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Logarithmic electroweak corrections to hadronic $Z + 1$ jet production at large transverse momentum

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Abstract:
We consider hadronic production of a Z boson in association with a jet and study one- and two-loop electroweak logarithmic corrections in the region of high Z-boson transverse momentum, $p_{\mathrm{T}} \gg M_Z$, including leading and next-to-leading logarithms. Numerical results for the LHC and Tevatron colliders are presented. At the LHC these corrections amount to tens of per cent and will be important for interpretation of the measurements.

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1 Introduction

Hadronic weak-boson production in association with jets forms one of the most important classes of Standard Model processes. Due to its relatively high cross section and clean experimental signature the $Z + 1$ jet production process can be used for a precise determination of the parton distribution functions [1]. In particular, the precise measurement of the Z-boson transverse momentum (p_{T}) distribution combined with theoretical predictions of correspondingly high accuracy will permit to extract the gluon distribution function with unprecendented accuracy at the Large Hadron Collider (LHC). In order to achieve desired precision (at the per cent level) radiative corrections should be included in the theoretical predictions.

The next-to-leading order QCD corrections to the total cross section and p_{T} distribution for hadronic $Z + 1$ jet production have been known for a relatively long time by now [2, 3]. Recent predictions for the Tevatron and the LHC can be found in Refs. [4] and [5], respectively. The size of these corrections can amount to several tens of per cent depending on the observable under consideration, including jet definition, as well as the renormalisation and factorisation scales.

On the basis of simple arguments regarding the strength of coupling constants one might expect electroweak (EW) corrections to be in general much smaller than QCD corrections. However, in the region $\sqrt{\hat{s}} \gg M_W \simeq M_Z$, where the center-of-mass energy $\sqrt{\hat{s}}$ of partonic scattering becomes much larger than the weak-boson mass scale, the EW corrections are strongly enhanced by logarithmic mass singularities. At $\mathcal{O}(\alpha^L)$ the leading logarithms (LLs) are of the form $\alpha^L \log^{2L}(\hat{s}/M_W^2)$, the next-to-leading logarithms (NLLs) are of the form $\alpha^L \log^{2L-1}(\hat{s}/M_W^2)$, etc. These EW logarithms originate from soft/collinear emission of virtual EW gauge bosons (γ, Z, W) off initial- or final-state particles. Owing to the finite weak-boson masses, the real emission of a soft/collinear Z or W boson can be observed as a separate process and hence does not need to be included in the definition of physical observables. Thus, in contrast to mass singularities in massless gauge theories such as QED or QCD, the EW mass singularities of virtual origin are not necessarily compensated by corresponding mass singularities from real weak-boson radiation.

Typically, at $\sqrt{\hat{s}} \simeq 1\,\mathrm{TeV}$ the EW logarithms are estimated to yield one-loop corrections of tens of per cent and two-loop corrections of a few per cent. At one loop it has been proven that the EW LLs and NLLs are universal and results applicable to arbitrary processes are available [6, 7]. For γW and ZW production at the LHC it was shown that these EW logarithmic corrections can become of the order of the QCD corrections [8]. Furthermore, in Ref. [9] the EW one-loop corrections to the γZ production process at the LHC were found numerically sizeable, with the logarithmic part providing a major contribution to the corrections. The resummation of the EW logarithms was discussed first in Ref. [10] with emphasis on four-fermion processes (for an early discussion see also Ref. [11]). General resummation prescriptions for the EW LLs and NLLs were derived in Refs. [12, 13] and Refs. [13, 14], respectively. For massless four-fermion processes the EW logarithms were resummed up to the NNLLs [15], and it was found that, at the TeV scale, the leading and subleading logarithms have similar size and alternating signs, which gives rise to large cancellations. These resummation prescriptions were confirmed by various two-loop calculations at the LL level [16, 17, 18] and NLL level [19, 20]. Only few two-loop results based on a diagrammatic calculation exist beyond the NLL level. In particular, the corrections to a massless fermionic form factor within an abelian massive gauge theory [21, 22] were evaluated up to N^{4}LL approximation and agreement with the earlier NNLL results was observed.

In this letter we investigate the $Z+1$ jet production at the LHC and the Tevatron. In particular, we focus on the high-p_{T} region, $p_{\mathrm{T}} \gg M_Z$, and discuss the impact of the one- and two-loop EW logarithms on the total cross section and the p_{T} distribution for this process. We also compare the one-loop logarithmic corrections with the results presented in Ref. [23], where the one-loop weak corrections were evaluated numerically.

