Detection of the heavy Higgs boson at $\gamma\gamma$ colliders

David Bowser-Chao* and Kingman Cheung**

* Center for Particle Physics, University of Texas at Austin, Austin, Texas 78712, USA
** Dept. of Physics & Astronomy, Northwestern University, Evanston, Illinois 60208, USA

ABSTRACT

We consider the possibility of detecting a heavy Higgs boson ($m_H > 2m_Z$) in proposed $\gamma\gamma$ colliders through the semi-leptonic mode $\gamma\gamma \rightarrow H \rightarrow ZZ \rightarrow q\bar{q}\ell^+\ell^-$. We show that due to the non-monochromatic nature of the photon beams produced by the laser-backscattering method, the resultant cross section for Higgs production is much smaller than the on-resonance cross section and generally decreases with increasing collider energy. Although continuum ZZ production is expected to be negligible, we demonstrate the presence of and calculate sizeable backgrounds from $\gamma\gamma \rightarrow \ell^+\ell^-Z$, $qq\bar{q}$, with $Z \rightarrow q\bar{q}$, $\ell^+\ell^-$, respectively, and $\gamma\gamma \rightarrow t\bar{t} \rightarrow b\bar{b}\ell^+\ell^-\nu\bar{\nu}$. This channel may be used to detect a Higgs of mass $m_H$ up to around 350 GeV at a 0.5 TeV $e^+e^-$ collider, assuming a nominal yearly luminosity of 10–20 fb$^{-1}$. 
I. INTRODUCTION

The next generation $pp$, $e^+e^-$, and $ep$ colliders will search for the Higgs boson [1] over a wide range of Higgs masses $m_H$. Detection of a heavy Higgs ($m_H > 2m_Z$) is considered feasible at the Superconducting Super Collider (SSC) or CERN Large Hadron Collider (LHC) through the gold-plated channel, $H \rightarrow ZZ \rightarrow \ell^+\ell^-\ell^+\ell^-$ [2], and also at the Next Linear $e^+e^-$ Colliders (NLC) through the decay of $H \rightarrow WW$, $ZZ \rightarrow (jj)(jj)$ [3]. Beyond the discovery of the Higgs boson and the determination of its mass, the measurement of its coupling to photons is desirable for the information afforded on the Yukawa coupling of the Higgs to the top quark [4] and the existence of ultraheavy new particles [5]. For a heavy Higgs, however, the branching ratio for $H \rightarrow \gamma\gamma$ is too small for accurate experimental measurement of the $H\gamma\gamma$ coupling at the SSC, LHC or NLC.

Proposed $\gamma\gamma$ colliders [6,7], based on the underlying next-generation linear $e^+e^-$ colliders, provide both a means of detecting a heavy Higgs [8] as well as obtaining information on the $H\gamma\gamma$ coupling, through direct Higgs production via photon-photon fusion. The Higgs may be detected through $H \rightarrow ZZ$, $W^+W^-$ and $t\bar{t}$. The latter two channels suffer from immense tree-level continuum backgrounds [9,10], and the top-quark has not yet been found, so we concentrate on the $H \rightarrow ZZ$ decay channel. Previous studies of Higgs resonance production [4,5,11–13] at $\gamma\gamma$ colliders, however, have noted the absence of a tree-level background to $\gamma\gamma \rightarrow H \rightarrow ZZ$. The signal is itself a one-loop process, but should be of order $\alpha$ less-suppressed than the production of $ZZ$ through box diagrams. Furthermore, the signal is almost entirely restricted to the interval $m_H \pm \Gamma_H$ in the invariant mass spectrum of $m(ZZ)$ because the Higgs width is very narrow for $m_H \lesssim 400$ GeV; for the visible decay channels, we can further reduce this background by requiring $m(ZZ)$ to be around the Higgs peak. If the continuum production represented the sole background, detection of the Higgs via $H \rightarrow ZZ$ would depend only on the total event rate. We shall demonstrate, however, the presence of additional backgrounds, which come into play when detection of the gauge boson pair is taken into account. The purpose of this paper is to investigate the feasibility of detecting
$H \rightarrow ZZ$ in light of these backgrounds, for the various decay modes of the gauge bosons.

