Jet photoproduction and the structure of the photon

B.W. Harris and J.F. Owens

Physics Department
Florida State University
Tallahassee, Florida 32306-4350, USA
(October 1997)

Abstract

Various jet observables in photoproduction are studied and compared to data from HERA. The feasibility of using a dijet sample for constraining the parton distributions in the photon is then studied. For the current data the experimental and theoretical uncertainties are comparable to the variation due to changing the photon parton distribution set.

PACS number(s): 12.38.Bx,13.60.Hb,14.70.Bh
I. INTRODUCTION

Large momentum transfer processes involving photons are useful for studying strong interaction dynamics, for studying hadronic structure, and for refining our understanding of the photon-hadron interaction itself. The photoproduction of hadronic jets is an example of such a process. In a previous paper [1], a complete $O(\alpha_s^3)$ calculation of jet photoproduction, based on the phase space slicing technique, was described and predictions were compared with data for several single jet and dijet observables. Since the completion of that analysis additional data have become available. The new data have increased statistics and have been analyzed using several different jet definitions. Results for several new distributions involving dijets are also now available. Therefore, it is appropriate to investigate the degree to which the new data are described using the program described in Ref. [1].

An important aspect of hard photon-hadron interactions is the existence of two different interaction mechanisms. The “direct” component corresponds to the case where the photon participates wholly in the hard scattering, giving its full energy to the underlying photon-parton subprocesses. The “resolved” component of the photon [2] corresponds to the case where the photon generates a shower in the initial state of nearly collinear quarks and gluons, one of which interacts with a parton from the initial state hadron. In this case it is convenient to define parton distributions in the photon and to sum the resulting collinear logarithms using the Altarelli-Parisi formalism appropriately extended to the photon case [3]. Information on the parton distributions in the photon is available from studies of the photon structure function $F_2^\gamma$, jet production in $e^+e^-$ two photon collisions, heavy quark production in $e^+e^-$ two photon collisions, and in jet photoproduction. For recent overviews, see, for example Refs. [4] and [5]. For the first two processes, the lowest order diagrams involve only quarks. On the other hand, the resolved component of heavy quark production in two photon collisions [6] and jet photoproduction both in lowest order include subprocesses with initial state gluons. Therefore, at least in principle, they should provide information on the gluon distribution in the photon which would complement that obtained from the other two
processes. Some of the issues discussed in this paper have been reviewed in [7].

The goal of this paper is to critically examine the current theoretical description of jet photoproduction and, as a result, assess the ability of the data to distinguish between different parametrizations of photon parton distributions. In Sec. II the predictions for inclusive single jet production are compared with the latest available data. A similar comparison for various dijet observables is presented in Sec. III. Conclusions and suggestions for further study are given in Sec. IV.

II. SINGLE JET INCLUSIVE DISTRIBUTIONS

If jet photoproduction is to serve as a reliable source of information on the parton distributions in the photon, then it is necessary to assess the quality of the theoretical description of the data. Differences between the theory and the data could be due to several sources, e.g., an inadequate description of the underlying hard scattering subprocesses or incorrect parton distributions in the initial state photon or hadron. In order to examine these issues, we first consider the inclusive single jet rapidity distributions.

The predictions for all of the jet observables in this paper have been generated with one or more variants of the cone algorithm for jet merging described in Ref. [1]. This algorithm depends on a parameter $R_{\text{sep}}$ [8] governing the merger of two partons that are widely separated but, none the less, lie within a cone of the prescribed radius $R$. This variable lies in the range $R \leq R_{\text{sep}} \leq 2R$. For $R_{\text{sep}} = 2R$ the algorithm reduces to that of Ref. [9] and the introduction of $R_{\text{sep}}$ has no effect. At the other extreme, $R_{\text{sep}} = R$, the algorithm applied to three parton final states is equivalent to the $k_T$ algorithm [10].

