{O}(\alpha_S)$ there are two type of diagrams that contribute to $\Delta_{SQCD} = \Delta_{SQCD}^{\text{loops}} + \Delta_{SQCD}^{\text{CT}},$ as shown in Fig. 4. The triangle diagram, with exchange of sbottoms, stops and gluinos, contributes to $\Delta_{SQCD}^{\text{loops}},$ whereas the bottom and top self-energy diagrams contribute to the counter-terms part $\Delta_{SQCD}^{\text{CT}}.$ The exact results in the on-shell scheme are summarized by,

$$\Delta_{SQCD}^{\text{loops}} = \frac{U_t}{D} H_t + \frac{U_b}{D} H_b,$$

$$\Delta_{SQCD}^{\text{CT}} = \frac{U_t}{D} \left( \frac{\delta m_t}{m_c} + \frac{1}{2} \delta Z_{L}^b + \frac{1}{2} \delta Z_{R}^t \right) + \frac{U_b}{D} \left( \frac{\delta m_b}{m_s} + \frac{1}{2} \delta Z_{L}^t + \frac{1}{2} \delta Z_{R}^b \right),$$

where,

$$D = (M_{H^+}^2 - m_t^2 - m_b^2) (m_t^2 \cot^2 \beta + m_b^2 \tan^2 \beta) - 4m_t^2 m_b^2,$$

$$U_t = (M_{H^+}^2 - m_t^2 - m_b^2) m_t^2 \cot^2 \beta - 2m_t^2 m_b^2,$$

$$U_b = (M_{H^+}^2 - m_t^2 - m_b^2) m_b^2 \tan^2 \beta - 2m_t^2 m_b^2,$$

$$H_t = -\frac{2\alpha_s}{3\pi} \frac{G_{ab}^*}{m_t \cot \beta} [m_t R_{1b}^{(t)} R_{1a}^{(b)*} (C_{11} - C_{12}) + m_b R_{2b}^{(t)} R_{2a}^{(b)*} C_{12}.$$
The GM case it does not make sense to consider the alternative limit of large 
considered again our 'large SUSY mass limit', and we have performed a sys-
tematic result, that is, the dominant term in this expansion for the par-
ticular choice of maximal mixing. Thus, for
\[ \theta_{b,t} \sim 45^\circ, \quad M_S^2 \equiv \frac{1}{2}(M_{b_1}^2 + M_{b_2}^2) \equiv \frac{1}{2}(M_{t_1}^2 + M_{t_2}^2) \quad \text{and} \quad R \equiv M_S/\tilde{M}_S \] we get:

\[ \Delta_{SQCD} = \frac{68}{3\pi} \left\{ -\frac{\tan \beta + \cot \beta}{\tan \beta} f_1(R) + \mathcal{O} \left( \frac{M_{EW}^2}{M_{SUSY}^2} \right) \right\} \]

This leading term does not vanish in the heavy SUSY particle limit and, therefore, 
there is no decoupling of stops, sbottoms and gluinos in the \( \Gamma(H^+ \to t\bar{b}) \) decay width.
to one-loop level. This can be seen clearly, for instance, for the simplest case of equal mass scales, $\mu = M_{\tilde{g}} = \tilde{M}_S$, where $f_1(R) = 1$. This leading term, when expressed in terms of an effective coupling of $H^+$ to $b\bar{t}$ is in agreement with the previous results of refs. [17,18] that were obtained in the zero external momentum approximation by using an effective Lagrangian approach. We see in this result the enhancement of $\Delta_{SQCD}$ by $\tan\beta$, so that this non-decoupling effect can be numerically important for large $\tan\beta$ values. As in the case of $h^0$, the sign of the SQCD correction is determined by the sign of $M_{\tilde{g}}$ and $\mu$. We have obtained similar results for the case of minimal mixing, as can be seen in [20].

Finally, in order to illustrate this non-decoupling behavior numerically, we present in Fig. 5 the $\Delta_{SQCD}$ correction as a function of a common SUSY mass scale $M_S = M_{\tilde{Q}} = M_{\tilde{U}} = M_{\tilde{D}} = M_{\tilde{g}} = A_B = A_t = \mu$. The Higgs mass has been fixed to $M_{H^+} = 250$ GeV, and several values of $\tan\beta$ have been considered. The fact that $\Delta_{SQCD}$ tends to a non-vanishing value for very large $M_S$ shows precisely this non-decoupling effect. The correction is quite sizeable, even for a very heavy SUSY spectrum. This is particularly noticeable for large $\tan\beta$.

