Charged Higgs Boson Production in Bottom-Gluon Fusion

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Abstract

We compute the complete next-to-leading order SUSY-QCD corrections for the associated production of a charged Higgs boson with a top quark via bottom-gluon fusion. The set of higher order corrections can be split into corrections in a general two Higgs doublet model and additional massive supersymmetric loop contributions. These supersymmetric contributions consist of the universal bottom Yukawa coupling corrections and non-factorizable diagrams. We find that over most of the relevant supersymmetric parameter space the Yukawa coupling corrections are sizeable, while the supersymmetric loop contributions remain negligible.

Higgs Physics at the LHC

Among the physics goals for the LHC the exploration of the nature of electroweak symmetry breaking is the most prominent and the one closest to current observations. The LEP precision measurements suggest the existence of a light Higgs boson which coincides with the prediction of the MSSM. It has been shown that the discovery of at least one Higgs boson in the MSSM will probably not pose any problem once the LHC collects a minimal integrated luminosity [1]. There is, however, only one way to tell the supersymmetric Higgs sector from its Standard Model counterpart: to discover the additional heavy Higgs bosons. While the chances of finding a heavy Higgs boson with a small value of $\tan\beta$ at the LHC are rather slim, the discovery of all heavy Higgs scalars in the large $\tan\beta$ regime is likely. The production of a charged Higgs boson in association with a top quark $gg \to \bar{b}tH^-$ and a subsequent decay into a tau lepton and a neutrino seems to be a promising search channel [2, 3, 4] and experimentally easier than for example the charged Higgs boson pair production [5]. Recently both LHC experiments have published detailed studies of this production channel with very promising results [6].

Bottom Parton Scattering

As a starting point we emphasize that the exclusive production channel $gg \to \bar{b}tH^-$ is perfectly consistent in the sense that it includes the squared matrix element to order $\alpha_s^2 y_{b,t}^2$, but beyond naive perturbation theory the integration over the final state bottom quark gives rise to possibly large logarithms [7]. The massive bottom propagator leads to an asymptotic transverse mass dependence $1/m_{T,b}$, so the infrared divergence is regularized by the bottom mass. For small transverse bottom momenta the differential partonic cross section approaches the asymptotic form $d\sigma/dp_{T,b} \propto p_{T,b}/m_{T,b}^2$. The transverse momentum of the bottom quark can be integrated out up to a factorization scale and yields a total cross section $\sigma_{\rm tot} \propto \log(\mu_F^2/m_b^2 + 1)$. Even though the logarithms $\log(p_{T,b}/m_b)$ which appear when we integrate the exclusive cross section are not divergent they can become large. Switching to a bottom parton description $bg \to tH^-$ corresponds to a resummation of these potentially large logarithms. This procedure relies on several approximations, which should be carefully examined.

First we assume that at leading order the intermediate bottom quark and therefore the outgoing bottom jet are collinear with the incoming gluon. This approximation will never be perfect, since the cutoff parameter m_b is only slightly smaller than the minimum observable transverse

