Production of $W^{\pm}H^{\mp}$ from heavy-quark fusion

Duane A. Dicus and Chung Kao
Center for Particle Theory, University of Texas at Austin, Austin, Texas 78712
(Received 22 September 1989)

The supersymmetric two-Higgs-doublet model is adopted to study the production of a charged Higgs boson ($H^{\mp}$) in association with a W boson ($W^{\pm}$) at future hadron colliders. The subprocess reactions $t\bar{t}\to W^{\pm}H^{\mp}$, $gt\to W^{\pm}H^{\mp}t$, $g\bar{t}\to W^{\pm}H^{\mp}\bar{t}$, and $gg\to W^{\pm}H^{\mp}t\bar{t}$ are combined to find the total rate.

I. INTRODUCTION

The standard model (SM) of electroweak interactions has been able to explain almost all experimental results to date, and it has been well confirmed especially with the discovery of the W and Z gauge bosons. Yet there is still a hole to be filled; that is, the spontaneous symmetry-breaking (SSB) sector. In the SM, the Higgs mechanism triggers SSB and only one Higgs doublet is required to generate masses of the fermions and gauge bosons. It leads to the appearance of a single neutral Higgs scalar after SSB. Some extensions of the SM have more Higgs multiplets and lead to additional physical spin-zero fields.

Although the Higgs-boson masses are *a priori* free parameters, unitarity constraints show that such masses must be less than 1 TeV or the SM is no longer a weak-coupling theory.[1] Therefore, at the energy of the CERN Large Hadron Collider (LHC) and, particularly, the Superconducting Super Collider (SSC), we should be able to discover the Higgs bosons or find the symmetry-breaking interactions becoming strong.

The production of $W^{\pm}H^{\mp}$ from top-quark and antiquark annihilation, $t\bar{t}\to W^{\pm}H^{\mp}$, have been studied[2] in the two-Higgs-doublet model popularized by supersymmetry. That work used heavy-quark distribution functions. It is known, however, that the contribution of single-Higgs-boson production from heavy-quark fusion, $Q\bar{Q}\to H^0$, overestimates the true cross section if the heavy quarks are simply treated as partons.[3] Therefore, it is necessary to evaluate $gg\to W^{\pm}H^{\mp}t\bar{t}$ and combine it with the processes $t\bar{t}\to W^{\pm}H^{\mp}$, $gt\to W^{\pm}H^{\mp}t$, and $g\bar{t}\to W^{\pm}H^{\mp}\bar{t}$ to obtain a precise cross section. It is also of interest to compare this "exact" cross section with the rates obtained from the $t\bar{t}$ reaction alone and from the $gg\to W^{\pm}H^{\mp}t\bar{t}$ reaction alone. The $t\bar{t}$ process is easy to calculate so it is useful to know if it is a good approximation to the total rate. On the other hand, if the top-quark mass is larger than the mass of the W, singularities in the matrix element (discussed below) leave the gg process as the only known approximation to the total rate. Recent experiments indicate the value of m_t is at least close to M_W, so the gg process may be the best we can do. For completeness we also compare our results with $W^{\pm}H^{\mp}$ production through gluon fusion into quark loops, $gg\to W^{\pm}H^{\mp}$.

In Sec. II the supersymmetric two-Higgs-doublet model will be briefly reviewed. Section III contains a discussion of the calculation and results from gluon fusion as well as exact top-quark fusion. A brief summary will be given in Sec. IV.

II. TWO-HIGGS-DOUBLET MODEL

In a general two-Higgs-doublet model, doublets with the vacuum expectation values v_1 and v_2 are needed to give mass to both up-type and down-type quarks as well as leptons and gauge bosons. After SSB, there remain five physical Higgs bosons: a pair of singly charged Higgs bosons $H^{\pm}$, two neutral scalars H_1^0 and H_2^0, and a neutral pseudoscalar H_3^0. The only constraint on v_1 and v_2 is $M_W^2=\frac{1}{2}g^2(v_1^2+v_2^2)$.

The couplings of the Higgs bosons to themselves and to fermions as well as gauge bosons are determined by six independent parameters: the four Higgs-boson masses, the ratio of vacuum expectation values $\tan\beta\equiv v_2/v_1$ $(0<\beta<\pi/2)$, and a mixing angle α between the weak and mass eigenstates of the neutral scalars H_1^0 and H_2^0 $(-\pi/2\le\alpha\le 0)$.

