# Quantum corrections for the MSSM Higgs couplings to SM fermions

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Higgs Yukawa couplings to down-type fermions receive, in the MSSM, supersymmetric quantum corrections that can be of order 1 for large values of $\tan\beta$, provided $|\mu| \sim M_{\mathrm{SUSY}}$. Therefore, a sensitive prediction for observables driven by any of these couplings can only be obtained after an all-order resummation of the large corrections. We perform this necessary step and show, as an example, the effect of the resummation on the computation of the $p\overline{p}$, $pp \to t\overline{b}H^- + X$ cross-section and on the branching ratio $\mathcal{BR}(b \to s\gamma)$ at the next-to-leading order.

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# 1  Introduction

The full experimental confirmation of the Standard Model (SM) still requires the finding of the Higgs boson. The last LEP results, suggesting a light Higgs of about 115 GeV [1], are encouraging, but we will have to wait for the upgraded Tevatron or the LHC to see this result either confirmed or dismissed. In any case, there is room for a extended Higgs sector of various kinds (extra doublets, singlets, even triplets).

The Higgs sector of the Minimal Supersymmetric Standard Model (MSSM), well-known nowadays (see [2], for instance), still deserves further studies. In particular, an interesting topic is the effect of supersymmetric corrections on the Yukawa interaction, because many production and decay channels are mediated by these couplings, some of them at the one-loop level (such as $H \to \gamma\gamma$ or $gg \to H$).

For large $\tan\beta$ values, one expects deviations of order 1 of the Yukawa couplings to down-type fermions from their tree-level values, due to gluino (SUSY-QCD) and, to a lesser extent, higgsino (SUSY-EW) radiative effects. This can be seen, for instance, in the computation of the one-loop correction to the $t \to bH^+$ partial decay rate [3], which grows with $\tan\beta$ as $(\alpha_{\mathrm{s,w}}/4\pi)\tan\beta$. With contributions of order $(\alpha_{\mathrm{s}}/4\pi)^n \tan^n\beta$ arising at higher orders in perturbation theory, the one-loop result can only be meaningful for small $\tan\beta$ values. Let us recall that large $\tan\beta$ scenarios, such as those derived from supersymmetric SO(10) models with unification of the top and bottom Yukawa couplings at high energies [4,5], have become more appealing since LEP searches for a light neutral Higgs boson, $h$, started to exclude the low-$\tan\beta$ region of the MSSM parameter space. The latest analyses rule out the MSSM for $\tan\beta$ in the range $0.52 < \tan\beta < 2.25$, even with maximal stop mixing [6].

In this talk we present the resummation of such corrections into the definition of the bottom Yukawa as a function of the bottom mass, restoring the reliability of the perturbative series for large $\tan\beta$. This "improved" formula for the Yukawa is then used in the evaluation of the $p\overline{p}$, $pp \to t\overline{b}H^- + X$ cross-section and of the branching ratio for $b \to s\gamma$ at the next-to-leading order (NLO), comparing the result with the case in which no resummation is made.

# 2  Resummation of SUSY corrections

Let us briefly explain how such large corrections could arise. The starting point is the MSSM superpotential. Supersymmetry constrains it to be holomorphic in the chiral superfields, implying that the left-handed components of down-type quarks and leptons only couple to the $H_1$ Higgs doublet, while the left-handed up-type quarks and leptons only couple to $H_2$. For the third-generation quarks, one has

$$\mathcal{L} = -h_b \, \overline{b}b \, H_1^0 - h_t \, \overline{t}t \, H_2^0 + \cdots \tag{1}$$

Soft-SUSY-breaking operators induce the forbidden couplings, $\bar{b}bH_2^0$ and $\bar{t}tH_1^0$, radiatively. After integrating out all R-odd particles in the MSSM, one obtains an effective two-Higgs-doublet model (2HDM) lagrangian

$$\mathcal{L}_{\text{eff}} = -(h_b + \Delta h_b^1)\,\bar{b}bH_1^0 - (0 + \Delta h_b^2)\,\bar{b}bH_2^0 + \cdots \tag{2}$$

