

Yukawa correction to top-quark production at the Fermilab Tevatron

A. Stange

Physics Department, Brookhaven National Laboratory, Upton, New York 11973

S. Willenbrock

*Physics Department, Brookhaven National Laboratory, Upton, New York 11973
and Fermi National Accelerator Laboratory, P.O. Box 500, Batavia, Illinois 60510*

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We calculate the correction to $q\bar{q} \rightarrow t\bar{t}$ of order $g^2 m_t^2/M_W^2$. This correction, proportional to the square of the Higgs-boson Yukawa coupling to the top quark, arises from loops of Higgs bosons and the scalar component of virtual vector bosons. The Yukawa correction to the total $t\bar{t}$ production cross section at the Fermilab Tevatron in the standard Higgs model is found to be much less than the theoretical uncertainty in the cross section. However, in a two-Higgs-doublet model, Yukawa couplings are generally enhanced. The Yukawa correction can increase the total $t\bar{t}$ production cross section in this model by as much as 20–35%, which is potentially observable at the Tevatron.

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I. INTRODUCTION

In the standard electroweak model, the strength of the weak interaction is comparable to that of the electromagnetic interaction at energies of the order of the weak-vector-boson masses. However, the sector of the weak interaction which generates the fermion masses m_f is characterized by a coupling of strength gm_f/M_W , where g is the weak coupling. This is the strength of the Yukawa coupling of the Higgs boson to the fermions, as well as that of real longitudinal vector bosons and the scalar component of virtual vector bosons to the fermions. If the top-quark mass is much greater than the W -boson mass, this Yukawa-strength coupling is sufficiently large that it could produce noticeable effects in processes involving the top quark.

The generation of the fermion masses, as well as the weak-vector-boson masses, is one of the outstanding puzzles of the electroweak theory. The top-quark coupling to the Higgs boson is a remnant of the mechanism which generates the top-quark mass. The observation of effects produced by this coupling is therefore a window on the top-quark mass-generating mechanism. The top quark could provide us with the first clue towards solving the mystery of mass generation.

In the standard Higgs model, a single Higgs doublet provides masses to the fermions and the weak vector bosons. Consequently, Yukawa couplings are fixed in terms of the corresponding fermion mass and the W -boson mass. However, in a two-Higgs-doublet model, the Higgs bosons which come from the doublet which provides mass to a given fermion generally have enhanced Yukawa couplings to that fermion. This results in a more pronounced effect of the fermion mass-generating mechanism on physical processes.

In this paper we investigate the Yukawa-strength correction to top-quark production at the Fermilab Tevatron. With the Main Injector, it may be possible to

detect a top quark as heavy as 250 GeV at the Tevatron. The dominant production mechanism for a heavy top quark ($m_t > 150$ GeV) at the Tevatron is quark-antiquark annihilation [1]; we therefore restrict our attention to this process. The diagrams which contribute to this correction are shown in Fig. 1. The dashed lines represent the Higgs boson and the unphysical scalar bosons, in Landau gauge, associated with the vector bosons. The Yukawa correction to the top-quark cross section is proportional to $g^2 m_t^2/M_W^2$, and dominates the ordinary weak correction, of $O(g^2)$, in the large m_t limit.

In the Landau gauge, the vector-boson propagator is entirely spin one, and the scalar component of the virtual vector boson is represented by the massless unphysical scalar boson (Goldstone boson). In this gauge, the interactions of Yukawa strength are isolated in the Higgs-boson and unphysical-scalar-boson couplings to the fermions. Real (but not virtual) longitudinal vector bosons also effectively couple to fermions with Yukawa strength.

The Yukawa correction to $t\bar{t}$ production in the standard Higgs model is discussed in Sec. II. In Sec. III we discuss the correction in a two-Higgs-doublet model, both in general and in the minimal supersymmetric model. We present our conclusions in Sec. IV. Analytic expressions for the loop corrections are given in an Appendix.

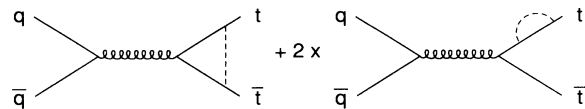


FIG. 1. Diagrams which contribute to the Yukawa correction to $q\bar{q} \rightarrow t\bar{t}$. The dashed lines represent the Higgs boson and the unphysical scalar Z and W bosons in Landau gauge.

