The Higgs boson sector is a crucial part of the Standard Model (SM) still escaping direct experimental verification. Once the scalar boson will be discovered either at LEP2, upgraded TEVATRON or at LHC, testing its properties will be a central issue at future linear colliders. In particular, an $e^+e^-$ collider with centre-of-mass (c.m.) energy $\sqrt{s} \simeq (300 \div 2000)\text{GeV}$ and integrated luminosity $O(100) \text{fb}^{-1}$ will allow an accurate determination of the mass, some couplings and parity properties of this new boson \[1, 2\]. Among other couplings, the interaction of scalars with the neutral electroweak gauge bosons, $\gamma$ and $Z$, are particularly interesting. Indeed, one can hope to test here some delicate feature of the Standard Model — the relation between the spontaneous symmetry breaking mechanism and the electroweak mixing of the two gauge groups $SU(2)$ and $U(1)$. In this respect, three vertices could be measured — $ZZH$, $\gamma\gamma H$ and $Z\gamma H$. While the $ZZH$ vertex stands in SM at the tree level, the other two contribute only at one-loop. This means that the $\gamma\gamma H$ and $Z\gamma H$ couplings could be sensitive to the contributions of new particles circulating in the loop.

Here, we discuss the case of an intermediate-mass Higgs boson, that is with $M_Z \lesssim m_H \lesssim 140$ GeV. A measurement of the $\gamma\gamma H$ coupling should be possible by the determination of the BR for the decay $H \to \gamma\gamma$, e.g. in the LHC Higgs discovery channel, $gg \to H \to \gamma\gamma$. Furthermore, at future photon-photon colliders\[4\], the precise measurement of the $\gamma\gamma H$ vertex looks realistic at the resonant production of the Higgs particle, $\gamma\gamma \to H$. To this end, the capability of tuning the $\gamma\gamma$ c.m. energy on the Higgs mass, through a good degree of the photons monochromaticity, will be crucial for not diluting too much the $\gamma\gamma \to H$ resonant cross section over the c.m. energy spectrum. Measuring the $Z\gamma H$ vertex is in general more complicated. Indeed, if one discusses the corresponding Higgs decay, the final states include the $Z$ decay products, jets or lepton pairs, where much heavier backgrounds are expected. Then, one can discuss the $H \to \gamma Z$ decay only for $m_H \gtrsim 115$ GeV (and $m_H \lesssim 140$), when the corresponding branching is as large as $O(10^{-3})$. Another possibility of measuring the $Z\gamma H$ vertex is given by collision processes. At electron-positron colliders, the corresponding channels are $e^+e^- \to \gamma H$ and $e^+e^- \to ZH$. However, in the $ZH$ channel the $Z\gamma H$ vertex contributes to the corresponding one-loop corrections, thus implying a large tree level background. The reaction $e^+e^- \to \gamma H$ has been extensively studied in the literature \[5, 6, 7\]. Unfortunately, the $e^+e^- \to H\gamma$ channel suffers from small rates, which are further depleted at large energies by the $1/s$ behavior of the

\[1\] Two further options are presently considered for a high-energy $e^+e^-$ linear collider, where one or both the initial $e^+/e^-$ beams are replaced by photon beams induced by Compton backscattering of laser light on the high-energy electron beams \[8\]. Then, the initial real photons could be to a good degree monochromatic, and have energy and luminosity comparable to the ones of the parent electron beam \[9\].
dominant s-channel diagrams. For example, \( \sigma_S \approx 0.05 \div 0.001 \) fb at \( \sqrt{s} \approx 500 \div 1500 \) GeV. We estimated the main background coming from the \( e^+e^- \to \gamma \bar{b}b \) process, and found it rather heavy: \( \sigma_B \approx 4 \div 0.8 \) fb for \( m_{\bar{b}b} = 100 \div 140 \) GeV, assuming a high resolution in the measurement of the invariant mass of \( \bar{b}b \) quark pair, i.e. \( \pm 3 \) GeV, and applying a minimum cut of \( 18^\circ \) on the angles \( [\gamma - \text{beams}] \) and \( [b(\bar{b}) - \text{beams}] \). Then, at \( \sqrt{s} = 1.5 \) TeV, one gets \( \sigma_B \approx 0.4 \div 0.07 \) fb. One can conclude that measuring the \( \gamma\gamma H \) vertex is not an easy task. Recently, the Higgs production in electron-photon collisions through the one-loop process \( e\gamma \to eH \) was analysed in detail. This channel will turn out to be an excellent mean to test both the \( \gamma\gamma H \) and \( Z\gamma H \) one-loop couplings with high statistics, without requiring a fine tuning of the c.m. energy.

