Probing Electroweak Precision Physics via boosted Higgs-strahlung at the LHC

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Introduction

Characterizing the properties of the Higgs boson is arguably the most concrete particle physics goal of our time. This is further motivated by the dearth of any signs of physics beyond the Standard Model (BSM) in LHC data so far. One well-motivated course of action in this situation is to probe heavy new physics outside the reach of direct searches via precise indirect measurements. A historic example of constraining high energy physics even beyond the energy coverage of a collider is the LEP experiment, which was able to probe scales up to the few TeV via indirect precision measurements although it ran at a much smaller collision energy.

As the Higgs boson could not be produced before the LHC experiment under controlled conditions, one might naively think that any measurement of interactions involving the Higgs boson is complementary to past measurements. However, an Effective Field Theory (EFT) perspective allows us to correlate measurements at different energy scales only on the basis of SM symmetry and perturbative mode instead of the full QCD accuracy for weakly coupled UV completions. A systematics limited result to the fractional cross-section deviation, \( \delta \sigma / \sigma_{SM} \) is only \( O(30\%) \) at high energies.

The specific process we are interested in here is Higgs-strahlung, \( pp \to Z(t^+t^-)h(b\bar{b}) \). Studying the \( h \to b\bar{b} \) mode instead of the \( h \to \gamma\gamma \) leads to a big enhancement in the rate but the Higgs-strahlung process still remains challenging with an \( O(50\%) \) signal-to-background ratio. Relating such a systematics limited result to the extraction of Higgs couplings can be at odds with the implicit assumption of perturbativity of the EFT expansion. We technically rely on the latter to perform proof-of-principle analyses and eventually full searches at ATLAS and CMS. As \( \delta \sigma / \sigma_{SM} \gtrsim 1 \) signals the breakdown of EFT validity for weakly coupled UV completions, a sensitivity to smaller values of \( \delta \sigma / \sigma_{SM} \) is essential. To gain such precision, we need high luminosities (at least 300 fb\(^{-1}\)) and advanced boosted Higgs tagging techniques which can reduce the ratio of the number of \( Zb\bar{b} \) to the SM \( Zh(b\bar{b}) \) events to an \( O(1) \) number as shown earlier in Refs. [13–15]. This work is, therefore, an example of a study at the “high energy-luminosity” frontier in the spirit of Ref. [16–18].

While adding the channel \( pp \to (H \to b\bar{b})(Z \to \nu\bar{\nu}) \) can further improve the limits on the effective operators we study [13], this channel is subjected to backgrounds and employs observables with larger systematic uncer-
tainties. We therefore leave an inclusion of this channel for future work.

As we will see, the leading high energy contribution to the $pp \to Zh$ process comes from the four contact interactions $hZ_{\mu L,R} \gamma^\mu \phi_{L,R}$ and $hZ_{\mu L,R} \gamma^\mu \phi_{L,R}$ that are present in the dimension-6 extended Lagrangian. Thus, although many more operators contribute to the $pp \to Zh$ process, the high energy limit isolates the four linear combinations of operators that generate the above contact terms. An interesting observation, first made in Ref. [18], is that the same four EFT directions (that the authors call “high energy primaries”) also control $W h$ and $W W / W Z$ production. The reason is that at high energies these four final processes correspond to the production of different components of the Higgs doublet due to the Goldstone Boson Equivalence Theorem [19]. They are therefore related by $SU(2)_L$ symmetry for $m^2 \gg m^2_Z$. Hence, although these four diboson processes may be very different from a collider physics point of view, they are intimately related by gauge symmetry, which stands at the heart of an EFT interpretation. This enables an elegant understanding of the connection of pseudo-observables in $WW$ production (such as TGCs) with those in $Zh$ production. It will also allow us to present our results in a combined way with the projections for $W Z$ production in Ref. [18].