As shown below, the non-monochromatic nature of the photon beams drastically reduces the Higgs production cross section from its on-resonance value. The “gold-plated” detection mode, where $H \rightarrow ZZ \rightarrow \ell^+\ell^-\ell^+\ell^-$, is thus marginal for all but very high collider luminosities because of the small branching fractions incurred. On the other hand, previous authors [5,13] have pointed out that the most abundant decay channel, where both $Z$ bosons decay hadronically, is obscured by the huge continuum background from $\gamma\gamma \rightarrow WW$ [9], where the $W$ bosons decay hadronically. The leptonic decay with neutrinos, where $H \rightarrow ZZ \rightarrow \ell^+\ell^-\nu\bar{\nu}$, has been considered [5,12,13], but we note that the leptonic decay of $\gamma\gamma \rightarrow W^+W^- \rightarrow \ell^+\nu\ell^-\bar{\nu}$ presents an overwhelming background, even after requiring the lepton pair to reconstruct a $Z$ boson. In the same vein, the decay of $WW \rightarrow q\bar{q}\ell\nu$ presents a large, detector-dependent background to the signature of $\gamma\gamma \rightarrow H \rightarrow ZZ \rightarrow q\bar{q}\nu\bar{\nu}$. Before any cuts, this background turns out to be about four orders of magnitude larger than the signal [9]. The viability of this channel crucially depends on the efficiency of a cut on $m(q\bar{q})$. A less important, though helpful, measure would be to additionally veto events with an isolated charged lepton. This mode furthermore precludes direct reconstruction of the Higgs mass.

This paper thus considers the feasibility of observing the remaining decay channel:

$$\gamma\gamma \rightarrow H \rightarrow ZZ \rightarrow \ell^+\ell^-q\bar{q},$$  \hspace{1cm} (1)$$

whereas limits in the reconstruction of either gauge boson lead to backgrounds from:

$$\gamma\gamma \rightarrow q\bar{q}Z \rightarrow q\bar{q}\ell^+\ell^-$, \hspace{1cm} (2)$$

$$\gamma\gamma \rightarrow \ell^+\ell^-Z \rightarrow \ell^+\ell^-q\bar{q},$$  \hspace{1cm} (3)$$

$$\gamma\gamma \rightarrow t\bar{t} \rightarrow b\bar{b}\ell^+\ell^-\nu\bar{\nu}_\ell,$$  \hspace{1cm} (4)$$

where the charged lepton and quark pairs are required to reconstruct the $Z$ mass to a given resolution (we note that process (4) is another potential background to $H \rightarrow ZZ \rightarrow q\bar{q}\nu\bar{\nu}$).
We will show that with selective acceptance cuts on the charged leptons and quarks, we can reduce these backgrounds to a manageable level.

The organization of this paper is as follows: we briefly describe the laser back-scattering method and the corresponding photon luminosity in Sec. II, following which Sec. III details the calculation of the signal and background processes. Finally, we discuss our results in Sec. IV, and summarize the conclusions of our analysis in Sec. IV.

II. LASER BACK-SCATTERING METHOD

Hard $\gamma\gamma$ collisions at $e^+e^-$ machines can be produced by directing a low energy (a few $eV$) laser beam at a very small angle $\alpha_0$, almost head to head, to the incident electron beam. Through 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. Further technical details may be found in Refs. [6]. Another possibility is to use the beamstrahlung effect [7] but this method produces photons mainly in the soft region [6], and depends critically on the beam structure [7]. For the production of a heavy Higgs we need a fairly high energy threshold of the two incoming photons. Therefore we shall limit our calculations to $\gamma\gamma$ collisions produced by the laser back-scattering method.

