Associated production of Intermediate Higgs or Z-boson

with $t\bar{t}$ pair in $\gamma\gamma$ collision

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ABSTRACT

Photon-photon linear colliders can be realized by laser back-scattering technique on the next generation linear $e^+e^-$ colliders. Here the associated productions of an intermediate mass Higgs (IMH) or Z-boson with $t\bar{t}$ pair in $\gamma\gamma$ collisions are studied. Since IMH is very unlikely to decay into $t\bar{t}$ pair, $t\bar{t}H$ production is the only direct channel to probe the Higgs-top Yukawa coupling in case of an IMH. $t\bar{t}Z$ production can be a potential background to $t\bar{t}H$ if the Higgs mass is close to $m_Z$. As an alternative to its parent $e^+e^-$ collider, $\gamma\gamma \rightarrow t\bar{t}H(Z)$ productions are compared with the corresponding productions in the $e^+e^-$ collisions.

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I. INTRODUCTION

Ever since the discoveries of $W$ and $Z$ bosons, the Standard Model (SM) has been tested to high accuracy. However, the spontaneous symmetry-breaking is not yet well understood, nor is there any experimental evidence to favour any particular symmetry-breaking model. The simplest model is the Minimal SM with a single neutral scalar Higgs boson to activate the Higgs mechanism. One of the major goals of the next generation $pp$, $e^+e^-$, $ep$ and the newly discussed $\gamma\gamma$ colliders is to look for the Higgs boson. For heavy Higgs ($m_H \gtrsim 2m_Z$) the $W^+W^-$ and $ZZ$ decay modes have been shown to be viable channels for detection in future $e^+e^-$ and $pp$ colliders. More troublesome is the intermediate mass Higgs, one must look for $H \rightarrow \gamma\gamma$, $WW^*$ and $ZZ^*$, and $b\bar{b}$ or $\tau\tau$ modes, and the feasibility depends sensitively on the resolution of the detectors.

On the other hand, the top-quark is very likely to exist because there must be an $SU(2)_L$ isospin partner to the $b$-quark. In the minimal SM the fermions acquire masses via their Yukawa coupling to the Higgs. At tree level the Higgs-boson couples to a fermion of mass $m_f$ with strength

$$g_{fH} = -\frac{i g m_f}{2m_W},$$

where $g$ is the $SU(2)$ gauge coupling. The coupling in Eq. (1) can be directly probed in the decay of the Higgs boson into a pair of fermions if kinematically allowed. To measure $g_{tH}$ directly by this method however we need a Higgs of mass $> 2m_t$. The top-quark mass is likely in the range $120\sim200$ GeV, so we need a Higgs-boson of mass greater than about 250 GeV to allow the decay into a $t\bar{t}$ pair. Consequently, if the Higgs mass lies within the intermediate mass range ($m_W < m_H < 2m_Z$), this method cannot be used to probe the $g_{tH}$ coupling directly. It can be probed indirectly in the decay of $H \rightarrow \gamma\gamma$ or $gg$ (see e.g. Refs. [4,6]) or the fusion of $\gamma\gamma \rightarrow H$ or $gg \rightarrow H$ through an internal top-quark loop; but it is likely to be affected by the presence of other heavy particles beyond the SM. An alternative direct probe is to use the associated production of a Higgs-boson with a $t\bar{t}$ pair at the $e^+e^-$ and $pp$ colliders. In principle, the same coupling can also be probed in the production process $e^+e^- \rightarrow t\bar{t}Z$, but the contribution
from the Higgs-exchange diagram is very small relative to the contributions from other diagrams unless the Higgs mass is above the $t\bar{t}$ threshold, so the $e^+e^- \rightarrow t\bar{t}Z$ production is very insensitive to the presence of an IMH. Therefore, in the case of IMH, $t\bar{t}Z$ production is not a good channel to probe the Higgs-top coupling, but rather a potential background to $t\bar{t}H$ production, especially if $m_H$ is close to $m_Z$.