2 Analytic results

The hadronic production of a Z boson with finite transverse momentum p_{T} in association with a jet receives contributions from the partonic subprocesses

$$q\bar{q} \to Zg, \qquad qg \to Zq, \qquad \bar{q}g \to Z\bar{q}. \tag{1}$$

The p_{T} distribution for the hadronic process $h_1 h_2 \to Z + j$ reads

$$\begin{aligned}\frac{\mathrm{d}\sigma^{h_1 h_2}}{\mathrm{d}p_{\mathrm{T}}} = \sum_{q=u,d,c,s,b} \int_0^1 \mathrm{d}x_1 \int_0^1 \mathrm{d}x_2\, \theta(x_1 x_2 - \hat{\tau}_{\min}) \Biggl\{ \Bigl[f_{h_1,q}(x_1) f_{h_2,\bar{q}}(x_2) + (1 \leftrightarrow 2) \Bigr] \frac{\mathrm{d}\hat{\sigma}^{q\bar{q}}}{\mathrm{d}p_{\mathrm{T}}} \\ + \Bigl[\Bigl(f_{h_1,q}(x_1) f_{h_2,g}(x_2) + f_{h_1,\bar{q}}(x_1) f_{h_2,g}(x_2) \Bigr) + (1 \leftrightarrow 2) \Bigr] \frac{\mathrm{d}\hat{\sigma}^{qg}}{\mathrm{d}p_{\mathrm{T}}} \Biggr\},\end{aligned} \tag{2}$$

where $f_{h,i}(x)$ are the parton distribution functions, $\hat{\tau}_{\min} = (p_{\mathrm{T}} + m_{\mathrm{T}})^2/s$ with $\sqrt{s}$ denoting the collider energy and $m_{\mathrm{T}} = \sqrt{p_{\mathrm{T}}^2 + M_Z^2}$. The p_{T} distribution for the unpolarized partonic subprocess $ij \to Zk$ reads

$$\frac{\mathrm{d}\hat{\sigma}^{ij}}{\mathrm{d}p_{\mathrm{T}}} = \frac{p_{\mathrm{T}}}{8\pi N_{ij} \hat{s} |\hat{t} - \hat{u}|} \left[\overline{\sum_{\mathrm{pol}}} |\mathcal{M}^{ij}|^2 + (\hat{t} \leftrightarrow \hat{u}) \right], \tag{3}$$

where $\hat{s} = (p_i + p_j)^2$, $\hat{t} = (p_i - p_Z)^2$, and $\hat{u} = (p_j - p_Z)^2$ are the Mandelstam invariants. The momenta $p_{i,j,k}$ of the partons are assumed to be massless, whereas $p_Z^2 = M_Z^2$ for the Z-boson momentum. In terms of x_1, x_2, p_{T}, and the collider energy $\sqrt{s}$, we have

$$\hat{s} = x_1 x_2 s, \qquad \hat{t} = \frac{M_Z^2 - \hat{s}}{2}(1 - \cos\theta), \qquad \hat{u} = \frac{M_Z^2 - \hat{s}}{2}(1 + \cos\theta), \tag{4}$$

where $\cos\theta = \sqrt{1 - 4p_{\mathrm{T}}^2 \hat{s}/(\hat{s} - M_Z^2)^2}$ corresponds to the cosine of the angle between the momenta p_i and p_Z in the partonic center-of-mass frame. The factor $1/N_{ij}$ in (3), with $N_{q\bar{q}} = N_{\mathrm{c}}^2$, $N_{qg} = N_{\mathrm{c}}(N_{\mathrm{c}}^2 - 1)$, and $N_{\mathrm{c}} = 3$, accounts for the initial-state colour average.

The unpolarised squared amplitude for the $q\bar{q} \to Zg$ process reads

$$\overline{\sum_{\mathrm{pol}}} |\mathcal{M}^{q\bar{q}}|^2 = 8\pi^2 \alpha \alpha_{\mathrm{S}} (N_{\mathrm{c}}^2 - 1) \frac{\hat{t}^2 + \hat{u}^2 + 2M_Z^2 \hat{s}}{\hat{t}\hat{u}} \left[A^{(0)} + \left(\frac{\alpha}{2\pi}\right) A^{(1)} + \left(\frac{\alpha}{2\pi}\right)^2 A^{(2)} \right], \tag{5}$$

where α and α_{S} are the electromagnetic and strong coupling constants, respectively. The amplitude for the qg-initiated process is easily obtained from (5) using crossing symmetry,

$$\overline{\sum_{\mathrm{pol}}}|\mathcal{M}^{qg}|^2 = - \overline{\sum_{\mathrm{pol}}}|\mathcal{M}^{q\bar{q}}|^2 \Bigg|_{\hat{s}\leftrightarrow\hat{u}} . \tag{6}$$

The tree-level contribution $A^{(0)}$ to (5) reads

$$A^{(0)} = \sum_{\lambda=\mathrm{L,R}} \left(I^Z_{q_\lambda}\right)^2 \qquad \text{with} \qquad I^Z_{q_\lambda} = \frac{c_{\mathrm{W}}}{s_{\mathrm{W}}} T^3_{q_\lambda} - \frac{s_{\mathrm{W}}}{c_{\mathrm{W}}}\frac{Y_{q_\lambda}}{2}, \tag{7}$$

where $T^3_{q_\lambda}$ and Y_{q_λ} are the weak isospin and hypercharge for left- ($\lambda = \mathrm{L}$) and right-handed ($\lambda = \mathrm{R}$) quarks, and we use the shorthands $c_{\mathrm{W}} = \cos\theta_{\mathrm{W}}$ and $s_{\mathrm{W}} = \sin\theta_{\mathrm{W}}$ for the weak mixing angle θ_{W}.