![Figure 5: $\Delta_{SQCD}$ in $H^+ \to t\bar{b}$ decay as a function of the common scale $M_{SUSY} = M_S$.](image)

As in the $h^0$ case, we have proved for the $H^+ \to t\bar{b}$ decay the independent decoupling of the gluinos and squarks whenever they are considered very heavy as compared to the electroweak scale and with a large gap among their masses. The decoupling of gluinos is much slower than the decoupling of squarks due again to the logarithmic dependence on the gluino mass. In fact, this very slow decoupling with the gluino mass is the responsible for the large size of $\Delta_{SQCD}$, specially for large $\tan\beta$. For instance, if $\tan\beta = 30$ and $M_{\tilde{g}} = 2$ TeV we get $\Delta_{SQCD} = -40\%$. Notice that the size
can be so large that the validity of the perturbative expansion can be questionable. We refer the reader to refs. [17,18] where this subject is studied and some techniques of resumation for a better convergence of the series are proposed.

7 SUSY-QCD corrections to $t \rightarrow W^+b$ in the decoupling limit

In this section we briefly comment on the SUSY-QCD corrections to $t \rightarrow W^+b$ at the one-loop level and to $O(\alpha_S)$, and we study them in the decoupling limit. These radiative corrections were studied in the context of the MSSM in ref. [36] and are known to be important for some regions of the MSSM parameter space. The standard QCD corrections are also known to be important and give a $\sim -10\%$ reduction in $\Gamma(t \rightarrow W^+b)$ [37]. The Feynman diagrams that contribute to the SQCD corrections are shown in Fig. 6. The size of the SQCD corrections has been estimated to range between $-5\%$ and $-10\%$ and are quite insensitive to $\tan \beta$ [38]. In contrast, the SUSY-Electroweak corrections that range between $-1\%$ and $-10\%$ are known to grow with $\tan \beta$ [38].

In order to analyze the decoupling limit in this observable we have chosen the simplest case with just one SUSY scale, $M_S$, which is considered very large as compared to the electroweak scale, $M_{EW}$,

$$M_Q = M_U = M_D = A_t = A_b = \mu = M_{\tilde{g}} = M_S \gg M_{EW}.$$ 

After performing an expansion of $\Delta_{SQCD}$ (we use here an analogous notation as in previous sections) in inverse powers of $M_S$ we have obtained the following result for the dominant contribution,

$$\Delta_{SQCD} = -\frac{\alpha_s}{3\pi} \frac{m_t^2}{M_S^2} \left( \frac{1}{6} + \frac{1}{24} (1 - \cot \beta)^2 + \frac{1}{6} (1 - \cot \beta) \right) + O\left( \frac{m_t M_W, M_W^2, ...}{M_S^2} \right)$$

From this result, we conclude that there is decoupling as $M_S$ becomes large in the SQCD corrections to the dominant top decay, $t \rightarrow W^+b$, and this decoupling which
behaves as \((m_t^2/M_S^2)\) is not delayed for large \(\tan \beta\) values. Indeed, we see in the previous equation that these corrections are not enhanced by \(\tan \beta\) factors. Thus, we do not expect relevant indirect signals from a heavy SUSY-QCD sector in this decay channel.

8 Indirect sensitivity to a heavy SUSY spectrum

In this section we shortly review the main results of ref. [25] where a set of optimal observables to search for indirect SUSY-QCD signals in Higgs bosons and top decays has been proposed. These observables (for the Higgs bosons case) are defined as ratios of Higgs branching ratios into quarks divided by the corresponding Higgs branching ratios into leptons, and are the most sensitive to the SUSY-QCD radiative corrections. This is a consequence from the fact that the corrections from heavy sbottoms, stops and gluinos do not decouple in the Higgs boson decays into quarks. We have already shown in the previous sections some of these non-decoupling corrections in the \(h^0 \to b\bar{b}\), and \(H^+ \to t\bar{b}\) decays. There are non-decoupling SUSY-QCD corrections in another Higgs bosons and top decays as well; for instance, in \(H^0 \to b\bar{b}\) and \(A^0 \to b\bar{b}\) decays, and in the top quark decay \(t \to H^+\bar{b}\). In addition, all these corrections are enhanced by \(\tan \beta\) factors and, therefore, they can provide sizeable contributions to the corresponding partial decay widths for large enough \(\tan \beta\) values, even for very heavy squarks and gluinos. The results of these SUSY-QCD non-decoupling contributions to all the relevant partial decay widths can be found in ref. [24], for arbitrary \(\tan \beta\) and \(M_A\) values and in the ‘large SUSY mass limit’.

