Unitarity-controlled resonances after Higgs discovery

Christoph Englert,\textsuperscript{1,*} Philip Harris,\textsuperscript{2,†} Michael Spannowsky,\textsuperscript{3,‡} and Michihisa Takeuchi\textsuperscript{4,§}

\textsuperscript{1}SUPA, School of Physics and Astronomy, University of Glasgow, Glasgow, G12 8QQ, United Kingdom
\textsuperscript{2}CERN, CH-1211 Geneva 23, Switzerland
\textsuperscript{3}Institute for Particle Physics Phenomenology, Department of Physics, Durham University, Durham, DH1 3LE, United Kingdom
\textsuperscript{4}Kavli IPMU (WPI), The University of Tokyo, Kashiwa, 277-8583, Japan

If the recently discovered Higgs boson’s couplings deviate from the Standard Model expectation, we may anticipate new resonant physics in the weak boson fusion channels resulting from high scale unitarity sum rules of longitudinal gauge boson scattering. Motivated by excesses in analyses of multi-leptons+missing energy+jets final states during run 1, we perform a phenomenological investigation of these channels at the LHC bounded by current Higgs coupling constraints. Such an approach constrains the prospects to observe such new physics at the LHC as a function of very few and generic parameters and allows the investigation of the strong requirement of probability conservation in the electroweak sector to high energies.

I. INTRODUCTION

After the discovery of the Higgs boson \cite{1} at the LHC and first preliminary tests of its coupling structure and strengths \cite{2,3}, a coarse-grained picture of consistency with the Standard Model (SM) has emerged. Resulting from Higgs quantum numbers, current constraints on the Higgs boson’s couplings, assuming a SM-value of the Higgs width or an upper limit on the Higgs’ coupling to electroweak gauge bosons indicate that the Higgs couplings to electroweak bosons agree with the SM expectation within \mathcal{O}(10\%) \cite{4,5}. This establishes the Higgs’ involvement in electroweak symmetry breaking and its role in the unitarization of massive longitudinal gauge boson scattering.

However, current constraints leave a lot of space for deviations from the SM-like implementation of electroweak symmetry breaking. In particular, small deviations from the SM Higgs coupling pattern are expected in a very broad class of models that explain the presence of the electroweak scale as a dimensional transmutation effect \cite{6,7}. In particular, these include composite Higgs scenarios where we expect new contributions from composite states analogous to the rho meson. Explicit examples have been discussed in the literature, mostly in the context of AdS/CFT duality, see e.g. \cite{8}.

Owing to the fact that any modification from the SM Higgs couplings explicitly introduces unitarity violation, novel resonant physics is likely to enter at a scale \( Q^2 \gg m_h^2 \) to conserve probability \cite{9} if we indeed deal with non-SM Higgs interactions. Weak boson scattering processes are theoretically well-motivated probes of such dynamics, correlating the size of the new physics effects with the deviation of the observed Higgs phenomenology from the SM.

Accessing longitudinal gauge boson scattering (which is highly sensitive to BSM effects) at the LHC in a phenomenologically useful way is difficult. Due to almost conserved light quark and lepton currents, weak boson fusion (WBF, for analyses see \cite{10,11}) is not too sensitive to modifications of the involved Higgs couplings. The Higgs exchange at energies \( m(VV) \gg m_h \) in a Higgs doublet model provides a destructive contribution to \( VVqq \) (\( V = W^\pm, Z \)) production. Thus, a \( \sim 10\% \) cross section excess at the LHC for inclusive WBF is mainly due to the smaller destructive Higgs contribution for smaller couplings, rather than diverging \( qq \to qVq \) processes getting tamed by the polynomial parton density function suppression at large parton energy fractions.

Nevertheless, it is important to realize that, if \( V_LV_L \to V_LV_L \) (Fig. 1) scattering violates the unitarity bound, the (leading order) electroweak sector becomes ill-defined, and there is no theoretically consistent interpretation of constraints and measurements even if the alternate hypothesis seems well-behaved \cite{12}.

