

NLO QCD+EW predictions for HV and HV +jet production including parton-shower effects

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ABSTRACT: We present the first NLO QCD+EW predictions for Higgs boson production in association with a $l\nu_l$ or l^+l^- pair plus zero or one jets at the LHC. Fixed-order NLO QCD+EW calculations are combined with a QCD+QED parton shower using the recently developed resonance-aware method in the POWHEG framework. Moreover, applying the improved MiNLO technique to $Hl\nu_l$ +jet and Hl^+l^- +jet production at NLO QCD+EW, we obtain predictions that are NLO accurate for observables with both zero or one resolved jet. This approach permits also to capture higher-order effects associated with the interplay of EW corrections and QCD radiation. The behavior of EW corrections is studied for various kinematic distributions, relevant for experimental analyses of Higgsstrahlung processes at the 13 TeV LHC. Exact NLO EW corrections are complemented with approximate analytic formulae that account for the leading and next-to-leading Sudakov logarithms in the high-energy regime. In the tails of transverse-momentum distributions, relevant for analyses in the boosted Higgs regime, the Sudakov approximation works well, and NLO EW effects can largely exceed the ten percent level. Our predictions are based on the POWHEG BOX RES+OpenLoops framework in combination with the Pythia 8.1 parton shower.

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1 Introduction

The discovery of the Higgs boson [1, 2] has opened the door to the direct experimental investigation of the Higgs and Yukawa sectors of the Standard Model. While present measurements of Higgs boson properties and interactions are consistent with the Standard Model [3], the full set of data collected during Run II and in subsequent runs of the LHC will provide more and more stringent tests of the mechanism of electroweak symmetry breaking.

In this context, the associated production of a Higgs and a vector boson, $pp \rightarrow HV$ with $V = W$ and Z , plays a prominent role. In spite of the fact that the total cross sections for these so-called Higgsstrahlung processes are subleading as compared to Higgs boson production via gluon fusion and vector-boson fusion, the possibility to reconstruct the full HV final state and the clean signatures that result from leptonically decaying vector bosons offer unique opportunities of testing Higgs boson interactions with vector bosons and heavy quarks (see refs. [4–6] and references therein). The associated HV production makes it possible to disentangle Higgs boson couplings to W and Z bosons from one another and to measure them in a broad kinematic range. In addition, the presence of the associated vector boson allows for an efficient suppression of QCD backgrounds. In particular, $pp \rightarrow HV$ is the most favorable channel for measurements of the $H \rightarrow b\bar{b}$ branching ratio, and thus for determinations of the bottom Yukawa coupling. In HV production with $H \rightarrow b\bar{b}$ decay, the boosted region, with Higgs boson transverse momentum above 200 GeV, plays a particularly important role, both in order to achieve an improved control of the QCD backgrounds [4] and for the sensitivity to possible anomalies in the HVV couplings. Higgsstrahlung processes permit also to probe invisible Higgs boson decays, both through direct measurements of $pp \rightarrow HZ$ with invisible Higgs decays and through indirect bounds based on measurements of the $H \rightarrow b\bar{b}$ branching ratio.

The accuracy of present and future measurements of HV production, at the level of both fiducial cross sections and differential distributions, calls for increasingly accurate theoretical predictions. The inclusion of higher-order QCD corrections is crucial, both for total rates and for a precise description of the QCD radiation that accompanies the production of the HV system. The role of QCD corrections can be particularly important in the boosted regime or in the presence of cuts and for observables that are sensitive to QCD radiation.

In general, in order to account for experimental cuts and observables, higher-order QCD and EW predictions should be available for arbitrary differential distributions, and experimental analyses require particle-level Monte Carlo generators where state-of-the-art theoretical calculations are matched to parton showers. Finally, when QCD and EW higher-order effects are both sizable, also their combination needs to be addressed.

Theoretical calculations for the associated-production processes are widely available in the literature. Among the numerous studies on HV production at next-to-leading order (NLO) QCD we quote here refs. [7–9]. Predictions for inclusive HZ and HW production at next-to-next-to-leading order (NNLO) QCD have first been obtained in refs. [10, 11] and are implemented in the `VH@NNLO` program [12]. Besides contributions of Drell-Yan (DY) type, where the Higgs boson results from an s -channel $V^* \rightarrow HV$ subtopology, Higgsstrahlung at NNLO QCD involves also extra $\mathcal{O}(\alpha_s^2)$ contributions where the Higgs boson

couples to heavy-quark loops. Such non-DY contributions arise via squared one-loop amplitudes in the $gg \rightarrow HZ$ channel [13] and through the interference of one-loop and tree amplitudes in the $qg \rightarrow HVq$ and crossing-related channels. Studies of possible anomalous coupling in the $gg \rightarrow HZ$ channel can be found in refs. [14, 15]. Heavy-quark loop contributions to the $gg \rightarrow HZ$ channel are known up to $\mathcal{O}(\alpha_s^3)$ in the limit of the mass of the bottom quark going to zero, and the mass of the top quark going to infinity [16]. Their impact, especially in the boosted regime, can be quite significant [17].

Fully differential NNLO calculations for HV production with off-shell vector-boson decays were first presented in refs. [18–20], including all DY contributions plus heavy-quark-loop contributions to $gg \rightarrow HZ$. More recently, a NNLO QCD calculation that includes also the small heavy-quark loop contributions in the $qg \rightarrow HVq$ channel and in the crossing-related $\bar{q}g$ and $q\bar{q}$ channels became available [21] and also $HV \rightarrow b\bar{b}V$ production with NNLO QCD corrections both in the production and in the decay part of the process [22].

Analytic resummations have been discussed in refs. [23–26], while leading-logarithmic resummation can be routinely achieved through the matching of NLO QCD calculations to parton showers (PS). The first NLO+PS generators in the MC@NLO [27] and POWHEG frameworks [28–30] have been presented in refs. [31] and [32], respectively. More recently, new generators that provide an NLO accurate description of HV and $HV + \text{jet}$ production became available. The first generator of this kind was presented in ref. [33] based on the MiNLO method [34, 35], while a simulation of $pp \rightarrow HZ + 0$ and 1 jet, based on the MEPS@NLO multijet merging technique [36, 37], was presented in ref. [38]. Concerning fermion loops, the POWHEG BOX generator of ref. [33] can account for all $\mathcal{O}(\alpha_s^2)$ NLO contributions of DY and non-DY type to $pp \rightarrow HV + \text{jet}$ and also for the finite $gg \rightarrow HZ$ loop-induced contributions, with the possibility of studying anomalous couplings in the “kappa framework”. A more general study, which uses an effective field theory approach and introduces generic six-dimensional operators, can be found in ref. [39].

The heavy-quark loop-mediated production $gg \rightarrow HZg$ was first studied in ref. [40]. More recently, the Sherpa generator of ref. [38] has included also NNLO-type squared quark-loop contributions in the $gg \rightarrow HZ$, $gg \rightarrow HZg$, and $qg \rightarrow HZq$ plus crossing-related channels. Lately, a NNLO+PS generator for $pp \rightarrow HW$ [41] that combines the NNLO QCD calculation of ref. [18] with the parton shower using the method of refs. [35, 42] was presented.

Electroweak corrections to $pp \rightarrow HV$, including off-shell W - and Z -boson decays, are known at NLO [43, 44] and are implemented in the parton-level Monte Carlo program HAWK [45]. These corrections are at the level of 5% for inclusive quantities, but in the high-energy regime they can reach various tens of percent due to the presence of Sudakov logarithms [46–53]. For this reason, especially in boosted searches, the inclusion of EW corrections is mandatory. An interesting aspect of these corrections in HV production is that they induce also a dependence on the Higgs sector, and in particular on the trilinear coupling λ_{HHH} . Thus, precise measurements of Higgsstrahlung processes can be exploited for setting limits on λ_{HHH} [54–57]. To date, none of the existing NLO+PS generators implement EW corrections.

In this paper, for the first time, we present NLO QCD and NLO EW calculations for the production of a Higgs boson in conjunction with a $\ell\nu_\ell$ or $\ell^+\ell^-$ leptonic pair, plus zero or one jet, at the LHC. While, for convenience, the above-mentioned processes will often be denoted as HV/HVj production (with $V = W^\pm$ and Z) in the rest of the paper, all the results we are going to present always correspond to the complete decayed final-state processes, with spin effects, off-shell and non-resonant contributions taken into account.

At NLO QCD we include the full set of $\mathcal{O}(\alpha_s)$ contributions to $pp \rightarrow HV$ and $\mathcal{O}(\alpha_s^2)$ contributions to $pp \rightarrow HVj$. Although terms of non-DY type are implemented in our codes, we have not included them in our simulations. In addition, we do not include NNLO-like loop-induced contributions to HZ plus 0 and 1 jet production.

Besides showing fixed-order NLO QCD+EW predictions at parton level for typical observables, we also present full NLO+PS simulations for HV and HVj production. To this end, we have implemented our NLO calculations for HV and HVj production into four separate codes (HW^\pm , $HW^\pm j$, HZ and HZj) in the POWHEG BOX framework. In this way, we have consistently combined the radiation emitted at NLO QCD+EW level with a QCD+QED parton shower. In this context, photon radiation from the charged leptons can lead to severe unphysical distortions of the Z - and W -boson line shapes, if not properly treated. This problem was first pointed out in the context of NLO QCD+PS simulations of off-shell top-quark production and decay, and was solved in the context of the POWHEG BOX framework by means of the so-called resonance-aware method [58]. The first application of this method and its variants, in the context of electroweak corrections, has appeared in refs. [59, 60]. In this paper, we exploit the flexibility of the resonance-aware method to perform a fully consistent NLO QCD+EW matching in the presence of non-trivial EW resonances. To this end, our NLO calculations and generators are implemented in the new version of the POWHEG BOX framework, known as POWHEG BOX RES. In this recent version, the hardest radiation generated by POWHEG preserves the resonance virtualities present at the underlying-Born level. At the same time, the resonance information can be passed on to the parton shower, which in turn preserves the virtualities of intermediate resonances of the hard process in subsequent emissions.

Similarly to what was done in ref. [33] for HVj production at NLO QCD, we have applied the improved MiNLO [34, 35] approach to HVj production in order to get a sample of events that has simultaneously NLO QCD accuracy for HV plus 0 and 1 jet. In the MiNLO framework, also the NLO EW corrections to HV and HVj production have been consistently combined in the same inclusive sample. This can be regarded as an approximate treatment of $\mathcal{O}(\alpha_s \alpha_{EM})$ corrections in observables that are very sensitive to QCD radiation and receive, at the same time, large EW corrections. Moreover, although we do not present a rigorous proof, based on considerations related to unitarity and factorization of soft and collinear QCD radiation, we will argue that our MiNLO predictions should preserve full NLO QCD+EW accuracy in the phase space with zero or one resolved jets. As we will see, this conclusion is supported by our numerical results.

While our NLO EW results are exact (apart from the treatment of photon-initiated contributions), we also present approximate NLO EW predictions in the so-called Sudakov limit, where all kinematic invariants are well above the electroweak scale. Specifically, based

on the general results of refs. [49, 61], we provide explicit analytic expressions for all logarithmic EW corrections to $pp \rightarrow HV + 0$ and 1 jet in next-to-leading-logarithmic (NLL) approximation. Based on the observed accuracy of the NLL Sudakov formulas, this approximation can be exploited both in order to speed up the evaluation of EW corrections at NLO and in order to predict the dominant EW effects beyond NLO.

All needed matrix elements for $pp \rightarrow HV + 0$ and 1 jet at NLO EW have been generated using the `OpenLoops` program [62, 63], which supports the automated generation of NLO QCD+EW scattering amplitudes for Standard Model processes [64–66]. The implementation in the `POWHEG BOX RES` framework was achieved exploiting the generic interface developed in ref. [67]. For what concerns NLO QCD corrections, on the one hand we implemented in-house analytic expressions for the virtual corrections. On the other hand, following the approach of ref. [33], for real-emission contributions we used `MadGraph4` [68] matrix elements, via the interface described in ref. [69].

The paper is organized as follows. In section 2 we introduce the various ingredients of HV and HVj production at NLO QCD+EW. In particular, in section 2.2 we present a schematic proof of the NLO QCD+EW accuracy of `MinLO` predictions for inclusive observables. Further technical aspects of the calculation as well as input parameters and cuts are specified in section 3. Fixed-order NLO QCD+EW predictions are discussed in section 4, while in sections 5 and 6 we present NLO+PS QCD+EW results for HV production and `MinLO` QCD+EW results for HVj production, respectively. The predictions of the NLO+PS HV and `MinLO`+PS HVj generators are compared in section 7. Our main findings are summarized in section 8. In the appendices we document the validation of EW corrections in HV production against `HAWK` (appendix A), detailed NLO EW formulas in the Sudakov approximation (appendix B), a reweighting approach that we employ in order to speed up the evaluation of EW corrections (appendix C), and technical aspects of the interface between the `POWHEG BOX RES` and `Pythia 8.1` (appendix D).

2 NLO QCD and EW corrections to HV and HVj production

In this section we describe the QCD and EW NLO corrections to the production of a Higgs boson in association with a $\ell\nu_\ell$ or $\ell^+\ell^-$ leptonic pair plus zero or one additional jets. For convenience, these Higgsstrahlung processes will be denoted as associated HV and HVj production, with $V = W^\pm$ or Z . However, all results presented in this paper correspond to the complete processes

$$\begin{aligned}
 pp &\rightarrow HW^+(j) \rightarrow H \ell^+ \nu_\ell(j), \\
 pp &\rightarrow HW^-(j) \rightarrow H \ell^- \bar{\nu}_\ell(j), \\
 pp &\rightarrow HZ(j) \rightarrow H \ell^+ \ell^-(j),
 \end{aligned}
 \tag{2.1}$$

including all spin-correlation and off-shell effects. The combination of HW^+/HW^+j and HW^-/HW^-j Higgsstrahlung will be denoted as HW/HWj production. In our calculations, we have considered only one leptonic generation, and all leptons are treated as massless.

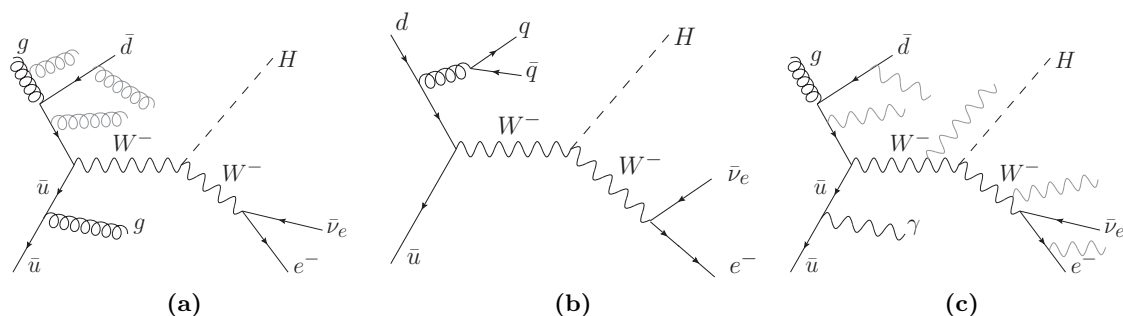


Figure 1. A sample of QCD (a, b), and EW (c), real radiation diagrams contributing to HW^-j production. While only one photon or gluon at a time is present at fixed order, for illustration purpose, in (a) and (c), we have shown various possible gluon and photon emissions.

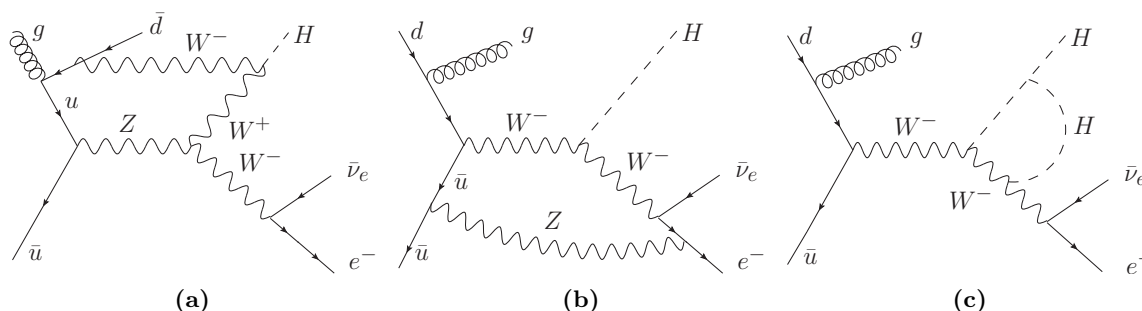


Figure 2. A sample of virtual EW diagrams contributing to HW^-j production.

2.1 NLO QCD+EW matrix elements

In this section we describe the various tree and one-loop amplitudes that have been assembled to form a NLO QCD+EW Monte Carlo program based on the POWHEG BOX framework [30].

Associated HV production proceeds through quark-antiquark annihilation at leading order, which corresponds to $\mathcal{O}(\alpha_{\text{EM}}^3)$. In HVj production, where the leading order corresponds to $\mathcal{O}(\alpha_s \alpha_{\text{EM}}^3)$, additional (anti)quark-gluon initiated processes contribute. All $\mathcal{O}(\alpha_s \alpha_{\text{EM}}^3)$ NLO QCD corrections to HV production have been computed analytically, since they simply affect the $Vq\bar{q}'$ vertex, and the calculation of the real and virtual corrections is trivial. In HVj production, the virtual $\mathcal{O}(\alpha_s^2 \alpha_{\text{EM}}^3)$ NLO QCD corrections have been computed analytically [70]. The color- and spin-correlated Born amplitudes and the real contributions at $\mathcal{O}(\alpha_s^2 \alpha_{\text{EM}}^3)$ have been computed using the automated interface [69] between the POWHEG BOX and MadGraph4 [68]. The real contributions involve tree diagrams with either an additional gluon or an external gluon replaced with a $q\bar{q}$ -pair. Example diagrams are shown in figure 1 (a, b).

The virtual EW corrections to HV and HVj production comprise loop amplitudes up to pentagon and hexagon configurations, respectively. Example diagrams for HVj production are shown in figure 2. All the internal resonances have been treated in the complex-mass scheme [71, 72] throughout.

As pointed out in the Introduction and illustrated in figure 2 (c), the virtual NLO EW amplitudes induce a dependence on the Higgs trilinear coupling λ_{HHH} . This dependence arises both from the bare virtual amplitudes and from the Higgs boson self-energy entering the Higgs boson wave-function renormalization. In view of the possibility of exploiting precision measurements of Higgsstrahlung processes for an indirect determination of λ_{HHH} , we allow λ_{HHH} to be set independently of the Higgs boson mass.¹

The real NLO EW corrections to HV and HVj production comprise QED radiation off all charged particles, i.e. they have an additional photon in the final state, as illustrated in figure 1 (c). Photon-induced real radiation contributions, where the photon is crossed to the initial state, are, on the other hand, not considered here, as they are suppressed by the small photon density in the proton. These corrections for HV production have been computed for the first time in ref. [44] and are included in the HAWK [45] Monte Carlo generator. Interestingly, they reach several percent for inclusive HW production, but remain at the 2% level when leptonic selection cuts are applied, and are negligible for HZ production [73]. For HVj production, photon-induced contributions enter already at Born level: however, they are of $\mathcal{O}(\alpha_{\text{EM}}^4)$ and thus formally subleading with respect to the $\mathcal{O}(\alpha_s \alpha_{\text{EM}}^3)$ leading order. Still, the NLO QCD corrections to these photon-induced processes are of $\mathcal{O}(\alpha_s \alpha_{\text{EM}}^4)$ and thus formally of the same order as the NLO EW corrections to the quark-antiquark and (anti)quark-gluon initiated channels in HVj production. Also not considered here are mixed QCD-EW bremsstrahlung contributions to HVj production at $\mathcal{O}(\alpha_s \alpha_{\text{EM}}^4)$. These tree-level contributions are finite and can easily be investigated separately. Similar contributions in the NLO EW corrections to V +jet production are known to yield relevant contributions only in jet observables at very large transverse momentum [65, 74]. Finally, also virtual QCD corrections to $HV\gamma$ production contribute formally at $\mathcal{O}(\alpha_s \alpha_{\text{EM}}^4)$ and are thus of the same perturbative order as the NLO EW corrections to HVj production. However, if a photon isolation is applied, as is done in this paper (see section 3.4), $HV\gamma$ production can be considered as a separate process and thus excluded from the definition of HVj production.

All the electroweak real and virtual corrections have been computed using a recent interface of the POWHEG BOX RES to OpenLoops [67].

In this study we combine NLO QCD and EW corrections in an additive way, i.e. corresponding perturbative contributions are simply added. At fixed order, an improved description can easily be obtained via a factorized ansatz, where differential NLO QCD cross sections are multiplied with relative EW correction factors. Such a multiplicative combination can be motivated from the factorization of soft QCD radiation and EW Sudakov logarithms, which can be tested comparing relative NLO EW corrections for HV and HVj production.

2.2 MiNLO approach for HVj production at NLO QCD+EW

In order to obtain an optimal description of QCD radiation, both in the hard and soft regime, all NLO QCD+EW calculations for $pp \rightarrow HVj$ have been performed using the “Multiscale improved NLO” (MiNLO) [34] method. This approach effectively resums logarithmic singularities of soft and collinear type to NLL accuracy, thereby ensuring a finite

¹The corresponding parameter can directly be set in the POWHEG BOX RES input file.