II. STANDARD HIGGS MODEL

The Yukawa correction to the $q\bar{q} \rightarrow t\bar{t}$ amplitude (see Fig. 1) is contained in the correction to the matrix element of the top-quark color current. The general form for this matrix element, consistent with current conservation, is

$$i\bar{u}(p_3)\Gamma^{\mu A}v(p_4) = -ig_s \left[\bar{u}(p_3)T^A\gamma^\mu v(p_4) - \left(\frac{gm_t}{2M_W} \right)^2 \frac{1}{(4\pi)^2} \bar{u}(p_3)T^A [V\gamma^\mu + Ti\sigma^{\mu\nu}q_\nu + A(\gamma^\mu q^2 - 2m_t q^\mu)\gamma_5 + P\gamma_5\sigma^{\mu\nu}q_\nu]v(p_4) \right],$$

where p_3 and p_4 are the outgoing momenta of the top quark and top antiquark, $q = p_3 + p_4$, the form factors V , T , A , and P are functions of $q^2 = s$, T^A is an SU(3) matrix [$\text{Tr}(T^A T^B) = \frac{1}{2}\delta^{AB}$], and we have factored out the couplings and loop factors. The chromoelectric-dipole form factor (P) is CP violating, and vanishes at one loop. The chromocharge (V), chromomagnetic (T), and chromoanapole (A) form factors are given in an Appendix.

The differential cross section for $q\bar{q} \rightarrow t\bar{t}$, summed over final and averaged over initial colors and spins, is

$$\frac{d\sigma}{dz} = \frac{8\pi}{9s^3} \alpha_s^2 \beta \left[(p_1 \cdot p_3)^2 + (p_2 \cdot p_3)^2 + m_t^2 p_1 \cdot p_2 - \frac{g^2 m_t^2}{64\pi^2 M_W^2} \text{Re}\{2V[(p_1 \cdot p_3)^2 + (p_2 \cdot p_3)^2 + m_t^2 p_1 \cdot p_2] + m_t s^2 T\} \right],$$

where p_1 and p_2 are the incoming momenta of the quark and antiquark, z is the cosine of the scattering angle between the quark and the top quark, and $\beta = (1 - 4m_t^2/s)^{1/2}$. The parity-violating chromoelectric-dipole (P) and chromoanapole (A) form factors do not contribute to the unpolarized cross section at one loop.

In Fig. 2 we show the percentage change in the cross section, as a function of the $t\bar{t}$ invariant mass, for $m_t = 150$ GeV and $M_H = 60, 200,$ and 800 GeV. The contributions of the Higgs boson and the unphysical scalar Z and W bosons are shown separately. The Higgs-boson contribution near threshold is large and positive for a relatively light Higgs boson. This is due to the Yukawa potential formed between the top quark and top antiquark by the exchange of the Higgs boson, a phenomenon that has been studied in $e^+e^- \rightarrow t\bar{t}$ near threshold [2–4]. The exchange of the unphysical scalar Z and W bosons does not lead to a similar effect, in the first case because the interaction is pseudoscalar, in the second because the interaction changes the top quark to a bottom quark. Exchange of the unphysical scalar Z boson produces a potential which vanishes in the nonrelativistic limit, so its contribution vanishes at threshold. The unphysical-scalar- W -boson contribution is negative and almost constant over the range shown. The Higgs-boson contribution is also negative far above threshold.

To obtain the Yukawa correction to the top-quark production cross section at the Tevatron, one convolutes the subprocess cross section with parton distribution functions. For completeness, one must also consider corrections of order $g^2 m_t^2/M_W^2$ to the parton distribution functions. The parton distribution functions are extracted from deep-inelastic scattering, Drell-Yan, and direct photon production [5]. Terms of order $g^2 m_t^2/M_W^2$ can potentially arise from the top-quark contribution to weak-vector-boson vacuum polarization. However, when

charged-current processes are expressed in terms of G_μ , these terms are absorbed into the coupling, since the same vacuum-polarization diagram occurs in muon decay. Neutral-current processes do have corrections of $O(g^2 m_t^2/M_W^2)$ via the ρ parameter [6],

$$\rho = 1 + 3 \frac{g^2 m_t^2}{64\pi^2 M_W^2},$$

but presently there are no neutral-current processes from which information on the parton distribution functions

