Yukawa Corrections to Single Top Quark Production at the Fermilab Tevatron in the Two-Higgs-Doublet Models

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ABSTRACT

We calculate Yukawa corrections of order $\frac{\alpha_{ew} M_t^2}{M_W^2}$ to single top quark production via $q\bar{q}' \rightarrow t\bar{b}$ at the Fermilab Tevatron in the two-Higgs-doublet models. In our calculation we also keep the terms proportional to $M_b^2 \tan^2 \beta$ since their effects may become rather important for large $\tan \beta$. The corrections can amount to more than a 15% reduction in the production cross section relative to the tree level result in the general two-Higgs-doublet model, and a 10% enhancement in the minimal supersymmetric model, which might be observable at a high-luminosity Tevatron.

PACS number: 14.80Dq; 12.38Bx; 14.80.Gt
1. Introduction

The top-quark physics has become a very active research area since the top quark was discovered by the CDF and D0 Collaborations at the Fermilab Tevatron[1]. It is timely to focus attention on directly investigating the properties of the top quark, especially its production mechanisms. With the increase in the number of top quark events at the Tevatron, the experimental errors are expected to be further reduced. With the next Tevatron run at $\sqrt{s} = 2.0$ TeV, one can expect about twenty times as much data as exist now. Thus, the comparison between the observed top quark production properties and more precise theoretical calculations will be an important probe for the possible existence of new physics. At the Tevatron top quarks are produced primarily via two independent mechanisms: The dominant production mechanism is the QCD pair production process $q\bar{q} \rightarrow t\bar{t}$[2]. Single top production via $W$-gluon fusion subprocess $g + W \rightarrow t\bar{b}$ [3] and the subprocess $q\bar{q}' \rightarrow t\bar{b}$ [4] are also important. These latter processes involve the electroweak interaction and, therefore, can probe the electroweak sector of the theory, in contrast to the QCD pair production mechanism. A recent analysis[5] of the process $q\bar{q}' \rightarrow t\bar{b}$ showed that it is potentially observable at the Tevatron with 2-3 $fb^{-1}$ of integrated luminosity. This process probes the top quark with a timelike $W$ boson, $q^2 > (M_t + M_b)^2$, while the $W$-gluon fusion process involves a spacelike $W$ boson, $q^2 < 0$, and these processes are therefore complementary. Moreover, in the Standard Model (SM), the process $q\bar{q}' \rightarrow t\bar{b}$ can be reliably predicted and the theoretical uncertainty in the cross section is only about a few percent due to QCD corrections[6]. Although the statistical error in the measured cross section for this process at the Tevatron will be about $\pm 30\%$ [5], a high-luminosity Tevatron would allow a measurement of the cross section with a statistical uncertainty of about 6%[6]. At this level of experimental accuracy a calculation of the radiative corrections is necessary. In Ref.[6] the QCD and Yukawa corrections to single top quark production $q\bar{q}' \rightarrow t\bar{b}$ have been calculated in the SM. While the QCD corrections were found to be quite large, the Yukawa corrections were found to be negligible. Since the SM weak corrections are expected to be comparable to the Yukawa corrections, they too should be negligible. Beyond the SM, the Yukawa corrections might be greatly enhanced, since more Higgs bosons with stronger couplings to top or bottom quarks are involved in some new physics models. Once the top quark mass is known precisely, these effects could be used as an indirect test for new physics beyond the SM; for example, the two-Higgs-doublet model(2HDM) and the minimal supersymmetric model(MSSM)[7]. At least, the data could be used to place restrictions on these models. Therefore, it is worthwhile to investigate single top quark production via $q\bar{q}' \rightarrow t\bar{b}$ in these models. In this paper we present
the calculation of the Yukawa corrections of order $\alpha_{ew}M_t^2/M_W^2$ to single top production at the Fermilab Tevatron in both the 2HDM and the MSSM. These corrections arise from the virtual effects of the third family (top and bottom) of quarks, neutral and charged Higgs bosons, and neutral and charged Goldstone bosons. We note that our calculations can be easily extended to the pseudo-Goldstone boson (PGB) corrections in technicolor models\cite{8} by substituting the virtual PGB’s in the technicolor models for the virtual Higgs bosons in the 2HDM and MSSM.