Next, we briefly comment on the proposed set of optimal observables. These are the following [25]:

\[
O_{h^0} \equiv \frac{B(h^0 \to b\bar{b})}{B(h^0 \to \tau^+\tau^-)}, \quad O_{H^0} \equiv \frac{B(H^0 \to b\bar{b})}{B(H^0 \to \tau^+\tau^-)},
\]

\[
O_{A^0} \equiv \frac{B(A^0 \to b\bar{b})}{B(A^0 \to \tau^+\tau^-)}, \quad O_{H^+} \equiv \frac{B(H^+ \to t\bar{b})}{B(H^+ \to \tau^+\nu)},
\]

and

\[
O_t \equiv \frac{B(t \to H^+b)}{B(t \to W^+b)},
\]

which complements the charged Higgs observable in the low \(M_A\) region. These observables are specially sensitive to indirect SUSY-QCD searches because of the following reasons:

- The SUSY-QCD non-decoupling corrections contribute to the numerator but not
to the denominator of these observables. The first will be considered as the search channel. The second, as the control channel;

- These corrections are maximized at large $\tan \beta$ and are sizeable enough as to be measurable;
- The production uncertainties are minimized in ratios;
- They will be experimentally accessible at LHC/Tevatron/Linear Colliders;
- They will allow to distinguish the MSSM Higgs sector from a general 2HDM of type II.

The non-decoupling SUSY-QCD contributions to these observables in the 'large SUSY mass limit' are the following:

\[
O_{h^o} = O_{h^o}^o \left[ 1 - \frac{2\alpha_s}{3\pi} \frac{M_{\tilde{g}}\mu}{M_{\tilde{g}}^2} (\tan \beta + \cot \alpha) \right]
\]

\[
O_{H^o} = O_{H^o}^o \left[ 1 - \frac{2\alpha_s}{3\pi} \frac{M_{\tilde{g}}\mu}{M_{\tilde{g}}^2} (\tan \beta - \tan \alpha) \right]
\]

\[
O_{A^o} = O_{A^o}^o \left[ 1 - \frac{2\alpha_s}{3\pi} \frac{M_{\tilde{g}}\mu}{M_{\tilde{g}}^2} (\tan \beta + \cot \beta) \right]
\]

\[
O_{H^+} = O_{H^+}^o \left[ 1 - \frac{2\alpha_s}{3\pi} \frac{M_{\tilde{g}}\mu}{M_{\tilde{g}}^2} (\tan \beta + \cot \beta) \right]
\]

\[
O_t = O_t^o \left[ 1 - \frac{2\alpha_s}{3\pi} \frac{M_{\tilde{g}}\mu}{M_{\tilde{g}}^2} (\tan \beta + \cot \beta) \right]
\]

where the leading terms, $O^o$, refer to the value of the observables without the SUSY particle contributions.

The numerical results for the SUSY-QCD corrections to the set of observables $O_{h^o}$, $O_{H^o}$, $O_{A^o}$, $O_{H^+}$ and $O_t$, together with a discussion in terms of the relevant parameters, $\tan \beta$, $M_A$, $\mu$ and $M_{\tilde{g}}$, can be found in ref. [25]. Here we concentrate on the results about their sensitivity with $\tan \beta$, and compare them with the theoretical uncertainties that would modify the prediction of the observables, previous to the SUSY-QCD corrections. In particular, the largest theoretical uncertainties coming from the errors in $m_b$, $m_t$, $\alpha_s$ and from neglecting the $O(\alpha_s^2)$ corrections have been included in this analysis.