Current analyses mostly focus on studying the impact of a subset of the 59 dimension-six operators (neglecting flavor structures) \cite{13} on Higgs physics in the on-and off-shell region. In this paper, we take a complementary approach and address the question of what to expect in WBF processes when unitarity is explicitly enforced by additional resonances in the TeV regime, following a strong-interaction paradigm.

If additional resonances in \( VV \) scattering are present, an identification will depend on their mass, width and

*Electronic address: christoph.englert@glasgow.ac.uk
†Electronic address: philip.coleman.harris@cern.ch
‡Electronic address: michael.spannowsky@durham.ac.uk
§Electronic address: michihisa.takeuchi@kcl.ac.uk

*In a general gauge the Goldstone contributions to the amplitude vanish in the chiral limit, signalling a vanishing contribution from longitudinal degrees of freedom at high invariant masses.
It is worth noting that similar sum rules cannot be formulated for $W^+W^-$ scattering, the sum rules read:

\begin{align}
g_{WWWW} &= g_{WW\gamma}^2 + \sum_i g_{WWZ_i}^2, \quad (1a) \\
4m_W^2 g_{WWWW} &= \sum_i 3m_i^2 g_{WWZ_i}^2 + \sum_i g_{WWH_i}^2, \quad (1b)
\end{align}

and for $WW \rightarrow ZZ$ (and crossed) scattering these are modified to

\begin{align}
g_{WWZZ} &= \sum_i g_{W_i}^2, \quad (1c) \\
2(m_W^2 + m_Z^2) g_{WWZZ} &= \sum_i \left(3m_i^2 - \frac{(m_Z^2 - m_W^2)^2}{m_t^2} \right) g_{W_i}^2 + \sum_i g_{WWH_i} g_{ZZH_i}. \quad (1d)
\end{align}

In both sums the index $i = 1$ refers to the SM $W$, $Z$ and Higgs bosons, respectively. It is important to realize that due to $SU(2)_L$ invariance (e.g. the absence of a quartic $Z$ interaction) the reasoning along the above lines does not apply to $ZZ \rightarrow ZZ$ scattering. In the high energy regime the Higgs exchange diagrams conspire

$$
\mathcal{M}(Z_LZ_L \rightarrow Z_LZ_L) \sim s + t + u = 4m_Z^2, \quad (2)
$$

* i.e. the scattering amplitude becomes independent of the center of mass energy. Hence, on the one hand, in scenarios where unitarity in $WW$ and $WZ$ scattering is enforced by iso-vectors, we do not expect new resonant structures in $pp \rightarrow 4\ell + 2j$. On the other hand if unitarity is conserved via the exchange of iso-scalar states, this channel will provide a phenomenological smoking gun. Obviously this is not a novel insight and under discussion in the context of e.g. Higgs portal scenarios. We will not investigate the $ZZ$ channel along this line in further detail.

For the purpose of this paper we start with a minimal, yet powerful set of assumptions, that can be reconciled in models that range from (perturbative and large $N$) AdS/CFT duality over SUSY to simple Higgs portal scenarios. We will focus on a vectorial realization of unitarity, assuming an electroweak doublet nature of the Higgs boson. This represents an alternative benchmark of new resonant physics involved in the mechanism of EWSB which has been largely ignored after the Higgs discovery so far.

The first rule of Eq. (1) is typically a consequence of gauge invariance while the second rule reflects the particular mechanism of EWSB. Similar sum rules exist for massive $q\bar{q} \rightarrow V_LV_L$ scattering, linking the Yukawa sector to the gauge sector. We are predominantly interested in a modified Higgs phenomenology in the standard WBF search channels. It is however important to note that the latter sum rules also predict new resonant

\footnote{See \cite{19} for a detailed discussion of WBF signatures in Higgs triplet scenarios.}
states in Drell-Yan type production \cite{21} (for a recent comprehensive discussion see also \cite{22}) or gluon fusion induced $VVjj$ production. For this analysis, gluon fusion events can efficiently be removed by imposing selection criteria \cite{23}; this process is neglected further on (see below).