HVj cross section in all regions of phase space, even when the extra jet becomes unresolved. In the MiNLO approach, NLL resummation is achieved by means of a CKKW scale setting [75, 76] for the strong coupling factors associated with each QCD vertex, together with an appropriate factorization-scale choice and NLL QCD Sudakov form factors. These are applied to all internal and external lines corresponding to the underlying-Born skeleton of each event. In addition, improving the MiNLO resummation as described in ref. [35], we have obtained a fully inclusive description of HV production with NLO QCD accuracy in all phase space regions. In other words, besides providing HVj kinematic distributions that are NLO accurate and also finite when the hardest jet goes unresolved, the improved MiNLO predictions for $pp \rightarrow HVj$ are NLO accurate also for distributions in inclusive variables such as the rapidity or the transverse momentum of the HV pair.

All NLO QCD+EW predictions for $pp \rightarrow HVj$ presented in this paper, both at fixed order and including matching to the parton shower, are based on the MiNLO approach, which is applied to all contributions of NLO QCD and NLO EW type. Technically, the MiNLO Sudakov form factors and scale choices are implemented at the level of $pp \rightarrow HVj$ underlying-Born events that correspond to so-called \bar{B} terms in the POWHEG jargon.² Note that the MiNLO procedure resums only logarithms associated with soft and collinear QCD singularities that result from the presence of QCD radiation at Born level, while QED radiation is not present at Born level. Thus there is no need to introduce NLO EW effects in the MiNLO Sudakov form factors. This implies that, in contrast to the case of NLO QCD, the NLO EW corrections to $pp \rightarrow HVj$ do not need to be matched to the MiNLO form factors. In practice, for what concerns the EW corrections, the MiNLO procedure is applied in a way that is equivalent to Born level.

For observables where QCD radiation is integrated out, the MiNLO improved NLO EW contributions assume the form

$$\frac{d\sigma_{HVj}^{\text{MiNLO EW}}}{d\Phi_{HV}} = \int d\Phi_j \bar{B}_{HVj}^{\text{EW}}(\Phi_{HV}, \Phi_j) \Delta(k_T(\Phi_j)), \quad (2.2)$$

where Φ_{HV} and Φ_j denote the factorized phase spaces of the HV system and the jet, respectively. The term $\bar{B}_{HVj}^{\text{EW}}(\Phi_{HV}, \Phi_j)$ includes $\mathcal{O}(\alpha_{\text{EM}})$ corrections³ of virtual and real type, and the latter are integrated over the corresponding emission phase space. The MiNLO approach is implemented through an implicitly understood CKKW scale choice for the α_s term in $\bar{B}_{HVj}^{\text{EW}}$, and through the NLL Sudakov form factor $\Delta(k_T(\Phi_j))$ in eq. (2.2). For later convenience, together with the Sudakov form factor, we introduce a corresponding emission kernel $K(\Phi_j)$ that is formally related to Δ via

$$\Delta(p_T) = \exp \left[- \int d\Phi_j K(\Phi_j) \Theta(k_T(\Phi_j) - p_T) \right]. \quad (2.3)$$

²Real-emission events of NLO QCD and NLO EW type are related to underlying-Born events of type $pp \rightarrow HVj$ via FKS mappings [77].

³Since Born contributions are part of the usual QCD \bar{B} term, in the \bar{B}^{EW} term we only include $\mathcal{O}(\alpha_{\text{EM}})$ corrections.

In the following, based on the factorization properties of soft and collinear QCD radiation, encoded in the kernel $K(\Phi_j)$, and using the unitarity relation

$$\int d\Phi_j K(\Phi_j) \Delta(k_T(\Phi_j)) = 1, \tag{2.4}$$

we will argue that the inclusive **MinLO** predictions of eq. (2.2) are not only NLO QCD accurate, but also NLO EW accurate. More precisely, we will prove (in a schematic way) that

$$\frac{d\sigma_{HVj}^{\text{MinLO EW}}}{d\Phi_{HV}} = \frac{d\sigma_{HV}^{\text{NLO EW}}}{d\Phi_{HV}} + \mathcal{O}(\alpha_{\text{EM}} \alpha_s), \tag{2.5}$$

where

$$\frac{d\sigma_{HV}^{\text{NLO EW}}}{d\Phi_{HV}} = \bar{B}_{HV}^{\text{EW}}(\Phi_{HV}). \tag{2.6}$$

We first demonstrate the Born-level version of eq. (2.5), which corresponds to

$$\frac{d\sigma_{HVj}^{\text{MiLO}}}{d\Phi_{HV}} = \frac{d\sigma_{HV}^{\text{LO}}}{d\Phi_{HV}} + \mathcal{O}(\alpha_s), \tag{2.7}$$

where **MiLO** denotes the Born (or LO) version of the **MinLO** approach. The above identity can be written as

$$\int d\Phi_j B_{HVj}(\Phi_{HV}, \Phi_j) \Delta(k_T(\Phi_j)) = B_{HV}(\Phi_{HV}) + \mathcal{O}(\alpha_s), \tag{2.8}$$

where B_{HVj} and B_{HV} are the Born counterparts of the $\bar{B}_{HVj}^{\text{EW}}$ and \bar{B}_{HV}^{EW} terms in eqs. (2.2) and (2.3). The meaning of eqs. (2.7) and (2.8) is that the **MinLO** approach at Born level guarantees LO accuracy for observables that are inclusive with respect to the extra jet. In order to demonstrate this property, we split the $pp \rightarrow HVj$ Born term B_{HVj} into an IR divergent and a finite part,

$$B_{HVj}(\Phi_{HV}, \Phi_j) = B_{HV}(\Phi_{HV})K(\Phi_j) + B_{HVj}^{\text{fin}}(\Phi_{HV}, \Phi_j). \tag{2.9}$$

Here the singularities associated with QCD radiation in the soft and collinear limits are factorized⁴ into the $pp \rightarrow HV$ Born term times the NLL kernel $K(\Phi_j)$, while the B_{HVj}^{fin} remainder is free from singularities. Thus, upon integration over the jet phase space, the B_{HVj}^{fin} remainder yields only $\mathcal{O}(\alpha_s)$ suppressed contributions with respect to B_{HV} , while using the unitarity relation (2.4) it is easy to show that the singular term in eq. (2.9) leads to eq. (2.8).

Thanks to the fact that applying the **MinLO** approach to NLO EW contributions is largely equivalent to applying **MinLO** at Born level, the NLO EW accuracy property (2.5) can be proven along the same lines as for the LO accuracy property (2.7). As sole additional ingredient, the NLO EW proof requires certain factorization properties of soft and collinear QCD radiation. More precisely, the factorization properties of eq. (2.9) must hold also in the presence of EW corrections, i.e.

$$\bar{B}_{HVj}^{\text{EW}}(\Phi_{HV}, \Phi_j) = \bar{B}_{HV}^{\text{EW}}(\Phi_{HV})K(\Phi_j) + \bar{B}_{HVj}^{\text{EW, fin}}(\Phi_{HV}, \Phi_j). \tag{2.10}$$

⁴In this schematic derivation we assume a simple factorization of multiplicative type, while the factorization of initial-state collinear singularities takes the form of a convolution.

Here the remainder $\bar{B}_{HVj}^{\text{EW,fin}}$ should be free from QCD singularities, so that it yields only $\mathcal{O}(\alpha_s)$ -suppressed contributions relative to \bar{B}_{HV}^{EW} , when the extra jet is integrated out. Based on this natural assumption, in full analogy with the LO case, we easily arrive at

$$\begin{aligned} \int d\Phi_j \bar{B}_{HVj}^{\text{EW}}(\Phi_{HV}, \Phi_j) \Delta(k_T(\Phi_j)) &= \bar{B}_{HV}^{\text{EW}}(\Phi_{HV}) \int d\Phi_j K(\Phi_j) \Delta(k_T(\Phi_j)) + \mathcal{O}(\alpha_{\text{EM}}\alpha_s) \\ &= \bar{B}_{HV}^{\text{EW}}(\Phi_{HV}) + \mathcal{O}(\alpha_{\text{EM}}\alpha_s), \end{aligned} \tag{2.11}$$

which is equivalent to the hypothesis (2.5).

In summary, based on unitarity and factorization properties of QCD radiation, we expect that the improved MiNLO procedure applied to NLO QCD+EW matrix elements for $pp \rightarrow HVj$ should preserve its full QCD+EW accuracy when the jet is integrated out. As we will see, this conclusion is well supported by our numerical findings in sections 4–7. Nevertheless, due to the schematic nature of the presented derivations and related assumptions, the above conclusions should be regarded as an educated guess that deserves further investigation.

2.3 Sudakov approximation at NLO EW

In the Sudakov high-energy regime, where all kinematic invariants are of the same order and much larger than the electroweak scale, the NLO EW corrections are dominated by soft and collinear logarithms of Sudakov type. Based on the general results of refs. [49, 61] we have derived analytic expressions for the NLO EW corrections to HV and HVj production in NLL approximation. Details and scope of this approximation are discussed in appendix B.

The Sudakov approximation at NLO provides us with qualitative and quantitative insights into the origin of the dominant NLO EW effects. Moreover, it can be easily extended to the two-loop level [51, 52], thereby opening the door to approximate NNLO EW predictions based on the combination of exact NLO EW corrections with Sudakov logarithms at two loops. From the practical point of view, the Sudakov approximation at NLO permits to obtain the bulk of the EW virtual corrections at much higher computational speed as compared to an exact NLO EW calculation.

In section 5, we will assess the quality of the Sudakov approximation⁵ through a detailed comparison against exact NLO EW corrections. Finally, in appendix C, we show how the NLL EW approximation can be used in order to speed up the Monte Carlo integration, while keeping full NLO EW accuracy in the final predictions.

3 Technical aspects and setup of the simulations

3.1 The POWHEG BOX RES framework at NLO QCD+EW

The QCD+EW NLO calculations for HV and HVj production have been matched to parton showers using the POWHEG method. To this end, we used the recently-released version of the POWHEG BOX framework, called POWHEG BOX RES. The major novelty of this

⁵As explained in more detail in appendix B, the Sudakov approximation is applied only to the virtual part of EW corrections, while real QED radiation is always treated exactly.

new version is the resonance-aware approach [58], which guarantees a consistent treatment of intermediate resonances at NLO+PS level. This is achieved by generating the hardest radiation in a way that preserves the virtuality of resonances present at the underlying-Born level. At the same time, the resonance information can be passed on to the parton shower, which in turn preserves the virtuality of intermediate resonances of the hard process in subsequent emissions. This method was introduced in order to address the combination of NLO QCD corrections with parton showers in the presence of top-quark resonances. However, since it is based only on general properties of resonances and infrared singularities, the resonance-aware approach is applicable also to the combination of EW corrections with QED parton showers. In fact, this method has already been applied in the context of electroweak corrections in refs. [59, 60].

In the POWHEG BOX RES jargon [58], a radiated parton (or photon) can be associated to one or more “resonances” present in the process, or to the “production” part, if it cannot be associated to a particular resonance. The POWHEG BOX RES framework automatically finds all the possible so-called “resonance histories” for a given partonic process. For the processes at hand, considering QED radiation, only two resonance histories are detected: a production history, where the photon can be emitted by any quark (both in the initial and in the final state), and a vector-boson decay history, where the photon is radiated off a final-state charged lepton, and the virtuality of the intermediate vector boson needs to be preserved. Soft photons that are radiated from a W resonance are attributed either to the production subprocess or to the W decay, consistently with the virtualities of the quasi-resonant W propagators “before” and “after” the photon emission.

The treatment of QED radiation was first introduced in the POWHEG BOX for the calculation of the EW corrections to Drell-Yan processes [59, 60, 78, 79]. In this context, leptons were considered as massive particles, and QED subtraction in the POWHEG BOX was implemented accordingly. In the study at hand, leptons are treated as massless, and we have implemented the treatment of photon radiation off massless charged particles (both leptons and quarks). To this end, we have adapted the QCD soft and virtual counterterms already present in the POWHEG BOX to the QED case. Moreover, we have computed a new upper-bounding function for the generation of photon radiation with the highest-bid method, as described in ref. [29].

By default, in the POWHEG BOX RES framework, only the hardest radiation out of all singular regions is kept, before passing the event to shower Monte Carlo programs like *Pythia* or *Herwig*. In this way, for each event, at most one of the decaying resonances (or the production part of the process) includes an NLO-accurate radiation. Moreover, in case of combined QCD and EW corrections, QED emission occurs in competition with the QCD one. The POWHEG BOX RES uses the highest-bid method to decide what kind of radiation (QED or QCD, initial- or final-state) is generated. Due to the larger center-of-mass energy available in the production stage, initial-state radiation is enhanced with respect to final-state radiation, and since the QCD coupling is larger than the QED one, initial-state quarks tend to radiate gluons rather than photons. Thus, QED emission from the decay of a resonance would hardly be kept at the Les Houches event (LHE) level, and the QED radiation would mainly be generated by the shower Monte Carlo program.

The resonance-aware formalism implemented in the `POWHEG BOX RES` framework offers the opportunity to further improve the `POWHEG` radiation formula. With this improvement, first introduced in ref. [80], radiation from each singular region is generated and, instead of keeping only the hardest overall one, the hardest from each resonance is stored. As a result, the LHE file contains a radiated particle for each decaying resonance, plus possibly one emission from the production stage. In this way NLO+LL accuracy is ensured for radiation off each resonance. The subsequent shower from each resonance generated by the Monte Carlo shower program has to be softer than each corresponding `POWHEG` radiation.⁶ All NLO+PS results presented in this paper are based on this multiple-radiation scheme.

As a final remark, we note that in the `POWHEG BOX RES` framework both the HV and HVj processes can be computed at NLO or NLO+PS level with only QCD corrections, with only EW corrections, or with combined NLO QCD+EW corrections.⁷

3.2 OpenLoops tree and one-loop amplitudes

All needed amplitudes at NLO EW have been generated with `OpenLoops` [62, 63] and implemented in the `POWHEG BOX RES` framework through the general interface introduced in ref. [67]. Thanks to the recursive numerical approach of ref. [62] combined with the `COLLIER` tensor reduction library [81], or with `CUTTOOLS` [82], the `OpenLoops` program permits to achieve high CPU performance and a high degree of numerical stability. The amplitudes employed for the EW corrections in this paper are based on the recently achieved automation of EW corrections in `OpenLoops` [64–66].

Within `OpenLoops`, ultraviolet and infrared divergences are dimensionally regularized in D dimensions. However, all ingredients of the numerical recursion are handled in four space-time dimensions. The missing $(4 - D)$ -dimensional contributions, called R_2 rational terms, are universal and can be restored from process-independent effective counterterms [83–85]. The implementation of the corresponding Feynman rules for the complete EW Standard Model in `OpenLoops` is largely based on refs. [86–89]. Relevant contributions for HV and HVj production have been validated against independent algebraic results in $D = 4 - 2\epsilon$ dimensions. UV divergences at NLO EW are renormalized in the on-shell scheme [90] extended to complex masses [71].

3.3 Input parameters, scales choices and other aspects of the setup

In our $pp \rightarrow HV(+\text{jet})$ simulations at NLO QCD+EW, we have set the gauge-boson masses and widths to the following values [91]

$$\begin{aligned} M_Z &= 91.1876 \text{ GeV}, & M_W &= 80.385 \text{ GeV}, \\ \Gamma_Z &= 2.4955 \text{ GeV}, & \Gamma_W &= 2.0897 \text{ GeV}. \end{aligned} \tag{3.1}$$

The latter are obtained from state-of-the-art theoretical calculations. Assigning a finite width to the Higgs boson in the final state would invalidate EW Ward identities: we then

⁶This multiple-radiation mode can be activated by setting the flag `allrad` to 1 in the input file.

⁷The flag `qed.qcd` controls this behavior in the input file. The values it can assume are: 0, to compute only QCD corrections, 1, to compute only EW corrections or 2, for both.

consider the Higgs boson as on shell with $\Gamma_H = 0$ and set its mass to $M_H = 125$ GeV. The top-quark mass and width are set respectively to $m_t = 172.5$ GeV and $\Gamma_t = 1.5083$ GeV. All other quarks and leptons are treated as massless. In the EW corrections, the top-quark contribution enters only at loop level, the dependence of our results on Γ_t is thus completely negligible.

For the treatment of unstable particles we employ the complex-mass scheme [71, 72], where finite-width effects are absorbed into complex-valued renormalized masses

$$\mu_k^2 = M_k^2 - i \Gamma_k M_k \quad \text{for } k = W, Z, t. \quad (3.2)$$

The electroweak couplings are derived from the gauge-boson masses and the Fermi constant, $G_\mu = 1.16637 \times 10^{-5} \text{ GeV}^{-2}$, and the electromagnetic coupling is set accordingly to

$$\alpha_{\text{EM}} = \left| \frac{\sqrt{2} s_w^2 \mu_W^2 G_\mu}{\pi} \right|, \quad (3.3)$$

where μ_W^2 and the squared sine of the weak mixing angle

$$s_w^2 = 1 - c_w^2 = 1 - \frac{\mu_W^2}{\mu_Z^2}, \quad (3.4)$$

are complex-valued.⁸

The absolute values of the CKM matrix elements are set to

$$|V^{\text{CKM}}| = \begin{array}{ccc} & d & s & b \\ \begin{array}{l} u \\ c \\ t \end{array} & \begin{pmatrix} 0.97428 & 0.2253 & 0.00347 \\ 0.2252 & 0.97345 & 0.0410 \\ 0.00862 & 0.0403 & 0.999152 \end{pmatrix} & & \end{array}. \quad (3.5)$$

Our default set of parton-distribution functions (PDF) is the NNPDF2.3_as.0119_qed set [92], that includes QED contributions to the parton evolution and a photon density.⁹ The value of the strong coupling constant corresponding to this PDF set is $\alpha_s(M_Z) = 0.119$.

Finally, in HV production, the renormalization and factorization scales are set equal to the invariant mass of the HV pair at the underlying-Born level,

$$\mu_R = \mu_F = M_{HV}, \quad M_{HV}^2 = (p_H + p_{\ell_1} + p_{\bar{\ell}_2})^2, \quad (3.6)$$

where ℓ_1 and ℓ_2 are the final-state leptons, while in $pp \rightarrow HVj$ the improved MiNLO [34, 35] procedure is applied, and the scales are set accordingly.

Predictions at NLO+PS generated with the POWHEG method are combined with the Pythia 8.1 QCD+QED parton shower using the ‘‘Monash 2013’’ tune [94]. Effects due to hadronization, multi-particle interactions and underlying events are not considered in this paper.

⁸By default we use the G_μ scheme throughout. However, in the POWHEG BOX RES framework, there is the option to evaluate the virtual EW corrections using α_{EM} computed in the G_μ scheme, and use the Thomson value $\alpha_{\text{EM}}(0) = 1/137.035999$ in the evaluation of the contribution due to photon radiation.

⁹It corresponds to the PDF set 244800, in the LHAPDF6 [93] numbering scheme.

3.4 Physics objects and cuts in NLO+PS simulations

In the following we specify the definition of physics objects and cuts that are applied in the phenomenological NLO+PS studies presented in sections 5–7.

All leptonic observables are computed in terms of dressed leptons, which are constructed by recombining the collinear photon radiation emitted within a cone (in the (y, ϕ) plane) of radius $R_{\gamma\ell} = 0.1$ from charged leptons, and the recombined photons are treated as unresolved particles. Observables that depend on the reconstructed vector bosons are defined by combining the momenta of the dressed charged leptons and the neutrino associated with their decay. The latter is taken at Monte Carlo truth level.

Jets are constructed with `FastJet` using the anti- k_T algorithm [95, 96] with $R = 0.5$. The jet algorithm is applied in a democratic way to QCD partons and non-recombined photons, with the exception of photons that fulfill the isolation criterion of ref. [97] with a cone of radius $R_0 = 0.4$ and a maximal hadronic energy fraction $\epsilon_h = 0.5$. The hardest of such isolated photons is excluded from the jet algorithm and is treated as resolved photon.

The following standard Higgsstrahlung cuts are applied. For every dressed charged lepton we require

$$p_T^\ell \geq 25 \text{ GeV}, \quad |y^\ell| \leq 2.5. \quad (3.7)$$

In HW/HWj production, we also impose

$$\cancel{E}_T \geq 25 \text{ GeV}, \quad (3.8)$$

where \cancel{E}_T is the transverse momentum of the neutrino that results from the W -boson decay at Monte Carlo truth level. In HZ/HZj production, the invariant mass of the dressed-lepton pair is required to satisfy

$$60 \text{ GeV} \leq M^{\ell^+\ell^-} \leq 140 \text{ GeV}. \quad (3.9)$$

Besides these inclusive selection cuts, we also present more exclusive results in the boosted regime. In this case, we impose the following additional cuts on the transverse momentum of the Higgs and vector bosons

$$p_T^H \geq 200 \text{ GeV}, \quad p_T^V \geq 190 \text{ GeV}. \quad (3.10)$$

Such a selection of events with a boosted Higgs boson improves the signal-over-background ratio in the $H \rightarrow b\bar{b}$ decay channel.

