Higgs or Neutral Vector Boson Production with a $W$ Pair in $\gamma\gamma$ Collisions

M. Baillargeon and F. Boudjema

Laboratoire de Physique Théorique ENSLAPP
Chemin de Bellevue, B.P. 110, F-74941 Annecy-le-Vieux, Cedex, France.

Abstract

Exploiting the fact that $W$ pair production in high-energy $\gamma\gamma$ collisions is very large, we use this process to trigger Higgs, $Z$ or photon radiation. We find that there are sizeable rising cross-sections for triple bosons production. At energies about $1 TeV$ the new mechanism for Higgs production becomes very competitive with the dominant Higgs production processes in $e^+e^-$ and $e\gamma$ reactions. The effect of different polarized photon spectra obtained through back-scattered laser light on the electron beam of a linear collider is investigated. We give a special attention to the search of the intermediate mass Higgs in $WWH$ production and discuss how to effectively suppress the backgrounds.

*On leave from Laboratoire de Physique Nucléaire, Université de Montréal, C.P. 6128, Succ. A, Montréal, Québec, H3C 3J7, Canada.
†URA 14-36 du CNRS, associé à l’E.N.S de Lyon, et au L.A.P.P. d’Annecy-le-Vieux.
1 Introduction

The ongoing intense activity in the physics potential offered by a high-energy $e^+e^-$ linear collider has stimulated a growing interest in the possibility of turning such a machine into a high-energy and high-luminosity $\gamma\gamma$ collider[1]. The large flux of very energetic photons is obtained through Compton backscattering of laser light on the single-pass electrons of the linear collider. Some of the principal attractions of running in this mode rest on the fact that this is a “democratic” means of producing all charged particles and that neutral scalar particles, notably the Higgs, can be produced as a resonance in the $s$-channel. The search for the Higgs in this mode has in fact gathered most attention.

For a TeV (or even a mid-TeV) $e^+e^-$ collider the $\gamma\gamma$ mode is a more efficient way of producing $W$ pairs. Indeed, the cross section for the process $\gamma\gamma \rightarrow W^+W^-$ is very large and does not decrease at high-energy due the spin-1 $t$-channel $W$ exchange[2]. Already at an effective $\sqrt{s}_{\gamma\gamma} \sim 400\text{GeV}$ the reaction reaches its plateau with a cross section of $\sim 80\text{pb}$. The asymptotic constant cross-section is $\sigma_{\text{asymp}}(\gamma\gamma \rightarrow W^+W^-) \sim 8\pi\alpha^2/M_W^2$. This is more than an order of magnitude larger than $W$ pair production in the usual $e^+e^-$ mode at the same centre-of-mass energy. The latter, as is known, decreases with energy, $\sigma_{\text{asymp}}(e^+e^- \rightarrow W^+W^-) \sim (\pi\alpha^2s_W^4/2s)\log(s/M_W^2)$. This also means that the currently discussed Next Linear $e^+e^-$ Collider operating at $500\text{GeV}$ with a yearly luminosity of $10-20\text{fb}^{-1}$ will produce about one million $W$-pairs in the $\gamma\gamma$ mode. One may then even contemplate using this reaction as a luminosity monitor. In this letter we exploit this process as a backbone reaction to which we “graft” one more additional boson. We will show that this reaction triggers a sizeable Higgs cross section and that $WWZ$ and $WW\gamma$ productions are even larger.

2 The use of a non-linear gauge at tree-level

The processes contributing to the triple boson production in $\gamma\gamma$ are shown in Fig. 1. Because of the relatively large number of diagrams an efficient way of calculating is almost mandatory. Calculating in the usual unitary gauge is rather awkward because of the cumbersome presence of the “longitudinal” mode (“$k_\mu k_\nu$” term) of the various $W$ propagators. A way out is to use a Feynman gauge. However, with the widespread choice of a linear gauge fixing term, this is done at the expense of having to deal with even more diagrams containing the unphysical Higgs scalars. An indisputable choice of gauge for photonic reactions or for processes involving a mixture of $W$’s and photons is to quantize with a non-linear gauge fixing term [3] and work with a parameter corresponding
to the ’t Hooft-Feynman gauge. For the processes at hand this means that one has the same number of diagrams as in the usual unitarity gauge save for the fact that we have no “longitudinal” mode to worry about and that there are no diagrams with unphysical scalars since the virtue of this choice is that the vertex with the photon, the $W$ and the unphysical Higgs field does not exist. Of course, one has to allow for small changes in the vertices which turn out to have an even more compact form than in the usual gauges. For instance, all diagrams where the quartic $WW\gamma\gamma$ vertex appears are identically zero when the two incident photons have opposite helicities ($J = \pm 2$). We foresee this choice to stand out for applications to $W$ dynamics at a future $\gamma\gamma$ collider as it has proved to be for one-loop weak bosons induced amplitudes for photonic processes [1].