As we have argued above, there is a clear motivation for studying the large $\tan\beta$ regime of the MSSM. If $\tan\beta$ is large, and after electroweak symmetry breaking, the $\Delta h_b^2$ term can induce corrections of order 1 to down-type fermion masses:

$$m_b = h_b\, v_1\, \left(1 + \Delta h_b^1/h_b + \Delta h_b^2/h_b\,\tan\beta\right), \tag{3}$$

or conversely, one can express the renormalized bottom Yukawa coupling as a function of the bottom mass through

$$h_b\, v_1 = \frac{m_b}{1 + \Delta h_b^1/h_b + \Delta h_b^2/h_b\,\tan\beta} \sim \frac{m_b}{1 + \Delta m_b}. \tag{4}$$

The set of quantum corrections included in (4) are universal, in the sense that they equally affect all amplitudes proportional to the bottom Yukawa. To derive (4), one matches the MSSM to a generic 2HDM at a scale $M_{\text{SUSY}}$ of the order of the relevant soft-SUSY-breaking parameters. Alternatively, within the MSSM, and using an on-shell renormalization scheme, these corrections are absorbed into the bottom mass counterterm, $\delta m_b^{\text{SUSY}} \sim -\Delta m_b^{\text{SUSY}}$.

The quantity $\Delta m_b$ is dominated by SUSY-QCD virtual effects,[1] and at the one-loop level can be cast into the simple expression [5]

$$\Delta m_b \sim \Delta m_b^{\text{SQCD}} = \frac{2\alpha_s}{3\pi}\,\mu\, m_{\tilde{g}}\,\tan\beta\, \text{I}(m_{\tilde{b}_1}, m_{\tilde{b}_2}, m_{\tilde{g}}), \tag{5}$$

where the function I is the limit of Passarino–Veltman's $C_0$ for vanishing external momenta.

An interesting property of $\Delta m_b$ is that it does not vanish for $M_{\text{SUSY}}/m_W \to \infty$. The SUSY-QCD contribution, for instance, evaluates to $\alpha_s/(3\pi)\tan\beta$ in this limit. This should never be understood as a non-decoupling behaviour of the MSSM, because the tree-level $h_b$ is not an observable. If the masses of both the SUSY partners and the non-standard higgses ($H$, $A$, $H^\pm$) become large, the SUSY radiative corrections to $h_b$ are cancelled out exactly by one-loop process-dependent corrections. If $m_A$ is not too large with respect to $m_W$, one can expect large deviations from the rule

$$g_{hbb}/g_{h\tau\tau} = m_b/m_\tau, \tag{6}$$

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[1]An analogous quantity, $\Delta m_\tau$, can be defined for the $\tau$-Yukawa. $\Delta m_\tau$, though, receives (generally) smaller SUSY-EW contributions.
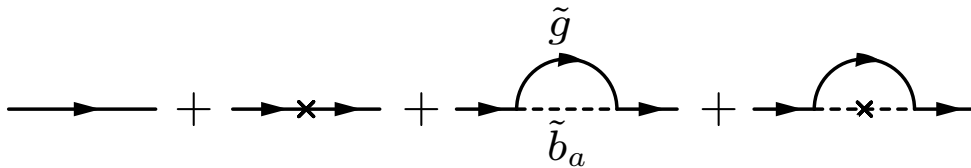
Figure 1: Complete set of Feynman diagrams contributing at order $\alpha_{\mathrm{s}}^n\tan^n\beta$ to the inverse bottom propagator, in SUSY-QCD with no virtual gluons. Dashed and solid internal lines represent sbottom quarks and gluinos, respectively. A cross denotes, in the second diagram, the insertion of a bottom mass-counterterm, and in the last diagram, the counterterm for the $\tilde{b}_L\tilde{b}_R H_2^0$ coupling.

which holds not only in the SM, but also in 2HDM of types I and II [2]. For the $H$, $A$, neutral higgses, eq. (6) can be violated even in the decoupling limit, their masses being of the order of $m_A$. This feature could help in distinguishing the MSSM Higgs sector from a generic type II 2HDM, specially if correlations among various Higgs couplings were checked.