4 Results for HV and HVj production at fixed NLO QCD+EW

In this section we present fixed-order NLO QCD+EW predictions for $pp \rightarrow HV$ and $pp \rightarrow HVj$ at 13 TeV. For HVj production the improved `miNLO` approach [34, 35] is applied. Higgs boson production in association with W and Z bosons is discussed in sections 4.1 and 4.2, respectively. Predictions based on exact NLO EW calculations (apart from photon-initiated contributions that have been neglected) are compared against the Sudakov NLL approximation (see appendix B), which includes virtual EW logarithms supplemented by an exact treatment of QED radiation.

The fixed-order results presented in this section are not subject to the cuts and definitions of section 3.4. No acceptance cut is applied, and differential observables are defined in terms of the momenta of the Higgs and vector bosons. The latter are defined in terms of the momenta of their leptonic decay products at the level of underlying-Born events, i.e. before the emission of NLO radiation. Photons and QCD partons are clustered in a fully democratic way using the anti- k_T algorithm with $R = 0.5$. Effectively this procedure corresponds to an inclusive treatment of QED radiation.

Besides total cross sections, we consider various differential distributions, focusing on regions of high invariant masses and transverse momenta, where EW corrections are enhanced by Sudakov logarithms. Such phase-space regions play an important role for experimental analyses of HV production in the boosted regime.

An in-depth validation of our fixed-order NLO EW results for HV production against the ones implemented in the public Monte Carlo program HAWK [44, 45], is presented in appendix A.

4.1 HW and HWj production

In this section we focus on NLO results for $pp \rightarrow HW$ and $pp \rightarrow HWj$. In table 1 we report inclusive NLO cross sections. In the case of HWj production, the improved MiNLO approach yields finite cross sections without imposing any minimum transverse momentum on the hardest jet. For comparison, we report also HWj MiNLO cross sections for the case where a minimum p_T of 20 GeV is required for the hardest jet. In this case, the MiNLO Sudakov form factor plays hardly any role, since it damps the cross section only at p_T of the order of a few GeV, i.e. far below the imposed cut. Thus, at fixed order, the MiNLO procedure only affects the choice of scales, as described in section 3.3. The EW corrections lower the inclusive NLO QCD cross section by roughly -7% for HW^\pm production and -5% for $HW^\pm j$ production, while they amount to only -2% when a resolved jet with $p_T^j > 20$ GeV is required in the $HW^\pm j$ calculation. Inclusive cross sections in the NLO QCD+NLL EW approximation differ by several percent from the exact NLO QCD+EW results. This is expected, since the NLL approximation is only valid in the high-energy regime.

In the following we investigate the impact of EW corrections and the validity of the NLL approximation in differential distributions for HW^- and HW^-j production. Results for $HW^+(j)$ production (not shown) are very similar.

In figure 3 we plot the invariant mass of the reconstructed HW^- pair, both for HW^- and HW^-j production. The three curves represent predictions at NLO QCD, NLO QCD+EW and in NLO QCD+NLL EW approximation. While EW corrections have a moderate impact on the total cross sections, they affect the tail of the M^{HW} distribution in a substantial way. At large M^{HW} we observe the typical Sudakov behavior, with increasingly large negative EW corrections that reach the level of -25% (-30%) for HV (HVj) production at 2 TeV. The Sudakov NLL approximation captures the bulk of these large EW corrections as expected. In the tail it agrees at the percent level with the exact result for both processes, while for moderate invariant masses it overestimates EW correction effects by up to 5%.

	HW^- NLO		HW^-j MiNLO	
selection	inclusive	inclusive	$p_T^j > 20$ GeV	
σ_{QCD} [fb]	59.25 ± 0.03	57.46 ± 0.02	26.720 ± 0.008	
$\sigma_{\text{QCD+EW}}$ [fb]	55.31 ± 0.02	55.3 ± 0.1	26.19 ± 0.04	
$\sigma_{\text{QCD+NLL EW}}$ [fb]	59.49 ± 0.01	59.6 ± 0.1	27.82 ± 0.04	
$\sigma_{\text{QCD+EW}}/\sigma_{\text{QCD}}$	0.93	0.96	0.98	
$\sigma_{\text{QCD+NLL EW}}/\sigma_{\text{QCD}}$	1.00	1.04	1.04	

	HW^+ NLO		HW^+j MiNLO	
selection	inclusive	inclusive	$p_T^j > 20$ GeV	
σ_{QCD} [fb]	93.24 ± 0.05	90.8 ± 0.2	42.2 ± 0.1	
$\sigma_{\text{QCD+EW}}$ [fb]	86.91 ± 0.02	86.2 ± 0.2	41.16 ± 0.09	
$\sigma_{\text{QCD+NLL EW}}$ [fb]	93.37 ± 0.02	93.0 ± 0.2	43.74 ± 0.09	
$\sigma_{\text{QCD+EW}}/\sigma_{\text{QCD}}$	0.93	0.95	0.98	
$\sigma_{\text{QCD+NLL EW}}/\sigma_{\text{QCD}}$	1.00	1.02	1.04	

Table 1. NLO total cross sections for HV (second column) and HVj (third and fourth column) production with $V = W^-$ (top) and $V = W^+$ (bottom), at a center-of-mass energy of $\sqrt{s} = 13$ TeV, at NLO QCD, NLO QCD+EW, and in the NLO QCD+NLL EW approximation. The HVj cross sections are based on the improved MiNLO procedure (third and fourth column). The effect of a cut of $p_T^j > 20$ GeV on the transverse momentum of the hardest jet in HVj production is shown in the last column. Listed uncertainties are due to Monte Carlo integration.

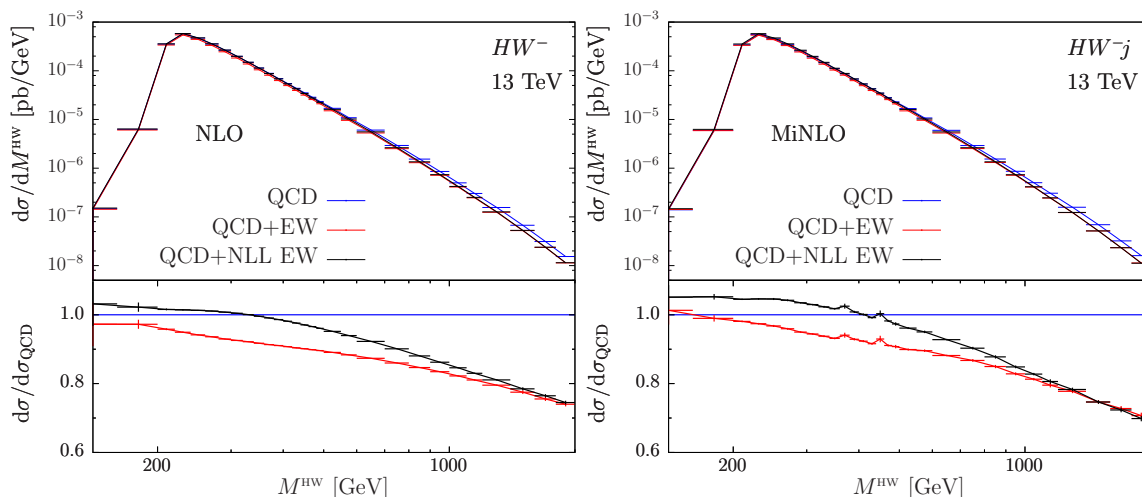


Figure 3. NLO predictions for the invariant mass of the HW^- pair in HW^- (left) and HW^-j (right) production. Shown are predictions at NLO QCD (blue), NLO QCD+EW (red) and at NLO QCD+NLL EW (black). In the lower panel we plot the ratio with respect to NLO QCD.

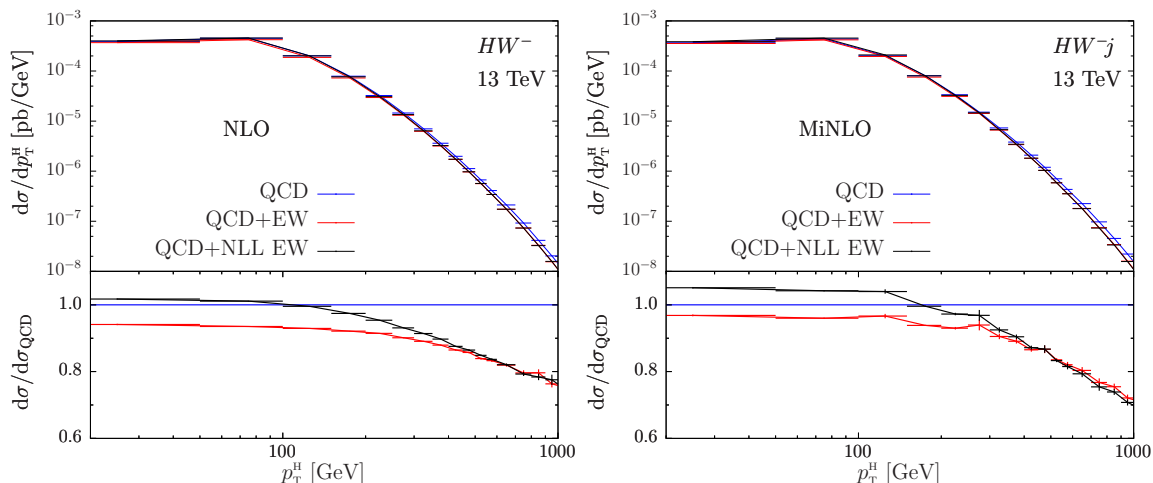


Figure 4. NLO predictions for the transverse momentum of the Higgs boson in HW^- (left) and HW^-j (right) production. Predictions and labels as in figure 3.

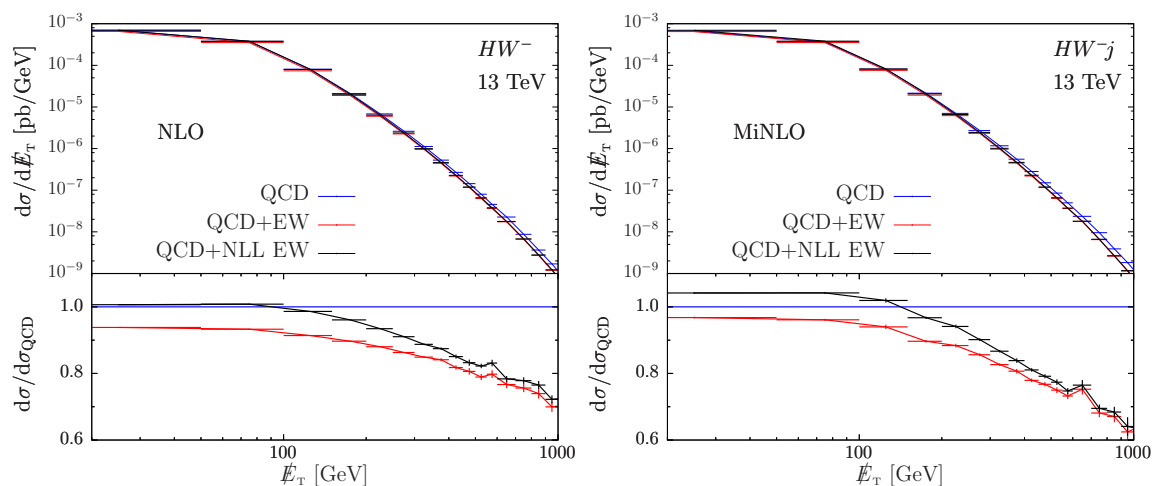


Figure 5. NLO predictions for the missing transverse momentum in HW^- (left) and HW^-j (right) production. Predictions and labels as in figure 3.

In figure 4 we investigate the transverse momentum of the Higgs boson. Also in this case EW corrections become negative and large in the tail, exceeding -20% in the TeV region. For both processes the Sudakov approximation agrees at the percent level with exact NLO EW results for $p_T^H > 300$ GeV.

The EW corrections have a sizable impact also on the missing transverse momentum distribution, shown in figure 5. Size and shape of these corrections are very similar to the ones observed for the Higgs boson p_T distribution.

In figure 6 we present HW^-j predictions for the distribution in the p_T of the leading jet. At low p_T^j (left plot) the MiNLO Sudakov form factor damps soft and collinear singularities at zero transverse momentum yielding finite cross sections below the Sudakov peak, which is located around 3 GeV. Concerning EW effects, the NLL approximation converges to the exact NLO results already for values of p_T^j around 200 GeV. In the region of mod-

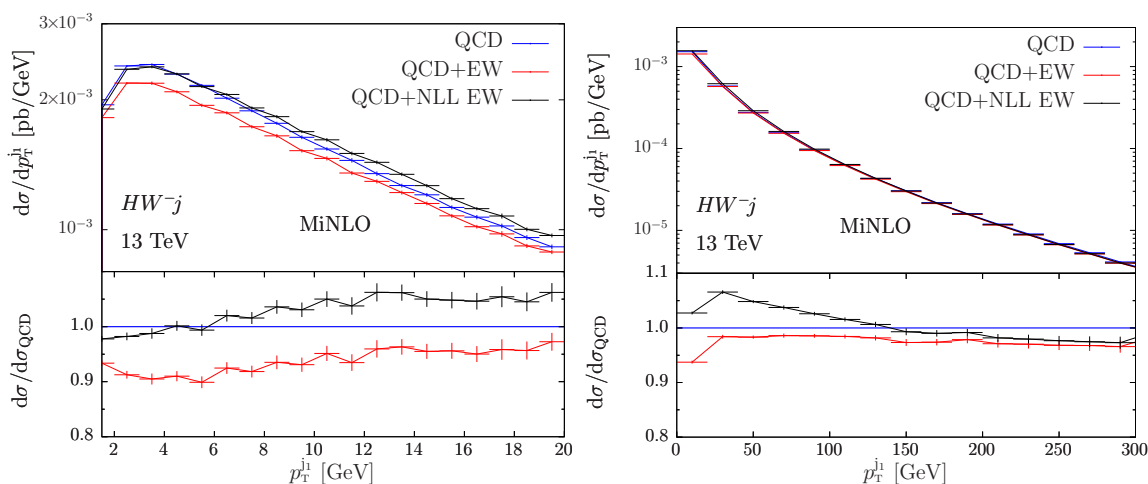


Figure 6. NLO predictions for the transverse momentum of the leading jet in HW^-j production for different p_T^{j1} ranges. The left plot corresponds to the first bin of the right plot. Predictions and labels as in figure 3.

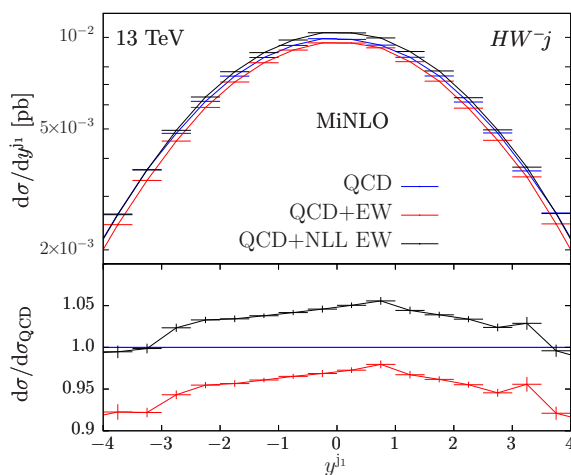


Figure 7. NLO predictions for the rapidity of the leading jet in HW^-j production. Same curves and labels as in figure 3.

erate transverse momentum, NLO EW corrections are nearly constant, and in the limit of vanishing jet- p_T they converge towards an EW K -factor that is very close to the one of the NLO QCD+EW calculation for the inclusive $pp \rightarrow HW^-$ cross section (see table 1). This observation is consistent with the theoretical considerations presented in section 2.2, namely with the fact that EW corrections are insensitive to soft and collinear QCD radiation, and that MiNLO predictions for HVj production preserve NLO QCD+EW accuracy when the extra jet is integrated out. In fact, in the inclusive distributions of figures 3–5, we observe that the EW corrections obtained from HW^- and HW^-j calculations are very similar, with small differences that can be attributed to NNLO effects.

Finally, in figure 7 we see that the EW corrections affect the rapidity of the leading jet in a rather uniform way over the whole phase space. We have observed a similar behavior of EW corrections in several angular distributions.

	<i>HZ</i> NLO		<i>HZj</i> MiNLO	
selection	inclusive	inclusive	$p_T^1 > 20$ GeV	
σ_{QCD} [fb]	25.551 ± 0.005	24.801 ± 0.009	11.720 ± 0.004	
$\sigma_{\text{QCD+EW}}$ [fb]	24.382 ± 0.008	24.59 ± 0.07	12.22 ± 0.02	
$\sigma_{\text{QCD+NLL EW}}$ [fb]	25.457 ± 0.008	25.84 ± 0.07	12.69 ± 0.01	
$\sigma_{\text{QCD+EW}}/\sigma_{\text{QCD}}$	0.95	0.99	1.04	
$\sigma_{\text{QCD+NLL EW}}/\sigma_{\text{QCD}}$	1.00	1.04	1.08	

Table 2. NLO total cross sections for *HZ* (second column) and *HZj* (third and fourth column) production at a center-of-mass energy of $\sqrt{s} = 13$ TeV. Predictions and labels as in table 1.

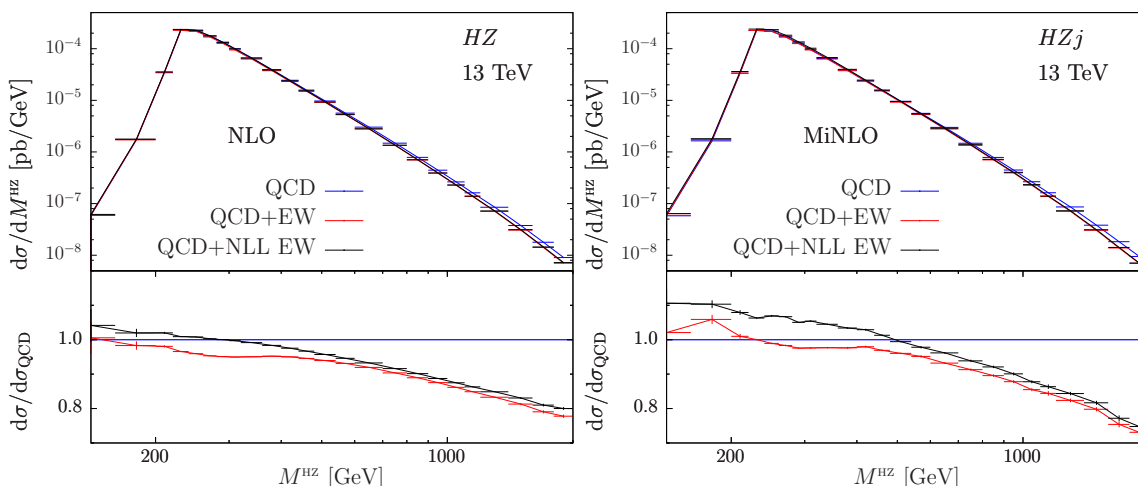


Figure 8. NLO predictions for the invariant mass of the *HZ* pair in *HZ* (left) and *HZj* (right) production. The three curves represent the QCD, QCD+EW and the QCD+NLL EW predictions. The lower panel displays ratios with respect to NLO QCD.

4.2 *HZ* and *HZj* production

In line with the discussion of *HW* and *HWj* production, we present in this section fixed-order results for *HZ* and *HZj* production. In table 2 we collect the inclusive cross sections at NLO QCD, NLO QCD+EW and NLO QCD+NLL EW. The EW corrections decrease the total NLO QCD cross section for *HZ* production by about 4%, and by about 1% for inclusive *HZj* production. In the presence of a jet threshold of 20 GeV, the EW corrections are positive and amount to about 4%.

In figures 8 and 9 we show distributions of the invariant mass of the reconstructed *HZ* pair and of the transverse momentum of the Higgs boson. Similarly as for *HW(j)* production, the inclusion of EW corrections is essential in the tails of these distributions, where the NLL Sudakov approximation agrees well with the exact NLO EW predictions.

5 Results for *HV* production at NLO+PS QCD+EW

In this section we present NLO QCD+EW predictions for *HV* production completed by the Pythia 8.1 QCD+QED parton shower using the “Monash 2013” tune [94]. All predictions

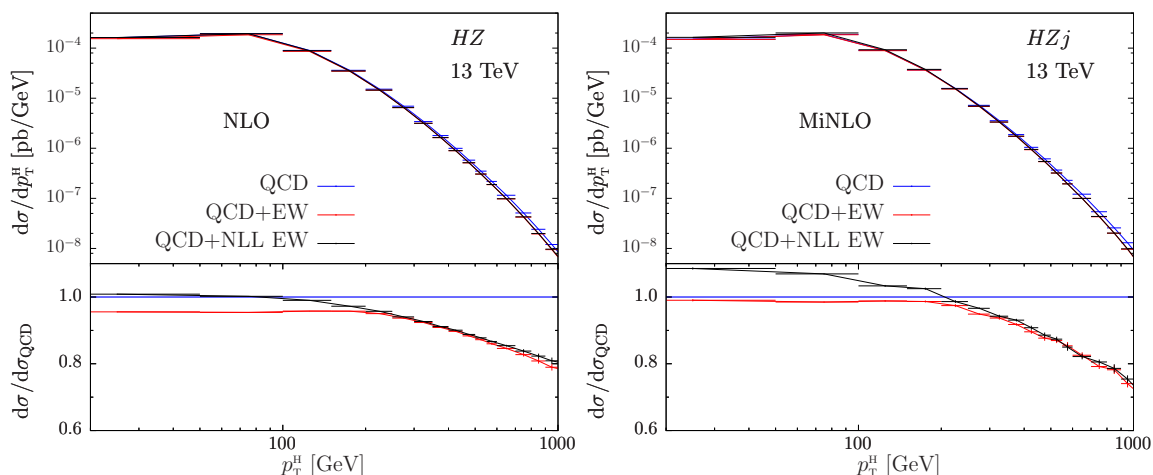


Figure 9. NLO predictions for the Higgs boson transverse momentum in HZ (left) and HZj (right) production. Same curves and labels as in figure 8.

are subject to the cuts and physics object definitions specified in section 3.4, and NLO EW corrections are treated exactly throughout, except for photon-initiated processes, that have been neglected. The NLL Sudakov approximation is only used in order to speed up the Monte Carlo integration, as detailed in appendix C.