A. Photon Luminosity

We use the energy spectrum of the back-scattered photon given by [6]

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

where

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

(5)

(6)
\( \xi = 4E_0\omega_0/m_e^2 \), \( \omega_0 \) is the energy of the incident laser photon, \( x = \omega/E_0 \) is the fraction of the incident electron’s energy carried by the back-scattered photon, and the maximum value \( x_{\text{max}} \) is given by

\[
x_{\text{max}} = \frac{\xi}{1 + \xi}.
\]

(7)

It is seen from Eq. (5) and (6) that the portion of photons with maximum energy grows with \( E_0 \) and \( \omega_0 \). A large \( \omega_0 \), however, should be avoided so that the back-scattered photon does not interact with the incident photon and create unwanted \( e^+e^- \) pairs. The threshold for \( e^+e^- \) pair creation is \( \omega_0 > m_e^2 \), so we require \( \omega_{\text{max}}\omega_0 \lesssim m_e^2 \). Solving \( \omega_{\text{max}}\omega_0 = m_e^2 \), we find

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

(8)

For the choice \( \xi = 4.8 \) one finds \( x_{\text{max}} \simeq 0.83 \), \( D(\xi) \simeq 1.8 \), and \( \omega_0 = 1.25(0.63) \) eV for a 0.5(1) TeV \( e^+e^- \) collider. Here we have assumed that the electron, positron and back-scattered photon beams are unpolarized. We also assume that, on average, the number of back-scattered photons produced per electron is 1 (i.e., the conversion coefficient \( k \) equals 1).

III. SIGNAL & BACKGROUND CALCULATION

In calculating the results presented below, we folded the photon distribution function into the hard scattering cross section \( \hat{\sigma} \) for each subprocess. The total cross section \( \sigma \) is given by

\[
\sigma(s) = \int_{\tau_{\text{min}}}^{x_{\text{max}}^2} d\tau \int_{\tau/x_{\text{max}}}^{x_{\text{max}}} dx_1 \frac{d\tau/x}{x_1} F_{\gamma/e}(x_1) F_{\gamma/e}(\tau/x_1) \hat{\sigma}(\hat{s} = \tau s),
\]

(9)

where

\[
\tau_{\text{min}} = \frac{(M_{\text{final}})^2}{s},
\]

(10)

and \( M_{\text{final}} \) is the sum of masses of the final state particles.
A. $\gamma\gamma \rightarrow H \rightarrow ZZ$

The contributing Feynman diagrams are depicted in Fig. 1 (the subsequent decay of the gauge bosons, $ZZ \rightarrow \ell^+\ell^- q\bar{q}$, is not shown). All massive charged particles contribute to the loop $[1]$; in particular, this process could be used to probe the Higgs sector with a heavy top quark $[4]$, or to detect the presence of new ultra-heavy fermions $[5]$ or non-SM $W'$ charged gauge bosons. In this paper, we assume only SM particles in the triangular loop; in particular, we have used a top-quark mass of 150 GeV.

The hard scattering cross section for $\gamma\gamma \rightarrow H \rightarrow ZZ$, for monochromatic photons with $m(\gamma\gamma) = \sqrt{s}$, is given by:

$$\hat{\sigma}(\hat{s}) = \frac{8\pi\Gamma(H^* \rightarrow \gamma\gamma)\Gamma(H^* \rightarrow ZZ)}{(\hat{s} - m_{H^*}^2)^2 + \Gamma_{H^*}^2 m_{H^*}^2},$$  \hspace{1cm} (11)

where the partial widths in the numerator are those of a virtual Higgs of mass $m_{H^*} = \sqrt{s}$. The on-resonance cross section $\sigma_0$ is simply given by $\sigma_0 = \hat{\sigma}(m_{H^*}^2)$. We have calculated the total cross section $\sigma(\gamma\gamma \rightarrow H \rightarrow ZZ)$ by folding in Eq. (11) with the photon luminosity functions as in Eq. (9).

It is instructive to apply the narrow width approximation to this cross section. For the range of $m_H$ considered here ($m_H < \sim 400$ GeV), the true width $\Gamma_H$ is quite small. Then the partial widths can be approximated by their on-shell values, obtaining:

$$\sigma(s) = \sigma_0 \left( \frac{\pi m_H \Gamma_H}{s} \right) \int_{m_H^2/(x_{\text{max}}^2 s)}^{x_{\text{max}}} F_{\gamma/e}(x_1) F_{\gamma/e}(\tau/x_1) \frac{dx_1}{x_1},$$  \hspace{1cm} (12)

where $\tau = m_H^2/s$ and $\sqrt{s}$ is the center-of-mass energy of the parent $e^+e^-$ collider. From Eq. (12) we anticipate that the total cross section should roughly decrease with increasing collider energy $\sqrt{s}$. However, as the maximum energy of the photon-photon collider ($x_{\text{max}}\sqrt{s}$) just passes $m_H$, $\sigma$ increases with $\sqrt{s}$ because of enhancement by the factor $\log(x_{\text{max}}^2 s/m_H^2)$ from the integration of $x_1$.

