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Signatures of anomalous Higgs couplings in angular asymmetries of $H \to Z\ell^+\ell^-$ and $e^+e^- \to HZ$

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Abstract

Parametrizing beyond Standard Model physics by the $SU(3) \times SU(2)_L \times U(1)_Y$ dimension-six effective lagrangian, we study the impact of anomalous Higgs couplings in angular asymmetries of the crossing symmetric processes $H \to Z\ell^+\ell^-$ and $e^+e^- \to HZ$. In the light of present bounds on $d = 6$ couplings, we show that some asymmetries can reveal BSM effects that would otherwise be hidden in other observables. The $d = 6$ $HZ\gamma$ couplings as well as (to a lesser extent) $HZ\ell\ell$ contact interactions can generate asymmetries at the several percent level, albeit having less significant effects on the di-lepton invariant mass spectrum of the decay $H \to Z\ell^+\ell^-$. The higher di-lepton invariant mass probed in $e^+e^- \to HZ$ can lead to complementary anomalous coupling searches at e^+e^- colliders.

Keywords: Higgs physics, dimension-six effective Lagrangian, Beyond Standard Model physics

1. Introduction, operators and couplings

The discovery of a light boson H with mass around 125 GeV at the LHC [1, 2] has opened a window to a new sector in the search for physics beyond the Standard Model (BSM). The new state is compatible with a SM Higgs, with the quantum numbers $J^P = 0^+$ being highly favoured by the data [3, 4]. The study of signal strengths of the new state has shown that the Higgs couplings are compatible with SM predictions. Evidence for BSM physics has proven to be more elusive than previously expected; the SM appears to be a good effective field theory (EFT) at the least up to the energies probed by the first run of LHC.

In the spirit of an EFT, the SM should be supplemented with all operators with dimension $d > 4$ constructed from its fields and compatible with its symmetry. In this work we adopt the linear realization of the $SU(2)_L \times U(1)_Y$ symmetry [5, 6]. The leading corrections to Higgs physics within this scheme arise from the dimension-six operators, that are suppressed by the large scale Λ characteristic of BSM physics, and generate anomalous Higgs boson couplings.

In the search for BSM physics in the flavour sector of the standard model, in particular in the case of flavour-changing neutral currents, dedicated observables were constructed from the angular distribution of the decay $B \to K^*\ell\ell$. The angular distribution of the decay $H \to Z\ell^+\ell^-$ offers similar possibilities that we exploit in this work.

The study of the decay $H \to Z\ell^+\ell^-$, with the on-shell Z also decaying into $\ell^+\ell^-$, has a long history. Its angular distributions were instrumental in the determination of the Higgs quantum numbers [3, 4], as suggested long ago (see e.g. Refs. [7, 8, 9, 10]). After the Higgs discovery, it has been suggested that the di-lepton mass distri-

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bution of the decay can reveal effects that would be hidden in the total decay width [11, 12, 13, 14]. Recently, the full angular distribution of the decay has been revisited in the framework of the EFT parametrization of BSM physics [15]. It has been shown that some angular asymmetries can reveal effects that would be hidden even in the di-lepton mass distributions.

More recently, we performed an extended study of the angular asymmetries of $H \to Z(\to \ell^+\ell^-)\ell^+\ell^-$ and of the crossing-symmetric process $e^+e^- \to HZ$ [16]. The latter process should be measured with precision at a high-energy e^+e^- collider such as the ILC [17] and should provide a clean way to extract Higgs couplings [18, 19, 20]. Our focus is on these asymmetries, their sensitivity to anomalous Higgs couplings, and the interplay between the asymmetries and the di-lepton mass distributions. Our aim here is to highlight some of the main findings of Ref. [16]; for additional details we refer to that reference.

In the massless lepton limit, the two processes are described by the same set of six form factors [7, 8, 11, 13, 15], albeit in different kinematic regimes, related by analyticity. Ignoring loop corrections and neglecting the lepton masses, the processes are governed by six independent angular functions of three independent angles among the four leptons. These functions can be expressed in terms of the six form factors that, in turn, can be written in terms of the couplings of the general $d = 6$ Lagrangian.