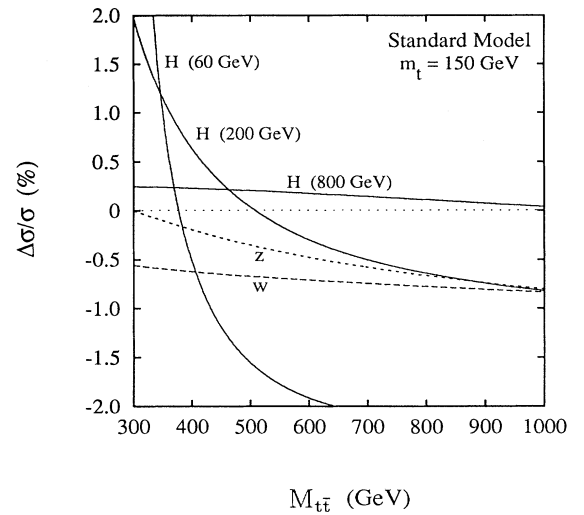


FIG. 2. Change in the cross section for $q\bar{q} \rightarrow t\bar{t}$, due to the Yukawa correction, as a function of the $t\bar{t}$ invariant mass, for $m_t = 150$ GeV. The contributions of the Higgs boson and the unphysical scalar Z and W bosons are shown separately.

is extracted. Thus there is no Yukawa correction to the parton distribution functions at this time.

To obtain the correction to the total top-quark production cross section at the Tevatron, we weight the correction to the subprocess cross section with the parton distribution functions and integrate over all $t\bar{t}$ invariant masses. The quark and antiquark distribution functions are evaluated at $x \sim 2m_t/\sqrt{S} \sim 0.1 - 0.2$, where they are well measured. We present our numerical results with the Morfin-Tung leading-order parton distribution functions [7]. The resulting change in the total cross section is shown in Fig. 3 for $m_t = 150, 200,$ and 250 GeV, as a function of the Higgs-boson mass. Because the parton distribution functions and the subprocess cross section fall off with increasing $t\bar{t}$ invariant mass, only the region within a few hundred GeV of threshold is numerically significant when we integrate over $M_{t\bar{t}}$. The positive contribution from the Higgs boson is largely compensated by the negative contribution from the unphysical scalar W boson, so the correction to the total cross section is much smaller than the typical correction at fixed invariant mass. The unphysical Z boson makes a negligible contribution. The contribution from a very light Higgs boson ($M_H < 100$ GeV) is reduced because the Yukawa enhancement is peaked very close to threshold, where the cross section is suppressed by phase space. The correction is largest for $M_H \approx 125$ GeV, but is at most $+2.4\%$ for $m_t < 250$ GeV.

The radiation of a real Higgs boson [8–10] or longitudinal vector boson [9,10] from a top quark is also a correction to top-quark production of Yukawa strength. We have calculated these processes and found that they are negligible at the Tevatron. The parton distribution functions are steeply falling at large masses, so the emission of an additional massive particle is suppressed.

The uncertainty in the top-quark total cross section at the Tevatron is due to the uncertainty in α_s and the

uncalculated next-to-next-to-leading order QCD correction [1,11]. The net uncertainty is estimated to be about $\pm 20\%$ [12], much greater than the Yukawa correction to the total cross section. We conclude that the Yukawa correction to the total cross section is unobservable at the Tevatron for $m_t < 250$ GeV.

Because the ordinary weak corrections to deep-inelastic scattering, Drell-Yan, and direct-photon production are not included in the extraction of the parton distribution functions, the ordinary weak correction to the $q\bar{q} \rightarrow t\bar{t}$ cross section does not represent a complete calculation of the ordinary weak correction to the hadronic $t\bar{t}$ cross section. The ordinary weak correction, as well as the Yukawa correction which we have evaluated, was calculated in Ref. [13], and numerical results were given for the correction to the $q\bar{q} \rightarrow t\bar{t}$ and $gg \rightarrow t\bar{t}$ cross sections. These results agree with ours in the region a few hundred GeV above threshold, although the threshold region is not shown in sufficient detail to allow a comparison. That study notes that there is a large negative correction to the cross sections at very large $t\bar{t}$ invariant masses. Since we do not find such a result from the Yukawa correction, it is presumably due to the ordinary weak correction, enhanced by a large infrared logarithm, $\ln^2 M_Z^2/s$. However, the large-invariant-mass region makes a negligible contribution to the total hadronic cross section. We question the estimate of a negative 10–20% correction to the $t\bar{t}$ production cross section at the CERN Large Hadron Collider (LHC) for $m_t = 200$ GeV quoted in that work. At the LHC and the Superconducting Super Collider (SSC) $gg \rightarrow t\bar{t}$ is the dominant top-quark production mechanism, so we cannot give a result for the Yukawa correction at these machines.