Remarkably enough, it can be shown that, in mass-independent renormalization schemes such as the $\overline{\mathrm{MS}}$, the whole set of SUSY-QCD corrections of the form[2] $\alpha_{\mathrm{s}}^n\tan^n\beta$ are resummed into the above definition for $h_b$ [7] in eq. (4). The proof involves the consideration of the perturbative series for the inverse bottom propagator, which can be used to determine the functional relation between $h_b$ and the bottom pole mass. In a first step one restricts the analysis to the set of diagrams with no gluons: only those in fig. 1 contribute at order $\alpha_{\mathrm{s}}^n\tan^n\beta$, higher-loop diagrams being suppressed either by inverse powers of $\tan\beta$ or by $m_b/M_{\mathrm{SUSY}}$ factors. Requiring the inverse propagator to vanish on-shell, one arrives at (4), apart from $1/\tan\beta$ and $m_b/M_{\mathrm{SUSY}}$ suppressed quantities. The full proof, that is, after allowing for diagrams containing virtual gluons, is more delicate. It requires, for instance, a careful analysis of the infrared behaviour of the extra diagrams, as $1/m_b$ mass singularities would invalidate the counting of $m_b/M_{\mathrm{SUSY}}$ powers used in the proof.

## 3   Prospects for $H^\pm$ searches at hadron colliders

As a first example of the use of eq. (4), we are going to consider the MSSM associated production of a charged Higgs boson, $H^\pm$, with top and bottom at hadron colliders, presenting results for both the Tevatron and the LHC [8]. From our point of view, the relevance of this channel is due to its ability to test the charged Higgs coupling to the third-generation quarks and leptons. Any information obtained about

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[2]The only exception being, for $n = 1$, the process-dependent one-loop effects that restore the SM low-energy limit of the theory for $m_A/m_W \to \infty$.
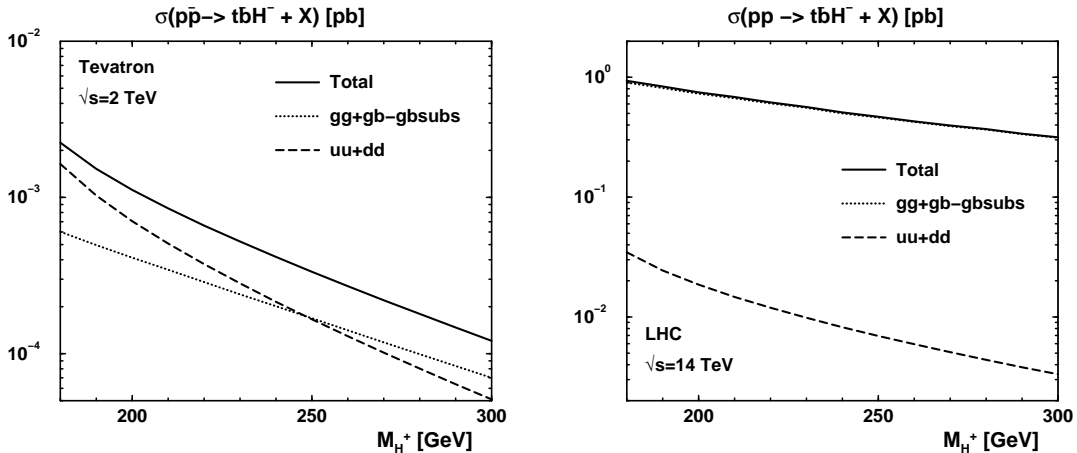
Figure 2: The $p\overline{p}$, $pp \to t\overline{b}H^- + X$ cross-section at the Tevatron Run II (left) and at the LHC (right), for $\mu = -200$ GeV, $\tan\beta = 30$ and $m_{\tilde{g}} = m_{\tilde{t}_1} = m_{\tilde{b}_1} = A_b = A_t = 500$ GeV. The dashed curve corresponds to the $q\overline{q}$ annihilation channel, the dotted curve to the sum of the $gg$-initiated and $gb$-initiated channels, after subtracting double counting in the $gb$ channel. The solid curve is the sum of all channels, $q\overline{q}$, $gg$ and $gb$.

these couplings could provide valuable hints on the exact nature of the Higgs sector.

After the LEP shutdown, charged Higgs searches concentrate on the Tevatron results. Both direct and indirect Tevatron analyses have been limited to the region $m_{H^+} < m_t - m_b$, placing constraints on the $m_{H^+}$–$\mathcal{BR}_{t \to bH^+}$ plane [9], which are usually translated to the $m_{H^+}$–$\tan\beta$ plane once the relevant MSSM parameters are fixed [10].