With $S^\pm$ being the unphysical Higgs bosons, the $W^\pm$-part of the linear gauge fixing condition

$$\mathcal{L}^{Gauge-Fixing}_{\text{linear}} = -\xi^{-1}|\partial_\mu W^{\mu+} + i\xi M_W S^+|^2$$

is replaced by the “constraint”

$$\mathcal{L}^{Gauge-Fixing}_{\text{non-linear}} = -\xi^{-1}|(\partial_\mu + ie A_\mu + ig \cos \theta_W Z_\mu)W^{\mu+} + i\xi M_W S^+|^2$$

where $\theta_W$ is the usual weak mixing angle. We have taken $\xi = 1$.

Having reduced the number of diagrams by the gauge-fixing choice we have eased our computational task by calculating the full helicity amplitudes. These were, in turn, fed into a Monte-Carlo event generator which preformed the phase space integrals. We have checked that our amplitudes were gauge invariant both analytically and numerically. We will take $M_W = 80.1\text{GeV}$, $M_Z = 91.18\text{GeV}$ and $\sin^2 \theta_W = 0.232$.

### 3 Behaviour of the cross sections

We first present our results for an ideal $\gamma\gamma$ collider to explicitly exhibit the interesting behaviour of the various cross sections with $\gamma\gamma$ centre-of-mass energy. We will then include the effect of more realistic luminosity spectra.
3.1 $\gamma\gamma \rightarrow W^+W^-\gamma$

For the $WW\gamma$ final state, a cut on the (final) photon energy is required. One may also prefer to take a cut on the transverse momentum of the photon. With a fixed cut $p_T^\gamma > 20$ GeV for all centre-of-mass energies, the cross section increases with energy. At 500 GeV one reaches a cross section of about 1.3 pb. This is about 1.6% of the $WW$ cross section at the same energy. The $J_Z = 0$ obtained when both photons have the same helicity slightly dominates over the $J_Z = 2$ (1.5 pb versus 1.1 pb). At $\sqrt{s_{\gamma\gamma}} = 2$ TeV the cross section with the same $p_T^\gamma > 20$ GeV cut reaches 3.7 pb. The logarithmic ($\log^2 s$) growth can be understood on the basis that this cross section can be factorised in terms of $\gamma\gamma \rightarrow WW$, which is constant at asymptotic $M_{WW}$ invariant masses, times the final state photon radiator which contains the logarithmic $s$ dependence. We note that this logarithmic increase only concerns the production of transverse $W$. When both $W$ are longitudinal ($W_LW_L$) the cross section decreases. This can also be traced back to the fact that $\gamma\gamma \rightarrow W_L^+W_L^-$ decreases with energy. We find that the $W_LW_L$ fraction of all $W$'s is about only a 1% at 500 GeV, 0.3% at 1 TeV and a mere 0.07% at 2 TeV. It must be noted that the bulk of the cross section occurs when all final particles are produced at very small angles: this is a typical example of multiparticle production in the very forward region. For instance increasing the $p_T^\gamma$ cut and at the same time imposing a pseudorapidity cut on the photon, the $WW\gamma$ yield, as shown in Table 1., drops considerably, especially at higher energies. The reduction is even more dramatic when we put an isolation cut between all the particles and forcing them away from the beam. With these strictures the cross section decreases with energy (See Table 1.)

| $\sqrt{s_{\gamma\gamma}}$(TeV) | 0.5 | 1 | 1.5 | 2 |
|-------------------------------|-----|---|-----|---|
| type of cut                   |     |   |     |   |
| 1.                           | 1254| 2469| 3195| 3678 |
| 2.                           | 1254| 1434| 1258| 1050 |
| 2. and 3.                    | 1235| 1373| 1159| 930  |
| 2. and 4.                    | 201 | 86 | 47 | 32 |

1. $p_T^\gamma > 20$ GeV  
2. $p_T^\gamma > 40$ GeV $\times \sqrt{s_{\gamma\gamma}}$(in TeV)  
3. $|y^\gamma| < 2$  
4. $\cos$(between any two particles) $< 0.8$

Table 1: Cross section for $\gamma\gamma \rightarrow W^+W^-\gamma$ (in fb) at different $\sqrt{s_{\gamma\gamma}}$ including various cuts.

While the $W_LW_L$ production is very much favoured in the $J_Z = 2$ mode, $W_TW_T$ and $W_TW_L$ productions (which are by far the largest contributions) are slightly more favoured
in the $J_Z = 0$ channel (see Fig. 2). This is the same behaviour as in the two body process $\gamma \gamma \rightarrow W^+W^-$. 

