

Adding Pseudo-Observables to the Four-Lepton Experimentalist's Toolbox

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ABSTRACT: The “golden” channel, in which the newly-discovered Higgs boson decays to four leptons by means of intermediate vector bosons, is important for determining the properties of the Higgs boson and for searching for subtle new physics effects. Different approaches exist for parametrizing the relevant Higgs couplings in this channel; here we relate the use of pseudo-observables to methods based on specifying the most general amplitude or Lagrangian terms for the HVV interactions. We also provide projections for sensitivity in this channel in several novel scenarios, illustrating the use of pseudo-observables, and analyze the role of kinematic distributions and (ratios of) rates in such $H \rightarrow 4\ell$ studies.

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1 Introduction

The discovery of a Standard Model (SM)-like Higgs boson (Higgs) at the CERN Large Hadron Collider (LHC) [1, 2] represents both the triumphant coda to one era, in which the SM was developed, tested, and confirmed, and the dawn of an exciting new era, in which physics beyond the SM (BSM) is probed, and eventually discovered. The “golden” channel, in which the Higgs decays to vector bosons (generally, Z bosons) that consequently decay to leptons, has played an important role in both the discovery of the Higgs and in the subsequent measurement of its couplings. It has therefore played an important role in both experimental [1–41] and theoretical [42–135] studies of the Higgs boson.

Measuring the couplings of the Higgs boson, and thus searching for the small deviations from SM predictions that may be the hallmark of BSM physics, requires a parameterization of these couplings. We can view this parameterization as being determined by:

- **Particles:** First we must specify the spectrum of particles that appear, either in the final state or as intermediate resonances.
- **Symmetries:** Next, we must determine the symmetries that the underlying theory obeys, either exactly or approximately.
- **Interactions:** Once the particles and symmetries have been specified, we can enumerate the possible interactions, all of which we expect to be present at some level. However, in general, there are additional considerations, such as the dimensionality of operators in an effective field theory (EFT) Lagrangian, which may allow us to consider only a finite number of terms.

In this paper we will consider parameterizations relevant to experimental studies of the four-lepton final state at the LHC. These parameterizations are based on the choices of

1. **Particles:** We consider only the SM degrees of freedom, with an additional (optional) heavy Z' boson, which is introduced as a BSM example of a concrete BSM scenario that allows for different couplings to electrons and muons. In particular, we do not consider the effects of light Z' s [28, 96, 98, 101, 103, 106, 108, 119, 120], dark matter, or other invisible or unidentified particles in the final state [136–144].
2. **Symmetries:** There have been many studies of Higgs couplings that make the (highly theoretically motivated) assumption that the Higgs is part of an $SU(2)$ doublet and the only allowed operators are part of the SM $SU(3) \times SU(2) \times U(1)$ gauge symmetry. We, however, make fewer theoretical assumptions as our aim is to facilitate relatively unbiased experimental measurements. If, as expected, potential interactions, which violate the above assumptions, are absent, we have obtained additional experimental support for our pre-existing viewpoint. Thus, we will assume only Lorentz invariance and QED and QCD gauge invariance in specifying possible interactions.
3. **Interactions:** Having made these choices of particles and symmetries, the allowed interactions are restricted. We will generally truncate the (infinite) list of possible interactions by considering operators only up to some finite mass dimension or (equivalently) amplitude structures only up to a specified power of momentum.

Experimental attempts to measure the tensor structure of Higgs to diboson couplings at the LHC (e.g., Ref. [22]) have largely used two formalisms. One is the specification of the most general $H \rightarrow VV$ amplitude allowed by the relevant symmetries [58, 73]. The second is to specify the interaction Lagrangian terms consistent with these same symmetries [84, 99]. Two important points are in order: (i) these approaches have different strengths and (ii) the approaches are nearly, though not completely, equivalent in the physics they parameterize, given certain reasonable physical assumptions.

Strengths of the amplitude approach include a greater ability to parametrize absorptive (*i.e.*, complex) contributions to the amplitude arising from loop-induced contributions from light particles running in the loop.¹ The Lagrangian approach, on the other hand, is more useful when utilizing standard computational tools, such as FeynRules [145, 146], CalcHEP [147, 148], CompHEP [149–151], and MadGraph [152–155]. It may also make the connection with more comprehensive physical frameworks more straightforward.

Recently, attention has been drawn to a third approach based on the use of pseudo-observables (PO) [156–158]. This approach, which was employed at the CERN Large Electron Positron (LEP) Collider [159, 160], represents a generalization of the amplitude approach and thus it shares its strengths. However, by demanding that the amplitude considered describe a transition between on-shell states, the PO method resolves several

¹ The magnitude of such contributions can be constrained using even conservative assumptions about the total Higgs width, as discussed, for example in Ref. [92].

potential ambiguities and, in particular, guarantees the gauge invariance of the quantities that are being measured. This property will become increasingly important as measurements increase in sensitivity, probing contributions from NLO effects in the SM [161–164], as well as potential new physics. The PO approach covers also BSM scenarios that affect the $H \rightarrow 4\ell$ channel without two intermediate Z or γ states. A simple example is a heavy Z' boson in $H \rightarrow ZZ' \rightarrow 4\ell$, but more complicated cases with additional heavy exotic intermediate states (possibly coupled in a non-universal way to SM fermions) are also covered by this approach.

Thus PO provide a useful third approach to parameterizing Higgs couplings. However, as there are also advantages to the amplitude and the Lagrangian frameworks (besides the important fact that their use by the experiments is well-established [39, 40]), it is useful to describe how the PO approach compares to these other methods. We do that in this paper. In Section 2 we discuss the theoretical relationship between various parameterizations. In particular, we specify the translation between the conventions used in each of these approaches in detail, so that results obtained in the context of one framework can be easily understood in the context of the other frameworks. In Section 3, we analyze the future sensitivity to Higgs pseudo-observables within specific benchmark scenarios with only two free parameters. We discuss in some detail the role of the different experimental inputs, namely the total number of events observed, the ratios between different channels and the kinematic distribution of the events. It is hoped that these results will be informative in their own right, but will also provide a useful template for future experimental studies.

2 Frameworks, Conventions, and Translation

2.1 The “Amplitude” Approach

In the “amplitude” approach [49, 58, 59, 73] that generally has been used in four-lepton studies, the most general form of the $1 \rightarrow 2$ vertex describing the decay of the putative Higgs boson to vector bosons is specified. (Of course, for the Standard Model Higgs the decay to Z bosons is the most important, though interesting phenomenology can result from decays to $Z\gamma^*$ and $\gamma^*\gamma^*$ [100].) The resulting expression can be used to obtain helicity amplitudes for the $H \rightarrow VV$ decay;² these can be combined with helicity amplitudes³ for the decay of the vector bosons to leptons to obtain the amplitude, and ultimately the differential cross section, for $gg \rightarrow H \rightarrow VV \rightarrow 4\ell$ (or processes where the Higgs is produced via other mechanisms). Following Ref. [22] (and references therein) we can write the general form of this amplitude as

$$A(H \rightarrow V_1 V_2) = \frac{i}{v} \left\{ \left[a_1^{V_1 V_2} + \frac{\kappa_1^{V_1 V_2} q_1^2 + \kappa_2^{V_1 V_2} q_2^2}{\left(\Lambda_1^{V_1 V_2} \right)^2} \right] m_{V_1}^2 \epsilon_1^* \cdot \epsilon_2^* \right. \quad (2.1)$$

² Here and below we suppress explicit indications that the vector bosons are, in general, off-shell.