2. Calculations

The tree-level Feynman diagram for single top quark production via $q\bar{q}' \rightarrow t\bar{b}$ is shown in Fig.1(a). The Yukawa corrections of order $\alpha_{ew}M_t^2/M_W^2$ to the process $q\bar{q}' \rightarrow t\bar{b}$ arise from the Feynman diagrams shown in Figs.1b-1h. In our calculations, we used dimensional regularization to control all the ultraviolet divergences in the virtual loop corrections and we adopted the on-mass-shell renormalization scheme\cite{9}. We also kept all the terms proportional to the product $M_b\tan \beta$ in the charged Higgs couplings to the third family of quarks since these effects may become rather important for large $\tan \beta$. We used the ’t Hooft-Feynman gauge for the propagators of virtual $W$ boson and Goldstone bosons. Including the $O(\alpha_{ew}M_t^2/M_W^2)$ Yukawa corrections, the renormalized amplitude for $q\bar{q}' \rightarrow t\bar{b}$ can be written as

$$M_{\text{ren}} = M_0 + \delta M_{\text{vertex}}$$

where $M_0$ is the tree-level matrix element, and $\delta M_{\text{vertex}}$ represents the $O(\alpha_{ew}M_t^2/M_W^2)$ Yukawa corrections arising from the self-energy diagrams Figs.1b-1e and vertex diagrams Figs.1f-1h. These are given by

$$M_0 = i \frac{g^2}{2} \frac{1}{\bar{s} - M_W^2} \bar{u}(p_2)\gamma_\mu P_L u(p_1)\bar{u}(p_3)\gamma^\mu P_L v(p_4)$$

and

$$\delta M_{\text{vertex}} = i \frac{g^2}{2} \frac{1}{\bar{s} - M_W^2} \bar{u}(p_2)\gamma_\mu P_L u(p_1)\bar{u}(p_3)[\gamma^\mu P_L \delta_1 + \gamma^\mu P_R \delta_2 + \bar{\Lambda}_L^\mu P_L + \bar{\Lambda}_R^\mu P_R] v(p_4)$$

where

$$\delta_1 = \left(\frac{1}{2} \delta Z_L^\mu + \frac{1}{2} \delta Z_R^\mu\right)_{\text{finite}} + f_1^L,$$

$$\delta_2 = f_1^R,$$

$$\bar{\Lambda}_L^\mu = f_2^L + f_3^L,$$

$$\bar{\Lambda}_R^\mu = f_2^R + f_3^R.$$
Here \( p_1 \) and \( p_2 \) denote the momentum of the incoming quarks \( q \) and \( q' \), while \( p_3 \) and \( p_4 \) are used for the outgoing \( t \) and \( b \) quarks, and \( \hat{s} \) is the center-of-mass energy of the subprocess. \( \delta Z_L^t \) and \( \delta Z_L^b \) are the wave-function renormalization constants, and \( f_i^{L,R} \) are form factors which are presented in Appendix A.

The renormalized differential cross section of the subprocess is

\[
\frac{d\hat{\sigma}}{d \cos \theta} = \frac{\hat{s} - M_t^2}{32\pi \hat{s}^2} \sum |M_{ren}|^2,
\]

(8)

where \( \theta \) is the angle between the top quark and incoming quark. Integrating this subprocess differential cross section over \( \cos \theta \) one finds

\[
\hat{\sigma} = \hat{\sigma}_0 + \Delta \hat{\sigma}
\]

(9)

where

\[
\hat{\sigma}_0 = \frac{g^4}{128\pi} \frac{\hat{s} - M_t^2}{\hat{s} - M_W^2} \left\{ \frac{2}{3}(\hat{s} - M_t^2)^2 + (\hat{s} - M_t^2)(M_t^2 + M_b^2) + 2M_t^2 M_b^2 \right\}
\]