The W -gluon fusion process, $Wg \rightarrow t\bar{b}$, is also a copious source of top quarks at the Tevatron, especially for large m_t [14]. However, once cuts are made in an effort to identify the signal, the rate for top-quark production via W -gluon fusion falls below that of $q\bar{q} \rightarrow t\bar{t}$ [15,16].

The full weak correction to $e^+e^- \rightarrow t\bar{t}$, including the Yukawa correction, is given in Ref. [17]. The correction is comparable to the correction found here for fixed $t\bar{t}$ invariant mass. The correction in the threshold region is studied in Ref. [18].

III. TWO-HIGGS-DOUBLET MODEL

The simplest extension of the Higgs sector of the standard electroweak model that does not conflict with experiment is a two-Higgs-doublet model [19]. The tree-level relation $\rho = M_W^2/(M_Z^2 \cos^2 \theta_W) = 1$ is satisfied automatically, as in the one-doublet model [20]. In order to avoid Higgs-boson-mediated tree-level flavor-changing neutral currents, it is necessary to require that fermions of a given electric charge receive their mass from only one of the Higgs doublets [21]. Since $M_W^2 = \frac{1}{4}g^2(v_1^2 + v_2^2)$, both v_1 and v_2 must be less than v , the vacuum-expectation value of the one-doublet model. The Yukawa couplings of a given fermion to the Higgs scalars which come from the doublet ϕ_i that provides mass to that fermion are proportional to m_f/v_i , and are therefore naturally enhanced.

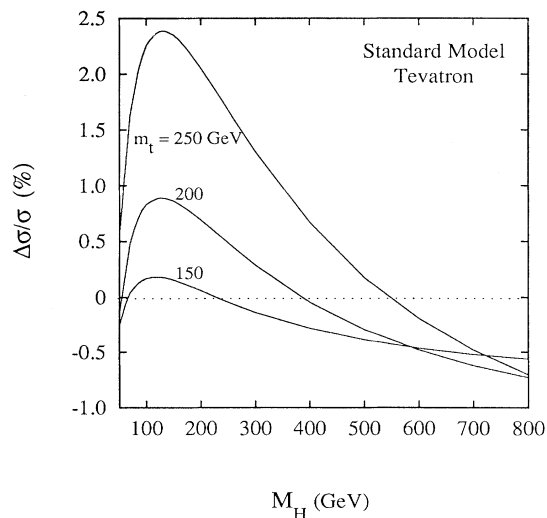


FIG. 3. Change in the total cross section, due to the Yukawa correction, for $p\bar{p} \rightarrow t\bar{t} + X$ at the Tevatron, vs the Higgs-boson mass.

Thus one may obtain an enhanced Yukawa correction to top-quark production in a two-Higgs-doublet model.

The two-Higgs-doublet model has five physical Higgs bosons, in addition to the unphysical scalar Z and W bosons. There are two scalars h and H , a pseudoscalar A , and two charged scalars H^\pm . With regard to their coupling to the top quark, the pseudoscalar and charged scalars are massive versions of the unphysical scalar Z and W bosons, respectively, but with the coupling factors given in Table I. It is conventional to choose ϕ_2 to give mass to the top quark, and to define $\tan\beta = v_2/v_1$; the top-quark Yukawa couplings are therefore enhanced for small β . The two scalar Higgs bosons have the same quantum numbers, and mix with a mixing angle α . The couplings of the unphysical scalar Z and W bosons are unchanged from the one-doublet model.

If β is sufficiently small, processes mediated by charged Higgs bosons may be enhanced such that they conflict with experiment. The strongest constraint appears to come from the lack of observation of $b \rightarrow s\gamma$ [22,23]. For small β , the upper limit on $b \rightarrow s\gamma$ requires a charged-Higgs-boson mass greater than several hundred GeV, depending on the model and the top-quark mass [23]. To avoid a large correction to the ρ parameter we must set $M_A \approx M_{H^\pm}$ [24], so the pseudoscalar Higgs boson must also be heavy.