### 3.2 $\gamma \gamma \rightarrow W^+W^-Z$

Contrary to the previous reaction one can calculate the total cross section. It exhibits an interesting behaviour at TeV energies. One notes that already at 1TeV the triple vector boson production is larger that top pair ($m_t \geq 130$GeV) and charged heavy scalars production (with $m_{H^\pm} \simeq 150$GeV) as shown in Fig. 3 which compares various process in $\gamma \gamma$ and $e\gamma$ collisions. At 2TeV the total $WWZ$ cross section is about 2.8pb and exceeds the total electron-positron pair production. This is a typical example of the increasing importance of multiparticle production in weak interactions at higher energies, purely within perturbation theory. The rising of the cross-section with the centre-of-mass energy is essentially from the very forward region due to the presence of the “non-annihilation” diagrams with the (spin-1) $W$ exchanges. A similar behaviour in $e^+e^-$ reactions is single vector boson production. What is certainly more interesting in $\gamma \gamma \rightarrow W^+W^-Z$ is the fact that it has a purely non-abelian origin. It may be likened to $gg \rightarrow ggg$ in QCD except that we do not need any infrared cut-off (the $W$ and $Z$ mass provide a natural cut-off).

The bulk of the cross section consists of both $W$ being transverse as is the case with the “parent” process $\gamma \gamma \rightarrow W^+W^-$. While the total cross-section is larger in the $J_Z = 0$ than in the $J_Z = 2$, the production of all three vector bosons being longitudinal occurs mainly in the $J_Z = 2$ channel and accounts for a dismal contribution. For instance, the ratio of $LLL/TTT$ (three longitudinal over three transverse) in the case of unpolarized beams amount to a mere 2per-mil at 500GeV and drops to 0.1per-mil at 2TeV. Nonetheless, the total production of longitudinal $Z$’s as compared to that of transverse $Z$ is not at all negligible. In fact, between 500GeV and 2TeV this ratio increases from about 23% to 32% (see Fig. 4). This is somehow counterintuitive as one expects the longitudinal states to decouple at high energies. The importance of $Z_L$ production (in association with $W_T^+W_T^-$) is, however, an “infrared” rather than an “ultraviolet” phenomenon in this reaction: the $Z$ is not energetic. First, one has to realize that the $\gamma \gamma \rightarrow W^+W^-Z$ amplitude is transverse in the momentum of the $Z$, $q$, as is the case with the photons in $\gamma \gamma \rightarrow W^+W^-\gamma$. With $k_1$ and $k_2$ being the momenta of the photons, the longitudinal polarization vector of the $Z$, with energy $E_Z$, writes:

\[E_Z = \frac{k_1 \cdot k_2}{E_1 E_2} q \]

\[\text{As opposed to the hypothetical surmise of large } W \text{ multiplicities due to topological effects at extremely high-energies. For a recent review see [3].}

\[\text{i.e., taking into account all polarization states of the } W.\]
\( \epsilon^L_\mu = \frac{1}{\sqrt{E^2_Z - M^2_Z}} \left( \frac{E_Z q_\mu}{M_Z} - \frac{M_Z}{\sqrt{s}} (k_1 + k_2)_\mu \right); \quad \epsilon^L \cdot \epsilon^L = -1 \) (3)

The transversality of the amplitude means that the leading ("ultraviolet" \( \propto E_Z \)) part does not contribute. Only the "infrared" part \( \propto M_Z \) does. This contribution should vanish in the limit of vanishing \( M_Z \). However, the amplitude, in analogy with what happens in \( WW\gamma \), has the infrared factor \( 1/E_Z \) and the "soft" term in equation (3) contributes. Furthermore, more importantly, the bulk of the cross section is from configurations where both \( W \) are transverse (see Fig. 4) and all three particles go down the beam. In the limit of vanishing masses this topology leads to collinear divergences. In this dominating configurations, in the exact forward direction, the longitudinal \( Z \) contributes maximally. At the same time, angular momentum conservation does not allow the \( Z \) to be transverse when all final particles are down the beam (with \( p_T = 0 \)) and both \( W \) are transverse. So the "maximal collinear enhancement" is not as operative for the transverse \( Z \) as it is for longitudinal \( Z \) when both \( W_T \) are at zero \( p_T \). However, as soon as one moves away from these singular configurations, the longitudinal \( Z \) does decouple and the "smooth" mass limit may be taken. This is well rendered in Table.2 which displays the ratio of \( Z_L/Z_T \) without any cut and with the inclusion of cuts. The most drastic of these cuts is when we impose angular separation cuts between the final particles and forcing them to be away from the beam, with the effect that the \( Z_L/Z_T \) decreases with energy and gets dramatically smaller.