³ We assume that Z and γ couple to leptons as in the Standard Model. We include also a possible contribution from a Z' boson. We allow its coupling to leptons to be an arbitrary sum of vector and axial vector contributions.

$$+ a_2^{V_1 V_2} f_{\mu\nu}^{*(1)} f^{*(2),\mu\nu} + a_3^{V_1 V_2} f_{\mu\nu}^{*(1)} \tilde{f}^{*(2),\mu\nu} \},$$

where ϵ_i is the polarization vector and q_i is the momentum of the gauge boson labelled “ V_i ” ($i = 1, 2$). The contribution to the amplitude corresponding to the field strength tensor is $f^{(i)\mu\nu} = \epsilon_i^\mu q_i^\nu - \epsilon_i^\nu q_i^\mu$, and its dual is given by $\tilde{f}^{(i)\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\alpha\beta} f^{(i),\alpha\beta}$. Symmetry considerations and gauge invariance require $\kappa_1^{ZZ} = \kappa_2^{ZZ}$ and $a_1^{Z\gamma}, a_1^{\gamma\gamma}, \kappa_1^{\gamma\gamma}, \kappa_2^{\gamma\gamma}, \kappa_1^{Z\gamma} = 0$.

We can proceed to find the helicity amplitudes that correspond to Eq. (2.1) using explicit expressions for the polarization vectors ϵ_1 and ϵ_2 . In the case of the four-charged-lepton final state, the relevant $H \rightarrow V_1 V_2$ sub-amplitudes are those where $V_1 V_2$ refers to $ZZ, Z\gamma, \gamma\gamma, ZZ',$ etc. It is important to note that terms in the general amplitude that vanish for on-shell vector bosons do not necessarily vanish for off-shell vector bosons. In particular, one can have a $q_\gamma^2 \epsilon_\gamma^* \epsilon_Z^*$ contribution, which, of course, vanishes when the photon is on-shell.

In general, the “constants” in Eq. (2.1) can be taken to be Lorentz invariant functions of q_1 and q_2 . (When the intermediate gauge bosons are identical, the function must also be invariant under the exchange of the labels V_1 and V_2 .) An arbitrary analytic Lorentz invariant function of momenta may be expressed as

$$f(q_1, q_2) = f_0 + \frac{1}{\Lambda^2} (f_{21} q_1^2 + f_{22} q_2^2 + f_{23} (q_1 + q_2)^2) + \mathcal{O}(\Lambda^{-4}), \quad (2.2)$$

where we assume that $f(q_1, q_2)$ can be expanded in a Taylor series. When V_1 and V_2 are produced in the decay of an on-shell boson, the situation considered in this work, we can set $(q_1 + q_2)^2 = m_H^2$, where m_H is the mass of the (Higgs) boson that decays to V_1 and V_2 . Thus we see that Eq. (2.1) is written to include the momentum dependence of the a_1^{VV} through $\mathcal{O}(\Lambda^{-2})$ terms, while only the leading constant is included for a_2^{VV} and a_3^{VV} . This apparent inconsistency is resolved when we realize that the amplitude structures associated with a_2^{VV} and a_3^{VV} involve two additional powers of momenta relative to the a_1^{VV} structure ($\epsilon_1^* \cdot \epsilon_2^*$). Therefore, if we take the a_i^{VV} , the κ_i , etc. to be true constants, the expression in Eq. (2.1) includes the full momentum dependence of the amplitude (in the case where the Higgs is on-shell) containing up to two powers of gauge boson momenta. We will see similar truncations in the cases of the “Lagrangian” and the “PO” approaches discussed below.

2.2 The “Lagrangian” Approach

The $H \rightarrow VV$ interactions can also be described in the Lagrangian formalism. For $H \rightarrow 4\ell$, the relevant interaction Lagrangian is

$$\begin{aligned} \mathcal{L} \supset & -\kappa_{1,ZZ} \frac{M_Z^2}{v} H Z_\mu Z^\mu - \frac{\kappa_{2,ZZ}}{2v} H Z_{\mu\nu} Z^{\mu\nu} - \frac{\kappa_{3,ZZ}}{2v} H Z_{\mu\nu} \tilde{Z}^{\mu\nu} \\ & + \frac{\kappa_{4,ZZ} M_Z^2}{M_H^2 v} \square H Z_\mu Z^\mu + \frac{2\kappa_{5,ZZ}}{v} H Z_\mu \square Z^\mu \\ & - \frac{\kappa_{2,Z\gamma}}{2v} H Z_{\mu\nu} F^{\mu\nu} - \frac{\kappa_{3,Z\gamma}}{2v} H Z_{\mu\nu} \tilde{F}^{\mu\nu} + \frac{\kappa_{5,Z\gamma} M_Z^2}{M_H^2 v} H Z_\mu \partial_\nu F^{\mu\nu} \\ & - \frac{\kappa_{2,\gamma\gamma}}{2v} H F_{\mu\nu} F^{\mu\nu} - \frac{\kappa_{3,\gamma\gamma}}{2v} H F_{\mu\nu} \tilde{F}^{\mu\nu} \end{aligned} \quad (2.3)$$

$$-\kappa_{1,ZZ'} \frac{M_Z^2}{v} H Z_\mu Z'^\mu,$$

where $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the QED field strength tensor; $Z_{\mu\nu} = \partial_\mu Z_\nu - \partial_\nu Z_\mu$ is defined by analogy. The dual field strength tensors \tilde{F} and \tilde{Z} are defined by (e.g.) $\tilde{F}_{\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\alpha\beta} F^{\alpha\beta}$, where ϵ is the totally antisymmetric Levi-Civita tensor. With the exception of interactions involving a Z' , this Lagrangian contains all operators up to dimension five (six, if one assumes that the operators involve the vacuum expectation value of the Higgs field). We note that there are five HZZ operators in Eq. (2.3), but fewer operators for $HZ\gamma$, $H\gamma\gamma$, and HZZ' . In the case of $HZ\gamma$ and $H\gamma\gamma$, this is due to the requirements provided by QED gauge invariance. For the Z' we have included only the lowest dimensional operator, as amplitudes involving the heavy Z' are suppressed by the large mass in the propagator. Had we chosen to consider a light Z' , we would need to include the additional HZZ' operators

$$\begin{aligned} \mathcal{L} \supset & -\frac{\kappa_{2,ZZ'}}{2v} H Z_{\mu\nu} Z'^{\mu\nu} - \frac{\kappa_{3,ZZ'}}{2v} H Z_{\mu\nu} \tilde{Z}'^{\mu\nu} \\ & + \frac{\kappa_{4,ZZ'} M_Z^2}{M_H^2 v} \square H Z_\mu Z'^\mu + \frac{2\kappa_{5,ZZ'}}{v} H Z_\mu \square Z'^\mu, \end{aligned} \quad (2.4)$$

the $HZ'\gamma$ operators

$$\mathcal{L} \supset -\frac{\kappa_{2,Z'\gamma}}{2v} H Z'_{\mu\nu} F^{\mu\nu} - \frac{\kappa_{3,Z'\gamma}}{2v} H Z'_{\mu\nu} \tilde{F}^{\mu\nu} + \frac{\kappa_{5,Z'\gamma} M_Z^2}{M_H^2 v} H Z'_\mu \partial_\nu F^{\mu\nu} \quad (2.5)$$

and the $HZ'Z'$ operators

$$\begin{aligned} \mathcal{L} \supset & -\kappa_{1,Z'Z'} \frac{M_Z^2}{v} H Z'_\mu Z'^\mu - \frac{\kappa_{2,Z'Z'}}{2v} H Z'_{\mu\nu} Z'^{\mu\nu} - \frac{\kappa_{3,Z'Z'}}{2v} H Z'_{\mu\nu} \tilde{Z}'^{\mu\nu} \\ & + \frac{\kappa_{4,Z'Z'} M_Z^2}{M_H^2 v} \square H Z'_\mu Z'^\mu + \frac{2\kappa_{5,Z'Z'}}{v} H Z'_\mu \square Z'^\mu. \end{aligned} \quad (2.6)$$

