1 High-Energy Physics Tools for Effective Field Theories[1]

1.1 Introduction

The discovery by the ATLAS and CMS collaborations of a Higgs boson boson with a mass of about 125 GeV [1, 2] has marked an important step forward in the study and the understanding of the electroweak symmetry breaking mechanism. Although the currently measured properties of this newly discovered boson seem to be compatible with the Standard Model expectation, the recent start of the second LHC experimental run has risen new hopes to detect phenomena beyond the Standard Model. In this context, present and future LHC data could be interpreted in an effective field theory framework where departures from the Standard Model are organized as a series expansion in the new physics energy scale Λ that is assumed to be large.

The leading effects implied by such an effective field theory description usually consist of dimension-six operators that are supplemented to the Standard Model Lagrangian, each of these being associated with a new interaction strength. The number of independent coefficients is usually large, but important classes of observables turn to only depend on a much smaller subset of parameters. The effective field theory approach is therefore testable and the results could be reinterpreted to constrain explicit new physics models. Consequently, the development of high-energy physics tools able to perform computations in the effective field theory context has been a very active field during the last decades. Recent progress addresses precision calculations for the production cross section of a single Higgs boson (see Section 1.2) and a pair of Higgs bosons (see Section 1.3), and for Higgs decays (see Section 1.4). In addition, machineries have been built so that we are now able to charaterize Higgs decays by means of pseudo-observables (see Section 1.5), as well as all Higgs properties by means of Monte Carlo event generation within the MADGRAPH5_aMC@NLO platform (see Section 1.6). Monte Carlo simulations can also be performed with the VBFNLO program, that includes both next-to-leading order QCD and higher-dimensional operator effects to vector boson fusion processes (see Section 1.7), and with the WHIZARD package at the leading-order accuracy (see Section 1.8). As all codes are using different conventions, we review the notation in each case.

An important aspect of the effective field theory approach is the freedom in the choice of the operator basis, so that a given effect could be modeled by several different combinations of operators at a fixed order in the effective energy scale expansion. This is related to the possibility of redefining the Standard Model fields in such a way that the Standard Model Lagrangian is unaltered, while certain combinations of dimension-six operators proportional to the Standard Model equations of motion can be eliminated up to subleading higher-dimensional effects. Different complete operator bases have been proposed in the past, and although each of them yields the same predictions, they present specific and different advantages. Existing calculations or tools are however often bound to a given basis choice, and it is desirable to be able to reuse results derived in the context of one basis in another basis. The ROSETTA platform (see Section 1.9) has been very recently released to close this gap.

1.2 HIGLU: Higgs Boson Production via Gluon Fusion

The program HIGLU [3,4] calculates the Higgs boson production cross section via gluon-fusion up to the next-to-next-to-leading order (NNLO) accuracy in QCD [5–15] and includes next-to-leading order (NLO) electroweak corrections [16–24] within the Standard Model, and up to the NNLO in QCD [5–15, 25–34] for the minimal supersymmetric extension of the Standard Model. While the supersymmetric electroweak corrections are unknown, the known genuine supersymmetric QCD corrections [35–45] are not included in the program. Starting from the Standard Model Higgs results, the contributions of dimension-six operators beyond the Standard Model are included up to the NNLO in QCD. The latter extension is based on the effective Lagrangian (in the heavy top limit for the Standard Model part) [46–49],

$$\mathcal{L}_{\text{eff}} = \frac{\alpha_s}{\pi}\left[\frac{c_t}{12}(1+\delta) + c_g\right] G_a^{\mu\nu} G^a_{\mu\nu} \frac{H}{v}, \tag{1}$$

where H denotes the Standard Model physical Higgs boson and v the vacuum expectation value of the neutral component of the Higgs field Φ. Moreover, the gluon field strength tensor is denoted by $G_a^{\mu\nu}$ and the strong coupling constant by α_s. The contributions of dimension-six operators are absorbed in the rescaling factor c_t for the top Yukawa coupling and the point-like coupling c_g. In other words, deviations of c_t and c_g from their Standard Model values $c_t = 1$ and $c_g = 0$ are understood to originate from dimension-six operators. The contribution of the chromomagnetic dipole operator [50, 51] is not included. The Lagrangian above includes QCD corrections via the δ parameter,

$$\delta = \delta_1 \frac{\alpha_s}{\pi} + \delta_2 \left(\frac{\alpha_s}{\pi}\right)^2 + \delta_3 \left(\frac{\alpha_s}{\pi}\right)^3, \tag{2}$$

[1] B. Fuks (ed.); F. Campanario, B. Chokoufe, R. Contino, F. Demartin, A. Falkowski, M. Ghezzi, A. Greljo, R. Gröber, C. Grojean, G. Isidori, W. Kilian, D. Marzocca, K. Mawatari, K. Mimasu, M. Mühlleitner, T. Ohl, G. Perez, M. Rauch, J. Reuter, F. Riva, R. Roth, V. Sanz, M. Sekulla, S. So Young, C. Speckner, M. Spira, J. Streicher, C. Weiss, M. Zaro, D. Zeppenfeld

whose three components read

$$\begin{aligned}\delta_1 &= \frac{11}{4}, \qquad \delta_2 = \frac{2777}{288} + \frac{19}{16}L_t + N_F\Big(\frac{L_t}{3} - \frac{67}{96}\Big),\\ \delta_3 =& \frac{897943}{9216}\zeta_3 - \frac{2761331}{41472} + \frac{209}{64}L_t^2 + \frac{2417}{288}L_t + N_F\left(\frac{58723}{20736} - \frac{110779}{13824}\zeta_3 + \frac{23}{32}L_t^2 + \frac{91}{54}L_t\right)\\ &+ N_F^2\left(-\frac{L_t^2}{18} + \frac{77}{1728}L_t - \frac{6865}{31104}\right),\end{aligned} \tag{3}$$

where $L_t = \log(\mu_R^2/m_t^2)$ with μ_R denoting the renormalization scale, N_F the number of active quark flavors and m_t the top quark pole mass.

The leading order (LO) cross section extended by the new physics contributions induced by the above effective Lagrangian is given by

$$\sigma_{LO}(pp \to H) = \sigma_0 \tau_H \frac{\mathrm{d}\mathcal{L}^{gg}}{\mathrm{d}\tau_H}, \tag{4}$$

where $\mathcal{L}^{gg}$ stands for the gluon-gluon partonic luminosity and where $\tau_H = m_h^2/S$, m_h being the Higgs-boson mass and S the hadronic center-of-mass energy. The σ_0 prefactor can be computed either in the non-linear case ($\sigma_0^{\rm NL}$) or in the SILH (strongly interacting light Higgs) framework ($\sigma_0^{\rm SILH}$),

$$\sigma_0^{\rm NL} = \frac{G_F\alpha_s^2}{288\sqrt{2}\pi}\left|\sum_Q c_Q A_{1/2}(\tau_Q) + 12c_g\right|^2 \qquad \text{and} \qquad \sigma_0^{\rm SILH} = \sigma_0^{\rm NL} - \frac{G_F\alpha_s^2}{288\sqrt{2}\pi}\left|\sum_Q (c_Q - 1)A_Q(\tau_Q) + 12c_g\right|^2, \tag{5}$$

where $\tau_Q = 4m_Q^2/m_h^2$ with m_Q being a generic quark mass, and where $G_F = (\sqrt{2}v^2)^{-1}$ stands for the Fermi constant. The LO form factors depend on the $A_{1/2}$ function that is given by

$$A_{1/2}(\tau) = \frac{3}{2}\tau\Big[1 + (1-\tau)f(\tau)\Big] \qquad \text{with} \qquad f(\tau) = \begin{cases} \arcsin^2\dfrac{1}{\sqrt{\tau}} & \text{for } \tau \geq 1\\ -\dfrac{1}{4}\left[\ln\dfrac{1+\sqrt{1-\tau}}{1-\sqrt{1-\tau}} - i\pi\right]^2 & \text{for } \tau < 1\,. \end{cases} \tag{6}$$

Rescaling factors c_Q have been introduced in Eq. (5) for all contributing quarks, *i.e.*, the top (c_t), bottom (c_b) and charm (c_c) quarks.

The Wilson coefficient c_g does not receive QCD corrections within the effective Lagrangian, but develops a scale dependence according to the renormalization group equation

$$\mu^2\frac{\partial c_g(\mu^2)}{\partial\mu^2} = -\left\{\beta_1\Big(\frac{\alpha_s(\mu^2)}{\pi}\Big)^2 + 2\beta_2\Big(\frac{\alpha_s(\mu^2)}{\pi}\Big)^3\right\}c_g(\mu^2), \tag{7}$$

with

$$\beta_0 = \frac{33 - 2N_F}{12}, \qquad \beta_1 = \frac{153 - 19N_F}{24} \qquad \text{and} \qquad \beta_2 = \frac{1}{128}\Big(2857 - \frac{5033}{9}N_F + \frac{325}{27}N_F^2\Big). \tag{8}$$

This renormalization group equation can be derived either from the scale-invariant trace anomaly term $\beta(\alpha_s)/\alpha_s G_a^{\mu\nu}G^a_{\mu\nu}$ [52–57], or from the scale dependence of the factor $(1+\delta)$ of the effective Lagrangian of Eq. (1), since both coefficients, $c_t(1+\delta)$ and c_g, have to develop the same scale dependence. The solution of the RGE for c_g up to the next-to-next-to-leading logarithmic (NNLL) level can be cast into the form

$$c_g(\mu^2) = c_g(\mu_0^2)\,\frac{\beta_0 + \beta_1\frac{\alpha_s(\mu^2)}{\pi} + \beta_2\Big(\frac{\alpha_s(\mu^2)}{\pi}\Big)^2}{\beta_0 + \beta_1\frac{\alpha_s(\mu_0^2)}{\pi} + \beta_2\Big(\frac{\alpha_s(\mu_0^2)}{\pi}\Big)^2}. \tag{9}$$

In order to compute the modified cross section up to NNLO, the mismatch of the individual terms of the effective Lagrangian of Eq. (1) with respect to the δ terms have been taken into account, the NNLL scale dependence of the Wilson coefficient c_g being properly included. In addition, the finite NLO quark mass term effects have been added at fixed NLO to the Standard Model contributions. This yields a consistent determination of the gluon-fusion cross section up to NNLO including the dimension-six operators that imply a rescaling of the top, bottom and charm Yukawa couplings and the point-like Hgg coupling parameterized by $c_g(\mu_R^2)$.

The present version 4.31 of HIGLU is linked to HDECAY (version 6.50) [58, 59] and allows one to choose the usual Standard Model Higgs input values in the separate input files `higlu.in` and `hdecay.in`. In addition, the rescaling

factors $c_{t,b,c}$ and the point-like Wilson coefficient $c_g(\mu_0^2)$ can be set, together with the corresponding input scale μ_0, in the `higlu.in` file. In this way, HIGLU provides a consistent calculation of the gluon-fusion cross section up to NNLO QCD including dimension-six operator effects. More detailed information about the input files `higlu.in` and `hdecay.in` can be found as comment lines at the beginning of the main FORTRAN files `higlu.f` and `hdecay.f` shipped with the program. Finally, we note that the pole masses of the bottom and charm quarks are computed from the $\overline{\rm MS}$ input values $\overline{m}_b(\overline{m}_b)$ and $\overline{m}_c(3\,\text{GeV})$ with N^{3}LO accuracy internally.