We show in Figs. 7-9 the results for three $M_A$ values in the large, medium and low $M_A$ region: $M_A = 500$, 250, and 100 GeV respectively. The central lines in these figures follow the predictions for the observables without the SUSY-QCD contribution, namely, $O^o$. The corresponding total theoretical uncertainties (evaluated from the partial uncertainties in quadrature) are shown as shadowed bands around the central values. The bold lines represent the SUSY-QCD corrected predictions for $\mu > 0$ (solid), and $\mu < 0$ (dashed, respectively). The observable $O_t$ appears in the lower left plot of Fig. 8 replacing $O_{H^+}$, since in this case $m_t > M_{H^+} + m_b$. 

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Figure 7: Predictions for $O_{h^0}$, $O_{H^0}$, $O_{H^+}$ and $O_{A^0}$ as a function of $\tan \beta$, for $M_A = 500$ GeV. The central lines are the corresponding predictions for $O^0$. The shadowed bands cover the total theoretical uncertainties. The bold lines represent the SUSY-QCD corrected predictions for $M_{\text{SUSY}} = M_{\tilde{g}} = |\mu|$. 
One can see from the figures that, first, the predictions for all the observables separate from the central values as $\tan \beta$ grows. The sign of the SUSY-QCD corrections is positive for $\mu < 0$, and negative for $\mu > 0$. The central values and their theoretical uncertainties are rather insensitive to $\tan \beta$ in the $H^0$ and $A^0$ channels. For the $h^0$, there is a very slight dependence at low $\tan \beta$, while for $H^+$ this dependence is very strong in the low $\tan \beta$ region and it softens for $\tan \beta > 15$. In any case, the predictions for all the Higgs observables without the SUSY-QCD contributions are $\tan \beta$ insensitive in the large $\tan \beta$ region. This is very different for the top observable: it depends strongly on $\tan \beta$ in all the parameter space.

Concerning the size of the SUSY-QCD corrections, we see that there is always a $\tan \beta$ region where these are larger than the theoretical uncertainties. For $M_A = 500$ GeV and $M_A = 250$ GeV, the observables $O_{A^0}$ and $O_{H^0}$ behave similarly and the
predictions surpass the shadowed band for \( \tan \beta > 5 \). For \( M_A = 100 \) GeV, the crossing still happens at \( \tan \beta > 5 \) for \( A^0 \), whereas for \( H^0 \) the bold lines lie outside the error band in the whole region \( 2 < \tan \beta < 50 \).

In the \( h^0 \) case the situation is qualitatively different. For \( M_A = 500 \) GeV, the SUSY-QCD correction is below the theoretical uncertainty for all values in the region \( 2 < \tan \beta < 50 \). This is a manifestation of the decoupling of this correction for large \( M_A \) values. For smaller values of \( M_A \), as \( M_A = 250, 100 \) GeV, the decoupling has not effectively operated yet and the SUSY-QCD corrections are sizeable for large enough \( \tan \beta \). In particular, for \( M_A = 250 \) (100) GeV, they are larger than the theoretical error band for \( \tan \beta > 15 \) (5 respectively).

Regarding the charged Higgs case, two different situations must be considered. For \( M_A = 100 \) GeV, where the decay into a top and a bottom is not kinematically allowed,
the observable $O_t$ is considered. Otherwise, for $M_A = 250$ GeV and $M_A = 500$ GeV, the relevant observable is $O_{H+}$. As can be seen in Fig. 9, the SUSY-QCD corrections in $O_t$ for $M_A = 100$ GeV are above the theoretical uncertainty for $\tan \beta > 15$. However, the central value prediction also depends strongly on $\tan \beta$ (and $M_A$) and it will be difficult to identify the effect of the SUSY-QCD corrections above the additional uncertainty related to the experimental errors on the measurement of $\tan \beta$ and $M_A$. For $O_{H+}$ and $M_A = 500, 250$ GeV both requirements, the correction being larger than the theoretical error and the leading contribution being insensitive to $\tan \beta$ are fulfilled for values larger than about 15.

In summary, we have seen that these observables have a high sensitivity to indirect SUSY-QCD signals.