Assuming the new physics sector to be characterized by the scale Λ, larger than the electroweak scale, the SM is supplemented with 59 independent $d = 6$ operators [5, 6]. This Lagrangian can be schematically cast as

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}}^{(4)} + \frac{1}{\Lambda^2}\sum_{k=1}^{59}\alpha_k O_k, \qquad (1)$$

where the α_k is the coupling of operator O_k. The effective Lagrangian implies a parametrization of anomalous Higgs interactions (contained in O_k) constrained by the SM gauge symmetry. In our expressions, we often employ the dimensionless coefficients $\widehat{\alpha}_k$ defined as

$$\widehat{\alpha}_k = \frac{v^2}{\Lambda^2}\alpha_k, \qquad (2)$$

where v is the classical Higgs vacuum expectation value. The coefficients $\widehat{\alpha}_k$ should be smaller than $O(1)$ for the EFT description to be applicable.

Here we employ the complete operator basis defined in Ref. [6], although different choices are possible and in use. In practice we only need to work with a subset of the 59 operators, since not all of them contribute at tree level to the processes of interest. Furthermore, assuming minimal-flavour violation to avoid tree-level flavour-changing neutral currents, flavour matrices of operators that involve a left-handed doublet and a right-handed singlet are fixed to be the same as in the SM Yukawa couplings. With this assumption, these operators are proportional to lepton masses and can be safely neglected.

The operators considered in this work are listed in Eqs. (3), (4) and (5) below. The notation and conventions follow those of Ref. [21] to which we refer for further details. The first two operators involve four Higgs doublets and are

$$\begin{aligned} O_{\Phi\Box} &= (\Phi^\dagger\Phi)\Box(\Phi^\dagger\Phi), \\ O_{\Phi D} &= (\Phi^\dagger D^\mu\Phi)^*(\Phi^\dagger D_\mu\Phi). \end{aligned} \qquad (3)$$

They modify the Higgs-gauge couplings and entail a redefinition of the Higgs field to preserve canonically normalized kinetic terms. Operators of the form $X^2\Phi^2$, where $X = W^I, B$, generate anomalous couplings of the Higgs to ZZ, γZ, and WW. They are, explicitly,

$$\begin{aligned} O_{\Phi W} &= (\Phi^\dagger\Phi)W^I_{\mu\nu}W^{I\mu\nu}, \\ O_{\Phi B} &= (\Phi^\dagger\Phi)B_{\mu\nu}B^{\mu\nu}, \\ O_{\Phi WB} &= (\Phi^\dagger\tau^I\Phi)W^I_{\mu\nu}B^{\mu\nu}, \\ O_{\Phi\widetilde{W}} &= (\Phi^\dagger\Phi)\widetilde{W}^I_{\mu\nu}W^{I\mu\nu}, \\ O_{\Phi\widetilde{B}} &= (\Phi^\dagger\Phi)\widetilde{B}_{\mu\nu}B^{\mu\nu}, \\ O_{\Phi\widetilde{W}B} &= (\Phi^\dagger\tau^I\Phi)\widetilde{W}^I_{\mu\nu}B^{\mu\nu}. \end{aligned} \qquad (4)$$

Finally, three operators involving two fermion fields, that yield contact $HZ\ell\ell$ interactions as well as modifications to gauge-boson couplings to leptons, should be taken into account:

$$\begin{aligned} O^{(1)}_{\Phi\ell} &= (\Phi^\dagger i\overleftrightarrow{D}_\mu\Phi)(\bar{\ell}\gamma^\mu\ell), \\ O^{(3)}_{\Phi\ell} &= (\Phi^\dagger i\overleftrightarrow{D}^I_\mu\Phi)(\bar{\ell}\gamma^\mu\tau^I\ell), \\ O_{\Phi e} &= (\Phi^\dagger i\overleftrightarrow{D}_\mu\Phi)(\bar{e}\gamma^\mu e). \end{aligned} \qquad (5)$$

As input parameters we employ G_F (the Fermi constant as measured in $\mu \to e\nu_\mu\bar{\nu}_e$ decay), the Z mass m_Z, the electromagnetic coupling α_{em}, and the Higgs mass m_H. We trade the Lagragian parameters g, g', the Higgs self-coupling λ, and the classical Higgs vacuum expectation value v for combinations of the former. Dimension six corrections to our input parameters must be taken into account and are discussed in detail in [16]. In particular, a four-fermion operator not listed above contributes to the redefinition of G_F and must be considered [22].