Recent studies of jet shapes [11], [12] find good agreement between data analyzed using the iterative cone algorithm and theory if the problem of merging overlapping jets is taken into account by varying the parameter $R_{\text{sep}}$ as a function of transverse energy and rapidity. The optimal values of $R_{\text{sep}}$ varied slowly with $E_T$ and $\eta_{\text{jet}}$ and values in the range of 1.3 – 1.4 were obtained for most of the region covered by the data.
In Fig. 1 the data for the single-jet inclusive cross section as a function of the jet pseudorapidity $\eta_{\text{jet}}$ integrated over $E_T^{\text{jet}} > E_T^{\text{min}}$ for $E_T^{\text{min}} = 14, 17, 21, \text{ and } 25$ GeV as measured by ZEUS [13] using an iterative cone jet finding algorithm with $R = 1$ are shown. These data are compared with the results of our next-to-leading calculation using the Aurenche-Fontannaz-Guillet (AFG) [14] photon distributions and the CTEQ4M [15] proton parton distributions. The three curves correspond to $R_{\text{sep}} = 2$ (top), $R_{\text{sep}} = 1.3$ (middle), and $R_{\text{sep}} = 1$ (bottom). The factorization and renormalization scales have both been chosen to be $E_T^{\text{max}}$ where $E_T^{\text{max}}$ is the largest transverse jet energy in the event. This comparison clearly shows that the theory underestimates the data in the region $\eta_{\text{jet}} > 0.5$ for the lowest $E_T^{\text{min}}$ set, with the discrepancy decreasing with increasing $E_T^{\text{min}}$. This result is consistent with that presented in [1]. One possibility for this discrepancy is that in the forward region fragments of the proton beam jet are being included in the jet cone and, therefore, the energy of the observed jet is increased. The effect of such a contribution would be decreased if a smaller cone radius was used to define the jet. Therefore, data have also been presented corresponding to a cone of 0.7. These are shown in Fig. 2 together with the corresponding theoretical predictions. The agreement is now much improved in the forward direction.

The use of a smaller cone size to define the jets appears to solve the problem of the jet excess in the forward $\eta_{\text{jet}}$ region. The agreement between theory and experiment shown in Fig. 2 increases the confidence one has in the theoretical description of jet photoproduction. However, the inclusive single-jet observables still involve integrations over unobserved jets with a corresponding loss of information. In order to further test the theoretical description it is necessary to look at observables which involve more than one jet. While there will still be integrations over unobserved jets or partons, dijet observables offer more ways to test the theory.
A fundamental requirement for utilizing jet photoproduction as a means of constraining photon parton distributions is that the underlying hard scattering subprocesses must be correctly described. A sensitive test of this is to study the dijet angular distribution in the dijet center of mass system. In lowest order, dijet production is described in terms of $2 \rightarrow 2$ subprocesses and the two jets will be back-to-back with balancing transverse momenta. The angular distribution of the dijet axis with respect to the beam is a fundamental prediction of the theory and its shape is only slightly modified when higher order effects are taken into account [1].

The top portion of Fig. 3 shows the inclusive dijet cross section as a function of the absolute value of the cosine of the angle between the dijet axis and the beam axis $|\cos \theta^*|$ with the dijet mass $M_{JJ}$ greater than 47 GeV. The data were measured by ZEUS [16] using their iterative cone algorithm with $R = 1$ and are compared with our next-to-leading order result for $R = 1$ with $R_{\text{sep}} = 2R$ (solid) and $R_{\text{sep}} = R$ (dot). This particular observable was found to be relatively insensitive to the choice of the photon parton distribution set. The predictions shown here were obtained using the Gordon-Storrow (GS) distributions [17]. The overall agreement is good, with the trend of the data being between the curves for the two values of $R_{\text{sep}}$. The good description of the angular distribution data suggests that the underlying subprocesses are correctly described.

Next, consider the dependence of the cross section on the mass of the dijet system. In lowest order the dijet mass is given by $M_{JJ} = \sqrt{x_\gamma x_p s}$ where $x_\gamma$ and $x_p$ are the momentum fractions carried by the partons from the photon and proton, respectively, and $s$ is the square of the photon-proton center of mass energy. Hence, the dijet mass distribution will be sensitive to the shapes of the parton distributions. Since $\cos \theta^*$ is being integrated over, the shape of the dijet mass distribution does not depend too sensitively on the shape of the angular distribution.

The lower part of Fig. 3 shows the same data as in the upper part, but integrated over
cos θ* and displayed versus M_{JJ}. Again, the agreement is good, lending support for the description of the overall production process.

We next consider the $E_T$ dependence of the highest $E_T$ jet in dijet events with the pseudorapidities of the two jets constrained to be in specified ranges. Data measured by ZEUS [16] are shown in Fig. 4 together with our theoretical results generated with the GS photon distributions and with both the renormalization and factorization scales equal to $E_T^{\text{max}}$. The events are symmetrized in $\eta_1$ and $\eta_2$ so there are two entries per event. Furthermore, the highest $E_T$ jet is required to have $E_{T_1} > 14$ GeV while the second jet can have an $E_T$ as low as 11 GeV. This asymmetric $E_T$ requirement avoids a potential problem with the definition of the dijet sample [1,18,19]. The lower curve of each pair has an additional cut of $x_\gamma > 0.75$ placed on it, where $x_\gamma$ is the momentum fraction of the parton from the photon and is reconstructed using [20]

$$x_\gamma = (E_{T_1} e^{-\eta_1} + E_{T_2} e^{-\eta_2})/2E_\gamma$$

with $E_\gamma$ the photon energy in the lab frame. This cut removes much of the resolved component, so the lower curve of each pair is dominated by the direct contribution. The jet definition used for these data is the $k_\perp$ algorithm [10]. It has several advantages over the cone algorithm, although an extensive discussion of these points is beyond the scope of this work. The interested reader is pointed towards the original literature [10] and recent detailed studies [21], [22] for details. For the three parton final states considered in this calculation, the $k_\perp$ algorithm corresponds to setting $R_{\text{sep}} = R = 1$.