### The high energy $V h$-amplitude in the SMEFT

Let us first study $V h$ production at high energy in the SMEFT where $V = W, Z$. Although we focus on $pp \to Zh$ production in the subsequent sections, here we keep the discussion more general considering also the $W h$ final state. We will see that $V h$ production at hadron colliders at high energies, isolates four independent directions in the full 59 dimensional space of dimension 6 operators. To derive this fact, consider first the vertices in the dimension 6 Lagrangian that contribute to the $f f \to Zh$ process in unitary gauge,

$$
\Delta \mathcal{L}_6 \supset \sum_f \delta g^2_f Z_{\mu L f} \bar{f} \gamma^\mu f + \delta g^2_{ud} (W_{\mu L}^+ \bar{u}_L \gamma^\mu d_L + h.c.)
$$

$$
g^h_{V V} h \left[ W_{\mu L}^+ W_{\mu L}^- + \frac{1}{2m^2_W} Z_{\mu L} Z_{\mu L} \right] + \delta g^h_{Z Z} h \frac{Z_{\mu L} Z_{\mu L}}{2m^2_W}
$$

$$
\sum_f \delta g^h_{Z Z} h \left[ \bar{Z}_{\mu L f} f \gamma^\mu f + \delta g^h_{W W} h \left( W_{\mu L}^+ \bar{u}_L \gamma^\mu d_L + h.c. \right) \right]
$$

$$
+ \kappa_{Z W} h \frac{A_{\mu W} Z_{\mu L} + \kappa_{W W} h \left( W_{\mu L}^+ \bar{u}_L \gamma^\mu d_L + h.c. \right)}{2m^2_W} Z_{\mu L} Z_{\mu L}.
$$

(2)

We are using the dimension 6 Lagrangian presented in Ref. [2] where any corrections to the SM vector propagators, i.e. the terms $V_{\mu L} V^\mu_{\nu L}$ and $V_{\mu L} F_{\mu \nu}$, have been traded in favor of the vertex corrections in the first two lines above and the TGCs. At high energies, i.e., for $s \gg m^2_Z$, we obtain for the amplitude $M(f f \to V_T L h)$,

$$
Z_T h : g^2_f \frac{2}{v} \hat{J}_f^2 \frac{m^2_Z}{s} \left[ 1 + \left( g^2_{Z f} - \kappa_{Z Z} \right) \frac{s}{2m^2_Z} \right],
$$

$$
Z_L h : g^2_f \frac{2}{v} \hat{J}_f^2 \frac{m^2_Z}{s} \left[ 1 + g^2_{Z f} \frac{s}{2m^2_Z} \right],
$$

$$
W_T h : g^2_f \frac{2}{v} \hat{J}_f^2 \frac{m^2_Z}{s} \left[ 1 + g^2_{W f} - \kappa_{W W} \right] \frac{s}{2m^2_W},
$$

$$
W_L h : g^2_f \frac{2}{v} \hat{J}_f^2 \frac{m^2_Z}{s} \left[ 1 + g^2_{W f} - \kappa_{W W} \right] \frac{s}{2m^2_W},
$$

(3)

where $g^2_f = g(T^f - Q_f s^2_{\phi_W})/c_{\phi_W}$, and $g^2_{W} = g/\sqrt{2}$. $J^\mu_f$ is the fermion current $f \gamma^\mu \bar{f}$, the subscript $L$ ($T$) denotes the longitudinal (transverse) polarization of the gauge boson, $q$ denotes its four momentum and $e$ the polarization vector.

We see that only the $g^2_{Z f}$ and $\kappa_{V V}$ couplings lead to an amplitude growing with energy. In the case of the $\kappa_{V V}$ couplings, the energy growth arises because of the extra powers of momenta in the $hV V$ vertex, whereas for the contact interaction, $g^2_{Z f}$, the energy growth is due to the fact that there is no propagator in the diagram involving this vertex. In fact for the latter interaction, the only difference in the amplitude with respect to the SM is the absence of the propagator. Thus, angular distributions are expected to be identical for BSM and SM production. Therefore, the only way to probe this interaction is through the direct energy-dependence of differential cross-sections.

On the other hand, the $\kappa_{V V}$ interactions contribute only to the transverse $V$ amplitude as a consequence of their vertex structure. Hence, they cannot interfere with the dominant longitudinal piece in the SM amplitude. As a result, after summing over all $V$-polarizations, the
leading piece in the high energy cross-section deviation, is controlled only by the couplings $g^{b}_{Zf}$ whereas the $\kappa_{VV}$ contribution is suppressed by an additional $\mathcal{O}(m_{Z}^{2}/\hat{s})$ factor.

