Non-decoupling effects of SUSY in the physics of Higgs bosons and their phenomenological implications

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Abstract. We consider a plausible scenario in the Minimal Supersymmetric Standard Model (MSSM) where all the genuine supersymmetric (SUSY) particles are heavier than the electroweak scale. In this situation, indirect searches via their radiative corrections to low energy observables are complementary to direct searches, and they can be crucial if the SUSY masses are at the TeV energy range. We summarize the most relevant heavy SUSY radiative effects in Higgs boson physics and emphasize those that manifest a non-decoupling behaviour. We focus, in particular, on the SUSY-QCD non-decoupling effects in fermionic Higgs decays, flavour changing Higgs decays and Yukawa couplings. Some of their phenomenological implications at future colliders are also studied.

INTRODUCTION

The search for SUSY particles via their indirect signals at colliders can be of great relevance in the future. The indirect SUSY searches via their effects on radiative corrections to low energy observables are complementary to direct searches, and can be crucial if the SUSY particles turn out to be too heavy as to be produced directly. Remember that it was the case in the past regarding the top quark indirect searches at LEP, via radiative corrections, which were an important guidance towards its final discovery at the Tevatron collider.

We consider here a plausible scenario in the MSSM where all the genuine SUSY particles are heavier than the electroweak scale $m_{EW}$. In this case, the most promising indirect SUSY signals come from SUSY non-decoupling effects in low energy observables. We shortly review here some of these most relevant effects and show that they will provide sizeable signals at future colliders even if the SUSY masses are as heavy as $M_{SUSY} \sim \mathcal{O}(TeV)$. We study the consequences for Higgs boson physics where the SUSY radiative corrections are known to be sizeable, and focus on the SUSY-QCD radiative corrections which are the dominant SUSY contributions.

The heavy SUSY scenario is defined by taking the soft SUSY breaking mass parameters of the MSSM very large as compared to the electroweak scale. For shortness, we choose here the simplest case where all these mass parameters are quasi-degenerate, and we refer to them by a generic common SUSY scale named $M_{SUSY}$. Concretely, the heavy SUSY scenario is implemented by, $M_{SUSY} \sim M_{\tilde{Q}} \sim M_{\tilde{U}} \sim M_{\tilde{D}} \sim M_{\tilde{g}} \sim |\mu| \gg m_{EW}$.

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where $M_{\tilde{U}}$, $M_{\tilde{D}}$, and $M_{\tilde{Q}}$ characterize the heavy squark masses, and $M_{\tilde{g}}$ the heavy gluino masses. The large $|\mu|$ values, although not needed to get heavy particles in the SUSY-QCD sector, guarantee (together with large values of the other relevant parameters in the SUSY-Electroweak sector, $M_{1,2}$) that the charginos and neutralinos are also heavy.

In addition, we have also considered the possibility that the extra (non-SM like) Higgs particles get also very heavy. This is called here the heavy Higgs sector scenario and is defined by taking the pseudoscalar mass very large, $m_{A} \gg m_{Z}$. The extra Higgs particles of the MSSM, $H^{0}$, $H^{\pm}$, and $A^{0}$ get very heavy, $m_{H^{0}} \sim m_{H^{\pm}} \sim m_{A} \gg m_{EW}$, while $h^{0}$ remains light with a mass close to its upper bound, $m_{h^{0}} \leq 135$ GeV, and with tree level interactions being as the Standard Model (SM) Higgs boson ones.

The outline of this talk is as follows. First we review the concepts of decoupling/non-decoupling of heavy particles. Next we review the non-decoupling effects of squark-gluino loops in flavour preserving Higgs decays and in flavour changing Higgs decays. Then, we summarize the results for the corresponding induced Yukawa effective couplings after the integration, at the one-loop level, of the squarks and gluinos in the path integral. Finally, we shortly review the main phenomenological consequences for future colliders.

**DECOUPLING/NON-DECOUPLING OF HEAVY PARTICLES**

Commonly, there are two different languages used to define the decoupling/non-decoupling behaviour of heavy particles. One uses the concept of the effective action (or, equivalently, the one-particle irreducible Green functions) that is generated after integration in the path integral of the heavy particles, and the other one studies the effects on the low energy observables that are induced from the heavy particles. To explain this more clearly we include next some illustrative examples.