| \( \sqrt{s_{\gamma\gamma}} \) (GeV) | \( \sigma \) (fb) | \( L/T \) | \( \sigma \) (fb) | \( L/T \) | \( \sigma \) (fb) | \( L/T \) | \( \sigma \) (fb) | \( L/T \) |
|---------------------------------|-------------|--------|-------------|--------|-------------|--------|-------------|--------|
| \( \sqrt{s_{\gamma\gamma}} = 500 \) | 428        | 24%    | 1443        | 27%    | 2195        | 30%    | 2734        | 32%    |
| \( \sqrt{s_{\gamma\gamma}} = 1000 \) | 368        | 19%    | 1025        | 18%    | 1321        | 19%    | 1465        | 19%    |
| \( \cos(W Z) < 0.8 \) | 164        | 18%    | 232        | 17%    | 186        | 17%    | 145        | 16%    |
| \( \cos(W beam) < 0.8 \) | 115        | 11%    | 140        | 8%     | 105        | 6%     | 77        | 5%     |
| \( \cos( "all" ) < 0.8 \) | 184        | 11%    | 1032       | 21%    | 1700       | 25%    | 2220       | 28%    |
| \( E_Z > 150 \) | 184        | 11%    | 1032       | 21%    | 1700       | 25%    | 2220       | 28%    |

Table 2: Cross section for \( \gamma\gamma \to W^+W^-Z \) and ratio of longitudinal over transverse \( Z \) \( (L/T) \) including various cuts. "all" means that we require the final particles to be separated and to be away from the beam by an angle \( \theta \) corresponding to \( \cos\theta < 0.8 \).

A more detailed study with exact analytical expressions for the helicity amplitudes exhibiting the above behaviour is left for a longer publication.

Footnote: In the limit \( M_V \to 0 \), added to the divergence in \( \gamma\gamma \to W^+W^- \), there is the collinear divergence when the \( Z \) and a \( W \) are collinear.
On the phenomenological side, the study of this reaction is important as, especially for \( M_H \sim M_Z \), it is a background to Higgs detection through \( WWH \) production to which we now turn.

### 3.3 \( \gamma\gamma \rightarrow W^+W^-H \)

While it is almost certain that the LHC/SSC will discover a Higgs if its mass is above \( 2M_Z \), the intermediate mass Higgs, \( IMH \), will be extremely difficult to track at these machines. This “mass gap” will be efficiently covered by the next \( e^+e^- \) linear collider where two complementary reactions are at work: the so-called Bjorken process \( e^+e^- \rightarrow ZH \) dominating at moderately low energies and the \( WW \) fusion process which starts to dominate above \( \sim 500\,\text{GeV} \)\footnote{For a recent comparison between the different modes of Higgs production in \( e^+e^- \) see, \cite{footnote}.}. One of the original motivations for a \( \gamma\gamma \) collider is to produce the scalar Higgs as a resonance. To set the stage, let us recall that, in the range: \( 90 < M_H < 140\,\text{GeV} \), and considering the dominant decay of the Higgs into b-quarks, the cross section \( \sigma(\gamma\gamma \rightarrow H \rightarrow b\bar{b}) \) is about \( \sim 50 - 60\,\text{fb} \), for an optimal set of cuts and parameters of the \( \gamma\gamma \) collider \cite{footnote}.

Another efficient mechanism for Higgs production in an \( e\gamma \) environment is through \( e\gamma \rightarrow \nu WH \)\cite{footnote}. Still, in the context of \( \gamma\gamma \) collisions, it has recently been suggested \cite{footnote} to look at the production of Higgs in association with a top pair in analogy to \( t\bar{t}H \) production in hadron machines. Unfortunately, the \( IMH \) yield does not exceeds \( 1 - 3\,\text{fb} \) (for \( m_t \leq 150\,\text{GeV} \)). The \( t\bar{t}H \) cross section decreases very slowly with \( \sqrt{s_{\gamma\gamma}} \).