However, assuming the Z' mass to be comparable to the cutoff of the theory, we can take all of the operators in Eqs. (2.4–2.6) to be higher dimensional and therefore ignore them.⁴ For *on-shell* Higgs studies (which are our focus here), the $(\square H)Z_\mu Z^\mu$ operator can be absorbed into the $HZ_\mu Z^\mu$ operator, though this will not be the case when the Higgs is off-shell [99]. With the caveat that the coefficients $\{\kappa_{i,VV}\}$ in Eq. (2.3) are real, there is a one-to-one mapping between the coefficients in Eqs. (2.3) and (2.1), as shown explicitly in Table 1 (see also Ref. [99]).

While we assume that Z bosons and photons couple to leptons as in the SM, we can make no such assumptions about the interaction of the Z' with leptons. Thus, to fully characterize $H \rightarrow 4\ell$ decays, we parameterize the Z' interaction to SM leptons with

$$\mathcal{L} \supset \sum_{\ell=L,R} g_{Z'}^\ell \bar{\ell} \gamma^\mu \ell Z'_\mu. \quad (2.7)$$

⁴ An alternate and perhaps clearer approach would be to integrate out the heavy Z' . We have chosen the approach above, with explicit $\kappa_{1,ZZ'}$ term in the interest of consistency with parameterizations already in use by the experiments.

2.3 The ‘‘Pseudo-Observables’’ Approach

The PO formalism described here is based on the specification of the pole structure of the amplitude for the entire $H \rightarrow 4\ell$ process. As we are considering an on-shell Higgs, and we single out the pole terms due to the exchange of the SM gauge bosons, the resulting amplitude is automatically gauge invariant, which is an advantage with respect to the ‘‘amplitude’’ approach described in Subsection 2.1, especially when considering the effects of higher order corrections [161].⁵ Furthermore, the PO formalism, like the amplitude formalism, can accommodate arbitrary loop corrections in a relatively straightforward way.

Following Ref. [156], the goal of $H \rightarrow 4\ell$ studies in the PO approach is to characterize in general terms the correlation function of the Higgs and two fermion currents,

$$\langle 0 | \mathcal{T} \{ H(0), J_\ell^\mu(x), J_{\ell'}^\nu(y) \} | 0 \rangle, \quad (2.8)$$

where $\ell, \ell' = e, \mu$. Assuming Lorentz invariance and reasonable assumptions about flavor conservation, the matrix element corresponding to this correlation function can be described by

$$\begin{aligned} \mathcal{A}_{n.c.} [X \rightarrow \ell^-(p_1) \ell^+(p_2), \ell'^-(p_3) \ell'^+(p_4)] &= i \frac{2m_Z^2}{v} \sum_{\ell=\ell_L, \ell_R} \sum_{\ell'=\ell'_L, \ell'_R} (\bar{\ell}' \gamma_\rho \ell) (\bar{\ell}' \gamma_\sigma \ell') \mathcal{T}^{\rho\sigma}(q_1, q_2) \\ \mathcal{T}^{\rho\sigma}(q_1, q_2) &= \left[F_1^{\ell\ell'}(q_1^2, q_2^2) g^{\rho\sigma} + F_3^{\ell\ell'}(q_1^2, q_2^2) \frac{q_1 \cdot q_2 g^{\rho\sigma} - q_2^\rho q_1^\sigma}{m_Z^2} \right. \\ &\quad \left. + F_4^{\ell\ell'}(q_1^2, q_2^2) \frac{\varepsilon^{\rho\sigma\alpha\beta} q_{V_2\alpha} q_{V_1\beta}}{m_Z^2} \right], \end{aligned} \quad (2.9)$$

where $q_1 = p_1 + p_2$ and $q_2 = p_3 + p_4$. Recognizing the presence of physical poles in the correlation function (2.8) due to the propagation of intermediate SM gauge bosons, we expand around these poles and define the PO directly from their residues:

$$\begin{aligned} F_1^{\ell\ell'}(q_1^2, q_2^2) &= \kappa_{ZZ} \frac{g_Z^\ell g_Z^{\ell'}}{P_Z(q_1^2) P_Z(q_2^2)} + \frac{\epsilon_{Z\ell}}{m_Z^2} \frac{g_Z^{\ell'}}{P_Z(q_2^2)} + \frac{\epsilon_{Z\ell'}}{m_Z^2} \frac{g_Z^\ell}{P_Z(q_1^2)} \\ &\quad + \Delta_1^{\text{SM}}(q_1^2, q_2^2), \end{aligned} \quad (2.10)$$

$$\begin{aligned} F_3^{\ell\ell'}(q_1^2, q_2^2) &= \epsilon_{ZZ} \frac{g_Z^\ell g_Z^{\ell'}}{P_Z(q_1^2) P_Z(q_2^2)} - \epsilon_{Z\gamma} \left(\frac{e g_Z^\ell}{q_2^2 P_Z(q_1^2)} + \frac{e g_Z^{\ell'}}{q_1^2 P_Z(q_2^2)} \right) \\ &\quad + \frac{\epsilon_{\gamma\gamma} e^2}{q_1^2 q_2^2} + \Delta_3^{\text{SM}}(q_1^2, q_2^2), \end{aligned} \quad (2.11)$$

$$\begin{aligned} F_4^{\ell\ell'}(q_1^2, q_2^2) &= \epsilon_{ZZ}^{\text{CP}} \frac{g_Z^\ell g_Z^{\ell'}}{P_Z(q_1^2) P_Z(q_2^2)} - \epsilon_{Z\gamma}^{\text{CP}} \left(\frac{e g_Z^\ell}{q_2^2 P_Z(q_1^2)} + \frac{e g_Z^{\ell'}}{q_1^2 P_Z(q_2^2)} \right) \\ &\quad + \frac{\epsilon_{\gamma\gamma}^{\text{CP}} e^2}{q_1^2 q_2^2}, \end{aligned} \quad (2.12)$$

⁵ If we were to consider contributions to high invariant mass four-lepton events from an off-shell Higgs boson [69, 88, 95, 99, 165–180], we would need to construct PO for the whole $gg \rightarrow 4\ell$ process, including both the ‘‘signal’’ $gg \rightarrow H \rightarrow VV \rightarrow 4\ell$ amplitude and the ‘‘background’’ $gg \rightarrow VV \rightarrow 4\ell$ amplitudes with which it interferes.

