Physics Letters B 292 (1992) 164–168
North-Holland

PHYSICS LETTERS B

Jet photoproduction at HERA in next-to-leading-order QCD

D. Bödeker [1]
Deutsches Elektronen-Synchrotron DESY, Notkestraße 85, W-2000 Hamburg 52, FRG

Received 15 April 1992; revised manuscript received 30 July 1992

The photoproduction of jets via direct photons at HERA in next-to-leading-order QCD is considered. For the one-jet inclusive cross section the scale dependence and the influence of the jet cone size is investigated. Jet transverse energy and rapidity distributions are discussed.

HERA will provide a unique opportunity to study the photoproduction of jets with large transverse energies [1–4] which will allow for a quantitative test of perturbative QCD.

In this process the electron emits a nearly on-shell photon which scatters on some parton in the proton producing a large transverse (to the beam direction) energy final state (fig. 1). Neglecting the transverse momentum of the photon the cross section is obtained by folding the photon–proton cross section with a Weizsäcker–Williams [5] distribution function. Besides this direct photon contribution there is also a contribution due to its quark and gluon content described by parton densities inside the photon (resolved photon). Since the distribution of the photon in the electron is harder than that of the quarks and gluons the direct contribution dominates at large transverse energy E_T and rapidity of the jet $\eta=-\ln(\tan\frac{1}{2}\theta)$ (positive η corresponds to the electron/photon direction). At HERA both contributions are of similar size for $E_T=35$ GeV, $\eta=-1$ and $E_T=45$ GeV, $\eta=-2$. The distinction between them is unambiguously defined only for the Born cross section because in higher orders (cf. refs. [6,7]) the photon can split up into a collinear quark–antiquark pair leading to a collinear singularity. This is already taken care of by the photon structure function and one has to use some factorization prescription to subtract this contribution.

[1] Supported by Bundesministerium für Forschung und Technologie, 06 OH 755.

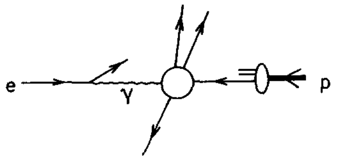

Fig. 1. Jet photoproduction in electron–proton collisions.

In leading order (LO) the jet photoproduction occurs via the $O(\alpha\alpha_s)$ subprocesses $\gamma q\to qg$ and $\gamma g\to q\bar{q}$. The one-jet inclusive cross section for $ep\to jet+X$ for a jet with given E_T and η, which is related to its ep CMS rapidity η^* by

$$\eta=\eta^*-\tfrac{1}{2}\ln\frac{E_p}{E_e}, \tag{1}$$

can be written as

$$\frac{d^2\sigma}{dE_T d\eta} = \frac{1}{8\pi}\frac{E_T}{s^2}\sum_i \int_{\bar\eta^*_{min}}^{\bar\eta^*_{max}} d\bar\eta^* \frac{f_{\gamma/e}(x_a)f_{i/p}(x_b)}{x_a x_b}\,\overline{|M|^2}, \tag{2}$$

where $\bar\eta^*$ is the ep CMS rapidity of the second parton with large transverse momentum and

$$\bar\eta^*_{min}=-\ln\left(\frac{\sqrt{s}}{E_T}-\exp(-\eta^*)\right),$$

$$\bar\eta^*_{max}=\ln\left(\frac{\sqrt{s}}{E_T}-\exp(\eta^*)\right), \tag{3}$$

0370-2693/92/$ 05.00 © 1992 Elsevier Science Publishers B.V. All rights reserved.

the sum runs over all partons from the proton and the color- and spin-averaged matrix element $\overline{|M|^2}$ is symmetrized in the final state momenta. By energy–momentum conservation all invariants including the momentum fractions x_a and x_b are determined by E_T, η^*, $\bar{\eta}^*$ and the ep CMS energy $\sqrt{s}$:

$$x_a = \frac{E_T}{\sqrt{s}} [\exp(\eta^*) + \exp(\bar{\eta}^*)] ,$$

$$x_b = \frac{E_T}{\sqrt{s}} [\exp(-\eta^*) + \exp(-\bar{\eta}^*)] . \qquad (4)$$

For the Born cross section one identifies jet and parton, so the transverse momentum p_T equals E_T (assuming massless partons), while in the next-to-leading order the jet may consist of two partons so that in general E_T and p_T are no longer equal.