Beyond the kinematical limit for $t \to bH^+$, apart from the one considered in this talk, there are two other promising production channels: pair production [11] and associated production with a $W$ boson [12]. The work presented here on $p\overline{p}$, $pp \to t\overline{b}H^- + X$ adds, with respect to previous analyses [13,14,15,16,17], a resummation of the leading —in powers of $\alpha_s\tan\beta$, for $\tan\beta \gtrsim 10$— SUSY radiative corrections (both strong and electroweak) and a estimation of the off-shell effects. It will be presented in full detail in [8], including a complete signal and background analysis.

## 3.1 Cross-section computation and results

At the parton level, the reactions $p\overline{p}$, $pp \to t\overline{b}H^- + X$ proceed through three main channels[3]: i) $q\overline{q} \to t\overline{b}H^-$, with $q = u, d$ (the $s$ contribution can be safely neglected), a channel only relevant to the Tevatron [14,15]; ii) $gg \to t\overline{b}H^-$, dominant at the LHC,

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[3]We shall omit the charge-conjugate process, $p\overline{p}$, $pp \to \overline{t}bH^+ + X$, for the sake of brevity. Including this process just amounts to multiplying our cross-section by a factor of 2.

and also at the Tevatron for increasing $H^\pm$ masses [14,15]. Since the bottom mass, $m_b$, is small with respect to the energy of the process, parton distribution functions (PDFs) for $b$-quarks have to be introduced, allowing for the resummation of collinear logs [18]. This provides an extra $bg \to tH^-$ channel contributing to the cross-section. Contrary to i) and ii), in this case the final state contains at most three high-$p_T$ $b$-quarks and, therefore, $bg$-initiated processes cannot be detected by using four high-$p_T$ $b$-tagging. Once a PDF for the $b$-quarks is used, there is some amount of overlap between $bg$- and $gg$-initiated amplitudes, which has to be removed [18,19]. To this end, we follow here the method described in ref. [19], straightforwardly translated to the $t\bar{b}H^-$ final state case (see also [14]). Figure 2 shows the relative relevance of the various channels to both the Tevatron Run II and the LHC, as explained above. The solid curve can be used to get a rough estimate of the reach of the process $p\bar{p}, pp \to t\bar{b}H^- + X$ in the search for a MSSM charged Higgs boson in these machines for the given parameters (see the caption).

The amplitude for $p\bar{p}, pp \to t\bar{b}H^- + X$ is, for large $\tan\beta$, approximately proportional to the Yukawa coupling of the bottom quark and, as explained above, receives supersymmetric quantum corrections that can be of order 1. An analysis of the reach of $p\bar{p}, pp \to t\bar{b}H^- + X$ in $H^\pm$ searches demands the appropriate inclusion and resummation (using (4)) of such corrections in the computation of the cross-section. For the present work we have also included the full off-shell SUSY-QCD and SUSY-EW corrections to the $H^+\bar{t}b$ vertex and to the fermion propagators, although $\Delta m_b$ in (4) is the only correction contributing at order $(\alpha_s/4\pi)^n \tan^n\beta$ and thus dominates for large $\tan\beta$. In fact, the approximation of neglecting vertex and propagator corrections in the cross-section, which we call "improved Born" approximation, is really justified in that region (see fig. 3).

We disregard virtual supersymmetric effects on the $gqq$ and $ggg$ vertices and on the gluon propagators. We expect those to be of order $(\alpha_s/4\pi)\cdot(\sqrt{s}/M_{\text{SUSY}})$, with no $\tan\beta$ enhancement, and thus suppressed both by a loop factor (any reasonable choice for $\alpha_s(Q)$ will be small) and by a MSSM form factor coming from the loop integrals. Therefore, we can neglect these contributions as we are only considering large $\tan\beta$ values. Besides, the cross-section for the signal is much smaller for $\tan\beta$ close to 1, so our approximation is well justified.

The only other source of potentially large radiative corrections is, of course, standard QCD. At least one group is currently addressing the NLO QCD correction to $p\bar{p}, pp \to bbH + X$, which can provide a good guess for the sign and size of the corrections in $p\bar{p}, pp \to t\bar{b}H^- + X$. In the meantime, we can parametrize our ignorance by using a K-factor ranging between 1.2 and 1.5 [22,16,20]. Once the exact value of K will be known, it will be easy to conveniently rescale our plots to take into account the effect of the gluon loops. The only QCD corrections we do incorporate to the cross-

section are those related to the running of $\alpha_{\mathrm{s}}(Q)$ and $\overline{m}_b(Q)$.[4] We choose to work with equal renormalization, $\mu_R$, and factorization, $Q$, scales fixed at $\mu_R = Q = m_t + m_{H^+}$.