Supersymmetry is one of the most compelling reasons for considering two-Higgs-doublet models. The supersymmetric standard model[4] (SUSY SM) requires at least two SU(2) Higgs doublets, both to generate masses for fermions and gauge bosons and to cancel triangle anomalies associated with the fermionic partners of the Higgs bosons. In the minimal version of the SUSY SM, β and the mass of the charged Higgs boson, M_H, are the only independent parameters needed to specify all Higgs-boson masses and couplings. The neutral-Higgs-boson masses and the mixing angle α are given by

$$\begin{aligned}
&m_3^2=M_H^2-M_W^2\ ,\\
&m_{1,2}^2=\tfrac{1}{2}\{m_3^2+M_Z^2\\
&\qquad\pm[(m_3^2+M_Z^2)^2-4m_3^2M_Z^2\cos^2 2\beta]^{1/2}\}\ ,\\
&\cos2\alpha=-\cos2\beta\left[\frac{m_3^2-M_Z^2}{m_1^2-M_Z^2}\right]\ ,\\
&\sin2\alpha=-\sin2\beta\left[\frac{m_1^2+m_2^2}{m_1^2-m_2^2}\right]\ ,
\end{aligned}\tag{1}$$

where M_H, m_1, m_2, and m_3 are the masses of the charged Higgs boson $H^{\pm}$, the scalars H_1^0 and H_2^0, and the pseudoscalar H_3^0, respectively. These relations imply that $M_H > M_W$ and $m_2 < M_Z < m_1$. One natural limit is $v_2 = v_1$. In this case, $\beta = -\alpha = \pi/4$, $m_2 = 0$ (at the tree level), and $m_1 \simeq m_3$. All couplings have the same strength as in the SM, except for H_2^0, which decouples from weak gauge bosons.

Generally speaking, m_1 is always about the same as M_H and $|\sin\alpha|$ becomes larger for larger $\tan\beta$ but smaller for larger M_H. For $M_H > 2M_W$, which is the region that contributes most of the cross section, m_3 is close to M_H, m_2 is close to M_Z (except when $\tan\beta$ is very near one), and α is almost independent of M_H. Given the values of M_H and $\tan\beta$, all the other parameters can be figured out with the above formulas.

III. PRODUCTION RATES

One reaction which produces $W^{\pm}H^{\mp}$ from gluon fusion is $gg \to W^{\pm}H^{\mp}$ through triangle and box quark loops. It yields a significant contribution and has been studied in Ref. 2. Its cross section will be compared with the contributions from other processes in this paper.

Gluons can also fuse to $W^{\pm}H^{\mp}$ through tree diagrams, $gg \to W^{\pm}H^{\mp}t\bar{t}$, as shown in Fig. 1 and this can be combined with $t\bar{t} \to W^{\pm}H^{\mp}$ (simple top-quark fusion), $gt \to W^{\pm}H^{\mp}t$, and $g\bar{t} \to W^{\pm}H^{\mp}\bar{t}$ to obtain an exact cross section. The 2nd and 10th–12th diagrams are similar to simple top-quark fusion, with the top-quark distribution function arising from gluons splitting into $t\bar{t}$. The top-quark distribution function contains information which is not present in the process $gg \to W^{\pm}H^{\mp}t\bar{t}$ such as the effects of the radiation of an arbitrary number of gluons (in the leading-log approximation). There is a formalism which allows us to combine these two approaches to obtain a cross section which contains the physics of both without double counting.[5,6] This formalism has thus far only been applied to $2 \to 2$ and $2 \to 3$ reactions; this is the first application to a $2 \to 4$ process. The technique of Ref. 6 is followed to evaluate the exact cross section, the top-quark mass is maintained and the bottom-quark mass is set to be zero.