Apart from SM contributions we include a single insertion of a dim-6 operator. We neglect $1/\Lambda^4$ terms in the square of the amplitudes (except for photon-pole enhanced terms, see Ref. [16]). The effects of the operators listed above can be summarized in the following effective Lagragian

$$\mathcal{L}_{\rm eff} \supset c^{(1)}_{ZZ}\, HZ_\mu Z^\mu + c^{(2)}_{ZZ} H\, Z_{\mu\nu}Z^{\mu\nu} + c_{Z\widetilde{Z}} H\, Z_{\mu\nu}\widetilde{Z}^{\mu\nu} + c_{AZ} H\, Z_{\mu\nu}A^{\mu\nu} + c_{A\widetilde{Z}} HZ_{\mu\nu}\widetilde{A}^{\mu\nu} + HZ_\mu \bar{\ell}\gamma^\mu (c_V + c_A\gamma_5)\,\ell + Z_\mu\bar{\ell}\gamma^\mu(g_V - g_A\gamma_5)\ell - g_{\rm em} Q_\ell A_\mu \bar{\ell}\gamma^\mu \ell. \quad (6)$$

The effective Lagrangian of HZZ interaction depends on the basis of $d = 6$ operators; the above Lagrangian is constructed from the complete and non-redundant operator basis of Ref. [6].

The effective couplings can be written in terms of the underlying dimension-six operators. We list here explicitly only three of them, namely the contact $HZ\ell\ell$ couplings $c_{V,A}$ and the CP-even anomalous $HZ\gamma$ coupling, since they enter the phenomenology discussed in the sequel. We have

$$\begin{aligned} c_V &= \sqrt{2}G_F\, m_Z\, \widehat{\alpha}^V_{\Phi\ell}, \\ c_A &= \sqrt{2}G_F\, m_Z\, \widehat{\alpha}^A_{\Phi\ell}, \\ c_{AZ} &= (\sqrt{2}G_F)^{1/2}\, \widehat{\alpha}_{AZ}, \end{aligned} \quad (7)$$

with

$$\begin{aligned} \widehat{\alpha}^V_{\Phi\ell} &= \widehat{\alpha}_{\Phi e} + \left(\widehat{\alpha}^{(1)}_{\Phi\ell} + \widehat{\alpha}^{(3)}_{\Phi\ell}\right), \\ \widehat{\alpha}^A_{\Phi\ell} &= \widehat{\alpha}_{\Phi e} - \left(\widehat{\alpha}^{(1)}_{\Phi\ell} + \widehat{\alpha}^{(3)}_{\Phi\ell}\right), \\ \widehat{\alpha}_{AZ} &= 2s_W\, c_W(\widehat{\alpha}_{\Phi W} - \widehat{\alpha}_{\Phi B}) + (s_W^2 - c_W^2)\widehat{\alpha}_{\Phi WB}, \end{aligned} \quad (8)$$

where s_W and c_W are respectively $\sin\theta_W$ and $\cos\theta_W$. The remaining couplings are given in detail in [16].

In the study of the phenomenological impact of the anomalous couplings one must take into account the constraints imposed on these couplings by the present LHC data, as well as EW precision data. We estimate [16] the following bounds for the contact $HZ\ell\ell$ couplings

$$\widehat{\alpha}^{V,A}_{\Phi\ell} \in [-5, 5] \times 10^{-3}. \quad (9)$$

Also, from the results of Ref. [23], one can deduce the following bounds within our conventions

$$\widehat{\alpha}_{AZ} \in [-1.3, 2.6] \times 10^{-2}. \quad (10)$$

We allow the above couplings to vary within these limits.