In section 5.1 we compare predictions at fixed-order NLO QCD+EW against corresponding predictions at the level of Les Houches events, which include only the hardest emission generated in the POWHEG BOX RES framework, and at NLO+PS level, where the full QCD+QED parton shower is applied. The effect of EW corrections is studied in section 5.2 in the case of fully showered NLO+PS simulations.

By default, at NLO+PS level, the full QCD+QED parton shower is applied, both for NLO QCD+EW and for pure NLO QCD simulations. Occasionally, we also present NLO QCD simulations with a pure QCD shower, where QED radiation is switched off. Such predictions are labeled “QCD (no QED shower)”.

The consistent combination of the NLO radiation to the parton shower requires the vetoing of shower emissions that are harder than the radiation generated in the POWHEG BOX RES framework. Since no standard interface is available in a multi-radiation scheme, we have implemented a dedicated veto procedure on the Pythia 8.1 showered events, as described in appendix D. This veto procedure is applied in case of NLO QCD+EW simulations. Instead, in case of NLO QCD simulations combined with the Pythia 8.1 QCD+QED shower, only QCD radiation is restricted by the POWHEG BOX RES hardest scale, while arbitrarily hard QED radiation can be generated by the shower.

We have verified that inclusive cross sections at NLO+PS QCD and NLO+PS QCD+EW agree within statistical uncertainties with the corresponding fixed-order results reported in tables 1 and 2. Thus, in the following we will focus on differential distributions.

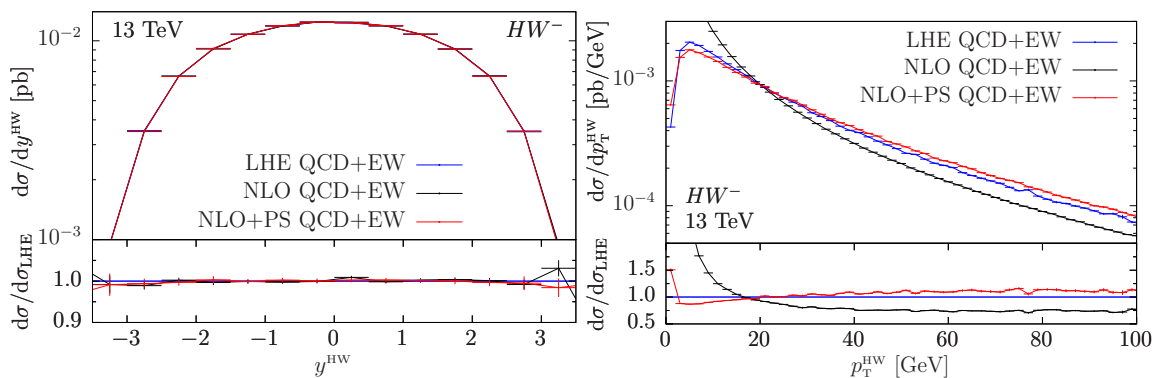


Figure 10. Rapidity (left) and transverse-momentum distribution (right) of the HW^- pair in HW^- production. Results at NLO QCD+EW are compared at fixed order, at the level of Les Houches events (LHE), and including also the full QCD+QED parton shower of Pythia 8.1 (NLO+PS). In the ratio plot results are normalized with respect to the LHE level prediction.

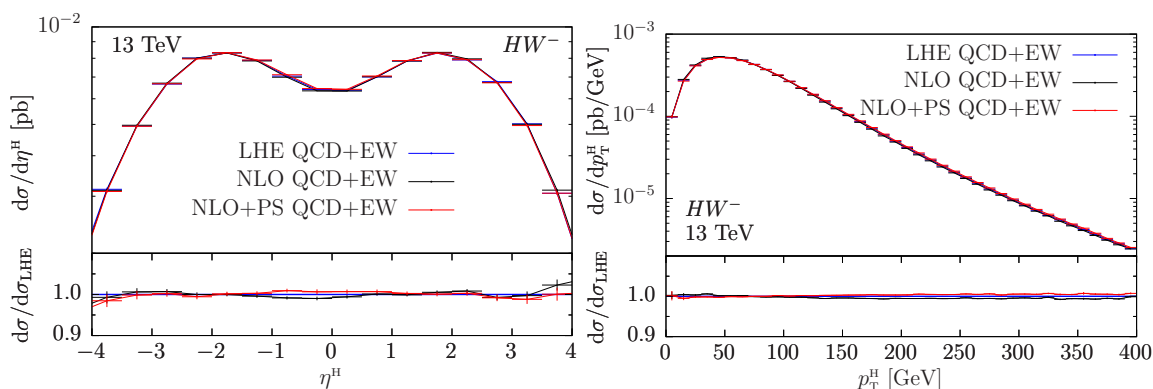


Figure 11. Pseudorapidity (left) and transverse-momentum distribution (right) of the Higgs boson in HW^- production. Same curves and labels as in figure 10.

5.1 From fixed NLO QCD+EW to NLO+PS QCD+EW

In this section we compare NLO QCD+EW predictions at fixed order with NLO+PS ones at LHE level and completed with the Pythia 8.1 shower. Since the various Higgsstrahlung processes behave in a very similar way, we will focus on HW^- production.

In figure 10 we plot the rapidity of the reconstructed HW^- pair, which is NLO accurate, and its transverse momentum, which is only LO accurate. Due to the inclusiveness of the rapidity of the HW^- pair, we find, as expected, very good agreement, within the integration errors, among the three predictions. The fixed-order curve for the transverse momentum displays the typical divergent behavior at low p_T . At LHE level, instead, the divergence is tamed by the Sudakov form factor. The effect of the parton shower is modest in the tail of this distribution, while at low p_T it slightly shifts the position of the Sudakov peak.

In figure 11 we plot the pseudorapidity and the transverse momentum of the Higgs boson: thanks to the inclusiveness of this variable, we find again very good agreement among the three predictions.

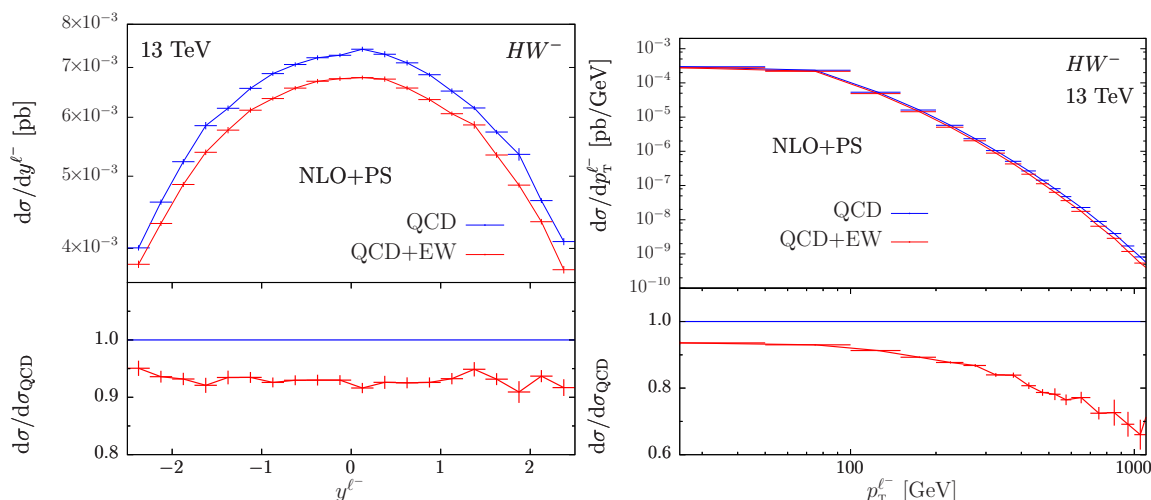


Figure 12. NLO+PS predictions for the distributions in the rapidity (left) and the transverse momentum (right) of the charged dressed lepton in HW^- production. Comparison between the full QCD+EW results and the QCD ones after the Pythia 8.1 QCD+QED shower.

5.2 Impact of the EW corrections in NLO+PS events

In this section we investigate EW correction effects at the level of fully showered NLO+PS predictions.

In figure 12 we show the rapidity (left) and the transverse momentum (right) of the charged dressed lepton in HW^- production. In the rapidity distribution, the impact of NLO EW effects is constant and amounts to about -7% . The shape of the p_T distribution, instead, changes drastically due to EW Sudakov logarithms in the high- p_T region, where differences with respect to the pure QCD predictions reach -30% around 1 TeV.

In figure 13 we plot the transverse mass of the reconstructed W^- boson

$$M_T^W = \sqrt{2 p_T^\ell \cancel{E}_T (1 - \cos \Delta\phi)}, \quad (5.1)$$

where $\Delta\phi$ is the azimuthal angle between the charged lepton and the missing transverse momentum. Similarly, as for the lepton rapidity, the EW corrections do not change the shape, but lower the differential cross section by about 7% with respect to the pure QCD corrections. If no QED shower is activated when Pythia 8.1 showers QCD-corrected events, the curve that is obtained is very similar to the QCD one, i.e. the impact of the QED shower is small for this distribution and no radiative tail can be observed.

In figure 14 we show the rapidity and the transverse momentum of the Higgs boson in the boosted regime, as defined by the cuts of eq. (3.10). The EW corrections have a constant negative impact around 10% on the rapidity distribution, and reach -25% around 1 TeV. Similar conclusions can be drawn for the rapidity and transverse momentum of the W^- boson.

We conclude this section by presenting kinematic distributions for HZ production in figures 15–17. In figure 15 we show the distribution in the rapidity and the transverse momentum of the dressed electron. The EW corrections give a constant contribution of

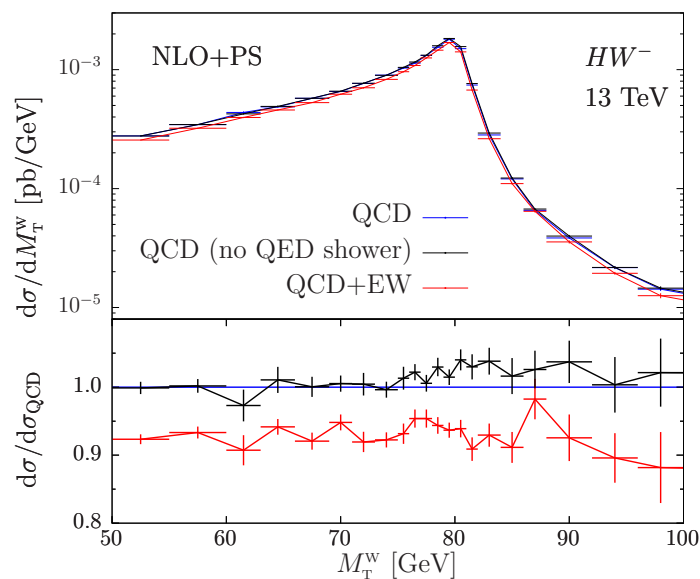


Figure 13. NLO+PS predictions for the transverse mass of the reconstructed W^- boson in HW^- production. Same curves and labels as in figure 12. To illustrate the effect of the QED shower, we also show results obtained by showering QCD-corrected events with the QED shower switched off in Pythia 8.1 (“no QED shower”).

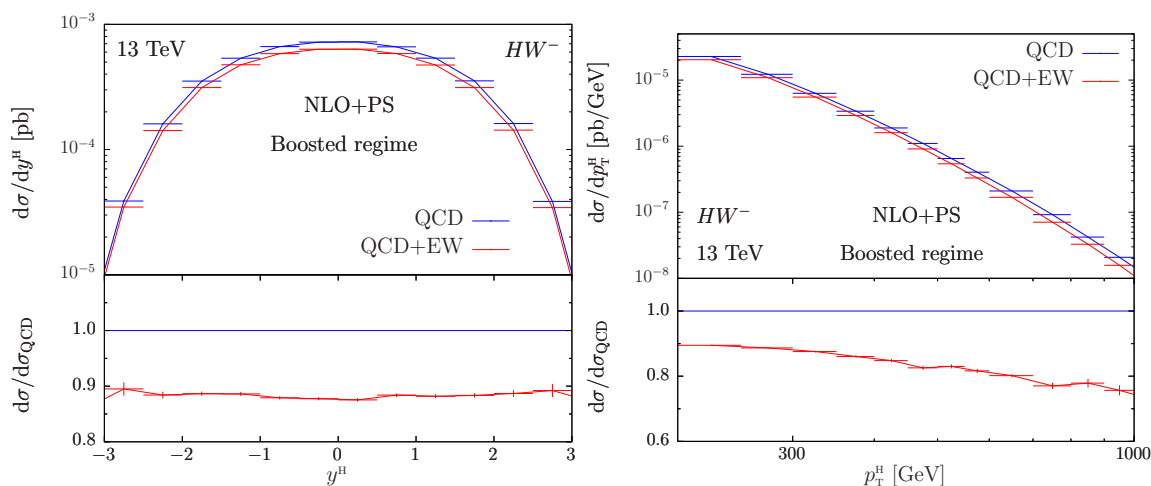


Figure 14. NLO+PS predictions for the rapidity (left) and the transverse-momentum distribution (right) of the Higgs boson in the boosted regime of eq. (3.10) for HW^- production. Same curves and labels as in figure 12.

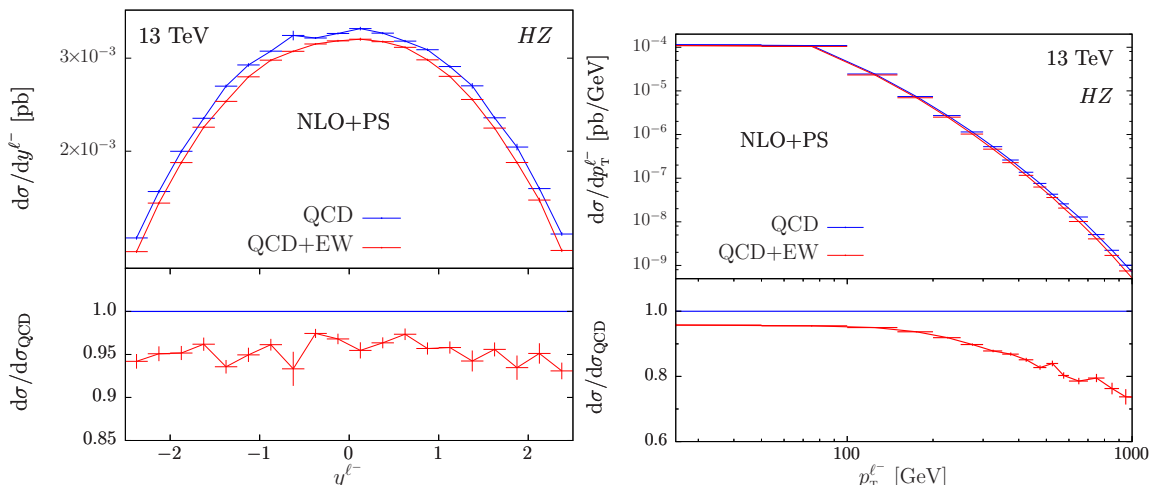


Figure 15. NLO+PS predictions for the rapidity (left) and for the transverse momentum (right) of the dressed electron in HZ production. Comparison between the full QCD+EW results and the QCD ones after the Pythia 8.1 QCD+QED shower.

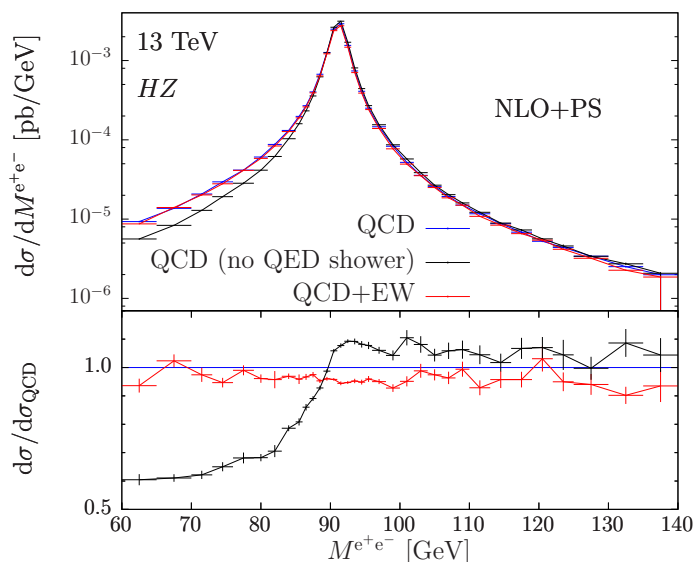


Figure 16. NLO+PS predictions for the invariant mass of the reconstructed leptonic pair in HZ production. Same curves and labels as in figure 15. For comparison, the result obtained by showering the QCD-corrected events without QED shower in Pythia 8.1 is also plotted.

about -5% in the plotted rapidity range, while in the high-energy tail of the p_T distribution the EW corrections decrease the differential cross section by roughly 30% due to Sudakov logarithms.

In figure 16 we plot the invariant mass of the reconstructed leptonic pair in the region around the Z resonance. In spite of the fact that the shape of the Z resonance is known to receive very large $\mathcal{O}(\alpha_{EM})$ radiative corrections (see e.g. refs. [79, 98]), NLO EW effects turn out to be almost constant and as small as -5% when we compare showered NLO+PS predictions at NLO QCD+EW versus NLO QCD. This is due to the fact that the bulk

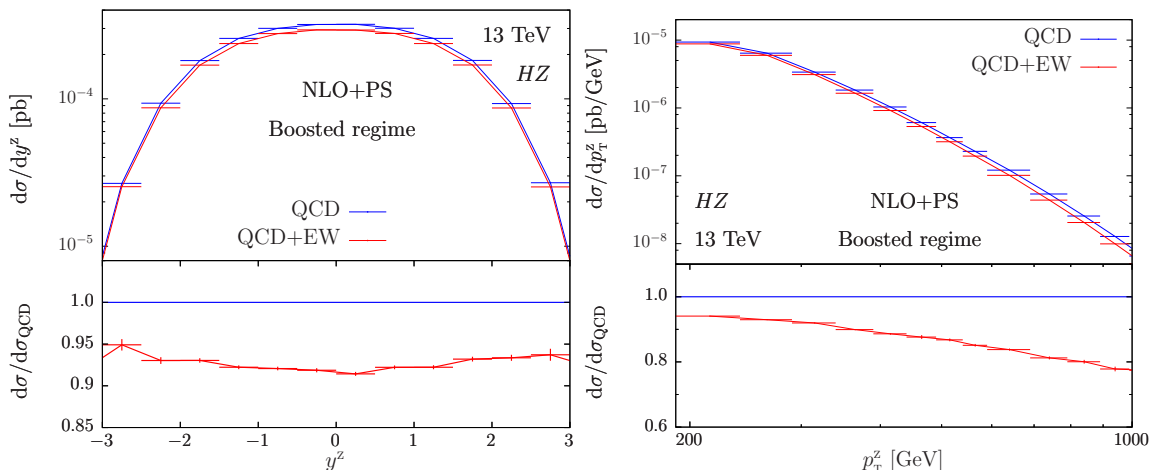


Figure 17. NLO+PS predictions for the rapidity (left) and the transverse momentum (right) of the Z boson for HZ production in the boosted regime. Same curves and labels as in figure 15.

of the $\mathcal{O}(\alpha_{\text{EM}})$ radiation is correctly described by the QED shower in `Pythia 8.1`. The importance of $\mathcal{O}(\alpha_{\text{EM}})$ radiation becomes evident when we switch off the QED shower (“no QED shower”) in the NLO QCD simulation. This results in a radiative tail with distortions of up to 40% in the region below the Z peak.

In HZ production, the momentum of the vector boson can be fully reconstructed. Thus, in figure 17 we display the rapidity and transverse-momentum distributions of the Z boson in the boosted regime, as defined by the cuts of eq. (3.10). These results are very similar to the ones obtained for the Higgs boson in HW^- production in figure 14. While EW corrections have a constant impact of about -8% on the rapidity distribution, the tail of the p_T distribution is dominated by large negative EW Sudakov logarithms, and we observe differences with respect to the pure QCD result of the order of -25% for $p_T \sim 1$ TeV.