The distribution function of a parton i (e.g., $i = t, g$ represents top quark, gluon, respectively) in hadron A may be denoted by $f_A^i(x,\mu)$, where x is the fractional momentum and μ is the energy scale. For a top quark, it is determined according to the Altarelli-Parisi equation, with the condition $f_A^t(x, \mu \le m_t) = 0$. For $\mu \ge m_t$, the distribution function to the lowest order in α_s is given by

$$\tilde{f}^{\,t}_A(x,\mu) = \frac{\alpha_s}{4\pi} \ln\frac{\mu^2}{m_t^2} \int_x^1 \frac{dy}{y} P_{tg}\left[\frac{x}{y}\right] f_A^g(y,\mu) \,, \tag{2}$$

where P_{tg} is the Altarelli-Parisi splitting function

$$P_{tg}(z) = z^2 + (1-z)^2 \,. \tag{3}$$

When the differential cross sections of the subprocesses have been evaluated, the corresponding total cross sections are obtained by integration after weighting by the appropriate set of parton distribution functions. The exact cross section for the production of $W^{\pm}H^{\mp}$ from the collision of hadrons A and B is

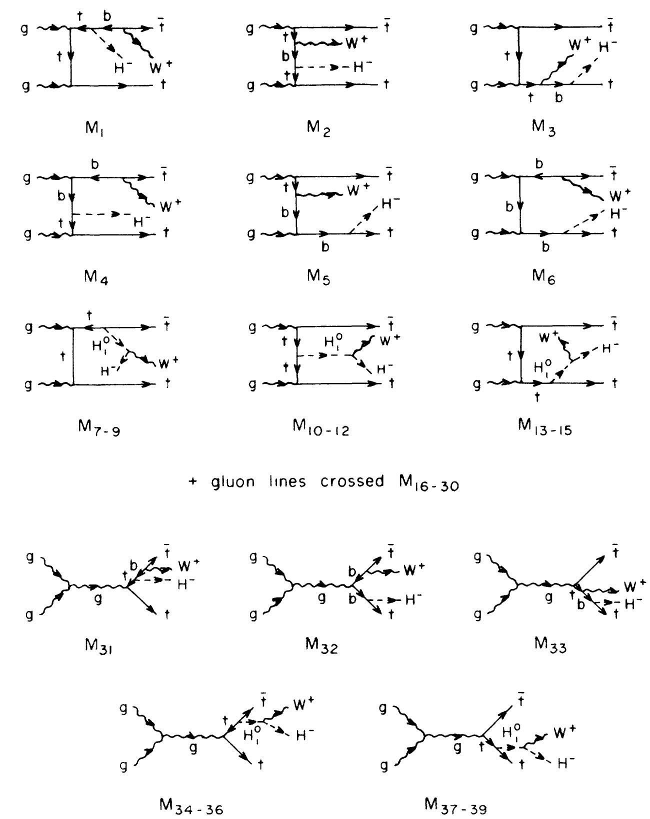

FIG. 1. Feynman diagrams for the process $gg \to W^+H^-t\bar{t}$, where H_i^0 are neutral Higgs boson, with $i = 1,2,3$. There are 39 diagrams.

$$\begin{aligned}\sigma_{AB\to W^{\pm}H^{\mp}} = {} & \bar{f}^{\,t}_A \otimes \hat{\sigma}^0_{t\bar{t}\to W^{\pm}H^{\mp}} \otimes \bar{f}^{\,\bar{t}}_B + \bar{f}^{\,\bar{t}}_A \otimes \hat{\sigma}^0_{\bar{t}t\to W^{\pm}H^{\mp}} \otimes \bar{f}^{\,t}_B + \bar{f}^{\,t}_A \otimes \hat{\sigma}^1_{tg\to W^{\pm}H^{\mp}t} \otimes f_B^g \\ & + f_A^g \otimes \hat{\sigma}^1_{gt\to W^{\pm}H^{\mp}t} \otimes \bar{f}^{\,t}_B + \bar{f}^{\,\bar{t}}_A \otimes \hat{\sigma}^1_{\bar{t}g\to W^{\pm}H^{\mp}\bar{t}} \otimes f_B^g + f_A^g \otimes \hat{\sigma}^1_{g\bar{t}\to W^{\pm}H^{\mp}\bar{t}} \otimes \bar{f}^{\,\bar{t}}_B \\ & + f_A^g \otimes \hat{\sigma}^2_{gg\to W^{\pm}H^{\mp}t\bar{t}} \otimes f_B^g \,, \end{aligned} \tag{4}$$