A sufficiently small value of β also yields a Yukawa coupling which is so strong that perturbation theory is unreliable. Consider the zeroth partial wave of the $t\bar{t}$ elastic scattering amplitude, keeping only the terms which are enhanced for small β [25]:

$$a_0 = -\frac{3G_F m_t^2 s}{8\pi\sqrt{2}\sin^2\beta} \left[\frac{\cos^2\alpha}{s - M_h^2} + \frac{\sin^2\alpha}{s - M_H^2} + \frac{\cos^2\beta}{s - M_A^2} \right].$$

Applying the unitarity condition $|\text{Re } a_0| < 1/2$ in the energy regime $M_{h,H}^2 \ll s \ll M_A^2$ implies [26]

$$\sin^2\beta > \frac{3G_F m_t^2}{4\pi\sqrt{2}}.$$

This gives $\beta > 0.21, 0.28, 0.36$ for $m_t = 150, 200, 250$ GeV, respectively. Although β may be less than these values in principle, we believe that they correspond to the largest Yukawa correction that is physically allowed.

In order to ascertain the maximum Yukawa correction to top-quark production at the Tevatron that one can expect in a two-Higgs-doublet model, we set β to its minimum value and set $M_{H^\pm} = M_A = 600$ GeV. The contributions of the Higgs bosons to the form factors are given

TABLE I. Factors associated with the top-quark Yukawa couplings in a two-Higgs-doublet model.

$ht\bar{t}$	$\frac{\cos\alpha}{\sin\beta}$
$Ht\bar{t}$	$\frac{\sin\alpha}{\sin\beta}$
$At\bar{t}$	$\frac{1}{\tan\beta}$
$H^\pm t\bar{b}$	$\frac{1}{\tan\beta}$

by the same expressions as in the one-doublet model, but with $M_H \rightarrow M_{h,H}$, $M_Z \rightarrow M_A$, and $M_W \rightarrow M_{H^\pm}$, and multiplied by the square of the associated coupling factor in Table I. Since the couplings of the unphysical scalar bosons are not enhanced, their contribution is negligible. As with the unphysical scalar Z boson in the one-doublet model, the pseudoscalar-Higgs-boson contribution is negligible, for all M_A . Recall that in the one-doublet model the positive contribution of the Higgs boson was largely compensated by the negative contribution of the unphysical scalar W boson. For small β , the charged Higgs boson plays the role of the unphysical scalar W boson; however, since it must be heavy, its (negative) contribution is suppressed. Thus, as long as both Higgs scalars are relatively light, or there is one light scalar whose coupling to the top quark is not suppressed by the mixing angle α , we can expect a large, positive Yukawa correction in the two-Higgs-doublet model for small β .

We show in Fig. 4 the maximum correction to the $t\bar{t}$ total cross section at the Tevatron, for $M_h = M_H$, as a function of the common scalar-Higgs-boson mass. The mixing angle α drops out for equal scalar masses. The correction is increased considerably over the standard-model correction, as large as +23%, +28%, +35% for $m_t = 150, 200, 250$ GeV. These corrections to the total cross section are potentially observable at the Tevatron, and could be used to set limits on the parameters of a two-Higgs-doublet model.

The Higgs sector of the minimal supersymmetric standard model is a special case of a two-Higgs-doublet model, and is specified by just two parameters (at the tree level), which we can take to be M_{H^\pm} and β [19]. For large M_{H^\pm} , $\alpha \rightarrow \beta - \frac{\pi}{2}$, so that for small β the H Yukawa coupling is enhanced while that of h approaches the one-doublet-model value. For large M_{H^\pm} , H is also heavy (it is nearly degenerate in mass with the charged Higgs

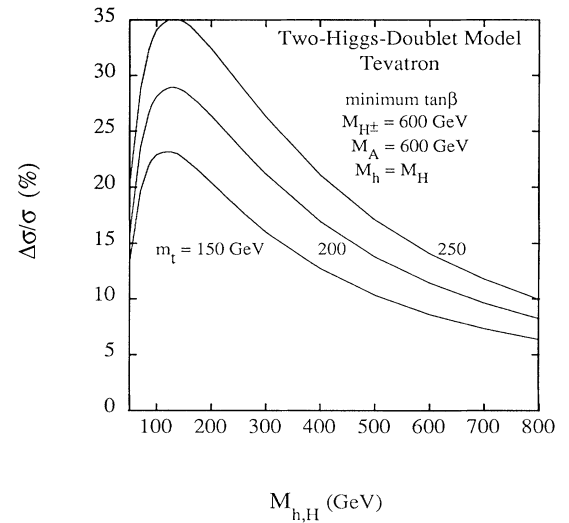


FIG. 4. Maximum change in the total cross section, due to the Yukawa correction in a two-Higgs-doublet model, for $p\bar{p} \rightarrow t\bar{t} + X$ at the Tevatron, vs the common scalar-Higgs-boson mass.

