Limitations and Opportunities of Off-Shell Coupling Measurements

Christoph Englert\textsuperscript{1,}\textsuperscript{*} and Michael Spannowsky\textsuperscript{2,}\textsuperscript{†}
\textsuperscript{1}SUPA, School of Physics and Astronomy, University of Glasgow, Glasgow G12 8QQ, United Kingdom
\textsuperscript{2}Institute for Particle Physics Phenomenology, Department of Physics, Durham University, Durham DH1 3LE, United Kingdom

Indirect constraints on the total Higgs width $\Gamma_h$ from correlating Higgs signal strengths with cross section measurements in the off-shell region for $p(g)p(g) \rightarrow 4\ell$ production have received considerable attention recently, and the CMS collaboration have published a first measurement. We revisit this analysis from a new physics and unitarity constraints perspective and conclude that limits on $\Gamma_h$ obtained in this fashion are not reliable unless we make model-specific assumptions, which cannot be justified at the current stage of the LHC programme. Relaxing the $\Gamma_h$ interpretation, we discuss the merits of high invariant mass cross section measurements in the context of Higgs CP analyses, higher dimensional operator testing, and resolved new physics in the light of electroweak precision constraints beyond effective theory limitations. Furthermore, we show that a rather model-independent LHC constraint can be obtained from adapting the $gg \rightarrow 4\ell$ analysis to the weak boson fusion channels at lower statistical yield.

I. INTRODUCTION

After the 2012 discovery\textsuperscript{[1, 2]}, the ATLAS and CMS collaborations have scrutinized the SM interpretation of the Higgs candidate within the boundaries of the currently available data. A strong resemblance of the particle’s properties with the SM Higgs expectation has emerged: it is likely to be a $CP$ even scalar boson and its “signal strengths”

$$ \mu_{i,j} = \frac{\sigma_{h,i} \times BR_j}{[\sigma_{h,i} \times BR_j]_{SM}} \sim \frac{\Gamma_i \Gamma_j}{\Gamma_{i,SM} \Gamma_{j,SM}} \Gamma_h $$

(1)

are in good agreement with the SM Higgs boson. $i,j$ in Eq. (1) refer to the different Higgs production and decay modes that have been observed so far. For fully inclusive measurements they can be related to the partial decay widths $\{\Gamma_i\}$. “Higgsistence” has mainly been established from gluon fusion, the largest Higgs production mechanism in the SM.

The apparent agreement of the measured quantities of Eq. (1) with the SM predictions highlights the question whether the discovered particle is indeed the Higgs boson as predicted by the SM.

On the one hand, unitarity largely constrains the bare couplings of massive fermions and gauge bosons to $CP$ even Higgs boson(s) in the SM. If the absolute values of the Higgs candidate’s couplings are close to the SM predictions, there will be little room left for resonant physics beyond the SM in e.g. the weak boson fusion (WBF) channels, which is a direct probe of longitudinal gauge boson scattering.

On the other hand, absolute values of couplings are difficult to infer at hadron colliders since signal strength measurements involve non-linear relations among the couplings and $\sigma \times BR$ phenomenology leaves the total Higgs width as a flat direction in coupling fits. This is usually overcome by making assumptions about the total Higgs width in these fits\textsuperscript{[2]}, or, alternatively, about the maximum coupling value of the Higgs candidate to gauge bosons\textsuperscript{[4]} which is determined by the Higgs’ gauge representation. The biases that are introduced in either of these approaches are far from being well-motivated at the current stage.

Assuming $\Gamma_h \sim \Gamma_{h,SM} \sim 4$ MeV skews coupling fits towards a parameter region that is oblivious of the Higgs bosons’ potential interplay with dark matter phenomenology\textsuperscript{[5]} and other phenomena that can be introduced via well-motivated portal-type interactions\textsuperscript{[6, 7]}.

Assumptions about the Higgs $SU(2)_L$ representation are usually limited to the 2’s due to the (accidental) custodial isospin symmetry that preserves $T \simeq 1$ in (multi-) Higgs doublet models. However, it is known that both current signal strength measurements and electroweak precision constraints can be accounted for in models with non-doublet Higgs fields\textsuperscript{[8]} and the complementing searches for Higgs exotics\textsuperscript{[9]} necessary to rule out such an option are not available yet.