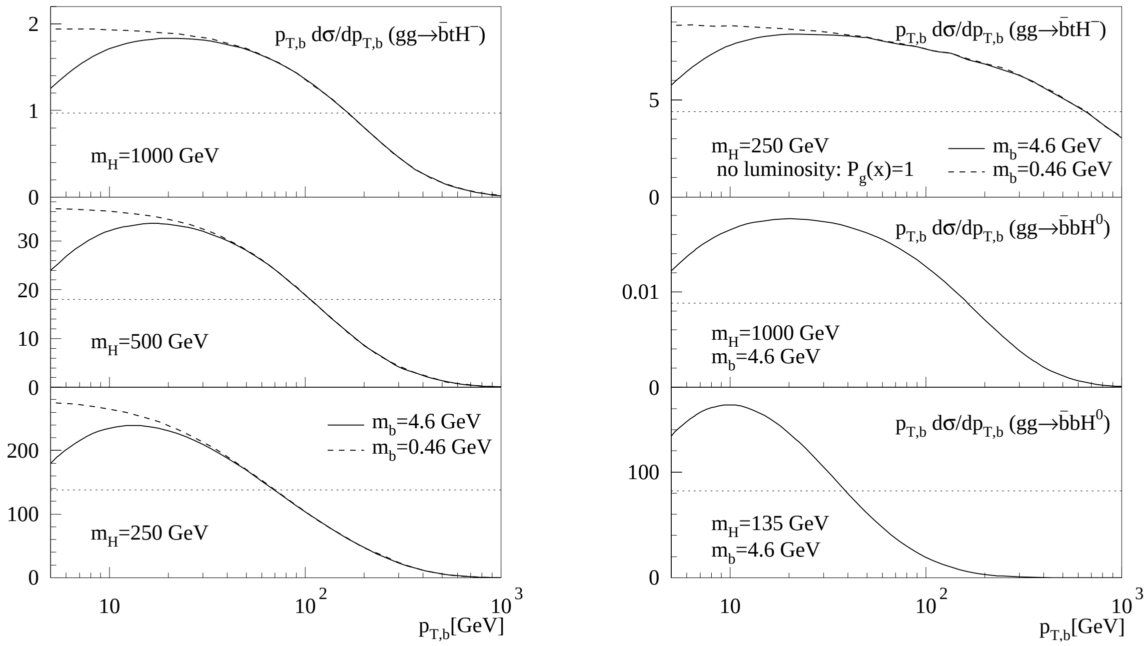

Figure 1: Left: the bottom transverse momentum distribution for exclusive The curves are given for the physical on-shell bottom mass $4.6\ \text{GeV}$ as well as for a smaller bottom mass as the infrared regulator. The thin dotted line indicates half the height of the plateau. Right: in the upper panel the same distribution for a heavy charged Higgs boson, but with the gluon luminosity set to unity. Below this in the two lower panels the same distribution for exclusive neutral Higgs boson production $gg \to \bar{b}bH$.

momentum at a collider. We show this collinear behavior for the exclusive charged Higgs boson production in Ref. [2].

After making sure that the collinear approximation describes the exclusive tree-level process we still have to determine if there are large logarithms to resum. While the $1/m_{T,b}$ behavior is by definition present in the matrix element, this is not necessarily true for the differential hadronic cross section $d\sigma/dp_{T,b}$. In Fig. 1 we show the $1/p_{T,b}$ behavior of the hadronic distributions. First of all we see how the zero bottom mass approximation breaks down when the transverse momentum is of the order of the bottom mass. If we replace the on-shell bottom mass with a smaller value the plateau extends to smaller transverse momentum, confirming the asymptotic behavior. The low end of the plateau in the transverse momentum spectrum, however, does not lead to large numerical effects. For those we have to focus on the high $p_{T,b}$ end. We see how the high $p_{T,b}$ end of the plateau roughly scales with the average mass in the final state. This coincides with the observation that the only scales allowed for the evaluation of total cross sections are external scales. They are typically chosen proportional to the average mass of the final state particles $\mu_F \sim Cm_{\text{av}}$ where the proportionality factor C is arbitrary. The curves shows that the naive choice $C \sim 1$ is not appropriate. This choice assumes large logarithms $\log(p_{T,b}^{\max}/m_b)$ which are resummed to values $\mu_F \sim m_{\text{av}}$ and it will therefore yield an overestimate of the total cross section.

Evaluating the expression for the exclusive asymptotic total cross section for two bottom masses we can determine the appropriate factorization scale μ_F. We obtain $185, 120, 80\ \text{GeV}$ for the three Higgs boson masses $1000, 500, 250\ \text{GeV}$, similar to our naive observation. This