9 Conclusions

In this review we have analyzed the one-loop SQCD corrections to the partial widths of $h^0 \to b\bar{b}$, $H^+ \to t\bar{b}$ and $t \to W^+ b$ decays, in the limit of large SUSY masses. In order to understand analytically the behavior of the SQCD corrections in this limit, we have performed expansions of the one-loop partial widths that are valid for large values of the SUSY mass parameters compared to the electroweak scale. We have shown that for the SUSY mass parameters and $M_A$ large and all of the same order, the SQCD corrections in $h^0 \to b\bar{b}$ decay decouple like the inverse square of these mass parameters, and the one-loop partial width $\Gamma(h^0 \to b\bar{b})$ tends to its SM value. In this case the effective low energy theory that one obtains after integrating out all the heavy non-standard modes of the MSSM is precisely the SM. However, if the mass parameters are not all of the same size, then this behavior is modified. In particular, if $M_A$ is of the order of the electroweak scale, then the SQCD corrections to the $\Gamma(h^0 \to b\bar{b})$ decay width do not decouple in the limit of large SUSY mass parameters. We have also presented and discussed here a similar non-decoupling SQCD correction to the $\Gamma(H^+ \to t\bar{b})$ decay width. We have also discussed similar SUSY non-decoupling effects that appear in other decay channels such as $H^0 \to b\bar{b}$, $A^0 \to b\bar{b}$ and $t \to H^+ b$.

More generally, it has been shown in ref. [24] that, in the limit of large SUSY mass parameters and $M_A$ of the order of the electroweak scale, the effective Higgs-quark-quark interactions that are generated from the explicit integration in the path integral of heavy squarks and heavy gluinos, at the one-loop level, are those of a more general 2HDM of type III [39], where both Higgs doublets couple to both top and bottom-type quarks. The particular values of the corresponding effective Yukawa couplings have also been computed [24]. These non-decoupling SQCD corrections are of phenomenological interest at present and future colliders. In particular we have shown a recently proposed set of optimal observables that can provide some clues in
the indirect search of a heavy SUSY spectrum at the next generation colliders [25].

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References

[1] H.E. Haber, G.L. Kane, Phys. Rep. 117 (1985) 75.

[2] J. F. Gunion, H. E. Haber, Nucl. Phys. B272, 1 (1986); B278, 449 (1986) [E: B402, 567 (1993)]; J. F. Gunion, H. E. Haber, G. Kane, S. Dawson, The Higgs Hunter’s Guide (Addison-Wesley, Reading, MA, 1990) [E: hep-ph/9302272].

[3] D. E. Groom et al. Eur. Phys. J. C15, 1 (2000).

[4] G.G. Hanson, Searches for new particles, talk presented at the XX International Symposium on Lepton and Photon Interactions at High energies, July 2001, Rome. Slides available at http://www.lp01.infn.it. For an updated version of the PDG see also http://pdg.lbl.gov/2001/sxxx.html.

[5] For an updated compilation of works, see Proceedings for 5th International Symposium on Radiative Corrections (RADCOR-2000), Carmel, CA, USA. September 11-15, 2000. Slides available at http://radcor2000.slac.stanford.edu/program.html.
[6] H. E. Haber and R. Hempfling, Phys. Rev. Lett. 66, 1815 (1991); Phys. Rev. D48, 4280 (1993). Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. 85, 1 (1991); Phys. Lett. B262, 54 (1991); J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B257, 83 (1991); Phys. Lett. B262, 477 (1991); R. Barbieri and M. Frigeni, Phys. Lett. B258, 167 (1991); Phys. Lett. B258, 395 (1991). For an updated study see: M. Carena, H. E. Haber, S. Heinemeyer, W. Hollik, C. E. M. Wagner and G. Weiglein, Nucl. Phys. B580, 29 (2000); and J. R. Espinosa and R.-J. Zhang, JHEP 0003, 026 (2000); Nucl. Phys. B586, 3 (2000), and references therein.

[7] A. Dabelstein, Nucl. Phys. B456 (1995) 25, hep-ph/9503443.

[8] J. A. Coarasa, R. A. Jiménez and J. Solà, Phys. Lett. B389, 312 (1996).

[9] R. A. Jiménez and J. Solà, Phys. Lett. B 389, 53-61 (1996) hep-ph/9511292.

[10] A. Bartl, H. Eberl, K. Hidaka, T. Kon, W. Majerotto and Y. Yamada, Phys. Lett. B 378, 167 (1996) hep-ph/9511385.

[11] J. A. Coarasa, D. Garcia, J. Guasch, R. A. Jiménez and J. Solà, Phys. Lett. B 425 (1998) 329.

[12] J. Guasch, R. A. Jiménez and J. Solà, Phys. Lett. B 360 (1995) 47.