(10)

is the tree-level result and the correction is

\[
\Delta \hat{\sigma} = \frac{g^4}{64\pi} \frac{\hat{s} - M_t^2}{\hat{s} - M_W^2} \left\{ \frac{1}{6}(\hat{s} - M_t^2)^2 + (\hat{s} - M_t^2)(M_t^2 + M_b^2) M_b^2 \right\} + \frac{1}{2} \hat{s} (\hat{s} - M_t^2)(M_b f_L^R - M_t f_L^B) + \frac{1}{2} \hat{s} (\hat{s} - M_t^2)(M_t f_L^R - M_b f_L^B) + \hat{s} M_t M_b (M_t f_L^R - M_b f_L^B) + \frac{1}{2} \hat{s} M_t M_b \delta_2 \right\}.
\]

(11)

The total hadronic cross section for the production of single-top-quark via \( qq' \) can be written in the form

\[
\sigma(s) = \sum_{i,j} \int dx_1 dx_2 \hat{\sigma}_{ij}(x_1 x_2 s, M_t^2, \mu^2) [f_i^A(x_1, \mu) f_j^B(x_2, \mu) + (A \leftrightarrow B)],
\]

(12)

where

\[
s = (P_1 + P_2)^2,
\]

(13)

\[
\hat{s} = x_1 x_2 s,
\]

(14)

\[
p_1 = x_1 P_1,
\]

(15)

and

\[
p_2 = x_2 P_2.
\]

(16)
Here $A$ and $B$ denote the incident hadrons and $P_1$ and $P_2$ are their four-momenta, while $i, j$ are the initial partons and $x_1$ and $x_2$ are their longitudinal momentum fractions. The functions $f^A_i$ and $f^B_j$ are the usual parton distributions. In our numerical calculations, we have used the CTEQ3L parton distribution functions for the tree level cross section, and CTEQ3M parton distribution functions[10] for the $O(\alpha_{\text{ew}} M_t^2/M_W^2)$ Yukawa corrections, as in ref.[6], to facilitate comparison. There is no Yukawa correction to parton distribution functions as pointed out in Ref.[11]. We will also compare our calculations with those calculated using the MRSG parton distribution functions[12] below. Finally, introducing the convenient variable $\tau = x_1 x_2$, and changing independent variables, the total cross section becomes

$$\sigma(s) = \sum_{i,j} \int \frac{d\tau}{\tau} \left( \frac{1}{s} \frac{dL_{ij}}{d\tau} \right)(\hat{s}\hat{\sigma}_{ij})$$

(17)

where $\tau_0 = (M_t + M_b)^2/s$. The quantity $dL_{ij}/d\tau$ is the parton luminosity, which is defined to be

$$\frac{dL_{ij}}{d\tau} = \int^1 \frac{dx_1}{x_1} [f^A_i(x_1, \mu)f^B_j(\tau/x_1, \mu) + (A \leftrightarrow B)]$$

(18)

3. Numerical results and discussions

In the following we present some numerical results for the Yukawa corrections to the total cross section for single top quark production via $q\bar{q}' \rightarrow t\bar{b}$ at the Fermilab Tevatron with $\sqrt{s} = 2$ TeV. In our numerical calculations we chose $M_Z = 91.188\,\text{GeV}$, $M_W = 80.33\,\text{GeV}$, $M_b = 5\,\text{GeV}$, $\alpha_{\text{ew}} = 1/128$, and $\mu = \sqrt{s}$. The Higgs masses $M_H, M_h, M_A, M_{H^+}$ and parameters $\alpha, \beta$ are not constrained in the general Two-Higgs-Doublet Model, but in the Minimal SUSY Model, relations[13] among these parameters are required by supersymmetry, leaving only two parameters free; e.g., $M_{H^+}$ and $\tan \beta$. Also, in the MSSM the charged Higgs mass is heavier than the $W$ mass due to the relation $M_{H^+}^2 = M_W^2 + M_A^2$. The experimental lower limit for the charged Higgs mass is 44.1GeV[14], independent of the additional parameters $\alpha$ and $\beta$. In our calculations we will use $M_{H^+} = 600\,\text{GeV}$ for 2HDM and $M_{H^+} > 100\,\text{GeV}$ for MSSM. The upper bound of $\tan \beta$; viz, $\tan \beta < 0.52\,\text{GeV}^{-1}M_{H^+}$, has been determined from data on $B \rightarrow \tau\nu X$[15]. The lower limits on $\tan \beta$ are $\tan \beta > 0.6$ from perturbative bounds [16] and $\tan \beta > 0.25$ (for $M_t = 175\,\text{GeV}$) from perturbative unitarity[16]. We will, therefore, limit the value of $\tan \beta$ to be in the range of 0.25 to 30.
Figure 2 shows the relative correction $\Delta \sigma/\sigma_0$ as a function of $M_H$ using the CTEQ3L parton distributions for the tree-level cross section $\sigma_0$ and the CTEQ3M parton distributions for the correction $\Delta \sigma$, as in ref.[6]. The solid curve corresponds to the 2HDM assuming $M_h = M_H$ and $M_t = 175\text{GeV}, M_{H^+} = M_A = 600\text{GeV}$ and $\tan \beta = 0.25$. The dotted curve corresponds to the SM for $M_t = 175$ GeV. For $M_h = M_H$ the corrections are independent of $\alpha$ in the 2HDM. The corrections in the SM are negligibly small, in agreement with Ref.[6]. However, in the 2HDM, the corrections can reduce the cross section by more than $-10\%$ for $M_h = M_H < 100\text{GeV}$, and for $M_h = M_H = 50\text{GeV}$ they can be as large as $-20\%$.