At hadron colliders, the $pp \rightarrow Vh$ process at high energies and at leading order are therefore controlled by the five contact interactions: $g^{b}_{Zf}$, with $f = u_{L}, u_{R}, d_{L}$ and $d_{R}$ and $g^{b}_{Wud}$. These five couplings correspond to different linear combinations of Wilson coefficients in any given basis. In Tab. I we show all operators in the “Warsaw” [1] and strongly-interacting light Higgs (SILH) [20] bases that generate these contact terms. As there are only four independent operators contributing to these five interactions in the Warsaw basis, there exists a basis independent constraint at the dimension-6 level,

$$g^{b}_{Wud} = g^{b}_{Z2d_{L}} - g^{b}_{Z2u_{L}}$$

leaving only the four independent $g^{b}_{Z2f}$ couplings.

In Table II, we show the linear combinations of Wilson coefficients contributing to the four $g^{b}_{Z2f}$ couplings in different EFT bases. The first row gives these directions in the Warsaw basis. The second row provides the aforementioned directions in the BSM Primary basis of Ref. [2], where the Wilson coefficients can be written in terms of already constrained pseudo-observables. It is clear in this basis that the directions to be probed by high energy $Vh$ production can be written in terms of the LEP (pseudo)observables. The couplings $\delta g^{Z}_{f}$ defined in Eq. (2) are strongly constrained by $Z$-pole measurements at LEP, whereas the anomalous TGCs, $\delta\kappa_{f}$ and $\delta g^{Z}_{f}$ (in the notation of Ref. [21]), were constrained by $WW$ production during LEP2.

For the physically motivated case where the leading effects of new physics can be parametrized by universal (bosonic) operators, the SILH Lagrangian provides a convenient formulation and we show the above directions in this basis in the third row of Table II. For this case one can again write the directions in terms of only the “oblique”/universal pseudo-observables, viz., the TGCs $\delta\kappa_{f}$ and $\delta g^{Z}_{f}$ and the Peskin-Takeuchi $S$-parameter [22] in the normalization of Ref. [23] (see e.g. Refs. [24, 25]). This is shown in the fourth row of Table II. As we already mentioned, upon using the Goldstone Equivalence Principle, one finds that the same 4 dimensional subspace of operators also controls the longitudinal $VV$ production at high energies. This space is defined in Ref. [18] in terms of the four high energy primaries which are linear combinations of the four $g^{b}_{Vf}$ couplings, as shown in the last row of Table II.

As it is not possible to control the polarization of the initial state partons in a hadron collider, the process can, in reality, only probe two of the above four directions. Taking only the interference term, these directions are

$$g^{Z}_{u} = g^{Z}_{2u_{L}} + \frac{g^{Z}_{u}}{g^{Z}_{u_{L}}} g^{Z}_{2u_{R}},$$

$$g^{Z}_{d} = g^{Z}_{2d_{L}} + \frac{g^{Z}_{d}}{g^{Z}_{d_{L}}} g^{Z}_{2d_{R}},$$

where $g^{Z}_{f}$ is defined below Eq. (3). Also, at a given energy, the interference term for the $pp \rightarrow Zh$ process is sensitive only to a linear combination of the up-type and down-type coupling deviations, i.e., to the direction,

$$g^{Z}_{p} = g^{Z}_{u} + \frac{E_{u_{L}}}{E_{u_{R}}} g^{Z}_{d},$$

where $E_{u,d}$ is the $uu$, $dd$ luminosity at a given partonic centre of mass energy. We find that the luminosity ratio changes very little with energy (between 0.65 and 0.59 if $\sqrt{s}$ is varied between 1 and 2 TeV). Thus, to a good approximation, $pp \rightarrow Zh$ probes the single direction in EFT space given by

$$g^{Z}_{p} = g^{Z}_{2u_{L}} - 0.76 g^{Z}_{2d_{L}} - 0.45 g^{Z}_{2u_{R}} + 0.14 g^{Z}_{2d_{R}},$$

where we have substituted the values for $g^{Z}_{f}$ and evaluated the luminosities at the energy $\hat{s} = (1.5 \text{ TeV})^{2}$.
Tab. II.

\[ g^h_{Zp} = 2\delta g_{Zu_L}^h - 1.52 g_{Zd_L}^h - 0.90 g_{Zu_R}^h + 0.28 g_{Zd_R}^h - 0.14 \delta \kappa - 0.89 \delta g_1^Z \]

\[ g^h_{Zp} = -0.14 (\delta \kappa - \delta \kappa_1) - 0.89 \delta g_1^Z - 1.3 W \] (8)

where the first line applies to the general case and the second line to the universal case.