2. Angular asymmetries of $H \to Z(\to \ell^+\ell^-)\ell^+\ell^-$ and $e^+e^- \to HZ(\to \ell^+\ell^-)$

We discuss here the angular structures of the differential decay amplitude of $H \to Z(\to \ell^+\ell^-)\ell^+\ell^-$ and of the cross-section of $e^+e^- \to HZ(\to \ell^+\ell^-)$. Here we discuss in some detail the decay case; the description of the scattering can be done in close analogy exploiting crossing symmetry.

Summing over spins of the final-state leptons, the four-fold differential decay width for the process $H(p_H) \to Z(p)(\to \ell^-(p_1)\ell^+(p_2))\ell^-(p_3)\ell^+(p_4)$ in the massless lepton limit can be written as a function of the di-lepton invariant mass squared $q^2 = (p_3 + p_4)^2$ and of three angles. We chose the angles $\theta_{1,2}$, the angles between the direction of $\boldsymbol{p}_1$ and $\boldsymbol{p}_3$ and the z-axis in the respective di-lepton rest frames, and the angle ϕ between the normals of the di-lepton decay planes. The expression for the differential decay width reads

$$\frac{d^4\Gamma}{dq^2 d\cos\theta_1 d\cos\theta_2 d\phi} = \frac{1}{m_H}\mathcal{N}(q^2)\,\mathcal{J}(q^2,\theta_1,\theta_2,\phi), \quad (11)$$

with the normalization factor

$$\mathcal{N}(q^2) = \frac{1}{2^{10}(2\pi)^5}\frac{1}{\sqrt{r}\,\gamma_Z}\lambda^{1/2}(1, r, s), \quad (12)$$

written in terms of the dimensionless variables

$$s = \frac{q^2}{m_H^2},\ \ r = \frac{m_Z^2}{m_H^2} \approx 0.53, \ \ \gamma_Z = \frac{\Gamma_Z}{m_H} \approx 0.020, \quad (13)$$

and the function $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc$. The maximum value of q^2 is $q^2_{\rm max} = (m_H - m_Z)^2 \approx (34.4\ \mathrm{GeV})^2$ which gives

$$0 \le s \le \frac{(m_H - m_Z)^2}{m_H^2} \approx 0.075. \quad (14)$$

The function $\mathcal{J}(q^2,\theta_1,\theta_2,\phi)$ has nine independent angular structures with coefficient functions J_1,...,J_9, which we write

$$\begin{aligned} &\mathcal{J}(q^2,\theta_1,\theta_2,\phi) \\ &= J_1(1 + \cos^2\theta_1\cos^2\theta_2 + \cos^2\theta_1 + \cos^2\theta_2) \\ &+ J_2\sin^2\theta_1\sin^2\theta_2 + J_3\cos\theta_1\cos\theta_2 \\ &+ (J_4\sin\theta_1\sin\theta_2 + J_5\sin 2\theta_1\sin 2\theta_2)\sin\phi \\ &+ (J_6\sin\theta_1\sin\theta_2 + J_7\sin 2\theta_1\sin 2\theta_2)\cos\phi \\ &+ J_8\sin^2\theta_1\,\sin^2\theta_2\,\sin 2\phi \\ &+ J_9\sin^2\theta_1\,\sin^2\theta_2\,\cos 2\phi. \end{aligned} \quad (15)$$