Figure 3 shows the relative correction $\Delta \sigma/\sigma_0$ in the MSSM, assuming $\tan \beta = 0.25$. Since the corrections are not sensitive to $M_{H^+}$, for $M_{H^+} > 400\text{GeV}$, we only present the results for $M_{H^+}$ in the range $100\text{GeV}$ to $400\text{GeV}$. In Fig.3 the solid curve corresponds to $M_t = 175$ GeV, again using CTEQ3L distributions for $\sigma_0$ and CTEQ3M distributions for $\Delta \sigma$. The dotted curve corresponds to $M_t = 175$ GeV but using the MRSG parton distributions. The dashed curve corresponds to $M_t = 200$ GeV and using CTEQ3L for $\sigma_0$ and CTEQ3M for $\Delta \sigma$. For a light charged Higgs, the corrections can be quite significant. For $M_{H^+} = 100\text{GeV}$ the corrections reach 9% for $M_t = 175\text{GeV}$ and 13% for $M_t = 200\text{GeV}$. The curve for $M_t = 175\text{GeV}$ has a peak at $M_{H^+} = 170\text{GeV}$ and the curve for $M_t = 200\text{GeV}$ has a peak at $M_{H^+} = 195\text{GeV}$, which is due to the fact that $M_b = 5\text{GeV}$ and the threshold for open top decay into a charged Higgs plus a bottom is crossed in these regions. If we change the top quark mass, we found that this region is also shifted correspondingly, which provides a check on our calculations, especially of the treatment of the threshold. From Fig.3 we see that the difference between the results obtained using CTEQ3 distributions and using MRSG distributions is negligibly small. We also found that the results using MRS(A') distributions [12] are almost the same as the MRSG results, and thus we did not present them.

In both Fig.2 and Fig.3, we used the minimal value (0.25) for $\tan \beta$. When $\tan \beta$ becomes larger, the corrections may drop rapidly since the dominant effects arise from the terms $\sim \alpha_{ew} \frac{M_b^2}{M_{W^\pm} \tan \beta}$. In Fig.4 we present the dependence of the relative correction, $\Delta \sigma/\sigma_0$, on the value of $\tan \beta$ using CTEQ3L for $\sigma_0$ and CTEQ3M for $\Delta \sigma$. The solid curve corresponds to the 2HDM assuming $M_t = 175\text{GeV}, M_{H^+} = M_A = 600\text{GeV}$ and $M_H = M_h = 65\text{GeV}$. The dotted curve corresponds to the MSSM assuming $M_t = 175\text{GeV}$ and $M_{H^+} = 100\text{GeV}$. The corrections are only significant for small $\tan \beta$ and are quite sensitive to $\tan \beta$ for $\tan \beta < 5$.