where $\bar{f}^{\,t}_A = f_A^t - \tilde{f}^{\,t}_A$ is the subtracted top-quark distribution function which must be used to avoid double counting when the gluon reactions are included. The superscripts denote the power of α_s. The distribution function of Duke and Owens set I (Ref. 7) have been used for both gluon-fusion reactions, and the code developed by Tung[8] with Duke and Owens set I as input has been employed for simple top-quark fusion as well as the exact top-quark fusion. The convolution $\otimes$ in Eq. (4) denotes the usual integration over the parton fractional momenta:

$$\sigma = \int_{M_W+M_H}^{\sqrt{s}} \frac{2M}{s} dM \int_{-Y}^{Y} dy \{[f_A^i(x_1,M^2) f_B^j(x_2,M^2) + f_A^i(x_2,M^2) f_B^j(x_1,M^2)]/(1+\delta_{ij})\}$$
$$\times \int_{-Z_0}^{Z_0} dz \left[\frac{d\hat{\sigma}}{dz} \right] , \tag{5}$$

where f may be $\bar{f}$ and M is the invariant mass of $W^{\pm}H^{\mp}$. Y is a rapidity cut and

$$z = \cos\theta, \quad x_{1,2} = \frac{M}{\sqrt{s}} e^{\pm y} ,$$
$$Z_0 = \min[\beta^{-1} \tanh(Y - |y|), 1] , \tag{6}$$

with

$$\beta = \left[\left[1 - \frac{M_W^2 + M_H^2}{M^2} \right]^2 - \frac{4M_W^2 M_H^2}{M^4} \right]^{1/2}$$
$$\times \left[1 + \frac{M_W^2 - M_H^2}{M^2} \right]^{-1} \tag{7}$$

The cross sections for producing $W^{\pm}H^{\mp}$ were calculated with the following values of the parameters: $\alpha = \frac{1}{128}$, $M_Z = 92.0$ GeV, $\sin^2\theta_W = 0.23$, and $Y = 2.5$. Four processes are separately presented: (1) the simple top-quark fusion (simply denoted by $t\bar{t}$), (2) the exact top-quark fusion (exact), (3) the gluon fusion via tree diagrams (gg tree) and (4) the gluon fusion through quark loops (gg loops). Since top-quark fusion and gluon fusion are presented for comparison with the full cross section [given by Eq. (4)] they are calculated using unsubtracted functions.

An interesting aspect has been found in all processes for the Higgs-boson mass considered: the amplitudes are all approximately proportional to $\cot\beta$, that is,

$$\sigma \propto \cot^2\beta , \tag{8}$$

where $\tan\beta \equiv v_2/v_1$. For every reaction, the total amplitude A with the subamplitudes A_i ($i = 1,2,3$) may be written as

$$A = c_1 A_1 + c_2 A_2 + c_3 A_3 \tag{9}$$

with

$$c_1 = \cot\beta ,$$
$$c_2 = \frac{\sin^2\alpha \cot\beta - \frac{1}{2}\sin 2\alpha}{D_1} + \frac{\cos^2\alpha \cot\beta + \frac{1}{2}\sin 2\alpha}{D_2} , \tag{10}$$
$$c_3 = \cot\beta / D_3 ,$$

where $D_i = M^2 - m_i^2$ and m_i equal to the masses of the two neutral scalar ($i = 1,2$) or the pseudoscalar ($i = 3$). It can be shown that

$$c_2 \simeq \cot\beta \left[\frac{\sin^2\alpha}{D_1} + \frac{\cos^2\alpha}{D_2} \right] \tag{11}$$

if

$$\frac{1}{\sin^2\beta} \frac{M^2 - M_Z^2}{M_H^2} \gg 1 . \tag{12}$$

The cross section is obtained by integrating over the invariant mass M, with $\sqrt{s} \ge M \ge M_H + M_W$. Condition (12) is obviously satisfied if $\tan\beta$ is small. If $\tan\beta$ is not small Eq. (8) is still satisfied over the important regions of integration as long as $M_H \le 1200$ GeV. Therefore, if $v_2 > v_1$ the production rate will be suppressed, while if $v_1 > v_2$ it will be enhanced. When the cross section is known at one value of $\tan\beta$ (e.g., $\tan\beta = 1.0$), the cross