The presence of unitarizing spin one resonances is tantamount to a modification of the 4-point gauge interactions when we choose the trilinear couplings to be SM-like. In higher dimensional and dual composite Higgs scenarios this fact is typically encoded in multiple definitions of the tree-level Weinberg angle and a resulting constraint from the $\rho$ parameter. The quartic gauge couplings are currently not well constrained and we use this freedom to saturate the above sum rules via a non-standard value of $g_{WVWV}$ and $g_{WVVZ}$. The deviations from the SM values, especially in the vicinity of the SM when $g_{WVZ} = g_{WVZ} = 0$ are small and well within the quartic coupling measurements' uncertainty.

\section{Results}

\subsection{Details of the simulation}

Using Eq. \cite{1}, we have a simple parameterization of new physics interactions in terms of mass and width of the new vector state, and Higgs coupling modification parameter. Since we do not specify a complete model we treat the extra boson widths as nuisance parameters. In concrete models the width can span a range from rather narrow to extremely wide. Masses are typically constrained by electroweak precision measurements. Since the sum rules give an independent prediction, we will not consider these corrections further.

We use a modified version of \texttt{VBFNLO} \cite{22} to simulate the weak boson fusion channel events for fully partonic final states inputting the relevant model parameters mentioned above. Since WBF can be identified as “double-DIS” we can efficiently include the impact of higher order QCD corrections on differential distributions by dynamically choosing the $t$-channel momentum transfer of the electroweak bosons as the factorization and renormalization scales \cite{23} irrespective of new resonant structures in the leptonic final state \cite{26}. We generate the gluon fusion contribution using again \texttt{VBFNLO}, but find that they are negligible for typical WBF requirements. As benchmarks we consider the following parameter points, defining $\alpha = g_H / g_{SM}$,

\begin{align*}
    m_{W',Z'} &= 700 \text{ GeV}, \quad \Gamma_{W',Z'} = 3 \text{ GeV}, \quad \alpha = 0.9, \\
    m_{W',Z'} &= 1000 \text{ GeV}, \quad \Gamma_{W',Z'} = 7 \text{ GeV}, \quad \alpha = 0.9, \\
    m_{W',Z'} &= 700 \text{ GeV}, \quad \Gamma_{W',Z'} = 10 \text{ GeV}, \quad \alpha = 0.5, \\
    m_{W',Z'} &= 1000 \text{ GeV}, \quad \Gamma_{W',Z'} = 30 \text{ GeV}, \quad \alpha = 0.5.
\end{align*}

\section{Projections for $2l + \cancel{E_T} + jj$ production}

For the analysis of the $2l + \cancel{E_T} + jj$ channel, we follow the event reconstruction depicted in Sec. \ref{sec:details_sim} where we require exactly 2 isolated leptons. We impose staggered cuts on the transverse momenta of both leptons, i.e.

\begin{align*}
    p_{T,l_1} &> 120 \text{ GeV}, \\
    p_{T,l_2} &> 80 \text{ GeV}.
\end{align*}

Additionally, for the two most forward jets with $p_T > 40$ GeV we impose a WBF selection of:

\begin{align*}
    y_{j_1} \times y_{j_2} &< 0, \\
    |y_{j_1} - y_{j_2}| &> 4.0, \\
    m_{j_1,j_2} &> 800 \text{ GeV}.
\end{align*}
The heavy resonance is reconstructed by requiring the transverse mass $m_T > 350$ GeV where

$$m_{T,2l}^2 = \left( \sqrt{m_T^2 - p_{T,miss}^2} + |p_{T,miss}| \right)^2 - \left( p_{T,ll} + p_{T,miss} \right)^2.$$  (6)

We show the results after each analysis step in Tab. II B. The WW channel is the most complicated final state in terms of background composition and final state reconstruction given the expected detector performance.