Note that in the discussion so far we have not considered the \( gg \to Zh \) production channel at hadron colliders [26–34]. While formally a higher order correction, after all the cuts are applied, this subprocess contributes an appreciable 15% of the total SM \( pp \to Zh \) cross-section in our analysis due to the top-threshold inducing boosted final states [31]. We find, however, that introduction of the EFT operators does not lead to a energy growing amplitude with this initial state, and thus this channel has a subdominant contribution to the EFT signal. While we fully include this contribution in our collider analysis, the introduction of this channel does not alter the discussion so far in an important way.

We now turn to the crucial issue of estimating the scale of new physics (and thus the cut-off for our EFT treatment) for a given size of the couplings, \( g^h_{Zf} \). This will also give us an idea of the new physics scenarios that our analysis can probe. As is clear from the operators in Tab. I, the \( g^h_{Zf} \) couplings arise from current-current operators that can be generated, for instance, by integrating out at tree-level a heavy \( SU(2)_L \) triplet (singlet) vector \( W^{a\alpha} (Z') \) that couples to SM fermion currents, \( F a \alpha \gamma_\mu F (f \gamma_\mu f) \) with a coupling \( g f \) and to the Higgs current \( h I_\alpha \partial_\mu H (h I_\alpha \partial_\mu H) \) with a coupling \( g_H \),

\[ g^h_{Zf} \sim \frac{g_H g f v^2}{\Lambda^2} \] (9)

where \( \Lambda \) is the mass of the vector and therefore the matching scale or cut-off of the low energy EFT. The coupling to the SM fermions can be universal if the heavy vector couples to them only via kinetic mixing with the SM gauge bosons. This results in a coupling of the heavy vector to the \( SU(2)_L \) and hypercharge currents given by \( g_W = g/2 \) and \( g_B = g' Y_f, Y_f \) being the SM hypercharge. As we want our results to be applicable to the universal case, we assume the coupling \( g f \) to be given by a combination of \( g_B \) and \( g_W \) to obtain,

\[ \begin{align*}
g^h_{Zu_L, d_L} & \sim \frac{g_B g^2 v^2}{2\Lambda^2}, \\
g^h_{Zu_R, d_R} & \sim \frac{g_B g^2 Y_{u_R, d_R} v^2}{\Lambda^2} \end{align*} \] (10)

and then further assume a weakly coupled scenario with \( g_H = 1 \) (note that this is a bit larger than the corresponding value \( g_H = g/(2c_{\theta_W}) \) for the SM \( hZZ \) coupling). In the above equation, we have ignored the smaller contributions from \( g_B \) to the left-handed couplings. For any set of couplings \( \{ g^h_{Zu_L}, g^h_{Zd_L}, g^h_{Zu_R}, g^h_{Zd_R} \} \), we evaluate the cut-off using Eq. (10) with \( g_H = 1 \) and take the smallest of the four values. It is clear that for strongly coupled scenarios with larger values of \( g_H \), the cut-off assumed in our analysis is smaller than necessary and thus our projected bounds will be conservative.

**Analysis**

In order to probe the reach of the high luminosity runs of the LHC in constraining the EFT directions in Tab. II, we optimize a hadron-level analysis to obtain maximum sensitivity to the BSM signal, which is well-pronounced in the high energy bins. To achieve this, we consider the \( Z(\ell^+ \ell^-)h \) production from a pair of quarks as well as from a pair of gluons. As far as the decay of the Higgs boson is concerned, we find that at an integrated luminosity of 300 fb\(^{-1}\), the diphoton mode yields less than 5 events at high energies (\( p_T > 150 \) GeV) and is thus not sensitive to the effects we want to probe. We thus scrutinize the \( h(bb)Z(\ell^+ \ell^-) \) final state where the dominant backgrounds are composed of \( Zb \) and the irreducible SM production of \( Zh \). Reducible contributions arise from \( Z+ \) jets production (c-quarks included but not explicitly tagged), where the light jets can be misidentified as b-jets, and the fully leptonic decay for \( t\bar{t} \). Instead of performing a standard resolved analysis, where one would demand two separate b-tagged jets, we demand a fat jet with a cone radius of \( R = 1.2 \). We employ the so-called BDRS approach [13] with minor modifications to maximize sensitivity. In a nutshell, this technique helps in discriminating boosted electroweak-scale resonances from large QCD backgrounds.