With the new possibility of $\gamma\gamma$ collisions at $e^+e^-$ colliders, the production process $\gamma\gamma \rightarrow t\bar{t}H$ offers another possible direct test of the $g_{ttH}$ coupling in addition to $t\bar{t}H$ production in $e^+e^-$ and hadronic collisions. The $\gamma\gamma$ collisions at $e^+e^-$ machines can be realized by shining a low energy (a few $eV$) laser beam at a very small angle $\alpha_0$, almost head to head, to the incident electron beam. By Compton scattering, there are abundant, hard back-scattered photons in the same direction as the incident electron, which carry a substantial fraction of the energy of the incident electron. Similarly, another laser beam can be directed onto the positron beam, and the resulting $\gamma$ beams effectively make a $\gamma\gamma$ collider. Actually, the second beam need not be positrons, but could also be electrons. For further technical details please see Ref. [19]. Another possibility is to use the beamstrahlung effect [20] but this method produces photons mainly in the soft region [19], and depends critically on the beam structure [20]. For the productions of $t\bar{t}H$ and $t\bar{t}Z$ we would need a high energy threshold of the beamstrahlung photons. Therefore we shall limit our calculations to $\gamma\gamma$ collisions produced by the laser back-scattering method.

The best signal for $t\bar{t}H$ production will be due to the dominant decay modes $H \rightarrow b\bar{b}$ and $t \rightarrow bW$, and therefore the signature is

$$\gamma\gamma \rightarrow t\bar{t}H \rightarrow b\bar{b}b\bar{b}WW.$$  \hspace{1cm} (2)

Since these rare events will only be searched for after the discovery of the Higgs-boson, backgrounds can be removed by using the constraints due to the $W$, $t$ and $H$ masses. Even so, $t\bar{t}Z$ production is a potential background, especially if the Higgs mass is close to the $Z$ mass; e.g., $|m_H - m_Z|$ less than a few GeV. Tagging the $b$ would be very helpful, but $b$-tagging efficiency is, so far, quite uncertain. If $b$-tagging has a high efficiency, then the $t\bar{t}Z$ background can be reduced substantially. There might be kinematic regions where $t\bar{t}H$ production dominates $t\bar{t}Z$ production in $\gamma\gamma$ collisions.
even though $t\bar{t}Z$ production is larger than $t\bar{t}H$ production in $e^+e^-$ collisions. For example, for $m_t = 150$ GeV and $m_H = 100$ GeV, $\sigma(e^+e^- \rightarrow t\bar{t}H) \simeq 2$ fb, and $\sigma(e^+e^- \rightarrow t\bar{t}Z) \simeq 5$ fb for $e^+e^-$ collisions at 1 TeV. However, it was found in Ref. [21] that $t\bar{t}$ production in $\gamma\gamma$ collisions realized by laser back-scattering is slightly larger than the direct $e^+e^- \rightarrow t\bar{t}$ production for $m_t \lesssim 130$ GeV at $\sqrt{s} = 0.5$ TeV; and at $\sqrt{s} = 1$ TeV the production of $\gamma\gamma \rightarrow t\bar{t}$ is much larger than $e^+e^- \rightarrow t\bar{t}$ for $m_t \sim 100 - 200$ GeV both with and without considering the threshold QCD effect. In the following we will explore how feasible the $\gamma\gamma$ collider is for $t\bar{t}H$ production, which will then directly probe the $g_{ttH}$ Yukawa coupling. In Sec. II we will present the calculation methods, which include the photon luminosity and subprocess cross sections. The results are discussed in Sec. III, and in Sec. IV the conclusions are summarized.

II. CALCULATION METHODS

A. Photon Luminosity

Using the laser back-scattering technique on an electron- or positron-beam abundant numbers of hard photons can be produced nearly in the same direction as the original beam. A low energy $\omega_0$ (a few $eV$) laser beam is directed onto the electron beam almost head to head. The energy $\omega$ of the scattered photon depends on its angle $\theta$ with respect to the incident electron beam and is given by

$$\omega = \frac{E_0(\frac{\xi}{1+\xi})}{1+(\frac{\theta}{\theta_0})^2},$$

where

$$\theta_0 = \frac{m_e}{E_0} \sqrt{1+\xi}, \quad \xi = \frac{4E_0\omega_0}{m_e^2},$$

and $E_0$ is the energy of the incident electron. Therefore, at $\theta = 0$, $\omega = E_0\xi/(1+\xi) = \omega_{\max}$ is the maximum energy of the back-scattered photon. The energy spectrum of the back-scattered photon, shown in Fig. [4], is given by [19]

$$F_{\gamma/e}(x) = \frac{1}{D(\xi)} \left[ 1 - x + \frac{1}{1-x} - \frac{4x}{\xi(1-x)} + \frac{4x^2}{\xi^2(1-x)^2} \right],$$