There are two major conclusions at this stage:

(i) Due to the departure of $\alpha < 1$, a continuum enhancement for the BSM signal over the expected electroweak $VVjj$ distribution is present. This excess is not big enough to be useful to constrain this scenario efficiently; this also applies to the novel non-resonant $t$- and $u$-channel contributions. When we approach the SM limit (as supported by current measurements) the signal contributions quickly decouple and the analysis loses sensitivity even for small widths. In this sense, the phase space region complementary to the on-shell Higgs region cannot be efficiently exploited phenomenologically. Deviations from the SM WBF hypothesis are typically of the order of 10%, which can easily be obstructed by additional experimental and theoretical systematics (see e.g. 31) neglected in this analysis. The gluon fusion contribution is highly suppressed and we do not include it in Fig. 3.

(ii) We therefore proceed to reconstruct the presence of $s$-channel resonances in a bump search sensitive to both the $WWjj$ and $ZZjj$ subprocesses. The $2\ell + \not{E}_T + jj$ final state, however, is also characterized by a relatively large fraction of missing energy, which substantially hampers a bump search, Fig. 3(a). This again becomes more severe when we turn to Higgs couplings in the vicinity of the SM expectation, see Fig. 2.
TABLE III: Results for 3 lepton search. The cross sections are given in femtobarn, corresponding to proton-proton collisions at \( \sqrt{s} = 14 \) TeV. Further details on the cuts can be found in the text.

| Sample | lepton cuts | WBF cuts | \( m_{T,3l} \) |
|--------|-------------|-----------|----------------|
| \((h \to WW)JJ GF\) | 0.03 | < 0.01 | < 0.01 |
| \(t\bar{t}+\text{jets}\) | 82.76 | 0.22 | 0.17 |
| \(WW+\text{jets}\) | 6.32 | 1.72 | 1.09 |
| \(WZ+\text{jets}\) | 0.47 | 0.07 | 0.04 |
| \(ZZ+\text{jets}\) | 0.64 | 0.12 | 0.06 |
| \(Z+\text{jets}\) | 0.08 | < 0.01 | < 0.01 |
| \(m_{W',Z'} = 700 \text{ GeV}, \alpha = 0.9\) | 6.37 | 1.84 | 1.24 |
| \(m_{W',Z'} = 1000 \text{ GeV}, \alpha = 0.9\) | 5.89 | 1.68 | 1.18 |
| \(m_{W',Z'} = 1500 \text{ GeV}, \alpha = 0.9\) | 5.80 | 1.67 | 1.13 |
| \(m_{W',Z'} = 2000 \text{ GeV}, \alpha = 0.9\) | 5.84 | 1.64 | 1.09 |
| \(m_{W',Z'} = 700 \text{ GeV}, \alpha = 0.5\) | 8.43 | 2.30 | 1.73 |
| \(m_{W',Z'} = 1000 \text{ GeV}, \alpha = 0.5\) | 6.85 | 1.96 | 1.41 |
| \(m_{W',Z'} = 1500 \text{ GeV}, \alpha = 0.5\) | 6.44 | 1.78 | 1.22 |
| \(m_{W',Z'} = 2000 \text{ GeV}, \alpha = 0.5\) | 6.36 | 1.77 | 1.17 |

TABLE II: Results for 4 lepton search. The cross sections are given in femtobarn, corresponding to proton-proton collisions at \( \sqrt{s} = 14 \) TeV. The t- and u-channel mass scales have no significant impact. Further details on the cuts can be found in the text.

| Sample | lepton cuts | WBF cuts | \( m_{4l} \) |
|--------|-------------|-----------|----------------|
| \(ZZ+\text{jets}\) | 0.25 | 0.074 | 0.054 |
| \(\alpha = 0.9\) | 0.23 | 0.075 | 0.053 |
| \(\alpha = 0.5\) | 0.24 | 0.078 | 0.058 |