The photon is known to have anomalous quark and gluon contents $[14]$. The Higgs can thus also be produced by the gluon-gluon fusion, where the contributing Feynman diagrams...
are similar to those of photon-photon fusion, but with only quarks in the triangular loop. Since we are searching for a heavy Higgs ($m_H > 2m_Z$), the threshold value for $\sqrt{\tau}$ (equal to $2m_Z/\sqrt{s}$) is about 0.1 and 0.4 at $\sqrt{s} = 1.5$ and 0.5 TeV, respectively. For this range of $\sqrt{\tau}$, the effective gluon-gluon luminosity is extremely small [10], so that the contribution from gluon-gluon fusion is negligible. Furthermore, gluon-photon fusion does not contribute at one-loop level because the Higgs is color neutral.

To obtain the total cross section for the decay channel $\gamma\gamma \rightarrow H \rightarrow ZZ \rightarrow \ell^+\ell^- q\bar{q}$, we simply multiply by the relevant branching fractions. To properly assess the effect of cuts on this mode, however, we included full spin-correlation of the $Z$-boson decay products. We thus alternately replaced the width $\Gamma(H^* \rightarrow ZZ)$ in Eq. (11) with $\Gamma(H^* \rightarrow Z_LZ_L)$ and $\Gamma(H^* \rightarrow Z_TZ_T)$, and generated the respective angular distribution of the decay products.

**B. $\gamma\gamma \rightarrow q\bar{q}Z$ and $\ell^+\ell^- Z$**

The contributing Feynman diagrams for the process $\gamma\gamma \rightarrow q\bar{q}Z$ are shown in Fig. 2(a). Those for $\gamma\gamma \rightarrow \ell^+\ell^- Z$ are obtained with the replacement of $q$ with $\ell$. Here again we have not depicted the subsequent decay of the gauge bosons into either a quark or charged-lepton pair. The expressions for Feynman amplitudes for the similar process $\gamma\gamma \rightarrow t\bar{t}Z$ have been presented in Ref. [15]. Those for $q\bar{q}Z$ and $\ell^+\ell^- Z$ are easily obtained by replacing the top-quark with the lighter quarks ($u, d, s, c, b$) and charged leptons ($e, \mu$), and substituting with the corresponding $g_{Zff}$ and $g_{\gamma ff}$ couplings. In our calculations of these backgrounds, we have set all quark and lepton masses to $m_f = m_q = m_l = 0.1$ GeV. In the next section, we will show that with our acceptance cuts the cross sections of these backgrounds are independent of the choice of $m_f$. We have also assumed that we can distinguish a tau from a muon, electron, or quark.

The processes of Eqs. (9) and (10) have the same final state particles $q\bar{q}\ell^+\ell^-$ as the signal. These backgrounds mimic the signal when the invariant-mass of the continuum $q\bar{q}$ or $\ell^+\ell^-$ pair is within detector resolution of the $Z$ mass. We use a rather conservative mass constraint
\[ |m(q\bar{q} \text{ or } \ell^+\ell^-) - m_Z| < 8 \text{ GeV}, \tag{13} \]

to reconstruct the “fake Z”. We have included full spin correlation for the subsequent decay of the real Z into either \( q\bar{q} \) or \( \ell^+\ell^- \) pair.

In contrast to the signal, we must consider the photon-gluon contribution to \( q\bar{q}Z \) production. However, the gluon distribution function inside photon has large uncertainties because of limited experimental data. We have employed two different parameterizations, Levy-Abramowicz-Charchula set III (LAC3) [16] and Drees-Grasie (DG) [17], for the photon structure functions. We have set the scale \( Q^2 = \hat{s}/4 \) for both the photon structure functions and \( \alpha_s \) (evaluated to second order), and \( \Lambda_4 \) to be 0.2 and 0.4 GeV for the LAC3 and DG sets, respectively.