6 Results for HVj production at NLO QCD+EW with `MinLO+PS`

In this section we study $pp \rightarrow HVj$ at NLO+PS accuracy in the `MinLO` approach, denoted in the following as `MinLO+PS`. Similarly as in the previous section, in section 6.1 we first compare NLO QCD+EW predictions obtained with `MinLO` at fixed order against corresponding results at the LHE level or including also the full QCD+QED parton shower. The effect of EW corrections is studied in section 6.2 in the case of fully showered `MinLO+PS` simulations. The cuts and physics object definitions of section 3.4 are applied throughout, and we do not impose any cut that requires the presence of jets.

6.1 From fixed-order `MinLO` to `MinLO+PS` at NLO QCD+EW

In figure 18 we analyze the rapidity and the transverse momentum of the reconstructed HW^- system. As a result of the `MinLO` prescription, the rapidity distribution, as well as any other inclusive observable, is finite. For the rapidity distribution we observe that the three predictions are very close to each other. At variance with figure 10, here fixed-order

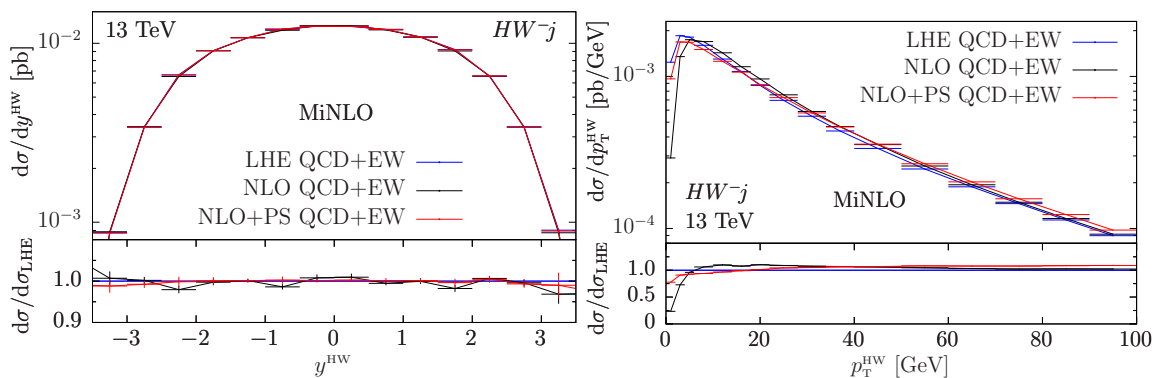


Figure 18. Distributions in the rapidity (left) and transverse momentum (right) of the reconstructed HW^- pair. Improved MiNLO results for $pp \rightarrow HW^-j$ at NLO QCD+EW are compared at fixed order, at the level of Les Houches events (LHE), and including also the full QCD+QED parton shower of Pythia 8.1 (NLO+PS). In the ratio plot results are normalized with respect to the LHE level prediction.

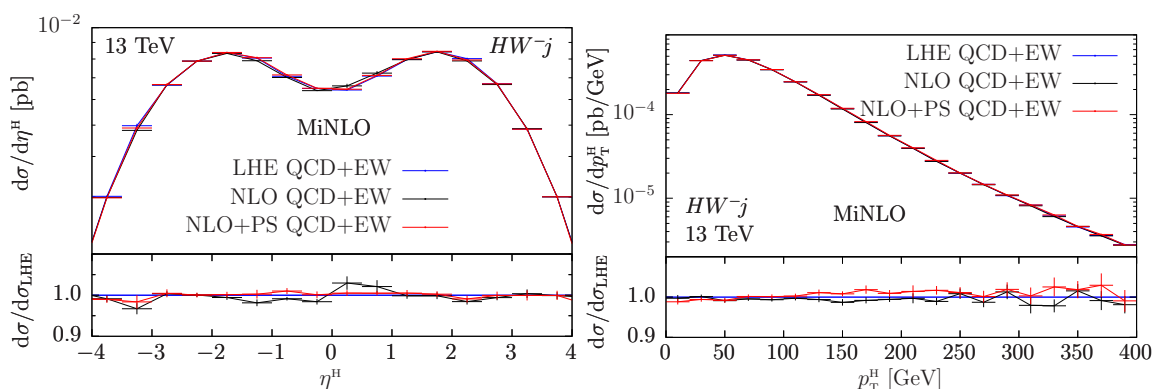


Figure 19. Distributions in the pseudorapidity (left) and transverse momentum (right) of the Higgs boson in the MiNLO improved $pp \rightarrow HW^-j$ simulation. Same curves and labels as in figure 18.

predictions for the p_T distribution are finite at small p_T , since soft and collinear divergences are suppressed by the MiNLO Sudakov form factors. Moreover, the NLO accuracy in the spectrum of the HW^- system leads to an improved agreement between fixed-order and NLO+PS results.

In figure 19 we show the pseudorapidity and the transverse-momentum distributions of the Higgs boson, finding again very good agreement among the three predictions.

We refrain from presenting results for HW^+j and HZj production since they behave qualitatively very similar as the results shown here for HW^-j production.

6.2 Impact of EW corrections at MiNLO+PS level

The impact of EW corrections at the level of fully showered MiNLO+PS predictions for HW^-j production is illustrated in figures 20–22.

For the distributions in the rapidity and transverse momentum of the Higgs boson in the boosted regime (figure 20) we find that the EW corrections induced by the boosted

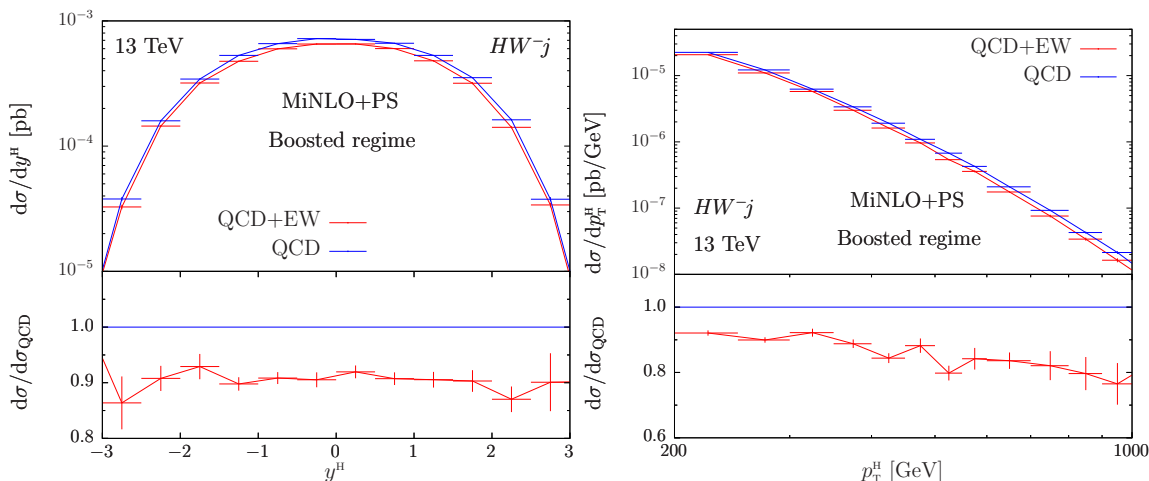


Figure 20. MiNLO+PS predictions for the rapidity (left) and transverse-momentum distribution (right) of the Higgs boson in the boosted regime, in HW^-j production. Comparison between the full QCD+EW results and the QCD ones after the Pythia 8.1 QCD+QED shower.

cut, $p_T^H \geq 200$ GeV, are nearly independent of y^H and around -10% , while they grow up to -20% and beyond when p_T^H enters the TeV regime. These results closely agree with the corresponding ones shown in figure 14 for the NLO+PS simulation of inclusive HW^- production. Consistently with the fixed-order findings discussed in section 4, also this observation supports the theoretical considerations of section 2.2, where we have argued that MiNLO improved predictions for HVj production should preserve NLO QCD+EW accuracy when the extra jet is integrated out. Also other inclusive observables, such as the distribution in the missing transverse momentum shown in figure 21, confirm this observation.

The EW corrections to the leading-jet p_T distribution, shown in figure 22, do not feature the standard Sudakov behavior. In this distribution, EW effects remain rather small, at the level of -5% , in the entire plotted range, i.e. from very low jet- p_T up to 400 GeV. This is not surprising, since a similar “non-Sudakov” behavior for inclusive jet spectra was already observed in refs. [65, 74] for the case of $V+jets$ production. Another important feature of figure 22 is that EW corrections are nearly constant in the region where the jet p_T approaches zero. Again, this confirms the considerations made in section 2.2 regarding the factorization of EW corrections in the presence of soft or collinear QCD radiation, and the NLO QCD+EW accuracy of inclusive MiNLO simulations. To be more precise, in the left panel of figure 22 we see that EW corrections effects are nearly constant at small p_T with the exception of the first bin. This effect can be attributed to photonic contributions to the jet transverse momentum, and to the fact that the Sudakov peak associated with the damping of QCD radiation is located well above the one associated with the damping of QED radiation. This mismatch tends to enhance the relative importance of QED radiation in the region between the QCD and QED Sudakov peaks. In any case, this effect cancels upon integration over the soft region of the jet spectrum. Thus, it should not spoil the expected NLO QCD+EW accuracy of inclusive MiNLO predictions.

We conclude this section by discussing the impact of NLO EW effects in HZj production, illustrated in figures 23 and 24. The distribution in the Z -boson p_T in the boosted

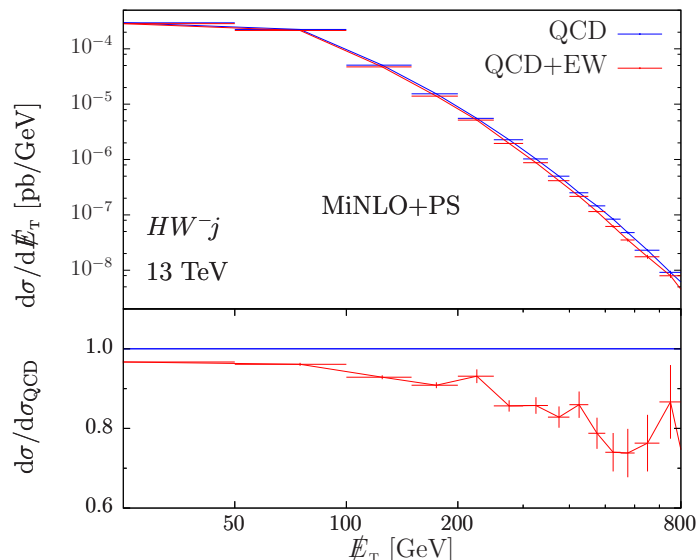


Figure 21. MiNLO+PS predictions for the missing transverse momentum in HW^-j production. Same curves and labels as in figure 20.

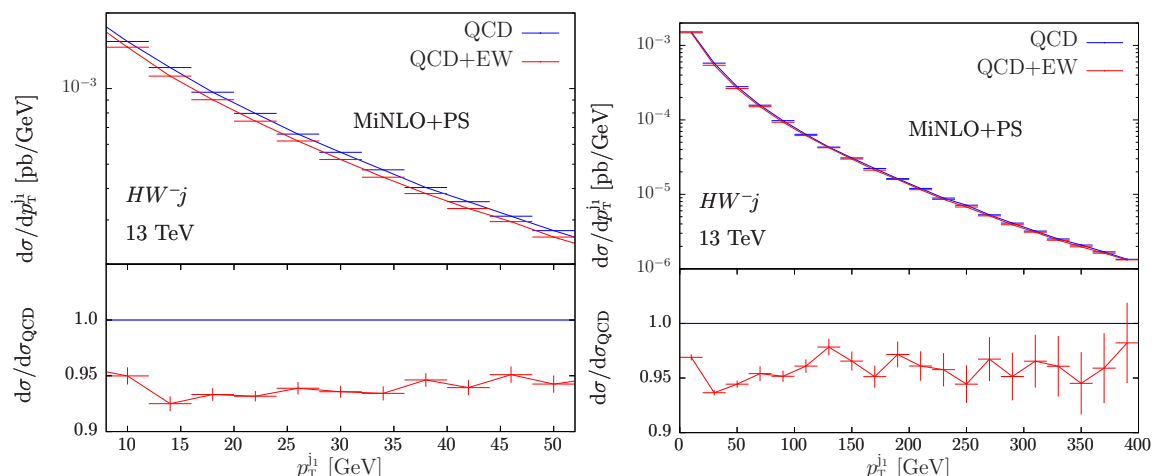


Figure 22. MiNLO+PS predictions for the transverse momentum of the leading jet in two transverse-momentum ranges in HW^-j production. Same curves and labels as in figure 20.

regime (figure 23) features the typical Sudakov EW behavior, with negative EW corrections that exceed -20% in the tail. In the leading-jet p_T distribution (figure 24) we observe relatively small and rather constant EW corrections. Both distributions behave similarly as the corresponding distributions for HW^-j production.

We refrain from showing further plots for HZj or HW^+j production, since EW correction effects are quite similar to the ones already discussed.

7 Comparison between the HV and HVj generators

In this section, we discuss and compare NLO+PS predictions for $pp \rightarrow HV$ against MiNLO+PS predictions for $pp \rightarrow HVj$, both at NLO QCD+EW accuracy. A similar com-

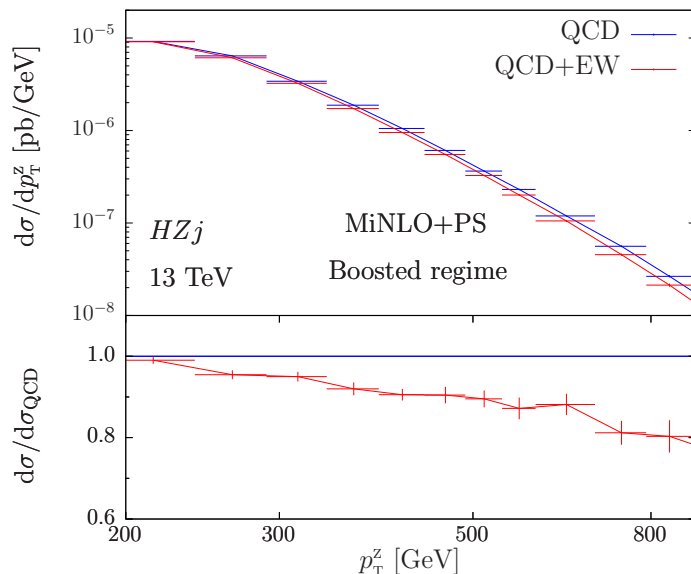


Figure 23. MiNLO+PS predictions for the transverse momentum of the Z boson for HZj production in the boosted regime. Same curves and labels as in figure 20.

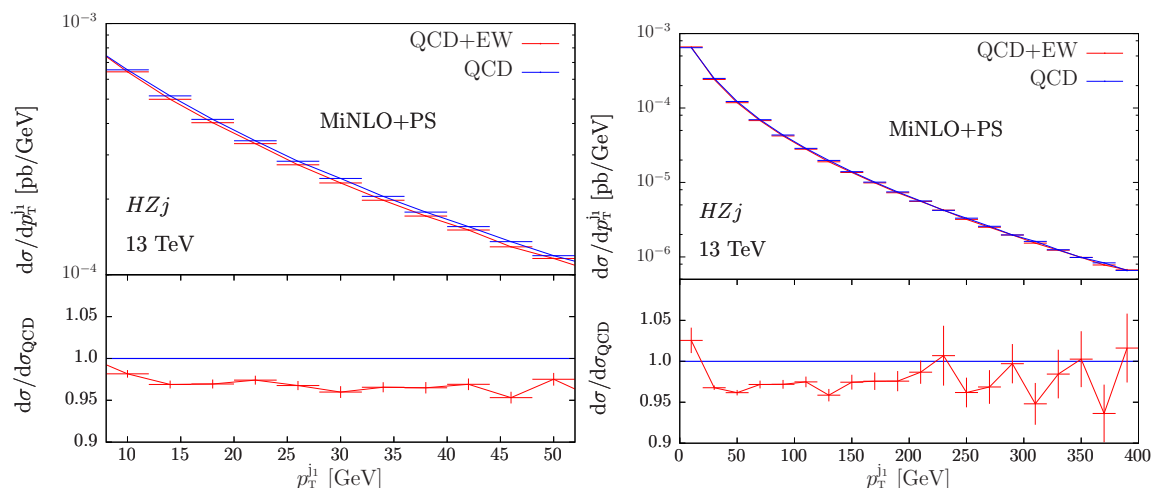


Figure 24. MiNLO+PS predictions for the transverse momentum of the leading jet in two transverse-momentum ranges in HZj production. Same curves and labels as in figure 20.

parison at NLO QCD accuracy was presented in ref. [33]. Since the various Higgsstrahlung processes behave in a very similar way, we will focus on the associated HW^- production. The comparison between the HV and HVj generators is motivated by the fact that the improved MiNLO prescription [35] applied to HVj production provides NLO accuracy also for inclusive HV quantities, i.e. for observables where the associated jet is not resolved. While this is well known at NLO QCD level, in section 2.2 we have argued that also NLO EW accuracy should be preserved when the jet is integrated out.

We also study the dependence of our results on scale variations. To this end we apply standard seven-point variations obtained by multiplying the central value of the

renormalization and factorization scales μ_R and μ_F , defined in eq. (3.6) for HV production, by the factors K_R and K_F , respectively, chosen among the seven pairs

$$(K_R, K_F) = \left(\frac{1}{2}, \frac{1}{2}\right), \left(\frac{1}{2}, 1\right), \left(1, \frac{1}{2}\right), (1, 1), (2, 1), (1, 2), (2, 2). \quad (7.1)$$

Scale-variation bands in the following plots are based on the envelope of the seven-point variations. In HVj production, where the scale setting is based on the improved MiNLO prescription, the scaling factors of eq. (7.1) are applied to the coupling constants at each interaction vertex and to the scale entering the Sudakov form factor.

For the fully inclusive NLO+PS and MiNLO+PS cross sections at NLO QCD+EW we find

$$\begin{aligned} \sigma_{HW^-}^{\text{NLO+PS}} &= 55.29_{-0.74}^{+0.80} \text{ fb}, & \sigma_{HW^-j}^{\text{MiNLO+PS}} &= 55.25_{-2.57}^{+1.25} \text{ fb}, \\ \sigma_{HZ}^{\text{NLO+PS}} &= 24.41_{-0.38}^{+0.27} \text{ fb}, & \sigma_{HZj}^{\text{MiNLO+PS}} &= 24.9_{-1.1}^{+0.6} \text{ fb}, \end{aligned} \quad (7.2)$$

where uncertainties correspond to scale variations. These results are well consistent, within statistics, to the corresponding ones reported in tables 1 and 2 at fixed-order NLO. Moreover, it turns out that, in the presence of EW corrections, cross sections obtained from HV and HVj simulations agree at the one-percent level, confirming again the expectation of inclusive NLO QCD+EW accuracy for MiNLO improved HVj simulations.

Scale variations are in general larger in HVj with respect to HV production. This is due to the fact that, in standard POWHEG BOX simulations, the scale associated to the emission of the hardest jet is kept fixed at the corresponding transverse momentum, while scale variations are applied only at the level of the so-called \bar{B} term, where QCD and QED radiation are integrated out. For this reason, scale variations in MiNLO+PS simulations provide a more realistic estimate of scale uncertainties associated with QCD radiation.

Figures 25–27 display differential distributions subject to the cuts of section 3.4. Red bands correspond to scale variations for HV and HVj production. We do not show the statistical uncertainties associated to the integration procedure on these bands, since they are much smaller than the width of the bands. When plotting instead the blue curves for the distributions computed at the central scales, we display the statistical uncertainties of the integration procedure as an error bar. The plots on the left-hand side show the uncertainty band for the HV process, while the ones on the right-hand-side show the uncertainty band for HVj production.

In figure 25 we display the rapidity distribution of the HW^- system. Since this inclusive quantity is predicted at NLO QCD accuracy by both simulations, we find very good numerical agreement between the two curves at NLO QCD+EW level. The uncertainty band is larger in the HW^-j case. This is due to the fact that for HW^- production there is no renormalization-scale dependence at LO, while in HW^-j such dependence is already present at leading order.

In figures 26 and 27 we compare the transverse momentum of the HW^- pair in two different p_T ranges. Here we observe significant differences due to the fact that this distribution is only computed at leading order in the HW^- simulation, while it is NLO accurate in the HW^-j case. Since we included also EW corrections, in our plots these differences

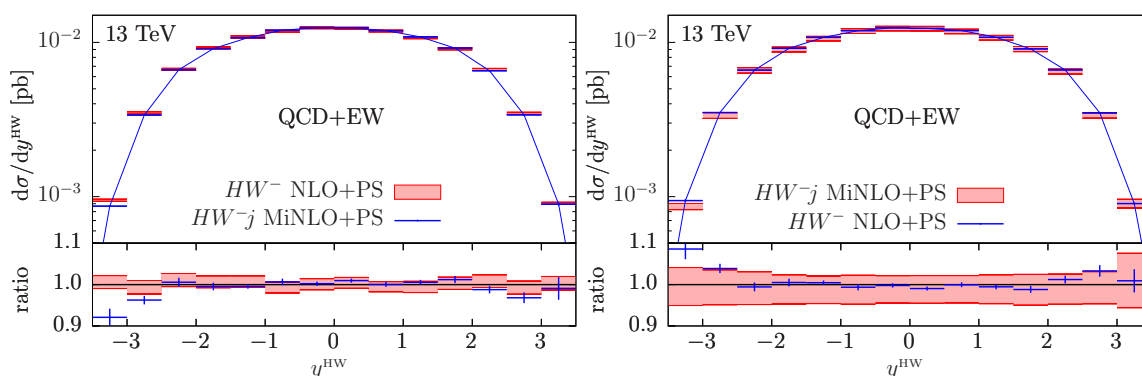


Figure 25. Comparison of NLO+PS and MiNLO+PS predictions for the distributions in the rapidity of the HW^- system in HW^- production. Corrections at NLO QCD+EW are included throughout. The red band is the envelope of the seven-point scale variations for the NLO+PS simulation, in the left panel, and for the MiNLO+PS one, in the right panel. The lower panels show the ratio plot with respect to the central-scale value of the band.