C. Projections for WBF 4l + jj production

The systematic shortcomings resulting from the missing transverse energy in WW final state are not present in the fully-reconstructible final state 4l + jj. We require exactly 4 leptons and follow Eqs. (4) and (5). Additionally, the four lepton mass is required to be \( m_{4l} > 350 \) GeV. The cut flow is depicted in Tab. IIB. The backgrounds are manageable, however, for the considered scenario there is no s-channel resonance and again the continuum enhancement is too small to provide solid discrimination from a non-SM realization of EWSB, if we compare the deviations of Tab. IIB to \( \mathcal{O}(10\%) \) expected experimental systematic uncertainties (see Fig. 3(c)). However, this channel remains a “golden channel” for an additional isoscalar resonance, and the comparison to WW and WZ analyses will allow to reach a fine-grained picture of the involved dynamics if resonances are discovered in either of the mentioned channels.

D. Projections for 3l + \( E_T + jj \) production

The 3l + \( E_T + jj \) “interpolates” between the previous analyses. There is no pollution from gluon fusion events (even if we allow a significant coupling of \( Z' \) to the fermion sector). Additionally, the major backgrounds of the 2l + \( E_T + jj \) can be completely removed through the requirement of exactly three isolated leptons with \( p_{T,l} > 15 \) GeV, with no charge requirement. We then require the cuts given in Eqs. (4) and (5) and for the lepton and WBF selection respectively. The signal is extracted following a final selection \( m_{T,3l} > 350 \) GeV, where

\[
\begin{align*}
    m_{T,3l}^2 & = \left[ \sqrt{m_{1l2l3}^2 + p_{T,1l2l3}^2} + |p_{T,\text{miss}}| \right]^2 \\
    & - |p_{T,1l2l3} + p_{T,\text{miss}}|^2.
\end{align*}
\]

(7)

Although a substantial amount of missing energy is present, the lepton-\( E_T \) system is highly correlated in this final state, allowing for recovery of most of the mass discrimination through Eq. (7), see Fig. 3(b).
As a result of the non-standard Higgs coupling, a large enhancement of the signal strength is present. This can be seen compared to the standard model background in Fig. 4.

E. Setting limits with $3l + \slashed{E}_T + jj$ production

Combining the analyses of the previous sections, we can see that the potential presence of new vector resonances for $\sim 10\%$ Higgs coupling deviations can be highly constrained with the $3l + \slashed{E}_T + jj$ channel. Although we believe that more advanced limit setting procedures that deal with full correlations can eventually be used to constrain iso-triplet states in the $2l + \slashed{E}_T + jj$ and $4l + jj$ final states, the $3l + \slashed{E}_T + jj$ provides the most direct avenue to constrain such a scenario.

We thus quote an expected significance using $3l + \slashed{E}_T + jj$ final states (Sec. II D) on the basis of mass, width and modified Higgs coupling strength in Figs. 5(a) and 5(b). The signal extraction is performed over a mass window of $0.3 \times m_{W'}$ in the transverse mass Eq. (7). The calculated significance follows from:

$$S = \frac{N(\text{BSM}) - N(\text{WBF,SM})}{\sqrt{N(\text{bkg,non-WBF}) + N(\text{WBF,SM})}},$$

where the individual $N$s refer to the signal counts at a given luminosity. Using this measure we can isolate a statistically significant deviation from the SM WBF distribution outside the Higgs signal region, taking into account the irreducible background in the WZ channel.

Already for a target luminosity of run 2 of 100/fb, a large parameter region can be explored in the $3l + \slashed{E}_T + jj$ channel. A crucial parameter in this analysis is the width of the additional resonance, which we take as a free parameter in our analysis. With an increasing width the signal decouples quickly, but stringent constraints can still be formulated at a high-luminosity LHC, especially if new physics gives rise to only a percent-level deformation of the SM Higgs interactions, see Fig. 6.

III. SUMMARY AND CONCLUSIONS

The search for new physics interactions after the discovery of the Higgs boson remains one of the main tar-

FIG. 5: Projections of the $3l + \slashed{E}_T + jj$ analysis for a small integrated luminosity of 100/fb.