The virtual electroweak corrections were evaluated in the high energy region where all invariants are much larger than the weak-boson mass scale,

$$|\hat{r}| \gg M_W^2 \sim M_Z^2 \qquad \text{for} \qquad \hat{r} = \hat{s}, \hat{t}, \hat{u}. \tag{8}$$

In this region, using the results of Refs. [6, 19, 14], we have computed the one- and two-loop electroweak logarithms of the type $\mathrm{L}^k_{\hat{r}} = \log^k(|\hat{r}|/M_W^2)$ to next-to-leading logarithmic accuracy. At L-loop level this approximation includes leading logarithms (LLs) and next-to-leading logarithms (NLLs), i.e. contributions of order $\mathrm{L}^{2L}_{\hat{s}}$ and $\mathrm{L}^{2L-1}_{\hat{s}}$, respectively. Contributions of order $\mathrm{L}^{2L-2}_{\hat{s}}$ are systematically omitted. In particular, for the angular-dependent logarithms, i.e. logarithms of the type $\mathrm{L}^{2L}_{\hat{r}}$ and $\mathrm{L}^{2L-1}_{\hat{r}}$ involving invariants $\hat{r} = \hat{t}, \hat{u}$, we use the approximation

$$\mathrm{L}^{2L}_{\hat{r}} = \mathrm{L}^{2L}_{\hat{s}} + 2L\log\left(\frac{|\hat{r}|}{\hat{s}}\right)\mathrm{L}^{2L-1}_{\hat{s}} + \mathcal{O}(\mathrm{L}^{2L-2}_{\hat{s}}), \qquad \mathrm{L}^{2L-1}_{\hat{r}} = \mathrm{L}^{2L-1}_{\hat{s}} + \mathcal{O}(\mathrm{L}^{2L-2}_{\hat{s}}). \tag{9}$$

In order to separate the logarithmic electroweak corrections into finite and infrared-divergent parts in an elegant way[1], it is useful to introduce a fictitious photon mass $\lambda = M_W$. In the present paper we restrict ourselves to the finite part of the corrections which is defined as the complete electroweak correction with $\lambda = M_W$. The remaining part, which is not considered here, depends on the gap between the (vanishing) photon mass and the weak-boson mass scale and is infrared divergent. This infrared-divergent part can be factorized and resummed separately [12] and needs to be combined with real photon radiation. Its size depends on the choice of cuts on the photon energies and angles.

[1]Using the infrared-evolution-equation (IREE) [12], the electroweak logarithmic corrections can be factorized in a SU(2)×U(1)-symmetric and finite part times a $\mathrm{U}(1)_{\mathrm{em}}$-symmetric and infrared-divergent part. The SU(2)×U(1)-symmetric part corresponds to the complete electroweak corrections with infrared divergences regularized by a fictitious photon mass $\lambda = M_W$. The $\mathrm{U}(1)_{\mathrm{em}}$-symmetric part contains the infrared divergences resulting from the gap between the (vanishing) photon mass and the weak-boson mass scale. Within the framework of the IREE these two parts factorize and exponentiate separately. At the LL level, the correctness of this IREE-approach has been completely confirmed by process-independent two-loop electroweak calculations based on different gauge-fixings [18, 19]. Recent two-loop calculations [19, 20, 22] have shown that the IREE-approach is applicable also beyond LL level.

The small contributions resulting from logarithms of the Z-W mass ratio are neglected. Mass-suppressed corrections of order $M_W^2/|\hat{r}|$ are not considered and, since the contributions from longitudinally polarised Z bosons are mass-suppressed in the high energy limit, the logarithmic corrections have been computed assuming that the external Z boson is transversely polarised. At one-loop level we obtain

$$A^{(1)} = -\sum_{\lambda=\mathrm{L,R}} I^Z_{q_\lambda}\left[I^Z_{q_\lambda} C^{\mathrm{ew}}_{q_\lambda}\left(\mathrm{L}^2_{\hat{s}} - 3\mathrm{L}_{\hat{s}}\right) + \frac{c_\mathrm{W}}{s^3_\mathrm{W}} T^3_{q_\lambda}\left(\mathrm{L}^2_{\hat{t}} + \mathrm{L}^2_{\hat{u}} - \mathrm{L}^2_{\hat{s}}\right)\right], \tag{10}$$

where $C^{\mathrm{ew}}_{q_\lambda} = Y^2_{q_\lambda}/(4c^2_\mathrm{W}) + C_{q_\lambda}/s^2_\mathrm{W}$ are the eigenvalues of the electroweak Casimir operator for quarks, with $C_{q_\mathrm{L}} = 3/4$ and $C_{q_\mathrm{R}} = 0$. The above result takes into account LLs and NLLs of soft/collinear origin as well as NLLs resulting from the running of the coupling constants.[2] The latter NLLs are, however, cancelled by the collinear logarithms that are associated with the final-state Z boson [6] and thus do not appear in the result (10). In the one-loop approximation and for the process under consideration the electroweak corrections can be split into electromagnetic and weak parts in a gauge-invariant way. The weak corrections (which result from virtual Z- and W-boson exchange) are easily obtained from (10) through the substitution $C^{\mathrm{ew}}_{q_\lambda} \to C^{\mathrm{ew}}_{q_\lambda} - Q^2_q$, where Q_q denotes the electromagnetic charge of the quark q.