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where

\[ D(\xi) = (1 - 4\frac{1}{\xi} - 8\frac{1}{\xi^2}) \ln(1 + \xi) + \frac{1}{2} + \frac{8}{\xi} - \frac{1}{2(1 + \xi)^2}, \]  

and \( x = \omega/E_0 \) is the fraction of the energy of the incident electron carried by the back-scattered photon. Therefore

\[ x_{\text{max}} = \frac{\omega_{\text{max}}}{E_0} = \frac{\xi}{1 + \xi}, \]  

is the maximum fraction of energy carried away by the back-scattered photon. From Eq. (6) and (7) the portion of photons with energy close to \( \omega_{\text{max}} \) grows with \( E_0 \) and \( \omega_0 \), and so does \( x_{\text{max}} \). However, we should not choose a large \( \omega_0 \), or the back-scattered photon will interact with the incident photon and create unwanted \( e^+e^- \) pairs. The threshold for \( e^+e^- \) pair creation is \( \omega_{\text{max}}\omega_0 > m_e^2 \), so we require \( \omega_{\text{max}}\omega_0 \lesssim m_e^2 \). Solving \( \omega_{\text{max}}\omega_0 = m_e^2 \), we find

\[ \xi = 2(1 + \sqrt{2}) \approx 4.8. \]  

For the choice \( \xi = 4.8 \) one finds \( x_{\text{max}} \approx 0.83 \), \( D(\xi) \approx 1.8 \), and

\[ \omega_0 = \frac{\xi m_e^2}{E_0}, \]

\[ = \begin{cases} 1.25 \text{eV} & \text{for a 0.5 TeV } e^+e^- \text{ collider} \\ 0.63 \text{eV} & \text{for a 1 TeV } e^+e^- \text{ collider.} \end{cases} \]  

Here we assume that the average polarization of the back-scattered photon is zero; i.e., an unpolarized \( \gamma \)-beam. We also assume that, on average, the number of the back-scattered photons produced per electron is 1, i.e., the conversion coefficient \( k \) is equal 1.

\textbf{B. Subprocesses}

For \( \gamma\gamma \to t\bar{t}H \) the contributing Feynman diagrams are shown in Fig. 2, in which the cross diagrams with the interchange of the two incoming photons are not shown. The Higgs can be radiated from any fermion-line, so each diagram is proportional to \( g_{t\bar{t}H} \) and the resulting cross section will then be proportional to \( g_{t\bar{t}H}^2 \). Consequently, this process directly probes the Higgs-top
coupling. The amplitudes for the contributing Feynman diagrams are given in Appendix A. The corresponding Feynman diagrams for $\gamma\gamma \rightarrow t\bar{t}Z$ can be derived from those in Fig. 2 by simply replacing the Higgs by the Z. These Feynman amplitudes are also given in Appendix A. The subprocesses $e^+e^- \rightarrow t\bar{t}H$ [16] and $t\bar{t}Z$ [18] have been previously calculated, and it is not necessary to repeat these formulas here. We, however, independently did the calculations and agree with the results in Refs. [16] and [18], respectively.

To obtain the total cross sections $\sigma$ we fold in the photon luminosity with the cross section $\hat{\sigma}$ for the subprocesses. The resulting total cross section $\sigma$ is

$$\sigma = \int_{x_{1\text{min}}}^{x_{1\text{max}}} \int_{x_{2\text{min}}}^{x_{2\text{max}}} F_{\gamma/e}(x_1) F_{\gamma/e}(x_2) \hat{\sigma}(\gamma\gamma \rightarrow t\bar{t}V) \text{at } \hat{s} = x_1 x_2 s \, dx_1 dx_2, \quad (10)$$

with the constraints

$$x_1, x_2 \leq x_{\text{max}}, \quad (11)$$

$$\frac{(2m_t + m_V)^2}{s} \leq x_1 x_2, \quad \text{where } V = H, Z. \quad (12)$$

Throughout the paper, $\sqrt{s}$ always refers to the center-of-mass energy of the parent $e^+e^-$ collider.