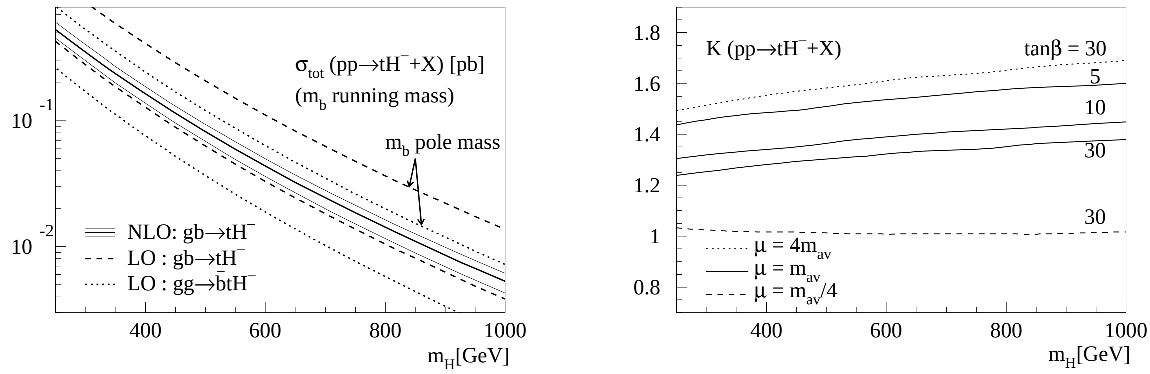

Figure 2: Left: the inclusive and exclusive production cross section $pp \to tH^-$. Both tree level results are also quoted using the (inappropriate) pole mass for the bottom Yukawa coupling. Right: the corresponding consistent K factors for the three values of $\tan\beta = 5, 10, 30$.

means that the appropriate factorization scale indeed scales with the average final state mass, but with $C \sim 1/3$. On the other hand we also emphasise that for associated charged Higgs boson and top quark production the bottom parton treatment is justified — with an appropriate choice of the bottom parton factorization scale. In the right panel of Fig. 1 we see that he asymptotic behavior with the gluon luminosity set to unity extends to much larger values. This means that the low scales observed in the left panel of Fig. 1 are due to the steeply falling gluon density which suppresses any large transverse momentum radiation of forward bottom jets.

To prove the universality of our argument we show the same distribution for the exclusive neutral Higgs boson production $gg \to b\bar{b}H$, which can be evaluated as partly [8] or completely [9] inclusive. The same reasoning as for the charged Higgs boson production applies in this case. First one shows that the bottom quarks are collinear. Then one determines an appropriate choice of the factorization scale. From the comparison of the two curves for a $1\ \mathrm{TeV}$ neutral and a $1\ \mathrm{TeV}$ charged Higgs boson we see that the behavior is very similar: the bottom parton description is valid, and the factorization scale should be chosen considerably below the respective final state mass. For a light neutral Higgs boson the asymptotic behavior only survives up to $p_{T,b} \lesssim 40\ \mathrm{GeV}$, which corresponds to a logarithmic enhancement $\log(p_{T,b}/m_b) \lesssim \log 8 \sim 2$.

The applicability of the bottom parton approach is very closely tied to the reason why the partly inclusive analyses are attractive: if the exclusive process exhibits a collinear final state bottom jet this jet is not likely to hit the detector, much less to be tagged. This means that the same feature which allows us to use the bottom parton approach makes it hard to utilize the exclusive process: the final state bottom jet is too collinear to be particularly useful. We therefore strongly advocate use of the inclusive cross section prediction, since the reliability of the cross section predictions will be significantly improved beyond naive perturbation theory.

Next-to-leading Order Results for a Two Higgs Doublet Model

To improve the theoretical cross section prediction and to reduce the theoretical uncertainty we compute the inclusive process $pp \to gb \to tH^-$ to next-to-leading order QCD. The corrections include virtual gluon loops as well as gluon radiation. The massive supersymmetric loops will be discussed in the next section. The strong coupling and the bottom Yukawa coupling are renormalized in the $\overline{\mathrm{MS}}$ scheme. This way α_s and $y_{b,t}$ both are running parameters, dependent on the renormalization scale μ_R. The numerical impact of the higher order contributions is