At two-loop level we obtain

$$\begin{aligned} A^{(2)} = \sum_{\lambda=\mathrm{L,R}} \Bigg\{ & \frac{1}{2}\left(I^Z_{q_\lambda} C^{\mathrm{ew}}_{q_\lambda} + \frac{c_\mathrm{W}}{s^3_\mathrm{W}} T^3_{q_\lambda}\right)\left[I^Z_{q_\lambda} C^{\mathrm{ew}}_{q_\lambda}\left(\mathrm{L}^4_{\hat{s}} - 6\mathrm{L}^3_{\hat{s}}\right)\right. \\ & \left. + \frac{c_\mathrm{W}}{s^3_\mathrm{W}} T^3_{q_\lambda}\left(\mathrm{L}^4_{\hat{t}} + \mathrm{L}^4_{\hat{u}} - \mathrm{L}^4_{\hat{s}}\right)\right] - \frac{T^3_{q_\lambda} Y_{q_\lambda}}{8 s^4_\mathrm{W}}\left(\mathrm{L}^4_{\hat{t}} + \mathrm{L}^4_{\hat{u}} - \mathrm{L}^4_{\hat{s}}\right) \\ & + \frac{1}{6} I^Z_{q_\lambda}\left[I^Z_{q_\lambda}\left(\frac{b_1}{c^2_\mathrm{W}}\left(\frac{Y_{q_\lambda}}{2}\right)^2 + \frac{b_2}{s^2_\mathrm{W}} C_{q_\lambda}\right) + \frac{c_\mathrm{W}}{s^3_\mathrm{W}} T^3_{q_\lambda} b_2\right]\mathrm{L}^3_{\hat{s}}\Bigg\}, \end{aligned} \tag{11}$$

where $b_1 = -41/(6c^2_\mathrm{W})$ and $b_2 = 19/(6s^2_\mathrm{W})$ are the one-loop β-function coefficients associated with the U(1) and SU(2) couplings, respectively. We note that at the two-loop level the purely weak corrections cannot be isolated from the complete electroweak corrections in a gauge-invariant way. The LLs as well as the angular-dependent subset of the NLLs in (11), i.e. all contributions of the form $\mathrm{L}^4_{\hat{r}}$ with $\hat{r} = \hat{s}, \hat{t}, \hat{u}$, have been derived from Ref. [19]. There, by means of a diagrammatic two-loop calculation in the spontaneously broken electroweak theory, it was shown that such two-loop terms result from the exponentiation of the corresponding one-loop corrections. The additional NLLs of the form $\mathrm{L}^3_{\hat{s}}$ in (11) have been obtained via a fixed order expansion of the general (process independent) resummed expression proposed in Ref. [14]. This NLL resummation [14] relies on the assumption that effects from spontaneous breaking of the SU(2)×U(1) symmetry can be neglected in the high energy limit.

[2]We note that in addition to the logarithmic corrections in (10) the partonic subprocesses involving b quarks receive also NLL corrections proportional to M_t^2/M_W^2, which originate from the Yukawa interaction in the heavy-quark doublet [6]. At partonic level these corrections can be relatively large. However, their contribution to the hadronic cross section is almost negligible owing to the smallness of the b-quark density of protons and anti-protons. We have checked that the correction resulting from these Yukawa terms does not exceed one per mille at the LHC.

3 Numerical results and discussion

In this section we discuss numerical effects of the one- and two-loop logarithmic EW corrections on the $Z+1$ jet cross sections at the LHC and the Tevatron. Results presented here are obtained for the following values of the electroweak parameters: $\alpha = \alpha(M_Z^2) = 1/127.9$, $s_{\mathrm{W}}^2 = 0.231$, $M_Z = 91.19\,\mathrm{GeV}$ and $M_W = M_Z\sqrt{1-s_{\mathrm{W}}^2}$. When discussing the (relative) statistical error, defined as $\Delta\sigma_{\mathrm{stat}}/\sigma = 1/\sqrt{N}$ with $N = \mathcal{L}\times\mathrm{BR}(Z \to l, \nu_l)\times\sigma_{LO}$, we include all leptonic decays of Z with the branching ratio $\mathrm{BR}(Z \to l, \nu_l) = 30.6\%$ and assume total integrated luminosity of $\mathcal{L} = 300\,\mathrm{fb}^{-1}$ and $\mathcal{L} = 11\,\mathrm{fb}^{-1}$ for the LHC [24] and the Tevatron [25], respectively.

To calculate the hadronic cross sections we use the leading order parton distribution functions MRST2001LO [26] at the factorisation scale $\mu_{\mathrm{F}}^2 = p_{\mathrm{T}}^2$. Similarly, the strong coupling constant α_{S} is evaluated (for five active flavours) at the scale $\mu^2 = p_{\mathrm{T}}^2$ and, consistently with the parton distribution analysis, we use the one-loop expression with $\alpha_{\mathrm{S}}(M_Z^2) = 0.13$ [26].