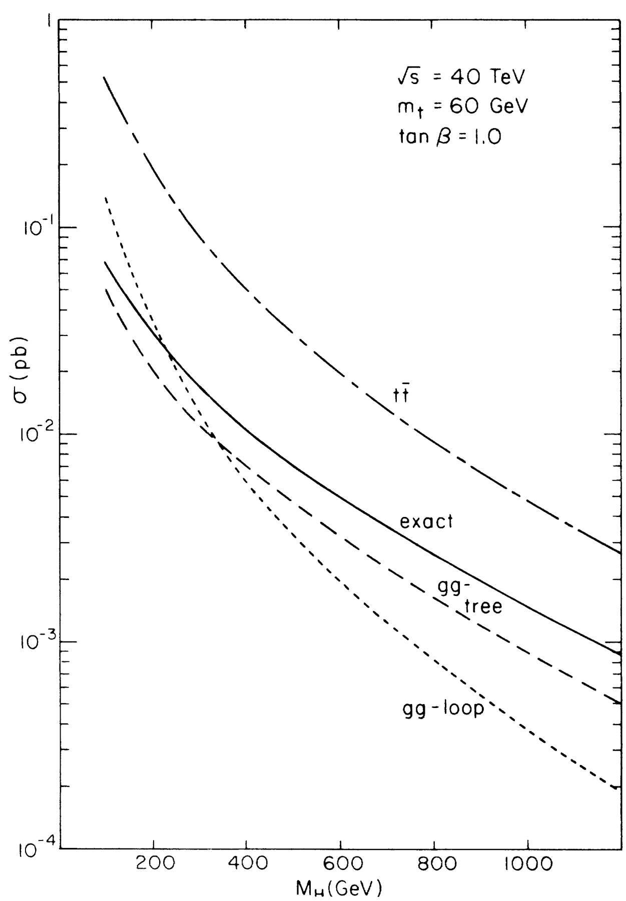

FIG. 2. Cross section for the process $pp \to W^{\pm}H^{\mp} + X$ (which is twice the contribution from either $pp \to W^+H^- + X$ or $pp \to W^-H^+ + X$), as a function of M_H at $\sqrt{s} = 40$ TeV with $m_t = 60$ GeV and $\tan\beta = 1.0$. The solid line is the cross section of exact top-quark fusion [Eq. (4)], the long-dashed–short-dashed line is the result for the simple top-quark fusion. The dashed line is the contribution from gluon fusion via tree diagrams and the dotted line is the result from gluon fusion through quark loops.

section at all values of $\tan\beta$ can be determined easily.

In Fig. 2 the cross sections are shown at the SSC energy ($\sqrt{s}=40$ TeV) with $m_t=60$ GeV and $\tan\beta=1.0$. If we simply treat the top quark as a patron, the contribution from $t\bar{t}\to W^{\pm}H^{\mp}$ overestimates the exact cross section by approximately a factor of 8 at $M_H=100$ GeV. It becomes better for a larger Higgs-boson mass but is still about three times the exact cross section at $M_H=1200$ GeV. The process of gluon fusion via tree diagrams, $gg\to W^{\pm}H^{\mp}t\bar{t}$, yields a reasonably accurate contribution for all values of M_H considered. At $M_H=100$ GeV it is about a factor of 1.3 times smaller than the exact cross section and 1.7 times smaller at $M_H=1200$ GeV. The contribution from gluon fusion via quark loops is large at a small Higgs-boson mass and becomes quite small for a large Higgs-boson mass.

Were v_2 much less than v_1, the production rate would be enhanced greatly. As shown in Fig. 3, at $\tan\beta=0.3$, the cross section of each process is about 11 times the corresponding process in Fig. 2. For example, the cross section of $t\bar{t}\to W^{\pm}H^{\mp}$ with $M_H=100$ GeV and $m_t=60$ GeV is 0.54 pb at $\tan\beta=1.0$, and has become 6.0 pb at $\tan\beta=0.3$. These numbers and the fact that all ratios among different contributions stay the same gives an empirical demonstration of Eq. (8) above.