where $g_Z^{\ell,\ell'}$ correspond to the well-measured LEP-I Z -pole PO,

$$\mathcal{A}(Z(\varepsilon) \rightarrow \ell^+ \ell^-) = i \sum_{\ell=\ell_L, \ell_R} g_Z^{\ell} \varepsilon_{\nu} \bar{\ell} \gamma^{\nu} \ell, \quad (2.13)$$

and $P_Z(q^2) = q^2 - m_Z^2 + im_Z \Gamma_Z$. Note the minus signs in Eqs. (2.11) and (2.12) occur because the relevant expressions contain a Q for the fermion charge in Ref. [156], which is equal to minus one for electrons and muons. In this decomposition around physical poles, we neglected further terms which are necessarily generated by local operators of dimension greater than 6 and thus strongly suppressed if the new physics scale is above the electroweak scale. The parameters g_Z^f , κ_{ZZ} , and ϵ_X are thus well-defined PO; they can be measured in experiments and their values can be calculated at any given order in perturbation theory in any underlying theory (or effective field theory).⁶

While the amplitude for $H \rightarrow 2e2\mu$ is given directly by Eq. (2.9), the decays $H \rightarrow 4e, 4\mu$ include an interference between two contributions, corresponding to the two possible assignments of the fermions inside the currents:

$$\begin{aligned} \mathcal{A} [h \rightarrow \ell(p_1) \bar{\ell}(p_2) \ell(p_3) \bar{\ell}(p_4)] &= \mathcal{A}_{n.c.} [h \rightarrow f(p_1) \bar{f}(p_2) f'(p_3) \bar{f}'(p_4)]_{f=f'=\ell} \\ &\quad - \mathcal{A}_{n.c.} [h \rightarrow f(p_1) \bar{f}(p_4) f'(p_3) \bar{f}'(p_2)]_{f=f'=\ell}. \end{aligned} \quad (2.14)$$

The tree-level connection between the PO and the parameters used in Eqs. (2.1) and (2.3) is given in Table 1. Of these parameters, only $\epsilon_{Ze_{L,R}}$ and $\epsilon_{Z\mu_{L,R}}$ can depend on the flavor of the final-state leptons. In the limit, where we neglect loop contributions from light states, κ_{ZZ} and ϵ_X are all real. The functions $\Delta_{1,3}^{\text{SM}}(q_1^2, q_2^2)$ describe loop-induced SM contributions, which cannot be described in terms of $D \leq 6$ effective operators; an explicit expression for $\Delta_3^{\text{SM}}(q_1^2, q_2^2)$ at one-loop can be found in Ref. [156], where it is shown that their size is small and thus can be neglected given experimental precision in the foreseeable future.

The tree-level matching of the operators in the Lagrangian approach to the PO, which is shown in Table 1, is obtained by computing the same decay amplitude and rearranging the terms so that the expansion around physical poles is recovered. For example, the $HZ_{\mu} \partial_{\nu} F^{\mu\nu}$ operator in Eq. (2.3) corresponds to an amplitude with a q_{γ}^2 in the numerator (from the \square operator). This q_{γ}^2 cancels the photon propagator in the denominator, so this operator contributes to ϵ_{Ze} and $\epsilon_{Z\mu}$. The contribution of the heavy Z' is evaluated in the limit $q_V^2 \ll m_{Z'}^2$ (consistently with our general assumptions). In such limit the Z' can easily be integrated out, generating a contribution to the contact terms with the flavor structure given by its couplings to fermions. Let us stress, however, that the Z' is just a specific example of a BSM scenario that generates flavor-dependent contact terms. The PO approach is not restricted in any sense to this particular case, in contrast to the Amplitude or Lagrangian approaches.

2.4 Global Approaches to Higgs Couplings

While our interest in this paper is in detailed parameterizations of BSM effects in $H \rightarrow 4\ell$, we note in passing that much work has gone into “global” studies of Higgs couplings to

⁶ Here we generically denote by ϵ_X the parameters $\epsilon_{ZZ, Z\gamma, \gamma\gamma, Zf}$ and $\epsilon_{ZZ, Z\gamma, \gamma\gamma}^{\text{CP}}$.

PO	Lagrangian parameter	Amplitude parameters
κ_{ZZ}	$-\kappa_{1,ZZ} - \kappa_{4,ZZ} - 2\kappa_{5,ZZ}$	$\frac{1}{2}a_1^{ZZ} + \frac{m_Z^2}{(\Lambda_1^{ZZ})^2}\kappa_1^{ZZ}$
ϵ_{ZZ}	$\kappa_{2,ZZ}$	a_2^{ZZ}
$\epsilon_{ZZ}^{\text{CP}}$	$\kappa_{3,ZZ}$	a_3^{ZZ}
$\epsilon_{Z\gamma}$	$\kappa_{2,Z\gamma}/2$	$a_2^{Z\gamma}$
$\epsilon_{Z\gamma}^{\text{CP}}$	$\kappa_{3,Z\gamma}/2$	$a_3^{Z\gamma}$
$\epsilon_{\gamma\gamma}$	$\kappa_{2,\gamma\gamma}$	$a_2^{\gamma\gamma}$
$\epsilon_{\gamma\gamma}^{\text{CP}}$	$\kappa_{3,\gamma\gamma}$	$a_3^{\gamma\gamma}$
ϵ_{Zf}	$-g_Z^f \kappa_{5,ZZ} + g_{Z'}^f \frac{\kappa_{1,ZZ'}}{2} \frac{m_Z^2}{m_{Z'}^2} - e \frac{\kappa_{5,Z\gamma}}{2} \frac{m_Z^2}{m_H^2}$	$g_Z^f \frac{\kappa_1^{ZZ}}{2(\Lambda_1^{ZZ})^2} \frac{m_Z^2}{m_Z^2} - g_{Z'}^f \frac{a_1^{ZZ'}}{2m_{Z'}^2} \frac{m_Z^2}{m_Z^2} - e \frac{\kappa_2^{Z\gamma}}{2(\Lambda_1^{Z\gamma})^2} \frac{m_Z^2}{m_Z^2}$

Table 1: In this table we provide a “dictionary”, allowing one to convert (at the tree level) between (i) the “amplitude” formalism of Eq. (2.1), (ii) the “Lagrangian” formalism of Eq. (2.3), and (iii) the “pseudo-observables” formalism of Eq. (2.9). The label f is a generic label for a fermion ($f = e_L, e_R, \mu_L, \mu_R$).

various final states, whether in the so-called “ κ ” or “signal-strength” formalism [181–188]) or in the context of SM effective theories [80, 168, 173, 189–250]. While such studies are clearly of value, we wish to emphasize that our goal here is quite different. Specifically, our goal is to allow a description of BSM effects in $H \rightarrow 4\ell$ that is as theoretically unbiased as possible using as much information as possible.

Generally, more global approaches to Higgs couplings make numerous additional assumptions and use less information about a specific process, such as Higgs to four leptons, than is experimentally available. As the assumptions in these studies tend to be theoretically reasonable, such work is extremely useful in providing a “big picture” view of what many different channels are telling us about the Higgs, which is complementary to our goal of extracting the most possible information from a single channel. To give an indication of what is possible in such analyses and to illustrate the use of the various formalisms described here, particularly the PO framework presented in Subsection 2.3, we present projections, in Section 3, for the sensitivity of LHC measurements in these frameworks.