Obviously, a model-independent constraint on $\Gamma_h$ (or $BR$(invisible) if a particular model leaves production modes unaltered) has a huge impact on BSM physics.\textsuperscript{[1]}

Hence, it is not surprising that the recent proposal by Caola and Melnikov\textsuperscript{[11]} that interprets off-shell cross section measurements of $pp \rightarrow 4\ell$\textsuperscript{[12]} as a probe of $\Gamma_h$ has received considerable attention\textsuperscript{[13–15]}.\textsuperscript{[2]}

\textsuperscript{1}See e.g. Ref.\textsuperscript{[10]} for a discussion of the invisible branching ratio measurements, e.g.\textsuperscript{[10]}, in relation with dark matter phenomenology.

\textsuperscript{2}Similar strategies\textsuperscript{[10]} have been proposed for $h \rightarrow \gamma\gamma$\textsuperscript{[16]}. 

\textsuperscript{*}Electronic address: christoph.engler@glasgow.ac.uk
\textsuperscript{†}Electronic address: michael.spannowsky@durham.ac.uk
CMS have presented first results [18] using this strategy, claiming \( \Gamma_h < 4.2 \times \Gamma_h^{\text{SM}} \) at 95% confidence level by injecting a global Higgs signal strength \( \mu \simeq 1 \). The strategy is sketched in Fig. 1 and we give a quick outline to make this work self-contained (for additional details see [11, 14, 18]):

As long as the narrow width approximation is applicable, the cross section for the process \( p(g)p(g) \rightarrow h \rightarrow ZZ \rightarrow 4\ell \) in the the Higgs on-shell region scales as

\[
\sigma_{h, g} \times \text{BR}(h \rightarrow ZZ \rightarrow 4\ell) \sim \frac{g^2_{gg h} g^2_{ZZ Z} \sqrt{s}}{\Gamma_h},
\]

where we denote the relevant couplings by \( g_X \). The dominant Feynman diagram in this phase space region is the triangle of Fig. 2 the continuum contribution from \( gg \rightarrow ZZ^* \) is highly suppressed and interference is negligible [12].

Since the Higgs width is anticipated to be a small parameter compared to the Higgs mass \( \Gamma_h/m_h \sim 10^{-4} \), we can expand the Higgs Breit-Wigner propagator \( D(s) = i/(s - m_h^2 + i\Gamma_h m_h) \) away from the peak region \( s \gg m_h^2 \)

\[
|D|^2 = \frac{1}{s^2} \left( 1 + \frac{m_h^4 \Gamma_h^2}{s^4 m_h^4} \right) + \mathcal{O}\left( \frac{s^4}{s^4} \right)
\]

which shows that the Higgs width parameter rapidly decouples from the scattering process for Higgs off-shell production. Therefore, the contribution from the triangle diagrams in Fig. 2 (neglecting interference for the moment) scales as

\[
d\sigma_h \sim \frac{g^2_{gg h}(\sqrt{s}) g^2_{hZZ}(\sqrt{s})}{s} \text{dLIPSxpdfs}.
\]

Now, if there is a direct correspondence between \( q_i(m_h) \) and \( q_i(\sqrt{s}) \), measuring the signal strength \( \mu \) in the off-shell and on-shell regions simultaneously allows us to set a limit on the width of the Higgs boson \( \Gamma_h \). More explicitly, for \( \Gamma_h > \Gamma_h^{\text{SM}} \), we need to have \( g^2_{gg h} g^2_{hZZ} > (g^2_{gg h} g^2_{hZZ})^{\text{SM}} \) to keep \( \mu = \mu^{\text{SM}} \) fixed, which in turn implies \( \sigma_h > \sigma_h^{\text{SM}} \). Fig. 1 validates this line of thought and qualitatively reflects the CMS analysis.

But how general is this approach, or put differently, how solid is a limit on \( \Gamma_h \) obtained this way once we include unknown new physics effects? And letting aside the interpretation in terms of a constraint on the Higgs width, what are the more general ramifications of a measurement of the gluon-fusion \( ZZ \) and \( WW \) cross section away from the Higgs mass peak?

It is the purpose of this letter to address these questions from a new physics perspective with a particular emphasis on probability conservation. First we interpret the outlined Higgs width measurement from a unitarity perspective, which paves the way to the formulation of a simple and transparent BSM counterexample. We analyse the interplay of new resolved physics contributions to \( gg \rightarrow VV^* \) to both Higgs and continuum \( ZZ, WW \) production in light of electroweak precision constraints and finally point out that, enforcing \( \mu \simeq \mu^{\text{SM}} \) the off-shell measurement provides additional statistical pull to constrain the Higgs' \( CP \) nature in the presence of higher

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We mainly focus on the final state \( e^+e^-\mu^+\mu^- \) in the following. Generalizing our results to full leptonic \( ZZ \) decays is straightforward due to negligible identical fermion interference.