(a) **Effective action/One-particle irreducible functions:**

Let us consider first the simple example of QED and study the effects of heavy electrons on the physics of photons at low energies. One starts with the classical action for photons and electrons in QED, $S_{QED} = -\frac{1}{4} \int dxF_{\mu\nu}F^{\mu\nu} + \int d\bar{\psi}(\not{D} - M_{e})\psi$, and study the effective action that is got after integration of the electron field at one-loop level, $\Gamma_{eff} = -\frac{1}{4} \int dxF_{\mu\nu}F^{\mu\nu} - \frac{e^{2}}{3(4\pi)^{2}} \Delta \int dxF_{\mu\nu}F^{\mu\nu} - \frac{e^{2}}{15(4\pi)^{2}} \int F_{\mu\nu}\partial^{2}F^{\mu\nu} + \ldots$ where $
abla \equiv \frac{2}{e} + \log 4\pi - \gamma - \log \frac{M_{e}^{2}}{\mu^{2}}$ contains the divergent piece in four dimensions, and an expansion in inverse powers of the heavy electron mass, $M_{e} \gg p$, has been performed. The remaining terms in this expansion are suppressed by higher inverse powers of $M_{e}$ and are not shown. Notice that the second term has the same structure as the kinetic term of photons and, therefore its effects are not physically observable since they can be absorbed by a photon wave function redefinition. The other terms, being proportional to inverse powers of $M_{e}$, vanish in the asymptotically large electron mass limit and, therefore, their effects decouple at low energies. This is referred to as decoupling of electrons a la Appelquist-Carazzone, since it follows the general behaviour described in the Decoupling Theorem by these authors [1].

The second example refers to the decoupling of SUSY particles, for the previously
introduced heavy SUSY scenario, in the low energy self-interactions of the electroweak
gauge bosons, $\gamma, Z, W$, at the one-loop level. By using this effective action language,
it has been shown in [2] that after the integration of all the SUSY particles (squarks,
sleptons, charginos and neutralinos) in the path integral formalism, and by considering
the large $M_{\text{SUSY}}$ limit, their effects are absorbed in redefinitions of the gauge boson
wave functions and weak boson masses or else they are suppressed by inverse powers
of $M_{\text{SUSY}}$. Therefore, the SUSY particles decouple. Similarly, it has been shown [2] that
the effects of the extra Higgs bosons, in the previously introduced heavy Higgs sector
scenario, also decouple in the $\gamma, Z, W$ self-interactions at one-loop level. The low energy
interactions left are then just those of the SM.

The third example is the decoupling of the extra Higgs particles in this same heavy
Higgs sector scenario, in the self-interactions of the lightest MSSM Higgs particle, $h_0$, at
the one-loop level. It has been shown in [3] that all the effects of the heavy Higgs bosons
are absorbed by a mass correction, $\Delta m_{h_0}$, or else they are suppressed by inverse powers
of the heavy mass $m_A$. Similarly, it has been shown [4] that the effects of heavy stops
also decouple in the $h_0$ self-interactions, at the one-loop level. These low energy
self-interactions, therefore, have the same structure as the SM Higgs boson self-interactions,
even at the one-loop level.

(b) Observables:
The first example is the non-decoupling of top quark loops in the partial decay width
$\Gamma(Z \rightarrow \bar{b}b)$, within the SM context. The result to one-loop level and in the large $m_t$
limit can be written as, $\Gamma(Z \rightarrow \bar{b}b) = \Gamma_0 \left(1 + a \alpha_s m_t^2 \right)$, where we have factorized out
the treel level result and $a$ is a numerical factor. We see clearly that the correction grows
quadratically with the top mass, indicating a non-decoupling behaviour.

The second example is the decoupling of squark-gluino loops in the partial decay width
$\Gamma(t \rightarrow W^+ b)$. The result to one-loop level and in the large $M_{\text{SUSY}}$ limit can be
written as, $\Gamma(t \rightarrow W^+ b) = \Gamma_0 \left(1 + b \alpha_s \frac{m_t^2}{M_{\text{SUSY}}} \right)$, where we have factorized again the
treel level result and $b$ is a numerical factor. For asymptotically large $M_{\text{SUSY}}$ values
the correction vanishes, indicating a decoupling behaviour.