No rapidity cuts are imposed for the purpose of this study, as their effect on the p_{T} distributions does not appear to be relevant, both at the LHC and the Tevatron. We checked that at the LHC a rapidity cut $|y_Z| < 4$ decreases the p_{T} distribution at $p_{\mathrm{T}} = 100\,\mathrm{GeV}$ only by less than 0.5% and for $p_{\mathrm{T}} \gg 100\,\mathrm{GeV}$ the effect of the rapidity cut becomes negligible.

The lowest order (LO) transverse momentum distribution of the outgoing Z boson at the LHC is shown in Fig. 1a. We also display seperately the lowest order contribution from the longitudinally polarised Z boson. Its smallness justifies to treat the radiative corrections to the unpolarised process as if the Z boson would be only transversally polarised. The next-to-leading order (NLO) and next-to-next-to-leading order (NNLO) results, i.e. one- and two-loop corrected distributions, are also presented.

The size of the corrections relative to the leading order p_{T} distribution is shown in Fig. 1b. These relative corrections increase with p_{T}. For high p_{T} they yield a negative contribution of tens of per cent to the p_{T} distribution in the one-loop case and a positive contribution of several per cent in the two-loop case. More precisely, for $1\,\mathrm{TeV} \le p_{\mathrm{T}} \le 2\,\mathrm{TeV}$, the logarithmic one-loop electroweak correction decreases the cross section by 24% to 39%, and the two-loop correction moves the result up by 3% to 8%. The biggest (positive) contribution to the two-loop result is given by the LL terms, with NLL terms accounting for about 50% of the size of the LL contribution at high p_{T} but with an opposite sign. This observation demonstrates that the control of NLL effects in the two-loop approximation is required to predict the cross section at a level better than 10% for high p_{T}. We note that for the one-loop result the EM contribution grows with p_{T} but even at the highest p_{T} considered it does not exceed 3-4% of the total electroweak one-loop logarithmic correction and thus is negligible.

The size of the one- and two-loop electroweak logarithmic corrections to the p_{T} distribution clearly makes them relevant for future measurements at the LHC. The significance of the corrections is also illustrated in Fig. 2 where we plot the relative one- and two-loop corrections to the total cross section obtained by integration over p_{T} with $p_{\mathrm{T}} > p_{\mathrm{T}}^{\mathrm{cut}}$ as a function of $p_{\mathrm{T}}^{\mathrm{cut}}$. Additionally, in the same figure we show the statistical error evaluated assuming the LHC luminosity and the branching ratio quoted above. The one-loop correction is evidently bigger than the statistical error in the entire range of p_{T} considered