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FIG. 2: Representative Feynman diagram topologies contributing to \( gg \rightarrow ZZ \) with leptonic \( Z \) boson decays in the SM and theories with extended fermionic sectors.
dimensional operators (unresolved new physics). We also discuss off-shell measurements in WBF in Sec. V.

As we will see, in order to gain qualitative control of new physics effects in the Higgs off-shell region we cannot rely on effective theory calculations for the SM spectrum. We consequently keep all quarks dynamical and include finite mass effects of the bottom and top quarks. Our work therefore extends beyond the assumptions of Ref. [19] which has discussed the impact of new operators to high invariant mass measurements in detail recently. We only focus on modified $ggh$ and $hZZ/hWW$ interactions and neglect QED contributions throughout; they are negligible for high invariant masses when both $Z$s are fully reconstructed, but can be sensitive to the presence of new physics when studied on the Higgs peak via $h \rightarrow Z\gamma^\ast, \gamma^\ast\gamma^\ast$ [20]. We will mainly focus our discussion on $\sqrt{s} = 8$ TeV; our results straightforwardly generalize to run II.

Computations have been performed and cross checked with a combination of FeynArts/FormCalc/LoopTools [21], HELAS [22], MADGRAPH/MadEvent [23], and VBFNLO [24]. We have checked our results against [13] and find very good agreement.

II. HIGGS WIDTH MEASUREMENTS FROM $gg \rightarrow VV$: A UNITARITY PERSPECTIVE

In Fig. 3 we show the individual contributions of $pp \rightarrow ZZ^\ast \rightarrow e^+e^-\mu^+\mu^-$ that result from the Feynman diagrams of Fig. 2. We also include a comparison of the full Higgs contribution with the low energy effective theory [25] as implemented in MADGRAPH/MadEvent [23], which shows large deviations when the absorptive parts of the top quark loop are solved (the corresponding Cutkosky cut [26] is included in Fig. 2). Obviously, a reliable analysis of the high invariant mass region in correlation with the on-shell part cannot be obtained by applying effective theory simplifications. The CMS analysis [18] focuses on $m(4\ell) \geq 330$ GeV.

It is known that the interference between the triangle and box diagrams is destructive [12] above the $2m_t$ threshold. This large interference effect becomes transparent when calculating the cross section for the process $q\bar{q} \rightarrow ZZ$ with massive quarks in the initial state. It

\[ \sigma \left[ \text{pb} \right] \]

\[ \sqrt{s} \text{[TeV]} \]

\[ t\bar{t} \rightarrow ZZ \text{ full} \]

\[ t\bar{t} \rightarrow ZZ \text{ Gauge} \]

\[ t\bar{t} \rightarrow ZZ \text{ Yukawa} \]

\[ \text{CMS exclusion} \]

FIG. 4: Unpolarized $t\bar{t} \rightarrow ZZ$ cross section as function of energy. We demonstrate unitarity cancellations between the gauge and Yukawa-type interactions (blue solid and dashed; the dashed line lies on top of the solid line), yielding a well-defined SM cross section (orange). We also show the parameter choice that corresponds to the CMS-like exclusion of $\Gamma_h \approx 5 \times \Gamma_h^{\text{SM}}$ based on the strategy outlined in [2] and the introduction.
volves a highly non-trivial cancellation between the gauge and Yukawa sector interactions in $q\bar{q} \rightarrow Z_L Z_L$ \cite{27}, and is part of the absorptive $gg \rightarrow ZZ$ amplitude according to the branch cuts shown in Fig. 3. Even though both contributions are gauge invariant under QCD transformations, they are related by weak gauge invariance, and only their coherent sum with SM-like couplings is well-defined.

In Fig. 4 this is demonstrated for the unpolarized $t\bar{t} \rightarrow ZZ$ cross section: The $s$ channel Yukawa couplings of $tLtRH+\text{h.c.}$ conspire via a coupling relation with the weak gauge interactions $g_{L}t_{L}Zt_{L} + g_{R}t_{R}Zt_{R}$ when $-t,-u \sim s$. A simple rescaling of one part of the amplitude is tantamount to unitarity violation in the Fermion-gauge interactions. This leaves a crucial question of the limit obtained in \cite{18}: Is the theory underlying the width constraint well-defined?