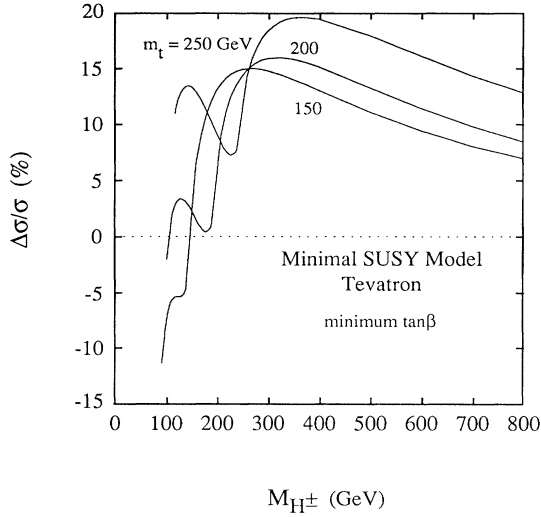


FIG. 5. Same as Fig. 4, but in the minimal supersymmetric two-Higgs-doublet model, and vs the charged-Higgs-boson mass.

boson), so its contribution is suppressed. Thus we do not find as large a Yukawa correction to top-quark production at the Tevatron in this model as in the generic two-Higgs-doublet model.

We show in Fig. 5 the maximum correction to the total top-quark production cross section at the Tevatron, as a function of the charged-Higgs-boson mass, in the minimal supersymmetric two-Higgs-doublet model, obtained by setting β to its minimum value for each top-quark mass. For a light charged Higgs boson, there is a cancellation between the contribution of the charged Higgs boson and the scalar Higgs bosons, as mentioned above. The dip at $M_{H^\pm} \approx m_t$ is due to the rapid variation of the charged-Higgs-boson form factor at the threshold for $t \rightarrow bH^+$. The implied lower limit on the charged-Higgs-boson mass from searches for $Z \rightarrow Z^*h$ at LEP is indicated in the figure [27]. The upper bound on $b \rightarrow s\gamma$ [22] implies $M_{H^\pm} > 500$ GeV [23,28], although the contribution of supersymmetric particles to this process is not included in this bound [29]. The correction to the total cross section does not exceed 20% for any value of M_{H^\pm} , and is therefore unlikely to be observable at the Tevatron, unless there is significant progress in the calculation of the total cross section. Furthermore, in the currently popular grand-unified supergravity models, $\tan\beta$ tends to be large rather than small [30].

The full weak correction to $e^+e^- \rightarrow t\bar{t}$ in a two-Higgs-doublet model, including the Yukawa correction, is given in Ref. [31], and the full weak correction near threshold is given in Ref. [32].

IV. CONCLUSIONS

We have calculated the correction of Yukawa strength, $g^2 m_t^2 / M_W^2$, to the total top-quark production cross section at the Fermilab Tevatron. In the standard Higgs

model, the correction is less than +2.4% for $m_t < 250$ GeV, much less than the theoretical uncertainty in the cross section.

The Yukawa correction to the total top-quark production cross section can be significantly enhanced in a two-Higgs-doublet model, if the doublet which generates the top-quark mass has a small vacuum-expectation value. Corrections as large as +23%, +28%, +35% for $m_t = 150, 200, 250$ GeV are obtained. Assuming a theoretical uncertainty in the cross section of $\pm 20\%$ from QCD, these corrections are potentially observable, and could be used as evidence for, or to place restrictions on, a two-Higgs-doublet model. A measurement of the total cross section with systematic and statistical errors less than $\pm 20\%$ would be required. If one specializes to the minimal supersymmetric two-Higgs-doublet model, one finds that the correction never exceeds 20%.

The top quark may provide a window on mass generation in the weak interaction. It is worthwhile to explore the different signatures of the top-quark mass-generating mechanism in top-quark production at the Tevatron. The calculation of the Yukawa correction to the total top-quark production cross section is a first step in this direction.

Note added in proof. The upper bound on the inclusive $b \rightarrow s\gamma$ branching fraction has recently been lowered to 5.4×10^{-4} [35]. This slightly increases the lower bound on the charged-Higgs-boson mass for a given value of $\tan\beta < 1$, and therefore does not alter the conclusions of this paper.