FIG. 6: Projections of the $3l + \slashed{E}_T + jj$ 95% confidence level contours for 100/fb (green), 500/fb (orange) and 3000/fb (red). The Higgs coupling deviation is $\alpha^2 = 0.95$.
gets of the LHC. Current constraints on Higgs couplings inferred from run 1 signal strength measurements, in particular in the ZZ channel, leave a lot of space for the appearance of new resonant phenomena at the TeV scale. These can, but not necessarily have to be iso-scalar degrees of freedom. To this end we have combined the observation of a SM-like Higgs boson with the appearance of new iso-vectorial degrees of freedom at the TeV scale. These are further corroborated by small excesses in similar and recent searches during run 1 [10]. Solely based on probability conservation, we provide predictions for the weak boson fusion channels, which are theoretically well motivated candidate processes to study resonant phenomena connected to unitarity and the anatomy of electroweak symmetry breaking. Our approach of saturating $W, Z$ unitarity sum rules with a single set of vector resonances as a function of vector boson mass and Higgs coupling deviation provides a complementary approach to singlet-extended Higgs sectors with a highly modified TeV-scale LHC phenomenology. While resonances and continuum excesses due to new $t-\mu$-channel contributions and a smaller destructive Higgs contribution at large multi-lepton mass might be challenging to observe in $2\ell+E_T+jj$ and $4\ell+jj$ production, we have shown that an analysis of $3\ell+E_T+jj$ production provides an excellent avenue to constrain or even observe the presence of such states over a broad range of mass and width scales. With comparably low integrated luminosity at the LHC, such an analysis captures complimentary and necessary information to pin down the very character of new physics for small deviations of the Higgs on-shell phenomenology, especially when results across the different WBF channels are combined.

Acknowledgments — CE is supported by the Institute for Particle Physics Phenomenology Associate scheme. We thank Sven Heinemeyer for helpful conversations. This research was supported in part by the European Commission through the “HiggsTools” Initial Training Network PITN-GA-2012-316704.

[1] F. Englert and R. Brout, Phys. Rev. Lett. 13 (1964) 321. P. W. Higgs, Phys. Lett. 12 (1964) 132 and Phys. Rev. Lett. 13 (1964) 508, G. S. Guralnik, C. R. Hagen and T. W. B. Kibble, Phys. Rev. Lett. 13 (1964) 585.

[2] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716 (2012) 1 [arXiv:1207.7214 [hep-ex]].

[3] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716 (2012) 30 [arXiv:1207.7235 [hep-ex]].

[4] M. Duhrsen, S. Heinemeyer, H. Logan, D. Rainwater, G. Weiglein and D. Zeppenfeld, Phys. Rev. D 70 (2004) 113009 [hep-ph/0406323]. B. A. Dobrescu and J. D. Lykken, JHEP 1302 (2013) 073 [arXiv:1210.3342 [hep-ph]]. P. Bechtle, S. Heinemeyer, O. Stal, T. Stefañiak and G. Weiglein, JHEP 1411 (2014) 039 [arXiv:1403.1582 [hep-ph]]. J. Ellis, V. Sanz and T. You, JHEP 1407 (2014) 036 [arXiv:1404.3667 [hep-ph]]. J. Ellis, V. Sanz and T. You, JHEP 1407 (2014) 036 [arXiv:1410.7703 [hep-ph]].

[5] D. Lopez-Val, T. Plehn and M. Rauch, JHEP 1310 (2013) 134 [arXiv:1308.1979 [hep-ph]]. C. Englert, A. Freitas, M. Mihleitner, T. Plehn, M. Rauch, M. Spira and K. Walz, J. Phys. G 41 (2014) 113001 [arXiv:1403.7191 [hep-ph]].

[6] G. F. Giudice, C. Grojean, A. Pomarol and R. Rattazzi, JHEP 0706 (2007) 045 [arXiv:0706.1864 [hep-ph]].