The J functions are in turn written in terms of six form-factors that we denote $H_i^{V,A}$ (with $i = 1, 2, 3$)

$$
\begin{aligned}
J_1 &= 2\,r\,s\left(g_A^2 + g_V^2\right)\left(|H_{1,V}|^2 + |H_{1,A}|^2\right),\\
J_2 &= \kappa\left(g_A^2 + g_V^2\right)[\kappa\left(|H_{1,V}|^2 + |H_{1,A}|^2\right)\\
&\quad + \lambda \mathrm{Re}\left(H_{1,V}H_{2,V}^* + H_{1,A}H_{2,A}^*\right)],\\
J_3 &= 32\,r\,s\,g_A\,g_V\,\mathrm{Re}\left(H_{1,V}\,H_{1,A}^*\right),\\
J_4 &= 4\kappa\,\sqrt{r\,s\,\lambda}\,g_A\,g_V\,\mathrm{Re}\left(H_{1,V}H_{3,A}^* + H_{1,A}H_{3,V}^*\right),\\
J_5 &= \frac{1}{2}\kappa\,\sqrt{r\,s\,\lambda}\left(g_A^2 + g_V^2\right)\mathrm{Re}\left(H_{1,V}H_{3,V}^* + H_{1,A}\,H_{3,A}^*\right),\\
J_6 &= 4\sqrt{r\,s}\,g_A\,g_V\,[4\kappa\,\mathrm{Re}\left(H_{1,V}H_{1,A}^*\right)\\
&\quad + \lambda\,\mathrm{Re}\left(H_{1,V}H_{2,A}^* + H_{1,A}H_{2,V}^*\right)],\\
J_7 &= \frac{1}{2}\sqrt{r\,s}\left(g_A^2 + g_V^2\right)[2\kappa\left(|H_{1,V}|^2 + |H_{1,A}|^2\right)\\
&\quad + \lambda\,\mathrm{Re}\left(H_{1,V}H_{2,V}^* + H_{1,A}H_{2,A}^*\right)],\\
J_8 &= 2\,r\,s\,\sqrt{\lambda}\left(g_A^2 + g_V^2\right)\mathrm{Re}\left(H_{1,V}H_{3,V}^* + H_{1,A}H_{3,A}^*\right),\\
J_9 &= 2\,r\,s\left(g_A^2 + g_V^2\right)\left(|H_{1,V}|^2 + |H_{1,A}|^2\right),
\end{aligned}
\tag{16}
$$

where $\kappa = 1 - r - s$. The form factors parametrize the amplitude $\mathcal{M}^{\mu}_{HZ\ell\ell}$ for the decay $H \to Z\ell\ell$ as follows

$$
\begin{aligned}
\mathcal{M}^{\mu}_{HZ\ell\ell} &= \frac{1}{m_H}\,\bar{u}(p_3, s_3)\Big[\gamma^{\mu}\,(H_{1,V} + H_{1,A}\,\gamma_5)\\
&+ \frac{q^{\mu}\not{p}}{m_H^2}\,(H_{2,V} + H_{2,A}\,\gamma_5)\\
&+ \frac{\epsilon^{\mu\nu\sigma\rho}\,p_{\nu}q_{\sigma}}{m_H^2}\,\gamma_{\rho}\,(H_{3,V} + H_{3,A}\,\gamma_5)\Big]v(p_4, s_4),
\end{aligned}
\tag{17}
$$

where $\epsilon_{0123} = +1$ and $q = p_3 + p_4$. An important observation is that SM tree level contributions appear solely in $H_1^{V,A}$, the form factors $H_{2,3}^{V,A}$ are suppressed by $1/\Lambda^2$. Therefore, we neglect higher powers of $H_{2,3}^{V,A}$ in Eq. (16).

The cross section for $e^+e^- \to HZ$ is governed by the same set of form factors, in another kinematic regime and related by analyticity to those of $H \to Z\ell^+\ell^-$. We can cast the cross section as[1]

$$
\frac{d\sigma}{d\cos\theta_1\,d\cos\theta_2 d\phi} = \frac{1}{m_H^2}\,\mathcal{N}_{\sigma}(q^2)\,\mathcal{J}(q^2, \theta_1, \theta_2, \phi),
\tag{18}
$$

where the new normalisation reads

$$
\mathcal{N}_{\sigma}(q^2) = \frac{1}{2^{10}(2\pi)^3}\,\frac{1}{\sqrt{r}\,\gamma_Z}\,\frac{\sqrt{\lambda(1, s, r)}}{s^2}.
\tag{19}
$$

[1]For the precise definition of the three angles in this case we refer to the appendix A.2 of Ref. [16].