At LHC energy ($\sqrt{s}=17$ TeV), all the production rates are reduced as shown in Fig. 4. With $M_H=100$ GeV and $m_t=60$ GeV, the cross section of $t\bar{t}$, exact, gg tree, and gg loop are approximately 3.8, 5.1, 6.9, and 3.7 times smaller than those at the SSC. They are even smaller for a larger Higgs-boson mass, and become about 12, 14, 21, and 14 times smaller at $M_H=1200$ GeV. Both $t\bar{t}$ and gg tree become worse approximations of the exact cross section, though the gg-tree contribution is still reasonably accurate. At $M_H=100$ GeV the ratios of $t\bar{t}$/exact and exact/gg tree are about 11 and 1.8 and become 3.6 and 2.5 at $M_H=1200$ GeV.

In Fig. 5 we attempt to show the cross sections as functions of the top-quark mass. However, there is a singularity in the $t\bar{t}$ process[2] when $M_H>m_t>M_W$ (or $M_H<m_t<M_W$). In this mass region, a pole develops in the t-channel, b-exchange graph where $t\to Wb$ and $\bar{t}b\to H$ (or $\bar{t}\to H\bar{b}$ and $\bar{b}t\to W$) are kinematically allowed. If the bottom-quark mass, which we have set to zero, is taken to be nonzero, then this can occur only for a range of values of $\hat{s}$, but these values are still within the physical region. This kinematic pole arises also in other processes. For example, it would occur in $t\bar{b}\to W^+Z$ for $\hat{s}>m_t^2(M_Z^2-M_W^2+m_t^2)/(m_t^2-M_W^2)$, if $m_t>M_W$. Including a width for the particle on pole is not a solution since one can think of situations where that particle is absolutely stable. The only way to avoid the pole seems to be to take into account how the heavy particles are produced so that they are never on mass shell—in this case

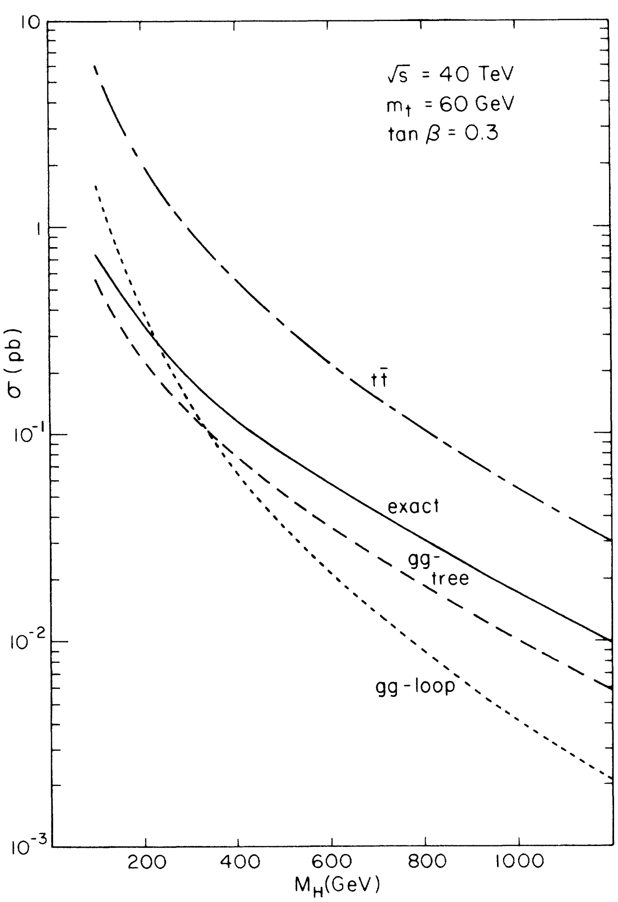

FIG. 3. Same as in Fig. 2, except $\tan\beta=0.3$.