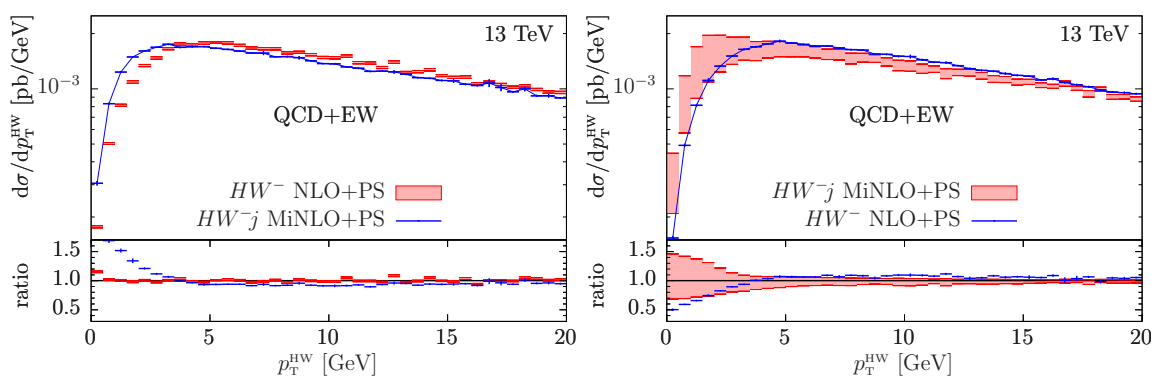


Figure 26. Comparison of NLO+PS and MiNLO+PS predictions for the distributions in transverse momentum of the HW^- system in HW^- production. Same curves and labels as in figure 25.

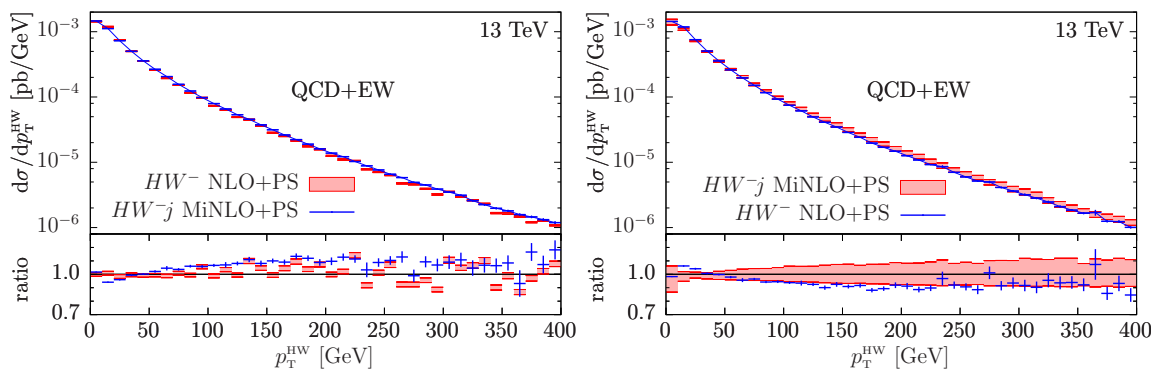


Figure 27. Same comparison as in figure 26, in a wider p_T^{HW} range.

are slightly more pronounced than in the pure QCD implementation of ref. [33]. The fact that such differences emerge in the region below the QCD Sudakov peak (figure 26) is consistent with the observation of enhanced EW effects in that region (figure 22) as discussed in section 6.2. We also note that the uncertainty band for the HW^- generator is smaller than the HW^-j one. This is due again to the fact that, at Born level, HW^- production does not depend upon α_s , while HW^-j production does, and this dependence amplifies the scale-variation band.

8 Summary and conclusions

In this paper we have presented the first NLO QCD+EW calculations for HV and HVj production, with $V = W^\pm, Z$, at the LHC. Specifically, we have considered complete Higgsstrahlung processes corresponding to Higgs boson production in association with off-shell $\ell\nu_\ell$ or $\ell^+\ell^-$ leptonic pairs plus zero or one jet. In addition to fixed-order predictions we have presented realistic simulations obtained by combining NLO QCD+EW calculations with a QCD+QED parton shower. This was achieved by means of the POWHEG BOX RES generator, a recent extension of the POWHEG BOX V2 framework, that allows for consistent NLO+PS simulations in the presence of resonances. In the case of HVj production, using the improved MiNLO approach, we have extended the applicability of NLO QCD+EW predictions to the full phase space, including regions where the hardest jet is unresolved. This is the first application of the MiNLO and POWHEG BOX RES approaches in combination with NLO EW corrections.

We have studied several kinematic distributions for HV and HVj production in proton-proton collisions at 13 TeV, and we have discussed predictions at fixed-order NLO, at the level of POWHEG BOX RES Les Houches events, and at NLO+PS level using Pythia 8.1. Particular care has been taken in combining the QCD+QED shower of Pythia 8.1 with the POWHEG BOX-generated events, since no standard interface is available, at present, to deal with multiple NLO emissions that can arise at production and decay level in resonant processes.

Electroweak corrections typically lower NLO+PS QCD predictions by 5 to 10% at the level of integrated cross sections and in angular distributions. We have observed quantitatively similar and rather constant EW corrections also in the jet- p_T spectrum, as well as in the reconstructed Z -mass and transverse W -mass in the vicinity of the corresponding resonances. In contrast, due to Sudakov logarithms, EW corrections can be much more sizable in the tails of transverse-momentum and invariant-mass distributions. For example, in the Higgs and vector-boson p_T distributions, EW corrections reach up to -25% around 1 TeV. In this respect, the HV and HVj Higgsstrahlung processes behave similarly, i.e. the emission of a jet does not have a sizable impact on EW corrections.

We have studied theoretical uncertainties associated with standard factor-two variations of the renormalization and factorization scales. In the context of the POWHEG formalism, scale variations are performed only at the level of the underlying-Born cross section, while the scale of the strong coupling constant associated with NLO radiation is kept fixed at the corresponding transverse momentum. Thus the resulting scale-variation bands are

typically smaller as compared to the ones obtained in fixed-order NLO calculations. In the total cross sections for HV and HVj production we have found scale uncertainties around 1-2% and 2-4%, respectively, while scale variations in kinematic distributions are typically at the 10% level.

Thanks to the improved `MinLO` prescription, simulations based on NLO QCD+EW matrix elements for HVj production can be applied to inclusive observables and compared against more conventional simulations based on NLO QCD+EW matrix elements for HV production. At NLO QCD, the observed agreement between HV and HVj predictions confirms that, as is well known, the improved `MinLO` approach guarantees NLO QCD accuracy also when the extra jet is integrated out. A similarly good level of agreement was found also at NLO QCD+EW level in a variety of observables. In this regard, based on unitarity and factorization properties of soft and collinear QCD radiation, we have sketched a proof of the fact that the improved `MinLO` approach, applied to QCD jet radiation computed with NLO QCD+EW matrix elements, should provide NLO QCD+EW accuracy in the full phase space.

All relevant matrix elements at NLO EW have been computed using a recent interface of the `POWHEG BOX RES` framework with the `OpenLoops` matrix-element generator. The other QCD amplitudes have been computed in part analytically and in part using the standard interface to `MadGraph4`. We have also presented simple analytic expression that approximate the virtual EW amplitudes in the Sudakov regime at next-to-leading-logarithmic accuracy. This approximation captures the bulk of EW corrections and reproduces exact NLO EW results with reasonable accuracy. Moreover it can be exploited in the combination of the reweighting approach that permits to speed up NLO QCD+EW simulations while providing full NLO EW accuracy in the final results.

The `POWHEG BOX RES` code together with the generators that we have implemented for HV and HVj production can be downloaded following the instructions at the `POWHEG BOX` web page: <http://powhegbox.mib.infn.it>

Acknowledgments

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A Validation of the fixed-order NLO EW corrections in HV production

In this section we compare our fixed-order NLO EW predictions for HW and HZ production with predictions obtained with the Monte Carlo program `HAWK` [45].

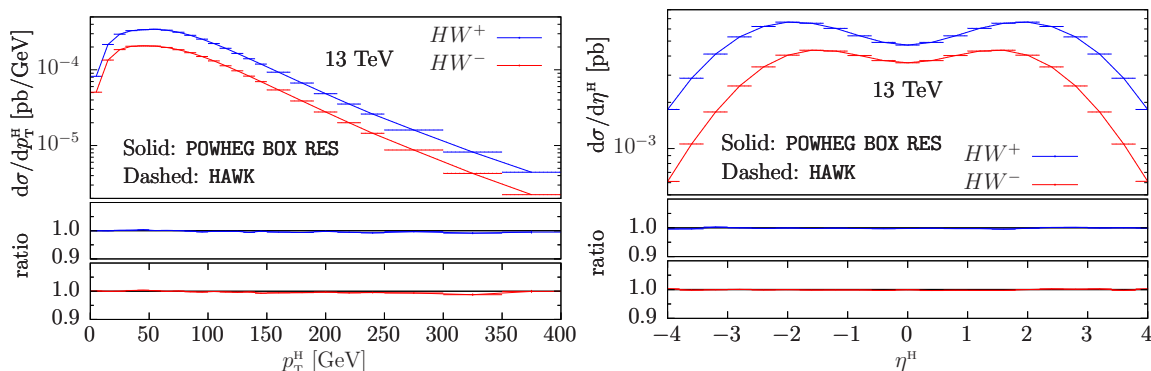


Figure 28. NLO EW predictions for the transverse momentum (left) and the pseudorapidity (right) of the Higgs boson in HW^\pm production. The POWHEG BOX RES and HAWK results are shown with solid and dashed lines respectively. The vertical bars (hardly visible) represent the statistical uncertainties associated to the Monte Carlo integration.

Setup for the comparison. In order to make a comparison between the results generated by HAWK and our results, we switched off photon-initiated contributions in HAWK, since these contributions are currently not included in the POWHEG BOX RES HV generators. Similarly, $b\bar{b}$ -initiated contributions have been discarded in the POWHEG BOX RES, since this sub-process is not included in HAWK. The CKM matrix elements have been set to

$$|V_{ud}^{\text{CKM}}| = |V_{cs}^{\text{CKM}}| = 0.974, \quad |V_{us}^{\text{CKM}}| = |V_{cd}^{\text{CKM}}| = \sqrt{1 - |V_{ud}^{\text{CKM}}|^2}, \quad (\text{A.1})$$

omitting mixing with third-generation quarks. The renormalization and factorization scales are set to the default values used in HAWK, i.e. to the sum of the Higgs and the vector boson masses

$$\mu_R = \mu_F = M_V + M_H, \quad V = W, Z. \quad (\text{A.2})$$

All other input parameters are chosen in accordance with section 3.3.

Photons are recombined with collinear charged leptons if $R_{\gamma\ell} < 0.1$, where $R_{\gamma\ell}$ is the angular separation variable in the (y, ϕ) plane. If more than one charged lepton is present in the final state, the eventual recombination is performed with the lepton having the smallest value of $R_{\gamma\ell}$. After photon recombination, we apply the following cuts on the charged dressed leptons

$$p_T^\ell > 20 \text{ GeV}, \quad |y^\ell| < 2.5, \quad (\text{A.3})$$

while for HW production we also require a missing transverse momentum of

$$\cancel{E}_T > 25 \text{ GeV}. \quad (\text{A.4})$$

Results. In figures 28 and 29 we compare NLO EW predictions obtained with POWHEG BOX RES (solid line) and HAWK (dashed line) for selected observables in HW^+ and HW^- production.

Figure 28 displays the Higgs boson transverse-momentum and pseudorapidity distributions. Within statistical uncertainties the two predictions fully overlap. As a further example, in figure 29 we plot the transverse momentum of the neutrino, i.e. the missing

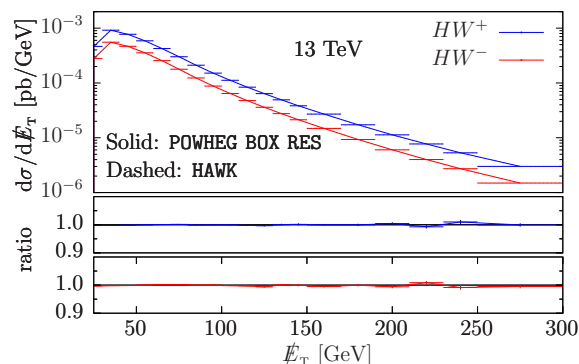


Figure 29. NLO EW predictions for the missing transverse momentum. Predictions and labels as in figure 28.

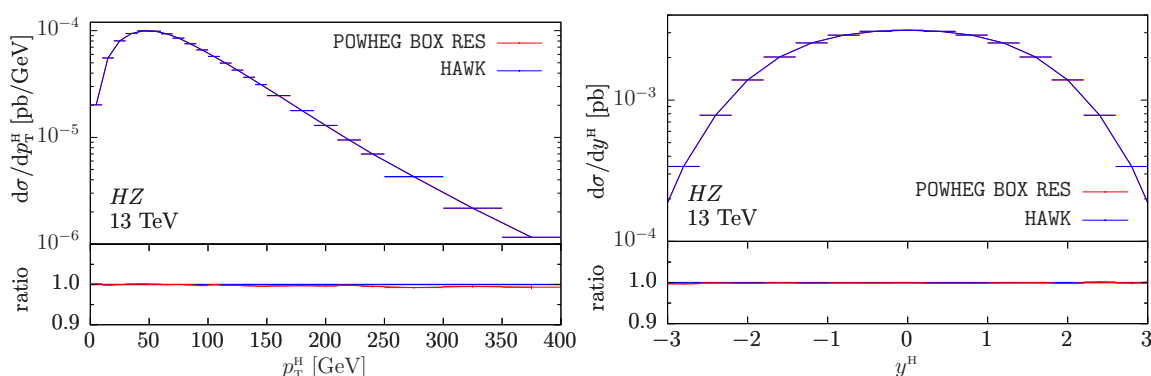


Figure 30. NLO EW predictions for the transverse momentum (left) and the rapidity (right) of the Higgs boson in HZ production. Comparison between the POWHEG BOX RES and the HAWK predictions.

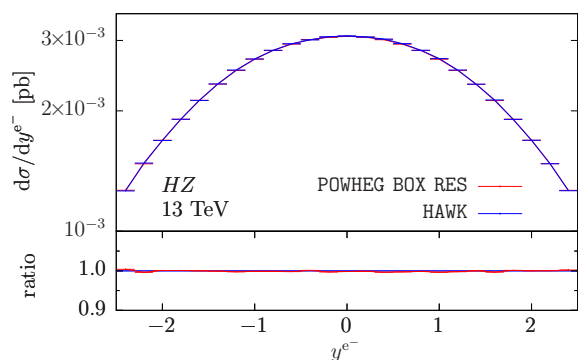


Figure 31. NLO EW predictions for the rapidity of the electron. Same curves and labels as in figure 30.

transverse momentum. Again, we observe perfect agreement between the fixed-order NLO POWHEG BOX RES and HAWK predictions, and a similar level of agreement was found in all considered observables.

As examples for the validation of HZ production, in figure 30 we present the transverse momentum and the rapidity of the Higgs boson, and in figure 31 the rapidity of the

produced electron. Again we find a perfect overlap between the POWHEG BOX RES and HAWK predictions, and the same level of agreement was found for all kinematic distributions that we have examined.

B The virtual EW Sudakov approximation

The calculation of EW virtual corrections is typically more complex than in the case of QCD. This is due to the nontrivial gauge-boson mass spectrum, the presence of Yukawa and scalar interactions, the fact that EW corrections enter also in leptonic vector-boson decays, as well as subtleties related to the treatment of unstable particles. For these reasons, the numerical evaluation of NLO EW virtual corrections can be time consuming. Motivated by this practical issue, in this appendix we present compact analytic formulas that provide a decent approximation of the bulk of NLO EW effects, based on the Sudakov approximation. Besides speeding up the numerics, this approximation provides also valuable insights into the origin of the bulk of the EW corrections.

The largest EW corrections originate in the Sudakov regime, where all kinematic invariants are of the same order and much larger than the electroweak scale. In this high-energy regime, the EW corrections are dominated by Sudakov logarithms [46, 48, 49, 99–101] of the form

$$L(s) = \frac{\alpha_{\text{EM}}}{4\pi} \log^2 \frac{s}{M_V^2}, \quad l(s) = \frac{\alpha_{\text{EM}}}{4\pi} \log \frac{s}{M_V^2}, \quad (\text{B.1})$$

i.e. by leading (LL) and next-to-leading logarithms (NLL) involving the ratio of the partonic center-of-mass energy \sqrt{s} to the electroweak-boson masses, $M_V = M_W, M_Z$. Sudakov EW logarithms originate from virtual gauge bosons that couple to one or two on-shell external particles in the soft and/or collinear limits.

General factorization formulas for LLs and NLLs that apply to any Standard Model process at one loop have been derived in refs. [49, 61, 102]. For a generic n -particle scattering processes with all particles φ_i and momenta p_i incoming¹⁰

$$\varphi_1(p_1) \varphi_2(p_2) \dots \varphi_n(p_n) \rightarrow 0, \quad (\text{B.2})$$

high-energy EW logarithms in one-loop matrix elements assume the general factorized form

$$\begin{aligned} \delta \mathcal{M}^{\varphi_1 \dots \varphi_n}(\{\lambda_i\}, p_1 \dots p_n) &= \sum_{\lambda_i} \delta \lambda_i \frac{\partial \mathcal{M}_0^{\varphi_1 \dots \varphi_n}}{\partial \lambda_i}(\{\lambda_i\}, p_1 \dots p_n) \\ &+ \sum_{\varphi_{1'} \dots \varphi_{n'}} \mathcal{M}_0^{\varphi_{1'} \dots \varphi_{n'}}(\{\lambda_i\}, p_1 \dots p_n) \delta_{\varphi_{1'} \dots \varphi_{n'}}^{\varphi_1 \dots \varphi_n}(\{\lambda_i\}, p_1 \dots p_n). \end{aligned} \quad (\text{B.3})$$

Here the first term is related to the running $\delta \lambda_i$ of the dimensionless coupling parameters in the Born amplitude, while the second term consists of process-independent correction factors $\delta_{\varphi_{1'} \dots \varphi_{n'}}^{\varphi_1 \dots \varphi_n}$ that contain all LL and NLL terms and multiply Born matrix elements for the process at hand. Note that the correction factors are matrices in SU(2) space. In general they act on one or two external particles, requiring the evaluation of SU(2)-transformed matrix elements $\mathcal{M}_0^{\varphi_{1'} \dots \varphi_{n'}}$.

¹⁰In the following we adopt the notation of refs. [49, 61, 102].

The logarithmic EW corrections of eq. (B.3) can be schematically split into five contributions

$$\delta\mathcal{M} = (\delta^{\text{LSC}} + \delta^{\text{SSC,n}} + \delta^{\text{SSC,\pm}} + \delta^{\text{C}} + \delta^{\text{PR}}) \mathcal{M}_0. \quad (\text{B.4})$$

The first three terms are due to double logarithms originating from soft-collinear gauge bosons exchanged between pairs of external legs. This gives rise to angular-independent LLs proportional to $L(s)$ (δ^{LSC}) and subleading angular-dependent logarithms of type $l(s) \log(|r_{kl}|/s)$, with $r_{kl} = (p_k + p_l)^2$. The latter are split into terms associated with neutral ($\delta^{\text{SSC,n}}$) and charged ($\delta^{\text{SSC,\pm}}$) soft-collinear gauge bosons. The remaining terms consist of single-logarithmic contributions from soft/collinear gauge bosons coupling to single external legs (δ^{C}) and from the usual renormalization-group evolution of dimensionless coupling parameters (δ^{PR}).

In the following, the general results of refs. [49, 61, 102] are applied to Higgsstrahlung processes.

B.1 NLL Sudakov approximation for HV and HVj production

In this section we discuss the application of the Sudakov approximation to resonant processes. We focus on vector-boson decays to leptons of the first generation, but our results are applicable also to μ and τ leptons. With u and d we denote generic up- and down-type quarks, with no assumptions on their generation, unless specified.