The alert reader might object at this stage that such a question, in fact, is also well-motivated for Higgs couplings measurements as performed by ATLAS and CMS \cite{28} when Higgs couplings are varied independently throughout the SM Lagrangian. This is certainly true if one would like to understand deviations from an electroweak precision point of view. However, the situation for Higgs $\sigma \times BR$ phenomenology is fundamentally different. The relevant scale at which couplings are evaluated is the Higgs mass and $\sigma \times BR$ phenomenology is manifestly free of UV problems to leading order in the electroweak perturbative series expansion.\footnote{This will dramatically change when the measurements of differential weak boson fusion distributions will be scrutinized at high precision \cite{29}.}

To address this question quantitatively, we show the zeroth partial wave projection as a function of the partonic center of mass energy in Fig. 5. Unitarity is violated when $a_{0} > 0.5$ \cite{30}, and to contextualize our $gg \rightarrow ZZ$ findings with the SM Higgs sector we also show curves for SM $W_L W_L$ scattering that violates unitarity at low scales if the Higgs contribution is neglected.

Indeed, the $gg \rightarrow Z_L Z_L$ scattering is sensitive to the coupling rescaling as can be seen from Fig. 5; however the partial wave does not get close to 0.5. The amplitude is sufficiently diluted by loop factors $16\pi^2 \sim 160$. Once this factor is resolved the unitarity bound becomes relevant. This, however, corresponds to a regime where the narrow width approximation is violated entirely.

Although the limit in this channels is not afflicted with probability non-conservation, it should be clear that the invoked rescaling leads to an ill-defined electroweak sector as demonstrated in Fig. 4; the triangle and box contributions remain intimately related. If high invariant mass measurements in the $gg \rightarrow ZZ$ channel yield a statistically significant increase over the SM, the interpretation in terms of a modified Higgs width becomes model-dependent.

### III. DECODERELATING ON-SHELL AND OFF-SHELL MEASUREMENTS IN BSM THEORIES

#### 1. BSM contributions to Higgs production

The interplay in the absorptive parts of $gg \rightarrow ZZ$ linked by unitarity in the high invariant mass regime and non-decoupling of top loops tells us that the limit setting procedure outlined in the introduction is based on a consistency argument for the electroweak sector and is very specific to masses that are generated through chirality-
changing interactions. This paves the way to construct a straightforward counterexample of the Higgs width measurement as outlined above.

Consider \( \phi \), a scalar 3 under \( \text{SU}(3)_C \), coupled to the Higgs sector via portal interactions (see, e.g., Ref. [31])

\[
\mathcal{L}_\phi = |D_\mu \phi|^2 - \tilde{m}_\phi^2 |\phi|^2 - \lambda |\phi|^4 |H|^2 + \ldots .
\]

When the Higgs field obtains its vacuum expectation value \( v \), the field \( \phi \) induces a contribution to single-Higgs production due to the interaction \( \lambda v |\phi|^2 h \), as shown in Fig. [6]. The physical mass \( m_\phi^2 = \tilde{m}_\phi^2 + \lambda v^2 \) is essentially a free parameter \( m_\phi^2 > 0 \).

The new contribution gives an additional and potentially large constructive or destructive contribution to \( gg \rightarrow h \), depending on the sign and size of \( \lambda \). To enforce SM-like signal strengths \( \mu \simeq 1 \) we need to introduce a compensating contribution to the Higgs width (this could be interpreted as a Higgs-portal dark matter realization) and we have \( \Gamma_h > \Gamma_h^{\text{SM}} \).

Due to the scalar and electroweak singlet nature of the new fields we only change the triangle Higgs production contribution while leaving the box \( gg \rightarrow ZZ \) contributions unaltered. Note, for this particle there is no unitarity relation between the boxes and triangles. We show the individual contributions of the scalar color triplet in Fig. [3], which allows us to compare their behaviour with the SM contributions. The scalar loops can easily be suppressed by two orders of magnitude, leaving absolute and interference contributions to the total hadronic cross section small for energetic events. This behaviour is qualitatively known from supersymmetric scenarios [32] but has also been discussed in non-supersymmetric models [31].