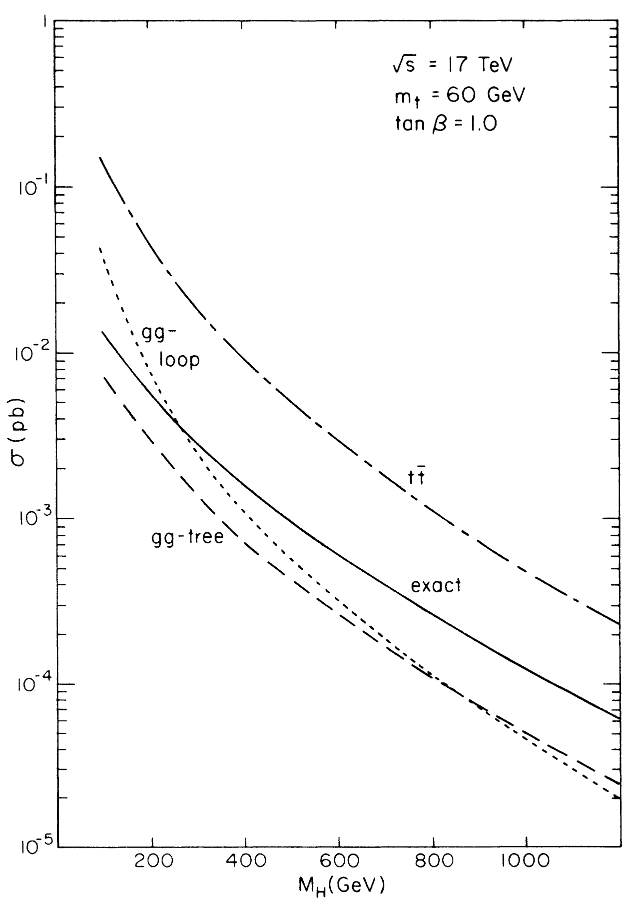

FIG. 4. Same as in Fig. 2, except $\sqrt{s}=17$ TeV.

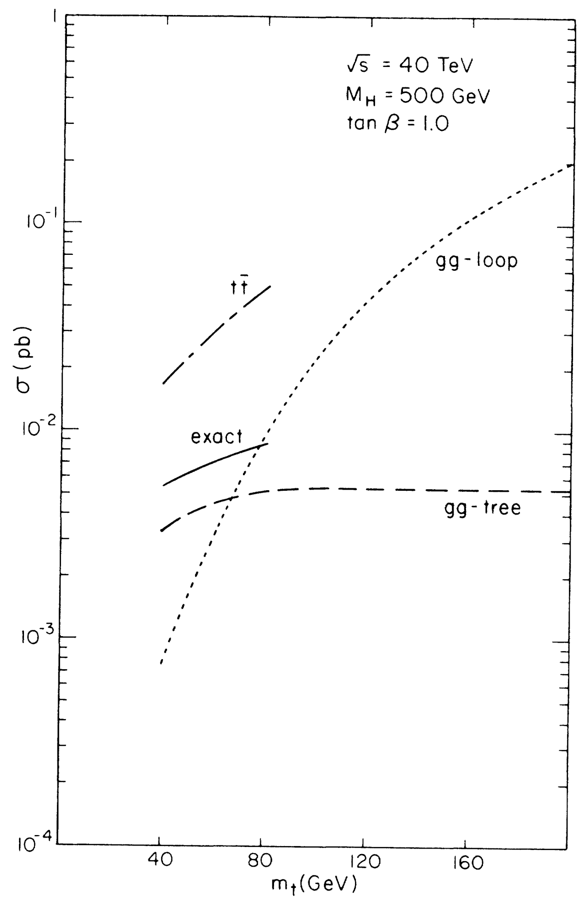

FIG. 5. Cross section for the process $pp \to W^{\pm}H^{\mp}+X$ as a function of the top-quark mass at $\sqrt{s}=40$ TeV with $M_H=500$ GeV and $\tan\beta=1.0$. The solid line is the cross section of exact top-quark fusion, the long-dashed–short-dashed line is the result for the simple top-quark fusion. The dashed line is the contribution from gluon fusion via tree diagrams and the dotted line is the result from gluon fusion through quark loops.

that means taking $gg \to t\bar{t}t\bar{t} \to W^{\pm}H^{\mp}t\bar{t}$ as the total cross section. It might be that the sharp rise in the $t\bar{t}$ cross section with increasing m_t is because of the approach of the singularity; thus the $t\bar{t}$ and exact production rates at $m_t \simeq 80$ GeV may not be accurate. The gg-tree contribution remains a reasonably good approximation to the exact cross section even up to a top-quark mass of 80 GeV and becomes flat for m_t larger than 100 GeV. The gg-loop contribution is the smallest when $m_t \le 60$ GeV and rises sharply for larger m_t.