Since the cross section for single top production can be reliably predicted in the SM [6] and the statistical error in the measurement of the cross section will be about 6% at a
high-luminosity Tevatron[6], these corrections may be observable; at least, interesting new constraints on these models can be established.

Note that in the MSSM, besides these Yukawa corrections arising from the Higgs sector, the supersymmetric (SUSY) corrections due to super particles (sparticles) should also be taken into account[17]. The dominant virtual effects of sparticles arise from supersymmetric QCD corrections of order $\alpha_s$ and the supersymmetric electroweak correction of order $\alpha_{ew} M_t^2 / M_W^2$ which arise from loops of charginos and neutralinos, the supersymmetry partners of Higgs and vector bosons. However, the anomalous magnetic moment for a spin 1/2 fermion vanishes in the SUSY limit[18] and away from the SUSY limit the cancellations have somewhat less effect. Therefore, in general one can expect the Yukawa corrections from the Higgs sector and the supersymmetric electroweak corrections from virtual charginos and neutralinos to cancel to some extent.

In conclusion, we calculated the Yukawa corrections of order $\alpha_{ew} M_t^2 / M_W^2$ to single top quark production via $q\bar{q} \rightarrow t\bar{b}$ at the Fermilab Tevatron in the general two-Higgs-doublet model and the minimal supersymmetric model. For favorable parameter values, the corrections can reduce the cross section by more than 15% in the general two-Higgs-doublet model and enhance the cross section by up to 10% in the minimal supersymmetric model. These effects may be observable at a high-luminosity Tevatron.

This work was supported in part by the U.S. Department of Energy, Division of High Energy Physics, under Grant No. DE-FG02-91-ER4086.

Appendix A

\begin{align}
\delta Z_t^L &= \frac{\alpha_{ew}}{16\pi M_W^2 s_W^2} \left\{ \sum_{i=H,h} M_t^2 \eta_i^2 \left[ -\frac{\Delta}{2} + F_1^{(tti)} + 2 M_t^2 (G_0^{(tti)} + G_1^{(tti)}) \right] \\
+ \sum_{i=A,G^0} M_t^2 \eta_i^2 \left[ -\frac{\Delta}{2} + F_1^{(tti)} + 2 M_t^2 (G_1^{(tti)} - G_0^{(tti)}) \right] \\
+ \sum_{i=H^+,G^+} \left[ 2 M_b^2 \lambda_i^2 \left( -\frac{\Delta}{2} + F_1^{(bbi)} \right) + 2 M_t^2 (M_t^2 \eta_i^2 + M_b^2 \lambda_i^2) G_1^{(bbi)} \right] \right\} \\
\delta Z_b^L &= \frac{\alpha_{ew}}{16\pi M_W^2 s_W^2} \left\{ \sum_{i=H,h} M_b^2 \lambda_i^2 \left[ -\frac{\Delta}{2} + F_1^{(bbi)} + 2 M_b^2 (G_0^{(bbi)} + G_1^{(bbi)}) \right] \\
+ \sum_{i=H^-} M_b^2 \lambda_i^2 \left[ -\frac{\Delta}{2} + F_1^{(b\bar{b}i)} + 2 M_b^2 (M_b^2 \eta_i^2 + M_t^2 \lambda_i^2) G_1^{(b\bar{b}i)} \right] \right\} \\ 
(A.1)
\end{align}
\[ + \sum_{i=A,G^0} M_i^2 \Delta_i^2 \left[ -\frac{\Delta}{2} + F_1^{(b)} + 2M_i^2 (G_1^{(b)} - G_0^{(b)}) \right] \\
+ \sum_{i=H^+,G^+} [2M_i^2 \eta_i^2 \left( -\frac{\Delta}{2} + F_1^{(b)} \right) + 4M_i^2 M_i^2 \eta_i \lambda_i G_0^{(b)} \]
\[ + 2M_i^2 (M_i^2 \eta_i^2 + M_i^2 \lambda_i) G_1^{(b)} \right] \} \] (A.2)

where \( \Delta \equiv \frac{1}{\epsilon} - \gamma_E + \log 4\pi \) with \( \gamma_E \) being the Euler constant and \( D = 4 - 2\epsilon \) is the space-time dimension. The functions \( F_n^{(ijk)} \), \( G_n^{(ijk)} \) are defined by

\[ F_n^{(ijk)} = \int_0^1 dy \eta^n \log \left[ \frac{M_i^2 y(y-1) + M_j^2 (1-y) + M_k^2 y}{\mu^2} \right], \] (A.3)

\[ G_n^{(ijk)} = - \int_0^1 dy M_i^2 y(y-1) + M_j^2 (1-y) + M_k^2 y. \] (A.4)