Over most of the $\eta_1$ and $\eta_2$ ranges shown in Fig. 4, the comparison between the theoretical results and the data appears to be good. The exception occurs when one or both of the jets has a negative pseudorapidity, with the discrepancy being worst if both jets have negative pseudorapidities. This point will be discussed in more detail below.

These same data are shown in Fig. 5 versus $\eta_2$ in bins of $\eta_1$ and integrated over $E_T > 14$ GeV. For each set of data, curves for three choices of the common renormalization and factorization scales are shown: $\mu = E_T^{\text{max}}/2, E_T^{\text{max}},$ and $2E_T^{\text{max}}$ all using the GS photon set.
In Fig. 6 these data are compared to the theoretical results obtained using the AFG [14], GS [17], and Glück-Reya-Vogt (GRV) [23] photon parton distributions all using $\mu = E_T^{\text{max}}$. From the results shown in Figs. 5 and 6, it can be seen that the variations due to the scale choice are comparable in magnitude to those resulting from the different choices of photon parton distributions. Furthermore, these variations are comparable to the experimental uncertainties on the data. However, it should be noted that the variation due to the change of scale leads to an overall shift in the normalization without an appreciable change in shape. Hence, even allowing for this scale uncertainty, constraints on the photon parton distributions are provided by the shapes of such observables. Generally speaking, the data are well described within the theoretical and experimental uncertainties. However, a clear discrepancy is apparent in the region of negative $\eta_1$ and $\eta_2$ (the left-most data point in the center and left-hand plots in these two figures).

In order to further examine this disagreement, the data are shown again in Fig. 7 and Fig. 8, the latter having the cut $x_\gamma > 0.75$ imposed. In each case, four curves are shown corresponding to the full result (solid line), the resolved component with only quarks in the photon (medium dashed line), the resolved component with only gluons in the photon (dotted line), and direct component (long dashed line). Several points are immediately clear. The negative pseudorapidity region is dominated by the direct component. Hence, no amount of variation of the photon parton distributions will bring the data and theory in line. Second, the gluon contribution is small nearly everywhere for the kinematics of the current data sample. Only for both rapidities in the forward direction (right-most region of the right-hand plots) is the gluon distribution in the photon making a significant contribution.

The HERWIG [24] Monte Carlo, which uses the parton-shower approach for intial-state and final-state QCD radiation including colour coherence and azimuthal correlations, both within and between jets, was also compared with the data. The results are shown in Fig. 9. The HERWIG results have been scaled upwards by a factor of 1.6, but the shape is in good agreement with the data. Next, Fig. 10 shows the ratio of the HERWIG results to the HERWIG results at the leading-order matrix element level, i.e., treating the two scattered
partons as being the jets with all showering and subsequent hadronization off. It can be seen that there is a correction which corresponds to a depletion of events in the negative rapidity region. This represents the effects of the showering and subsequent jet reconstruction slightly shifting the energies of the jets. Near the edge of the kinematic region covered by the data this results in a decrease in the cross section. This decrease makes the full HERWIG results steeper in this region and gives a better description of the data. Hence, it appears that the origin of the inability of the next-to-leading-order calculation to describe this region properly is that the single parton branching that occurs in a three-body final state does not give a sufficient description of the full showering.

An estimate of the amount of showering correction coming from the next-to-leading-order terms can be obtained by looking at the ratio of the results of our calculation using the full next-to-leading-order matrix elements to those obtained using the leading-order matrix elements only. This ratio shows to what extent the showering correction is modeled by the three-body final states and is shown in Fig. 11. In the region where the discrepancy exists, the ratio is relatively flat. This shows that the shapes of the leading-order and next-to-leading-order results are very similar and this, in turn, implies that the next-to-leading-order calculation will not give a correct description of the data in that region. Note that as \( \eta_2 \) is further decreased the ratio rapidly increases. This phenomenon has been anticipated in dijet production at the Tevatron [25] and is due to a kinematic suppression of the two-body contribution relative to the three-body contribution near the edge of phase space. Note, however, that this effect occurs outside the region covered by the data in this case.