In the following we will present some other interesting examples, within the MSSM,
where the SUSY particles of the SUSY-QCD sector do not decouple. These are the
fermionic decay widths of the Higgs bosons, both flavour preserving and flavour violat-
ing, and the effective low energy Yukawa interactions.

**NON-DECOUPLING OF SUSY LOOPS IN HIGGS DECAYS**

We choose the Higgs decays into quarks, $H \rightarrow q\bar{q}$, because, on one hand, these are
the dominant decays in a large region of the MSSM parameter space and, on the
other hand, the associated Yukawa couplings, $\lambda_{Hqq}$, enter in relevant Higgs production
processes. We study in particular $h^0 \rightarrow \bar{b}b$ and $H^+ \rightarrow t\bar{b}$, because $h^0$ will be probably
the first to be detected, and because the charged Higgs provides itself an unambiguous
signal of physics beyond SM. As we have said, we focus here on the SUSY radiative
corrections from loops of squarks and gluinos. These being $\delta(\alpha_s)$ are the dominant
SUSY corrections.

There are two types of one-loop diagrams contributing to $h^0 \to bb^*$, the triangular vertex diagrams with two sbottoms and one gluino in the internal propagators of the triangle, and the self-energy diagrams for the external bottom legs with one sbottom and one gluino in the internal propagators. These SUSY-QCD (SQCD) corrections and the genuine QCD corrections modify the tree level result $\Gamma_0$ by, $\Gamma(h^0 \to bb^*)(1 + 2\Delta_{QCD} + 2\Delta_{SQCD})$, and are known to be quite sizeable, $2\Delta_{QCD} \simeq -50\%$, $|2\Delta_{SQCD}| \leq 50\%$. The result of $\Delta_{SQCD}$ is [5],

$$\Delta_{SQCD} = \frac{\alpha_s}{\pi} g_{h^0b\bar{b}} \frac{\mu_y \cos \beta}{\sin \alpha} \left[ m_b (R_{2b}^{(b)\ast} R_{2a}^{(b)}) + R_{1a}^{(b)\ast} R_{1b}^{(b)\ast} \right] C_{11} + M_q (R_{2b}^{(b)\ast} R_{1a}^{(b)}) +$$

$$+ R_{1b}^{(b)\ast} R_{2a}^{(b)}) C_0 \left( m_b^2, m_{b}\mu_y^2, m_b^2, M_b^2, M_b^2, M_b^2 \right) + \Sigma^b_{S}(m_b^2) - 2m_b^2 \left[ \Sigma^b_{t}(m_b^2) + \Sigma^b_{t}(m_b^2) \right]$$

Here $R_{ij}^{(b)}$ are the rotation matrices that relate the interaction eigenstates and the sbottom mass eigenstates. Notice that the trilinear soft breaking parameter $A_t$ and $\mu$ enter in the $g_{h^0b\bar{b}}$ couplings (not shown explicitly here, for brevity) and in $L - R$ squark mixing. $M_q$ enters both in external factors and in the one-loop integrals for the vertex corrections, $C_0, C_{11}$, and for the self-energies $B_0, B_1, M_q$ enters in the integrals. The relevant $\tan \beta$ parameter enters in $L - R$ squark mixing. Similar results have been obtained for $H^+ \to t\bar{b}$ from loops of stops, sbottoms, and gluinos [6].

In order to study the consequences of the heavy SUSY scenario on the previous Higgs boson decays, we perform expansions of the integrals and mixing angles appearing in $\Delta_{SQCD}$, which are valid for $M_{SUSY} \gg m_{EW}$, and get the two first terms being $O\left( \frac{m_{EW}^2}{M_{SUSY}^2} \right)^n$ with $n = 0, 1$ respectively [7]. We give here the result for the leading term, $n = 0$, that is valid for all $\tan \beta$ values,

$$\Delta_{SQCD} = \frac{\alpha_s}{\pi} \left\{ \frac{-\mu_y M_q}{M_{SUSY}^2} \left( \tan \beta + \cot \alpha \right) f_1(R) + O\left( \frac{m_{EW}^2}{M_{SUSY}^2} \right) \right\} \text{ where, } M_{SUSY}^2 \equiv \frac{1}{2}(M_{b_1}^2 + M_{b_2}^2), R \equiv M_q/M_{SUSY}, f_1(1) = 1.$$ We see clearly that for large $M_q \sim M_{SUSY} \sim \mu \sim M_{SUSY}$, the correction gives a non-vanishing constant and therefore the heavy squarks and gluinos do not decouple. We also see that this non-decoupling contribution is enhanced at large $\tan \beta$, which is in agreement with the result in the zero external momentum approximation and large $\tan \beta$ limit of [8]. Finally, if we also consider the heavy Higgs sector scenario, $m_A \gg m_Z$, where $\cot \alpha = -\tan \beta - 2 m_Z^2 \tan \beta \cos 2\beta / m_A^2 + \ldots$, we see that the leading term cancels and the correction decouples as $\sim m_Z^2 / m_A^2$, recovering the SM result, as expected.