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APPENDIX

Below we give the form factors for the matrix element of the top-quark current. We separately list the contribution from the Higgs boson and the unphysical scalar Z and W bosons, of mass M_Z and M_W ('t Hooft-Feynman gauge). We set M_Z and M_W to zero (Landau gauge) in the numerical calculations, as described in the Introduction. In a two-Higgs-doublet model, the pseudoscalar-Higgs-boson (A) and the charged-Higgs-boson (H^\pm) form factors are given by the unphysical scalar Z - and W -boson form factors, with $M_Z \rightarrow M_A$ and $M_W \rightarrow M_{H^\pm}$, and multiplied by the square of the associated coupling factor in Table I. The same is true of the two scalar Higgs bosons (h, H) of the two-Higgs-doublet model with respect to the Higgs-boson form factor.

The form factors are written in terms of the usual one-, two-, and three-point scalar loop integrals A_0 , B_0 , and C_0 [33]. The integrals were evaluated with the code *FF* [34], whose notation we have adopted.

1. Higgs boson (scalar Higgs bosons)

$$\begin{aligned}
V(s - 4m_t^2)^2 &= (s - 4m_t^2)^2 \left[\frac{1}{2m_t^2} [M_H^2 B_0(m_t^2, M_H^2, m_t^2) - A_0(M_H^2) + A_0(m_t^2)] + (4m_t^2 - M_H^2) B'_0(m_t^2, M_H^2, m_t^2) \right] \\
&+ [A_0(M_H^2) - A_0(m_t^2)] [-2(s - 4m_t^2)] + B_0(m_t^2, M_H^2, m_t^2) [sM_H^2 - 8sm_t^2 - 16m_t^2 M_H^2 + 32m_t^4] \\
&+ B_0(m_t^2, m_t^2, s) \left[-\frac{1}{2}s^2 + sM_H^2 + 10sm_t^2 + 8m_t^2 M_H^2 - 32m_t^4 \right] \\
&+ C_0(m_t^2, m_t^2, M_H^2, s, m_t^2, m_t^2) [4s^2 m_t^2 + sM_H^4 - 32sm_t^4 + 12sm_t^2 M_H^2 + 8m_t^2 M_H^4 - 48m_t^4 M_H^2 + 64m_t^6] \\
&+ \frac{1}{2}s(s - 4m_t^2),
\end{aligned}$$

$$\begin{aligned}
T(s - 4m_t^2)^2 \frac{1}{2m_t} &= [A_0(M_H^2) - A_0(m_t^2)] \frac{1}{2m_t^2} (s - 4m_t^2) + B_0(m_t^2, M_H^2, m_t^2) \left[-s \frac{M_H^2}{2m_t^2} + 2s + 5M_H^2 - 8m_t^2 \right] \\
&+ B_0(m_t^2, m_t^2, s) \left[-\frac{3}{2}s - 3M_H^2 + 6m_t^2 \right] + C_0(m_t^2, m_t^2, M_H^2, s, m_t^2, m_t^2) [-3sM_H^2 - 3M_H^4 + 12m_t^2 M_H^2] \\
&- \frac{1}{2}(s - 4m_t^2),
\end{aligned}$$

$A = 0.$

2. Unphysical scalar Z boson (pseudoscalar Higgs boson)

$$\begin{aligned}
V(s - 4m_t^2)^2 &= (s - 4m_t^2)^2 \left[\frac{1}{2m_t^2} [M_Z^2 B_0(m_t^2, M_Z^2, m_t^2) - A_0(M_Z^2) + A_0(m_t^2)] - M_Z^2 B'_0(m_t^2, M_Z^2, m_t^2) \right] \\
&+ [A_0(M_Z^2) - A_0(m_t^2)] [-2(s - 4m_t^2)] + B_0(m_t^2, M_Z^2, m_t^2) [sM_Z^2 - 16m_t^2 M_Z^2] \\
&+ B_0(m_t^2, m_t^2, s) \left[-\frac{1}{2}s^2 + sM_Z^2 + 2sm_t^2 + 8m_t^2 M_Z^2 \right] \\
&+ C_0(m_t^2, m_t^2, M_Z^2, s, m_t^2, m_t^2) [sM_Z^4 + 4sm_t^2 M_Z^2 + 8m_t^2 M_Z^4 - 16m_t^4 M_Z^2] + \frac{1}{2}s(s - 4m_t^2),
\end{aligned}$$