[7] C. Englert, J. Jaeckel, V. V. Khoze and M. Spannowsky, JHEP 1304 (2013) 060 [arXiv:1301.4224 [hep-ph]]. M. Heikinheimo, A. Racicoppi, M. Raidal and C. Spethmann, Phys. Lett. B 726 (2013) 781 [arXiv:1307.7146]. J. D. Clarke, R. Foot and R. R. Volkas, JHEP 1402 (2014) 123 [arXiv:1310.8042 [hep-ph]]. A. Farzinnia and J. Ren, arXiv:1405.0498 [hep-ph].

[8] G. Cacciapaglia, C. Csaki, G. Marandella and J. Terning, JHEP 0702 (2007) 036 [hep-ph/0611358]. J. Galloway, B. McElrath, J. McRaven and J. Terning, JHEP 0911 (2009) 031 [arXiv:0908.0532 [hep-ph]].

[9] J. M. Cornwall, D. N. Levin and G. Tiktopoulos, Phys. Rev. Lett. 30 (1973) 1268 [Erratum-ibid. 31 (1973) 572]. J. M. Cornwall, D. N. Levin and G. Tiktopoulos, Phys. Rev. D 10 (1974) 1145 [Erratum-ibid. D 11 (1975) 972].

[10] V. D. Barger, K. -m. Cheung, T. Han and D. Zeppenfeld, Phys. Rev. D 44 (1991) 2701 [Erratum-ibid. D 48 (1993) 5444]. J. Bagger, V. D. Barger, K. -m. Cheung, J. F. Gunion, T. Han, G. A. Ladinsky, R. Rosenfeld and C. - P. Yuan, Phys. Rev. D 52 (1995) 3878 [hep-ph/9504426]. D. L. Rainwater and D. Zeppenfeld, Phys. Rev. D 60 (1999) 113004 [Erratum-ibid. D 61 (2000) 099901] [hep-ph/9906218]. N. Kauer, T. Plehn, D. L. Rainwater and D. Zeppenfeld, Phys. Lett. B 503 (2001) 13 [hep-ph/0012351]. C. Englert, B. Jager, M. Worek and D. Zeppenfeld, Phys. Rev. D 80 (2009) 035027 [arXiv:0810.4861 [hep-ph]].

[11] A. Ballestrero, D. B. Franzosi, L. Oggero and E. Maina, JHEP 1203 (2012) 031 [arXiv:1112.1171 [hep-ph]]. P. Borel, R. Franceschini, R. Rattazzi and A. Walzer, JHEP 1206 (2012) 122 [arXiv:1202.1904 [hep-ph]]. A. Freitas and J. S. Gainer, Phys. Rev. D 88 (2013) 1, 017302 [arXiv:1212.3598].

[12] C. Englert and M. Spannowsky, Phys. Rev. D 90, no. 5, 055003 (2014) [arXiv:1405.0285 [hep-ph]]. A. Biekotter, A. Knochel, M. Kraemer, D. Liu and F. Riva, [arXiv:1406.7320 [hep-ph]].

[13] B. Grzadkowski, M. Iskrzynski, M. Misiak and J. Rosiek, JHEP 1010 (2010) 085 [arXiv:1008.4884 [hep-ph]].

[14] A. Birkedal, K. Matchev and M. Perelstein, Phys. Rev. Lett. 94 (2005) 191803 [hep-ph/0412278].

[15] C. Csaki, C. Grojean, H. Murayama, L. Pilo and J. Terning, Phys. Rev. D 69 (2004) 055006 [hep-ph/0305237]. C. Csaki, C. Grojean, L. Pilo and J. Terning, Phys. Rev. Lett. 92 (2004) 181802 [hep-ph/0308038]. C. Csaki, J. Hubisz and P. Meade, [hep-ph/0510275].

[16] S. Chatrchyan et al. [CMS Collaboration], Phys. Rev. D 90, no. 3, 032006 (2014) [arXiv:1404.5801 [hep-ex]]. G. Aad et al. [ATLAS Collaboration], [arXiv:1503.03290]
[hep-ex]. S. Chatrchyan et al. [CMS Collaboration], CMS PAS HIG-14-008. S. Chatrchyan et al. [CMS Collaboration], CMS EXO-12-012.

