Charm Photoproduction Via Fragmentation

M. Cacciari, M. Greco

Charm photoproduction via fragmentation

Matteo Cacciari[1,*], Mario Greco[2,3]

[1]DESY, Deutsches Elektronen-Synchrotron, D-22607 Hamburg, Germany (e-mail: cacciari@desy.de)
[2]Dipartimento di Fisica, Università di Roma III, Roma, Italy
[3]INFN, Laboratori Nazionali di Frascati, Frascati, Italy (e-mail: greco@lnf.infn.it)

Abstract. The next-to-leading open charm production in γp collisions is calculated within the perturbative fragmentation functions formalism, to allow resummation of $\alpha_s \log(p_T^2/m^2)$ terms. In the large p_T region ($p_T > m$) the result is consistent with the fixed order NLO calculation, small discrepancies being found for very large p_T and at the edge of phase space. The two approaches differ in the definition and the relative contribution of the direct and resolved terms, but essentially agree on their sum. The resummation is found to lead to a reduced sensitivity to the choice of the renormalization/factorization scale.

1 Introduction

The inclusive photoproduction of heavy quarks, namely the process

$$\gamma + H \to Q + X,$$

has been evaluated up to next-to-leading order (NLO) in recent years. Ellis and Nason have first evaluated the QCD corrections to the direct $\gamma H \to QX$ process [1]. Their calculation has been successively confirmed by Smith and van Neerven [2]. The NLO corrections for parton-parton scattering, needed for the resolved part, have been produced by Nason et al. [3] and also by Beenakker et al. [4].

More recently Frixione et al. have finally written computer codes [5, 6] capable of integrating the NLO formulas thereby producing phenomenologically useful cross sections. These codes will be used to perform comparisons with our results.

Within the context of NLO calculations of heavy quark photoproduction, it is also worth reminding the calculation of inelastic J/ψ photoproduction which has recently been presented by M. Krämer et al. [7].

All the above fixed-order perturbative calculations rely on the fact that the mass of the heavy quark acts as a cutoff for infrared singularities. The splitting processes which involve a heavy quark are therefore finite order by order in perturbation theory: there is no need to subtract singularities when handling a heavy quark line. This has induced the authors of [1, 3] to adopt a modified version [8] of the $\overline{MS}$ subtraction scheme in which the heavy quark effects decouple in processes involving momenta much smaller than the heavy quark mass. This also implies that the heavy quark does not contribute to the evolution of the running coupling and of the structure functions, and therefore the absence of subprocesses initiated by an intrinsic heavy flavour generated by the structure function. All the effects of the heavy quark are therefore contained in the kernel cross section only.

The NLO one particle inclusive differential distribution will however contain terms of the kind $\alpha_s \log(p_T^2/m^2)$ which, in the large p_T limit, will become large and will spoil the perturbative expansion of the cross section. This is reflected in a large sensitivity to the choice of the renormalization/factorization (r/f) scales and hence in a large uncertainty in the theoretical prediction. Similar potentially large logarithms, of the kind $\alpha_s \log(Q/m)$, Q being the photon virtuality, do appear in NLO calculations of heavy quark electroproduction total cross sections. The resummation of these logs is being considered in a series of papers [9].

In [10] the problem of large $\alpha_s \log(p_T^2/m^2)$ terms was tackled in the hadroproduction case by introducing the technique of the perturbative fragmentation functions (PFF) through which these terms were resummed to all orders and the cross section was shown to display a milder scale sensitivity.

More in detail, the PFF approach is based on the assumption that when the momenta involved are much larger than their mass, the heavy flavours behave as if they were massless. Technically, mass terms in the kernel cross sections are suppressed by powers of m/p_T. This means that one can go back to the usual $\overline{MS}$ scheme and treat all the quarks as massless, thereby subtracting also the

singularities of the heavy quark lines and introducing a fragmentation function to absorb the final state divergencies. The key point is that the massiveness of the heavy quark can be exploited to calculate in perturbative QCD the initial state conditions for the heavy quark fragmentation functions at a scale of the order of its mass [11]. Therefore the PFF approach relies on the same input parameters of the perturbative calculation, but allows a resummation of large logarithms of (p_T/m). However the neglecting of the mass terms will of course make the PFF result not accurate in the region $p_T \sim m$. In particular, the PFF approach is not able to calculate the total cross section, which is obtained from the fixed-order calculation.