1.3 HPAIR: Higgs Boson Pair Production via Gluon Fusion

The program HPAIR, available at

http://tiger.web.psi.ch/hpair/,

calculates the Higgs boson pair production cross section via gluon fusion up to the NLO in QCD [60] in the limit of heavy top quarks within the Standard Model and the quark-loop induced contributions in the minimal supersymmetric extension of the Standard Model. While genuine supersymmetric QCD and electroweak corrections are unknown, known subleading NLO top mass effects [61–64] and NNLO QCD corrections [65–67] are not included in the program. Starting from the Standard Model Higgs result in the heavy top mass limit, the contributions of dimension-six operators beyond the Standard Model are included up to the NLO in QCD [68]. The latter extension is based on an effective phenomenological Lagrangian given by

$$\mathcal{L}_{\rm eff} = -m_t \bar{t}t\left[c_t\frac{H}{v}+c_{tt}\frac{H^2}{2v^2}\right] - \frac{1}{6}c_3\left[\frac{3m_h^2}{v}\right]H^3 + \frac{\alpha_s}{\pi}G_a^{\mu\nu}G^a_{\mu\nu}\left[c_g\frac{H}{v}+c_{gg}\frac{H^2}{2v^2}\right], \tag{10}$$

As in Section 1.2, the contributions of the dimension-six operators are absorbed in the rescaling factors c_t for the top Yukawa coupling and c_3 for the trilinear Higgs coupling, *i.e.*, deviations of c_t and c_3 from the Standard Model expectation of $c_t = c_3 = 1$ originate from dimension-six operators. The remaining couplings c_{tt}, c_g and c_{gg} are novel contributions purely arising from dimension-six operators and not present in the Standard Model Lagrangian. Integrating out the heavy top quark loops one arrives at the effective Lagrangian

$$\mathcal{L}_{\rm eff} = \frac{\alpha_s}{\pi}G_a^{\mu\nu}G^a_{\mu\nu}\left\{\frac{H}{v}\left[\frac{c_t}{12}\left(1+\frac{11}{4}\frac{\alpha_s}{\pi}\right)+c_g\right] + \frac{H^2}{v^2}\left[\frac{c_{tt}-c_t^2}{24}\left(1+\frac{11}{4}\frac{\alpha_s}{\pi}\right)+\frac{c_{gg}}{2}\right]\right\}, \tag{11}$$

describing the Higgs couplings to gluons when only the leading QCD corrections for the single Higgs interactions are kept (compared to Eq. (1)).

The partonic LO cross section extended by beyond the Standard Model contributions induced by the above Lagrangian is given by

$$\hat{\sigma}_{LO}(gg\to HH) = \int_{\hat{t}_-}^{\hat{t}_+}{\rm d}\hat{t}\,\frac{G_F^2\alpha_s^2(\mu_R)}{256(2\pi)^3}\left[|C_\Delta F_1+F_2|^2+|c_t^2 G_\Box|^2\right], \tag{12}$$

where the Mandelstam variables are defined by

$$\hat{s}=Q^2\,,\qquad \hat{t}=m_h^2-\frac{Q^2(1-\beta\cos\theta)}{2}\qquad\text{and}\qquad \hat{u}=m_h^2-\frac{Q^2(1+\beta\cos\theta)}{2}, \tag{13}$$

with $\beta=\sqrt{1-4m_h^2/Q^2}$ and with Q denoting the invariant mass of the Higgs boson pair. Moreover, the integration bounds at $\cos\theta=\pm1$ read

$$\hat{t}_\pm = m_h^2-\frac{Q^2(1\mp\beta)}{2}. \tag{14}$$

The form factors F_1 and F_2 can be cast into the form

$$F_1 = c_t F_\Delta + \frac{2}{3}c_\Delta \qquad\text{and}\qquad F_2 = c_t^2 F_\Box + c_{tt}F_\Delta - \frac{2}{3}c_\Box, \tag{15}$$

with the explicit expressions of F_Δ, $F_\Box$ and $G_\Box$ being available in Ref. [69, 70]. In the limit of heavy top quarks they simplify to

$$F_1^{\rm lim} = \frac{2}{3}\Big(c_t+c_\Delta\Big)\,,\qquad F_2^{\rm lim} = \frac{2}{3}\Big(-c_t^2+c_{tt}-c_\Box\Big), \tag{16}$$

after introducing the abbreviations c_Δ and $c_\Box$. The latter read, together with the C_Δ variable of Eq. (12),

$$C_\Delta \equiv \lambda_{hhh}\frac{m_Z^2}{Q^2-m_h^2+im_h\Gamma_h}\,,\qquad c_\Delta\equiv 12c_g\qquad\text{and}\qquad c_\Box\equiv -12c_{gg}, \tag{17}$$

with m_Z denoting the Z-boson mass, Γ_h the Higgs boson width, and the trilinear Higgs coupling being normalized as

$$\lambda_{hhh} = \frac{3m_h^2c_3}{m_Z^2}. \tag{18}$$

The partonic cross section has then to be convoluted with the gluon density in the proton in order to obtain the hadronic cross section.

The present version 2.00 of HPAIR needs the usual Standard Model Higgs input values that are provided in a separate input file `hpair.in`. In addition, this file also allows one for choosing the anomalous dimension-six factors c_t, c_{tt}, c_g, c_{gg} and c_3. Alternatively the dimension-six coefficients $\bar{c}_H$, $\bar{c}_u$, $\bar{c}_6$ and $\bar{c}_g$ can be provided within the usual SILH basis [71], starting from the effective SILH Lagrangian

$$\Delta\mathcal{L}_6^{\rm SILH} \supset \frac{\bar{c}_H}{2v^2}\partial_\mu(\Phi^\dagger\Phi)\partial^\mu(\Phi^\dagger\Phi) + \frac{\bar{c}_u}{v^2}y_t\Big(\Phi^\dagger\Phi\bar{q}_L\Phi^c t_R + \text{h.c.}\Big) - \frac{\bar{c}_6}{6v^2}\frac{3m_h^2}{v^2}\Big(\Phi^\dagger\Phi\Big)^3 + \bar{c}_g\frac{g_s^2}{m_W^2}\Phi^\dagger\Phi G^a_{\mu\nu}G_a^{\mu\nu}, \quad (19)$$

where $y_t = \sqrt{2}m_t/v$ denotes the top quark Yukawa coupling, g_s the strong coupling constant and m_W the W-boson mass. The connection to the Lagrangian of Eq. (10) is given by

$$c_t = 1 - \frac{\bar{c}_H}{2} - \bar{c}_u\,, \qquad c_{tt} = -\frac{1}{2}(\bar{c}_H + 3\bar{c}_u)\,, \qquad c_3 = 1 - \frac{3}{2}\bar{c}_H + \bar{c}_6\,, \qquad c_g = c_{gg} = \bar{c}_g\frac{4\pi}{\alpha_2}, \quad (20)$$

with $\alpha_2 = \sqrt{2}G_F m_W^2/\pi$. Contrary to the strongly interacting non-linear case, in the SILH case products of dimension-six coefficients are not taken into account in accordance with a consistent expansion of the physical observable up to linear dimension-six terms. In the SILH case, this also requires the expansion of the LO cross section of Eq. (12) up to linear dimension-six terms. The NLO terms are treated accordingly in the HPAIR program that allows for both options, *i.e.*, the non-linear and SILH cases. In this way, HPAIR provides a consistent calculation of the gluon-fusion cross section up to NLO QCD including the effects of dimension-six operators. More detailed information about the input file `hpair.in` can be found as comment lines at the beginning of the main FORTRAN file `hpair.f` shipped with the program.

1.4 eHDECAY, Higgs Decays in the Effective Lagrangian Approach

In Ref. [72], we have presented the program eHDECAY, a FORTRAN code based on a modification of the program HDECAY [58, 59], and discussed the relation between the non-linear and the linear effective Lagrangian approaches. The program can be downloaded from

http://www.itp.kit.edu/~maggie/eHDECAY/.

In eHDECAY, the full list of leading bosonic operators of the Higgs effective Lagrangian has been implemented, both for a linear and a non-linear realization of the electroweak symmetry and for two benchmark composite Higgs models called MCHM4 [73] and MCHM5 [74]. All the relevant QCD corrections have been included and we detail in the following the importance of higher-order QCD corrections and of mass effects on the corrections. We also show how to consistently include electroweak corrections whenever it is possible, focusing on the example of the Higgs boson decay into two gluons. As the leading part of the QCD corrections in general factorizes with respect to the expansion in the number of fields and derivatives of the effective Lagrangian, they can be included by taking over the results from the Standard Model. The electroweak corrections, on the contrary, require dedicated computations that are only partly available at present. They are consequently only implemented up to higher orders in a v^2/f^2 expansion in the framework of the linear Lagrangian, where $f = \Lambda/g_\star$ with $g_\star$ being a typical new physics coupling strength. They are valid, for instance, in the SILH framework [71] in which the deviations from the Standard Model are small. On different lines, higher-order QCD corrections and mass effects are included also in the non-linear implementation.

Denoting by c_ψ the modifications of the Higgs couplings to fermions with respect to the Standard Model and by c'_{gg} the effective Higgs coupling to gluons, the related non-linear effective Lagrangian reads

$$\Delta\mathcal{L}_{\rm NL} = -\sum_{\psi=u,d,\ell} c_\psi m_\psi\bar{\psi}\psi\frac{h}{v} + \frac{c'_{gg}}{2}G^a_{\mu\nu}G_a^{\mu\nu}\frac{h}{v}\,. \quad (21)$$

The decay rate into gluons implemented in eHDECAY in the framework of the non-linear Lagrangian is then given by

$$\begin{aligned}\Gamma_{gg}^{\rm NL} =& \frac{G_F\alpha_s^2 m_h^3}{4\sqrt{2}\pi^3}\Bigg[\Bigg|\sum_{q=t,b,c}\frac{c_q}{3}A_{1/2}(\tau_q)\Bigg|^2 c_{\rm eff}^2\,\kappa_{\rm soft} + 2\,\text{Re}\Bigg\{\sum_{q=t,b,c}\frac{c_q}{3}A^*_{1/2}(\tau_q)\frac{2\pi c_{gg}}{\alpha_s}\Bigg\}c_{\rm eff}\,\kappa_{\rm soft} + \Bigg|\frac{2\pi c_{gg}}{\alpha_s}\Bigg|^2\kappa_{\rm soft} \\ &+ \frac{1}{9}\sum_{q,q'=t,b} c_q\,A^*_{1/2}(\tau_q)c_{q'}\,A_{1/2}(\tau_{q'})\kappa^{\rm NLO}(\tau_q,\tau_{q'})\Bigg]\,.\end{aligned} \quad (22)$$