In this talk we discuss the prospects of the \( e\gamma \to eH \) reaction in setting experimental bounds on the value of the anomalous \( Z\gamma H \) coupling. For this analysis we use a model independent approach, where \( \dim = 6 \) \( SU(2) \times U(1) \) invariant operators are added to the SM Lagrangian. In realistic models extending the SM, these operators contribute in some definite combinations. However, if one discusses the bounds on the possible deviations from the standard-model one-loop Higgs vertices, this approach can give some general insight into the problem. These anomalous operators contribute to all the three vertices \( \gamma\gamma H \), \( Z\gamma H \) and \( ZZH \), with only the first two involved in the \( e\gamma \to eH \) reaction. As we mentioned above, the anomalous contributions to the \( \gamma\gamma H \) vertex can be bounded through the resonant \( \gamma\gamma \to H \) reaction, or by measuring the total rate of the discussed reaction, \( e\gamma \to eH \). Here we discuss the case where the anomalous contributions to the \( \gamma\gamma H \) vertex have been well tested in some other experiment, and one would like to get limitations just on the anomalous contributions to the \( Z\gamma H \) vertex. We underline that this specific case is out of the \( \gamma\gamma \) domain.

All the results presented in this talk were obtained with the help of CompHEP package [9].

**The reaction \( e\gamma \to eH \) : main features**

In [8], we presented the complete analytical results for the helicity amplitudes of the \( e\gamma \to eH \) process (see also [10]). The total rate of this reaction is rather high. For \( m_H \) up to about 400 GeV, one finds \( \sigma > 1 \) fb.

The main strategy to enhance the \( Z\gamma H \) vertex effects (depleted by the \( Z \) propagator) with respect to the dominant \( \gamma\gamma H \) contribution consists in requiring a final electron tagged at large angle. E.g., for \( p_T^e > 100 \) GeV, \( Z\gamma H \) is about 60% of \( \gamma\gamma H \), and \( Z\gamma H \) gives a considerable fraction of the total production rate, which is still sufficient to guarantee investigation (about 0.7 fb).

The main irreducible background to the process \( e\gamma \to eH \to \bar{b}b \) comes from the channel \( e\gamma \to \bar{b}b \). A further source of background is the charm production through \( e\gamma \to ec\bar{c} \), when the \( c \) quarks are misidentified into \( b \)'s. The cut \( \theta_{b(c) - \text{beam}} > 18^\circ \) (between each \( b(c) \) quark and both the beams) reduces the signal and background at a comparable level. Numerically the \( e\gamma \to ec\bar{c} \) “effective rate” is less than \( 1/3 \) of the \( e\gamma \to \bar{b}b \) rate. A further background, considered in [8], is the resolved \( e\gamma(g) \to \bar{b}b(ec\bar{c}) \) production, where the photon interacts via its gluonic content. It was found \(< 0.01 \) fb at \( \sqrt{s} = 500 \) GeV with our standard cuts \( p_T^e > 100 \) GeV and \( \theta_{b(c) - \text{beam}} > 18^\circ \).