$$\begin{aligned}
T(s - 4m_t^2)^2 \frac{1}{2m_t} &= [A_0(M_Z^2) - A_0(m_t^2)] \frac{1}{2m_t^2} (s - 4m_t^2) + B_0(m_t^2, M_Z^2, m_t^2) \left[-s \frac{M_Z^2}{2m_t^2} + 5M_Z^2 \right] \\
&+ B_0(m_t^2, m_t^2, s) \left[\frac{1}{2}s - 3M_Z^2 - 2m_t^2 \right] + C_0(m_t^2, m_t^2, M_Z^2, s, m_t^2, m_t^2) [-sM_Z^2 - 3M_Z^4 + 4m_t^2 M_Z^2] \\
&- \frac{1}{2}(s - 4m_t^2),
\end{aligned}$$

$A = 0.$

3. Unphysical scalar W boson (charged Higgs boson)

$$\begin{aligned}
V(s - 4m_t^2)^2 &= (s - 4m_t^2)^2 \left[\frac{1}{2m_t^2} [(m_t^2 + M_W^2) B_0(0, M_W^2, m_t^2) - A_0(M_W^2)] + (m_t^2 - M_W^2) B'_0(0, M_W^2, m_t^2) \right] \\
&+ A_0(M_W^2) [-2(s - 4m_t^2)] + B_0(0, M_W^2, m_t^2) [sM_W^2 - 3sm_t^2 - 16m_t^2 M_W^2] \\
&+ B_0(0, 0, s) \left[-\frac{1}{2}s^2 + sM_W^2 + 7sm_t^2 + 8m_t^2 M_W^2 - 8m_t^4 \right] \\
&+ C_0(0, 0, M_W^2, s, m_t^2, m_t^2) [s^2 m_t^2 + sM_W^4 - 3sm_t^4 + 10sm_t^2 M_W^2 + 8m_t^2 M_W^4 - 16m_t^4 M_W^2 + 8m_t^6] \\
&+ \frac{1}{2}s(s - 4m_t^2),
\end{aligned}$$

$$\begin{aligned}
T(s - 4m_t^2)^2 \frac{1}{2m_t^2} &= A_0(M_W^2) \frac{1}{2m_t^2} (s - 4m_t^2) + B_0(0, M_W^2, m_t^2) \left[-\frac{sM_W^2}{2m_t^2} + \frac{1}{2}s + 5M_W^2 + m_t^2 \right] \\
&\quad + B_0(0, 0, s) \left[-\frac{1}{2}s - 3M_W^2 - m_t^2 \right] \\
&\quad + C_0(0, 0, M_W^2, s, m_t^2, m_t^2) [-2sM_W^2 - sm_t^2 - 3M_W^4 + 2m_t^2M_W^2 + m_t^4] - \frac{1}{2}(s - 4m_t^2), \\
As(s - 4m_t^2) &= A_0(M_W^2) \frac{1}{2m_t^2} (s - 4m_t^2) + B_0(0, M_W^2, m_t^2) \left[-\frac{sM_W^2}{2m_t^2} + \frac{1}{2}s + M_W^2 + m_t^2 \right] \\
&\quad + B_0(0, 0, s) \left[-\frac{1}{2}s + M_W^2 - m_t^2 \right] \\
&\quad + C_0(0, 0, M_W^2, s, m_t^2, m_t^2) [-sm_t^2 + M_W^4 - 2m_t^2M_W^2 + m_t^4] + \frac{1}{2}(s - 4m_t^2).
\end{aligned}$$

The first line of each expression for V is the contribution from top-quark wave function renormalization, with

$$\begin{aligned}
B'_0(m_t^2, M_B^2, m_t^2) &= \frac{1}{m_t^2} \left[\frac{1}{x_+^B - x_-^B} \left(\frac{(x_+^B)^2}{x_-^B} \ln x_-^B - \frac{(x_-^B)^2}{x_+^B} \ln x_+^B \right) - 1 \right], \\
B'_0(0, M_B^2, m_t^2) &= \frac{1}{m_t^2} \left[r_W \ln \frac{r_W}{r_W - 1} - 1 \right],
\end{aligned}$$

where $r_B = M_B^2/m_t^2$ ($B = H, Z, W$), and

$$x_{\pm}^B = \frac{1}{2}[r_B \pm (r_B^2 - 4r_B)^{1/2}].$$

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