The form factors \( f_{ij}^{L,R} \) are obtained by

\[ f_{ij}^{L,R} = \frac{\alpha_{ew}}{16\pi M_W^2 s_W^2} \sum_{j=1}^6 f_{ij}^{L,R} \] (A.5)

\( f_{ij}^{L,R} \) are given by

\[ f_{11}^{L} = 2M_b^2 \xi_{ij} \lambda_j \lambda_i \bar{c}_{24} \] (A.6)

\[ f_{21}^{L} = M_b M_b^2 \xi_{ij} \lambda_j [\eta_i (2c_{23} + c_{12} - 2c_{22}) + \lambda_i (c_{12} + 2c_{22} + 2c_{21} + c_{11} - 4c_{23})] \] (A.7)

\[ f_{31}^{L} = M_b M_b^2 \xi_{ij} \lambda_j [\eta_i (-c_{12} - 2c_{22}) + \lambda_i (-c_{11} + 2c_{22} - 2c_{23} + c_{12})] \] (A.8)

\[ f_{11}^{R} = 2M_t M_b \xi_{ij} \lambda_j \eta_i \bar{c}_{24} \] (A.9)

\[ f_{21}^{R} = M_b \xi_{ij} \lambda_j [M_b^2 \lambda_i (2c_{23} + c_{12} - 2c_{22}) \]
\[ + M_b^2 \eta_i (-c_{12} + 2c_{22} + 2c_{21} + c_{11} - 4c_{23})] \] (A.10)

\[ f_{31}^{R} = M_b \xi_{ij} \lambda_j [M_b^2 \lambda_i (-c_{12} - 2c_{22}) + M_b^2 \eta_i (-c_{11} + 2c_{22} - 2c_{23} + c_{12})] \] (A.11)

\[ f_{12}^{L} = 2M_b^2 \xi_{ij} \lambda_i \lambda_j \bar{c}_{24} \] (A.12)

\[ f_{22}^{L} = -M_t M_b^2 \xi_{ij} \lambda_j [\eta_i (2c_{23} + c_{12} - 2c_{22}) \]
\[ - \lambda_i (-c_{12} + 2c_{22} + 2c_{21} + c_{11} - 4c_{23})] \] (A.13)

\[ f_{32}^{L} = -M_t M_b^2 \xi_{ij} \lambda_j [\eta_i (-c_{12} - 2c_{22}) - \lambda_i (-c_{11} + 2c_{22} - 2c_{23} + c_{12})] \] (A.14)

\[ f_{12}^{R} = -2M_t M_b \xi_{ij} \lambda_j \eta_i \bar{c}_{24} \] (A.15)

\[ f_{22}^{R} = M_b \xi_{ij} \lambda_j [M_b^2 \lambda_i (2c_{23} + c_{12} - 2c_{22}) \]
\[ - M_b^2 \eta_i (-c_{12} + 2c_{22} + 2c_{21} + c_{11} - 4c_{23})] \] (A.16)
\[ f_{32}^R = M_b \xi_{ij} \lambda_j [M_b^2 \lambda_i (-c_{12} - 2c_{22}) - M_b^2 \eta_i (-c_{11} + 2c_{22} - 2c_{23} + c_{12})] \] (A.17)

\[ f_{13}^L = -2M_t^2 \xi_{ij} \eta_{ij} \eta_i \bar{c}_{24} \] (A.18)

\[ f_{23}^L = -M_t \xi_{ij} \eta_{ij} [M_t^2 \eta_i (2c_{21} + 2c_{22} - 4c_{23} + c_{12} - c_{11} - c_0)] + M_t^2 \lambda_i (2c_{23} - 2c_{22} + c_{12}) \] (A.19)