We will see that using this approach will allow us to reduce the ratio of \( Zbb \) to SM \( Zh \) events from about 40 to an \( O(1) \) number with about 40 SM events still surviving at 300 fb\(^{-1}\). This shows that the kind of analysis performed here would not be possible at integrated luminosities smaller than 300 fb\(^{-1}\).

The BDRS approach recombines jets using the Cambridge-Aachen (CA) algorithm [35, 36] with a significantly large cone radius to contain all the decay products of the resonance. One then works backwards through the jet clustering and stops when a significant mass drop, \( m_{j_1} < \mu m_j \) with \( \mu = 0.66, (m_j \) being the mass of the fatjet) occurs for a not too asymmetric splitting ,

\[ \min\left(\frac{p_{T,j_1}^2}{m_j^2}, \frac{p_{T,j_2}^2}{m_j^2}\right) \Delta R_{j_1,j_2} > y_{\text{cut}}, \]

with \( y_{\text{cut}} = 0.09 \). If this condition is not met, the softer subjet, \( j_2 \) is removed and the subjets of \( j_1 \) are tested for the aforementioned criteria to be satisfied in an iterative process. The algorithm stops as soon as one obtains two subjets, \( j_1 \) and \( j_2 \) abiding by the mass drop condition.
To improve the resonance reconstruction, the technique considers a further step called filtering. In this step, the constituents of \( j_1 \) and \( j_2 \) are again recombined using the CA algorithm with a cone radius \( R_{\text{filt}} = \min(0.3,R_{b\bar{b}}/2) \). Only the hardest three filtered subjets are retained to reconstruct the resonance. In the original work of Ref. [13], this resonance is the SM-like Higgs and thus the two hardest filtered subjets are \( b \)-tagged. In our work, we find that the filtered cone radius \( R_{\text{filt}} = \max(0.2,R_{b\bar{b}}/2) \) works better in removing the backgrounds.\(^1\) The filtering step greatly reduces the active area of the initial fatjet.

We use the FeynRules [37] and UFO [38] toolkits to implement the signal contributions (we will comment on the effect of including squared dimension 6 interactions as compared to interference-only terms below). Both signal and background processes are generated including the full decay chain with \texttt{MG5_aMC@NLO} [39], at leading order. For the gluon initiated part of the SM and BSM \( Zh \) production, we employ the \texttt{FeynArts/FormCalc/LoopTools} [40, 41] framework and the decays are performed using \texttt{MadSpin} [42, 43]. We shower and hadronize the samples using \texttt{Pythia 8} [44, 45] and perform a simplified detector analysis.

Because our ultimate goal is to look for new physics effects in high energy bins, we generate the \( Zh \), \( Zbb \) and \( t\bar{t} \) samples with the following cuts: \( p_T,(b,\bar{b}) > 15 \text{ GeV} \), \( \Delta R_{b\bar{b}} > 0.2 \), \( \Delta R_{\ell\ell} > 0.15 \). \( 70 \text{ GeV} < m_{\ell\ell} < 110 \text{ GeV} \), \( 95 \text{ GeV} < m_{b\bar{b}} < 155 \text{ GeV} \) and \( p_T,\ell \) > 150 GeV. The former two processes are generated upon merging with an additional matrix element (ME) parton upon using the MLM scheme [46]. For the \( Z+jets \) process, we generate the samples without the invariant mass cuts on the jets; we further merge the sample up to three ME partons.

To account for higher order QCD corrections for the \( q\bar{q} \)-initiated \( Zh \) process, we apply a bin-by-bin (in \( M_{Zh} \), the invariant mass of the filtered double \( b \)-tagged fat jet and the reconstructed \( Z \)-boson) \( K \)-factor reweighting to the NLO-accurate distribution both for the SM background and the EFT signal using Ref. [47]. For the \( gg \) initiated \( Zh \) process, we consider a conservative NLO \( K \)-factor of 2 [32]. For the \( Zbb \) and \( Z+jets \) processes, flat \( K \)-factors of 1.4 (computed within \texttt{MG5_aMC@NLO}) and 0.91 [48] are applied.