3 LHC Projections

3.1 Benchmark scenarios

While the PO are directly related to physical properties of the Higgs decay amplitudes, the generic expectation is that explicit new physics models would contribute to some combination of PO. For this reason, the experimental analysis should be performed with the most general case possible in mind. In order to reduce the number of independent parameters, the best option is to impose relations due to specific symmetries [156].

With this important caveat in mind, but also with an eye to studying the sensitivity of the LHC and other future colliders for measuring Higgs PO in $H \rightarrow 4\ell$ decays, let us list

the following six simplified scenarios, each of which has, conveniently, only two independent PO:

1. **Longitudinal vs. transverse:** (κ_{ZZ} vs. ϵ_{ZZ}), where all other $\epsilon \rightarrow 0$.
2. **CP admixture:** (κ_{ZZ} vs. $\epsilon_{ZZ}^{\text{CP}}$), where all other $\epsilon \rightarrow 0$.
3. **Linear EFT-inspired [114]:** (κ_{ZZ} vs. $\epsilon_{Z\ell_R}$), where $\epsilon_{Z\ell_L} = 2\epsilon_{Z\ell_R}$, $\epsilon_{Ze_{L,R}} = \epsilon_{Z\mu_{L,R}}$, and other $\epsilon \rightarrow 0$ (or κ_{ZZ} vs. some other combination of contact terms).
4. **Flavor universal contact terms:** ($\epsilon_{Z\ell_L}$ vs. $\epsilon_{Z\ell_R}$), where $\epsilon_{Ze_{L,R}} = \epsilon_{Z\mu_{L,R}}$, $\kappa_{ZZ} = 1$, and other $\epsilon \rightarrow 0$.
5. **Flavor non-universal vector contact terms:** (ϵ_{Ze_R} vs. $\epsilon_{Z\mu_R}$), where $\epsilon_{Ze_L} = \epsilon_{Ze_R}$, $\epsilon_{Z\mu_L} = \epsilon_{Z\mu_R}$, $\kappa_{ZZ} = 1$, and other $\epsilon \rightarrow 0$.
6. **Flavor non-universal axial contact terms:** (ϵ_{Ze_R} vs. $\epsilon_{Z\mu_R}$), where $\epsilon_{\ell_L} = -\epsilon_{Z\ell_R}$, $\kappa_{ZZ} = 1$, and other $\epsilon \rightarrow 0$.

Scenarios 1 and 2 have been the focus of much of the existing work in four-lepton phenomenology and experiment, whereas the rest of cases have received much less attention. Scenarios 1, 2, and 3 include κ_{ZZ} , which only affects the overall normalization of all $H \rightarrow 4\ell$ channels and which consequently can only be probed through its effect in the total rates. This in turn requires the implicit assumption that other BSM effects affecting Higgs production are absent. For this reason in the next subsections we focus on the last three scenarios, where kinematic distributions and ratios of rates also provide interesting information on all PO under study. This study will serve as an example of the use of the PO formalism. These scenarios all involve contact operators that couple the Higgs to an intermediate boson and two leptons. They could arise, *e.g.*, via a heavy Z' that is integrated out, but also via other mechanisms that might not even involve two intermediate bosons. In scenario 4, this interaction is of the same strength for electrons and muons; the helicity structure of the coupling gives us two independent parameters. In scenarios 5 and 6, we fix the helicity structure of the couplings to be vector or axial, respectively, and make the couplings to electrons and muons the two independent parameters. The violation of lepton-flavor universality makes these two scenarios qualitatively different from the rest, both conceptually and in practice, as we will see in the phenomenological analyses below.

3.2 Projections Using Rates and Kinematics

The ATLAS collaboration has recently published the first $h \rightarrow 4\ell$ analysis in the PO framework using 36.1 fb^{-1} of data at 13 TeV [251] (see also [252]). They work in the (lepton flavor universal) benchmark scenarios 3 and 4, using the binned (m_{12}, m_{34}) invariant mass distribution as experimental input. Let us stress that they do not work with normalized distributions, and thus they are not only sensitive to the effects of the PO in the shapes but also in the total rates. As mentioned above, this assumes SM-like Higgs production, which allows them to probe κ_{ZZ} in scenario 3.

In this section we expand the scope of this ATLAS search and go one step beyond, highlighting the importance of lepton flavor universality tests. Namely, we split the (m_{12}, m_{34}) histogram, shown in Fig. 3 (right) of [251], into four categories based on the lepton flavor: $4e$, 4μ , $2e2\mu$, and $2\mu2e$.⁷ Apart from that, we closely follow the ATLAS analysis of Ref. [251] and obtain projections for our scenarios 4, 5 and 6.

Let us first explain our simulation procedure. The signal events ($pp \rightarrow h \rightarrow 4\ell$) are generated using the `HiggsPO UFO` model from Ref. [164] within the `MadGraph5_aMC@NLO` framework [256]. Reliable gluon-gluon fusion production kinematics is obtained with the leading order matrix element and parton shower jet merging, while the normalization factor ($K_F = 2.32$) is taken from the best higher-order QCD prediction [253]. Subsequent showering and hadronisation effects are simulated with `Pythia 6` [254], while the detector effects are simulated with `Delphes 3` [255]. Event samples are generated for enough points in the PO parameter space allowing for the reconstruction of the quadratic dependence of any observable (to be discussed below). The dominant background, coming from $pp \rightarrow ZZ^*(\gamma^*) \rightarrow 4\ell$, follows a similar simulation pipeline. We estimate the NLO QCD effects for the background by computing the K -factor in the signal region ($K_F = 1.3$), and using it to rescale the LO simulation.

We closely follow the event selection of the ATLAS search. The signal region is defined as follows. Four leptons ($\ell = e, \mu$) are selected to make a lepton quadruplet. Electrons are required to have $E_T > 7$ GeV and $\eta < 2.47$, while muons satisfy $p_T > 5$ GeV and $\eta < 2.7$. Jets are reconstructed with the anti- k_T algorithm and considered when $p_T > 30$ GeV and $\eta < 4.5$. A lepton quadruplet consisting of two pairs of same flavor opposite-charge leptons is required, with the p_T cuts in the quadruplet of 20 GeV, 15 GeV and 10 GeV for the leading leptons. Quadruplets with same flavor opposite-charge lepton's invariant mass below 5 GeV are discarded. The opposite sign same flavor lepton pair closest to the Z -boson mass is referred as the leading dilepton, with invariant mass, m_{12} , required to be between 50 GeV and 106 GeV. The sub-leading dilepton invariant mass, m_{34} , is required to be in the range 12 GeV to 115 GeV. As a validation of our recast procedure, we correctly reproduce the ATLAS expected signal and background events in Fig. 3 (right) of Ref. [251].