We also considered the possibility of having the electron beam longitudinally polarized. In table 1 results for the \( e \) polarization dependence of the total cross section, its \( \gamma\gamma H \), \( Z\gamma H \) and BOX components, and their interference pattern are collected for \( m_H = 120 \) GeV. This de-

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2In [8], we define different contributions: ‘\( \gamma\gamma H \)’ and ‘\( Z\gamma H \)’, related to the \( \gamma\gamma H \) and \( Z\gamma H \) vertices respectively, and a ‘BOX’ contribution. The separation of the rate into these three parts corresponds to the case where the Slavnov-Taylor identities for the ‘\( \gamma\gamma H \)’ and ‘\( Z\gamma H \)’ Green functions just imply the transversality with respect to the incoming photon momentum.

3we also assume a 10% probability of misidentifying a \( c \) quark into a \( b \).
dependence turns out to be very sensitive to the electron polarization. For instance, assuming $P_e = -1$ ($P_e = +1$) the total cross section increases (decreases) by about 94\% at $\sqrt{s} = 500\text{GeV}$. For $P_e = +1$, there is a strong destructive interference between the terms $\gamma\gamma H$ and $Z\gamma H$.

Important improvements in the $S/B$ ratio can be obtained by exploiting the final-electron angular asymmetry in the signal. Indeed, the final electron in $e\gamma \rightarrow eH$ moves mostly in the forward direction. On the other hand, we found that in the $e\gamma \rightarrow e\bar{b}b$ background the final electron angular distribution, although not completely symmetric, is almost equally shared in the forward and backward direction with respect to the beam. In table 2, we report, apart from the total rate, the forward cross sections. It could also be convenient to measure the FB asymmetry, that has the advantage of being free from possible uncertainties on the absolute normalization of the cross sections.

The main conclusion obtained in [8] is the following. With a luminosity of 100 fb$^{-1}$, at $\sqrt{s} = 500\text{GeV}$, one expects an accuracy as good as about 10\% on the measurement of the $Z\gamma H$ effects assuming the validity of Standard Model.

**Anomalous vertices**

There are two pairs of $dim = 6$ operators, CP-even and CP-odd respectively, giving anomalous contributions to the process $e\gamma \rightarrow eH$ [11]:

$$\mathcal{L}^{eff} = d \cdot \mathcal{O}_{UW} + d_B \cdot \mathcal{O}_{UB} + \bar{d} \cdot \bar{\mathcal{O}}_{UW} + \bar{d}_B \cdot \bar{\mathcal{O}}_{UB},$$

$$\mathcal{O}_{UW} = \frac{1}{v^2} \left( |\Phi|^2 - \frac{v^2}{2} \right) \cdot W^{i\mu\nu} W_i^{\mu\nu}, \quad \mathcal{O}_{UB} = \frac{1}{v^2} \left( |\Phi|^2 - \frac{v^2}{2} \right) \cdot B^{\mu\nu} B_{\mu\nu},$$

$$\bar{\mathcal{O}}_{UW} = \frac{1}{v^2} |\Phi|^2 \cdot \bar{W}^{i\mu\nu} \bar{W}_{i\mu\nu}, \quad \bar{\mathcal{O}}_{UB} = \frac{1}{v^2} |\Phi|^2 \cdot B^{\mu\nu} \bar{B}_{\mu\nu},$$

where $\bar{W}^{i\mu\nu} = \epsilon_{\mu\nu\rho\sigma} W^{i\rho\sigma}$ and $\bar{B}^{\mu\nu} = \epsilon_{\mu\nu\rho\sigma} B^{\rho\sigma}$. In these formulas $\Phi$ is the Higgs doublet and $v$ is the electroweak vacuum expectation value.