On the other hand, the part of the linear Lagrangian that contributes to the Higgs boson decay into two gluons reads, in the SILH basis,

$$\Delta\mathcal{L}_6^{\rm SILH} \supset \frac{\bar{c}_H}{2v^2}\partial^\mu(\Phi^\dagger\Phi)\partial_\mu(\Phi^\dagger\Phi) + \frac{\bar{c}_g g_s^2}{m_W^2}\Phi^\dagger\Phi G^a_{\mu\nu}G_a^{\mu\nu} + \Big(\frac{\bar{c}_u y_u}{v^2}\Phi^\dagger\Phi\bar{q}_L\Phi^c u_R + \frac{\bar{c}_d y_d}{v^2}\Phi^\dagger\Phi\bar{q}_L\Phi d_R + \text{h.c.}\Big)\,, \quad (23)$$

where $\bar{c}_H$, $\bar{c}_g$, $\bar{c}_u$ and $\bar{c}_d$ are the Wilson coefficients of the corresponding dimension-six effective operators and $y_{u,d}$ the Yukawa couplings of the up- and down-type quarks, respectively. In this case, the Higgs boson decay rate into gluons implemented in eHDECAY is given by

$$\begin{aligned}\Gamma_{gg}^{\rm SILH} =& \frac{G_F \alpha_s^2 m_h^3}{4\sqrt{2}\pi^3}\left[\frac{1}{9}\sum_{q,q'=t,b,c}(1-\bar{c}_H-\bar{c}_q-\bar{c}_{q'})A_{1/2}^*(\tau_{q'})A_{1/2}(\tau_q)c_{\rm eff}^2\,\kappa_{\rm soft}\right.\\ &+2\,{\rm Re}\left\{\sum_{q=t,b,c}\frac{1}{3}A_{1/2}^*(\tau_q)\frac{16\pi\bar{c}_g}{\alpha_2}\right\}c_{\rm eff}\,\kappa_{\rm soft}+\left|\sum_{q=t,b,c}\frac{1}{3}A_{1/2}(\tau_q)\right|^2 c_{\rm eff}^2\,\kappa_{\rm ew}\,\kappa_{\rm soft}\\ &\left.+\frac{1}{9}\sum_{q,q'=t,b}(1-\bar{c}_H-\bar{c}_q-\bar{c}_{q'})A_{1/2}^*(\tau_q)A_{1/2}(\tau_{q'})\kappa^{\rm NLO}(\tau_q,\tau_{q'})\right].\end{aligned} \tag{24}$$

In both Eq. (22) and Eq. (24), $\tau_q = 4m_q^2/m_h^2$ and the loop function $A_{1/2}(\tau)$ is defined as in Eq. (6). We use the pole masses for the top, bottom and charm quark masses and α_s is computed up to N^{3}LO at a scale fixed to m_h and for $N_F = 5$ active flavors. The QCD corrections have been taken into account up to N^{3}LO in QCD in the limit of heavy loop-particle masses so that the effect from low-energy gluon radiation, given by the coefficient $\kappa_{\rm soft}$, factorizes in this limit. The corrections from high-energy gluon and quark exchange (with virtuality $q^2 \gg m_t^2$) are encoded in the coefficient $c_{\rm eff}$. Namely, for $m_h \ll 2m_t$, the top quark can be integrated out, leading to the effective five-flavor Lagrangian

$$\mathcal{L}_{\rm eff} = -2^{1/4} G_F^{1/2} C_1 G^a_{\mu\nu} G_a^{\mu\nu} h\,, \tag{25}$$

where the dependence on the top quark mass m_t is included in the coefficient function C_1. We then have

$$\kappa_{\rm soft} = \frac{\pi}{2m_h^4}\,{\rm Im}\left\{\Pi^{GG}(q^2=m_h^2)\right\} \qquad \text{and} \qquad c_{\rm eff} = -\frac{12\pi\, C_1}{\alpha_s(m_h)}\,, \tag{26}$$

with the vacuum polarization $\Pi^{GG}(q^2)$ being induced by the gluon operator. The N^{3}LO expressions for C_1 and for ${\rm Im}\{\Pi^{GG}\}$ have been given in Refs. [47, 48, 75, 76] and Ref. [77], respectively. In particular, the NLO expressions for $\kappa_{\rm soft}$ and $c_{\rm eff}$ read [5, 78, 79]

$$\kappa_{\rm soft}^{\rm NLO} = 1 + \frac{\alpha_s}{\pi}\left(\frac{73}{4} - \frac{7}{6}N_F\right) \qquad \text{and} \qquad c_{\rm eff}^{\rm NLO} = 1 + \frac{\alpha_s}{\pi}\frac{11}{4}\,, \tag{27}$$

in agreement with the low-energy theorem [80–82]. The additional mass effects at NLO [8] in the top and bottom quark loops are taken into account by the function $\kappa^{\rm NLO}(\tau_q,\tau_{q'})$ in the last lines of Eqs. (22) and (24). This function quantifies the difference between the NLO QCD corrections for the top (bottom) contribution taking into account finite mass effects in the loop, and the result for the top (bottom) contribution in the limit of a large loop-particle mass.

Higher order corrections are large, with in particular the N^{3}LO QCD corrections increasing the total width by almost up to 90%. The mass effects at NLO QCD are relevant for the bottom loop where they amount to 8%, while they are negligible for the top loop. To be consistent with the non-linear approach, no electroweak corrections have been included in the non-linear parameterization, as a perturbative expansion in powers of v^2/f^2 is not possible in the general case. In contrast, in the linear parameterization, we have included both QCD and electroweak corrections. The treatment of the QCD corrections is in accordance with the non-linear case, while the electroweak corrections [16–23] are taken into account by the factor $\kappa_{\rm ew}$ of Eq. (24). In the Standard Model, they can be factorized in the particular case of the Higgs boson decaying into two gluons. However, this factorization is not valid in the general case. In particular, when the contributions from the effective Lagrangian are taken into account, modified Higgs couplings absent at the leading order can appear and spoil the factorization. Hence, a consistent inclusion of the electroweak corrections is only possible in the linear-Lagrangian case and up to higher orders in v^2/f^2. More precisely, the Standard Model electroweak corrections can be added to get a result that includes the leading $\mathcal{O}(v^2/f^2)$ and the next-to-leading $\mathcal{O}(\alpha/4\pi)$ corrections as well as the mixed $\mathcal{O}[(\alpha_s/4\pi)^5(\alpha/4\pi)]$ contributions, but that neglects terms of $\mathcal{O}[(\alpha/4\pi)(v^2/f^2)]$, $\mathcal{O}[((v^2/f^2)^2]$ and $\mathcal{O}[(\alpha/4\pi)^2]$, where α is generically meant as the electroweak expansion parameter.

For the future we are planning to add to eHDECAY contributions from the fermionic effective operators and those from the effective Lagrangian to the set of input observables m_W, m_Z and G_F. Moreover, in accordance with the Higgs Cross Section Working Group Internal Note of Ref. [83], a switch from the pole-mass scheme to the $\overline{\rm MS}$ scheme for the fermionic masses is under consideration, as well as the implementation of an interface for the so-called 'Higgs basis' (Ref. to the WG2 chapter) parameterization of the effective Lagrangian (see also Section 1.9).

1.5 Implementation of Higgs Pseudo-Observables in the Universal FEYNRULES Output

With the upcoming LHC Run-II, Higgs physics is entering a precision era. This would allow us to look for new physics effects not only in the overall signal strengths, but also in kinematical distributions. In this perspective, the so called 'κ-framework' [84] is insufficient and needs to be extended. The natural extension is the so-called Higgs pseudo-observables

Process	$h \to bb$	$h \to \tau\tau$	$h \to cc$	$h \to \mu\mu$
Pseudo-observables	$\kappa_b,\ \lambda_b^{\rm CP}$	$\kappa_\tau,\ \lambda_\tau^{\rm CP}$	$\kappa_c,\ \lambda_c^{\rm CP}$	$\kappa_\mu,\ \lambda_\mu^{\rm CP}$

Table 1: Pseudo-observables relevant for Higgs decays into two fermions.

Process	Pseudo-observables
$h \to \gamma\gamma$	$\kappa_{\gamma\gamma},\ \lambda_{\gamma\gamma}^{\rm CP}$
$h \to Z\gamma$	$\kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP}$
$h \to \gamma 2\nu$	$\kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP}$
$h \to \gamma 2\ell$	$\kappa_{\gamma\gamma},\ \lambda_{\gamma\gamma}^{\rm CP},\ \kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP}$
$h \to Z2\ell$	$\kappa_{ZZ},\ \epsilon_{ZZ},\ \epsilon_{ZZ}^{\rm CP},\ \kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP},\ \epsilon_{Z\ell_L},\ \epsilon_{Z\ell_R}$
$h \to 2\ell 2\ell'$	$\kappa_{ZZ},\ \epsilon_{ZZ},\ \epsilon_{ZZ}^{\rm CP},\ \kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP},\ \kappa_{\gamma\gamma},\ \lambda_{\gamma\gamma}^{\rm CP},\ \epsilon_{Z\ell_L},\ \epsilon_{Z\ell_R},\ \epsilon_{Z\ell'_L},\ \epsilon_{Z\ell'_R}$
$h \to 4\ell$	$\kappa_{ZZ},\ \epsilon_{ZZ},\ \epsilon_{ZZ}^{\rm CP},\ \kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP},\ \kappa_{\gamma\gamma},\ \lambda_{\gamma\gamma}^{\rm CP},\ \epsilon_{Z\ell_L},\ \epsilon_{Z\ell_R}$
$h \to \bar{\ell}\ell 2\nu$	$\begin{cases} \kappa_{ZZ},\ \epsilon_{ZZ},\ \epsilon_{ZZ}^{\rm CP},\ \kappa_{Z\gamma},\ \lambda_{Z\gamma}^{\rm CP},\ \epsilon_{Z\ell_L},\ \epsilon_{Z\ell_R},\ \epsilon_{Z\nu} \\ \kappa_{WW},\ \epsilon_{WW},\ \epsilon_{WW}^{\rm CP},\ \epsilon_{W\ell},\ \phi_{W\ell} \end{cases}$
$h \to \bar{\ell}\ell' 2\nu$	$\kappa_{WW},\ \epsilon_{WW},\ \epsilon_{WW}^{\rm CP},\ \epsilon_{W\ell},\ \phi_{W\ell},\ \epsilon_{W\ell'},\ \phi_{W\ell'}$

Table 2: Pseudo-observables relevant for Higgs decays into a gauge-boson pair, and into a three-body or four-body system. We denote by ℓ an electron, a muon or a tau and ν indicates any of the three neutrino species.

framework introduced in Refs. [85, 86]. A detailed discussion about this formalism is presented elsewhere in this Yellow Report 4. Here we limit ourself to summarize its main features in order to illustrate its implementation via the Universal FeynRules Output model HiggsPO that is available at

http://www.physik.uzh.ch/data/HiggsPO.

The pseudo-observables are a finite set of parameters that are experimentally accessible, well-defined from the point of view of a quantum field theory, and that characterize possible deviations from the Standard Model in processes involving the Higgs boson in great generality. More precisely, the Higgs pseudo-observables are defined from a general decomposition (based on analyticity, unitarity, and crossing symmetry) of on-shell amplitudes involving the Higgs boson and a momentum expansion following the assumption of no new light particles in the kinematical regime where the decomposition is assumed to be valid. A further key assumption of the pseudo-observables formalism is that the Higgs boson is a spin zero resonance with a narrow width, such that new physics effects in production and decay factorize.

Due to their simple kinematics, two-body Higgs decays (into a $\gamma\gamma$ and an $\bar{f}f$ pair) can be parameterized by only two pseudo-observables each, one describing the CP-conserving amplitude and another for the CP-violating one. Unless the polarization of the final states can be measured, only a combination of the two is experimentally accessible. Three-body and four-body Higgs decays (such as into $2\ell\gamma$, 4ℓ and $2\ell 2\nu$ systems) have a more complicated kinematics. In this case the pseudo-observables are defined from the expansion of the on-shell amplitudes around the known physical poles due to the propagation of intermediate Standard Model electroweak gauge bosons. General amplitude decomposition and pseudo-observables definitions can be found in Ref. [85] for the Higgs decays and in Ref. [86] for Higgs boson production in the vector-boson fusion mode and in association with a gauge boson.

The pseudo-observables are defined directly at the amplitude level. As such, on the one hand they can be computed from a specific Lagrangian (to a specific order in perturbation theory), and on the other hand the pseudo-observables are directly connected to S-matrix elements, providing a direct link to physical observables [85]. They are thus particularly well suited for the analysis of experimental data with the matrix-element method [87]. In order to use the pseudo-observable decomposition for precision studies, it is important to account for the long-distance contributions due to soft and collinear photon emission (*i.e.*, the leading QED radiative corrections). These represent a universal correction factor that can be implemented, by means of appropriate convolution functions or, equivalently, by showering algorithms in Monte Carlo simulations, irrespective of the specific short-distance structure of the amplitude. It has been shown that inclusion of such correction in $h \to 2e2\mu$ decays recovers the complete NLO Standard Model predictions within an accuracy of about 1% [88].