The threshold energy for the reaction is given by $\sqrt{q_{\rm th}^2} = (m_H + m_Z) \approx 217$ GeV which gives, in units of m_H^2, the minimal s value

$$
s_{\rm th} = q_{\rm th}^2/m_H^2 \approx 2.98.
\tag{20}
$$

The form factors are therefore probed at much higher energies, which leads to non-trivial phenomenological consequences in comparison with $H \to Z\ell^+\ell^-$.

Integrating over the three angles, the differential decay rate and the total cross section are proportional to the combination $4J_1 + J_2$. Angular asymmetries can be constructed to give us access to the information contained in other J functions. Here we show results for two of them, written explicitly for the case of the decay $H \to Z\ell^+\ell^-$. Following the notation of Ref. [16] we write

$$
\begin{aligned}
\mathcal{A}_{\phi}^{(3)} &= \frac{1}{d\Gamma/dq^2}\int_0^{2\pi} d\phi\ \mathrm{sgn}(\cos\phi)\,\frac{d^2\Gamma}{dq^2 d\phi}\\
&= \frac{9\,\pi}{32}\,\frac{J_6}{4J_1 + J_2},
\end{aligned}
\tag{21}
$$

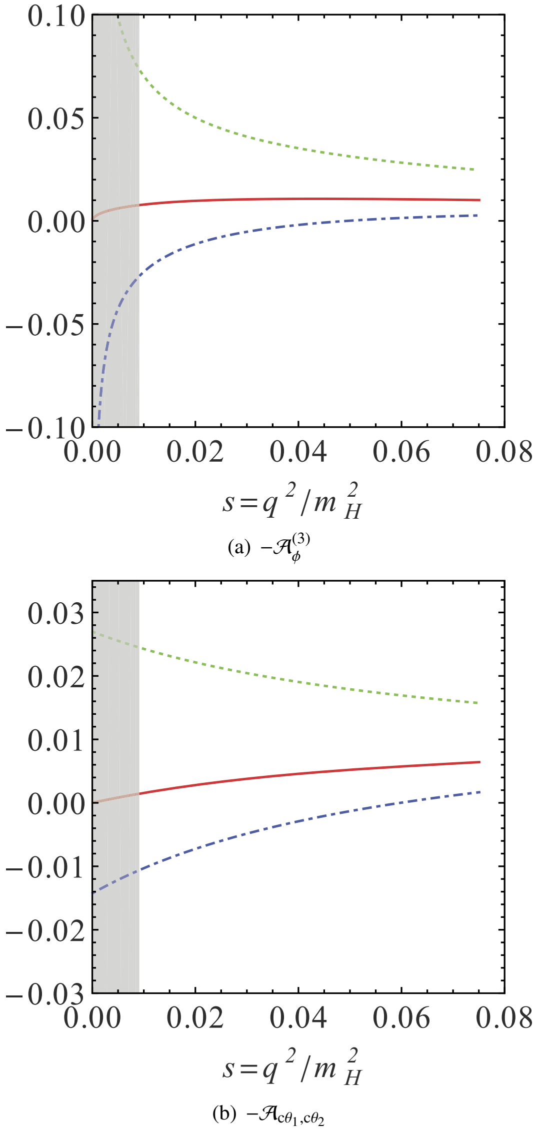
$$
\begin{aligned}
\mathcal{A}_{c\theta_1,c\theta_2} &= \frac{1}{d\Gamma/dq^2}\\
&\times \int_{-1}^{1} d\cos\theta_1\,\mathrm{sgn}(\cos\theta_1)\int_{-1}^{1} d\cos\theta_2\,\mathrm{sgn}(\cos\theta_2)\\
&\quad \frac{d^3\Gamma}{dq^2 d\cos\theta_1 d\cos\theta_2} = \frac{9}{16}\,\frac{J_3}{4J_1 + J_2}.
\end{aligned}
\tag{22}
$$

The sign function is $\mathrm{sgn}(\pm|x|) = \pm 1$. The asymmetries for $e^+e^- \to HZ$ can be written in close analogy to the ones above.