\[ f_{33}^L = -M_t \xi_{ij} \eta_{ij} [M_t^2 \eta_i (2c_{22} - 2c_{23} + 3c_{12} - c_{11} + c_0) - M_b^2 \lambda_i (2c_{22} + c_{12})] \] (A.20)

\[ f_{13}^R = -2M_t M_b \xi_{ij} \eta_{ij} \lambda_i \bar{c}_{24} \] (A.21)

\[ f_{23}^R = -M_t^2 M_b \xi_{ij} \eta_{ij} [\eta_i (2c_{23} - 2c_{22} + c_{12}) + \lambda_i (2c_{21} + 2c_{22} - 4c_{23} + c_{12} - c_{11} - c_0)] \] (A.22)

\[ f_{33}^R = -M_t^2 M_b \xi_{ij} \eta_{ij} [-\eta_i (2c_{22} + c_{12}) + \lambda_i (2c_{22} - 2c_{23} + 3c_{12} - c_{11} + c_0)] \] (A.23)

\[ f_{14}^L = 2M_t^2 \xi_{ij} \eta_{ij} \eta_i \bar{c}_{24} \] (A.24)

\[ f_{24}^L = -M_t \xi_{ij} \eta_{ij} (M_t^2 \eta_i (4c_{23} - 2c_{21} - 2c_{22} + 3c_{12} - 3c_{11} - c_0) + M_t^2 \lambda_i (2c_{23} - 2c_{22} + c_{12})] \] (A.25)

\[ f_{34}^L = -M_t \xi_{ij} \eta_{ij} [M_t^2 \eta_i (2c_{23} - 2c_{22} + c_{12} + c_{11} + c_0) - M_b^2 \lambda_i (2c_{22} + c_{12})] \] (A.26)

\[ f_{14}^R = -2M_t M_b \xi_{ij} \eta_{ij} \lambda_i \bar{c}_{24} \] (A.27)

\[ f_{24}^R = -M_t^2 M_b \xi_{ij} \eta_{ij} [\eta_i (2c_{22} - 2c_{23} - c_{12}) + \lambda_i (2c_{21} + 2c_{22} - 4c_{23} + 3c_{11} - 3c_{12} + c_0)] \] (A.28)

\[ f_{34}^R = -M_t^2 M_b \xi_{ij} \eta_{ij} [\eta_i (2c_{22} + c_{12}) + \lambda_i (2c_{22} - 2c_{23} - c_{12} - c_{11} - c_0)] \] (A.29)

\[ f_{15}^L = -M_t^2 M_b^2 \eta_i \lambda_j (c_0 + 2c_{12} - c_{11}) \] (A.30)

\[ f_{25}^L = 2M_t M_b^2 \eta_i \lambda_j (c_{12} + c_{23}) \] (A.31)

\[ f_{35}^L = 2M_t M_b^2 \eta_i \lambda_j (c_{12} + c_{23}) \] (A.32)

\[ f_{15}^R = M_t M_b \eta_i \lambda_j [2\bar{c}_{24} + \tilde{\delta}(c_{12} + c_{23}) - M_t^2 (c_0 - c_{21} + c_{12} + c_{23})] \] (A.33)

\[ f_{25}^R = 2M_t^2 M_b \eta_i \lambda_j (c_0 - c_{21}) \] (A.34)

\[ f_{35}^R = 2M_t^2 M_b \eta_i \lambda_j (c_{12} + c_{23}) \] (A.35)

\[ f_{16}^L = M_t^2 M_b^2 (c_0 + c_{11}) \] (A.36)

\[ f_{26}^L = 2M_t M_b^2 (c_{12} + c_{23}) \] (A.37)

\[ f_{36}^L = 2M_t M_b^2 (c_{12} + c_{22}) \] (A.38)

\[ f_{16}^R = M_t M_b [2\bar{c}_{24} + \tilde{\delta}(c_{12} + c_{23}) - M_t^2 (c_{12} + c_{23} - c_{21} - 2c_{11} - c_0)] \] (A.39)

\[ f_{26}^R = -2M_t^2 M_b (c_0 + 2c_{11} + c_{21}) \] (A.40)