To quantitatively understand why the scalars are suppressed at large invariant masses, let us consider the ratio of the off-shell \( gg \rightarrow h \) subamplitudes for scalars and tops (assuming \( m_\phi = m_t = y_t v/\sqrt{2} \lambda = y_t \) for simplicity):

\[
y_t \frac{\mathcal{M}_\phi}{\mathcal{M}_t} = \frac{1 + 2m_t^2 C_0(s, m_t)}{(s - 4m_t^2) C_0(s, m_t) - 2},
\]

where \( C_0(s, m_t^2) \) denotes the characteristic scalar three-point function following the Passarino-Veltman reduction [34]. The \( \phi \)-induced amplitude is suppressed \( \sim s^{-1} \), leading to a dominant behaviour of the top loops at large momenta. This means that, even though we have a modified Higgs phenomenology at around \( m_h \simeq 125 \text{ GeV} \) it is exactly the decoupling of the Higgs width according to Eq. [3] which renders the high invariant mass measurement insensitive to modifications of \( \Gamma_h \).

There is an interesting possibility when we consider larger \( \phi \) masses and larger couplings \( \lambda \). For invariant masses \( s^2 \geq 4m_\phi^2 \) we can have a sizeable constructive interference of the \( \phi \) diagrams with the top loops and as a result the cross section for large \( m(4\ell) \) rises again and we recover the qualitative behaviour of Ref. [11]. For these parameter choices, however, we find that the excess is smaller than expected for rescalings of \( g_{\text{ggh}} g_{ZZh} \) to keep \( \mu \simeq 1 \), Tab. [4]. Similar effects show up for light spectra \( m_\phi \lesssim 2m_t \), where this interference in destructive and the high invariant mass search region has a slightly smaller cross section although \( \Gamma_h/\Gamma_h^{\text{SM}} \approx 1 \), outside the current CMS exclusion.

In total, it is well possible to achieve \( \Gamma_h \gg \Gamma_h^{\text{SM}} \) without modifying the high invariant mass regime of \( pp \rightarrow 4 \ell \) and without running into unitarity issues as mentioned above. If such a contribution can be present the Higgs width is an essentially unconstrained parameter, at least for a measurement as outlined in [11, 18].

Even though Eq. [5] is a toy model to demonstrate the limitations of total Higgs width measurements in the \( gg \rightarrow 4 \ell \) channel, color triplets of this form appear in any supersymmetric BSM scenario and our argument has a broad validity, see e.g. Refs. [33, 37] for a discussion of squark contributions to Higgs production from gluon fusion in the MSSM. If the extra scalars are charged under flavour, e.g. they are top partners, exclusion will remain difficult [38] for the SUSY chimney regions (note, there are two chimney regions where one can hide 170 GeV and 70 GeV scalars). Despite being color charged, they could exist as stable particles on collider lifetimes when SUSY is relaxed [39]. Quite naturally, details quickly become highly model-dependent. By fixing \( m_\phi \) we can map \( \Gamma_h = 4.2 \times \Gamma_h^{\text{SM}} \) onto \( \lambda \) and obtain \( \sigma/\sigma^{\text{SM}} \) in Tab. [4]. When the \( m_\phi \) masses become heavy we start approaching an effective theory limit, which quickly decouples unless we allow non-perturbative couplings as \( m_\phi \) is not generated via the Higgs mechanism. This is also visible in Fig. [6] and we recover the qualitative behaviour of Ref. [11] also in this model as alluded to above. It is important to note,

| \( m_\phi \) (GeV) | \( \mu \) (h peak) | \( \Gamma_h/\Gamma_h^{\text{SM}} \) | \( \sigma/\sigma^{\text{SM}} \) [\( m(4\ell) \geq 330 \text{ GeV} \)] | 
|---|---|---|---|
| 70 | \( \simeq 1.0 \) | \( \simeq 5 \) | \(-2\% \) |
| 170 | \( \simeq 1.0 \) | \( \simeq 4.7 \) | \(+80\% \) |
| 170 | \( \simeq 1.0 \) | \( \simeq 1.7 \) | \(+6\% \) |

*We impose the cut set used by CMS [18] without the MELA cut [32].

TABLE I: Results for a single triplet scalar [5] giving the correlation between \( \mu, \Gamma_h/\Gamma_h^{\text{SM}} \) and high invariant mass cross section \( \sigma \) for the CMS selection cuts.

FIG. 6: New Feynman diagram topologies to Higgs production via gluon fusion arising from Eq. [5].
however, that the interpretation of the measurement is still far more complicated. If we imagine to become sensitive to the SM tail distribution within a small error whilst observing a SM-like Higgs peak phenomenology, the mass constraints to decorrelate the on- and off-shell region are weakened and heavier coloured bosons in this channel and scenario become essentially unconstrained.