We first test the power of a cut-based analysis. In doing so, we construct the fatjets with a cone radius of \( R = 1.2 \), having \( p_T > 80 \text{ GeV} \) and rapidity, \( |y| < 2.5 \) using \texttt{FastJet} [49]. We isolate the leptons (\( e,\mu \)) by demanding that the total hadronic activity around a cone radius of \( R = 0.3 \) must be less than 10% of its \( p_T \) and the leptons are required to have \( p_T > 20 \text{ GeV} \) and \( |y| < 2.5 \). All non-isolated objects are considered while constructing the fatjets. In selecting our events, we consider only those with exactly two isolated leptons having opposite charge and same flavour (OSSF). Moreover, we demand the invariant mass of the pair of leptons to lie in the range \([80 \text{ GeV}, 100 \text{ GeV}] \) in order to reconstruct the \( Z \)-peak. The reconstructed \( Z \) is required to be boosted with \( p_T > 160 \text{ GeV} \) and the separation between the two isolated leptons is required to be \( \Delta R > 0.2 \). In reconstructing the Higgs boson, we demand that each event has at least one fatjet containing no less than two \( B \)-meson tracks with \( p_T > 15 \text{ GeV} \). The minimum transverse momentum of the fatjet is required to be \( p_T > 110 \text{ GeV} \). After satisfying the mass drop and filtering criteria, we require exactly two subjets after the former step and at least two subjets after filtering. We proceed with \( b \)-tagging the two hardest subjets. We choose a \( b \)-tagging efficiency of 70% and a misidentification rate for light jets of 2%. After the filtering and \( b \)-tagging steps, we require events with exactly two \( b \)-tagged subjets, which are well-separated from the isolated leptons; \( \Delta R(b_i,\ell_j) > 0.4 \) for both leptons \( \ell_{1,2} \) and \( b \)-tagged subjets \( b_i \). We reconstruct the Higgs boson by requiring its invariant mass to lie in the range \([115 \text{ GeV}, 135 \text{ GeV}] \).

In order to further reduce the backgrounds, we demand both the reconstructed \( Z \) and the Higgs bosons to have \( p_T > 200 \text{ GeV} \). The \( t\bar{t} \) background can be removed almost entirely by requiring \( \sum E_T < 30 \text{ GeV} \). The cut-flow affecting the most dominant background \( Zbb \) and the SM

| Cuts                                                                 | \( Zbb \) | \( Zh \) (SM) |
|---------------------------------------------------------------------|----------|--------------|
| At least 1 fat jet with 2 \( B \)-mesons with \( p_T > 15 \text{ GeV} \) | 0.16     | 0.41         |
| 2 OSSF isolated leptons                                           | 0.41     | 0.50         |
| \( 80 \text{ GeV} < M_{\ell\ell} < 100 \text{ GeV} \), \( p_T,\ell > 160 \text{ GeV} \), \( \Delta R_{\ell\ell} > 0.2 \) | 0.85     | 0.89         |
| At least 1 fat jet with 2 \( B \)-meson tracks with \( p_T > 110 \text{ GeV} \) | 0.95     | 0.98         |
| 2 Mass drop subjets and \( \geq 2 \) filtered subjets             | 0.87     | 0.93         |
| 2 \( b \)-tagged subjets                                           | 0.38     | 0.41         |
| \( 115 \text{ GeV} < m_b < 135 \text{ GeV} \)                     | 0.25     | 0.49         |
| \( \Delta R(b_i,\ell_j) > 0.4 \), \( E_T < 30 \text{ GeV} \), \( |y| < 2.5 \), \( p_T,h,Z > 200 \text{ GeV} \) | 0.48     | 0.70         |

\(^1\) The criteria \( R_{\text{filt}} = \max(0.2,R_{b\bar{b}}/2) \) followed by \( R_{\text{filt}} = \min(0.3,R_{b\bar{b}}/2) \) hardly changes the results.
Zh channel, is summarized in Table III.