Concerning the method employed to compute the squared matrix elements, we have made intensive use of the package CompHEP [24], for both the signal and background processes. Although CompHEP is only able to deal with tree-level calculations, we have managed to add the supersymmetric corrections to the $t\overline{b}H^-$ vertex and to the fermion propagators in the following way: first, we have modified CompHEP's Feynman rules to allow for the most general off-shell $t\overline{b}H^-$ vertex, then we have let CompHEP reckon the squared matrix elements and dump the result into REDUCE code. At this point, we have inserted expressions for the coefficients of the off-shell $t\overline{b}H^-$ vertex that include the one-loop off-shell supersymmetric corrections to the vertex itself and to the off-shell fermion propagators and fermionic external lines.[5] Only half the renormalization of an internal fermion line has to be included, the other half being associated to the $gqq$ vertex. This procedure has allowed us to estimate the relative size of the off-shell effects in the signal cross-section, which never exceeds the few per cent level.

In fig. 3, we compare the above various approximations to the $p\overline{p} \rightarrow t\overline{b}H^- + X$ cross-section at the Tevatron Run II. The curves correspond to the total cross-section, as a function of $\tan\beta$, for a centre-of-mass energy of 2 TeV and a charged Higgs mass of 250 GeV. The tree-level result is given by the dotted line; it grows almost quadratically with $\tan\beta$. After including the $\overline{\mathrm{MS}}$ off-shell one-loop supersymmetric corrections in the $t\overline{b}H^-$ vertex, in the internal fermion propagators and in the external fermion lines, one obtains the dashed line. For the chosen parameters, that is $\mu = -200$ GeV, $m_{\tilde{g}} = m_{\tilde{t}_1} = m_{\tilde{b}_1} = A_b = A_t = 500$ GeV, the overall correction turns out to be positive, as it is driven by the SUSY-QCD correction in (4), which is positive for negative $\mu$. The resummation of the order $\alpha_{\mathrm{s}}^n \tan^n\beta$ supersymmetric corrections further increases the result up to the top solid curve, labelled "improved" $\overline{\mathrm{MS}}$. The effect is not dramatic because $\mu$ is sizeably smaller than the rest of the relevant soft-SUSY-breaking parameters, namely the gluino and sbottom masses. To illustrate the possibility of a suppression of the cross-section due to virtual supersymmetric effects, we also plot the resummed result for the same parameters but taking $\mu = 200$ GeV and $A_t = -500$ GeV, which corresponds to the bottom (red) solid line. It does not differ much from the tree-level because of a partial cancellation of the correction due to the effect of the resummed high-order terms in (4). Finally, the dot-dashed curve is obtained by just replacing the $\overline{m}_b(Q)$ in the tree-level approximation to the cross-section with $\overline{m}_b(Q)/(1 + \Delta m_b)$, as suggested by (4). This "improved" tree-level

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[4]In the $t \rightarrow bH^+$ decay rate, this actually accounts for most of the QCD virtual effects (for $Q = m_t$) [23].

[5]We shall not write down here the analytic expressions for the renormalized vertex and propagators. They can easily be derived by just generalizing previous on-shell calculations, such as those for $t \rightarrow bH^+$ [3].
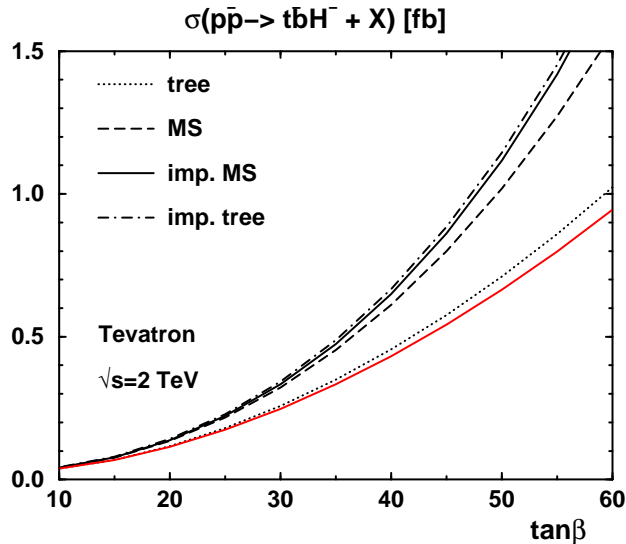
Figure 3: Various approximations to $\sigma(p\overline{p} \to t\overline{b}H^- + X)$ at the Tevatron Run II, as a function of $\tan\beta$, for a charged Higgs mass of 250 GeV. The remaining MSSM parameters are set as in fig. 2. Shown are the tree-level (dotted line), the one-loop $\overline{\text{MS}}$ (dashed line), the improved or resummed $\overline{\text{MS}}$ (top solid line and bottom red solid line, the latter for $\mu = 200$ GeV and $A_t = -500$ GeV) and the improved tree-level (dot-dashed line) results.