The number of events obtained in the each bin is proportional to the corresponding squared matrix element, which is a quadratic function of the PO. Therefore, we introduce matrices, X_{ij} , for every bin to express the number of events as

$$\frac{N}{N^{SM}} = \sum_{j \geq i} X_{ij} \epsilon_i \epsilon_j, \quad (3.1)$$

where N^{SM} is the number of expected events in the SM and where we introduced the vector $\epsilon \equiv (\epsilon_{Ze_L}, \epsilon_{Ze_R}, \epsilon_{Z\mu_L}, \epsilon_{Z\mu_R})^T$. The latter only contains contact terms because we focus on the benchmark scenarios 4, 5 and 6. The extension to a more general case with more PO is straightforward. In order to determine the X matrices, we need to run the simulation for different values of the Higgs PO as the input parameters. This needs to

⁷The channels $2e2\mu$ and $2\mu2e$ are split according to which lepton pair comes from the onshell Z boson, or more precisely, which one has an invariant mass closest to the Z mass.

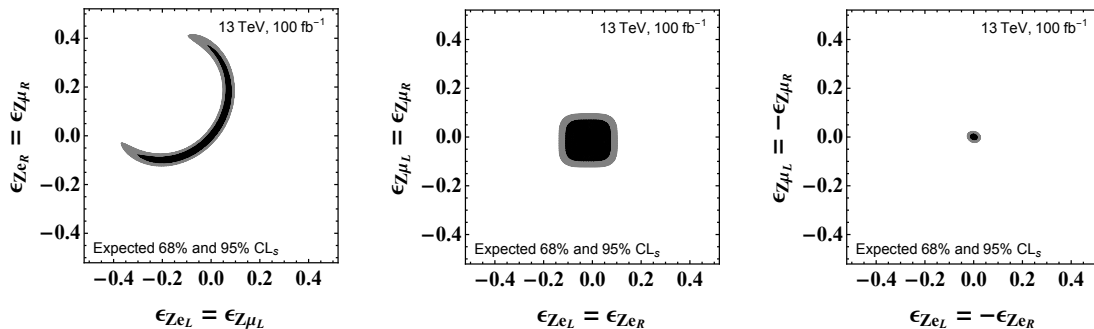


Figure 1: Projection of the ATLAS analysis [251] to 100 fb^{-1} ($\sqrt{s} = 13 \text{ TeV}$) for scenarios 4, 5, and 6. Black (gray) regions represent 68% (95%) expected CL intervals.

be done ten times at least, since this is the number of independent parameters of a real, symmetric, 4×4 matrix. Furthermore, we divided the points in the $m_{12} - m_{34}$ plane into five bins, as was done in the ATLAS search [251]. After we calculate the X matrices, we can compare “exotic PO” with the “SM Higgs” by constructing the appropriate likelihood function. Although we will work with a fix luminosity of 100 fb^{-1} , let us note that the X_{ij} coefficients do not scale with the luminosity.

The events are categorised by four decay channels, $h \rightarrow 4e$, $h \rightarrow 4\mu$, $h \rightarrow 2e2\mu$ and $h \rightarrow 2\mu 2e$, each further separated into five bins. The likelihood function is given by the product of Poisson probabilities for all bins and all categories:

$$L(\epsilon) = \prod_n \frac{\exp(-\mu_n)(\mu_n)^{N_n^{\text{exp}}}}{N_n^{\text{exp}}!}, \quad (3.2)$$

where the index n refers to a specific bin and category, and

$$\mu_n = N_n^{\text{bkg,SM}} + N_n^{\text{sig,SM}} \epsilon^T X_n \epsilon. \quad (3.3)$$

The quantities $N_n^{\text{bkg,SM}}$ and $N_n^{\text{sig,SM}}$ represent the number of events per category and bin for background and signal events within the SM, respectively. N_n^{exp} is number of events per bin and category obtained in the experiment. For the purpose of this study, we neglect systematic uncertainties.

Using the likelihood function in Eq. (3.2), we construct a profile likelihood ratio statistic, $\lambda(\epsilon) \equiv -2 \log(L(\epsilon)/L(\hat{\epsilon}))$, to obtain 68% and 95% expected confidence levels for 100 fb^{-1} assuming the SM. Our results for scenarios 4, 5 and 6, are shown in Fig. 1 from left to right, respectively. In the following, we will try to break down these limits and understand the impact of the overall normalization, lepton universality rate ratios and event kinematics separately.

3.3 Projections Using Relative Normalization and Kinematics

In this section, we study again the sensitivity of future LHC data to Higgs pseudo-observables using benchmark scenarios 4–6. However, in contrast to the analysis in the

previous subsection, we will not include the information contained in the overall normalization common to all $H \rightarrow 4\ell$ flavor channels. Moreover, we will not use the binned (m_{12}, m_{34}) invariant mass distribution but will instead use the unbinned one with full event-by-event information. This approach is similar to that followed in, *e.g.*, the CMS search described in Ref. [22].

In addition to presenting results derived using only kinematic information, we will also present results where the discovery significance shown results from both kinematical information and “relative normalization”, *i.e.*, the relative number of $4e$, 4μ , and $2e2\mu$ events. These analyses help us to understand how much of the sensitivity to exotic PO points is due to kinematical information, how much is due to changes in the ratio of various event types, and how much is due to the overall normalization.

To ease the comparison with the results of the previous subsection we work again with a fixed luminosity of 100 fb^{-1} at 13 TeV. Figs. 2–4 show discovery projections for the non-standard contact terms, $\epsilon_{Z\ell}$, present in scenarios 4–6. These projections are obtained using the Matrix Element Method (MEM) [257–265], in essentially the same manner as in Ref. [92]. Let us discuss below the technical details.

Since the likelihood ratio, the essential quantity to calculate in the MEM, reduces to the ratio of squared matrix elements between different coupling hypotheses, our analysis is built on such “discriminants”, specifically

$$\mathcal{D}_{\text{SM}} = \ln \frac{|\mathcal{M}_{\text{H}}(p|m_H = M)|^2}{|\mathcal{M}_{\text{ZZ}}|^2}, \quad (3.4)$$

and

$$\mathcal{D}_{\text{exo}} = \ln \frac{|\mathcal{M}_{\text{H}}(p|m_H = M)|^2}{|\mathcal{M}_{\text{exo}}(p|m_{\text{exo}} = M)|^2}, \quad (3.5)$$

where M is the invariant mass of the four lepton system, $|\mathcal{M}_{\text{H}}(p|m_H = M)|^2$ is the squared matrix element for the four-lepton momenta, p , under the hypothesis that the four leptons are produced by the decay of an SM Higgs with mass M (and only tree level couplings), $|\mathcal{M}_{\text{exo}}(p|m_{\text{exo}} = M)|^2$ is the squared matrix element for the four-lepton momenta, p , under the hypothesis that the four leptons are produced by the decay of a Higgs boson with the exotic (non-SM) PO indicated by the axes of the corresponding plot, and $|\mathcal{M}_{\text{ZZ}}|^2$ is the squared matrix element for the four-lepton momenta p under the hypothesis that the four leptons are produced by (leading order) SM $q\bar{q} \rightarrow 4\ell$ processes.