C. \( \gamma\gamma \to t\bar{t} \)

The contributing Feynman diagrams are shown in Fig. 2(b). The top-quark pair decays into \( b\bar{b}WW \); if the \( W \) bosons decay leptonically into \( \ell^+\ell^-\nu\bar{\nu}_\ell (\ell = e, \mu) \), this process could be mistaken as signal if both \( m(q\bar{q}) \) and \( m(\ell^+\ell^-) \) are in the vicinity of the Z mass. We have chosen \( m_t = 150 \) GeV for illustrative purposes. Even with a possible doubling due to smaller top-quark mass or enhancement from threshold effects [10], this process turns out to be negligible due to the additional branching fractions for the \( W \) decays and the simultaneous Z-mass constraints of Eq. (13) on both \( m(q\bar{q}) \) and \( m(\ell^+\ell^-) \). We have included full spin correlations in the decays of \( t\bar{t} \) and the subsequent decays of \( W^+W^- \).

IV. RESULTS & DISCUSSION

The dependence of the cross section \( \hat{\sigma} \) for \( \gamma\gamma \to H \to ZZ \) on the center-of-mass energies \( \sqrt{s_{\gamma\gamma}} \) of the incoming photons is shown in Fig. 3(a) for a range of Higgs masses \( m_H \) from 200–500 GeV, assuming monochromatic photons; this figure agrees with the results in Ref. [3]. As discussed above, however, laser back-scattering produces photons with a spread in energies,
so that we must fold $\sigma_0$ in with the effective photon-photon luminosities to obtain the actual cross section. We show the dependence of the total cross section $\sigma(e^+e^- \rightarrow \gamma\gamma \rightarrow H \rightarrow ZZ)$ on the center-of-mass energies $\sqrt{s_{e^+e^-}}$ of the parent $e^+e^-$ collider for the same range of $m_H$ in Fig. 3(b). As previously discussed, the resultant cross sections are greatly reduced from their on-resonance values $\sigma_0$, and generally fall with $\sqrt{s_{e^+e^-}}$. The cross sections, however, increase in the small range of $\sqrt{s_{e^+e^-}} \gtrsim m_H$, which is most apparent for $m_H \gtrsim 400$ GeV. We also show the dependence of the actual cross section on the Higgs-boson mass $m_H$ for $\sqrt{s_{e^+e^-}} = 0.5$ and 1 TeV in Fig. 3(c). We can see that the maximum cross section occurs at slightly above $m_H = 200$ GeV, then falls quite sharply with further increase in $m_H$. It should be noted that the smallness of these cross sections precludes the use of the gold-plated mode $H \rightarrow ZZ \rightarrow \ell^+\ell^-\ell^+\ell^-$, for which the optimal case of $m_H = 200$ GeV at $\sqrt{s_{e^+e^-}} = (0.3)0.5$ TeV and an integrated luminosity of 100 fb$^{-1}$ would yield only about (13) 10 events. Furthermore, the drop in direct Higgs production with increasing collider energies suggests that associated production of the Higgs with a $t\bar{t}$ pair \cite{14} or $W^+W^-$ pair \cite{17} would be more appropriate for $\sqrt{s_{e^+e^-}} \gtrsim 1.5$ TeV.

With only the mass constraints of Eq. (13), the $t\bar{t}$ background of Eq. (4), as remarked above, turns out to be very small, totalling about $\sim 0.08$ fb for $m_t = 150$ GeV. The major backgrounds come from the processes of Eqs. (3) and (3), which total to about 30–90 fb for the fermion masses $m_f = 100$ MeV to 0.5 MeV. The cross section for $\gamma\gamma \rightarrow \ell^+\ell^-Z$ is approximately ten times larger than $\gamma\gamma \rightarrow q\bar{q}Z$ because of the difference between the branching fractions of $Z$ into quarks and charged leptons, respectively, even after including the contribution to the latter process from photon-gluon fusion. We expect, due to our conservative cuts imposed on the lepton pair invariant mass, that these backgrounds could in practice be smaller (in accordance with actual detector resolution for $m(\ell^+\ell^-)$).