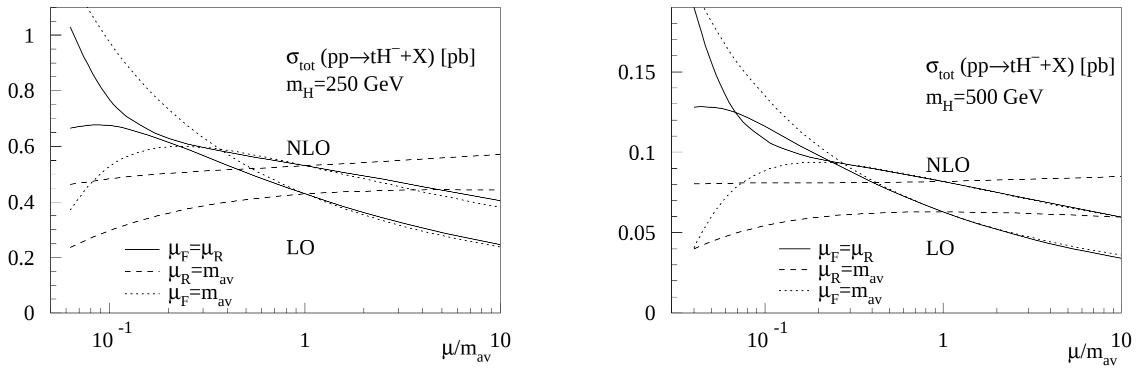

Figure 3: The variation of the total inclusive cross section $pp \to tH^-$ as a function of the renormalization and factorization scales, around the central value $\mu = m_{\rm av}$. The lower end of the curves corresponds to $\mu \sim 10$ GeV.

shown in Fig. 2. The leading order results are given for the running bottom mass as well as for the bottom pole mass in the Yukawa coupling. We want to stress, however, that the pole mass Yukawa coupling always yields a huge overestimate of cross sections and should generally not be used [10]. The corrections described by the K factor are perturbatively well under control, ranging from $+30\%$ to $+40\%$ for $\tan\beta = 30$ and Higgs boson masses between 250 and 1000 GeV.

The QCD corrections are flavor blind and proportional to the Born coupling structure $y_{b,t}^2$, which as a function of $\tan\beta$ is either dominated by the top quark or by the bottom quark Yukawa coupling. However, the shift in the consistent bottom Yukawa coupling absorbs another factor $y_{b,2-\text{loop}}^2/y_{b,1-\text{loop}}^2 \sim 0.84$, while the top Yukawa coupling is essentially stable. The difference between the three curves for the K factor comes from the running Yukawa coupling, which is dominantly bottom for large values of $\tan\beta$. The consequence is a larger K factor for smaller values of $\tan\beta$.

In Fig. 3 we see that the leading order dependence of the cross section on the factorization scale becomes large only once the bottom factorization scale comes close to the bottom mass, where it has to vanish for $\mu_F \to m_b$. To next-to-leading order the scale dependence stays flat even for very small factorization scales. Assuming that light flavor quark initiated processes are suppressed at the LHC the purely gluon initiated exclusive process $gg \to tH^- + X$ dominates for factorization scales $\mu_F \to m_b$. This way the next-to-leading order inclusive calculation interpolates between the inclusive and the exclusive results. At the one-loop level the next-to-leading order inclusive cross section approaches the exclusive tree level result in the limit of no large logarithms, where the enhancement through the resummation disappears.

The dominant theoretical uncertainty comes from the unknown renormalization scale of the strong coupling. A small renormalization scale yields a larger strong coupling. Identifying both scales inherently leads to a cancellation. Moreover, if we evaluate the cross section for very small values $\mu/m_{\rm av} \lesssim 0.1$ the next-to-leading order prediction increases rapidly. Physically this is not a problem, since the scales have to be very small, which is certainly not appropriate for the renormalization scale. We know that for small scales the dependence on the logarithms $\log(\mu_F/m_b)$ and $\log(\mu_R/m_H)$ largely cancels. However, terms proportional to $\log(\mu_F/m_b)\times\log(\mu_R/m_H)$ become large. One way to look at this effect is that the unphysically small renormalization scale gives a large negative prefactor for the factorization scale dependence, namely $\log(\mu_R^2/m_H^2)$. This dominates the factor in front of $\log(\mu_F/m_b)$, which for more appropriate renormalization scales is small and positive instead.