In order to use the pseudo-observable framework in experimental analyses, it is convenient to have a tool capable of generating signal events for specific values of the pseudo-observables. Such a tool has been developed in the context of the Higgs decays and is publicly available on the above webpage that also includes a detailed user manual. The imple-

mentation of pseudo-observables in electroweak Higgs boson production including NLO QCD corrections is underway and will be available at the same Internet link. We have implemented into FEYNRULES [89] (version 2.3.1) a HIGGSPO model, whose source file is named `HPO.fr`, by means of an effective Lagrangian generating the corresponding vertices at the tree level. We stress that such a Lagrangian should not be considered as a specific model or an effective field theory description, but rather as an auxiliary tool to be used (at tree level only) in order to reproduce the correct Higgs decay amplitude decomposition in terms of pseudo-observables that is valid beyond the tree level.

The HIGGSPO FEYNRULES model as implemented in the `HPO.fr` file is exported to the Universal FeynRules Output (UFO) format [90], which can then be used within the MADGRAPH5_aMC@NLO [91] and SHERPA [92] event generators. The general idea is that after simulating Higgs boson production with a dedicated Monte Carlo generator, the Higgs boson can be decayed at the parton-level with the HIGGSPO model and the partonic events can then be passed to a general purpose event generator for subsequent parton showering and hadronization (such as PYTHIA [93]). This last step will automatically account for the important radiative corrections [88]. We note that it is very practical to use the MADSPIN module [94] of MADGRAPH5_aMC@NLO to decay the Higgs boson with the HIGGSPO model on the fly. We stress again that our FEYNRULES implementation only consists of a set of effective interactions that generate exactly the scattering amplitude of interest at the tree level and is supposed to be used for this purpose only. It should not be used as a Lagrangian for arbitrary process and beyond the tree level.

Here we summarize all the Higgs decay processes implemented in the HIGGSPO model, together with the associated Higgs pseudo-observables accessible in the parameter card. The Higgs pseudo-observables relevant for describing the Higgs decays into two fermions are shown in Table 1. The coupling strengths have been assigned an interaction order `YUK` $= 1$. The pseudo-observables relevant for all other Higgs decays are presented in Table 2 and the coupling strengths have been related to an interaction order `HPO` $= 1$. Translation of the pseudo-observable language to the Higgs basis defined in Chapter (Ref. to the WG2 chapter) is also available in Refs. [85, 86].

1.6 Higgs and Beyond the Standard Model Characterisation in the MADGRAPH5_aMC@NLO Framework

The HIGGS CHARACTERISATION (HC) implementation provides a complete framework, based on an effective field theory description, that allows for the study of the Higgs boson properties in a consistent, systematic and accurate way. The HC [95] follows the general strategy outlined in Ref. [96], and has been implemented in a complete simulation chain from Lagrangian to hadron-level events, especially in the FEYNRULES/MADGRAPH5_aMC@NLO framework. The HC effective Lagrangian features bosons $X(J^P)$ with various spin-parity assignments ($J^P = 0^+$, 0^-, 1^+, 1^- and 2^+) and has been implemented in terms of the Standard Model mass eigenstates into FEYNRULES [89], whose output files are interfaced with various event generators [90, 97]. All model files are publicly available online,

http://feynrules.irmp.ucl.ac.be/wiki/HiggsCharacterisation.

Later on, the HC framework has been extended as a BEYOND THE STANDARD MODEL CHARACTERISATION (BSMC) framework [98] where the effective Lagrangian has been constructed as above, but starting from the Higgs basis Lagrangian (Ref. to the WG2 chapter) where all possible effective operators of dimension up to six are written in terms of the Standard Model mass eigenstates. In particular, fermionic operators different from the Yukawa interactions are now included. However, instead of imposing the realization of the electroweak symmetry that relates the Wilson coefficients, the latter have been kept independent. Due to the lack of manifest $SU(2)_L \times U(1)_Y$ invariance, the BSMC Lagrangian is associated with a larger number of independent coefficients compared to more traditional bases such as the Warsaw [99] or SILH [71, 100] bases. As for the HC case, the BSMC Lagrangian has been implemented into FEYNRULES and is available online,

http://feynrules.irmp.ucl.ac.be/wiki/BSMCharacterisation.

A corresponding implementation of the dimension-six Lagrangian above the weak scale, where $SU(2)_L \times U(1)_Y$ is an exact symmetry, has been also achieved [101] and has overlapping as well as complementary features with respect to the HC and BSMC Lagrangians. It is available at

http://feynrules.irmp.ucl.ac.be/wiki/HEL.

There are several advantages in having a first principle implementation in terms of an effective Lagrangian which can be automatically interfaced to event generators. First and most important, all relevant production and decay modes can be studied within the same model, and the corresponding processes automatically generated within minutes. Second, it is straightforward to modify the model implementation to extend it further in case of need, by adding further interactions, for example of higher-dimensions in energy. Finally, higher-order effects can be easily accounted for, by generating multi-jet merged samples or computing NLO corrections within automatic frameworks. As accumulated data has been increasing, the last point became more and more important. The detailed demonstrations and analyses for all the main production modes of the Higgs boson (gluon fusion, vector boson fusion, Vh associated production and $t\bar{t}h$ production) as well as th production at NLO accuracy in QCD have been done recently [102–104], within the MADGRAPH5_aMC@NLO program [91]. In the following, we focus, for the sake of the example, on a spin-0 (Higgs) boson that is denoted by X_0, and compute several differential distributions for some production processes and for different benchmark scenarios.

Scenario	HC parameter choice	Scenario	HC parameter choice
0^+(GF, SM)	$\kappa_{Hgg/Htt} = 1$ $(c_\alpha = 1)$	0^+(VBF,SM)	$\kappa_{SM} = 1$ $(c_\alpha = 1)$
0^-(GF)	$\kappa_{Agg/Att} = 1$ $(c_\alpha = 0)$	0^+(VBF,HD)	$\kappa_{HZZ,HWW} = 1$ $(c_\alpha = 1)$
$0^\pm$(GF)	$\kappa_{Hgg,Agg/Htt,Att} = 1$ $(c_\alpha = 1/\sqrt{2})$	0^-(VBF,HD)	$\kappa_{AZZ,AWW} = 1$ $(c_\alpha = 0)$
		$0^\pm$(VBF,HD)	$\kappa_{HZZ,HWW,AZZ,AWW} = 1$ $(c_\alpha = 1/\sqrt{2})$

Table 3: Benchmark scenarios for X_0 production in the gluon fusion and $t\bar{t}X_0$ (GF) and in the vector boson fusion (VBF) channel.

In the HC framework, this means that the only assumptions are that the 125 GeV resonance found at the LHC corresponds to a spin-0 state, that no other new state coupled to such a resonance exists below the cutoff scale Λ and that new physics is dominantly described by the lowest dimensional operators. We thus include all effects steming from the complete set of dimension-six operators allowed by the Standard Model gauge symmetry. The effective interaction Lagrangian is given by

$$\begin{aligned}\mathcal{L}_0 =& \Big\{ -\sum_{f=t,b,\tau} \bar{\psi}_f \big(c_\alpha \kappa_{Htt} g_{Htt} + i s_\alpha \kappa_{Att} g_{Att}\, \gamma_5\big)\psi_f + c_\alpha \kappa_{\rm SM} \big[\frac{1}{2} g_{HZZ}\, Z_\mu Z^\mu + g_{HWW}\, W^+_\mu W^{-\mu}\big] \\ & - \frac{1}{4}\big[c_\alpha \kappa_{H\gamma\gamma} g_{H\gamma\gamma}\, A_{\mu\nu}A^{\mu\nu} + s_\alpha \kappa_{A\gamma\gamma} g_{A\gamma\gamma}\, A_{\mu\nu}\widetilde{A}^{\mu\nu}\big] - \frac{1}{2}\big[c_\alpha \kappa_{HZ\gamma} g_{HZ\gamma}\, Z_{\mu\nu}A^{\mu\nu} + s_\alpha \kappa_{AZ\gamma} g_{AZ\gamma}\, Z_{\mu\nu}\widetilde{A}^{\mu\nu}\big] \\ & - \frac{1}{4}\big[c_\alpha \kappa_{Hgg} g_{Hgg}\, G^a_{\mu\nu}G^{a,\mu\nu} + s_\alpha \kappa_{Agg} g_{Agg}\, G^a_{\mu\nu}\widetilde{G}^{a,\mu\nu}\big] - \frac{1}{4}\frac{1}{\Lambda}\big[c_\alpha \kappa_{HZZ}\, Z_{\mu\nu}Z^{\mu\nu} + s_\alpha \kappa_{AZZ}\, Z_{\mu\nu}\widetilde{Z}^{\mu\nu}\big] \\ & - \frac{1}{2}\frac{1}{\Lambda}\big[c_\alpha \kappa_{HWW}\, W^+_{\mu\nu}W^{-\mu\nu} + s_\alpha \kappa_{AWW}\, W^+_{\mu\nu}\widetilde{W}^{-\mu\nu}\big] \\ & - \frac{1}{\Lambda} c_\alpha \big[\kappa_{H\partial\gamma}\, Z_\nu \partial_\mu A^{\mu\nu} + \kappa_{H\partial Z}\, Z_\nu \partial_\mu Z^{\mu\nu} + \kappa_{H\partial W}\big(W^+_\nu \partial_\mu W^{-\mu\nu} + h.c.\big)\big]\Big\}\, X_0\,, \end{aligned} \tag{28}$$

where $A_{\mu\nu}$ ($\widetilde{A}_{\mu\nu}$), $Z_{\mu\nu}$ ($\widetilde{Z}_{\mu\nu}$) and $W^\pm_{\mu\nu}$ ($\widetilde{W}^\pm_{\mu\nu}$) represent the reduced field strength tensors of the physical electroweak bosons, and $G^a_{\mu\nu}$ and $\widetilde{G}^a_{\mu\nu}$ the gluon field strength tensor and its dual. This parameterization allows for the easy recovery of the Standard Model case by fixing the dimensionless parameters κ_i to appropriate values, since the dimensionful couplings $g_{Xyy'}$ are set to their Standard Model values. For instance, we enforce that $g_{H\gamma\gamma} = 47\alpha/18\pi$ and $g_{Agg} = \alpha_s/2\pi$. Moreover, the model description includes the interactions of a 0^- state typical of supersymmetric or of generic theories with two Higgs doublets, and enables the CP-mixing between the 0^+ and 0^- states via a mixing angle α whose sine and cosine are denoted by s_α and c_α. In the Standard Model, $c_\alpha = 1$ and $s_\alpha = 0$.