[13] J. A. Coarasa, D. Garcia, J. Guasch, R. A. Jiménez and J. Solà, Eur. Phys. J C2 (1998) 373.

[14] M. Carena, S. Mrenna and C. E. M. Wagner, Phys. Rev. D60, 075010 (1999); Phys. Rev. D62, 055008 (2000).

[15] F. Borzumati, G. R. Farrar, N. Polonsky and S. Thomas, Nucl. Phys. B555, 53 (1999).

[16] K.S. Babu and C. Kolda, Phys. Lett. B451, 77 (1999).

[17] M. Carena, D. Garcia, U. Nierste and C. E. Wagner, Nucl. Phys. B 577 (2000) 88 hep-ph/9912516.

[18] H. Eberl, K. Hidaka, S. Kraml, W. Majerotto and Y. Yamada, Phys. Rev. D 62 (2000) 055006 hep-ph/9912463.

[19] H. E. Haber, M. J. Herrero, H. E. Logan, S. Peñaranda, S. Rigolin and D. Temes, Phys. Rev. D 63 (2001) 055004 hep-ph/0007006.
[20] M. J. Herrero, S. Peñaranda, D. Temes, *SUSY-QCD decoupling properties in H+ → t̅b decay* 2001, [hep-ph/0105097](http://arxiv.org/abs/hep-ph/0105097), FTUAM 00/20, IFT-UAM/CSIC 00/42, KA-TP-10-2001. To appear in Phys. Rev. D (2001).

[21] H. E. Haber, M. J. Herrero, H. E. Logan, S. Peñaranda, S. Rigolin and D. Temes, [hep-ph/0102169](http://arxiv.org/abs/hep-ph/0102169), “Decoupling properties of MSSM particles in Higgs and top decays”. Invited talk given by M. J. H. at the RADCOR-2000 symposium, Carmel CA, USA, 11-15 September, 2000. Slides available at: [http://radcor2000.slac.stanford.edu/program.html](http://radcor2000.slac.stanford.edu/program.html).

[22] J. Guasch, W. Hollik, S. Peñaranda, *Distinguishing Higgs models in H → b̅b/H → τ+τ−* 2001, [hep-ph/0106027](http://arxiv.org/abs/hep-ph/0106027).

[23] M. Carena, H. E. Haber, H. E. Logan, S. Mrenna, *Distinguishing a MSSM Higgs Boson from the SM Higgs Boson at a Linear Collider* 2001, [hep-ph/0106116](http://arxiv.org/abs/hep-ph/0106116).

[24] A. Dobado, M. J. Herrero, D. Temes, *Effective Higgs-quark-quark couplings from a heavy SUSY spectrum* 2001, [hep-ph/0107147](http://arxiv.org/abs/hep-ph/0107147), FTUAM 01/12, IFT-UAM/CSIC 01/18.

[25] A. M. Curiel, M. J. Herrero, D. Temes and J. F. de Troconiz, *Optimal observables to search for indirect SUSY-QCD signals in Higgs bosons decays*, [hep-ph/0106267](http://arxiv.org/abs/hep-ph/0106267), FTUAM 01/13, IFT-UAM/CSIC 01/19.

[26] A. Dobado, M. J. Herrero and S. Peñaranda, Eur. Phys. J. C 7 (1999) 313 [hep-ph/9710313](http://arxiv.org/abs/hep-ph/9710313); In Barcelona 1997, *Quantum effects in the minimal supersymmetric standard model* 266-286, World Scientific, ed. J. Solà [hep-ph/9711441](http://arxiv.org/abs/hep-ph/9711441); Eur. Phys. J. C 12 (2000) 673 [hep-ph/9903211](http://arxiv.org/abs/hep-ph/9903211); Eur. Phys. J. C 17 (2000) 487 [hep-ph/0002134](http://arxiv.org/abs/hep-ph/0002134).

[27] W. Hollik and S. Peñaranda, *Yukawa coupling quantum corrections to the self couplings of the lightest MSSM Higgs boson*, [hep-ph/0108245](http://arxiv.org/abs/hep-ph/0108245), KA-TP-20-2001.