To perform the analysis, we obtain the 2-D histograms (“templates” or, in statistical language, probability mass functions) in dimensions of \mathcal{D}_{SM} and \mathcal{D}_{exo} using parton-level signal and background events. The SM normalization per template is applied as is appropriate, which is derived using the SM $H \rightarrow 4\ell$ and $q\bar{q} \rightarrow 4\ell$ processes.⁸ The parton-level histograms were obtained using a background sample of 140 thousand SM $q\bar{q} \rightarrow 4\ell$ events and a signal sample which, after selection cuts, consists of 560 thousand SM $H \rightarrow 4\ell$ events. The SM Higgs and $q\bar{q} \rightarrow 4\ell$ events were obtained using MG5_aMC@NLO [256]; to obtain the vast set of samples that model all the $\gtrsim 150$ coupling points, a matrix-element-based

⁸ As mentioned above, we assume throughout this analysis that the exotic coupling point has the same cross section as the SM, *i.e.*, we do not use the overall normalization in our analysis.

event reweighting was applied, thus enabling a creation of exotic signal templates from the SM Higgs events (see, *e.g.*, Ref. [266]). We conservatively account for mass resolution effects, as in Ref. [92], by including all events within a “window” of 6 GeV around the Higgs boson mass, which we take to be 125 GeV. Separate histograms were produced for each four-lepton final-state flavor, *i.e.*, $4e$, 4μ , and $2e2\mu$. This allows us both to consider the effects of interference, which arise between the amplitudes for different lepton pairings in the $4e$ and 4μ final states (see, *e.g.*, Ref. [79]), and to consider the effects of flavor non-universal anomalous PO in scenarios 5 and 6.

As the \mathcal{D}_{exo} discriminant depends on the chosen coupling point, we need to obtain templates for the three final states, each of different lepton flavor, and for the three event types “SM Higgs”, “exotic PO”, and “background”, for all 7×7 coupling points that we evaluate for each scenario: a total of $3 \times 3 \times 7 \times 7 = 441$ templates. For scenario 4, we construct templates at additional coupling points to make the plot of the expanded (“zoomed-in”) region (Fig. 2b). Having the necessary templates, we (a) produce the distributions of log-likelihood⁹ values for 50000 4-lepton pseudo-experiment events per coupling point, with the number of expected signal and background events given by the SM Higgs and background event yields, respectively, from Ref. [92] and (b) produce the distributions of log-likelihood values for 50000 4-lepton pseudo-experiment events per coupling point, with the number of expected $4e$, 4μ , and $2e2\mu$ events in the ratio predicted for the PO point, but with the total number of $4e$, 4μ , and $2e2\mu$ events set to the SM value.

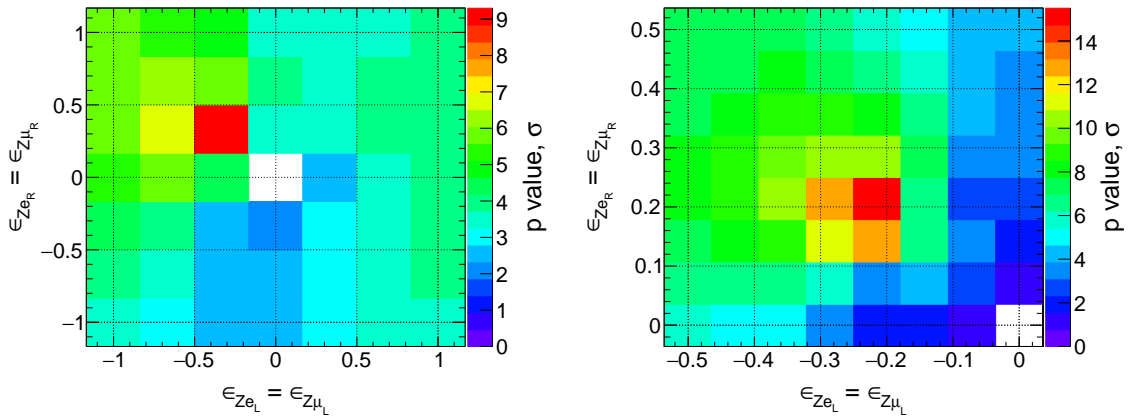
Our immediate aim, for each exotic PO point, and for each of these two analyses, was to determine the mean separation power between the “SM Higgs” and the “exotic PO”. This is done by evaluating the distance in σ_{SM} between the means of corresponding log-likelihood distributions. Such a separation criterion can be loosely thought of as a mean p value and is defined here as the following:

$$\text{p value} = \frac{\langle \text{exotic PO} \rangle - \langle \text{SM Higgs} \rangle}{\sigma_{SM}}, \quad (3.6)$$

where $\langle \text{exotic PO} \rangle$ ($\langle \text{SM Higgs} \rangle$) is the mean of the log-likelihood distribution when the exotic coupling point (the SM Higgs) is the true hypothesis. To be conservative, we included 5% normalization (yield) uncertainties for each of the several different (independent) sources that are typical to experiments, namely (i) the total event count due to theoretical knowledge, (ii) the total event count due to partonic density mismodeling, (iii) the total event count due to generic detector inefficiencies, (iv) the expected number of exotic PO events, (v) the expected number of the SM Higgs events, and (vi) the expected number of background events. These uncertainties effectively smear (widen) the log-likelihood distributions, thus increasing the σ_{SM} denominator in Eq. (3.6). In the limit of large numbers of expected events, as for the 100 fb^{-1} considered in Figs. 2–4, this quantity corresponds to the Z value¹⁰ of the expected discovery significance and can be scaled in a straightforward way to obtain predictions for other luminosities. The results of the analyses of type (a)

⁹ The probability mass functions, which are the above-described templates, times the Poisson probability density functions are used to evaluate the likelihoods.

¹⁰By “ Z value” we mean, essentially, the number of “ σ ” describing the statistical significance of a result.



(a) A full-range version

(b) A zoomed-in version

Figure 2: These figures show the expected discovery sensitivity in sigma (with 100 fb^{-1} at 13 TeV) for the hypothesis described in scenario 4, where, in addition to the SM coupling, the PO $\epsilon_{Z\ell_L} = \epsilon_{Ze_L} = \epsilon_{Z\mu_L}$ and $\epsilon_{Z\ell_R} = \epsilon_{Ze_R} = \epsilon_{Z\mu_R}$ take on non-zero values. These quantities are specified on the x and y axes; the SM is reflected by the point (0,0). The figure on the right, Fig. 2b, gives a zoomed-in version of the second quadrant of Fig. 2a. The results are obtained using only the kinematical distribution of the events via the procedure specified in Subsection 3.3. Including information about the relative normalization of the various $H \rightarrow 4\ell$ channels has a relatively small impact here.

are shown in the left panels of Figs. 2-4 and depend on only kinematical properties of the events, while the results of analyses of type (b) are shown in the right panels Figs. 2-4 and depend on both kinematical information and the relative number of events in each flavor channel that one would obtain from various PO points.

3.4 Discussion of Results

The comparison of the results obtained using rates and kinematical information with the results obtained using only kinematical information in our benchmark scenarios shows that the overall normalization of all $H \rightarrow 4\ell$ channels provides by far the main sensitivity to the contact terms. However it is important to note that kinematic distributions and ratios of rates do provide additional information that allows one to exclude regions of parameter space with reasonably large nonstandard terms that do not affect the overall normalization much due to accidental cancellations. This is why we do not find a perfect ring-shape region in Fig. 1 for scenario 4 (the impact on scenario 5 and 6 is actually much larger).