Before focussing on the very high-energy effects by imposing cuts on \(M_{Z\ell}\), we find that the ratio of cross-section between SM Zh and Zbb is \(\sim 0.32\). A multivariate implementation at this level strengthens this ratio further. In order to be quantitative, we impose looser cuts on the aforementioned variables \(70 \text{ GeV} < m_{\ell\ell} < 110 \text{ GeV}, p_T^{\ell\ell} > 160 \text{ GeV}, \Delta R_{\ell\ell} > 0.2, p_T^{\text{fatjet}} > 60 \text{ GeV}, 95 \text{ GeV} < m_h < 155 \text{ GeV}, \Delta R_{b_{\text{t}}}, \ell \geq 0.4 \) and \(E_T < 30 \text{ GeV}\). Because \(Z+\)jets and \(\ell\ell\) are much less significant than \(Zbb\), we train the boosted decision trees only with the SM \(q\bar{q}\)-initiated Zh and Zbb samples using the following variables: \(p_T\) of the two isolated leptons, \(\Delta R\) between pairs of \(b\)-subjets and isolated leptons, between the two isolated leptons and between the hardest two \(b\)-subjets in the Higgs fatjet, the reconstructed \(Z\)-boson mass and its \(p_T\), \(\Delta \Phi\) separation between the fatjet and the reconstructed \(Z\)-boson, \(E_T\), mass of the reconstructed Higgs jet and its \(p_T\) of the two \(b\)-tagged filtered subjets, the ratio of their \(p_T\) and the rapidity of the Higgs jet. We ensure that we do not have variables which are \(\sim 100\%\) correlated but we retain all other variables. Because our final distribution of interest is the invariant mass of the \(Zh\)-system, we do not consider it as an input variable. We use the TMA [50] framework to train our samples and always ensure that the Kolmogorov-Smirnov statistic is at least of the order \(\sim 0.1\) in order to avoid overtraining of the samples [51]. We find that the aforementioned ratio increases to \(\sim 0.54\) upon using the boosted decision tree algorithm showing that a further optimisation of the cut-based analysis was necessary. Finally, we test all our samples with the training obtained from the SM \(q\bar{q}\) initiated Zh and the Zbb samples.

To distinguish between the EFT signal and the irreducible SM Zh(bb) background we utilise the growth of the EFT cross-section at high energies. The effects are readily seen in the \(M_{Z\ell}\) distribution, our observable of interest. In Fig. 1 we show the differential distribution with respect to this variable for the EFT signal as well as the different backgrounds for an integrated luminosity of 300 fb\(^{-1}\). For the EFT signal we take a point that can be excluded in our analysis but is well within the LEP allowed region. We see that the EFT cross-section keeps growing with energy, but much of this growth is unphysical at energies above the cut-off, \(i.e., M_{Z\ell} > \Lambda\), where \(\Lambda\) is the cut-off evaluated as described below Eq. (10) and shown by a vertical line in Fig. 1. For \(M_{Z\ell} < \Lambda\), the EFT deviations are never larger than an \(\mathcal{O}(1)\) factor with respect to the SM background as expected on general grounds. Note, however, that even for \(M_{Z\ell} < \Lambda\) the fractional deviation can be a few tens of a percent even though the underlying anomalous couplings, \(g^{h}_{2\ell}\), are per-mille to percent level.

To make full use of the shape deviation of the EFT signal with respect to the background, we perform a binned log likelihood analysis assuming a 5% systematic error. The likelihood function is taken to be the product of Poisson distribution functions for each bin with the mean given by the number of events expected for a given BSM point. To account for the 5% systematic error we smear the mean with a Gaussian distribution. To obtain the projection for the 95\% CL exclusion curve we assume that the observed number of events agrees with the SM.
Discussion

Considering only the SM-BSM interference term, we find the per-mille level bounds,

\begin{align}
  g^h_{Zp} & \in [-0.003, 0.003] \quad (300 \text{ fb}^{-1}) \\
  g^h_{Zp} & \in [-0.001, 0.001] \quad (3000 \text{ fb}^{-1}).
\end{align}  

Using Eq. (10) the above bounds can be translated to a lower bound on the scale of new physics given by 2.4 TeV (4.4 TeV) at 300 fb\(^{-1}\) (3000 fb\(^{-1}\)). One can now compare the above projections with existing LEP bounds by turning on the LEP observables contributing to \(g^h_{Zp}\) in Eq. (8) one by one. This is equivalent to assuming that there are no large cancellations in Eq. (8) so that each individual term is bounded by Eq. (11). The results are shown in Tab. IV. We see that our projections are much stronger than the LEP bounds for the TGCs \(\delta g^2_Z\) and \(\delta \kappa_Z\), and comparable in the case of the \(Z\)-pole observables \(\delta g^Z_1\), that parametrize the deviations of the \(Z\) coupling to quarks.