We then make use of MADGRAPH5_AMC@NLO to generate both a numerical code and events, at the NLO accuracy in QCD, for the production of an X_0 state plus two jets in the gluon fusion channel automatically. This is achieved by issuing the following commands (with the `/ t` syntax forbidding diagrams containing top loops),

```
> import model HC_NLO_X0-heft
> generate p p > x0 j j / t [QCD]
> output
> launch
```

where the `-heft` suffix in the model name refers to the corresponding model restriction. As a result, all the amplitudes featuring the Higgs–gluon effective vertices in the heavy-top limit are generated, including corrections up to the NLO in QCD. Analogous commands can be issued to generate events related to the production of an X_0 state plus zero and one jet. After the `launch` command, one can modify the `param_card.dat` file to change the values of the Λ, κ and c_α parameters. Similarly, a code and events can be generated in the vector boson fusion case (left, with the $$ sign forbidding diagrams with $W^\pm$ or Z bosons in the s-channel as they are included in Vh production) and in the VX_0 associated production mode (right) by typing in

```
> import model HC_NLO_X0
> generate p p > x0 j j $$ w+ w- z / a [QCD]
```

```
> generate p p > x0 l+ vl [QCD]
> add process p p > x0 l- vl~ [QCD]
> add process p p > x0 l+ l- / a [QCD]
```

as well as in the $t\bar{t}X_0$ production mode with the command

```
> generate p p > x0 t t~ [QCD]
```

Furthermore, the X_0 and the top quark decays are subsequently performed starting from the event files generated as above with the help of the MADSPIN [94] package, following the procedure described in Ref. [105] and that allows one to keep spin correlations.

We show in Figure 1 (left and central panels) a few kinematical distributions, including theoretical uncertainties, for the benchmark scenarios presented in Table 3. The parton densities are evaluated by using the NNPDF2.3 (LO/NLO)

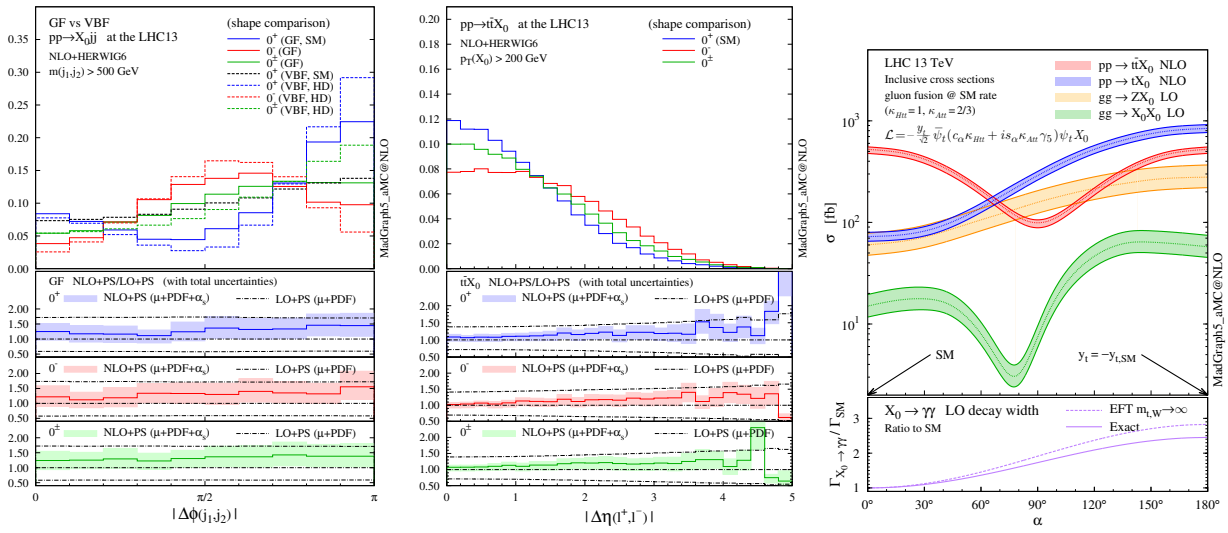

Fig. 1: Normalized kinematical distributions at the LHC, running at a center-of-mass energy of 13 TeV. We show the azimuthal difference between the two tagging jets for $pp \to X_0 jj$ after imposing a $m_{jj} > 500$ GeV selection (left) and the opening angle between the leptons (center) for the dileptonic decay channel in $pp \to t\bar{t}X_0$ after enforcing a $p_T(X_0) > 200$ GeV selection. For each scenario, the bottom panels give the ratio of the NLO results matched to parton showers to the LO results matched to parton showers, together with the total theoretical uncertainties. In the right panel of the figure, we show NLO (or loop-induced LO) cross sections, presented with the associated scale uncertainties, for $t\bar{t}X_0$ and t-channel tX_0 (ZX_0 and X_0X_0) production as a function of the CP-mixing angle α. The κ_{Htt} and κ_{Att} parameters have been set to reproduce the Standard Model gluon-fusion cross section for every value of α. The ratio of the $X_0 \to \gamma\gamma$ partial decay width to the corresponding Standard Model value is also shown in the lower panel of the figure.

parameterization [106] and the central value μ_0 for the renormalization (μ_R) and factorisation (μ_F) scales is set to $H_T/2$, m_W and $\sqrt[3]{m_T(t)\, m_T(\bar{t})\, m_T(X_0)}$ in the gluon fusion, vector boson fusion and $t\bar{t}X_0$ production channel respectively. Uncertainties coming from missing higher orders are estimated by varying μ_R and μ_F independently by a factor of two up and down around μ_0. Finally, in the right panel of the figure, we present the dependence of the th and $t\bar{t}X_0$ production cross sections on the CP-mixing angle α. The nature of the top quark Yukawa coupling also affects the loop-induced Higgs coupling to gluons and photons. In oreder to maintain the Standard Model gluon-fusion production cross section, the rescaling parameters are set to $\kappa_{Htt} = 1$ and $\kappa_{Att} = 2/3$. The LO cross sections for the loop-induced HZ [107] and HH [108] production processes via gluon fusion are also shown as references.

1.7 Anomalous Couplings in VBFNLO

Next-to-leading order QCD predictions including anomalous couplings effects can be studied for several processes with the flexible Monte Carlo program VBFNLO [109, 110], available at

https://www.itp.kit.edu/vbfnlo.

The ensemble of implemented processes includes Higgs, single and double vector boson production via vector boson fusion (VBF), WH production, as well as double and triple vector boson (plus jet) production. The Standard Model QCD-induced background for double vector boson production in association with two jets is also available at the NLO accuracy. Furthermore, anomalous HVV coupling effects are also included in the gluon-induced contributions to diboson (plus jet) production as well as in the gluon fusion processes $gg \to Hjj \to VVjj$. Although these processes are all one-loop induced and hence computed at the leading-order accuracy, the full top- and bottom-quark mass dependence is retained.

In the VBFNLO-3.0 β release, an interface compliant with the Binoth Les Houches Accord (BLHA) [111, 112] has been added for all VBF processes including fully leptonic decays, which allows for Monte Carlo studies at NLO in QCD including the full functionality of event generators like HERWIG 7 [113, 114]. The K-matrix unitarization procedure has been implemented for the two dimension-eight operators $\mathcal{O}_{S,0} + \mathcal{O}_{S,2}$ and $\mathcal{O}_{S,1}$ that are defined in Eq. (32) below. The strength of the anomalous triple and quartic gauge boson couplings can be set in the file `anomV.dat`. They are parameterized using an effective Lagrangian, as described in Refs. [115–118],

$$\mathcal{L}_{\text{eff}} = \sum_i \frac{f_i}{\Lambda^n} \mathcal{O}_i^{n+4} \,, \tag{29}$$

where $n + 4$ signifies the dimension of the operator $\mathcal{O}_i$. VBFNLO then defines anomalous gauge couplings in terms of

the coefficients f_i/Λ^n of the dimension-six and dimension-eight operators. The full list of implemented operators can be found in the Appendix 1 of the VBFNLO manual [119]. The explicit form of the included $\mathcal{CP}$-even dimension-six operators is given by

$$\begin{aligned} \mathcal{O}_W &= (D_\mu\Phi)^\dagger \widehat{W}^{\mu\nu}(D_\nu\Phi), \qquad \mathcal{O}_B = (D_\mu\Phi)^\dagger \widehat{B}^{\mu\nu}(D_\nu\Phi), \qquad \mathcal{O}_{WWW} = \mathrm{Tr}\Big[\widehat{W}_{\mu\nu}\widehat{W}^{\nu\rho}\widehat{W}_\rho{}^{\mu}\Big], \\ \mathcal{O}_{WW} &= \Phi^\dagger \widehat{W}_{\mu\nu}\widehat{W}^{\mu\nu}\Phi, \qquad \mathcal{O}_{BB} = \Phi^\dagger \widehat{B}_{\mu\nu}\widehat{B}^{\mu\nu}\Phi. \end{aligned} \tag{30}$$

The building blocks for these operators (following the notation of Refs. [116, 120]) are defined by

$$\widehat{W}_{\mu\nu} = igT_a W^a_{\mu\nu}, \qquad \widehat{B}_{\mu\nu} = ig'Y B_{\mu\nu}, \qquad D_\mu = \partial_\mu + igT_a W^a_\mu + ig'Y B_\mu, \tag{31}$$

where g and g' are the $SU(2)_L$ and $U(1)_Y$ gauge couplings, and T_a and Y the generators of the $SU(2)$ group in the fundamental representation and the hypercharge operator, respectively. The $\mathcal{CP}$-odd part of the Lagrangian is obtained replacing the field strength tensor with the corresponding dual field strength tensors. The dimension-eight operators that are supported are taken from Ref. [117] although slightly different normalizations for the field strength tensors are employed in this work, $\widehat{W}_{\mu\nu} = T_a W^a_{\mu\nu}$ and $\widehat{B}_{\mu\nu} = B_{\mu\nu}$. The conversion factors for the coupling strengths f_i relating Ref. [117] to our implementation can be found in the Appendix 1 of the VBFNLO manual. The dimension-eight operators that are supported can be split into three categories, namely operators depending on the gauge-covariant derivative $D_\mu\Phi$ only,

$$\begin{aligned} \mathcal{O}_{S,0} &= \Big[(D_\mu\Phi)^\dagger D_\nu\Phi\Big]\Big[(D^\mu\Phi)^\dagger D^\nu\Phi\Big], \qquad \mathcal{O}_{S,1} = \Big[(D_\mu\Phi)^\dagger D^\mu\Phi\Big]\Big[(D_\nu\Phi)^\dagger D^\nu\Phi\Big], \\ \mathcal{O}_{S,2} &= \Big[(D_\mu\Phi)^\dagger D_\nu\Phi\Big]\Big[(D^\nu\Phi)^\dagger D^\mu\Phi\Big], \end{aligned} \tag{32}$$

operators depending on both $D_\mu\Phi$ and the electroweak field strength tensors $\widehat{W}_{\mu\nu}$ and $\widehat{B}_{\mu\nu}$, *e.g.*,

$$\mathcal{L}_{M,0} = \mathrm{Tr}\Big[\widehat{W}_{\mu\nu}\widehat{W}^{\mu\nu}\Big]\Big[(D_\rho\Phi)^\dagger D^\rho\Phi\Big], \tag{33}$$

and operators depending only on the electroweak field strength tensors $\widehat{W}_{\mu\nu}$ and $\widehat{B}_{\mu\nu}$, such as,

$$\mathcal{L}_{T,0} = \mathrm{Tr}\Big[\widehat{W}_{\mu\nu}\widehat{W}^{\mu\nu}\Big]\mathrm{Tr}\Big[\widehat{W}_{\rho\sigma}\widehat{W}^{\rho\sigma}\Big]. \tag{34}$$

Anomalous Higgs boson HVV coupling parameters are controlled in VBFNLO via the file `anom_HVV.dat`, where three different parameterizations are available. The latter consist of the Wilson coefficients associated with the subset of dimension-six operators introduced above that contribute to the Higgs couplings, the parametrization used by the L3 collaboration that is defined in Ref. [121], and a parameterization based on anomalous couplings in the mass basis that is thus expressed in terms of the field strength and dual field strength tensors of the W and Z bosons [122]. The relationships between these three parameterizations are discussed in more detail on the VBFNLO webpage.