The cleanest signature of W^+H^-, is given by $H^- \to \bar{t}b \to W^-b\,(+\text{ soft } b) \to l^+l^- + \text{jets} + \not{p}_T$ with the associated W^+ decaying leptonically: $W^+ \to l^+ + \not{p}_T$. The signature of W^-H^+ is the charge conjugation of these modes.

If the charged Higgs boson has a small mass $M_H \simeq 300$ GeV, with $m_t=60$ GeV, and $\tan\beta=0.3$ the cross sections for the $t\bar{t}$, exact, gg-tree, and gg-loop reactions at the SSC will be about 1, 0.2, 0.13, and 0.15 pb. With a yearly integrated luminosity at the SSC of 10^4 pb^{-1}/yr and the branching ratio folded in, there will be approximately 120, 24, 15, and 18 events/yr, respectively. The production rates will be larger for smaller Higgs-boson mass and bigger top-quark mass and will be greatly enhanced if $\tan\beta \ll 1.0$.

IV. SUMMARY

The production of a charged Higgs boson ($H^{\mp}$) together with a W boson ($W^{\pm}$) has been investigated in the two-Higgs-doublet model at the SSC and the LHC via $pp \to W^{\pm}H^{\mp}+X$. The SUSY SM relations have been used to reduce a large number of free parameters to three: m_t, M_H, and $\tan\beta$. The cross sections of four processes, simple top-quark fusion ($t\bar{t}$), exact top-quark fusion (exact), and gluon fusion through tree diagrams (gg tree), as well as quark loops (gg loop), are presented. The SSC produces the larger rates by a factor of 4–20 than the LHC for the different processes if identical luminosities are assumed. The $t\bar{t}$ cross section overestimates the exact contribution, while the gg-tree contribution is a reasonably accurate approximation, which agrees with previous analyses of $2\to2$ and $2\to3$ process. Because of a singularity in the region $M_H > m_t > M_W$, the gg-tree process becomes the only clue for the exact cross section when $m_t \gtrsim 80$ GeV. Evidence is accumulating that the top-quark mass is large,[9] probably bigger than 78 GeV; the analysis at lower values of m_t is necessary to give us confidence in the gluon-fusion process at m_t values where the exact method of determining the cross section [Eq. (4)] cannot be used. The gg-loop process produces large signal rates for $m_t > 100$ GeV. The production rates become larger for smaller Higgs-boson mass and bigger top-quark mass and it will be greatly enhanced if $\tan\beta \ll 1.0$, which might make the detection of a SUSY charged Higgs boson easier than that of a SM neutral Higgs boson.

ACKNOWLEDGMENTS

The computing resources were provided by the University of Texas Center for High Performance Computing. C.K. would like to acknowledge the Graduate School of the University of Texas at Austin for financial support. This work was supported in part by the U. S. Department of Energy under Contract No. DE-FG05-ER8540200.

[1] D. A. Dicus and V. S. Mathur, Phys. Rev. D **7**, 3111 (1973); M. Veltman, Acta Phys. Pol. **B8**, 475 (1977); B. W. Lee, C. Quigg, and H. Thacker, Phys. Rev. D **16**, 1519 (1977); J. van der Biji and M. Veltman, Nucl. Phys. **B231**, 205 (1984).

[2] D. A. Dicus, J. L. Hewett, C. Kao, and T. G. Rizzo, Phys. Rev. D **40**, 787 (1989).

[3] D. A. Dicus and S. S. D. Willenbrock, Phys. Rev. D **39**, 751 (1989).

[4]J. F. Gunion and H. E. Haber, Nucl. Phys. **B272**, 1 (1986); **B278**, 449 (1986).
[5]R. M. Barnett, H. E. Haber, and D. E. Soper, Nucl. Phys. **B306**, 697 (1988).
[6]F. I. Olness and W.-K. Tung, Nucl. Phys. **B308**, 813 (1988).
[7]D. Duke and J. Owens, Phys. Rev. D **30**, 49 (1984).
[8]J. C. Collins and W.-K. Tung, Nucl. Phys. **B278**, 934 (1986).
[9]P. Sinervo, in Proceedings of the Fourteenth International Symposium on Lepton and Photon Interactions, Stanford, California, 1989 (unpublished).