[17] A. Alboteanu, W. Kilian and J. Reuter, JHEP 0811 (2008) 010 [arXiv:0806.4145 [hep-ph]].

[18] T. Binoth and J. J. van der Bij, Z. Phys. C 75 (1997) 17 [hep-ph/9608245]. M. Bowen, Y. Cui and J. D. Wells, JHEP 0703 (2007) 036 [hep-ph/0701035]. C. Englert, T. Plehn, D. Zerwas and P. M. Zerwas, Phys. Lett. B 703 (2011) 298 [arXiv:1106.3007 [hep-ph]]. E. Weihs and J. Zurita, JHEP 1202 (2012) 041 [arXiv:1110.5909 [hep-ph]].

[19] A. Alboteanu, W. Kilian and J. Reuter, JHEP 0811 (2008) 010 [arXiv:0806.4145 [hep-ph]].

[20] S. Godfrey and K. Moats, Phys. Rev. D 81 (2010) 075026 [arXiv:1003.3053 [hep-ph]]. R. Killick, K. Kumar and H. E. Logan, Phys. Rev. D 88 (2013) 033015 [arXiv:1305.7236 [hep-ph]]. C. Englert, E. Re and M. Spannowsky, Phys. Rev. D 88 (2013) 035024 [arXiv:1306.6228 [hep-ph]]. C. W. Chiang, A. L. Kuo and K. Yagyu, JHEP 1310 (2013) 072 [arXiv:1307.7526 [hep-ph]].

[21] S. Godfrey and K. Moats, Phys. Rev. D 81 (2010) 075026 [arXiv:1003.3053 [hep-ph]].

[22] D. Pappadopulo, A. Thamm, R. Torre and A. Wulzer, JHEP 1409 (2014) 060 [arXiv:1402.4431 [hep-ph]].

[23] V. D. Barger, R. J. N. Phillips and D. Zeppenfeld, Phys. Lett. B 346 (1995) 106 [hep-ph/9412276]. J. R. Andersen, K. Arnold and D. Zeppenfeld, JHEP 1006 (2010) 091 [arXiv:1001.3822 [hep-ph]]. J. R. Andersen, C. Englert and M. Spannowsky, Phys. Rev. D 87 (2013) 1, 015019 [arXiv:1211.3011 [hep-ph]].

[24] K. Arnold, M. Bahr, G. Bozzi, F. Campanario, C. Englert, T. Figy, N. Greiner and C. Hackstein et al., Comput. Phys. Commun. 180 (2009) 1661 [arXiv:0811.4559 [hep-ph]].

[25] B. Jager, C. Oleari and D. Zeppenfeld, JHEP 0607 (2006) 015 [hep-ph/0603177]. G. Bozzi, B. Jager, C. Oleari and D. Zeppenfeld, Phys. Rev. D 75 (2007) 073004 [hep-ph/0701105].

[26] C. Englert, B. Jager and D. Zeppenfeld, JHEP 0903 (2009) 060 [arXiv:0812.2564 [hep-ph]].

[27] C. W. Chiang, A. L. Kuo and K. Yagyu, JHEP 1310 (2013) 072 [arXiv:1307.7526 [hep-ph]].

[28] M. L. Mangano, M. Moretti, F. Piccinini, R. Pittau and A. D. Polosa, JHEP 0307 (2003) 001 [hep-ph/0206293].

[29] ATLAS collaboration, ATL-PHYS-PUB-2013-004.

[30] M. Cacciari, G. P. Salam and G. Soyez, JHEP 0804, 063 (2008) [arXiv:0802.1189 [hep-ph]].

[31] S. Willenbrock and G. Valencia, Phys. Lett. B 259 (1991) 373. R. G. Stuart, Phys. Lett. B 262 (1991) 113. M. Nowakowski and A. Pilaftsis, Z. Phys. C 60 (1993) 121 [hep-ph/9305321]. U. Baur and D. Zeppenfeld, Phys. Rev. Lett. 75 (1995) 1002 [hep-ph/9503344].