Since the pure operators for anomalous gauge boson couplings might lead to a violation of tree-level unitarity within the energy range of the LHC, special care has to be taken to avoid this unphysical behaviour. Within VBFNLO, we have opted for using the following form factors, all of them depending on Λ, the characteristic scale where the form factor effects become relevant. Equivalently, introducing these form factors is like restricting the validity range of the effective field theory rather than implementing a sharp cutoff. For HVV vertices, two different form factors can be chosen as described in Refs. [122, 123].

$$F_1 = \frac{\Lambda^2}{|q_1|^2+\Lambda^2}\,\frac{\Lambda^2}{|q_2|^2+\Lambda^2}, \qquad F_2 = -2\,\Lambda^2\,C_0\Big(q_1^2, q_2^2, (q_1+q_2)^2, \Lambda^2\Big). \tag{35}$$

where q_i are the momenta of the vector bosons and C_0 is the scalar one-loop three point function in the notation of Ref. [124]. For triple and quartic gauge couplings the form factor takes the form

$$F = \left(1+\frac{s}{\Lambda^2}\right)^{-p}, \tag{36}$$

for each phase space point, where s is a universal scale identified with the invariant mass squared of the produced bosons. Finally, for γjj production in VBF, the form factor reads

$$F = \left(1+\frac{|q_1|^2}{\Lambda^2}+\frac{|q_2|^2}{\Lambda^2}+\frac{|q_3|^2}{\Lambda^2}\right)^{-p}, \tag{37}$$

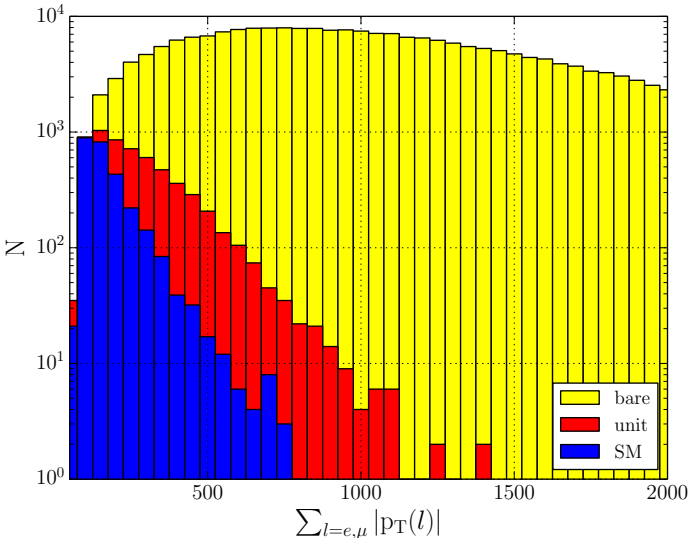

Fig. 2: Events generated by WHIZARD when a dimension-eight coupling $F_{S,0} = 480$ TeV^{-4} (the Wilson coefficient associated with the $\mathcal{O}_{S,0}$ operator of Eq. (32)) is added to the Standard Model and for the process $pp \to e^+\mu^+\nu_e\nu_\mu jj$ at a center-of-mass energy of 14 TeV and for an integrated luminosity of $\mathcal{L} = 1000$ fb^{-1}. Event selection requires a dijet invariant mass $M_{jj} > 500$ GeV, a rapidity separation between the jets of $\Delta y_{jj} > 2.4$, and the two jets must have a transverse momentum $p_T^j > 20$ GeV and a pseudorapidity $|\eta_j| < 4.5$. In addition, leptons are required to have a transverse momentum $p_T^l > 20$ GeV. Red and yellow histograms represent the naive and unitarized effective field theory predictions, respectively, while the pure Standard model expectation is shown as a blue histogram for comparison.

with q_1^2, q_2^2 and q_3^2 as the invariant masses squared of the three vector bosons involved in the VVV vertex. On the VBFNLO webpage, we provide the 'Formfactor Calculation Tool for aGC' which gives the maximum value of the scale Λ which is allowed by unitarity. The value is determined by calculating on-shell VV scattering and computing the zeroth partial wave of the amplitude. As unitarity criterion, the absolute value of the real part of the zeroth partial wave has to be below 0.5 [125]. Each channel in $VV \to VV$ scattering (with $V = W/Z/\gamma$) is checked individually, while additionally channels with the same electrical charge for the VV system are combined [126]. We recall that the definition of the partial wave expansion in Ref. [126] differs from ours by a factor of 2.

Finally, the K-matrix unitarization procedure has been implemented for the two dimension-eight operators $\mathcal{O}_{S,0}+\mathcal{O}_{S,2}$ and $\mathcal{O}_{S,1}$ by using the relations to the operators $\mathcal{O}_4$ and $\mathcal{O}_5$ of the electroweak chiral Lagrangian and the procedure worked out in Ref. [127]. This method guarantees the preservation of unitarity when either of these operators is used in the study of anomalous quartic gauge couplings. In contrast to using form factors, no additional input parameters need to be set. The anomalous contributions are automatically suppressed at the energy scale where unitarity would be violated without unitarization. After this energy scale is reached, the anomalous contributions are kept at a finite value, representing the maximally possible anomalous contribution. We refer to Refs. [118, 127, 128] for details of the K-matrix unitarization procedure and its implementation.

1.8 Event generation with WHIZARD

WHIZARD [129] is a multipurpose event generator for hadron and lepton colliders. It has a highly optimized internal matrix element generator, O'MEGA, for the recursive computing of tree-level amplitudes for almost arbitrary theories [130]. In the QCD case, it additionally uses the color flow formalism [131]. WHIZARD has its own parton shower implementation [132], with both a p_T-ordered shower and an analytic shower. Recently, automated FKS subtraction and POWHEG matching for NLO QCD corrections (using external virtual matrix elements) have been implemented [133, 134]. For external theories beyond the Standard Model, interfaces to the packages SARAH and FEYNRULES [135] are available, and support for the UFO file format allowing for arbitrary Lorentz and color structures is currently under way.

With respect to Higgs Effective Field Theories and higher-dimensional operators, WHIZARD supports the whole set of bosonic dimension-six operators in the Warsaw basis [99]. For vector-boson scattering at the LHC, usually dimension-eight operators in the coupled system of electroweak gauge and Higgs bosons are as important as dimension-six operators as they can be generated at tree-level in some new physics models. These electroweak dimension-eight operators have been implemented in WHIZARD, both as plain operators and also with their interplay with new resonances in the electroweak sector. The description of new physics contributions with a low energy effective energy is however only valid up to an

a priori unknown scale Λ. At energies above Λ, the effective field theory will lead to unphysical predictions. As an example, the sum of the norm of the transverse momenta of all leptons produced in the process $pp \to e^+\mu^+\nu_e\nu_\mu$ is shown on Figure 2 when the dimension-eight operator $\mathcal{O}_{S,0}$ of Eq. (32) is included. The number of events generated with a naive effective field theory description (yellow histogram) will largely overshoot a physically possible distribution, because S-matrix unitarity is violated within the experimental energy reach. A selection to avoid energy regions where the theoretical description breaks down is not possible in general, due to the inability to reconstruct the invariant mass for some final states. Using the T-matrix scheme, a unitarization prescription for high-energy regions of the phase space [127, 128, 136], avoid the unphysical high number of generated events by WHIZARD. Instead, the number of events are saturated in every isospin-spin channel to satisfy S-matrix unitarity (red histogram). In order to simplify unitarization, WHIZARD does not use the basis of Ref. [137] as this basis does not respect isospin symmetry which makes calculation of the channels in need for unitarization much easier.

On different lines, fermionic dimension-six operators are implemented for the top quark sector as (form-factor regularized) anomalous couplings [138, 139].

1.9 ROSETTA

Different complete and non-redundant operator bases of dimension-six effective operators have been proposed in the literature, the most popular choices including the Warsaw basis [99], the SILH basis [71, 100] and the beyond the Standard Model primaries basis [140–142]. It is however cumbersome to express any experimental result in a basis-independent manner. Different bases may be convenient for particular applications, either because they facilitate the comparison with a given class of theories or simply because different experimental analyses look more transparent in a specific basis. The ROSETTA package [98] has been designed to explicitly solve such problems by allowing for a straightforward translation between different effective field theory languages. In addition to translating, another important goal of the ROSETTA program is to provide a platform for communication with Monte Carlo event generators, no matter which basis is chosen. To achieve this, ROSETTA contains an implementation of the Higgs basis (Ref. to the WG2 chapter) and is connected to the BSMC Lagrangian introduced in Section 1.6. More precisely, the output format of ROSETTA has been tuned so that the translation maps an effective field theory Lagrangian given in a specific basis to the BSMC Lagrangian and generates an output file that is compatible with the BSMC implementation into FEYNRULES [89]. As a consequence, any high-energy physics tool that is interfaced to FEYNRULES can be employed within the context of any basis of dimension-six operators that is included in ROSETTA. A full description of the program including detailed example usage and information on how to create a user-defined basis class can be found in the manual [98].

The most basic functionality of ROSETTA is to map a chosen set of input parameters (the Wilson coefficients in a specific basis choice) onto the BSMC coefficients such that the output can be employed within tools relying on a BSMC Lagrangian description. In addition, the user may define his/her own map to the BSMC coefficients (or to any other basis implementation) and proceed with event generation using the related FEYNRULES implementation. This highlights one of the key features of ROSETTA, the possibility to easily define one's own input basis and directly use it in the context of many programs via the translation functionality. An example of this would be the eHDECAY program (see Ref. [100] and Section 1.4) for which ROSETTA provides an interface to calculate the Higgs width and branching fractions including effective field theory effects in any basis. The strength of this approach is that it is much simpler than developing from scratch new modules for existing tools in the context of a new basis. Moreover, ROSETTA not only enables translation into the BSMC Lagrangian, but also allows for translations into any of the other bases included in the package that are currently the Higgs, Warsaw and SILH bases. Translations between these three bases in any direction are possible, so that the addition of a new basis by the user only requires the specification of translation rules to any one of the three core bases. One is subsequently able to indirectly translate the new basis into any of the other two bases, as well as into the BSMC Lagrangian.

The latest release of ROSETTA can be obtained from
`http://rosetta.hepforge.org`.
The package contains a PYTHON executable named `translate`, an information file named `README` and two directories, a first folder (named `Cards`) collecting example input files and a second folder (named `Rosetta`) including the source code of ROSETTA. Each basis is implemented as a class in its own PYTHON module, where the coefficients and the required inputs are declared. Moreover, the input format must respect conventions inspired by the Supersymmetry Les Houches Accord (SLHA) [143, 144] and the related SLHA block structure is defined in the class. Translations are defined as member functions of the class with a special syntax to identify the target basis. Several utility functions may also be implemented to either calculate the values of parameters declared as dependent as a function of those declared as independent or modify the values of the Standard Model input parameters as a function of the effective operator coefficients.

The `translate` executable takes as input an SLHA-style parameter file with the coefficients of the dimension-six operators associated with a particular basis. Information on the format of such an input file can be found in the manual [98]. Depending on the basis implementation, some basic input quantities, such as Standard Model input constants and particle masses may be required in addition to the effective field theory coefficients. The execution of the `translate` script

from a shell yields the generation of an output parameter file where all parameters are this time the coefficients of the dimension-six operators associated with a specified new basis, the default choice being the BSMC Lagrangian. The tool can be used by typing in

```
./translate PARAMCARD.dat OPTION
```

where `PARAMCARD.dat` is the name of the SLHA-style input file and `OPTIONS` stands for optional arguments. These can range from specifying the output file name (`-output`, `-o`) or target basis into which to translate (`-target`, `-t`) to invoking the eHDECAY (`-ehdecay`, `-e`) interface to additionally generate an SLHA decay block for the Higgs as part of the output file. A particularly useful option is the `-flavour` or `-f` one which allows users to specify the treatment of the flavor structure relevant for fermionic operators. This can take the values `general`, `diagonal` and `universal` depending on the desired degree of simplicity. A more complicated example of usage might read

```
./translate myinput.dat -t warsaw -f universal -e -o myoutput.dat
```

which would read the `myinput.dat` file and translate it to the Warsaw basis, assuming the universal flavor structure which is designed to map easily to the Minimal Flavour Violation assumption [145]. The eHDECAY interface will also be called, writing in the output file, specified to be `myoutput.dat` an additional SLHA block containing Higgs decay information.