**III. RESULTS & DISCUSSION**

We have used the energy spectrum of back-scattered photons shown in Fig. 1. With the choice of $\xi = 4.8$ a large fraction (> 50%) of the photons have energies greater than $0.5E_0$ and the spectrum peaks at the end point $x_{\text{max}} \simeq 0.83$. We shall consider, typically, $m_t > 120$ GeV, which satisfies the CDF 95% confidence level bound of $\gtrsim 91$ GeV [22], and the range $m_H \sim 60 - 140$ GeV. The energy threshold for $t\bar{t}H$ production will then be at least 300 GeV. The corresponding threshold value for $z = \sqrt{x_1 x_2} = \sqrt{s}/\hat{s}$ is 0.6, 0.3 and 0.15 for 0.5, 1 and 2 TeV $e^+e^-$ colliders, respectively. At these high values of $z$, the gluon content inside the “resolved” photon is negligible (see Fig. 1(b) and 1(c) of Ref. [21]). Therefore, we need only to consider direct $\gamma\gamma$ collisions.

In Fig. 3 we show the total cross sections as a function of the center-of-mass energy $\sqrt{s}$ of the parent $e^+e^-$ collider for $m_H = 90$ GeV and $m_t = 120$ and 150 GeV. We have ignored beamstrahlung
and bremsstrahlung of the initial states in the calculations of $e^+e^- \rightarrow t\bar{t}H$, $t\bar{t}Z$. For $e^+e^-$ collisions, both $t\bar{t}H$ and $t\bar{t}Z$ productions reach a maximum between about $\sqrt{s} = 500$ GeV to 750 GeV, and then fall gradually as $\sqrt{s}$ increases further. There are two main factors for this feature: one is the phase space factor and another is the $1/s$ factor in the $s$-channel $\gamma$ or $Z$ propagator. As $\sqrt{s}$ first increases from 500 GeV, the phase space factor increases; but as $\sqrt{s}$ increases further, the increase in phase space factor is offset by the $1/s$ decrease of the propagator. Roughly, $t\bar{t}Z$ production in $e^+e^-$ collisions is about a factor of 2 to 5 larger than production of $t\bar{t}H$. Consequently, if $m_H$ is close to $m_Z$ the $t\bar{t}H$ signal could be difficult to identify due to the potential $t\bar{t}Z$ background.

On the other hand, the cross sections for $\gamma\gamma \rightarrow t\bar{t}H$ and $t\bar{t}Z$ start off very small at $\sqrt{s} = 500$ GeV because they are very limited by the luminosity function $F_{\gamma/e}(x)$ at this energy. But both increase gradually as $\sqrt{s}$ increases, because a growing range of $x$ is available and there is no propagator contributing a factor $1/s$ as in the corresponding case of $e^+e^-$ collisions. However, $\sigma(\gamma\gamma \rightarrow t\bar{t}H)$ begins to flatten out after $\sqrt{s} = 1.5$ TeV. For both values of $m_t$, $t\bar{t}H$ production is larger than $t\bar{t}Z$ production at lower energies. But $t\bar{t}Z$ increases above $t\bar{t}H$ at about $\sqrt{s} = 1(2)$ TeV for $m_t = 120(150)$ GeV. For $m_t = 150$ GeV we have a $t\bar{t}H$ signal larger than the potential $t\bar{t}Z$ background for the entire range of $\sqrt{s}$ from 0.5 to 2 TeV. This is an important advantage of $\gamma\gamma$ collisions over $e^+e^-$ collisions for directly probing the $g_{tHH}$ Yukawa coupling. For $\sqrt{s}$ from 0.5 TeV to about 1.1 TeV, the $e^+e^- \rightarrow t\bar{t}H$ cross sections are larger than the $\gamma\gamma \rightarrow t\bar{t}H$ cross sections. However, for this range of $\sqrt{s}$, the potential background from $t\bar{t}Z$ production is also greater in $e^+e^-$ collisions. As $\sqrt{s}$ increases further $\gamma\gamma$ collisions are more advantageous both because $\sigma(\gamma\gamma \rightarrow t\bar{t}H)$ is larger and because there is a smaller $t\bar{t}Z$ background.