\[ f_{36}^R = 2M_t^2 M_b (c_{12} + c_{23}) \] (A.41)
where the sums over \(i = H^+, G^+, j = H, h\) for \(f_{i1}^{L,R}\) and \(f_{i3}^{L,R}\), \(i = H^+, G^+, j = A, G^0\) for \(f_{i2}^{L,R}\) and \(f_{i4}^{L,R}\), and \(j = H, h\) for \(f_{i5}^{L,R}\) are implied. The functions \(c_{ij}\) defined as

\[
c_{ij} = c_{ij}(-p_3, k, M_b, M_i, M_j) \quad \text{for} \quad f_{i1}^{L,R}, f_{i2}^{L,R}
\]
\[
c_{ij} = c_{ij}(-p_3, k, M_t, M_j, M_i) \quad \text{for} \quad f_{i3}^{L,R}, f_{i4}^{L,R}
\]
\[
c_{ij} = c_{ij}(-p_3, -p_4, M_t, M_j, M_b) \quad \text{for} \quad f_{i5}^{L,R}
\]
\[
c_{ij} = c_{ij}(-p_3, -p_4, M_t, M_A, M_b) \quad \text{for} \quad f_{i6}^{L,R}
\]

are the three-point Feynman integrals\[19\]. The constants \(\eta, \lambda, \xi\) are defined as

\[
\eta_{H^+} = \eta_A = \cot \beta, \quad \eta_{G^+} = \eta_{G^0} = 1, \quad (A.43)
\]
\[
\eta_H = \sin \alpha / \sin \beta, \quad \eta_h = \cos \alpha / \sin \beta, \quad (A.44)
\]
\[
\lambda_{H^+} = \lambda_A = \tan \beta, \quad \lambda_{G^+} = \lambda_{G^0} = 0, \quad (A.45)
\]
\[
\lambda_H = \cos \alpha / \cos \beta, \quad \lambda_h = -\sin \alpha / \cos \beta, \quad (A.46)
\]
\[
\xi_{H+H} = -\xi_{G+h} = \sin(\beta - \alpha), \quad (A.47)
\]
\[
\xi_{H+h} = \xi_{G+H} = -\cos(\beta - \alpha), \quad (A.48)
\]
\[
\xi_{H+A} = \xi_{G+G^0} = 1, \quad \xi_{H+G^0} = \xi_{G+H} = 0 \quad (A.49)
\]
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Figure Captions

Fig.1 Feynman diagrams: (a) tree-level, (b)-(e) self-energies, (f)-(h) vertex corrections.

Fig.2 The relative correction $\Delta \sigma/\sigma_0$ as a function of $M_H$ using the CTEQ3L parton distributions for $\sigma_0$ and the CTEQ3M parton distributions for $\Delta \sigma$. The solid curve corresponds to the 2HDM assuming $M_h = M_H, M_t = 175$ GeV, $M_{H^+} = M_A = 600$GeV and $\tan \beta = 0.25$. The dotted curve corresponds to the SM for $M_t = 175$ GeV.

Fig.3 The relative correction $\Delta \sigma/\sigma_0$ as a function of $M_{H^+}$ in the MSSM, assuming $\tan \beta = 0.25$. The solid curve corresponds to $M_t = 175$ GeV using the CTEQ3L parton distributions for $\sigma_0$ and the CTEQ3M parton distributions for $\Delta \sigma$. The dotted curve corresponds to $M_t = 175$ GeV using the MRSG parton distributions. The dashed curve corresponds to $M_t = 200$ GeV using the CTEQ3L distributions for $\sigma_0$ and the CTEQ3M distributions for $\Delta \sigma$.

Fig.4 The relative correction $\Delta \sigma/\sigma_0$ as a function of $\tan \beta$ using the CTEQ3L distributions for $\sigma_0$ and the CTEQ3M distributions for $\Delta \sigma$. The solid curve corresponds to the 2HDM assuming $M_t = 175$GeV, $M_{H^+} = M_A = 600$GeV and $M_H = M_h = 65$GeV. The dotted curve corresponds to the MSSM assuming $M_t = 175$GeV and $M_{H^+} = 100$GeV.