In next-to-leading order (NLO) one gets the $O(\alpha\alpha_s^2)$ interference terms of the Born and one-loop diagrams for the leading order subprocesses. For $\gamma q \to qg$ they have been calculated in ref. [6]. For $\gamma g \to q\bar{q}$ they had to be calculated because by crossing their result for $\gamma q \to qg$ one cannot take care of the so-called π^2 terms. By keeping the imaginary parts in the arguments of the logarithms I could reproduce the matrix element for $\gamma q \to qg$ in agreement with ref. [6].

Furthermore one has to take into account the real $O(\alpha\alpha_s^2)$ $2 \to 3$ processes

(i) $\gamma q \to qgg$,
(ii) $\gamma g \to q\bar{q}g$,
(iii) $\gamma q \to qq'\bar{q}'$,
(iv) $\gamma q \to qq\bar{q}$.

The corresponding matrix elements were obtained in ref. [6]. I confirmed their results up to some misprints which where also noticed in ref. [7].

For the phase space integration of the $2 \to 3$ processes one has to choose some jet definition to decide whether two of the produced partons count as one jet or two separate ones. The definition adopted here is the one used in ref. [8]. It is invariant under boosts along the beam direction and therefore well suited for an asymmetric machine as HERA. Two partons with transverse energies E_{T_1}, E_{T_2} (all partons are assumed massless so these are equal to their transverse momenta), rapidities η_1, η_2 and azimuthal angles ϕ_1, ϕ_2 are combined to one jet if they both fit into a jet cone with radius R in η–ϕ-space around (η, ϕ) i.e.

$$(\eta_i - \eta)^2 + (\phi_i - \phi)^2 < R^2 \quad (i = 1, 2) . \qquad (5)$$

The cone direction (η, ϕ) is chosen such that

$$\frac{1}{E_{T_1} + E_{T_2}} [E_{T_1}(\eta_1, \phi_1) + E_{T_2}(\eta_2, \phi_2)] = (\eta, \phi) . \qquad (6)$$

This defines the jet direction as (η, ϕ), the jet transverse energy is given by $E_T = E_{T_1} + E_{T_2}$, which for small R becomes equal to its transverse momentum.

The phase space integration of the $2 \to 3$ processes contains soft and collinear singularities. By extracting the singular denominators of the integrand and setting the kinematic variables to their singular values except in the denominator, as described in ref. [8] for hadron–hadron collisions, one gets terms with the same singular behavior as the integrand. These terms are subtracted from the integrand which in this way becomes regular on the whole phase space so that it can be integrated numerically. They are then integrated in $d = 4 - 2\epsilon$ dimensions and the singularities, which appear as poles in ϵ, are cancelled against those of the virtual corrections. The remaining singularities from collinear radiation by the incoming photon/parton from the proton were factorized and absorbed by a renormalization of the structure functions. For both photon and proton the $\overline{\rm MS}$ prescription [9] has been used. For the plots the HMRS set B [10] parton distributions for the proton and the two-loop running coupling for both LO and NLO have been used, the number of flavors has been set equal to four.