The $Z\gamma H$ anomalous terms in the helicity amplitudes of $e\gamma \rightarrow eH$ are

$$\frac{4\pi\alpha}{M_Z(M_Z^2 - t)} \sqrt{-\frac{t}{2}} \left\{ d_{\gamma Z} [(u - s) - \sigma \lambda (u + s)] - i\bar{d}_{\gamma Z} [\lambda (u - s) + \sigma(u + s)] \right\},$$

where $s$, $t$, and $u$ are the Mandelstam kinematical variables, $\sigma/2 = \pm 1/2$ and $\lambda = \pm 1$ are the electron and photon helicities, respectively. The anomalous couplings contribute in the combinations $d_{\gamma Z} = d - d_B$ and $\bar{d}_{\gamma Z} = \bar{d} - \bar{d}_B$.

Note that there is no interference between CP-odd terms (couplings $\bar{d}$) and any triangle terms in the SM amplitude, although the interference with the BOX amplitude is nonvanishing. This is due to the real value of the SM triangle amplitudes for $M_H < 2M_W$ and $M_H < 2m_{top}$ (one can neglect contributions of light quark loops), while the box diagrams contribute with complex numbers. However, the BOX amplitude gives a rather small contribution. So, the cross section depends only marginally on the sign of the CP-odd couplings $d_{\gamma Z}$. On the contrary, for CP-even terms, the interference is relevant, and the dependence on the $d_{\gamma Z}$ coupling is not symmetrical with respect to the SM point $d_{\gamma Z} = 0$. This effect, generally speaking, decreases the sensitivity to the CP violating anomalous couplings.

**Bounds on the anomalous $Z\gamma H$ coupling**

As we mentioned, we assume that the anomalous contributions to the $\gamma\gamma H$ vertex are already tightly bounded somehow, by using the data from other experiments, and one wants to get limitations on the anomalous contributions to the $Z\gamma H$ vertex. So, we discuss the case where the anomalous operators contribute in the combinations $d - d_B = d_{\gamma Z}$ and $\bar{d} - \bar{d}_B = \bar{d}_{\gamma Z}$ only.
The bounds on the anomalous $d_{\gamma Z}$ and $\bar{d}_{\gamma Z}$ couplings presented below, have been computed by using the requirement that no deviation from the SM cross section is observed at the 95% CL:

$$N_{\text{anom}}(\kappa) < 1.96 \cdot \sqrt{N_{\text{tot}}(\kappa)}, \quad \kappa = d_{\gamma Z}, \bar{d}_{\gamma Z},$$

$$N_{\text{tot}}(\kappa) = \mathcal{L}_{\text{int}} \cdot [\sigma_S(\kappa) + \sigma_B], \quad N_{\text{anom}}(\kappa) = \mathcal{L}_{\text{int}} \cdot [\sigma_S(\kappa) - \sigma_S(0)].$$

Here, by $\sigma_S(\kappa)$ we denote the cross section of the signal reaction $e\gamma \to eH \to e\bar{b}b$ with the anomalous contributions, so $\sigma_S(0)$ is the standard-model cross section. Then, by $\sigma_B$ we denote the total cross section of the background processes $e\gamma \to e\bar{b}b$, $e\gamma \to e\bar{c}c$ (with 10% probability of misidentifying a $c$ quark into a $b$ quark), and the corresponding contributions from the resolved photons. $\mathcal{L}_{\text{int}}$ is the integrated luminosity. The results of these calculations are collected in the last two columns in table 2, where we present the bounds on the anomalous couplings $d_{\gamma Z}$ and $\bar{d}_{\gamma Z}$, attainable with an integrated luminosity of 100 fb$^{-1}$. We find, that these bounds do not depend significantly on the electron angular cut while the $S/B$ ratio is improved up to $\sim 1$. Nevertheless, we stress that the measurement of the relative final-electron angular asymmetry could be more convenient, if the systematic errors connected with the absolute normalization of the cross sections are relevant. Furthermore, one can see from the table 2, that the strongest limitation on the CP-even coupling (at the level $|d_{\gamma Z}| \lesssim 0.0015$) is obtained from a left-handed polarized electron beam. On the other hand, in the CP-odd case, the right-handed polarized electron beam gives the best performance, with bounds at the level of $|\bar{d}_{\gamma Z}| \lesssim 0.005$. In the latter case, the violation of the strong destructive interference between the $Z\gamma H$ and $\gamma\gamma H$ terms by the anomalous terms compensates the decrease in statistics. As a result, we get better bounds. Even in the CP-even case, a right-handed polarized electron beam could add some valuable information to the ‘left-handed’ data. Indeed, due to the rather large interference between the CP-even anomalous terms with the SM amplitudes, the experimental data (obtained at integrated luminosity 100 fb$^{-1}$) will give two limitation intervals. One interval includes the SM point $d_{\gamma Z} = 0$, and the other a disconnected range $0.019 < d_{\gamma Z} < 0.022$. This is due to the asymmetry of the cross section with respect to the SM point $d_{\gamma Z} = 0$. Hence, an experiment with a right-handed polarized $e$ beam could resolve these discrete uncertainty.