For the leading-order kinematics of HV and HVj production at particle level we use the notation

$$P_1(p_1) P_2(p_2) \rightarrow H(p_3) V(k) \left(P_6(p_6) \right) \rightarrow H(p_3) \ell_1(p_4) \bar{\ell}_2(p_5) \left(P_6(p_6) \right), \quad (\text{B.5})$$

where $P_i = q, \bar{q}, g$ are generic partons, and $k = p_4 + p_5$ is the off-shell momentum of the decaying vector boson. In HV production, P_6 is not present, and the two incoming partons are always a quark-antiquark pair, while for HVj production an extra gluon can appear, both as an initial-state parton or in the final state.

In order to apply the Sudakov approximation of refs. [49, 61, 102], the Higgsstrahlung processes (2.1) need to be factorized into separate parts associated with the production and decay of the vector boson. This is achieved in a gauge-invariant way by using the leading-pole approximation (LPA) [103, 104], which corresponds to the leading term of a systematic expansion in Γ_V/M_V . At leading order, the LPA for Higgsstrahlung processes reads

$$\mathcal{M}_{0, \text{LPA}}^{P_1 P_2 \rightarrow H \ell_1 \bar{\ell}_2(P_6)} = \frac{1}{k^2 - M_V^2 + i \Gamma_V M_V} \sum_{\lambda=0, \pm 1} \mathcal{M}_0^{P_1 P_2 \rightarrow H V_\lambda(P_6)} \mathcal{M}_0^{V_\lambda \rightarrow \ell_1 \bar{\ell}_2}, \quad (\text{B.6})$$

where factorized matrix elements for vector-boson production and decay on the r.h.s. are summed over the physical polarizations of the vector boson. The propagator in eq. (B.6) depends on the off-shell vector-boson momentum k , while, in the matrix elements on the r.h.s., an on-shell projected momentum k' must be used in order to ensure gauge invariance. This can be achieved with a mapping that, conserving energy and momentum, projects on shell the V and H momenta and rescales accordingly the momenta of the decay products.

In our implementation, we employ such a mapping by keeping fixed the angles formed by the vector boson and by one of the leptons.

In general, in leading-pole approximation, three types of NLO EW corrections need to be considered: factorizable corrections to the production and decay parts, and non-factorizable corrections that connect production and decay. However, the latter are typically quite small [105–108]. Moreover, vector-boson decay do not involve Sudakov EW logarithms. Thus, only the production part receives Sudakov EW corrections, i.e.

$$\delta\mathcal{M}_{\text{LPA}}^{P_1 P_2 \rightarrow H \ell_1 \bar{\ell}_2(P_6)} = \frac{1}{k^2 - M_V^2 + i \Gamma_V M_V} \sum_{\lambda} \delta\mathcal{M}^{P_1 P_2 \rightarrow H V_{\lambda}(P_6)} \mathcal{M}_0^{V_{\lambda} \rightarrow \ell_1 \bar{\ell}_2}, \quad (\text{B.7})$$

and $\delta\mathcal{M}^{P_1 P_2 \rightarrow H V_{\lambda}(P_6)}$ as well as its decay counterpart need to be computed for both transversely- and longitudinally-polarized vector bosons. In the framework of refs. [49, 61, 102], tree amplitudes with longitudinal vector bosons need to be related to corresponding amplitudes with Goldstone bosons using the Goldstone-boson equivalence theorem [109–111]

$$\mathcal{M}_0^{V_L^{a_1} \dots V_L^{a_m} \varphi_1 \dots \varphi_n} = \prod_{k=1}^m i^{(1-Q_{V^{a_k}})} \mathcal{M}_0^{\Phi_{a_1} \dots \Phi_{a_m} \varphi_1 \dots \varphi_n} + \mathcal{O}(ME^{d-1}), \quad (\text{B.8})$$

where $V_L^{a_i}$ are the longitudinal gauge bosons, Φ_{a_i} the corresponding Goldstone bosons, φ_i are the fermions and scalars in the process, M and E are typical scale masses and energies involved in the process, d is the mass dimension of the matrix element and $Q_{V^{a_k}}$ is the electric charge of the vector boson V^{a_k} .

In the following sections we present analytic results for all relevant NLL EW correction factors. These formulae contain group-theoretical quantities such as the electric charge Q of the external particles, their weak isospin T^a , or the electroweak Casimir operator C^{ew} . Their values can be found in appendix B of ref. [102]. For the sine and cosine of the Weinberg angle, we use the shorthand s_w and c_w , respectively.

Large logarithms of the light-fermion masses do not need to be included since we use the G_{μ} scheme, which incorporates the running of the electromagnetic coupling up to the EW scale, and we regularize QED infrared singularities of virtual type at the EW scale by using an effective photon mass $\lambda = M_W$. This approach effectively corresponds, in logarithmic approximation, to the combination of virtual EW corrections with the emission of real photons up to transverse momenta of the order of M_W .

In the framework of the Sudakov NLL approximation, the Sudakov limit is applied only to virtual EW effects, while real QED radiation is treated exactly. More precisely, FKS-subtracted real-emission matrix elements are treated exactly, while only the finite part of the integrated FKS terms, defined via $\overline{\text{MS}}$ subtraction of the IR poles at the scale $\mu = \mu_R$, is included. Concerning IR singularities, this $\overline{\text{MS}}$ subtraction is consistent with the cancellation of virtual QED singularities through the above mentioned $\lambda = M_W$ regularization approach. However, as far as QED logarithms are concerned, we note that we do not apply a fully consistent matching of the (regularized) virtual contributions to real QED radiation. In fact, the former are effectively cut off at the scale M_W , while the latter are subtracted at the scale $\mu = \mu_R$. This implies missing logarithmic terms of order

$\alpha_{\text{EM}} \ln(\mu_R/M_W)$. Nevertheless, as demonstrated by our numerical results, such uncontrolled logarithms do not jeopardize the accuracy of the Sudakov approximation at high energies.

B.2 HW and HWj production

Here we focus on HW^- production and we first consider the partonic process

$$d_L(p_1) \bar{u}_L(p_2) \rightarrow H(p_3) W^-(k) \rightarrow H(p_3) e_L^-(p_4) \bar{\nu}_e(p_5), \quad (\text{B.9})$$

which involves only left-chiral quarks and leptons. Matrix elements for the charge-conjugated process

$$\bar{d}_L(p_1) u_L(p_2) \rightarrow H(p_3) W^+(k) \rightarrow H(p_3) e_L^+(p_4) \nu_e(p_5), \quad (\text{B.10})$$

as well as crossing-related matrix elements corresponding to permutations of the initial-state quarks, can be easily obtained from the ones for the processes (B.9). For the crossed production process

$$d_L(p_1) \bar{u}_L(p_2) H(-p_3) W_\lambda^+(-k) \rightarrow 0, \quad (\text{B.11})$$

the Born amplitudes in the high-energy limit read

$$\mathcal{M}_0^{d_L \bar{u}_L HW_T^+} = \frac{e^2}{\sqrt{2}s_w^2} M_W V_{ud}^{\text{CKM}} \frac{A_{T-}}{q^2}, \quad (\text{B.12})$$

$$\mathcal{M}_0^{d_L \bar{u}_L HW_L^+} = \mathcal{M}_0^{d_L \bar{u}_L H\phi^+} = \frac{e^2}{2\sqrt{2}s_w^2} V_{ud}^{\text{CKM}} \frac{A_{L-}}{q^2}, \quad (\text{B.13})$$

where $q = p_1 + p_2$. Transverse and longitudinal gauge-boson polarization states are denoted as $\lambda = \pm 1 \equiv T$ and $\lambda = 0 \equiv L$, respectively, and

$$A_{T-} = -i \bar{v}_L(p_2) \gamma^\mu u_L(p_1) \epsilon_\mu^T(-k), \quad (\text{B.14})$$

$$A_{L-} = -i \bar{v}_L(p_2) \gamma^\mu u_L(p_1) (-k + p_3)_\mu. \quad (\text{B.15})$$

For the decay of the polarized W^- boson we have

$$\mathcal{M}_0^{W_\lambda^- e^+ \nu_e} = -i \frac{e}{\sqrt{2}s_w} \bar{u}_L(-p_4) \gamma^\mu v_L(-p_5) \epsilon_\mu^\lambda(k). \quad (\text{B.16})$$

For the crossed HW^-j production process

$$d_L(p_1) \bar{u}_L(p_2) H(-p_3) W^+(-k) g(-p_6) \rightarrow 0, \quad (\text{B.17})$$

the polarized Born amplitudes at high energy read

$$\mathcal{M}_0^{d_L \bar{u}_L HW_T^+ g} = \frac{e^2}{\sqrt{2}s_w^2} M_W V_{ud}^{\text{CKM}} g_s t^a \frac{A'_{T-}}{q^2}, \quad (\text{B.18})$$

$$\mathcal{M}_0^{d_L \bar{u}_L HW_L^+ g} = \mathcal{M}_0^{d_L \bar{u}_L H\phi^+} = \frac{e^2}{2\sqrt{2}s_w^2} V_{ud}^{\text{CKM}} g_s t^a \frac{A'_{L-}}{q^2}, \quad (\text{B.19})$$

where g_s is the strong coupling, t^a is the color matrix, and

$$A'_{T-} = -i \bar{v}_L(p_2) \left[\gamma^\mu \frac{\not{p}_1 - \not{p}_6}{(p_1 - p_6)^2} \gamma^\nu + \gamma^\nu \frac{\not{p}_6 - \not{p}_2}{(p_6 - p_2)^2} \gamma^\mu \right] u_L(p_1) \epsilon_\mu^T(-k) \epsilon_\nu(-p_6), \quad (\text{B.20})$$

$$A'_{L-} = -i \bar{v}_L(p_2) \left[\gamma^\mu \frac{\not{p}_1 - \not{p}_6}{(p_1 - p_6)^2} \gamma^\nu + \gamma^\nu \frac{\not{p}_6 - \not{p}_2}{(p_6 - p_2)^2} \gamma^\mu \right] u_L(p_1) (-k + p_3)_\mu \epsilon_\nu(-p_6). \quad (\text{B.21})$$

The gluon-initiated processes can be obtained via appropriate crossing transformations.

Sudakov correction factors for HW production. In the following, we list explicit results for the various corrections factors of eq. (B.4) in the case of the HW production process (B.11), using the Mandelstam invariants $s = (p_1 + p_2)^2$, $t = (p_1 - p_3)^2$ and $u = (p_1 - k)^2$. For transversely polarized W bosons we obtain

$$\begin{aligned}
 \delta^{\text{LSC}} &= -\frac{1}{2}L(s) [2C_q^{\text{ew}} + C_\Phi^{\text{ew}} + C_W^{\text{ew}}] + l(s) \log \frac{M_Z^2}{M_W^2} [(I_{d_L}^Z)^2 + (I_{u_L}^Z)^2 + (I_H^Z)^2 + (I_W^Z)^2] \\
 &\quad + \delta_H^{\text{LSC,h}}, \\
 \delta^{\text{SSC,n}} &= 2l(s) \left(R_{d_L W^+} \log \frac{|u|}{s} - R_{u_L W^+} \log \frac{|t|}{s} \right), \\
 \delta^{\text{SSC,\pm}} &= 2l(s) s_W \left[\log \frac{|t|}{s} \left(\frac{I_{u_L}^Z}{2c_W} - \frac{Q_u}{2s_W} + \frac{I_{d_L}^Z}{s_W^2 c_W} \right) - \log \frac{|u|}{s} \left(\frac{I_{d_L}^Z}{2c_W} - \frac{Q_d}{2s_W} + \frac{I_{u_L}^Z}{s_W^2 c_W} \right) \right], \\
 \delta^{\text{C}} &= l(s) \left[3C_q^{\text{ew}} + 2C_\Phi^{\text{ew}} + \frac{1}{2}b_{WW}^{\text{ew}} - \frac{3}{4s_W^2} \frac{m_t^2}{M_W^2} \right] \\
 &\quad + \frac{\alpha_{\text{EM}}}{4\pi} \left[\left(\frac{3}{4s_W^2} \frac{m_t^2}{M_W^2} + T_{WW} \right) \log \frac{m_t^2}{M_W^2} + \left(\frac{1}{24s_W^2} - 2C_\Phi^{\text{ew}} \right) \log \frac{M_H^2}{M_W^2} \right], \\
 \delta^{\text{PR}} &= \frac{\alpha_{\text{EM}}}{4\pi} \left[\frac{5}{12s_W^2} \log \frac{M_H^2}{M_W^2} - \left(\frac{9 + 6s_W^2 - 32s_W^4}{18s_W^4} + T_{WW} - \frac{3}{4s_W^2} \frac{m_t^2}{M_W^2} \right) \log \frac{m_t^2}{M_W^2} \right] \\
 &\quad + l(s) \left(-\frac{3}{2}b_{WW}^{\text{ew}} + 2C_\Phi^{\text{ew}} - \frac{3}{4s_W^2} \frac{m_t^2}{M_W^2} \right), \tag{B.22}
 \end{aligned}$$

where $L(s)$ and $l(s)$ are defined in eq. (B.1), the factors $\delta_H^{\text{LSC,h}}$ and T_{WW} are defined, respectively, in eqs. (3.26)–(3.29) and (5.36)–(5.37) of ref. [102], and $C_q^{\text{ew}} = C_{d_L}^{\text{ew}} = C_{u_L}^{\text{ew}}$. The coefficient $R_{\phi_1\phi_2}$ is related to the charge and to the weak isospin of the scattering particles via

$$R_{\phi_1\phi_2} = Q_{\phi_1} Q_{\phi_2} + I_{\phi_1}^Z I_{\phi_2}^Z. \tag{B.23}$$

Note that the parameter-renormalization term, δ^{PR} , receives contributions from the renormalization of the (dimensionless) electric charge e and Weinberg angle θ_W , as well as from the renormalization of the W -boson mass in eq. (B.18). This is due to the fact that Born matrix elements for transversely-polarized vector bosons are mass suppressed.¹¹

For longitudinally polarized W bosons, we obtain the Sudakov correction factors

$$\begin{aligned}
 \delta^{\text{LSC}} &= -L(s) (C_q^{\text{ew}} + C_\Phi^{\text{ew}}) + l(s) \log \frac{M_Z^2}{M_W^2} [(I_{d_L}^Z)^2 + (I_{u_L}^Z)^2 + (I_H^Z)^2 + (I_{\phi^+}^Z)^2] \\
 &\quad + \delta_H^{\text{LSC,h}} + \delta_{\phi^\pm}^{\text{LSC,h}}, \\
 \delta^{\text{SSC,n}} = \delta^{\text{SSC,\pm}} &= 2l(s) \left[\log \frac{|t|}{s} (iI_{H^+}^Z I_{d_L}^Z - R_{u_L\phi^+}) - \log \frac{|u|}{s} (iI_{H^+}^Z I_{u_L}^Z - R_{d_L\phi^+}) \right],
 \end{aligned}$$

¹¹Note that certain aspects of the derivation of the general Sudakov EW formulas of refs. [49, 102] are not applicable to mass-suppressed processes. Nevertheless, as one can verify by comparison against the exact EW corrections, such approximate formulae provide a decent approximation. This is also due to the fact that, being mass suppressed, transversely-polarized contributions have a minor impact at high energies.

$$\begin{aligned}
 \delta^{\text{C}} &= l(s) \left[3C_q^{\text{ew}} + 4C_{\Phi}^{\text{ew}} - \frac{3}{2s_w^2} \frac{m_t^2}{M_W^2} \right] \\
 &\quad + \frac{\alpha_{\text{EM}}}{4\pi} \left[\frac{3}{2s_w^2} \frac{m_t^2}{M_W^2} \log \frac{m_t^2}{M_W^2} + \left(\frac{1}{8s_w^2} - 2C_{\Phi}^{\text{ew}} \right) \log \frac{M_H^2}{M_W^2} \right], \\
 \delta^{\text{PR}} &= -b_{WW}^{\text{ew}} l(s) + \frac{\alpha_{\text{EM}}}{4\pi} \left(\frac{5}{6s_w^2} \log \frac{M_H^2}{M_W^2} - \frac{9 + 6s_w^2 - 32s_w^4}{18s_w^4} \log \frac{m_t^2}{M_W^2} \right), \tag{B.24}
 \end{aligned}$$

where the group theoretical quantities involving the charged Goldstone boson ϕ^- arise from the Goldstone-boson equivalence theorem, and the explicit expression for the SU(2) β -function coefficient b_{WW}^{ew} can be found in appendix B of ref. [102].

The correction factors of eqs. (B.22) and (B.24) are equally valid for the HW^- and HW^+ production processes in eqs. (B.9) and (B.10). Corresponding results for processes with the initial-state quarks interchanged are easily obtained by swapping the Mandelstam variables t and u .

Sudakov correction factors for HWj production. The Sudakov correction factors for HWj production are quite similar to the ones for HW production. In fact, the presence of an extra SU(2) \times U(1) singlet gluon has only indirect effects of kinematic type on the Sudakov EW corrections. In particular, the δ^{LSC} , δ^{C} and δ^{PR} factors of eqs. (B.22) and (B.24) are directly applicable to HWj production without any modification. In contrast, the $\delta^{\text{SSC,n}}$ and $\delta^{\text{SSC,\pm}}$ factors need to be generalized by including extra angular-dependent logarithms of type $\log(|r_{12}|/s)$ associated with vector-boson exchange between the initial-state quarks. For $d\bar{u} \rightarrow HW^-$ this kind of logarithms vanishes due to $r_{12} = s$. However, in the case of $d\bar{u} \rightarrow HW^-g$, they need to be taken into account, since they give rise to non-vanishing contributions via crossing transformations of the type $r_{12} \leftrightarrow r_{16}$, which correspond to the case of quark-gluon initial states.

In the transverse case ($\lambda = T$) the SSC correction factors become

$$\begin{aligned}
 \delta^{\text{SSC,n}} &= 2l(s) \left(\log \frac{|r_{1k}|}{s} R_{d_L W^+} - \log \frac{|r_{2k}|}{s} R_{u_L W^+} - \log \frac{|r_{12}|}{s} R_{d_L u_L} \right), \\
 \delta^{\text{SSC,\pm}} &= 2l(s) s_w \left[\log \frac{|r_{23}|}{s} \left(\frac{Q_d}{2s_w} - \frac{I_{d_L}^Z}{2c_w} \right) + \log \frac{|r_{2k}|}{s} \frac{I_{d_L}^Z}{s_w^2 c_w} \right. \\
 &\quad \left. - \log \frac{|r_{13}|}{s} \left(\frac{Q_u}{2s_w} - \frac{I_{u_L}^Z}{2c_w} \right) - \log \frac{|r_{1k}|}{s} \frac{I_{u_L}^Z}{s_w^2 c_w} \right], \tag{B.25}
 \end{aligned}$$

while for longitudinal W^- bosons ($\lambda = L$) they read

$$\begin{aligned}
 \delta^{\text{SSC,n}} &= 2l(s) \left[-\log \frac{|r_{12}|}{s} R_{u_L d_L} + \log \frac{|r_{1k}|}{s} R_{d_L \phi^+} - \log \frac{|r_{2k}|}{s} R_{u_L \phi^+} \right. \\
 &\quad \left. + iI_{H_X}^Z \left(\log \frac{|r_{13}|}{s} I_{d_L}^Z - \log \frac{|r_{23}|}{s} I_{u_L}^Z + \log \frac{|r_{3k}|}{s} I_{\phi^+}^Z \right) \right], \tag{B.26} \\
 \delta^{\text{SSC,\pm}} &= 2l(s) \left[\log \frac{|r_{23}|}{s} R_{d_L \phi^+} - \log \frac{|r_{13}|}{s} R_{u_L \phi^+} - iI_{H_X}^Z \left(\log \frac{|r_{1k}|}{s} I_{u_L}^Z - \log \frac{|r_{2k}|}{s} I_{d_L}^Z \right) \right],
 \end{aligned}$$

where

$$r_{1k} = (p_1 - k)^2, \quad r_{2k} = (p_2 - k)^2, \quad r_{3k} = (p_3 + k)^2. \tag{B.27}$$

The above correction factors are directly applicable to HW^+j production as well, while processes with the initial partons exchanged require the swap $r_{13} \leftrightarrow r_{23}$ and $r_{1k} \leftrightarrow r_{2k}$.

B.3 HZ and HZj production

One of the main differences between the Sudakov EW corrections for HZ and HW production is due to the fact that Z bosons couple to both left- and right-handed currents. As a consequence, for the HZ production process

$$q(p_1) \bar{q}(p_2) \rightarrow H(p_3) Z(k) \rightarrow H(p_3) e^+(p_4) e^-(p_5), \quad (\text{B.28})$$

the squared Born amplitude in LPA reads

$$\left| \mathcal{M}_{0, \text{LPA}}^{q\bar{q} \rightarrow H e^+ e^-} \right|^2 = \frac{1}{(k^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2} \sum_{\kappa, \kappa'} \left| \sum_{\lambda} \mathcal{M}_0^{q\kappa \bar{q}\kappa' \rightarrow H Z \lambda} \mathcal{M}_0^{Z \lambda \rightarrow e_{\kappa'}^+ e_{\kappa}^-} \right|^2, \quad (\text{B.29})$$

where we have explicitly indicated the incoherent sums over the chiralities of external quarks and leptons (κ, κ'), as well as the coherent sum over intermediate vector-boson helicities. We thus need the Born elements for the production and decay of both transverse and longitudinal Z bosons, for different fermion chiralities.