In LO and NLO the cross section depends on the renormalization scale μ and the factorization scale M_p of the proton, while the NLO also depends on the factorization scale for the photon M_γ. If not stated otherwise, they are all set equal to ξE_T, where we choose ξ to vary between $\frac{1}{8}$ and 4. In fig. 2 the ξ-dependence of $d^2\sigma/dE_T\,d\eta$ is shown for $E_T = 45$ GeV, $\eta = -1$ and jet radius $R = 1$ at HERA energy. The Born cross section (dotted line) varies by a factor >2 in the ξ-range under consideration. If one adds the $O(\alpha\alpha_s^2)$ corrections (full line) the scale dependence is hardly reduced. The LO and NLO results differ by a nearly constant value, which, of course, depends on the jet radius. However, since the LO cross section is independent of M_γ, there are no terms compensating the variation of the $O(\alpha\alpha_s^2)$ corrections with M_γ. This is demonstrated by the dash dotted line in fig. 2 where

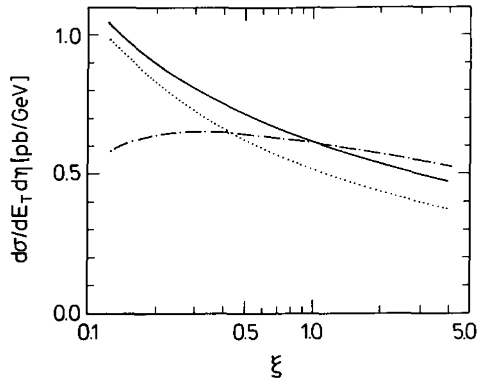

Fig. 2. Scale dependence of $d^2\sigma/dE_T\,d\eta$: $\xi=\mu/E_T=M_p/E_T$. Dotted line: LO, full line: NLO with $M_\gamma/E_T=\xi$, dash-dotted line: NLO with $M_\gamma=E_T$ fixed. $E_T=45$ GeV, $\eta=-1$, $R=1$, $\sqrt{s}=314$ GeV.

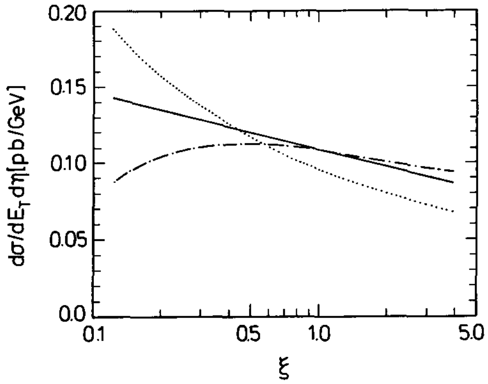

Fig. 3. The same as fig. 2 for $E_T=60$ GeV.

$M_\gamma=E_T$ is fixed and only μ and M_p are varied. The sensitivity to the choice of ξ is drastically reduced compared to the Born cross section, the cross section varies by about $\simeq 20\%$ for $\frac{1}{8}<\xi<4$. There is a maximum near $\xi\simeq\frac{1}{2}$, a behavior which was also observed in ref. [11] for $p\bar{p}$ collisions. At $E_T=45$ GeV, $\eta=-1$ the direct photon contribution is dominant but the resolved one is of the same order of magnitude. Since there the NLO cross section for the resolved part is monotonically increasing with ξ [3], one may expect the sum of direct and resolved contribution to be less sensitive to the choice of the scales.

At $E_T=60$ GeV (see fig. 3) the Born cross section shows a weaker scale dependence and by including the $O(\alpha\alpha_s^2)$ corrections it is further reduced. Again for fixed M_γ the result is rather stable at $\xi\simeq\frac{1}{2}$, so in all other plots $\mu=M_p=M_\gamma=\frac{1}{2}E_T$ has been used.

The M_γ-dependence is compensated in each order in α_s by the resolved photon contribution. Simply adding the $O(\alpha\alpha_s)$ resolved part will reduce the dependence on M_γ. On the other hand the corresponding hard scattering process is of the order α_s^2 and so the dependence on the renormalization scale μ is even larger than for the direct contribution. Therefore the inclusion of the LO resolved contribution will not help to reduce the scale dependence without taking into account the higher order corrections.