For the case of $H^+ \to t\bar{b}$, we have found similar results [9]. In particular the result for the leading term, $n = 0$, valid for all $\tan \beta$ values, reads,

$$\Delta_{SQCD} = \frac{\alpha_s}{\pi} \left\{ \frac{-\mu_y M_q}{M_{SUSY}^2} \left( \tan \beta + \cot \beta \right) f_1(R) + O\left( \frac{m_{EW}^2}{M_{SUSY}^2} \right) \right\} \text{ where, } M_{SUSY}^2 \equiv \frac{1}{2}(M_{q_1}^2 + M_{q_2}^2), \bar{q} = i, \bar{b}, R \equiv M_q/M_{SUSY}, f_1(1) = 1,$$ and it agrees again with the result of the effective approach, [10], at large $\tan \beta$. We find again the same non-decoupling behaviour as in $h^0 \to bb^*$ but the correction is numerically larger. For instance, $\Delta_{SQCD} \simeq -40\%$ for $\tan \beta = 30, M_{SUSY} = 1$ TeV. In addition, we find that the next to leading order terms, with $n = 1$, are smaller than $1\%$ for $M_{SUSY} \geq 300$ GeV.
FLAVOUR CHANGING NEUTRAL HIGGS DECAYS FROM SQUARK-GLUINO LOOPS

Flavour Changing Neutral Current (FCNC) processes are ideal to look for indirect SUSY signals, or any other radiative effects from possible physics beyond the SM, since the SM predicts negligible rates for these processes. These FC interactions are absent at the tree-level in the SM and in the MSSM, but they can be generated at the one-loop level and lead to sizable contributions in the MSSM, specially at large tan β.

We study here the Flavour Changing Neutral Higgs Boson Decays (FCHD) generated from squark-gluino loops in the MSSM \([11]\). They are proportional to \(\partial^i (\alpha_S)\) and therefore dominate the SUSY corrections. We focus on the particular decays, \(h_0, H_0, A_0 \to b \bar{s}, s \bar{b}\) and \(H_0, A_0 \to t \bar{c}, c \bar{t}\) in the non-minimal flavour scenario where there is squark mixing from misalignment between quark and squark mass matrices, which constitutes the most general case in the MSSM. Our study is devoted to the second and third generation quarks because the squark mixing between these two generations is the less constrained experimentally.

Once the quark mass matrices are diagonalized, and assuming that FC squark mixing is significant only in LL entries, the squark mass matrices for the up and down sector can be written as follows,

\[
M_{\tilde{u}}^2 = \begin{pmatrix}
M_{L,c}^2 & m_{c}X_c & \Delta_{LL}^d & 0 \\
m_{c}X_c & M_{R,c}^2 & 0 & 0 \\
\Delta_{LL}^d & 0 & M_{L,t}^2 & m_{t}X_t \\
0 & 0 & m_{t}X_t & M_{R,t}^2
\end{pmatrix}, \quad M_{\tilde{d}}^2 = \begin{pmatrix}
M_{L,s}^2 & m_{s}X_s & \Delta_{LL}^d & 0 \\
m_{s}X_s & M_{R,s}^2 & 0 & 0 \\
\Delta_{LL}^d & 0 & M_{L,b}^2 & m_{b}X_b \\
0 & 0 & m_{b}X_b & M_{R,b}^2
\end{pmatrix}
\]

where,

\[
M_{L,q}^2 = M_{Q,q}^2 + m_q^2 + \cos2\beta(T_q^2 - Q_q s_W^2)m_Z^2; \quad M_{R,(c,t)}^2 = M_{\tilde{U},(c,t)}^2 + m_{c,t}^2 + \cos2\beta Q_t s_W^2 m_Z^2; \\
M_{R,(s,b)}^2 = M_{\tilde{D},(s,b)}^2 + m_{s,b}^2 + \cos2\beta Q_b s_W^2 m_Z^2; \quad X_{c,t} = m_{c,t} (\lambda c_t - \mu \cot\beta), \\
X_{s,b} = m_{s,b} (\lambda_{s,b} - \mu \tan\beta), \quad \Delta_{LL}^d = \lambda M_{L,c} M_{L,t}, \quad \Delta_{LL}^d = \lambda M_{L,s} M_{L,b}; \quad 0 \leq \lambda \leq 1
\]