For a reasonably large renormalization scale almost the entire scale variation is driven by the renormalization scale. This effect is well known from supersymmetric particle production at the LHC: for processes mediated by a strong coupling at tree level, the scale variation is an appropriate measure for the theoretical uncertainty. On the other hand, for weakly interacting particles produced in Drell–Yan type processes, the leading order scale variation is dominated by the factorization scale and is not a good measure for the theoretical uncertainty. In the process considered here the remaining theoretical uncertainty can be estimated to be $\lesssim 20\%$ for a central choice of scales.

Next-to-leading Order Results with Supersymmetry

Even though the Standard Model with a two doublet Higgs sector is a perfectly well-defined renormalizable theory, we are particularly interested in the MSSM version of this model. There the number of free tree level parameters in the Higgs sector is reduced to two, which are usually chosen to be the pseudoscalar mass m_A and $\tan\beta$. All next-to-leading order corrections to the total cross section coming from supersymmetric loop diagrams we include in a supersymmetric correction factor $K_{\rm SUSY} = (\sigma_{\rm SUSY} + \sigma_{\rm NLO})/\sigma_{\rm NLO}$

At one-loop order the off-diagonal entry in the sbottom mass matrix can connect a left handed with a right handed bottom quark. Even though in the final result we neglect the bottom mass we do have to take into account this contribution to the bottom mass counter term. This mass counter terms Δm_b modifies the relation between the bottom mass and the bottom Yukawa coupling [11, 12]. The authors of Ref. [12] have shown that this correction is the leading term in powers of $\tan\beta$, where the charged Higgs boson search is promising. The reason why this contribution is usually referred to as non-decoupling is that for large supersymmetric particle masses in the loop *and for a large trilinear mass parameter A_b or higgsino mass parameter μ*, the correction to the Yukawa coupling does not vanish. This is well understood, since at the one-loop level it couples the 'wrong' Higgs doublet to the bottom quarks. To estimate how good the leading $\tan\beta$ approximation given by Δm_b is, we also compute the whole set of MSSM loop diagrams. The result for two different Higgs boson masses is shown in Fig. 4(a). To simplify the presentation we choose a diagonal line in the mSUGRA parameter space: the scalar and gaugino mass scales are identified $m_{\rm SUGRA} = m_0 \equiv m_{1/2}$. For the Δm_b corrections the sign of the higgsino mass parameter is crucial: for $\mu < 0$ we find $\Delta m_b < 0$, which enhances the cross section. For the opposite sign of μ the Δm_b corrections to the production cross section are negative. The supersymmetric corrections apart from the Δm_b corrections are negligible in comparison with the Δm_b terms. This is a feature of the large value of $\tan\beta$ and is even more pronounced for $\tan\beta = 50$ in Fig. 4(b). We note, however, that the picture changes significantly once we do not run the higgsino mass parameter $|\mu|$ to large values, together with the other heavy supersymmetric masses. In that case the Δm_b corrections decouple as shown in Fig. 4(c). Moreover, for a value $\tan\beta = 10$ the Δm_b correction drops below a $\pm 2\%$ effect, becoming even smaller than the explicit MSSM loop corrections. We note, however, that choosing large values for $\tan\beta$ and $|\mu|$ can in principle lead to almost arbitrarily large Δm_b effects, only limited by unitarity constraints.

Heavy particle loops contribute to both the running strong coupling $\alpha_s(\mu_R)$ and the third generation Yukawa coupling $y_{b,t}(\mu_R)$. They give rise to supersymmetric counter terms and can