In Fig. 4, we plot the variation of total cross sections with Higgs mass $m_H$ for the range 60–140 GeV and $m_t = 150$ GeV at $\sqrt{s} = 1$ and 2 TeV. Of expected, $t\bar{t}H$ production in both $e^+e^-$ and $\gamma\gamma$ collisions decreases with increasing $m_H$, simply because less phase space is available. In contrast $\gamma\gamma \rightarrow t\bar{t}Z$ is independent of $m_H$ and the effect of $m_H$ on $e^+e^- \rightarrow t\bar{t}Z$ is negligible since the Higgs-exchange diagram is insignificant for this range of $m_H$.

In Fig. 5 we show the dependence of the total cross sections on the top-quark mass $m_t$ for $m_H = 90$ GeV at $\sqrt{s} = 1$ and 2 TeV. Two factors dominate the main features of these curves:
phase space and $g_{ttH}$ coupling. The coupling $g_{ttH}$ grows linearly with increasing $m_t$. Therefore both $e^+e^-$, $\gamma\gamma \rightarrow t\bar{t}Z$ decrease as $m_t$ increases simply because less phase space becomes available, and the effect is more pronounced at the smaller value of $\sqrt{s}$. On the other hand, at $\sqrt{s} = 1$ and 2 TeV both $e^+e^-$, $\gamma\gamma \rightarrow t\bar{t}H$ productions are enhanced as $m_t$ increases because the increase in the coupling $g_{ttH}$ is more important than the decreasing phase space. At $\sqrt{s} = 1$ TeV, $\gamma\gamma \rightarrow t\bar{t}H$ however begins to flatten out after about $m_t=160$ GeV.

IV. CONCLUSION

The Next Linear Colliders (NLC) will have center-of-mass energies from 0.5 to 2 TeV. The Yukawa coupling $g_{ttH}$ can be probed directly via $t\bar{t}H$ production, although there is a potential background from $t\bar{t}Z$ production if $m_H$ lies close to $m_Z$. For $\sqrt{s}$ from 0.5 to 1 TeV, the $e^+e^- \rightarrow t\bar{t}H$ cross section (1.5–3 fb) is larger than $\sigma(\gamma\gamma \rightarrow t\bar{t}H)$, but so is the potential background: $\sigma(e^+e^- \rightarrow t\bar{t}Z) \sim 3$–6 fb. For $\sqrt{s} \gtrsim 1$ TeV $\gamma\gamma$ collisions provide a better approach than $e^+e^-$ collisions since the cross section ($\sim 1.5$–2 fb) is larger and there is less potential background from $t\bar{t}Z$ production ($\sim 0.5$–2.5 fb). For a yearly luminosity of 10 fb$^{-1}$, $m_t = 150$ GeV and $m_H = 90$ GeV we have about 14(22) $t\bar{t}H$ events and about 6(21) $t\bar{t}Z$ events in $\gamma\gamma$ collisions realized by the laser back-scattering method at 1(2) TeV $e^+e^-$ colliders.

Added Note: after completing this work, we came across a paper by E. Boos et al. [23]. On the part of overlapping, our results agree with theirs.
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APPENDIX A

This appendix gives the formulas for the Feynman amplitudes of the subprocesses $\gamma\gamma \rightarrow t\bar{t}H$ and $\gamma\gamma \rightarrow t\bar{t}Z$ (see Fig. 2). Defining some notations as follows:

$$D^t(k) = \frac{1}{k^2 - m_t^2},$$  \hspace{1cm} (A1)
$$g^V(t) = g^V_v(t) + g^V_a(t)\gamma^5, \quad \text{where } V = \gamma, Z.$$  \hspace{1cm} (A2)

The couplings are given by

$$g^Z_v(f) = g_Z(\frac{T_3}{2}f - Q_f),$$
$$g^Z_a(f) = -g_Z T_3 f,$$
$$g^\gamma_v(f) = g\sin\theta_W Q_f,$$
$$g^\gamma_a(f) = 0,$$  \hspace{1cm} (A3)

where $g_Z = g/\cos\theta_W$ and $f$ is a fermion.