The functions J_3 and J_6 are sensitive to the contact $HZ\ell\ell$ interaction and to CP-even anomalous $HZ\gamma$ coupling, whose expressions are given in Eqs. (7) and (8). In the remainder we discuss briefly how these anomalous couplings impact the asymmetries in two specific cases. A thorough discussion of the asymmetries and couplings in both $H \to Z\ell^+\ell^-$ and $e^+e^- \to HZ$ is found in Ref. [16].

First, let us study the asymmetries of Eqs. (21) and (22) in a scenario where we consider the impact of $\widehat{\alpha}_{AZ}$; all other anomalous couplings are set to zero for the moment. The results of this exercise are shown in Fig. 1. The asymmetries are particularly sensitive to this anomalous coupling, in Fig. 1(a), $|\mathcal{A}_{\phi}^{(3)}|$ reaches $\sim 10\%$ for lower values of q^2, compared to an almost zero asymmetry in the SM. This relatively high sensitivity is, in the case of $\mathcal{A}_{\phi}^{(3)}$, due to the photon pole.

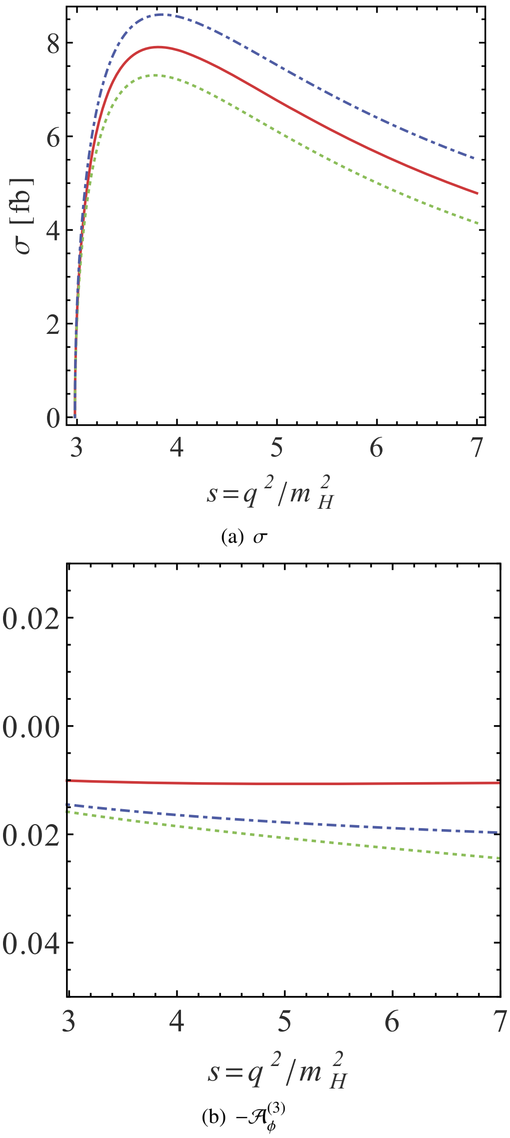
Another case of interest is the impact of the contact $HZ\ell\ell$ interactions in the cross section and asymmetries



(a) $-\mathcal{A}_{\phi}^{(3)}$

(b) $-\mathcal{A}_{c\theta_1,c\theta_2}$

Figure 1: (a) $-\mathcal{A}_{\phi}^{(3)}$, (b) $-\mathcal{A}_{c\theta_1,c\theta_2}$. Three scenarios are considered. The red solid-line is the SM case. The dot-dashed blue line corresponds to $\widehat{\alpha}_{AZ} = -1.3 \times 10^{-2}$, whereas the dotted green line corresponds to $\widehat{\alpha}_{AZ} = 2.6 \times 10^{-2}$. The gray band excludes the region $0 \leq q^2 \leq (12\ \text{GeV})^2$ where the decay $(Z^*, \gamma^*) \to \ell^+\ell^-$ is dominated by hadronic resonances.