With the advent of NLO-accurate Monte Carlo event generation software, it is important that ROSETTA remains flexible enough to eventually provide compatibility with this new generation of tools. The future development plans of the program are connected to the recent progresses that have been made on the theory side both in implementing the linear dimension-six description in the FEYNRULES framework [146] and in calculating the renormalization group evolution of the full set of operators and their mutual mixing [147–150]. In the former case, ROSETTA can simply be extended to provide an output compatible with the eventual NLO model implementation, analogously to the BSMC Lagrangian. The latter case of evaluating the renormalization group running effects, while being a slightly separate issue, highlights a key feature of ROSETTA, given that the calculation of these effects has only been performed in the Warsaw basis and that ROSETTA could allow for the application of these results in any desired basis.

References

[1] ATLAS Collaboration, G. Aad *et al.*, Phys.Lett. **B716**, 1 (2012), 1207.7214. (1)

[2] CMS Collaboration, S. Chatrchyan *et al.*, Phys.Lett. **B716**, 30 (2012), 1207.7235. (1)

[3] M. Spira, (1995), hep-ph/9510347. (1)

[4] M. Spira, Nucl. Instrum. Meth. **A389**, 357 (1997), hep-ph/9610350. (1)

[5] A. Djouadi, M. Spira, and P. M. Zerwas, Phys. Lett. **B264**, 440 (1991). (1, 5)

[6] S. Dawson, Nucl. Phys. **B359**, 283 (1991). (1)

[7] D. Graudenz, M. Spira, and P. M. Zerwas, Phys. Rev. Lett. **70**, 1372 (1993). (1)

[8] M. Spira, A. Djouadi, D. Graudenz, and P. M. Zerwas, Nucl. Phys. **B453**, 17 (1995), hep-ph/9504378. (1, 5)

[9] R. Harlander and P. Kant, JHEP **12**, 015 (2005), hep-ph/0509189. (1)

[10] C. Anastasiou, S. Bucherer, and Z. Kunszt, JHEP **10**, 068 (2009), 0907.2362. (1)

[11] S. Catani, D. de Florian, and M. Grazzini, JHEP **05**, 025 (2001), hep-ph/0102227. (1)

[12] R. V. Harlander and W. B. Kilgore, Phys. Rev. **D64**, 013015 (2001), hep-ph/0102241. (1)

[13] R. V. Harlander and W. B. Kilgore, Phys. Rev. Lett. **88**, 201801 (2002), hep-ph/0201206. (1)

[14] V. Ravindran, J. Smith, and W. L. van Neerven, Nucl. Phys. **B665**, 325 (2003), hep-ph/0302135. (1)

[15] C. Anastasiou and K. Melnikov, Nucl. Phys. **B646**, 220 (2002), hep-ph/0207004. (1)

[16] A. Djouadi and P. Gambino, Phys. Rev. Lett. **73**, 2528 (1994), hep-ph/9406432. (1, 5)

[17] K. G. Chetyrkin, B. A. Kniehl, and M. Steinhauser, Phys. Rev. Lett. **78**, 594 (1997), hep-ph/9610456. (1, 5)

[18] K. G. Chetyrkin, B. A. Kniehl, and M. Steinhauser, Nucl. Phys. **B490**, 19 (1997), hep-ph/9701277. (1, 5)

[19] U. Aglietti, R. Bonciani, G. Degrassi, and A. Vicini, Phys. Lett. **B595**, 432 (2004), hep-ph/0404071. (1, 5)

[20] U. Aglietti, R. Bonciani, G. Degrassi, and A. Vicini, Two-loop electroweak corrections to Higgs production in proton-proton collisions, in *TeV4LHC Workshop: 2nd Meeting Brookhaven, Upton, New York, February 3-5, 2005*, 2006, hep-ph/0610033. (1, 5)

[21] G. Degrassi and F. Maltoni, Phys. Lett. **B600**, 255 (2004), hep-ph/0407249. (1, 5)

[22] S. Actis, G. Passarino, C. Sturm, and S. Uccirati, Phys. Lett. **B670**, 12 (2008), 0809.1301. (1, 5)

[23] S. Actis, G. Passarino, C. Sturm, and S. Uccirati, Nucl. Phys. **B811**, 182 (2009), 0809.3667. (1, 5)

[24] C. Anastasiou, R. Boughezal, and F. Petriello, JHEP **04**, 003 (2009), 0811.3458. (1)

[25] R. P. Kauffman and W. Schaffer, Phys. Rev. **D49**, 551 (1994), hep-ph/9305279. (1)

[26] A. Djouadi, M. Spira, and P. M. Zerwas, Phys. Lett. **B311**, 255 (1993), hep-ph/9305335. (1)

[27] M. Spira, A. Djouadi, D. Graudenz, and P. M. Zerwas, Phys. Lett. **B318**, 347 (1993). (1)

[28] S. Dawson, A. Djouadi, and M. Spira, Phys. Rev. Lett. **77**, 16 (1996), hep-ph/9603423. (1)

[29] M. Muhlleitner and M. Spira, Nucl. Phys. **B790**, 1 (2008), hep-ph/0612254. (1)

[30] R. Bonciani, G. Degrassi, and A. Vicini, JHEP **11**, 095 (2007), 0709.4227. (1)

[31] C. Anastasiou, S. Beerli, S. Bucherer, A. Daleo, and Z. Kunszt, JHEP **01**, 082 (2007), hep-ph/0611236. (1)

[32] U. Aglietti, R. Bonciani, G. Degrassi, and A. Vicini, JHEP **01**, 021 (2007), hep-ph/0611266. (1)

[33] R. V. Harlander and W. B. Kilgore, JHEP **10**, 017 (2002), hep-ph/0208096. (1)

[34] C. Anastasiou and K. Melnikov, Phys. Rev. **D67**, 037501 (2003), hep-ph/0208115. (1)

[35] R. V. Harlander and M. Steinhauser, Phys. Lett. **B574**, 258 (2003), hep-ph/0307346. (1)

[36] R. Harlander and M. Steinhauser, Phys. Rev. **D68**, 111701 (2003), hep-ph/0308210. (1)

[37] R. V. Harlander and M. Steinhauser, JHEP **09**, 066 (2004), hep-ph/0409010. (1)

[38] R. V. Harlander and F. Hofmann, JHEP **03**, 050 (2006), hep-ph/0507041. (1)

[39] G. Degrassi and P. Slavich, Nucl. Phys. **B805**, 267 (2008), 0806.1495. (1)

[40] G. Degrassi and P. Slavich, JHEP **11**, 044 (2010), 1007.3465. (1)

[41] G. Degrassi, S. Di Vita, and P. Slavich, JHEP **08**, 128 (2011), 1107.0914. (1)

[42] G. Degrassi, S. Di Vita, and P. Slavich, Eur. Phys. J. **C72**, 2032 (2012), 1204.1016. (1)

[43] C. Anastasiou, S. Beerli, and A. Daleo, Phys. Rev. Lett. **100**, 241806 (2008), 0803.3065. (1)

[44] M. Muhlleitner, H. Rzehak, and M. Spira, PoS **RADCOR2009**, 043 (2010), 1001.3214, [,42(2010)]. (1)

[45] M. Muhlleitner, H. Rzehak, and M. Spira, MSSM Higgs boson production via gluon fusion, in *Physics at the LHC2010. Proceedings, 5th Conference, PLHC2010, Hamburg, Germany, June 7-12, 2010*, pp. 415–417, 2010. (1)

[46] K. G. Chetyrkin, B. A. Kniehl, and M. Steinhauser, Phys. Rev. Lett. **79**, 2184 (1997), hep-ph/9706430. (1)

[47] K. G. Chetyrkin, B. A. Kniehl, and M. Steinhauser, Nucl. Phys. **B510**, 61 (1998), hep-ph/9708255. (1, 5)

[48] M. Kramer, E. Laenen, and M. Spira, Nucl. Phys. **B511**, 523 (1998), hep-ph/9611272. (1, 5)

[49] M. Spira, Fortsch. Phys. **46**, 203 (1998), hep-ph/9705337. (1)

[50] D. Choudhury and P. Saha, JHEP **08**, 144 (2012), 1201.4130. (1)

[51] C. Degrande, J. M. Gerard, C. Grojean, F. Maltoni, and G. Servant, JHEP **07**, 036 (2012), 1205.1065, [Erratum: JHEP03,032(2013)]. (1)

[52] C. G. Callan, Jr., Phys. Rev. **D2**, 1541 (1970). (2)

[53] K. Symanzik, Commun. Math. Phys. **18**, 227 (1970). (2)

[54] S. R. Coleman and R. Jackiw, Annals Phys. **67**, 552 (1971). (2)

[55] R. J. Crewther, Phys. Rev. Lett. **28**, 1421 (1972). (2)

[56] M. S. Chanowitz and J. R. Ellis, Phys. Lett. **B40**, 397 (1972). (2)

[57] M. S. Chanowitz and J. R. Ellis, Phys. Rev. **D7**, 2490 (1973). (2)

[58] A. Djouadi, J. Kalinowski, and M. Spira, Comput. Phys. Commun. **108**, 56 (1998), hep-ph/9704448. (2, 4)

[59] A. Djouadi, M. M. Muhlleitner, and M. Spira, Acta Phys. Polon. **B38**, 635 (2007), hep-ph/0609292. (2, 4)

[60] S. Dawson, S. Dittmaier, and M. Spira, Phys. Rev. **D58**, 115012 (1998), hep-ph/9805244. (3)

[61] J. Grigo, J. Hoff, K. Melnikov, and M. Steinhauser, Nucl. Phys. **B875**, 1 (2013), 1305.7340. (3)

[62] R. Frederix *et al.*, Phys. Lett. **B732**, 142 (2014), 1401.7340. (3)

[63] F. Maltoni, E. Vryonidou, and M. Zaro, JHEP **11**, 079 (2014), 1408.6542. (3)

[64] J. Grigo, J. Hoff, and M. Steinhauser, Nucl. Phys. **B900**, 412 (2015), 1508.00909. (3)

[65] D. de Florian and J. Mazzitelli, Phys. Rev. Lett. **111**, 201801 (2013), 1309.6594. (3)

[66] D. de Florian and J. Mazzitelli, JHEP **09**, 053 (2015), 1505.07122. (3)

[67] J. Grigo, K. Melnikov, and M. Steinhauser, Nucl. Phys. **B888**, 17 (2014), 1408.2422. (3)

[68] R. Grober, M. Muhlleitner, M. Spira, and J. Streicher, JHEP **09**, 092 (2015), 1504.06577. (3)

[69] E. W. N. Glover and J. J. van der Bij, Nucl. Phys. **B309**, 282 (1988). (3)

[70] T. Plehn, M. Spira, and P. M. Zerwas, Nucl. Phys. **B479**, 46 (1996), hep-ph/9603205, [Erratum: Nucl. Phys.B531,655(1998)]. (3)