The squark mass eigenstates, \(\tilde{q}_\alpha\), and the interaction eigenstates, \(\tilde{q}_\alpha'\), are related by, \(\tilde{q}_\alpha' = \sum R_{q\alpha}^{(q)} \tilde{q}_\beta\), where, \(\tilde{q}_\alpha' = (\tilde{c}_L, \tilde{c}_R, \tilde{t}_L, \tilde{t}_R)\), \(\tilde{q}_\alpha = (\tilde{s}_L, \tilde{s}_R, \tilde{b}_L, \tilde{b}_R)\), \(\tilde{u}_\alpha = (\tilde{u}_1, \tilde{u}_2, \tilde{u}_3, \tilde{u}_4)\) and \(\tilde{d}_\alpha = (\tilde{d}_1, \tilde{d}_2, \tilde{d}_3, \tilde{d}_4)\).

We present next our results for the partial decay widths, containing the one-loop corrections in terms of the form factors \(F_{L,R}^{q\bar{q}}(H_a)\) associated to each decay \(H_a \to q \bar{q}'\), with \(H_a = h_0, H_0, A_0\), and defined by \(iF = -ig\bar{q}_q (p_1) (F_{L,R}^{q\bar{q}}(H)_P L + F_{R}^{q\bar{q}}(H)_P R) v_{q'} (p_2) H (p_3)\). Similarly to the flavour preserving case, there are two types of one-loop diagrams that contribute, the triangular vertex loop diagrams and the FC self-energies of the external quarks. We show here only \(F_{L}^{hs}(H_a)\) as an illustrative example (the rest can be found in \([11]\)).

\[
F_{L}^{hs}(H_a) = -\frac{8 m_{\tilde{b}} m_{H_a}}{i g} \frac{2 \alpha_c}{3 \pi} \left( m_{b} R_{3\alpha}^{d(d)} R_{1\beta}^{(d)*} (C_{11} - C_{12}) + m_{s} R_{4\alpha}^{d(d)} R_{2\beta}^{(d)*} C_{12} + M_{s} R_{3\alpha}^{d(d)} R_{1\beta}^{(d)*} C_{0} \right) (m_{b}^2, m_{b}^2, m_{s}^2, m_{s}^2, m_{s}^2, m_{s}^2, M_{d}^2, M_{d}^2, M_{d}^2) + \frac{G_{\text{F}} s_W}{i g} \frac{m_{\tilde{b}}}{m_{\tilde{b}}^2 - m_{b}^2} \kappa_{L}^{d(d)} \left[ m_{b} (\Sigma_{R}^{bs}(m_{b}^2) + \right.
\]

\[
\left. m_{s} (\Sigma_{R}^{bs}(m_{s}^2) + \right) \]
\[ \Gamma(H_o \to b\bar{s} + s\bar{b}) \text{ as a function of the MSSM parameters. The corresponding fixed values are } \mu = 1500 GeV, M_0 = 600 GeV, M_{\tilde{g}} = 300 GeV, A = 200 GeV, m_A = 250 GeV, \tan \beta = 35, \lambda = 0.5. \]

\[ \text{FIGURE 1. } \]

\[ \text{FIGURE 2. } Br(H_o \to b\bar{s} + s\bar{b}) \text{ as a function of } \lambda \text{ for the selected MSSM parameters of fig. 1.} \]

\[ + \sum_{R_L}^{bs} (m_b^2) + m_{\tau} (\sum_{L}^{bs} (m_b^2) + \sum_{R}^{bs} (m_b^2)) \right) + \frac{G_{H\tilde{b}\tilde{b}}}{g^2} \frac{1}{m_Z^2 - m_b^2} \kappa^L \left[ m_b^2 \sum_{L}^{bs} (m_b^2) + \right. \]

\[ + m_b m_{\tau} (\sum_{R}^{bs} (m_b^2) + \sum_{L}^{bs} (m_b^2)) + m_b^2 \sum_{R}^{bs} (m_b^2) \]}