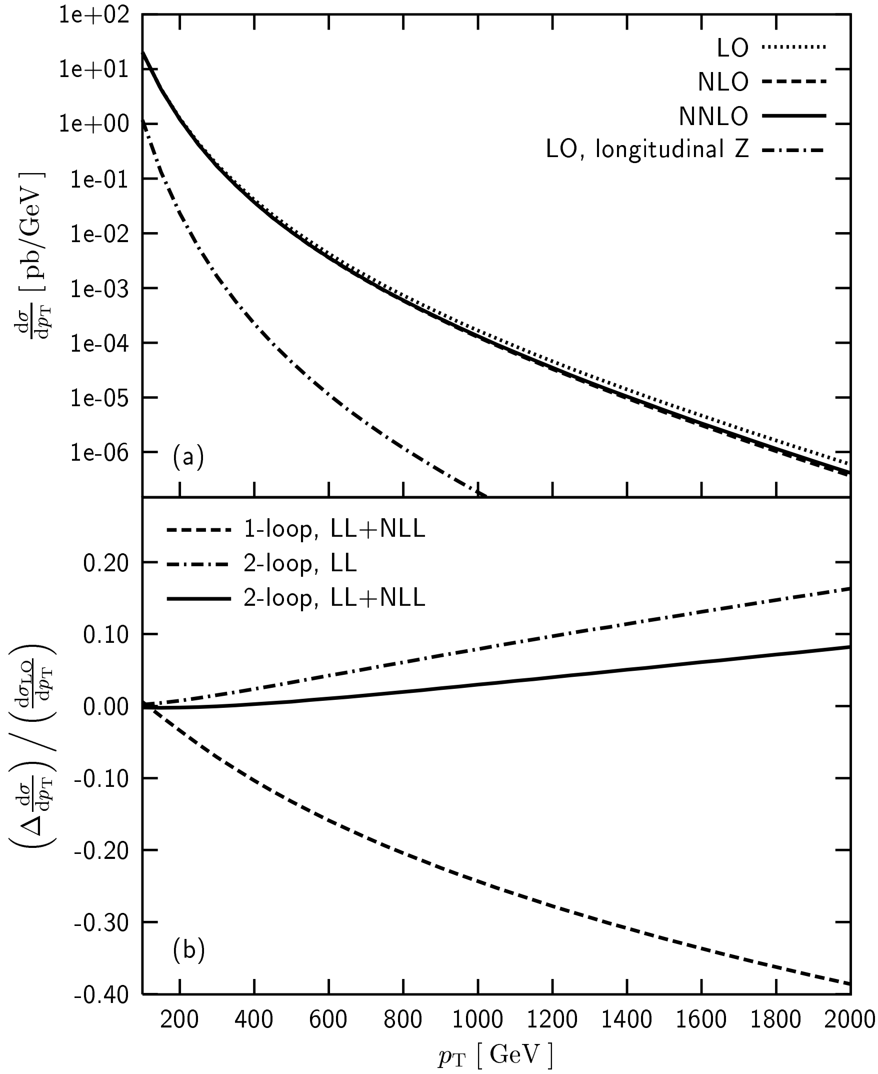

Figure 1: (a) Transverse momentum distribution for $pp \to Zj$ at $\sqrt{s} = 14\,\text{TeV}$: LO (dotted), NLO (dashed) and NNLO (solid) result for unpolarised Z and LO contribution from longitudinally polarised Z (dash-dotted). (b) Relative electroweak correction to the lowest order unpolarised p_T distribution for $pp \to Zj$ at $\sqrt{s} = 14\,\text{TeV}$: 1-loop LLs+NLLs (dashed), 2-loop LLs (dash-dotted) and 2-loop LLs+NLLs (solid).

here. The two-loop correction is larger than (or comparable to) the anticipated statistical error for values of p_T up to about 1500 GeV. In fact, the correction from the LL term is significantly larger, being, however, strongly reduced by the NLL correction. This demonstrates that at the two-loop level at least leading- and next-to-leading logarithms will be relevant for a reliable prediction. Of course, a full estimate of the experimental error requires including systematic effects what leads to an increase in the error value. Yet, Fig. 2 indicates relevance of the electroweak corrections for the precise analysis of experimental data.

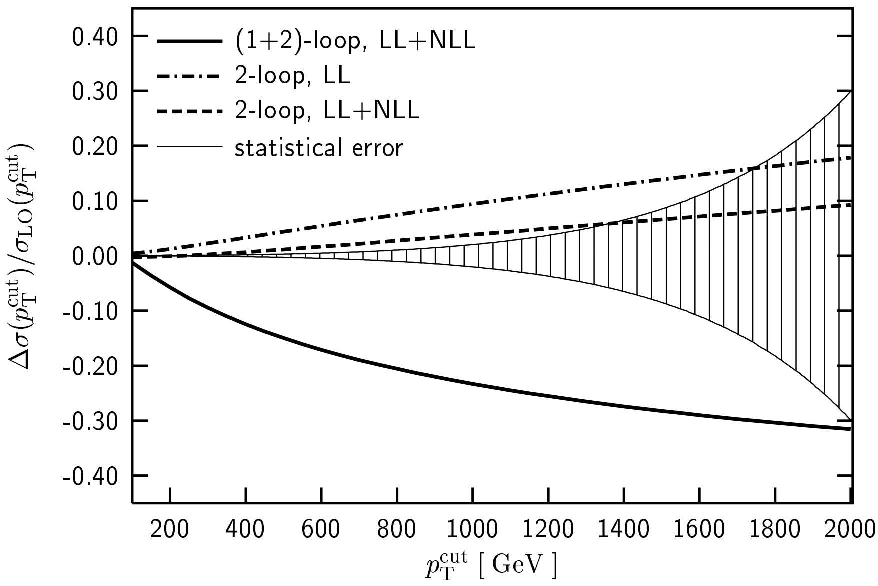

Figure 2: Relative electroweak correction and statistical error for the unpolarised integrated cross section for $pp \to Zj$ at $\sqrt{s} = 14\,\mathrm{TeV}$ as a function of $p_{\mathrm{T}}^{\mathrm{cut}}$: (1+2)-loop LL+NLL (solid), 2-loop LL (dash-dotted) and 2-loop LL+NLL (dashed) correction and statistical error (shaded region) with respect to the lowest order cross section.

We also compared the lowest order p_{T} distribution and one-loop logarithmic correction for the LHC with the results of Ref. [23]. In their calculations the authors of Ref. [23] introduced rapidity cuts, chose the scale of running α_{S} to be $\mu^2 = \hat{s}$, and while evaluating the one-loop corrections did not consider the electromagnetic part but only the weak contributions. Taking into account these differences we find a good agreement for the lowest order p_{T} distribution and one-loop corrections, at the level of a per cent for the latter. This shows that the main part of the one-loop EW correction is given by the logarithmic contribution.

Results for the $Z + 1$ jet production at the Tevatron are shown in Fig. 3a,b and in Fig. 4 for the p_{T} distribution and the integrated cross section, respectively. Contrary to the LHC case, at the Tevatron the one- and two-loop logarithmic electroweak corrections to the cross section do not bear much significance for the precise measurement, as illustrated in Fig. 4. Comparing our results with those of Ref. [23] we note a disagreement for the lowest order p_{T} distribution by roughly a factor of ten at high p_{T} values.