constitutes a fairly good approximation to the complete resummed $\overline{\text{MS}}$ result (top solid curve). Similar conclusions apply for $pp \to t\overline{b}H^- + X$ at the LHC.

## 4   $b \to s\gamma$ and supersymmetry with large $\tan\beta$

The computation of the $b \to s\gamma$ branching ratio at the NLO in the MSSM is clearly a complicated matter [25,26], and completely general expressions have not yet been derived. Nevertheless, it turns out that the leading $\alpha_{\text{s}}^n \tan^{n+1}\beta$ corrections can be calculated and resummed to all orders in perturbation theory by replacing the tree-level Yukawa of the bottom quark by eq. (4) in the Wilson coefficients of the leading-order (LO) computation [27].

To show how this procedure works, let us start by identifying the dominant one-loop Feynman diagrams. In a type II 2HDM, the diagram contributing to the highest $\tan\beta$ power is that in fig. 4, with the exchange of a virtual charged Higgs in the loop, which is proportional to $h_b \cdot h_t \cos\beta$ and, therefore, of order $\tan^0\beta$. Substituting a chargino by the charged Higgs, one obtains the leading diagram in the MSSM, which is proportional to $h_b \cdot h_t$ and of order $\tan\beta$. We have checked, using the formulae in [25], that the leading —in $\tan\beta$— NLO corrections follow from the replacement
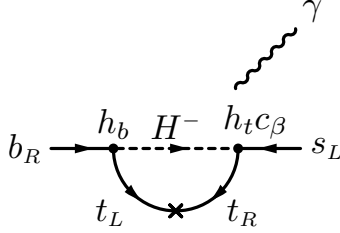
Figure 4: Dominant, in powers of $\tan\beta$, LO contribution to the $b \to s\gamma$ branching ratio in the 2HDM.

$h_b \to -h_b\Delta m_b$ (expanding and truncating at first order in $\alpha_{\rm s}$) in the LO contributions associated to the above-mentioned diagrams. These NLO corrections are to be associated with the insertion of the counterterm for the bottom Yukawa in the LO, and thus have their origin in a one-loop diagram. The resulting effect is of order $\alpha_{\rm s}\tan^2\beta$ for the chargino loop, and $\alpha_{\rm s}\tan\beta$ for the $H^-$ loop.

There is one additional source of $\tan\beta$-enhanced corrections in the charged Higgs diagrams, which is not related to the bottom mass counterterm [25]: while the tree-level $H^+\bar{t}_R s_L$ vertex is suppressed by $1/\tan\beta$, this suppression is absent at the one-loop level, so that the NLO charged-Higgs contribution to $\mathcal{BR}(b \to s\gamma)$ is $\tan\beta$-enhanced with respect to the LO one. No enhancement occurs in higher-loop diagrams, which are suppressed either by $1/\tan\beta$ powers or by $m_b/M_{\rm SUSY}$ factors [7].

Now that all sources of $\tan\beta$-enhanced terms have been identified, it is an easy matter to proceed with the improvement of the NLO expressions, that is, to resum all terms of order $\alpha_{\rm s}^n\tan^{n+1}\beta$. Just replace, in the LO expressions, $h_b$ by $h_b/(1 + \Delta m_b)$, and remove double counting, i.e. $-h_b\Delta m_b$ terms, in the NLO formulae (see ref. [27] for a detailed description of the procedure). This is enough to extend the validity of the calculation presented in ref. [25] to large values of $\tan\beta$.