Notice that the results for the form factors are finite, as expected, since renormalization is not needed in this process. The dependences on the MSSM parameters, \( m_A, \tan \beta, \mu, M_{\tilde{g}}, M_0, A \), are similar to the flavour preserving cases studied before. Here we assume universal squark masses \( M_0 \) and trilinear couplings, \( A \). The dependency on the crucial flavour mixing parameter \( \lambda \) enters in the rotation matrices and in the physical squark masses. We show in figs. 1 and 2 the numerical results for the case \( H_o \to \bar{s}b + b\bar{s} \) as a function of the six previous MSSM parameters and \( \lambda \). We see clearly that the branching ratio can be quite sizable. Similar results are found for \( A_o \) and \( h_o \) and also for the decays into \( t\bar{c} \) and \( t\bar{t} \).

Regarding the large \( M_{\text{SUSY}} \) behaviour of these corrections, we find the following result for the leading \( G(M_{\text{SUSY}}^{-1})^0 \) term of the expansion, valid for all \( \tan \beta \) values,

\[ F^{bs}_{L,R}(\{h_o,H_o,A_o\}) = \frac{\alpha_s}{8\pi} \frac{m_{\tilde{b}}}{m_W} \left[ \sin \alpha \frac{\cos \alpha}{\cos \beta} - i \tan \beta \right] (\tan \beta + \{ \cot \alpha, - \tan \alpha, \cot \beta \}) \frac{\mu M_{\tilde{b}}}{M_0} F(\lambda) \]

where, \( F(\lambda) = \frac{2}{\lambda^2} (\lambda + 1) \ln(\lambda + 1) + (\lambda - 1) \ln(1 - \lambda) - 2\lambda \). Similar formulas are found for the u-sector \[ \text{[1]} \]. Again, there is no decoupling with \( M_{\text{SUSY}} \) and the effect is enhanced at large \( \tan \beta \). Decoupling appears only in the case of the light Higgs if \( m_A \gg m_Z \), since \( \cot \alpha \to - \tan \beta \), and thus \( F^{bs}_{L,R}(h_o) \to 0 \), recovering the SM result.

This SUSY non-decoupling behaviour explains the large FCHD rates found:

\[ \text{BR}(H_0,A_0 \to b\bar{s} + s\bar{b}) \leq 0.2, \text{BR}(h_o \to b\bar{s} + s\bar{b}) \leq 0.01, \text{BR}(H_0,A_0 \to t\bar{c} + c\bar{t}) \leq 5 \times 10^{-5}, \]
for $\lambda \leq 0.6$ that will certainly produce interesting phenomenological SUSY signals at colliders.

**NON-DECOUPLING OF SUSY LOOPS IN EFFECTIVE HIGGS-QUARK-QUARK INTERACTIONS**

Here we present the results for the effective Yukawa couplings generated from squark-gluino loops [12]. We start with the relevant terms of the MSSM classical action, containing the free part and the interaction terms, $S = S_0[H] + S_0[g] + S[H,q] + S_0[\bar{q}] + S[\bar{q}, g, q]$, and compute the one-loop effective action, at order $\alpha_S$, that is generated after integration of squarks and gluinos in the path integral formalism, $\mathcal{L}_{\text{eff}} = \int [d\bar{Q}] [dQ][d\bar{q}] e^{i(S_0[\bar{q}] + S_0[\bar{q}]) + S(H,q)}$. We then perform a large $M_{\text{SUSY}}$ expansion and get the corresponding effective lagrangian, defined by, $\Delta \mathcal{L}_{\text{eff}} = \int dx \Delta \mathcal{L}_{\text{eff}}(H,q)$. After redefining properly the quark wave functions and quark masses we get,

$$
\Delta \mathcal{L}_{\text{eff}}(H,q) = \Delta \mathcal{L}_0 + \Delta \mathcal{L}_1 + \Delta \mathcal{L}_2,
$$

with,

$$
\Delta \mathcal{L}_0 = \frac{g_s^2}{4\pi} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} (q \bar{q}) (H \bar{q} q),
$$

and,

$$
\Delta \mathcal{L}_1 = \frac{g_s^2}{4\pi} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} (q \bar{q}) (H \bar{q} q),
$$

and,

$$
\Delta \mathcal{L}_2 = \frac{g_s^2}{4\pi} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} (q \bar{q}) (H \bar{q} q),
$$

The important result is that new $Hq\bar{q}$ couplings emerge of 2HDMIII type. These are precisely the non-decoupling effects from the squark-gluino loops. On the other hand, these explicit expressions of the effective Yukawa couplings, which are valid for all $\tan \beta$ values, have interesting applications for phenomenology as we will show next.