In fig. 4 $d^2\sigma/dE_T\,d\eta$ is plotted versus the jet cone size R. For the NLO one sees the typical behavior known from $p\bar{p}$ collisions [11] which is approximately given by

$$\frac{d^2\sigma}{dE_T\,d\eta}\simeq a+b\ln R+cR^2\,. \tag{7}$$

Fig. 5 shows the E_T-distribution of $d^2\sigma/dE_T\,d\eta$ down to $E_T=20$ GeV. For $E_T>40$ GeV the $O(\alpha\alpha_s^2)$ corrections are smaller than 20%. Below 40 GeV the full $O(\alpha\alpha_s^2)$ cross section becomes much larger than the LO result. The direct and resolved photon contribution have a different angular distribution [2] so it is of interest how the $O(\alpha\alpha_s^2)$ corrections modify the rapidity spectrum. In fig. 6 $d^2\sigma/dE_T\,d\eta$ is plotted versus η. The form of the distribution is not changed, the full $O(\alpha\alpha_s^2)$ cross section is slightly more peaked than the LO result.

The direct jet photoproduction cross section for fixed photon energy has previously been calculated in NLO in ref. [7]. Here the method of "parton res-

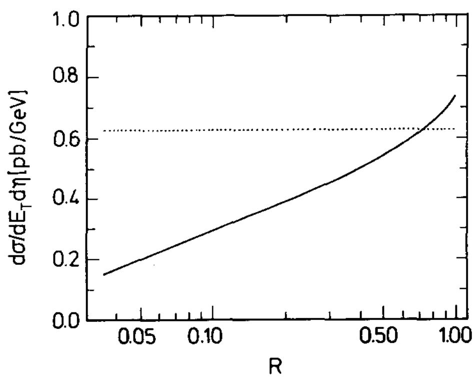

Fig. 4. Dependence of $d^2\sigma/dE_T\,d\eta$ on the jet cone size R. Dotted line: LO, full line: NLO. $E_T=45$ GeV, $\eta=-1$, $\sqrt{s}=314$ GeV.

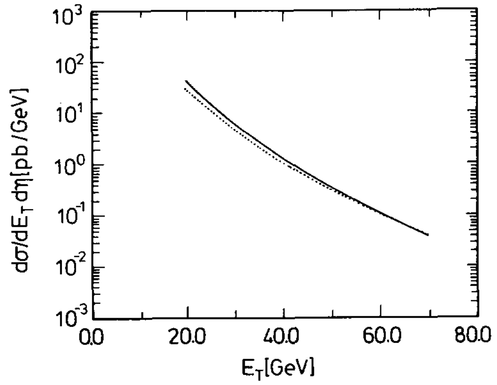

Fig. 5. $d^2\sigma/dE_T\,d\eta$ versus the transverse energy E_T. Dotted line: LO, full line: NLO. $\eta=-1$, $R=1$, $\sqrt{s}=314$ GeV.

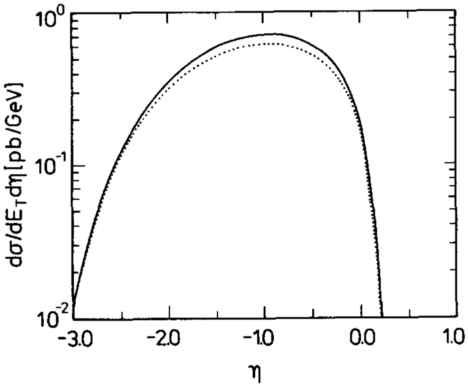

Fig. 6. $d^2\sigma/dE_T\,d\eta$ versus the jet rapidity η. Dotted line: LO, full line: NLO. $E_T=45$ GeV, $R=1$, $\sqrt{s}=314$ GeV.