Taking, once again, advantage of the large \( WW \) cross section, we propose to search for Higgs in association with a \( W \) pair. We find that for a Higgs mass of \( 100\,\text{GeV} \) we obtain a cross section of about \( 20\,\text{fb} \) at \( \sqrt{s_{\gamma\gamma}} = 500\,\text{GeV} \). The \( WWH \) cross section quickly rises to yield \( \simeq 400\,\text{fb} \) at \( 2\,\text{TeV} \) (for \( m_H = 100\,\text{GeV} \)). The importance of this mechanism at TeV energies is best illustrated, by contrasting it with top pair production (see Fig. 3). For \( m_H = 100\,\text{GeV} \) and \( m_t = 130\,\text{GeV} \) (consistent with present indirect LEP limits), the two process have the same threshold energy and lead to the same final state (\( IMH \) decays predominantly into \( b\bar{b} \)). While at \( \sqrt{s_{\gamma\gamma}} \simeq 500\,\text{GeV} \) top pair production is almost two-orders of magnitude larger than \( WWH \), the latter which is a third order process is twice as large at \( 2\,\text{TeV} \). Nonetheless, the \( WWZ \) cross section is about an order-of-magnitude larger than the “\( IMH-WWH \)” for all centre-of-mass energies.

Comparing at the same \( \sqrt{s_{\gamma\gamma}} \) and \( \sqrt{s_{e\gamma}} \) centre-of-mass, in the \( IMH \) case, the cross sections for \( WWH \) start becoming larger than those of \( e\gamma \rightarrow \nu WH \) for energies around \( 700\,\text{GeV} \).
At lower energies the $e\gamma$ mode benefits from a larger phase space (see Fig. 3).

In Fig. 5 we contrast the various mechanism of Higgs production in an $e^+e^-$ environment in the $e^+e^-$, $\gamma\gamma$, and $e\gamma$ modes before folding with the luminosity spectra. In the IMH case, taking for illustration $M_H = 80\text{GeV}$, at $500\text{GeV}$, $\sigma(\gamma\gamma \rightarrow WWH) \approx 30\text{fb}$ which is by only factor 2 smaller than $\sigma(e\gamma \rightarrow \nu WH)$ and a factor 3.3 compared to the dominant $WW$ fusion process in $e^+e^-$. On the other hand, $\sigma(\gamma\gamma \rightarrow WWH)$ is larger than all the $VVH$ ($WWH, ZZH, ZH\gamma$) processes in $e^+e^-$ by at least a factor 3. Higgs production from top bremsstrahlung ($t\bar{t}H$ final state), either in $e^+e^-$ or $\gamma\gamma$ is abysmally small. At $1\text{TeV}$ our process becomes very comparable to $e\gamma \rightarrow \nu WH$ and is only about a factor 2 smaller than the dominant $WW$ fusion process in $e^+e^-$. Nonetheless, the fact that in $\sigma(\gamma\gamma \rightarrow WWH)$, unlike the $WW$ fusion in $e^+e^-$ or the corresponding one at $e\gamma$, all final particles can be observed or reconstructed (hence alleviating the lack in energy constraints) makes this reaction worth considering especially at a $\text{TeV} \gamma\gamma$ collider. But of course, this statement tacitly assumes an ideal monochromatic $\gamma\gamma$ collider. We will now turn to more realistic photon luminosity spectra. Before so doing, it is worth pointing out that an almost equal number of $H$ is produced in the $J_Z = 0$ or the $J_Z = 2$ with both $W$ being essentially transverse.

4 Inclusion of the photon luminosity spectrum

Thus far, two-photon processes at $e^+e^-$ have exploited the “Weiszäcker-William” spectrum, which is essentially a “soft-photons” spectrum. The $\gamma\gamma$ luminosity peaks for very small fraction of the invariant $\gamma\gamma$ mass $\sqrt{s_{\gamma\gamma}}$, i.e., for $\tau = s_{\gamma\gamma}/s_{e^+e^-} \ll 1$. The laser scheme on the other hand permits to transmit a very large proportion of the energy of the electron ($E_e$) to the “collider” photon by shining a low-frequency ($\omega_0$) laser beam at a glancing angle. With $x$ being the reduced invariant mass of the original $e\gamma_{\text{laser}}$ system: $x \simeq 4E_e\omega_0/m_e^2$, the maximum energy fraction, $y_{\text{max}} = \omega_{\text{max}}/E_e$, that the colliding photon can take is $y_{\text{max}} = x/(x + 1)$. This occurs for photons produced in the exactly forward ($e^-$) direction. One then has to tune the energy of the laser so that one gets the highest $x$. However, this $x$ can not be arbitrarily large, otherwise one reaches the threshold for $e^+e^-$ creation by the interaction of the laser beam and the converted photon. This occurs for $x = x_0 \simeq 4.8$ and translates into $\sqrt{s}_{\text{max}} \approx 0.83$. Our analysis is based on taking this value of $x_0$ for all $e^+e^-$ energies, which means that one has to take different laser frequencies for different $\sqrt{s}$ colliders.