**HOW TO LOOK FOR INDIRECT SUSY SIGNALS**

Finally, we shortly review the main phenomenological consequences for future colliders. In order to do that, we have chosen a set of observables,

$$
\frac{BR(H^0 \rightarrow b\bar{b})}{BR(H^+ \rightarrow t\bar{t})}, \frac{BR(A^0 \rightarrow b\bar{b})}{BR(A^+ \rightarrow t\bar{t})}, \frac{BR(H^0 \rightarrow t\bar{t})}{BR(H^+ \rightarrow t\bar{t})}, \frac{BR(t \rightarrow H^+ b)}{BR(t \rightarrow W^+ b)}, \frac{BR(t \rightarrow H^0 b)}{BR(t \rightarrow W^+ b)},
$$

that are optimal for various reasons [13]. First, the dominant SUSY-QCD corrections do not decouple in the numerator. The SUSY-EW corrections do not decouple either (in the numerator and the denominator), but are smaller [8,14]. For example, for $\tan \beta = 30$ and universal large $M_{\text{SUSY}}, \Delta_{\text{SEW}} \simeq 8\%$, $\Delta_{\text{SQC}} \simeq -40\%$. There are sizeable SUSY-QCD corrections at large $\tan \beta$ in these observables and the ratios allow to cancel the production uncertainties and to minimize the systematic errors using the leptonic decays as control channels. Another interesting feature of these observables is that they will be experimentally accessible at LHC/TeVatron/Linear Colliders, being the SUSY corrections larger than the expected precision.
We show the sensitivity to SUSY-QCD corrections with \( \tan \beta \) in fig. 3(a). The central line corresponds to the value of the observable without SUSY-QCD corrections, while the other two lines are our predictions for the observables, including the SUSY-QCD correction. The central band corresponds to the theoretical uncertainty in the observable coming from the experimental uncertainty in the SM parameters, which is dominated by \( \Delta m_b \). We can see that we are sensitive to the SUSY-QCD corrections for most of the \( \tan \beta \) parameter space, having higher sensitivity at large \( \tan \beta \) [13, 14]. In fig. 3(b) we plot the predictions for the observables including the SUSY-QCD corrections as a function of \( m_{A^0} \) and \( \tan \beta \). The solid contour lines follow the points in the \((m_{A^0}, \tan \beta)\) plane with constant value of the corresponding observable, denoted here by \( O \). The shaded area represents the region where the corrections are smaller than the mentioned theoretical uncertainty. The long (short) dashed lines join the points where an experimental resolution of 50\% (20\%) is required to achieve a meaningful measurement. The regions above these dashed lines fulfill the required sensitivity to the SUSY-QCD corrections. We see that the SUSY effects are visible in most of the parameter space even for the pessimistic case where only a 50\% resolution would be achieved. As already said before, there is greater sensitivity at large \( \tan \beta \). For the light Higgs, to distinguish the MSSM and SM will only be possible at low \( m_A \).

**CONCLUSIONS**

In summary, we find a non-decoupling behaviour of squark-gluino loops in \( h^0 \to \tilde{b} \tilde{b} \) and \( H^+ \to t \bar{b} \) and also in FCHD, \( h^0, H^0, A^0 \to b \bar{s}, s \bar{b} \) and \( H^0, A^0 \to t \bar{c}, c \bar{t} \). We compute the low energy effective theory by integrating out squarks and gluinos in the path integral and find a 2HDM of type III, giving the explicit expressions to \( O(\alpha_S) \) for the effective...
$Hqq$ Yukawa couplings. Finally, we study the sensitivity to these non-decoupling effects at colliders with a set of optimal observables, concluding that they are promising to look for indirect SUSY signals, even if the SUSY spectrum is very heavy.

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