olution parameters" (cf. ref. [12]) has been used instead of the subtraction method [8,13] as in the present calculation in order to make the phase space integral of the $2\to3$ processes finite. By replacing the Weizsäcker–Williams distribution function by a deltafunction the present calculation may be compared with ref. [7]. To obtain their numerical results these authors eliminated the collinear initial state singularity of the photon by demanding that there is no photon remnant jet, i.e. no jet with rapidity larger than $\eta_{\text{cut}}=2$. This requires a precise definition since one has to allow the soft gluons to be radiated in any direction in order to cancel the soft singularities of the virtual corrections. Although not explicitly stated in ref. [7] the authors proceed as follows [14]: Quarks are required to have $\eta<\eta_{\text{cut}}$ while gluons with $\eta>\eta_{\text{cut}}$ and $E<\frac{1}{2}\delta_s\sqrt{\hat{s}}$ are included, where E is the gluon energy in the parton CMS frame and δ_s is the "parton resolution parameter" defining the soft region of the three-particle phase space, which is possible since at $O(\alpha\alpha_s^2)$ the collinear emission of gluons is not singular. In this way δ_s gets a physical meaning and the cross section depends on δ_s. For large values of η_{cut}, however, this dependence is small. To compare with their results I have also included quarks with $\eta>\eta_{\text{cut}}$ and $E<\frac{1}{2}\delta_s\sqrt{\hat{s}}$, the collinear singularity is removed by the $\overline{\text{MS}}$ counterterm. In order to be close to the jet definition of ref. [7] two partons i, j are combined to one jet if $(\eta_i-\eta_j)^2+(\phi_i-\phi_j)^2<1$ while the jet direction is defined by eq. (6). Using $\mu=M_p=M_\gamma$ and $\delta_s=0.05$ I then find good agreement with the numerical results of ref. [7].

To summarize, the full $O(\alpha\alpha_s^2)$ jet-photoproduction cross section for the direct photon at HERA has been calculated. At very high transverse energies the scale dependence is reduced by including the $O(\alpha\alpha_s^2)$ corrections, at medium E_T the cross section is quite sensitive to the factorization of the collinear singularity of the photon. The dependence on the jet cone size shows the expected behavior. For an appropriate choice of the renormalization and factorization scales, which was suggested by requiring stability against their variation, the corrections are small except at small transverse energies, where the resolved photon gives the dominant contribution to the jet photoproduction.

I thank W. Buchmüller and G.A. Schuler for helpful discussions.

References

[1] Z. Kunszt and W.J. Stirling, in: Proc. HERA Workshop, Vol. 1 (1987), ed. R.D. Peccei.
[2] H. Baer, J. Ohnemus and J.F. Owens, Z. Phys. C 42 (1989) 657.
[3] G. Kramer and S.G. Salesch, in: Proc. Workshop Physics at HERA, Vol. 1 (1991), eds. W. Buchmüller and G. Ingelman.
[4] D. Bödeker, in: Proc. Workshop Physics at HERA, Vol. 1 (1991), eds. W. Buchmüller and G. Ingelman.
[5] C.F. von Weizsäcker, Z. Phys. 88 (1934) 612;
E.J. Williams, Kgl. Danske Vidensk. Selskab. Mat.-Fiz. Medd. 13 (1935) N4.

[6] P. Aurenche, R. Baier, A. Douiri, M. Fontannaz and D. Schiff, Nucl. Phys. B 286 (1987) 553.
[7] H. Baer, J. Ohnemus and J.F. Owens, Phys. Rev. D 40 (1989) 2844.
[8] S.D. Ellis, Z. Kunszt and D.E. Soper, Phys. Rev. D 40 (1989) 2188.
[9] B. Curci, W. Furmanski and R. Petronzio, Nucl. Phys. B 175 (1980) 27;
L. Baulieu, E.G. Floratos and C. Kounnas, Nucl. Phys. B 166 (1980) 321;
J.C. Collins and D.E. Soper, Nucl. Phys. B 194 (1982) 445.
[10] P. Harriman, A. Martin, R. Roberts and J. Stirling, Phys. Rev. D 42 (1990) 798.
[11] S.D. Ellis, Z. Kunszt and D.E. Soper, Phys. Rev. Lett. 64 (1990) 2121.
[12] W.T. Giele and E.W.N. Glover, report FERMILAB-Pub-91/100-T.
[13] R.K. Ellis, D.A. Ross and A.E. Terrano, Nucl. Phys. B 178 (1981) 421.
[14] J. Owens, private communication.