To achieve a higher degree of monochromaticity of the spectrum, polarization is essential. Instead of writing uninspiring lengthy formulae for the polarized luminosity spectra, we
prefer to refer to Fig. 6 (see also [1]). It shows that the hardest spectrum is arrived at by choosing the circular polarization of the laser ($P_c$) and the mean helicity of the electron ($\lambda$) to be opposite. For the photon mode of the collider this means $2\lambda P_c = 2\lambda' P'_c = -1$ ($'$ are for the opposite arm of the photon collider). Fig. 6 shows the case where both lasers are tuned to have a right-handed circular polarization ($P_c = P'_c = +1$). This has the added advantage that the high-energy photons are produced mostly with the same helicity therefore giving a $J_Z = 0$ dominated environment, for short we will refer to this as the “0-dom.” case. The $J_Z = 2$ tail almost disappears for $\sqrt{s} > 0.7$. For some processes where the $J_Z = 2$ is dominant, or if one wants to compare the $J_Z = 2$ and the $J_Z = 0$ on an “equal basis”, one would also like to isolate the $J_Z = 2$ at the expense of the $J_Z = 0$ spectrum. We point out that this could be easily achieved by flipping both the electron and laser polarizations of one of the arms only while maintaining $2\lambda P_c = -1$ (for a maximum of monochromaticity). In this case, the $J_Z = 0$ and $J_Z = 2$ spectra in Fig. 6 have to be interchanged, for short “2-dom.”. The spectrum one expects in case of no polarization is rather flat, with a slight hump in the “mid-range” $\sqrt{s} \simeq 0.2 - 0.5$. For processes which increase with energy, as with the three processes we have studied, it is best to choose the hardest spectrum arrived at through oppositely-handed $e, \gamma_{\text{laser}}$. This also helps in sensibly reducing standard processes which at the “partonic” ($\gamma\gamma$) level drop as $1/s$. We will discuss the effect of the luminosity spectrum in the reactions we have studied for the case of polarized beams and in the case of no polarization.

5 Folding with the luminosity spectra

We illustrate the effect of different luminosity spectra by concentrating on the IMH search through $WWH$. This will lead us to consider $WWZ$ production which is the most obvious background for $M_H \sim M_Z$. With $\sqrt{s_{ee}} = 500\text{GeV}$, the inclusion of the spectra changes the $WWH$ yield significantly due to the fact that the maximum $\sqrt{s_{\gamma\gamma}} \simeq 400\text{GeV}$ leaves a small phase space for the IMH. Even when we choose the polarization of the primary beams to give the peaked $J_Z = 2$-dominating spectrum (“2-dom”), the cross section does not exceed 4fb and is therefore almost two orders of magnitude below the $WW$ fusion process in the $e^+e^-$ mode and an order of magnitude smaller than $\nu WH$ production in the $e\gamma$ mode (See Fig. 7a). The situation is much more favourable at 1TeV. Up to $M_H \simeq 300\text{GeV}$ this mode produces almost twice as many Higgses as the conventional Bjorken process. For $M_H = 100\text{GeV}$ and choosing a setting which gives a “0-dom”, we obtain $\sim 37.5\text{fb}$ (compared to 37.2 in the “2-dom”) and $26.3\text{fb}$ with no polarization for the primary beams. The advantage of a polarized spectrum is undeniable. $\nu WH$ (in $e\gamma$) and
Table 3: Cross section in fb for three-boson (and $t\bar{t}$) productions with $M_H = 100\,\text{GeV}$ and $m_t = 150\,\text{GeV}$. The $WW\gamma$ includes a $p_T$ cut of 20 GeV at 500 GeV and 40 GeV at 1 TeV. “non-top” means all Higgs events with Higgs decaying into $bb$ and where the simultaneous $Wb$ invariant mass has been applied as explained in the text. “direct” means that we have not taken into account top pairs produced through the gluons inside the photon, i.e., the “resolved” photons contribution has not been considered.

$H\nu\nu$ (in $e^+e^-$) are respectively about 2 and 5 times larger in the IMH case. A comparison between the variety of Higgs production modes in the NLC(1TeV) environment is shown in Fig. 7b which clearly brings out the importance of $WWH$.

The effect of switching between different polarization settings is even more drastic in the case of $WWZ$. The largest cross sections are in the “0-dom.” case. At $\sqrt{s_{ee}} = 500\,\text{GeV}$ we find $\int \sigma(WWZ) \sim 55\,\text{fb}$ which is twice as large as the non-polarized case (24.2fb). Note that, when choosing the “0-dom” the $WWZ$ yield is larger than in the conventional $e^+e^-$ mode ($\sigma(e^+e^- \rightarrow W^+W^-Z)_{\sqrt{s_{ee}}=500\,\text{GeV}} \sim 40\,\text{fb}$). At 1 TeV the $WWZ$ reaches $\sim 470\,\text{fb}$ in the “0-dom” and is slightly smaller (410) in the “2-dom”.