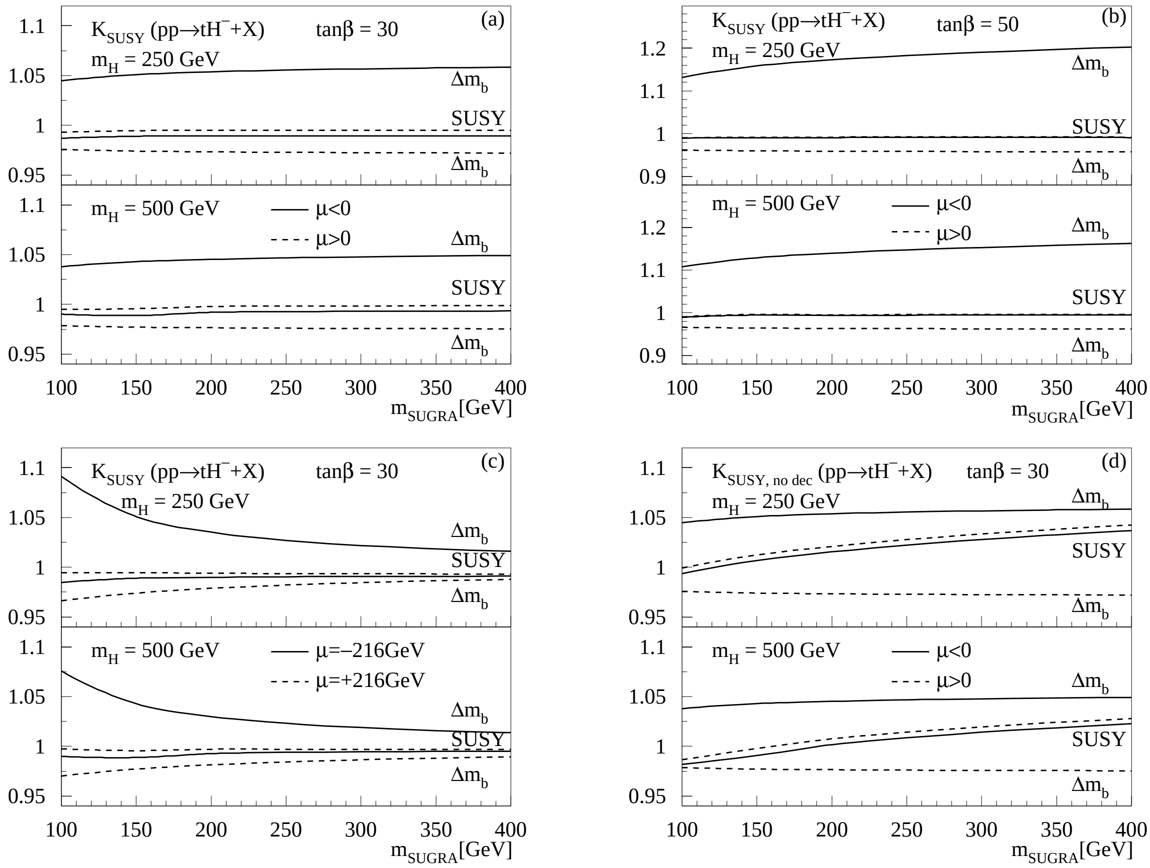

Figure 4: The dependence of the total cross section $pp \to tH^- + X$ on supersymmetric loop contributions. The mass scale is defined as $m_{\rm SUGRA} = m_0 \equiv m_{1/2}$: (a) corrections for $\tan\beta = 30$ and with a running higgsino mass parameter; (b) same as (a), but with $\tan\beta = 50$; (c) same as (a), but with μ fixed at its value for $m_{\rm SUGRA} = 150$ GeV; (d) same as (a), but without decoupling the heavy spectrum from the running Yukawa coupling.

thereby yield a logarithmic divergence $\log(m_{\rm heavy}/\mu_R)$ in the cross section. On the other hand we use Standard Model measurements for these observables, which means that their running has to be governed by the light particle beta function. The contributions from heavy particles to their beta function has to be explicitly cancelled, and as expected this decoupling absorbs all logarithmically divergences in the one-loop cross section. We show the (misleading) result one would get without decoupling the heavy particles from the running Yukawa coupling in Fig. 4(d).

Summary

For the inclusive process $pp \to tH^-$ we show why the bottom parton approach is valid and gives a numerically reliable prediction for the cross section. The one-loop contributions hugely improve the theoretical uncertainty of the leading order cross section prediction to $\lesssim 20\%$. The over-all corrections to the total cross section in the two Higgs double model range between $+30\%$ and $+40\%$ for Higgs boson masses between 250 and 1000 GeV for the average final state mass scale choice.