4 Conclusions

Experiments at the LHC will for the first time explore energies in the range well beyond 1 TeV. In this region, where $\hat{s} \gg M_W^2 \sim M_Z^2$, the electroweak radiative corrections become important due to large logarithms of $\hat{s}/M_W^2$. In the present paper we have studied their impact on hadronic Z-boson production at large transverse momentum. We have calculated the one- and two-loop electroweak corrections in logarithmic approximation, in-

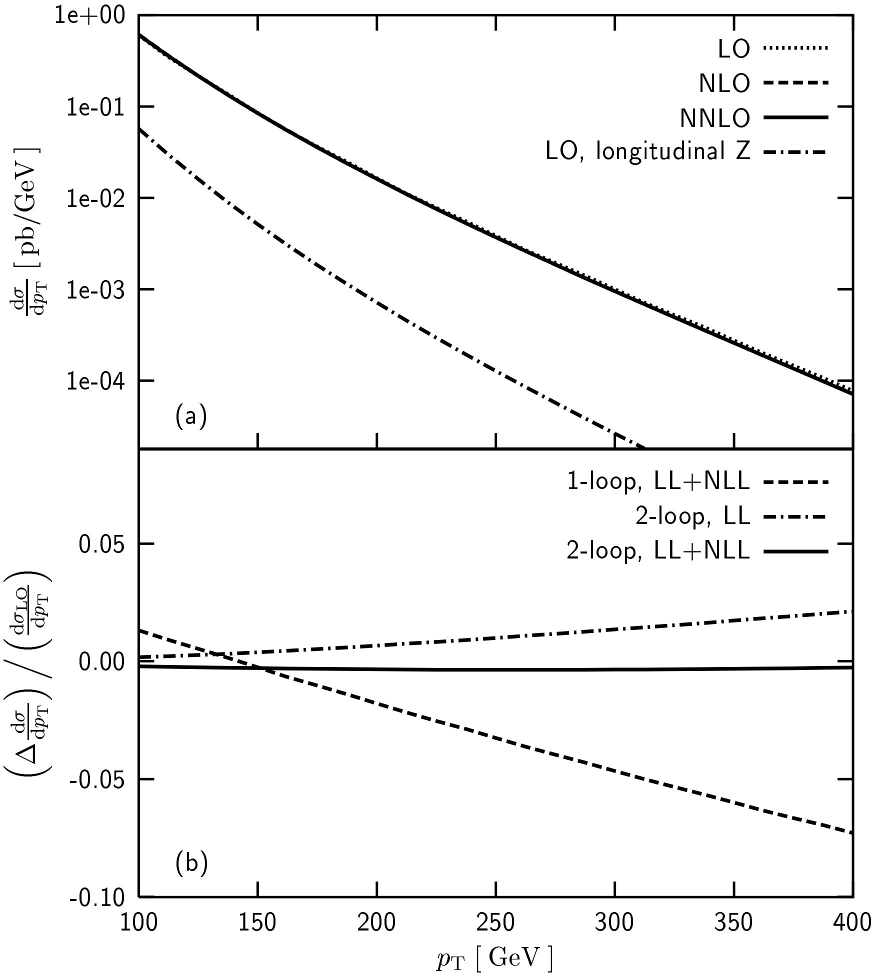

Figure 3: (a) Transverse momentum distribution for $p\bar{p} \to Zj$ at $\sqrt{s} = 2\,\text{TeV}$: LO (dotted), NLO (dashed) and NNLO (solid) result for unpolarised Z and LO contribution from longitudinally polarised Z (dash-dotted). (b) Relative electroweak correction to the unpolarised lowest order p_T distribution for $p\bar{p} \to Zj$ at $\sqrt{s} = 2\,\text{TeV}$: 1-loop LLs+NLLs (dashed), 2-loop LLs (dash-dotted) and 2-loop LLs+NLLs (solid).

cluding leading and next-to-leading logarithmic terms. These corrections increase rapidly with p_T. At the LHC, for values of p_T beyond 1 TeV the one-loop terms amount to several tens of per cent, reaching up to 40 per cent at 2 TeV. These results are in good agreement with those based on an explicit one-loop calculation [23]. To fully control the prediction with a precision that may be reached in the final stage of LHC operation, even the dominant two-loop terms must be included. Significant compensations between leading and next-to-leading two-loop logarithms are observed. The combined two-loop effect of several per cent is comparable to the anticipated statistical error at the LHC. Contrary to the LHC, at the Tevatron the one- and two-loop logarithmic electroweak corrections do not have numerical significance.