2 Scales in the photoproduction process

As well known, when calculating the next-to-leading order correction to a direct photoproduction process (i.e. the $\mathcal{O}(\alpha_{em}\alpha_s^2)$ terms) one finds that singular terms associated with the electromagnetic vertex do appear. They are due to the emission of collinear light (massless) partons from the incoming photon and they are a signal for the presence of a non-perturbative region where the photon splits into quarks and gluons before interacting with the partons in the hadronic target. This problem is solved in very much the same way as the hadroproduction case: the collinear singularity is subtracted at a given scale and factored into a structure function. The photoproduction process gets therefore splitted into two pieces.

Above the aforementioned scale the photon couples with pointlike coupling to a quark which then undergoes a hard scattering. This term is called the point-like or direct contribution. Below the factorization scale the photon is regarded as a hadron, and its structure function gives the probability for it emitting a parton which subsequently participates in the hard scattering process. Hence the name of hadronic or resolved photon component.

The photon structure functions will of course depend upon the momentum scale μ_γ at which the collinear singularities of the photon leg are subtracted. Neither the pointlike nor the hadronic components are separately independent of μ_γ, because the subtracted term in the direct component is related to the definition of the photon structure function in the resolved component.

Let us now consider the generic expression for the heavy-quark production at large p_T via fragmentation due to an on-shell photon colliding with a hadron H. Expliciting the various scale dependencies in the process, we write the $\mathcal{O}(\alpha_{em}\alpha_s^2)$ cross section in the following form

$$d\sigma_{\gamma H \to QX} \sim \sum_{ik} \int dx\, dz\, F_H^i(x, \mu_F) D_k^Q(z, \mu_F)$$

$$\times d\hat{\sigma}_{\gamma i \to k}(x, z, \alpha_s(\mu_R), \mu_R, \mu_F, \mu_\gamma)$$

$$+ \sum_{ijk} \int dx_1\, dx_2\, dz\, F_\gamma^i(x_1, \mu_F', \mu_\gamma) F_H^j(x_2, \mu_F') D_k^Q(z, \mu_F')$$

$$\times d\hat{\sigma}_{ij \to k}(x_1, x_2, z, \alpha_s(\mu_R'), \mu_R', \mu_F', \mu_\gamma) \quad (1)$$

with, at the NLO here considered,

$$d\hat{\sigma}_{\gamma i \to k}(\alpha_s(\mu_R), \mu_R, \mu_F, \mu_\gamma)$$

$$= \alpha_{em}\alpha_s(\mu_R) d\hat{\sigma}_{\gamma i \to k}^{(0)} + \alpha_{em}\alpha_s^2(\mu_R) d\hat{\sigma}_{\gamma i \to k}^{(1)}(\mu_R, \mu_F, \mu_\gamma)$$

$$d\hat{\sigma}_{ij \to k}(\alpha_s(\mu_R), \mu_R, \mu_F, \mu_\gamma)$$

$$= \alpha_s^2(\mu_R) d\hat{\sigma}_{ij \to k}^{(0)} + \alpha_s^3(\mu_R) d\hat{\sigma}_{ij \to k}^{(1)}(\mu_R, \mu_F, \mu_\gamma). \quad (2)$$

Here μ_R and μ_R' are renormalization scales, μ_F and μ_F' are factorization scales for collinear singularities arising from strong interactions, and μ_γ is a factorization scale for collinear singularities arising from the electromagnetic vertex. All the structure and fragmentation functions obey the usual Altarelli-Parisi evolution equations

$$\frac{dF^i(x, \mu_F)}{d\log\mu_F^2} = \frac{\alpha_s(\mu_F)}{2\pi} \int_x^1 \frac{dy}{y} P_{ki}(x) F^k(y/x, \mu_F) \quad (3)$$

with the exception of the photon structure function F_γ^i which also has an inhomogeneous evolution in μ_γ:

$$\frac{\partial F_\gamma^i(x, \mu_F, \mu_\gamma)}{\partial \log \mu_\gamma^2} = \frac{\alpha_{em}}{2\pi} e_i^2 [x^2 + (1 - x)^2] + \mathcal{O}(\alpha_s). \quad (4)$$