For the crossed process

$$q(p_1) \bar{q}(p_2) H(-p_3) Z_{\lambda}(-k) \rightarrow 0, \quad (\text{B.30})$$

the Born matrix elements in the Sudakov limit read

$$\mathcal{M}_0^{q\kappa \bar{q}\kappa' H Z T} = \frac{e^2 M_Z}{s_w c_w} I_{q\kappa}^Z \frac{A_{TZ}^{\kappa}}{q^2}, \quad (\text{B.31})$$

$$\mathcal{M}_0^{q\kappa \bar{q}\kappa' H Z L} = i \mathcal{M}_0^{q\kappa \bar{q}\kappa' H \chi} = \frac{e^2}{2 s_w c_w} I_{q\kappa}^Z \frac{A_{LZ}^{\kappa}}{q^2}. \quad (\text{B.32})$$

Here and in the following we keep track of the quark chirality $\kappa = R, L$ in the group-theoretical quantities, while $q = p_1 + p_2$, and

$$A_{TZ}^{\kappa} = -i \bar{v}_{\kappa}(p_2) \gamma^{\mu} u_{\kappa}(p_1) \epsilon_{\mu}^T(-k), \quad (\text{B.33})$$

$$A_{LZ}^{\kappa} = -i \bar{v}_{\kappa}(p_2) \gamma^{\mu} u_{\kappa}(p_1) (-k + p_3)_{\mu}. \quad (\text{B.34})$$

The interchange of the initial-state quarks modifies only the spinor part, without changing the structure of the matrix element. The Z -decay matrix element for generic λ reads

$$\mathcal{M}_0^{Z \lambda e_{\kappa'}^- e_{\kappa}^+} = -i e I_{e_{\kappa'}}^Z \bar{u}_{\kappa'}(-p_5) \gamma^{\mu} v_{\kappa'}(-p_4) \epsilon_{\mu}^{\lambda}(k). \quad (\text{B.35})$$

For HZj production the quark-initiated process is given by

$$q(p_1) \bar{q}(p_2) \rightarrow H(p_3) Z(k) g(p_6) \rightarrow H(p_3) e^+(p_4) e^-(p_5) g(p_6). \quad (\text{B.36})$$

The matrix elements for the production of a transverse and a longitudinal Z boson are similar to the ones for HZ production, with the insertion of a gluon

$$\mathcal{M}_0^{q\kappa\bar{q}\kappa HZ_T g} = \frac{e^2 M_Z}{s_w c_w} g_s t^a I_{q\kappa}^Z \frac{A'_{TZ}}{q^2}, \quad (\text{B.37})$$

$$\mathcal{M}_0^{q\kappa\bar{q}\kappa HZ_L g} = i \mathcal{M}_0^{q\kappa\bar{q}\kappa H\chi g} = \frac{e^2}{2s_w c_w} g_s t^a I_{q\kappa}^Z \frac{A'_{LZ}}{q^2}, \quad (\text{B.38})$$

where the spinor parts are given by

$$A'_{TZ}{}^\kappa = -i \bar{v}_\kappa(p_2) \left[\gamma^\mu \frac{\not{p}_1 - \not{p}_6}{(p_1 - p_6)^2} \gamma^\nu + \gamma^\nu \frac{\not{p}_6 - \not{p}_2}{(p_6 - p_2)^2} \gamma^\mu \right] u_\kappa(p_1) \epsilon_\mu^T(-k) \epsilon_\nu(-p_6), \quad (\text{B.39})$$

$$A'_{LZ}{}^\kappa = -i \bar{v}_\kappa(p_2) \left[\gamma^\mu \frac{\not{p}_1 - \not{p}_6}{(p_1 - p_6)^2} \gamma^\nu + \gamma^\nu \frac{\not{p}_6 - \not{p}_2}{(p_6 - p_2)^2} \gamma^\mu \right] u_\kappa(p_1) (-k + p_3)_\mu \epsilon_\nu(-p_6). \quad (\text{B.40})$$

Related amplitudes with an initial-state gluon can be obtained via crossing symmetry.

Sudakov correction factors for HZ production. In the following, we present Sudakov EW correction factors for the HZ production process (B.30) for generic initial-state quark chirality ($\kappa = R, L$) and flavour ($q = u, d$).

For the transverse case they read

$$\begin{aligned} \delta^{\text{LSC}} &= -\frac{1}{2} L(s) [2C_{q\kappa}^{\text{ew}} + C_\Phi^{\text{ew}} + C_{ZZ}^{\text{ew}}] + l(s) \log \frac{M_Z^2}{M_W^2} [2(I_{q\kappa}^Z)^2 + (I_H^Z)^2] + \delta_H^{\text{LSC,h}}, \\ \delta^{\text{SSC,n}} &= 0, \\ \delta_u^{\text{SSC},\pm} &= \delta_{\kappa L} l(s) \frac{F_T^{\text{SSC}}}{I_{u\kappa}^Z}, \\ \delta_d^{\text{SSC},\pm} &= -\delta_{\kappa L} l(s) \sum_{u_i=u,c} |V_{u_i d}^{\text{CKM}}|^2 \frac{F_T^{\text{SSC}}}{I_{d\kappa}^Z}, \\ \delta^{\text{C}} &= \frac{\alpha_{\text{EM}}}{4\pi} \left[\left(\frac{3}{4s_w^2} \frac{m_t^2}{M_W^2} + T_{ZZ} \right) \log \frac{m_t^2}{M_W^2} + \left(\frac{M_Z^2}{24s_w^2 M_W^2} - 2C_\Phi^{\text{ew}} \right) \log \frac{M_H^2}{M_W^2} \right] \\ &\quad + l(s) \left[3C_{q\kappa}^{\text{ew}} + 2C_\Phi^{\text{ew}} + \frac{1}{2} b_{ZZ}^{\text{ew}} - \frac{3}{4s_w^2} \frac{m_t^2}{M_W^2} \right], \\ \delta^{\text{PR}} &= l(s) \left[-b_{WW}^{\text{ew}} + \rho_{q\kappa} \frac{s_w}{c_w} b_{AZ}^{\text{ew}} + 2C_\Phi^{\text{ew}} - \frac{1}{2} b_{ZZ}^{\text{ew}} - \frac{3}{4s_w^2} \frac{m_t^2}{M_W^2} \right] \\ &\quad + \frac{\alpha_{\text{EM}}}{4\pi} \left\{ \left(\frac{5}{6s_w^2} + \frac{5\rho_{q\kappa}}{6c_w^2} - \frac{5M_Z^2}{12s_w^2 M_W^2} \right) \log \frac{M_H^2}{M_W^2} \right. \\ &\quad \left. - \left[\frac{9 + 6s_w^2 - 32s_w^4}{18s_w^2} \left(\frac{1}{s_w^2} + \frac{\rho_{q\kappa}}{c_w^2} \right) + T_{ZZ} - \frac{3}{4s_w^2} \frac{m_t^2}{M_W^2} \right] \log \frac{m_t^2}{M_W^2} \right\}. \quad (\text{B.41}) \end{aligned}$$

Note that the charged-current SSC contributions take a different form for up- and down-type quarks, and

$$F_T^{\text{SSC}} = -\frac{c_w(1+c_w^2)}{2s_w^3} \left[\log \frac{|t|}{s} + \log \frac{|u|}{s} \right]. \quad (\text{B.42})$$

The PR corrections depend on T_{ZZ} , defined in eqs. (5.36)–(5.37) of ref. [102], and

$$\rho_{q\kappa} = \frac{Q_q - T_{q\kappa}^3}{T_{q\kappa}^3 - Q_q s_w^2}. \quad (\text{B.43})$$

Similarly as for the W case, they receive contributions from the renormalization of e , c_w and M_Z . The renormalization of the Weinberg angle also affects the couplings $I_{q\kappa}^Z$ of the quark to the Z boson.

For longitudinally polarized Z bosons we obtain

$$\begin{aligned}
\delta^{\text{LSC}} &= -L(s) [C_{q\kappa}^{\text{ew}} + C_{\Phi}^{\text{ew}}] + 2l(s) \log \frac{M_Z^2}{M_W^2} [(I_{q\kappa}^Z)^2 + (I_H^Z)^2] + \delta_H^{\text{LSC,h}} + \delta_\chi^{\text{LSC,h}}, \\
\delta^{\text{SSC,n}} &= 0, \\
\delta_u^{\text{SSC,\pm}} &= \delta_{\kappa L} l(s) \frac{F_L^{\text{SSC}}}{I_{u\kappa}^Z}, \\
\delta_d^{\text{SSC,\pm}} &= -\delta_{\kappa L} l(s) \sum_{u_i} |V_{u_i d}^{\text{CKM}}|^2 \frac{F_L^{\text{SSC}}}{I_{d\kappa}^Z}, \\
\delta^{\text{C}} &= \frac{\alpha_{\text{EM}}}{4\pi} \left[\frac{3}{2s_w^2} \frac{m_t^2}{M_W^2} \log \frac{m_t^2}{M_W^2} + \left(\frac{M_Z^2}{8s_w^2 M_W^2} - 2C_{\Phi}^{\text{ew}} \right) \log \frac{M_H^2}{M_W^2} \right] \\
&\quad + l(s) \left[3C_{q\kappa}^{\text{ew}} + 4C_{\Phi}^{\text{ew}} - \frac{3}{2s_w^2} \frac{m_t^2}{M_W^2} \right], \\
\delta^{\text{PR}} &= \frac{\alpha_{\text{EM}}}{4\pi} \left(\frac{1}{s_w^2} + \frac{\rho_{q\kappa}}{c_w^2} \right) \left[\frac{5}{6} \log \frac{M_H^2}{M_W^2} - \frac{9 + 6s_w^2 - 32s_w^4}{18s_w^2} \log \frac{m_t^2}{M_W^2} \right] \\
&\quad + l(s) \left[-b_{WW}^{\text{ew}} + \rho_{q\kappa} \frac{s_w}{c_w} b_{AZ}^{\text{ew}} \right], \tag{B.44}
\end{aligned}$$

where

$$F_L^{\text{SSC}} = -\frac{c_w}{s_w^3} \left[\log \frac{|t|}{s} + \log \frac{|u|}{s} \right]. \tag{B.45}$$

The above correction factors (B.41) and (B.44) are $t \leftrightarrow u$ invariant and thus directly applicable also to processes with exchanged initial-state quarks.

Sudakov correction factors for HZj production. Similarly as for HWj production, also in the case of HZj production only the SSC correction factors need to be generalized.

For the transverse case we get

$$\begin{aligned}
\delta^{\text{SSC,n}} &= -2R_{q\kappa q\kappa} l(s) \log \frac{|r_{12}|}{s}, \\
\delta_u^{\text{SSC,\pm}} &= \frac{l(s) \delta_{\kappa L}}{I_{uL}^Z} \left\{ -\frac{1}{s_w^2} \left[\frac{1}{2} s_w c_w \left(\log \frac{|r_{13}|}{s} + \log \frac{|r_{23}|}{s} \right) + I_{dL}^Z \log \frac{|r_{12}|}{s} \right] \right. \\
&\quad \left. + \frac{c_w^3}{s_w} \left(\log \frac{|r_{1k}|}{s} + \log \frac{|r_{2k}|}{s} \right) \right\} + 2c_w^3 \left(\frac{Q_u}{s_w} - \frac{I_{uL}^Z}{c_w} \right) \log \frac{|r_{3k}|}{s}, \\
\delta_d^{\text{SSC,\pm}} &= \frac{l(s) \delta_{\kappa L}}{I_{dL}^Z} \left\{ \frac{1}{s_w^2} \sum_{u_i} |V_{u_i d}^{\text{CKM}}|^2 \left[\frac{1}{2} s_w c_w \left(\log \frac{|r_{13}|}{s} + \log \frac{|r_{23}|}{s} \right) - I_{uL}^Z \log \frac{|r_{12}|}{s} \right] \right. \\
&\quad \left. + \frac{c_w^3}{s_w} \left(\log \frac{|r_{1k}|}{s} + \log \frac{|r_{2k}|}{s} \right) \right\} + 2c_w^3 \left(\frac{Q_d}{s_w} - \frac{I_{dL}^Z}{c_w} \right) \log \frac{|r_{3k}|}{s}, \tag{B.46}
\end{aligned}$$

while for the longitudinal case,

$$\delta^{\text{SSC,n}} = -2l(s) \left[R_{q\kappa q\kappa} \log \frac{|r_{12}|}{s} + (iI_{HX}^Z)^2 \log \frac{|r_{3k}|}{s} \right],$$

$$\begin{aligned}
 \delta_u^{\text{SSC},\pm} &= -\frac{l(s)\delta_{\kappa L}}{s_w^2 I_{u_L}^Z} \left[I_{d_L}^Z \log \frac{|r_{12}|}{s} + \frac{R_{u_L\phi^-}}{iI_{H_X}^Z} \log \frac{|r_{3k}|}{s} \right. \\
 &\quad \left. + \frac{c_w}{2s_w} \left(\log \frac{|r_{13}|}{s} + \log \frac{|r_{1k}|}{s} + \log \frac{|r_{23}|}{s} + \log \frac{|r_{2k}|}{s} \right) \right], \\
 \delta_d^{\text{SSC},\pm} &= \frac{l(s)\delta_{\kappa L}}{s_w^2 I_{d_L}^Z} \left\{ -\log \frac{|r_{3k}|}{s} \frac{R_{d_L\phi^-}}{iI_{H_X}^Z} + \sum_i |V_{u_i d}^{\text{CKM}}|^2 \left[-I_{u_L}^Z \log \frac{|r_{12}|}{s} \right. \right. \\
 &\quad \left. \left. + \frac{c_w}{2s_w} \left(\log \frac{|r_{13}|}{s} + \log \frac{|r_{1k}|}{s} + \log \frac{|r_{23}|}{s} + \log \frac{|r_{2k}|}{s} \right) \right] \right\}, \quad (\text{B.47})
 \end{aligned}$$

where r_{1k} , r_{2k} and r_{3k} are defined in eq. (B.27). Due to the fact that these expressions are symmetric under the $r_{13} \leftrightarrow r_{23}$ and $r_{1k} \leftrightarrow r_{2k}$ permutations, they also hold for processes with exchanged initial quarks.

C Fast evaluation of the virtual electroweak corrections

The evaluation of EW virtual corrections can be relatively time demanding, in particular for HVj production, and the POWHEG BOX RES framework disposes of a few options to speed up this part of the calculation.

Fixed-order NLO results. If one is interested in fixed-order results, the POWHEG BOX RES code can be run twice in the following way:

- In the first high-statistics run, the user sets the flag `select_EW_virt` to 0 in the input file, thus including only the QCD part of the virtual contribution, which are fast to be evaluated. This has the advantage that the bulk of the inclusive cross section is computed with high statistics at a reduced computational cost.
- Then, the code is run again with lower statistics by using the same importance-sampling integration grids generated in the first run, and computing only the missing EW part of the virtual contributions. This is done by setting the flags `virtonly` and `qed_qcd` to 1 in the input file. In this way, if the flag `select_EW_virt` is set to 1, the complete virtual corrections are included. If instead the user is interested in obtaining the Sudakov approximated results, the code can be run by setting `select_EW_virt` to 2.

Finally, the kinematic distributions obtained in the two previous steps should be combined, by summing them together.

Les-Houches-level Monte Carlo events. If one is interested in generating Monte Carlo events at the Les Houches level, then the code can be run with some approximation of the virtual contributions (or even with the virtual corrections set to zero). At the end, the generated events need only to be reweighted, using the POWHEG BOX RES reweighting feature, with the full virtual corrections activated. In this way, when the event is generated, the use of an approximated virtual contribution (whose evaluation could be requested several times

per event) considerably speeds up the code. Once the event is generated, the reweighting procedure calls the full virtual contribution only once per event. This reduces the running time in a drastic way.

The different options available are the following:

- The user can generate the events omitting virtual contributions of QCD and EW kind. This is obtained by setting the POWHEG BOX RES flag `novirtual` to 1 in the input file. Since the inclusive cross section used to generate the weight associated to the event is computed without the finite part of the virtual corrections, the weight associated with a single event can be very different with respect to the weight obtained after the reweighting procedure applied on that event. This can give rise to statistical fluctuations in the kinematic distributions, that would need a higher number of events to be smoothed.
- The user can generate the events including only the QCD part of the virtual contributions, that are quite fast to be evaluated. This is done by setting the POWHEG BOX RES flag `select_EW_virt` to 0 in the input file. The difference with the previous case is that an important part of the virtual corrections is included in the calculation of the inclusive cross section, and the results after reweighting tend to be smoother.
- The best option is to include the QCD part of the virtual corrections together with their EW NLL approximation. This is achieved by setting `select_EW_virt` to 2. This option is the one we have used to generate the events analyzed in section 5. Since the EW NLL approximation of the virtual corrections captures most of the dominant Sudakov logarithms, running the code with this setting generates events whose weight is very similar to the final weight associated to each event after reweighting.

The default value for the `select_EW_virt` flag is 1, which corresponds to the inclusion of exact virtual EW contributions at all stages.

D Interface to Pythia 8.1 and the veto procedure

In order to generate realistic event samples at NLO+PS accuracy, including both QCD and QED corrections, the radiation of QCD partons and photons generated at the LHE level by the POWHEG BOX RES framework has to be completed by a Monte Carlo showering program. This is achieved through a dedicated interface that feeds the LH events to `Pythia 8.1`. The initialization requires the following instructions:

```
pythia.readString("SpaceShower:pTmaxMatch = 1");
pythia.readString("TimeShower:pTmaxMatch = 1");
```

Photon radiation off quarks and leptons is activated with:

```
pythia.readString("TimeShower:QEDshowerByL = on");
pythia.readString("TimeShower:QEDshowerByQ = on");
pythia.readString("SpaceShower:QEDshowerByQ = on");
pythia.readString("TimeShower:QEDshowerByGamma = off");
```

The last instruction prevents photons from further splitting into fermion-antifermion pairs.

In our analysis we do not include hadronization or underlying-event effects, and we consider the Higgs boson as stable.

The veto procedure. In the following we discuss the veto procedure that is applied in order to guarantee a consistent combination of QCD and QED radiation generated at LHE level with subsequent parton-shower emission. Since we apply the multi-radiation mode described in section 3.1, each LH event generated by the POWHEG BOX RES framework can be accompanied by both QCD and QED radiation. Radiation of QCD type arises only at the “production” level, while photon radiation can come both from “production” and from the charged leptons that arise from the decays of Z and W resonances.

For QCD radiation, the standard veto shower implemented in any Monte Carlo program is used. In practice, the highest transverse momentum of the radiation (of QCD or QED type) generated at the “production” level by the POWHEG BOX RES framework is passed to the shower Monte Carlo program through the variable `scalup`, in the Les Houches interface.

For what concerns QED radiation, since the Les Houches interface does not provide a standard mechanism to veto radiation from resonance decays, we have implemented a dedicated veto procedure. The POWHEG BOX RES events can have up to two photons at LHE level, one associated with the production part, and one with the decay part of the process, and the shower Monte Carlo has to be instructed to veto, separately at the level of production and decay, any photon with transverse momentum higher than the hardness of the emissions produced by the POWHEG BOX RES framework.

To this end, we first scan the Les Houches event to identify the photons that have been generated by the POWHEG BOX RES framework, determining if their mother belongs to the production or to the decay products.¹² We then shower the event with `Pythia 8.1`, restricting QCD radiation by means of the `scalup` variable as discussed above, and identifying the extra photons produced by the shower algorithm.

For photons that are generated by the shower at the production level we apply a similar veto procedure as for QCD radiation:

1. we compute the transverse momentum of each photon produced by the shower, at production level, and store its maximum value p_T^{\max} for the event at hand;
2. if p_T^{\max} is greater than `scalup` the event is vetoed. This procedure effectively amounts to requiring that, at production level, the shower does not generate any QED radiation with transverse momentum greater than the radiation of QCD or QED type produced by the POWHEG BOX RES.
3. Since, in order to ensure momentum conservation, the Monte Carlo reshuffling procedure during shower generation slightly modifies the momenta of the particles, we also check that p_T^{\max} does not exceed the hardness of the LH photon after reshuffling. If this happens, the event is vetoed.

¹²In HVj production, photons radiated by the hardest final-state quark, already present at the LH event level, are considered as coming from the production stage.

For photons associated to the resonance decay, we proceed as follows:

1. if no photon is present at the LHE level, this means that the POWHEG BOX RES has not been able to generate radiation harder than the minimum value of 10^{-3} GeV, set as a minimum for the transverse-momentum of photon radiation.¹³ In this case, any shower QED radiation harder than 1 MeV is vetoed in the decay.
2. If instead a photon is already present, we compute its transverse momentum with respect to the lepton emitter in the center-of-mass frame of the mother resonance, and store this value in $p_{T,\text{rel}}^{\text{max}}$. In HZ and HZj production, at Les Houches level, it is not possible to know if the photon has been emitted by the lepton or by the antilepton, and $p_{T,\text{rel}}^{\text{max}}$ is set to the minimum value between the two relative transverse momenta. We then veto the event if, among the photons produced at decay level, the maximum relative transverse momentum is greater than $p_{T,\text{rel}}^{\text{max}}$.

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¹³The parameter that fixes the square of this value is `rad_ptsqmin_em = 1e-6`.

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