The net result is that one must include a showering correction in order to describe the data points with the most negative pseudorapidity values. On the other hand, this region is dominated by the direct component and, therefore, is not critical for the purpose of constraining the parton distributions in the photon.
IV. CONCLUSION

As new data have accumulated, the level of understanding of jet photoproduction has increased. When a cone size of $R = 0.7$ is used to define the experimentally measured jets, the agreement between theory and experiment for inclusive single-jet observables improves significantly. Furthermore, the dijet mass distribution and the dijet angular distribution are well described by a next-to-leading-logarithm calculation. The pseudorapidity dependence of the dijet cross section in the region where both jets have negative pseudorapidities is not well described, however. This has been shown to be due to the lack of parton showering in the next-to-leading-logarithm calculation by comparing to HERWIG results with and without parton showering.

The theoretical description of the data in the regions dominated by the resolved component appears to be good. There is thus reason to be optimistic that the jet photoproduction data will soon reach a point where they can be used to constrain photon parton distributions in conjunction with data from two photon collisions from $e^+e^-$ facilities. However, with the current data the experimental and theoretical uncertainties are comparable to the variations due to changing the photon parton distribution set. Furthermore, the contribution of the gluon distribution in the photon is rather small over the kinematic region currently covered. The contribution increases as the pseudorapidities of both jets are increased. If the data region can be extended, then the sensitivity to the gluon distribution can be increased. Alternatively, if the photon energy can be increased while holding the jet $E_T$ fixed, then $x_\gamma$ will be decreased, thereby increasing the gluon contribution. This would require a data sample with larger values of $y = E_\gamma/E_e$ or higher beam energies.

**Note added.** After completion of this work, we received [26] which contains additional comparisons with ZEUS data as well as new comparisons with data from H1.
REFERENCES


Figure Captions

Fig. 1 The single-jet inclusive cross section as a function of $\eta_{\text{jet}}$ integrated over $E_T^{\text{jet}} > E_T^{\text{min}}$ for $E_T^{\text{min}} = 14, 17, 21, \text{ and } 25 \text{ GeV}$ as measured by ZEUS [13] using $R = 1$ compared with our next-to-leading order result for $R = 1$ with $R_{\text{sep}} = 2R$ (top), $R_{\text{sep}} = 1.3R$ (middle), $R_{\text{sep}} = R$ (bottom).

Fig. 2 The single-jet inclusive cross section as a function of $\eta_{\text{jet}}$ integrated over $E_T^{\text{jet}} > E_T^{\text{min}}$ for $E_T^{\text{min}} = 14, 17, 21, \text{ and } 25 \text{ GeV}$ as measured by ZEUS [13] using $R = 0.7$ compared with our next-to-leading order result for $R = 0.7$ with $R_{\text{sep}} = 2R$ (top), $R_{\text{sep}} = 1.3R$ (middle), $R_{\text{sep}} = R$ (bottom).

Fig. 3 The dijet inclusive cross section as a function of $|\cos \theta^*|$ integrated over $M_{JJ} > 47 \text{ GeV}$, and $M_{JJ}$ integrated over $|\cos \theta^*| < 0.8$ as measured by ZEUS [16] compared with our next-to-leading order result.

Fig. 4 The dijet inclusive cross section as a function of $E_T$ in bins of $\eta_1$ and $\eta_2$ as measured by ZEUS [16] compared with our next-to-leading order result.

Fig. 5 The dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$ integrated over $E_T > 14 \text{ GeV}$ as measured by ZEUS [16] compared with our next-to-leading order result for various renormalization-factorization scale choices.

Fig. 6 The dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$ integrated over $E_T > 14 \text{ GeV}$ as measured by ZEUS [16] compared with our next-to-leading order result for various photon parton distribution sets.

Fig. 7 The dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$ integrated over $E_T > 14 \text{ GeV}$ for all $x_\gamma$ as measured by ZEUS [16] compared with our next-to-leading order result. The curves shown correspond to the full result (solid line), the resolved component with only quarks in the photon (medium dashed line), resolved component with only gluons in the photon (dotted line), and direct component (long dashed line).
**Fig. 8** The dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$ integrated over $E_T > 14$ GeV for $x_\gamma > 0.75$ as measured by ZEUS [16] compared with our next-to-leading order result. The curves shown correspond to the full result (solid line), the resolved component with only quarks in the photon (medium dashed line), and direct component (long dashed line).

**Fig. 9** The dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$ integrated over $E_T > 14$ GeV as measured by ZEUS [16] compared with 1.6 times the HERWIG [24] result.

**Fig. 10** The ratio of the HERWIG results to the HERWIG results at the leading order matrix element level, i.e. treating the two scattered partons as being the jets with all showering and subsequent hadronization off, for the dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$.

**Fig. 11** The ratio of the full next-to-leading order results to those obtained using only the leading order matrix elements for the dijet inclusive cross section as a function of $\eta_2$ in bins of $\eta_1$. 
FIG. 1.
FIG. 2.
FIG. 3.
FIG. 4.
FIG. 11.