[71] G. F. Giudice, C. Grojean, A. Pomarol, and R. Rattazzi, JHEP **06**, 045 (2007), hep-ph/0703164. (4, 7, 12)

[72] R. Contino, M. Ghezzi, C. Grojean, M. Mühlleitner, and M. Spira, Comput. Phys. Commun. **185**, 3412 (2014), 1403.3381. (4)

[73] K. Agashe, R. Contino, and A. Pomarol, Nucl. Phys. **B719**, 165 (2005), hep-ph/0412089. (4)

[74] R. Contino, L. Da Rold, and A. Pomarol, Phys. Rev. **D75**, 055014 (2007), hep-ph/0612048. (4)

[75] Y. Schroder and M. Steinhauser, JHEP **0601**, 051 (2006), hep-ph/0512058. (5)

[76] K. Chetyrkin, J. H. Kuhn, and C. Sturm, Nucl.Phys. **B744**, 121 (2006), hep-ph/0512060. (5)

[77] P. Baikov and K. Chetyrkin, Phys.Rev.Lett. **97**, 061803 (2006), hep-ph/0604194. (5)

[78] T. Inami, T. Kubota, and Y. Okada, Z.Phys. **C18**, 69 (1983). (5)

[79] K. Chetyrkin, B. A. Kniehl, and M. Steinhauser, Phys.Rev.Lett. **79**, 353 (1997), hep-ph/9705240. (5)

[80] J. R. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl.Phys. **B106**, 292 (1976). (5)

[81] M. A. Shifman, A. Vainshtein, M. Voloshin, and V. I. Zakharov, Sov.J.Nucl.Phys. **30**, 711 (1979). (5)

[82] B. A. Kniehl and M. Spira, Z.Phys. **C69**, 77 (1995), hep-ph/9505225. (5)

[83] A. Denner *et al.*, Report No. LHCHXSWG-INT-2015-006, 2015 (unpublished). (5)

[84] LHC Higgs Cross Section Working Group, A. David *et al.*, (2012), 1209.0040. (5)

[85] M. Gonzalez-Alonso, A. Greljo, G. Isidori, and D. Marzocca, Eur. Phys. J. **C75**, 128 (2015), 1412.6038. (6, 7)

[86] A. Greljo, G. Isidori, J. M. Lindert, and D. Marzocca, (2015), 1512.06135. (6, 7)

[87] J. S. Gainer, J. Lykken, K. T. Matchev, S. Mrenna, and M. Park, The Matrix Element Method: Past, Present, and Future, in *Community Summer Study 2013: Snowmass on the Mississippi (CSS2013) Minneapolis, MN, USA, July 29-August 6, 2013*, 2013, 1307.3546. (6)

[88] M. Bordone, A. Greljo, G. Isidori, D. Marzocca, and A. Pattori, Eur. Phys. J. **C75**, 385 (2015), 1507.02555. (6, 7)

[89] A. Alloul, N. D. Christensen, C. Degrande, C. Duhr, and B. Fuks, Comput. Phys. Commun. **185**, 2250 (2014), 1310.1921. (7, 12)

[90] C. Degrande *et al.*, Comput. Phys. Commun. **183**, 1201 (2012), 1108.2040. (7)

[91] J. Alwall *et al.*, JHEP **07**, 079 (2014), 1405.0301. (7)

[92] T. Gleisberg *et al.*, JHEP **02**, 007 (2009), 0811.4622. (7)

[93] T. Sjostrand, S. Mrenna, and P. Z. Skands, Comput. Phys. Commun. **178**, 852 (2008), 0710.3820. (7)

[94] P. Artoisenet, R. Frederix, O. Mattelaer, and R. Rietkerk, JHEP **03**, 015 (2013), 1212.3460. (7, 8)

[95] P. Artoisenet *et al.*, JHEP **11**, 043 (2013), 1306.6464. (7)

[96] N. D. Christensen *et al.*, Eur. Phys. J. **C71**, 1541 (2011), 0906.2474. (7)

[97] P. de Aquino, W. Link, F. Maltoni, O. Mattelaer, and T. Stelzer, Comput. Phys. Commun. **183**, 2254 (2012), 1108.2041. (7)

[98] A. Falkowski *et al.*, Eur. Phys. J. **C75**, 583 (2015), 1508.05895. (7, 12)

[99] B. Grzadkowski, M. Iskrzynski, M. Misiak, and J. Rosiek, JHEP **10**, 085 (2010), 1008.4884. (7, 11, 12)

[100] R. Contino, M. Ghezzi, C. Grojean, M. Mühlleitner, and M. Spira, JHEP **1307**, 035 (2013), 1303.3876. (7, 12)

[101] A. Alloul, B. Fuks, and V. Sanz, JHEP **04**, 110 (2014), 1310.5150. (7)

[102] F. Maltoni, K. Mawatari, and M. Zaro, Eur. Phys. J. **C74**, 2710 (2014), 1311.1829. (7)

[103] F. Demartin, F. Maltoni, K. Mawatari, B. Page, and M. Zaro, Eur. Phys. J. **C74**, 3065 (2014), 1407.5089. (7)

[104] F. Demartin, F. Maltoni, K. Mawatari, and M. Zaro, Eur. Phys. J. **C75**, 267 (2015), 1504.00611. (7)

[105] S. Frixione, E. Laenen, P. Motylinski, and B. R. Webber, JHEP **04**, 081 (2007), hep-ph/0702198. (8)

[106] R. D. Ball *et al.*, Nucl. Phys. **B867**, 244 (2013), 1207.1303. (9)

[107] B. Hespel, F. Maltoni, and E. Vryonidou, JHEP **06**, 065 (2015), 1503.01656. (9)

[108] B. Hespel, D. Lopez-Val, and E. Vryonidou, JHEP **09**, 124 (2014), 1407.0281. (9)

[109] K. Arnold *et al.*, Comput. Phys. Commun. **180**, 1661 (2009), 0811.4559. (9)

[110] J. Baglio *et al.*, (2014), 1404.3940. (9)

[111] T. Binoth *et al.*, Comput. Phys. Commun. **181**, 1612 (2010), 1001.1307, [,1(2010)]. (9)

[112] S. Alioli *et al.*, Comput. Phys. Commun. **185**, 560 (2014), 1308.3462. (9)

[113] M. Bahr *et al.*, Eur. Phys. J. **C58**, 639 (2008), 0803.0883. (9)

[114] J. Bellm *et al.*, (2015), 1512.01178. (9)

[115] W. Buchmuller and D. Wyler, Nucl. Phys. **B268**, 621 (1986). (9)

[116] K. Hagiwara, S. Ishihara, R. Szalapski, and D. Zeppenfeld, Phys. Rev. **D48**, 2182 (1993). (9, 10)

[117] O. J. P. Eboli, M. C. Gonzalez-Garcia, and J. K. Mizukoshi, Phys. Rev. **D74**, 073005 (2006), hep-ph/0606118. (9, 10)

[118] C. Degrande *et al.*, Monte Carlo tools for studies of non-standard electroweak gauge boson interactions in multi-boson processes: A Snowmass White Paper, in *Community Summer Study 2013: Snowmass on the Mississippi (CSS2013) Minneapolis, MN, USA, July 29-August 6, 2013*, 2013, 1309.7890. (9, 11)

[119] K. Arnold *et al.*, (2011), 1107.4038. (10)

[120] K. Hagiwara, R. Szalapski, and D. Zeppenfeld, Phys. Lett. **B318**, 155 (1993), hep-ph/9308347. (10)

[121] L3, P. Achard *et al.*, Phys. Lett. **B589**, 89 (2004), hep-ex/0403037. (10)

[122] T. Figy and D. Zeppenfeld, Phys. Lett. **B591**, 297 (2004), hep-ph/0403297. (10)

[123] V. Hankele, G. Klamke, D. Zeppenfeld, and T. Figy, Phys. Rev. **D74**, 095001 (2006), hep-ph/0609075. (10)

[124] G. Passarino and M. J. G. Veltman, Nucl. Phys. **B160**, 151 (1979). (10)

[125] V. D. Barger, K.-m. Cheung, T. Han, and R. J. N. Phillips, Phys. Rev. **D42**, 3052 (1990). (11)

[126] G. J. Gounaris, J. Layssac, and F. M. Renard, Phys. Lett. **B332**, 146 (1994), hep-ph/9311370. (11)

[127] A. Alboteanu, W. Kilian, and J. Reuter, JHEP **11**, 010 (2008), 0806.4145. (11, 12)

[128] W. Kilian, T. Ohl, J. Reuter, and M. Sekulla, Phys. Rev. **D91**, 096007 (2015), 1408.6207. (11, 12)

[129] W. Kilian, T. Ohl, and J. Reuter, Eur. Phys. J. **C71**, 1742 (2011), 0708.4233. (11)

[130] M. Moretti, T. Ohl, and J. Reuter, (2001), hep-ph/0102195. (11)

[131] W. Kilian, T. Ohl, J. Reuter, and C. Speckner, JHEP **10**, 022 (2012), 1206.3700. (11)

[132] W. Kilian, J. Reuter, S. Schmidt, and D. Wiesler, JHEP **04**, 013 (2012), 1112.1039. (11)

[133] C. Weiss, B. C. Nejad, W. Kilian, and J. Reuter, PoS **EPS-HEP2015**, 466 (2015), 1510.02666. (11)

[134] B. C. Nejad, W. Kilian, J. Reuter, and C. Weiss, PoS **EPS-HEP2015**, 317 (2015), 1510.02739. (11)

[135] N. D. Christensen, C. Duhr, B. Fuks, J. Reuter, and C. Speckner, Eur. Phys. J. **C72**, 1990 (2012), 1010.3251. (11)

[136] W. Kilian, T. Ohl, J. Reuter, and M. Sekulla, (2015), 1511.00022. (12)

[137] O. J. P. Eboli, M. C. Gonzalez-Garcia, and S. M. Lietti, Phys. Rev. **D69**, 095005 (2004), hep-ph/0310141. (12)

[138] F. Bach and T. Ohl, Phys. Rev. **D86**, 114026 (2012), 1209.4564. (12)

[139] F. Bach and T. Ohl, Phys. Rev. **D90**, 074022 (2014), 1407.2546. (12)

[140] R. S. Gupta, A. Pomarol, and F. Riva, Phys. Rev. **D91**, 035001 (2015), 1405.0181. (12)

[141] E. Masso, JHEP **10**, 128 (2014), 1406.6376. (12)

[142] A. Pomarol, Higgs Physics, in *2014 European School of High-Energy Physics (ESHEP 2014) Garderen, The Netherlands, June 18-July 1, 2014*, 2014, 1412.4410. (12)

[143] P. Z. Skands *et al.*, JHEP **07**, 036 (2004), hep-ph/0311123. (12)

[144] B. C. Allanach *et al.*, Comput. Phys. Commun. **180**, 8 (2009), 0801.0045. (12)

[145] G. D'Ambrosio, G. F. Giudice, G. Isidori, and A. Strumia, Nucl. Phys. **B645**, 155 (2002), hep-ph/0207036. (13)

[146] C. Degrande, B. Fuks, K. Mawatari, K. Mimasu, and V. Sanz, (2016), in preparation. (13)

[147] E. E. Jenkins, A. V. Manohar, and M. Trott, JHEP **10**, 087 (2013), 1308.2627. (13)

[148] E. E. Jenkins, A. V. Manohar, and M. Trott, JHEP **01**, 035 (2014), 1310.4838. (13)

[149] R. Alonso, E. E. Jenkins, A. V. Manohar, and M. Trott, JHEP **04**, 159 (2014), 1312.2014. (13)

[150] B. Henning, X. Lu, and H. Murayama, (2014), 1412.1837. (13)