Considering the large $WWZ$ yield, $b$ tagging is almost necessary for the IMH search. Another dangerous background, even for the case of $b$-tagging is due to top pair production: $\gamma\gamma \rightarrow t\bar{t} \rightarrow W^+W^-bb$. For instance, at $\sqrt{s_{ee}} = 500\,\text{GeV}$ this is about two-orders of magnitude larger than $WWH_{\rightarrow bb}$. Fortunately, one can eliminate this huge contamination by rejecting all those $WWH$ events where the simultaneous cuts on the invariant mass of the two $Wb$ is such that the $Wb$ does not reconstruct the top mass (within 15 GeV)

$$m_t - 15\,\text{GeV} < M_{W^+b} < m_t + 15\,\text{GeV} \quad \text{and} \quad m_t - 15\,\text{GeV} < M_{W^-\bar{b}} < m_t + 15\,\text{GeV}$$

or

$$m_t - 15\,\text{GeV} < M_{W^+\bar{b}} < m_t + 15\,\text{GeV} \quad \text{and} \quad m_t - 15\,\text{GeV} < M_{W^-b} < m_t + 15\,\text{GeV}$$

(4)
The reason we try both combinations $W^+b$ or $W^+b'$ is that we do not want to rely on charge identification, for the $b$ especially, which necessarily entails a reduction in the $b$ sample (and hence our signal). A good vertex detector should be sufficient. In carrying the vetoing in our Monte-Carlo sample we made the Higgs decay isotropically in its rest frame. The effective loss at 500GeV is about a mere 0.3fb while at 1TeV, where we have a “healthy” cross section, the percentage loss is only about 4% for all choices of the polarization. Table 3. shows the cross sections taking a Higgs mass of $M_H = 100$GeV with $Br(H \rightarrow b\bar{b}) \sim 80\%$ and the cut of equation 4, assuming $m_t = 150$GeV. Once the “faked” top events have been dealt with, the $WWZ \rightarrow b\bar{b}$ do not bury the signal (for $M_H \sim M_Z \pm 10$GeV) as Table 3 shows. These $WWZ$ can be further reduced by judiciously switching the “2-dom.” setting, both at 500GeV and at 1TeV. Although at the former energy the event rate is probably too small to be useful, at 1TeV, in the “2-dom.”, we have, after including the cuts and the branching fractions into $b$, 30fb of signal compared to 60fb from $WWZ$. With one $W$ at least, decaying into jets and not taking into account decays into $\tau$’s, the number of $WWH$ with the contemplated integrated luminosity of $\mathcal{L} = 60fb^{-1}$ will be about 1400 events. Even if one allows for an overall efficiency of 50% this is a very important channel to look for the Higgs. There is one background which we have not considered. It concerns the $W^+W^-b\bar{b}$ final state with $b\bar{b} \rightarrow W^+W^-$ as a sub-process. We expect this to be very negligible once one puts a high $p_T$ cut on both $b$’s and require $m_{bb} \sim M_H$. We will give a more detailed analysis of all these processes and a more thorough discussion on background elimination in a longer forthcoming paper.

To conclude, we have shown that this new mechanism of Higgs production in a $\gamma\gamma$ mode of $\sim 1TeV$ $e^+e^-$ collider is a very promising prospect. The oft discussed intermediate mass Higgs production, as a narrow resonance in $\gamma\gamma$ collisions, relies on a spectrum which is peaked around the Higgs mass in a $J_Z = 0$ dominated setting. The extensive study in [8] finds that with $\int \mathcal{L}_{ee} = 10fb^{-1}$, one expects between about 500 Higgs events for $M_H \sim M_Z$ to about 600 events for $M_H \sim 140$GeV. For the same $\int \mathcal{L}_{ee} = 10fb^{-1}$ this is about 2 – 3 times more than what we get with $WWH$ at $\sqrt{s_{ee}} = 1$TeV. However, the resonance scheme means that the available $\gamma\gamma$ invariant mass covers a very narrow, and in the case of the $IMH$, low range of energies. Hence while allowing a precise study of the $H\gamma\gamma$ coupling it forbids the study of a wealth of interesting processes in the $\gamma\gamma$ mode of the NLC. Higgs detection through $WWH$ at 1TeV will be one aspect among a variety of studies of weak processes ($WW, ZZ, WWZ,...$ etc) in a $\gamma\gamma$ environment.