Finally, let us rewrite the obtained bounds on the anomalous couplings in terms of the mass scale of the New Physics. We introduced the anomalous couplings as dimensionless parameters in the effective Lagrangian using the electroweak scale $v \approx 250$ GeV. If these anomalous terms appear as contributions of new particles in the $Z\gamma H$ loop with the mass $M_{\text{new}}$, then one gets

$$d_{\gamma Z}, \bar{d}_{\gamma Z} \sim (v/M_{\text{new}})^2.$$

By using this relation, one obtains the bounds $M_{\text{new}} > 6.2$ TeV in the CP-even case and $M_{\text{new}} > 3.5$ TeV in the CP-odd case.

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Table 1: Interference pattern between the \( \gamma\gamma Z \), \( Z\gamma H \) and boxes contributions versus the \( e\)-beam polarizations, for \( p_T^e > 100\) GeV. \( \gamma \)-beam is unpolarized.

| \( m_H = 120 \) GeV | \( \sqrt{s} = 500 \) GeV | \( \sigma(e\gamma \rightarrow eH) \) fb | Interference terms |
|---------------------|------------------|-----------------|-----------------|
|                     | Total            | \( |\gamma H|^2 \) | \( |Z\gamma H|^2 \) | \( |BOX|^2 \) |
| \( P_e = 0 \)       | 0.705            | 0.296           | 0.166           | 0.0532 |
| \( P_e = -1 \)      | 1.37             | 0.296           | 0.199           | 0.106  |
| \( P_e = +1 \)      | 0.0405           | 0.296           | 0.133           | 0.0004 |

Table 2: Bounds on the anomalous \( Z\gamma H \) couplings for different \( e\)-beam polarizations. The cross sections are shown for signal process and for background, where all contributions are collected, from \( e\gamma \rightarrow ebb \) and \( e\gamma \rightarrow ecc \) (assuming 10\% of misidentifying a \( c \) quark into a \( b \)) and from the resolved photons. The cuts \( \theta_{b\text{-beam}} > 18^\circ \), \( p_T^e > 100 \) GeV and \( m_H - 3 \text{GeV} < m_{bb(c\bar{c})} < m_H + 3 \text{GeV} \) are applied.

| \( m_H = 120 \) GeV | \( \sqrt{s} = 500 \) GeV | \( \sigma_S/\sigma_B \) fb | Bounds at 100 fb\(^{-1}\) |
|---------------------|------------------|-----------------|-----------------|
|                     | no \( \theta_e \) cut | \( \theta_e < 90^\circ \) | \( d_{\gamma Z} \) | \( \tilde{d}_{\gamma Z} \) |
| \( P_e = 0 \)       | 0.530/0.881      | 0.425/0.483     | (-0.0025, 0.004) | (-0.006, 0.006) |
| \( P_e = -1 \)      | 1.03/1.307       | 0.820/0.719     | (-0.0015, 0.0015) | (-0.006, 0.006) |
| \( P_e = +1 \)      | 0.0249/0.472     | 0.0245/0.270    | (-0.007, 0.004) | (-0.005, 0.005) |