[28] H. E. Haber and Y. Nir, Nucl. Phys. B335, 363 (1990); H. E. Haber, in Proceedings of the US–Polish Workshop, Warsaw, Poland, September 21–24, 1994, edited by P. Nath, T. Taylor, and S. Pokorski (World Scientific, Singapore, 1995) pp. 49–63. H. E. Haber, *Nonminimal Higgs sector: The decoupling limit and its phenomenological implications* (1994), [hep-ph/9501320](http://arxiv.org/abs/hep-ph/9501320).

[29] T. Appelquist and J. Carazzone, Phys. Rev. D11, 2856 (1975).

[30] E. Braaten and J. P. Leveille, Phys. Rev. D22, 715 (1980); N. Sakai, Phys. Rev. D22, 2220 (1980); T. Inami and T. Kubota, Nucl. Phys. B179, 171 (1981).
[31] L. J. Hall, R. Rattazzi and U. Sarid, Phys. Rev. D50, 7048 (1994); R. Hempfling, Phys. Rev. D49, 6168 (1994); M. Carena, M. Olechowski, S. Pokorski and C. E. M. Wagner, Nucl. Phys. B426, 269 (1994).

[32] D. M. Pierce, J. A. Bagger, K. Matchev and R.-J. Zhang, Nucl. Phys. B491, 3 (1997).

[33] S. Heinemeyer, W. Hollik and G. Weiglein, Eur. Phys. J. C16, 139 (2000).

[34] W. Hollik, in Precision Tests of the Standard Electroweak Model, edited by P. Langacker (World Scientific, Singapore, 1995), p. 37–116.

[35] A. Mendez and A. Pomarol, Phys. Lett. B252, 461 (1990); C. S. Li and R. Oakes, Phys. Rev. D43, 855 (1991); A. Djouadi and P. Gambino, Phys. Rev. D51, 218 (1995).

[36] A. Dabelstein, W. Hollik, R. A. Jimenez, C. Junger and J. Sola, Nucl. Phys. B454, 75 (1995).

[37] M. Jezabek and J. H. Kühn, Nucl. Phys. B314, 1 (1989); ibid B320, 20 (1989); C. S. Li, R. J. Oakes and T. C. Yuan, Phys. Rev. D43, 3759 (1991); A. Denner and T. Sack, Z. Phys. C46, 653 (1990); Nucl. Phys. B358, 46 (1991).

[38] D. Garcia, W. Hollik, R. A. Jimenez and J. Sola, Nucl. Phys. B427, 53 (1994); J. M. Yang and C. S. Li, Phys. Lett. 320, 117 (1994).

[39] W. S. Hou, Phys. Lett. B296, 179 (1992); D. Chang, W. S. Hou and W. Y. Keung, Phys. Rev. D48, 217 (1993); D. Atwood, L. Reina and A. Soni, Phys. Rev. D55, 3156 (1997).
\[ \Delta_{SQCD} \]

\[
\begin{align*}
\tan \beta &= 8 \\
\tan \beta &= 30 \\
\tan \beta &= 50 \\
M_A &= M_{\text{gluino}} = \mu = A_b = 200 \text{ GeV}
\end{align*}
\]

- exact formula
- analytic formula
- large \( M_S \) expansion
\[ M_0 = M_D = M_{\tilde{u}} = A_b = A_1 = \mu = 1 \text{ TeV} \]
\[ M_{H^*} = 250 \text{ GeV} \]
$\Delta_{SQCD}$ vs $M_g$ (GeV) with different values of $\tan\beta$:

- $\tan\beta = 8$
- $\tan\beta = 30$
- $\tan\beta = 40$

$M_A = M_S = 200$ GeV
\( \Delta_{\text{SQCD}} \) vs. \( M_S \) (GeV)

- \( M_A = 500 \text{ GeV} \)
- \( M_A = 300 \text{ GeV} \)
- \( M_A = 200 \text{ GeV} \)

\[ \tan \beta = 8 \]

- **Exact result**
- **Large \( M_S \) expansion**

The plot shows the dependence of \( \Delta_{\text{SQCD}} \) on \( M_S \) for different values of \( M_A \). The curves indicate how the correction \( \Delta_{\text{SQCD}} \) varies with the scale parameter \( M_S \) at fixed \( M_A \) and \( \tan \beta \).
$\Delta_{SQCD}$ vs $M_A$ (GeV) for $\tan\beta = 8$, $M_S = 200$ GeV and $M_S = 500$ GeV. The solid line represents the exact result, while the dashed line represents the large $M_S$ expansion.