Because the production of $f\bar{f}Z$, where $f = \ell, q$, would be collinearly divergent in the beam direction for $m_f$ equal to zero, we expect a cut on the angles between the decay products and the beam direction to be efficacious; some cut of this nature would of course be required due to limited detector coverage. This expectation is clearly borne out by Fig. 3.
which shows the dependence of the integrated cross sections $\sigma(z < z_{\text{cut}})$ on $z_{\text{cut}}$, where $z$ is defined by

$$z = \max \{|\cos \theta_i|\}, \quad i = q, \bar{q}, \ell^+\ell^-$$

(14)

It can be seen that the background is rather sharply-peaked in the forward and backward regions. We therefore require

$$z_{\text{cut}} = 0.95$$

(15)

which corresponds to requiring an angle of about $8^\circ$ from the beam-pipe. This cut has the effect of drastically reducing the backgrounds of Eq. (2) and (3) but only mildly affecting the signal and the $t\bar{t}$ background of Eq. (4). As described above, we have taken the fermion-masses $m_f = 0.1 \text{ GeV}$ in presenting the background calculation of Eq. (2) and (3). We have confirmed that the results with the above cut on $z$ are independent of the fermion-mass $m_f$ for $m_f < \sim 1 \text{ GeV}$.

Because all the decay products can be detected after requiring separation from the beam-pipe by the cut on $z$ above, we can reconstruct the total invariant mass of the quark and charged-lepton pairs. In Fig. 5 we show the dependence of the differential cross section $d\sigma/dm(ZZ)$ on the invariant mass of the $ZZ$ pair, after the acceptance cuts of $z < 0.95$ and the mass constraint of Eq. (13). The Higgs-boson signal forms a sharp peak for $m_H \lesssim 300 \text{ GeV}$ and a less sharp but still distinguishable peak for $m_H = 350$ and $400 \text{ GeV}$. We define the signal as the cross section under the Higgs peak in the interval

$$m_H \pm \Gamma, \quad \text{where } \Gamma = \max(\Gamma_{\text{resolution}}, \Gamma_H),$$

(16)

and we take a rather conservative $\Gamma_{\text{resolution}} = 10 \text{ GeV}$. The same cut is applied to each background process. With better resolution in $m(ZZ)$, of course, the background can be further reduced.

Cross sections for the signal and backgrounds for various Higgs masses and collider energies are presented in Table I. We assume that the Higgs-boson signal can be discovered
if there are six or more events under the peak and the significance \( S/\sqrt{B} \) of this signal is greater than four. The significances for integrated luminosities of 10, 20, 50 and 100 fb\(^{-1}\) are presented in Table [I]. It is clear that luminosity, rather than collider energy \( \sqrt{s_{e^+e^-}} \), is the key to detecting a heavy Higgs through photon-photon fusion production. Furthermore, detection of lower Higgs masses \((m_H \lesssim 350 \text{ GeV})\) is easier in this mode at lower collider energies. With a luminosity of 10–20 fb\(^{-1}\), a \( \sqrt{s_{e^+e^-}} = 0.5 \text{ TeV} \) collider could detect Higgs bosons for \( m_H \) up to around 350 GeV. A higher collider energy would necessitate a correspondingly higher luminosity. A luminosity of 100 fb\(^{-1}\) would be required to probe Higgs masses up to 400 GeV.

\[ \text{V. CONCLUSION} \]

We have shown that direct production of a heavy Higgs at \( \gamma\gamma \) colliders offers a feasible discovery mode, provided that the energy of the underlying \( e^+e^- \) collider is not too much greater than \( m_H \). Due to the backgrounds principally arising from the immense production of \( W^+W^- \) pairs, the optimal detection channel is \( H \to ZZ \to q\bar{q}\ell^+\ell^- \). Although the continuum background \( \gamma\gamma \to ZZ \) to this process is minimal, large backgrounds arise from the production of \( \gamma\gamma \to q\bar{q}Z, \ell^+\ell^-Z \). An angular cut requiring the quark and charged-lepton pairs to be away from the beam-pipe can drastically reduce these backgrounds. With this acceptance cut and the Z-mass constraints on the quark and charged-lepton pairs, we should be able to discover a heavy Higgs-boson signal for \( m_H \) up to 300 GeV at a \( \sqrt{s_{e^+e^-}} = 0.5 \text{ TeV} \) \( e^+e^- \) collider, assuming a nominal yearly luminosity of 10 fb\(^{-1}\), and up to 350 GeV for a luminosity of 20 fb\(^{-1}\). The upper limit of applicability of this mode is \( m_H \lesssim 400 \text{ GeV} \), for which one requires a luminosity of the order 100 fb\(^{-1}\).