As mentioned before, the large sensitivity of the overall normalization to contact terms comes however at a price: the implicit assumption that other BSM effects affecting Higgs production (or the SM-like pseudo-observable κ_{ZZ}) are absent. On the other hand, the relative normalization or the kinematic distribution of the events do not require this assumption. Our results in Figs. 2-4 show that these normalized observables provide a significant sensitivity to the contact terms in a large part of the parameter space for all

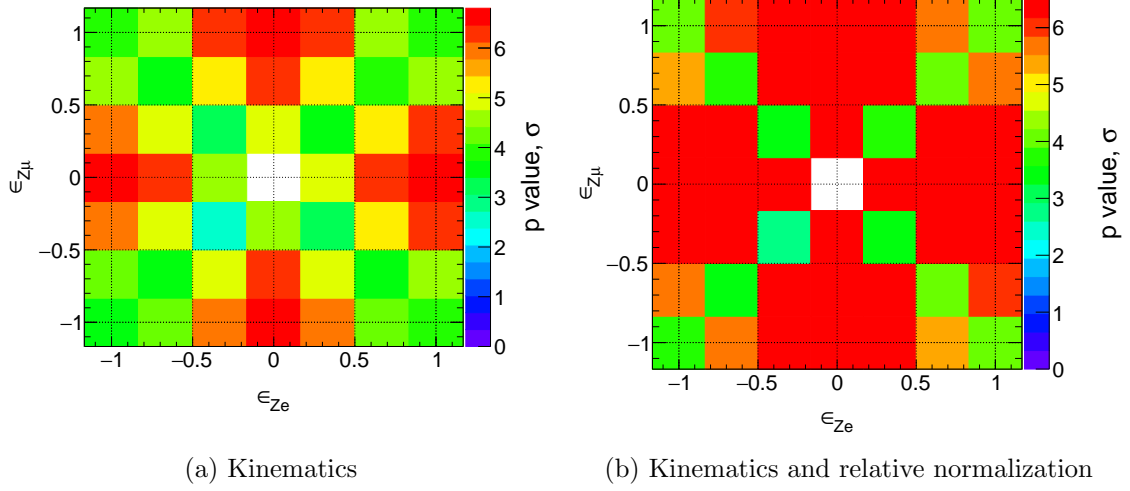


Figure 3: The left figure is the same as Fig. 2a, except that in this case we are considering the benchmark scenario 5, in which the anomalous PO are $\epsilon_{Ze} = \epsilon_{Ze_L} = \epsilon_{Ze_R}$ and $\epsilon_{Z\mu} = \epsilon_{Z\mu_L} = \epsilon_{Z\mu_R}$. The right figure is obtained also for scenario 5 but including not only kinematics but also information about the relative normalization of the various $H \rightarrow 4\ell$ flavor channels.

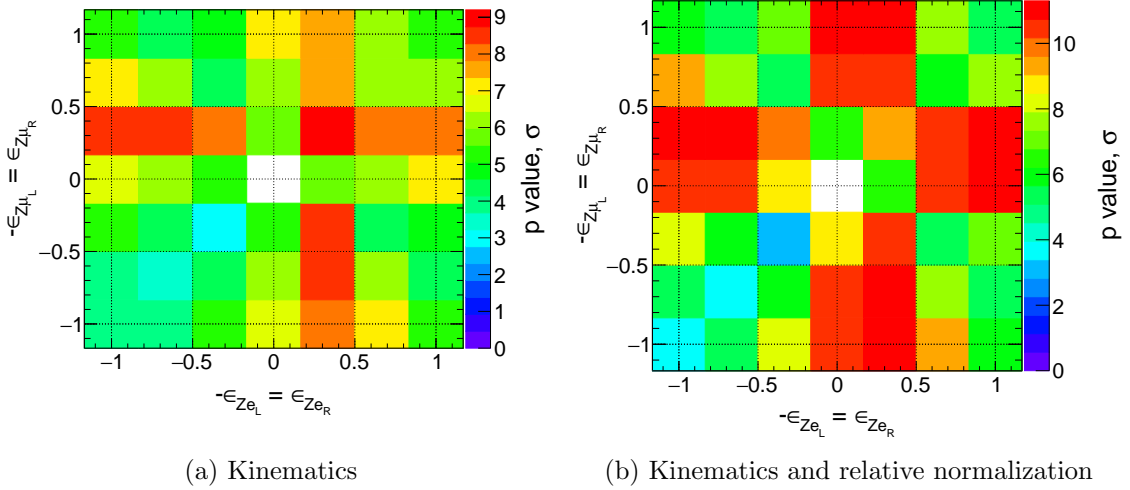


Figure 4: This figure is the same as Fig. 2a, except that here we are considering scenario 6, in which the anomalous PO are $\epsilon_{Ze_R} = -\epsilon_{Ze_L}$ and $\epsilon_{Z\mu_R} = -\epsilon_{Z\mu_L}$.

of the tested scenarios with 100 fb^{-1} , indicating that these points are discoverable (or excludable) with the current LHC Run 2 data set.

Particularly strong sensitivity in distinguishing SM Higgs couplings from contact terms PO points occurs when the contributions to the amplitude from the contact terms interfere destructively with the SM Higgs amplitude [92], as can be seen in Fig. 1 (left) or Fig. 2 (left) for $\epsilon_{Z\ell_L} \sim -\epsilon_{Z\ell_R} \sim -0.23$. In scenarios 5 and 6, where lepton-flavor universality violation is present, the comparison of the different $H \rightarrow 4\ell$ channels represents also a very

sensitive probe, as long as couplings to electrons and to muons are not too similar. This latter effect is clearly seen comparing the left and right plots in Figs. 3 and 4.

It is easy to note that the separation power for different coupling points obeys certain symmetries: the swapping of electrons and muons (x with y), in Figs. 3 and 4, does not change the sensitivity for the scenarios 5 and 6, while the flip of any coupling sign, in Fig. 3, gives an approximately symmetric result.

4 Conclusions

Studies of the Higgs boson in the golden, four-lepton, final state remain important in LHC Run 2 and beyond. A useful toolkit for describing the probed Higgs-to-diboson (and, more generally, $HZ\ell\ell$) interactions exists and contains three major methods for parameterizing couplings. These consist in the specification of the most general $H \rightarrow VV$ amplitude compatible with the symmetries assumed (the amplitude approach), the specification of the most general $H \rightarrow VV$ Lagrangian terms consistent with these same symmetries (the Lagrangian approach), and the use of the most general decomposition of the $H \rightarrow 4\ell$ on-shell amplitude in terms of pseudo-observables (the PO approach). The latter can be considered a generalisation of the two other methods (widely adopted so far in the experimental analyses) able, in particular, to cover also more general BSM frameworks.

In this paper we have compared the use of PO to the other approaches. We showed how to translate between the conventions used by the ATLAS and CMS experiments and in selected theoretical papers, and the parameterization of PO for the four-lepton case specified in Ref. [156]. We have also provided projections for the LHC Run 2 sensitivity to departures from SM couplings in this channel, and we have analyzed the role of kinematic distributions and (ratios of) rates in such $H \rightarrow 4\ell$ studies. In doing so, we have both demonstrated the use of PO and illustrated their relationship to other parametrization methods.

Our work here has aimed at expanding the experimental and phenomenological toolbox for studies of the four-lepton final state. This channel, which provides a treasure trove of information about the underlying coupling of the Higgs to bosons (and leptons), will play an justifiably important role in the LHC physics program for years to come. It is our hope that subsequent studies of the golden channel will lead to a deeper understanding of the Higgs and, perhaps, help inaugurate a golden age of BSM discovery in particle physics.

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