For the universal case, the EFT directions presented in Table II can be visualized in the \(\delta \kappa_Z - \hat{S}\) vs. \(\delta g^Z_1\) plane as shown in Fig. 2 for the interesting class of models where \(W = Y = 0\) [18]. The flat direction related to the \(pp \rightarrow Zh\) interference term, i.e., \(g^h_{Zp} = 0\), Eq. (7), is shown by the dashed blue line, where the direction \(g^h_{Zp}\) is now given by the second line of Eq. (8). The grey shaded area shows the allowed region after the LEP II bounds [53] from the \(e^+e^- \rightarrow W^+W^-\) process are imposed. The results of this work are shown in blue (light (dark) blue for results at 300 (3000) fb\(^{-1}\)). To understand the shape of the blue bands, note that along the dashed line, the SM-BSM interference term vanishes. If the interference was the only dominant effect, the projected allowed region would be a band along this direction. The BSM squared term thus plays a role in determining the shape of the blue region. To the left of the dashed blue line, the squared and the interference terms have the same sign while there is a partial cancellation between these two terms on the right hand side of the dashed line. This results in the curvature of the blue band with stronger bounds to the left of the dashed line and weaker bounds to its right.

We see that, as we move further from the origin, the effect of the squared term becomes more pronounced. This is expected, as along the dashed line, the interference term is accidentally zero, even for energies below the cut-off, and thus, the parametrically sub-dominant squared term is larger. To achieve a partial cancellation between these two terms one needs to deviate more and more from the dashed line. While EFT validity has been carefully imposed to derive our bounds, the fact that the interference term vanishes along the flat direction and the squared term becomes important, does imply that for weakly coupled UV completions our bounds are susceptible to \(O(1)\) dimension 8 deformations in this direction. In the orthogonal direction shown by the dotted line, on the other hand, our projections are more robust and not sensitive to such effects.

As we have emphasized already, \(VV\) production constrains the same set of operators as the \(Vh\) production. In Fig. 2, we also show the projected bound from the \(WZ\) process at 300 fb\(^{-1}\) obtained in Ref. [18]. When both these bounds are combined, only the purple region remains. At 3000 fb\(^{-1}\), this region shrinks further to the green region shown in Fig. 2. Thus, we see a drastic reduction in the allowed LEP region is possible by investigating \(pp \rightarrow Zh\) at high energies.

Conclusions

As hints for new physics beyond the SM remain elusive with the LHC entering a new energy territory, model-independent approaches based on the assumption of no additional light propagating degrees of freedom are gaining ground. The power of effective field theory is that theoretical correlations between independent measurements can be exploited to formulate tight constraints on the presence of new physics, solely based on the SM symmetries and matter content.
The high precision measurements performed during the LEP era are therefore the driving forces behind combined constraints early in the LHC program. To enter new territory, the LHC has to push beyond the LEP sensitivity for interactions that relate the phenomenology at both collider experiments. The Higgs boson, as arguably the most significant TeV scale degree of freedom, can be placed at the core of such a program, that will naturally involve LHC measurements at high luminosity.

In this paper, we focussed on the impact of associated Higgs production that provides complementary information to the diboson production modes observed at LEP2, which determine the precision of the associated coupling constraints. Using a dedicated investigation of expected signal and backgrounds, we find that the LHC will ultimately be able to improve the sensitivity expected from LEP measurements. Our results are summarised in Eq. (11), Tab. IV and Fig. 2. Higgs-strahlung is also complementary to diboson production at LHC investigated in Ref. [18]. Combining Higgs-strahlung measurements with diboson results in the high energy limit will allow us to drastically improve the sensitivity to the underlying new physics parameters in an unparalleled way.

Both high energies and luminosities are crucial for a study like ours. Potentially even higher new physics scales can thus be probed at the High Energy LHC or other future colliders.

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