**We have not tried to cut the $t\bar{t}$ by demanding that $m_{bb} = M_H \pm 10$GeV’, as the cut above is very efficient. Moreover, based on our previous analysis of $WWH$ in $e^+e^-$ [7], the $Wb$ cut was by far more efficacious.
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Figure Captions

Fig. 1 Feynman graphs contributing to $\gamma\gamma \to WWH, WWZ$ or $WW\gamma$. These have to be properly symmetrised. One then counts 4 diagrams of type (a) and (d), and 2 of types (b) and (c).

Fig. 2 $\gamma\gamma \to WW\gamma$ cross section as a function of $\sqrt{s_{\gamma\gamma}}$ with $p_T^\gamma > 20$GeV, for different initial and final polarization states. The subscript $0$ and $2$ refer respectively to an initial state with $J_Z = 0$ and $J_Z = \pm 2$. TT is for both W being transverse, TL when only one is transverse and LL when both are longitudinal.

Fig. 3 Typical processes in $\gamma\gamma$ and $e\gamma$ reactions, with $m_t = 130$GeV, $M_{H^+} = 150$GeV. In $\gamma\gamma \to W^+W^-\gamma$ the cut is $p_T^\gamma > 20$GeV. $e^+e^-\text{cut}$ represents $\gamma\gamma \to e^+e^-$ with $|\cos(\gamma e)| < 0.8$. The Higgs masses (100GeV and 200GeV) are indicated by the subscripts in $WWH$ and $WH\nu$.

Fig. 4 WWZ cross section for different combinations of final polarizations as a function of $\sqrt{s_{\gamma\gamma}}$, with unpolarized photons. $Z_T$ ($Z_L$) is the transverse (longitudinal) Z yield for any W helicity state. Also shown is the ratio of longitudinal over transverse Z ($Z_L/Z_T$) summed over all polarization states of the W’s.

Fig. 5 A comparison of Higgs production cross-sections in $e^+e^-\gamma\gamma$ and $e\gamma$ reactions at 500GeV (5.a) and 1TeV (5.b) before folding with any photon luminosity spectrum. When the initial state is not specified, it should be understood as an $e^+e^-$ process.

Fig. 6 The $\gamma\gamma$ spectrum for different choices of the primary electron longitudinal ($\lambda$) and photon circular ($P_c$) polarization with $x_0 = 4.82$. ’ are for the opposite arm of the collider. The figure also shows the $J_Z = 0$ and $J_Z = 2$ part in the case of $P_c = P_c' = +1$.

Fig. 7 Comparison between different Higgs production mechanisms at a 500GeV (a) and a future 1TeV (b) $e^+e^-$ machine in the three modes of the collider. No beam polarization effects are included apart from $\gamma\gamma \to WWH$ ($WWH_{pol}$) at 1TeV where we also show the effect of a “$J_Z = 0$-dominated” setting (see text).

References

[1] I.F. Ginzburg, G.L. Kotkin, V.G. Serbo and V.I. Telnov, Nucl. Instrum. Methods 205 (1983) 47; I.F. Ginzburg, G.L. Kotkin, S.L. Panfil, V.G. Serbo and V.I. Telnov, ibid 219 (1984) 5; V.I. Telnov, ibid A294 (1990) 72.

[2] For a list of references on this process, see G. Bélanger and F. Boudjema, Phys. Lett. B288 (1992) 210.
[3] M. Gavela, G. Girardi, C. Malleville and P. Sorba, Nucl. Phys. B193 (1981) 257.

[4] M. Baillargeon and F. Boudjema, Phys. Lett. B272 (1991) 158 and references therein. A variant of the non-linear gauge fixing presented here has been, very recently, been employed by G. Jikya for the calculation of $\gamma\gamma \rightarrow ZZ$, IHEP-preprint, IHEP-97/93.

[5] A. Ringwald, CERN-preprint, CERN-TH-6862-93, April 1993.

[6] M. Baillargeon and F. Boudjema, in preparation.

[7] M. Baillargeon et al., CERN-preprint, CERN-TH-6932/93, June 1993.

[8] D.L. Borden, D.A. Bauer, D.O. Caldwell, SLAC-Preprint, SLAC-PUB-5715/UCSB-HEP-92-01, January 1992. See also F. Richard, Orsay-Preprint, LAL-91-62, Nov. 1991.

[9] K. Hagiwara, I. Watanabe and P.M. Zerwas, Phys. Lett. B278 (1992) 187. See also E. Boos et al., DESY-Preprint 91-114, Oct. 1991.

[10] E. Boos et al., Z. Phys. C56 (1992) 487. K. Cheung, Phys. Rev. D47 (1993) 3750.

[11] A. Djouadi, J. Kalinowski and P. Zerwas, Z. Phys. C54, 255 (1992).

[12] G. Bélanger and F. Boudjema, Phys. Lett. B288 (1992) 201.