The amplitudes for the $\gamma(p_1)\gamma(p_2) \rightarrow t(q_1)\bar{t}(q_2)H(k_1)$ are given by

$$\mathcal{M}^{(a)} = -\frac{gm_t}{2m_W}D^t(q_1 + k_1)D^t(p_2 - q_2)$$
$$\times \bar{u}(q_1)(\hat{q}_1 + \not{k}_1 + m_t)f(p_1)g^\gamma(t)(\hat{p}_2 - \not{q}_2 + m_t)f(p_2)g^\gamma(t)v(q_2),$$  \hspace{1cm} (A4)

$$\mathcal{M}^{(b)} = -\frac{gm_t}{2m_W}D^t(q_1 - p_1)D^t(p_2 - q_2)$$
$$\times \bar{u}(q_1)f(p_1)g^\gamma(t)(\hat{q}_1 - \not{p}_1 + m_t)(\hat{p}_2 - \not{q}_2 + m_t)f(p_2)g^\gamma(t)v(q_2),$$  \hspace{1cm} (A5)

$$\mathcal{M}^{(c)} = -\frac{gm_t}{2m_W}D^t(q_1 - p_1)D^t(q_2 + k_1)$$
$$\times \bar{u}(q_1)f(p_1)g^\gamma(t)(\hat{q}_1 - \not{p}_1 + m_t)f(p_2)g^\gamma(t)(-\not{q}_2 - \not{k}_1 + m_t)v(q_2).$$  \hspace{1cm} (A6)
with the addition of the terms by interchanging $\gamma(p_1) \leftrightarrow \gamma(p_2)$. The Feynman amplitudes for $\gamma(p_1)\gamma(p_2) \to t(q_1)\bar{t}(q_2)Z(k_1)$ are given by

$$M^{(a)} = -D^t(q_1 + k_1)D^t(p_2 - q_2) \times \bar{u}(q_1)f(k_1)g^Z(t)(\not{q}_1 + \not{k}_1 + m_t)f(p_1)g^\gamma(t)(\not{p}_2 - \not{q}_2 + m_t)f(p_2)g^\gamma(t)v(q_2),$$  \hspace{1cm} (A7)

$$M^{(b)} = -D^t(q_1 - p_1)D^t(p_2 - q_2) \times \bar{u}(q_1)f(p_1)g^\gamma(t)(\not{q}_1 - \not{p}_1 + m_t)f(k_1)g^Z(t)(\not{p}_2 - \not{q}_2 + m_t)f(p_2)g^\gamma(t)v(q_2),$$  \hspace{1cm} (A8)

$$M^{(c)} = -D^t(q_1 - p_1)D^t(q_2 + k_1) \times \bar{u}(q_1)f(p_1)g^\gamma(t)(\not{q}_1 - \not{p}_1 + m_t)f(p_2)g^\gamma(t)(\not{q}_2 - \not{k}_1 + m_t)f(k_1)g^Z(t)v(q_2),$$  \hspace{1cm} (A9)

with the addition of the terms by interchanging $\gamma(p_1) \leftrightarrow \gamma(p_2)$. 


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FIGURES

FIG. 1. The energy spectrum $E_{\gamma/e}(x)$ of the back-scattered photon versus the energy fraction $x$ of the incident electron being carried away by the back-scattered photon.

FIG. 2. Contributing Feynman diagrams for the process $\gamma\gamma \to t\bar{t}H$. Cross diagrams by interchanging the two incoming photons are not shown.

FIG. 3. Total cross sections versus center-of-mass energy of the parent $e^+e^-$ collider, for $m_H = 90$ GeV and $m_t = (a) 120$, (b) 150 GeV. The subprocesses $\gamma\gamma \to t\bar{t}H$ (solid), $t\bar{t}Z$ (dashed); $e^+e^- \to t\bar{t}H$ (dotted), $t\bar{t}Z$ (dash-dotted) are shown.

FIG. 4. Total cross sections versus the mass $m_H$ of Higgs-boson for $m_t = 150$ GeV, and $\sqrt{s} = (a)$ 1 TeV, and (b) 2 TeV. The subprocesses $\gamma\gamma \to t\bar{t}H$ (solid), $t\bar{t}Z$ (dashed); $e^+e^- \to t\bar{t}H$ (dotted), $t\bar{t}Z$ (dash-dotted) are shown.

FIG. 5. Total cross sections versus the top-quark mass $m_t$ for $m_H = 90$ GeV and $\sqrt{s} = (a)$ 1 TeV, and (b) 2 TeV. The subprocesses $\gamma\gamma \to t\bar{t}H$ (solid), $t\bar{t}Z$ (dashed); $e^+e^- \to t\bar{t}H$ (dotted), $t\bar{t}Z$ (dash-dotted) are shown.