Even though we have limited our discussion to $ZZ \to 4\ell$, the findings of this sections straightforwardly generalize to $ZZ \to 2\ell2\nu$ and $WW$.

2. **BSM contributions to continuum $ZZ$ and $WW$ production**

Our previous example shows that new contributions to $gg \to h$ can significantly loosen the bounds on the Higgs width interpretation. In a similar fashion we can imagine a situation where $\Gamma_h \neq \Gamma_{h,SM}^{\text{SM}}$ and the correlation of Eq. (2) and (3) is changed by new contributions to continuum $ZZ$ production. Such effects are expected in composite Higgs scenarios with vector-like quarks [10] and typically have non-trivial and non-diagonal electroweak interactions in the extended flavour sector. In realistic models [11] such sectors can be quite large, and the phenomenology becomes non-transparent especially when the different mass scales are resolved and effective field theory simplifications cannot be applied.

To keep our discussion as transparent as possible we will focus on a minimal, anomaly-free toy model of vector-like quarks:

$$-\mathcal{L} = m_Q Q_L^\dagger Q_R^\nu + m_d d_L^\dagger d_R^\nu + m_u u_L^\dagger u_R^\nu + \text{h.c.}$$

$$+ y_d^Q (Q_L H)d_R^\nu + y_u^Q (Q_L H\sigma^2 H^\dagger)u_R^\nu + \text{h.c.}$$

$$+ y_d^u (Q_R H)d_L^\nu + y_u^u (Q_R H\sigma^2 H^\dagger)u_L^\nu + \text{h.c.},$$

where $Q_L^Q, Q_R^Q = (3, 2, -\frac{1}{2})$, $d_L^d, d_R^d = (3, 1, -1)$, $u_L^u, u_R^u = (3, 1, 0)$ under SU(3)$_C \times$SU(2)$_L \times$U(1)$_Y$, i.e. we choose lepton hypercharge quantum numbers for simplicity. Depending on the relative size of the Yukawa coupling $y_i$ we can dial between the modifications in the triangle and box contributions, while the box contributions become sensitive to flavour changing charged and neutral current interactions that follow from diagonalizing Eq. (4). In this sense, Eq. (7) reflects the qualitative features of more realistic models and allows us to correlate the off-peak cross section with oblique parameter constraints.

Depending on the size of (non-)diagonal couplings we can, in principle, induce a new scale at large invariant $ZZ$ or $WW$ masses, see Figs. 2 and 4. In case the masses of these extra fermions are dominated by vectorial mass terms, oblique electroweak constraints [42] are largely insensitive to their presence. For $y_i \equiv 0$ we generically find $|S| \sim 10^{-3}$, $T = 0$ and $|U| \sim 10^{-2}$ over a broad range of parameter choices by explicit calculations. For pure vector-like terms we find for the CMS search region described in Ref. [18] (neglecting again MELA cuts)

$$\Delta_{WW}^{SM} \frac{\sigma_{BSM}^{SM} - \sigma_{SM}^{SM}}{\sigma_{SM}^{SM}} \simeq -3.8\%$$

even for light physical fermion mass scales

$$(m_{d,1}, m_{d,2}, m_{u,1}, m_{u,2}) = (300, 200, 400, 300) \text{ GeV},$$

which are already under pressure from direct search exclusion limits. The fermions quickly decouple for larger masses, owing to the dimension eight structure of the resulting effective theory. We can also define a high invariant $WW \to \ell\ell'2\nu$ mass region that is characterized by a large value of the transverse mass

$$m_T^2 = (p_T + \not{p_T})^2 - (p_T + \not{p_T})^2,$$

and we choose $m_T \geq 300 \text{ GeV}$ in the following in addition to $p_T(\ell) \geq 10 \text{ GeV}$, $|y_i| \leq 2.5$, $\Delta R(\ell') \geq 0.2$ and $p_T(\ell') \geq 20 \text{ GeV}$. As can be seen in Fig. 7, $WW$ has a qualitatively different sensitivity to this particular model class. We find in this region

$$\Delta_{WW}^{SM} \frac{\sigma_{BSM}^{SM} - \sigma_{SM}^{SM}}{\sigma_{SM}^{SM}} \simeq -3.8\%.$$

For our example mass point Eq. (9).