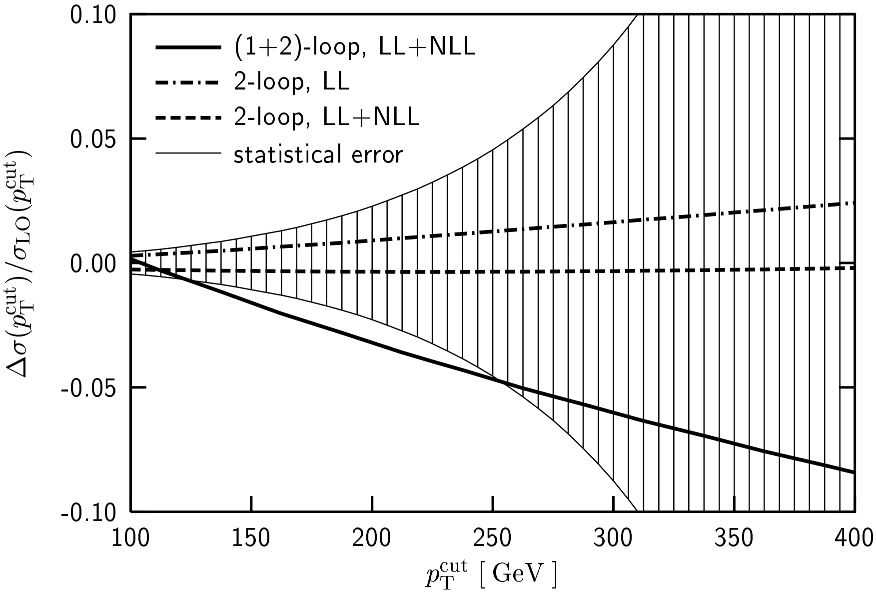

Figure 4: Relative electroweak correction and statistical error for the unpolarised cross section for $p\bar{p} \to Zj$ at $\sqrt{s} = 2\,\mathrm{TeV}$ as a function of $p_{\mathrm{T}}^{\mathrm{cut}}$: (1+2)-loop LL+NLL (solid), 2-loop LL (dash-dotted) and 2-loop LL+NLL (dashed) correction and statistical error (shaded region) with respect to the lowest order cross section.

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References

[1] S. Catani *et al.*, hep-ph/0005025.

[2] R. K. Ellis, G. Martinelli and R. Petronzio, Nucl. Phys. B **211** (1983) 106;
P. B. Arnold and M. H. Reno, Nucl. Phys. B **319**, 37 (1989) [Erratum-ibid. B **330**, 284 (1990)];
R. J. Gonsalves, J. Pawlowski and C. F. Wai, Phys. Rev. D **40** (1989) 2245.

[3] W. T. Giele, E. W. N. Glover and D. A. Kosower, Nucl. Phys. B **403** (1993) 633 [hep-ph/9302225].

[4] J. Campbell and R. K. Ellis, Phys. Rev. D **65** (2002) 113007 [hep-ph/0202176].

[5] J. Campbell, R. K. Ellis and D. L. Rainwater, Phys. Rev. D **68** (2003) 094021 [hep-ph/0308195].

[6] A. Denner and S. Pozzorini, *Eur. Phys. J. C* **18** (2001) 461 [hep-ph/0010201]; *Eur. Phys. J. C* **21** (2001) 63 [hep-ph/0104127].

[7] S. Pozzorini, *doctoral thesis, Universität Zürich, 2001*, hep-ph/0201077.

[8] E. Accomando, A. Denner and S. Pozzorini, *Phys. Rev. D* **65** (2002) 073003 [hep-ph/0110114].

[9] W. Hollik and C. Meier, Phys. Lett. B **590** (2004) 69 [hep-ph/0402281].

[10] J. H. Kühn and A. A. Penin, hep-ph/9906545.

[11] P. Ciafaloni and D. Comelli, Phys. Lett. B **476** (2000) 49 [hep-ph/9910278].

[12] V. S. Fadin *et al.*, *Phys. Rev. D* **61** (2000) 094002 [hep-ph/9910338].

[13] J. H. Kühn, A. A. Penin and V. A. Smirnov, *Eur. Phys. J. C* **17** (2000) 97 [hep-ph/9912503].

[14] M. Melles, *Phys. Rev. D* **63** (2001) 034003 [hep-ph/0004056]; *Phys. Rev. D* **64** (2001) 014011 [hep-ph/0012157]; *Phys. Rev. D* **64** (2001) 054003 [hep-ph/0102097]; *Phys. Rept.* **375** (2003) 219 [hep-ph/0104232]; *Eur. Phys. J. C* **24** (2002) 193 [hep-ph/0108221].

[15] J. H. Kühn *et al.*, *Nucl. Phys. B* **616** (2001) 286 [hep-ph/0106298].

[16] M. Melles, *Phys. Lett. B* **495** (2000) 81 [hep-ph/0006077].

[17] M. Hori, H. Kawamura and J. Kodaira, *Phys. Lett. B* **491** (2000) 275 [hep-ph/0007329].

[18] W. Beenakker and A. Werthenbach, *Phys. Lett. B* **489** (2000) 148 [hep-ph/0005316]; *Nucl. Phys. B* **630** (2002) 3 [hep-ph/0112030].

[19] A. Denner, M. Melles and S. Pozzorini, Nucl. Phys. B **662** (2003) 299 [hep-ph/0301241].

[20] S. Pozzorini, Nucl. Phys. B **692** (2004) 135 [hep-ph/0401087].

[21] B. Feucht, J. H. Kühn and S. Moch, Phys. Lett. B **561** (2003) 111 [hep-ph/0303016].

[22] B. Feucht, J. H. Kühn, A. A. Penin and V. A. Smirnov, hep-ph/0404082.

[23] E. Maina, S. Moretti and D. A. Ross, hep-ph/0403050.

[24] V. A. Mitsou, hep-ph/0004161.

[25] Department of Energy Review Committee Report on the Run II Luminosity Performance of the Fermilab Tevatron, October 28-31 2002, http://www-bd.fnal.gov/doereview02/Run_II_lum_review_Final.pdf

[26] A. D. Martin, R. G. Roberts, W. J. Stirling and R. S. Thorne, Phys. Lett. B **531** (2002) 216 [hep-ph/0201127].