\[ \text{ACKNOWLEDGMENTS} \]

We are grateful for helpful discussions with D. A. Dicus, C. Kao, and A. L. Stange. This work was supported by the U. S. Department of Energy, Division of High Energy
Physics, under Grants DOE-FG02-91-ER40684 (K. C.) and DOE-FG05-85ER40200 (D. B.-C.). Computing resources were provided in part by the University of Texas Center for High Performance Computing.
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FIGURES

FIG. 1. Contributing Feynman diagrams for the signal process $\gamma\gamma \to H \to ZZ$. The cross diagram of the incoming photons is not shown.

FIG. 2. Contributing Feynman diagrams for the background processes (a) $\gamma\gamma \to q\bar{q}(\ell\bar{\ell})Z$, and (b) $\gamma\gamma \to t\bar{t}$. Cross diagrams of the incoming photons are not shown.

FIG. 3. The dependence of the total cross section for the signal $\gamma\gamma \to H \to ZZ$ on (a) $\sqrt{s_{\gamma\gamma}}$ for monochromatic photons with $m_H = 200, 300, 400$, and $500$ GeV, (b) $\sqrt{s_{e^+e^-}}$ of the parent $e^+e^-$ collider for the same range of $m_H$, and (c) $m_H$ for $\sqrt{s_{e^+e^-}} = 0.5$ and $1$ TeV.

FIG. 4. The integrated cross section $\sigma(z < z_{\text{cut}})$ versus $z_{\text{cut}}$, where $z = \max\{ |\cos \theta_i| \}$, $i = q, \bar{q}, \ell^+\ell^-$, for the signal and the total background (indicated by the dashed curve) with $m_H = 200, 300, 400$ GeV and $m_t = 150$ GeV at $\sqrt{s_{e^+e^-}} = 0.5$ TeV. The only cut employed is $|m(q\bar{q} \text{ or } \ell^+\ell^-) - m_Z| < 8$ GeV. The curves extend to $z_{\text{cut}} = 0.999$.

FIG. 5. The differential cross section $d\sigma/dm(ZZ)$ versus the invariant mass $m(ZZ)$ for the signal with various Higgs masses for $m_t = 150$ GeV at $\sqrt{s_{e^+e^-}} = 0.5$ TeV, after the acceptance cut of $z < 0.95$ and the $Z$-mass constraints of $|m(q\bar{q} \text{ or } \ell^+\ell^-) - m_Z| < 8$ GeV. The sum of the backgrounds is indicated by the dashed curve.
TABLE I. Cross sections in fb for the signal ($\gamma\gamma \rightarrow H \rightarrow ZZ \rightarrow q\bar{q}\ell^+\ell^-$, $\ell = e, \mu$) and the backgrounds, for various values of $m_H$, at $\sqrt{s} = 0.5$, 1 and 1.5 TeV $e^+e^-$ colliders. The cross sections already include the branching fraction to the $q\bar{q}\ell^+\ell^-$ final state. The acceptance cuts are $z < 0.95$, $|m(\ell\ell \text{ or } qq) - m_Z| < 8$ GeV, and $m_H - \Gamma < m(ZZ) < m_H + \Gamma$, where $\Gamma$ is defined in Eq. (16).