To understand how far oblique correction constraints limit the size of novel electroweak degrees of freedom, we can use the above toy model and take the mass spectrum of Eq. (9) as a baseline. Then we change the “chiral” components by increasing $y_i \geq 0$ until we reach the
boundary of the $S,T$ exclusion ellipse [42]. For the resulting parameter point we obtain cross sections analogous to Eqs. (8) and (11)

$$\Delta_{ZZ}^{+h} \simeq -4.3\%, \quad \Delta_{WW}^{+h} \simeq -3.7\%. \quad (12)$$

This example demonstrates that we can expect that large BSM contributions to the continuum $ZZ$ production are highly limited by electroweak precision constraints, especially because they also link to non-trivial gauge representations under $SU(3)_C$. The $ZZ$ and $WW$ production modes are directly correlated with $S,T,U$ and Higgs vev-induced mass terms in Eq. (7) quickly introduce a tension with electroweak precision constraints. They are known to have a significant impact on $H \to \gamma\gamma$ [44] and can be resolved at precision experiments [45].

IV. HIGHER DIMENSIONAL OPERATORS: IMPROVING $CP$ SENSITIVITY

As a final application of the off-shell measurements we discuss the impact of higher dimensional operators [46,47] on such a measurement. A recent and comprehensive analysis has been presented in Ref. [19]. Here we limit ourselves to the investigation of $CP$ properties, keeping all $m_t, m_b$ dependencies, and show that differential information in the high invariant mass regime can be used to add a statistically independent measurement to the $CP$ hypothesis test as already performed by ATLAS and CMS [48].

FIG. 8: $\cos \theta_1$ (top panels) and $\cos \theta^*$ distributions for the inclusive peak region and in the search region defined by CMS in Ref. [18]. $\alpha$ denotes the admixture of $CP$ even vs. $CP$ odd, signal strengths $\mu \sim 1$ in the peak region are enforced via $\Gamma_h \neq \Gamma_{SM}^h$. 

\begin{align*}
\Delta_{ZZ}^{+h} &\simeq -4.3\%, \quad \Delta_{WW}^{+h} \simeq -3.7\%. \quad (12)
\end{align*}
For this purpose, we focus on interactions
\[ \mathcal{L} \supset \sum_{V=Z,W^+} c_{\alpha, V} g_V m_V V^\mu h + \frac{c_{\alpha, V}}{m_V} \tilde{V}^{\mu\nu} V_{\mu\nu} A \] (13)
and define the physical Higgs boson as a linear combination of \( C_P \) even and odd states,
\[ X = \cos \alpha h + \sin \alpha A. \] (14)
We fix the signal strength for different angles \( \alpha \) by changing \( \Gamma_h \) accordingly and focus in the following on the two angles
\[ \cos \theta_1 = \frac{p(e^+) \cdot p_X}{\sqrt{p^2(e^+)} p^2_X} \bigg|_{Z \to e^+e^-}, \] (15)
\[ \cos \theta^* = \frac{p(Z \to e^+e^-) \cdot b}{\sqrt{p^2(Z \to e^+e^-)} b^2} |_{X}, \] (16)
where \( \ldots |_R \) refers to the rest frame \( R \) in which the angle is defined. \( p_X(X) = p_\mu(e^+) + p_\mu(e^-) + p_\mu(\mu^+) + p_\mu(\mu^-) \) coincides in the on-shell region with the Higgs boson’s rest frame, and \( b \) is an arbitrary three-vector along the positive beam direction. As defined, \( \cos \theta^* \) correlates the production mechanism with the resonance’s decay products by projecting onto the beam-component of the 4-lepton system. While \( \cos \theta^* \) is known to be flat, \( \cos \theta_1 \) is sensitive to the \( C_P \) properties of the Higgs boson when produced in the on-shell region, see Figs. and Ref. [49]. As can be seen, on top of a cross section increase due to the higher dimensional operator structure [19], there is complementary information in the spin/\( C_P \) observables.\(^5\)

V. OFF-SHELL MEASUREMENTS IN WEAK BOSON FUSION

The potentially unknown loop contributions that can decorrelate the on-shell and off-shell region in gluon fusion are not present in weak boson fusion, assuming indeed a \( C_P \) even SM-like Higgs boson. In these channels, the method of Ref. [11] becomes largely model-independent except for a potential asymmetric deviation of the \( WWh \) and \( ZZh \) couplings. This directly links to the \( T \) parameter and a deviation at tree level is expected to be small.