(a) σ

(b) $-\mathcal{A}_{\phi}^{(3)}$

Figure 2: (a) $d\Gamma/ds$, (b) $-\mathcal{A}_{\phi}^{(3)}$. Three scenarios with the same $\widehat{\alpha}_{\Phi\ell}^{V}$ coupling are considered. The red solid-line is the SM case. The dotted green line corresponds to $(\widehat{\alpha}_{\Phi\ell}^{V}, \widehat{\alpha}_{\Phi\ell}^{A}) = (5, 5) \times 10^{-3}$, whereas the dot-dashed blue line to $(\widehat{\alpha}_{\Phi\ell}^{V}, \widehat{\alpha}_{\Phi\ell}^{A}) = (5, -5) \times 10^{-3}$.

of $e^+e^- \to HZ$. Fig. 2(a) shows that the total cross section is rather sensitive to the axial contact coupling. Such a sensitivity is a consequence of the higher values of s probed in the scattering and is not observed in the decay rate of $H \to Z\ell^+\ell^-$. The asymmetry of Fig. 2(b) can reach a few percent and is mainly sensitive to the vector contact coupling $\alpha_{\Phi\ell}^{V}$. We remark that the pattern of contributions to the different asymmetries can be well understood with the help of approximated analytic expressions given in Ref. [16].

3. Conclusions

- We identify several angular asymmetries, which are indeed very sensitive to anomalous couplings.
- Within the presently allowed range of the anomalous $HZ\gamma$ interaction strength, $\widehat{\alpha}_{AZ}$, modifications of angular asymmetries of $\mathcal{O}(1)$ and even larger relative to the SM value are still possible indicating sensitivity to multi-TeV scales.

- Anomalous $HZ\ell\ell$ contact interactions have smaller effects. This is mainly because we find that their size is already tightly constrained by existing data, in agreement with the constraints derived in Ref. [23]. The effects of the contact $HZ\ell\ell$ interactions in the angular asymmetries of $H \to Z\ell^+\ell^-$ were previously investigated in Ref. [15]. While we formally agree with their results, we find significantly smaller asymmetries, since the typical values of $\widehat{\alpha}^V_{\Phi\ell}$ adopted in that paper are about a factor of four larger than those allowed in the present analysis.
- At present, the CP-odd $d = 6$ couplings are not strongly constrained by data. We showed, in Ref. [16], that CP-odd asymmetry $\mathcal{A}^{(1)}_\phi$ can reach the few percent level in both in $H \to Z\ell^+\ell^-$ decay and $e^+e^- \to HZ$ Higgs production. In $H \to Z\ell^+\ell^-$ an asymmetry-zero may occur. However, for allowed values of the CP-odd couplings the asymmetry that can display this zero is never large.
- Most interesting asymmetries are small in absolute terms, reaching at most 10%, and often much less, because they are suppressed by the small vector $Z\ell\ell$ coupling.
- Overall, the process $e^+e^- \to HZ$ seems better suited than $H \to Z\ell^+\ell^-$ for the study of anomalous $HZ\ell\ell$ contact interactions due to the higher di-lepton invariant masses. This is particularly true for the contributions of $\widehat{\alpha}^A_{\Phi\ell}$ (as well as of $\widehat{\alpha}_{ZZ}$) to the total cross section, where 15% percent modifications are possible, which is illustrated in Fig. 2(a). On the other hand, $H \to Z\ell^+\ell^-$ provides better sensitivity to the anomalous $HZ\gamma$ coupling due to the photon-pole enhancement, as seen in Fig. 1(a).

In Ref. [16] we provided an estimate of SM loop effects, which suggests that the loop effects are small compared to the present bounds on $d = 6$ contributions. Obviously, a realistic extraction of $d = 6$ couplings from high-statistics data requires the inclusion of the SM loop contributions, which were computed in the past [24, 25, 26, 27].

Finally, our results show that the experimental detection of angular asymmetries will be challenging even with the planned higher statistics up-grades of the LHC.

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