Two kinds of supersymmetric corrections appear in addition: the on-shell renormalization of the bottom quark mass alters the relation between the bottom mass and the bottom Yukawa coupling. These Δm_b corrections are the leading supersymmetric one-loop corrections with respect to powers of $\tan\beta$. Their effect on the total cross section in a simple mSUGRA model we estimate to stay below $\pm 5\%$ for $\tan\beta = 30$ and below $\pm 20\%$ for $\tan\beta = 50$. Because the charged Higgs boson searches are most promising in the large $\tan\beta$ regime the remaining explicit supersymmetric loop diagrams only contribute on a negligible few percent level.

Bibliography

[1] ATLAS TDR, report CERN/LHCC/99-15 (1999); CMS TP, report CERN/LHCC/94-38 (1994); T. Plehn, D. Rainwater, and D. Zeppenfeld, Phys. Lett. B **454**, 297 (1999) and Phys. Rev. D **61**, 093005 (2000).

[2] T. Plehn, arXiv:hep-ph/0206121.

[3] A. C. Bawa, C. S. Kim, and A. D. Martin, Z. Phys. C **47**, 75 (1990); D. P. Roy, Phys. Lett. B **277**, 183 (1992); and Phys. Lett. B **459**, 607 (1999). F. Borzumati, J. L. Kneur, and N. Polonsky, Phys. Rev. D **60**, 115011 (1999); A. Belyaev, D. Garcia, J. Guasch, and J. Sola, arXiv:hep-ph/0203031.

[4] L. G. Jin, C. S. Li, R. J. Oakes, and S. H. Zhu, Eur. Phys. J. C **14**, 91 (2000); S. H. Zhu, arXiv:hep-ph/0112109. For all diagrams included in both calculations the results agree within the uncertainties from the scheme dependence in the top mass renormalization.

[5] S. S. Willenbrock, Phys. Rev. D **35**, 173 (1987); A. Krause, T. Plehn, M. Spira, and P. M. Zerwas, Nucl. Phys. B **519**, 85 (1998); O. Brein and W. Hollik, Eur. Phys. J. C **13**, 175 (2000);

[6] K. A. Assamagan and Y. Coadou, Acta Phys. Polon. B **33**, 707 (2002); D. Denegri *et al.*, report CMS-NOTE-2001-032, arXiv:hep-ph/0112045.

[7] J. C. Collins and W. K. Tung, Nucl. Phys. B **278**, 934 (1986); M. A. Aivazis, J. C. Collins, F. I. Olness, and W. K. Tung, energies," Phys. Rev. D **50**, 3102 (1994); F. I. Olness and W. K. Tung, Nucl. Phys. B **308**, 813 (1988); R. M. Barnett, H. E. Haber, and D. E. Soper, Nucl. Phys. B **306**, 697 (1988); J. C. Collins, Phys. Rev. D **58**, 094002 (1998).

[8] C. S. Huang and S. H. Zhu, Phys. Rev. D **60**, 075012 (1999); J. Campbell, R. K. Ellis, F. Maltoni, and S. Willenbrock, arXiv:hep-ph/0204093.

[9] D. A. Dicus and S. Willenbrock, Phys. Rev. D **39**, 751 (1989); D. Dicus, T. Stelzer, Z. Sullivan, and S. Willenbrock, Phys. Rev. D **59**, 094016 (1999).

[10] E. Braaten and J. P. Leveille, Phys. Rev. D **22**, 715 (1980). S. G. Gorishnii, A. L. Kataev, and S. A. Larin, Sov. J. Nucl. Phys. **40**, 329 (1984) [Yad. Fiz. **40**, 517 (1984)]; A. L. Kataev and V. T. Kim, Mod. Phys. Lett. A **9**, 1309 (1994); K. Melnikov, Phys. Rev. D **53**, 5020 (1996).

[11] L. J. Hall, R. Rattazzi and U. Sarid, Phys. Rev. D **50**, 7048 (1994); M. Carena, M. Olechowski, S. Pokorski, and C. E. Wagner, Nucl. Phys. B **426**, 269 (1994); A. Belyaev, D. Garcia, J. Guasch, and J. Sola, Phys. Rev. D **65**, 031701 (2002).

[12] M. Carena, D. Garcia, U. Nierste, and C. E. Wagner, Nucl. Phys. B **577**, 88 (2000).