Furthermore, the weak boson fusion topology allows us to suppress gluon fusion contributions using forward tagging jets in opposite detector hemispheres with large invariant mass and rapidity gap [50]. By imposing an additional central jet veto [51], the gluon fusion events are almost entirely removed from the sample [52] and the impact on a correlation of the on- and off-shell regions will be unaffected by unknown physics beyond the SM as a consequence.

In Fig. 9 we show the result of such an analysis at NLO QCD [24, 53] (we choose a common rescaling of \( g_{ZZh} \) and \( g_{WWh} \) to achieve \( \mu \approx 1 \) in the on-peak region). The selection cuts are identical to CMS’ choice for the \( Z \) reconstruction and lepton selection. We lower the 4\( \ell \) mass cut to \( m(4\ell) \geq 130 \text{ GeV} \) to increase the statistics as much as possible. In addition, we employ typical WBF cuts such as outlined above
\[ p_T(j) > 20 \text{ GeV}, \quad \Delta R(jj) \geq 0.6, \quad |y_j| < 4.5, \quad \Delta y(jj) \geq 4.5, \quad y_{j1} \times y_{j2} < 0, \quad m(jj) \geq 800 \text{ GeV}, \] (17)
and a jet veto
\[ |y_j^{\text{veto}}| < 2.5, \quad p_T^{\text{veto}}(j) > 50 \text{ GeV}, \quad \Delta y(j^{\text{veto}}) > 0.3. \] (18)
The leptons need to be well separated from the jets \( \Delta R(\ell) \geq 0.6 \) and need to fall inside the tagging jets’ rapidity gap. We furthermore reject events with \( m(4\ell) \geq 2 \text{ TeV} \) to avoid picking up sensitivity from the region of phase space where the off-shell modification probes the unitarity-violating regime.

Obviously, when performed in the WBF channel (our reasoning also applies to the \( WW \) channel), we observe a similar behaviour [12], however, at a much smaller cross section \( \sigma(\text{WBF}) \approx 0.04 \text{ fb at } 14 \text{ TeV} \) slightly over light lepton flavours \( \ell = e, \mu \) [24]. Nonetheless such a measurement can be used to obtain a fairly model-independent measurement of the total Higgs width following [11] at large integrated luminosity, especially when statistically independent information from multiple WBF channels is combined.

\(^5\)Not included in Fig. [5] is the WBF contribution that can give rise to an additional \( \sim 10\% \) effect. We have checked the angular distributions with a modified version of VBFNLO and find no significant impact on the quoted results.
VI. SUMMARY AND CONCLUSIONS

After the Higgs discovery with a mass of $m_h \simeq 125$ GeV and TeV scale naturalness under siege, the total Higgs width is one of the most sensitive parameters to light physics beyond the standard model with a relation to the electroweak scale. A model-independent constraint on $\Gamma_h$ would have a huge impact on BSM physics. Correlating on- and off-shell Higgs production in $gg \to ZZ$, $WW$ and understanding cross section measurements in terms of a total Higgs width limit, however, only applies to the SM. Injecting the SM hypothesis into a global Higgs fit however will always yield much tighter constraints.\cite{4}

The on- and off-peak continuum regions can be decorrelated in gluon fusion by new degrees of freedom, which link to a modified Higgs phenomenology; large BSM effects in continuum $gg \to ZZ$, $WW$ seem unlikely given existing electroweak precision constraints.

If deviations at large invariant masses for $VV$ final states are observed in the future, unitarity of the scattering amplitude dictates that the interpretation of $\Gamma_h > \Gamma_h^{\text{SM}}$ will need to involve model-specific assumptions, a fact that is unavoidable in hadron collider physics.

The cross section measurement can be used to constrain momentum-dependent modifications of Higgs couplings that underlie the modelling of $\text{spin/CP}$ testing and the general limit setting procedure of higher dimensional operators.

Applying the strategy of Ref.\cite{11} to WBF allows us to formulate a constraint that is largely free of the gluon fusion shortcomings, however at considerably smaller cross sections.

A precise model-independent constraint on the total Higgs width $\Gamma_h \simeq \Gamma_h^{\text{SM}}$, if not a measurement of $\Gamma_h$ is challenging from a statistics and systematics point of view and probably remains the remit of a future theoretically and experimentally clean linear collider environment. At, e.g., a 250 GeV $e^+e^-$ machine the combined investigation of associated and WBF Higgs production, and exclusive final states $H \to b\bar{b}$, $WW$ allows us to formulate constraints on $\Gamma_h$ in the 10% range\cite{54}.

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