In the commonly-used photon density parametrizations, μ_γ is usually kept equal to μ_F, so that the term given in (4) becomes a correction to the usual Altarelli-Parisi equation (the so-called inhomogeneous term).

The cross section (1) is of course independent of the renormalization/factorization scales at the perturbative order at which it is calculated. Indeed μ_R cancels between $\alpha_s(\mu_R)$ and the explicit dependence of the $\sigma^{(1)}$ kernels; μ_F cancels between the structure/fragmentation functions and again the NLO kernel dependencies; finally, μ_γ cancels between the pointlike component and the photon structure function evolution in the resolved component. This last cancellation, in particular, prompts for the need to consider both the pointlike and the resolved component when evaluating photoproduction cross section. Being the factorization procedure entirely arbitrary, none of the two components has any physical meaning, only their sum being observable.[1]

3 The PFF approach

We will consider now in detail the photoproduction of heavy quarks in the framework of PFF. Then, following for instance ref. [13], we can rewrite the multidifferential cross sections for the photoproduction process

$$\gamma(P_1) + H(P_2) \to Q(P) + X$$

as follows. The resolved part reads

$$\frac{Ed^3\sigma^{res}}{d^3P} = \frac{1}{\pi S} \sum_{ijk} \int_{1-V+VW}^1 \frac{dz}{z^2} \int_{VW/z}^{1-(1-V)/z} \frac{dv}{1-v} \int_{VW/zv}^1 \frac{dw}{w}$$

$$\times F_H^i(x_1, \mu_F) F_H^j(x_2, \mu_F) D_k^Q(z, \mu_F)$$

---

[1]For a detailed discussion on the cancellation of the μ_γ dependence in the photoproduction of jets, see [12]

$$\times \left[\frac{1}{v} \left(\frac{d\sigma^0(s, v)}{dv} \right)_{ij \to k} \delta(1 - w) \right.$$

$$\left. + \frac{\alpha_s^3(\mu_R)}{2\pi} K_{ij \to k}(s, v, w; \mu_R, \mu_F) \right] \tag{5}$$

having defined the hadron-level quantities

$$V = 1 + \frac{T}{S} \quad W = \frac{-U}{S + T} \tag{6}$$

with $S = (P_1 + P_2)^2$, $T = (P_1 - P)^2$ and $U = (P_2 - P)^2$. Similarly are defined the parton-level ones s, v e w. In terms of the momentum fractions x_1, x_2 and z it holds

$$s = x_1 x_2 S \quad x_1 = \frac{VW}{zvw} \quad x_2 = \frac{1 - V}{z(1 - v)}. \tag{7}$$

Notice that we have abandoned any distinction between the photon and the hadron factorization scales μ_γ and μ_F. In the following we'll also usually identify the renormalization and the factorization scales, i.e. $\mu_R = \mu_F = \mu$.

The expression for the pointlike contribution amounts to rewriting (5) with the constraint $F_\gamma^i(x, \mu_F) = \delta(1 - x)$ with kernel cross sections for γ scattering. By exchanging the integrations over v and w we can write

$$\frac{Ed^3\sigma^{\text{point}}}{d^3P} = \frac{1}{\pi S} \sum_{ik} \int_{1 - V + VW}^1 \frac{dz}{z^2} \int_{VW/(z - 1 + V)}^1 \frac{dw}{w}$$

$$\times \int \frac{dv}{1 - v} F_H^i(x, \mu_F) D_k^Q(z, \mu_F)$$

$$\times \left[\frac{1}{v} \left(\frac{d\sigma^0(s, v)}{dv} \right)_{\gamma i \to k} \frac{VW}{zw} \delta \left(v - \frac{VW}{zw} \right) \delta(1 - w) \right.$$

$$\left. + \frac{\alpha_s^2(\mu_R)}{2\pi} K_{\gamma i \to k}(s, v, w; \mu_R, \mu_F) \frac{VW}{zw} \delta \left(v - \frac{VW}{zw} \right) \right] \tag{8}$$

with

$$s = xS \quad x = \frac{1 - V}{z - VW/w}, \tag{9}$$

and we have dropped the integration limits for v, being this integration constrained by the δ-function.