The quantitative effect of the improvement can be assessed from fig. 5, where we compare the NLO theoretical prediction for $\mathcal{BR}(b \to s\gamma)$ of ref. [25] with (solid line) and without (dashed line) including the all-order resummation of the dominant $\tan\beta$-enhanced radiative corrections, for typical values of the supersymmetric parameters. We use a negative value of $A_t = -500$ GeV at low energies, with $m_{\tilde{g}}A_t < 0$. This choice is inspired in supergravity models, where this sign relation holds unless the boundary value of $A_t$ at the high-energy input scale is one order of magnitude larger than the gaugino soft-supersymmetry-breaking mass parameters [4,28]. For the experimental measurement of the $b \to s\gamma$ branching ratio, we use the combined result of CLEO [29] and ALEPH [30], $\mathcal{BR}(b \to s\gamma) = (3.14 \pm 0.48) \times 10^{-4}$. Owing to cancellations among the various contributions, the relative size of the effect turns out to be sizeable only for $\mu A_t < 0$: for the set of parameters used in fig. 5, it decreases the NLO result by 20% at $\tan\beta \simeq 30$.
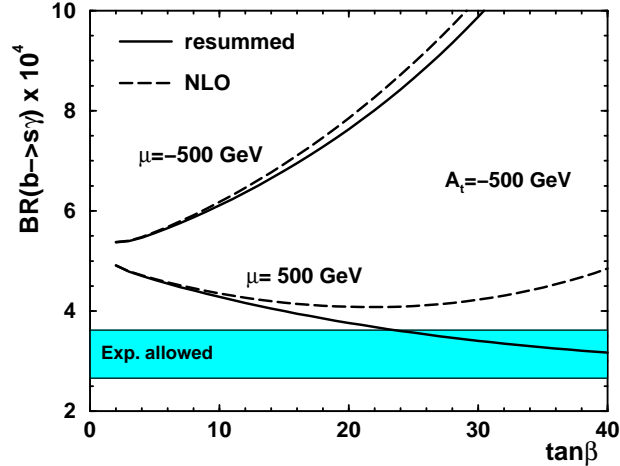
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Figure 5: Comparison of the NLO theoretical prediction for $\mathcal{BR}(b \to s\gamma)$ of ref. [25] before (solid line) and after (dashed line) performing a resummation of the dominant supersymmetric corrections of order $\alpha_{\mathrm{s}}^n \tan^{n+1}\beta$, as explained in the text. The charged-Higgs boson mass is 200 GeV and the light stop mass is 250 GeV. The values of $\mu$ and $A_t$ are indicated in the plot, while $M_2$, the gluino, heavy-stop and down-squark masses are set at 800 GeV.

Notice that, with our sign conventions, positive values of $\mu$ are necessary in order to obtain correct values for $\mathcal{BR}(b \to s\gamma)$, even after considering higher-order effects, within minimal supergravity models, for which, as explained in the above paragraph, the sign of $A_t$ at low energies tends to be negative. This is in contradiction with the results of ref. [31]. We believe sign errors in the charged Goldstone and Higgs couplings to stop and down-like squarks in the published version of [25] are at the origin of this discrepancy (see [27]).[6]

## 5   Conclusions

Motivated by the latest LEP analyses ruling out $\tan\beta$ values around 1 [6] and by large $\tan\beta$ supersymmetric SO(10) models, we have analysed the Yukawa couplings of down-type fermions in the MSSM, which receive potentially large supersymmetric corrections when $\tan\beta$ is large. We claim that, in observables where these couplings are relevant, a resummation of the leading —in powers of $\tan\beta$— subset of corrections is needed if the observables are to be computed for large $\tan\beta$. We perform such resummation, which is essentially independent of the particular process under consideration, and explore its numerical impact on two exemplary processes: the $p\bar{p}$, $pp \to t\bar{b}H^- + X$ cross-section and the branching ratio of $b \to s\gamma$.

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[6]The authors of [25] have independently detected these sign errors, and posted a corrected version of the paper to the hep-ph archive.

# Acknowledgments

# Note added

After presentation of this talk, and concerning the $b \to s\gamma$ branching ratio, similar formulae for the improvement of the NLO computation in the MSSM with large $\tan\beta$ have been given in [32].

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