| $m_H$ | $\gamma\gamma \rightarrow H$ | $\gamma\gamma \rightarrow \ell^+\ell^-$ | $\gamma\gamma \rightarrow q\bar{q}Z$ | $\gamma g \rightarrow q\bar{q}Z$ | $\gamma\gamma \rightarrow t\bar{t}$ |
|-------|-----------------|-----------------|-----------------|-----------------|-----------------|
|       | $\sqrt{s} = 0.5$ TeV | $\sqrt{s} = 1$ TeV | $\sqrt{s} = 1.5$ TeV |
|       | DG              | LAC3            | DG              | LAC3            | DG              | LAC3            |
| $\Gamma = \infty$ | —               | 1.7             | 0.13            | 0.0043          | 0.016           | 0.064           |
| 200   | 1.6             | 0.29            | 0.021           | 0.0016          | 0.0055          | 0.018           |
| 250   | 1.4             | 0.22            | 0.016           | 0.00042         | 0.0017          | 0.0099          |
| 300   | 0.92            | 0.13            | 0.0093          | 6.8$\times 10^{-5}$ | 0.00028         | 6.7$\times 10^{-4}$ |
| 350   | 0.38            | 0.12            | 0.0088          | 1.2$\times 10^{-5}$ | 4.2$\times 10^{-5}$ | 0.0            |
| 400   | 0.099           | 0.070           | 0.0051          | 6.4$\times 10^{-7}$ | 1.4$\times 10^{-6}$ | 0.0            |
| $\Gamma = \infty$ | —               | 0.82            | 0.059           | 0.016           | 0.042           | 0.080           |
| 200   | 0.52            | 0.084           | 0.0061          | 0.0038          | 0.0093          | 0.020           |
| 250   | 0.53            | 0.084           | 0.0061          | 0.0022          | 0.0055          | 0.012           |
| 300   | 0.42            | 0.055           | 0.0040          | 0.00087         | 0.0023          | 0.0027          |
| 350   | 0.21            | 0.063           | 0.0045          | 0.00056         | 0.0016          | 0.0011          |
| 400   | 0.11            | 0.071           | 0.0052          | 0.00038         | 0.0011          | 1.2$\times 10^{-4}$ |
| $\Gamma = \infty$ | —               | 0.40            | 0.030           | 0.019           | 0.045           | 0.045           |
| 200   | 0.23            | 0.038           | 0.0027          | 0.0036          | 0.0081          | 0.013           |
| 250   | 0.23            | 0.037           | 0.0027          | 0.0024          | 0.0055          | 0.0068          |
| 300   | 0.20            | 0.025           | 0.0018          | 0.0012          | 0.0027          | 0.0023          |
| 350   | 0.11            | 0.032           | 0.0023          | 0.0010          | 0.0024          | 3.7$\times 10^{-4}$ |
| 400   | 0.060           | 0.037           | 0.0026          | 0.00086         | 0.0021          | 9.6$\times 10^{-5}$ |
TABLE II. Cross sections for the signal \((S)\) and the total background \((B)\) in fb, and the significance \((S/\sqrt{B})\) of the signal, for \(m_H = 200, 250, 300, 350, 400\) GeV at \(\sqrt{s} = 0.5, 1\) and 1.5 TeV \(e^+e^-\) colliders with various integrated luminosities.

| \(m_H\) | \(S\)  | \(B\)  | \(S/\sqrt{B}\) |
|---------|-------|-------|----------------|
|         | (a) \(\sqrt{s} = 0.5\) TeV |       |                |
| 200     | 1.6   | 0.34  | 8.7 12 19 27   |
| 250     | 1.4   | 0.25  | 8.9 12 20 28   |
| 300     | 0.92  | 0.14  | 7.8 11 17 25   |
| 350     | 0.38  | 0.13  | 3.3 4.7 7.4 11 |
| 400     | 0.099 | 0.075 | 1.1 1.6 2.6 3.6|
|         | (b) \(\sqrt{s} = 1\) TeV |       |                |
| 200     | 0.52  | 0.12  | 4.7 6.7 11 15  |
| 250     | 0.53  | 0.11  | 5.1 7.1 11 16  |
| 300     | 0.42  | 0.064 | 5.3 7.4 12 16  |
| 350     | 0.21  | 0.070 | 2.5 3.5 5.6 7.9|
| 400     | 0.11  | 0.077 | 1.3 1.8 2.8 4.0|
|         | (c) \(\sqrt{s} = 1.5\) TeV |       |                |
| 200     | 0.23  | 0.062 | 2.9 4.1 6.5 9.2|
| 250     | 0.23  | 0.052 | 3.2 4.5 7.1 10 |
| 300     | 0.20  | 0.032 | 3.5 5.0 7.9 11 |
| 350     | 0.11  | 0.037 | 1.8 2.5 4.0 5.7|
| 400     | 0.060 | 0.042 | 0.9 1.3 2.1 2.9|