We recall that the PFF approach to large-p_T heavy quark production assumes that the heavy quark is produced via fragmentation of partons (charm included) which have been produced in a massless way in the hard scattering. The fragmentation functions D_k^Q can be evaluated in perturbative QCD at a scale of the order of the heavy quark mass and subsequently evolved up to the desired factorization scale μ_F through the Altarelli-Parisi equations.

The lowest order massless kernel cross sections $d\sigma_{ij \to k}^0$ and $d\sigma_{\gamma i \to k}^0$ (of order α_s^2 and $\alpha_{em}\alpha_s$ respectively) have been long known in literature. The next-to-leading terms for massless partons scattering $K_{ij \to k}$ (of order α_s^3) and $K_{\gamma i \to k}$ (of order $\alpha_{em}\alpha_s^2$) have been instead produced in recent years. Indeed Aversa et al. [14] have given the

explicit expressions for the NLO massless parton-parton scattering processes, $K_{ij \to k}$, while Aurenche et al. [15] have instead calculated the NLO corrections to the γ-massless parton scattering process, $K_{\gamma i \to k}$.

For the numerical evaluation of (5) and (8) we have used the computer programs originally developed by those authors [14, 15] to perform the convolution of the scattering kernels with the structure and fragmentation functions. Some modifications had however to be implemented to ensure the numerical convergence of the integrals. The codes had actually been designed to handle light hadrons fragmentation functions. These functions typically tend rapidly to zero for increasing values of the argument z, and can therefore easily be numerically integrated. This is no more true when treating heavy quarks due to the singular behaviour of the D_Q^Q fragmentation function in $z = 1$. In [10] a so called "pole subtraction" procedure was implemented. In this case we have instead modified the fragmentation function by convoluting it with a "regulator function" $D_{rf}(x; \alpha, \beta)$ given by

$$D_{rf}(x; \alpha, \beta) = A(1 - x)^\alpha x^\beta \tag{10}$$

α and β have to be chosen such that $D_{rf}(x)$ is peaked very near $x = 1$ (we have taken $\alpha = 1$ and $\beta = 1000$), and A is the normalization factor:

$$A^{-1} = \int_0^1 dx \, D_{rf}(x; \alpha, \beta) = B(\beta, \alpha + 1) \tag{11}$$

B being the Euler beta-function. The $D_{rf}(x; \alpha, \beta)$ goes to zero in $x = 1$ fast enough to allow the numerical integration. We have explicitly checked that the numerical result is compatible - within errors - with the "exact" one given by the pole subtraction method.

A further modification of the codes concerns the effect of the heavy quark mass m on the renormalization/factorization scales and on the kinematical boundaries of the integrals. Indeed the heavy quark mass has been taken explicitly into account by using

$$\mu_{\text{ref}} = \sqrt{p_T^2 + m^2} \tag{12}$$

as the central choice for the f/r scale and also by calculating the S, T and U invariant with massive kinematics. Given the heavy quark transverse momentum p_T and its scattering angle in the center of mass frame θ this amounts to write for the heavy quark quadrimomentum P

$$P = \left(\sqrt{\frac{p_T^2}{\sin^2\theta} + m^2}; p_T, 0, \frac{p_T \cos\theta}{\sin\theta} \right). \tag{13}$$

These simple kinematical considerations have of course some clear effect in the low-p_T region, making the numerical integration easier, but in no way allow one to reproduce the result of the full massive calculation, being the $\mathcal{O}(m/p_T)$ terms in the kernel cross sections still missing.

4 Numerical results and conclusions

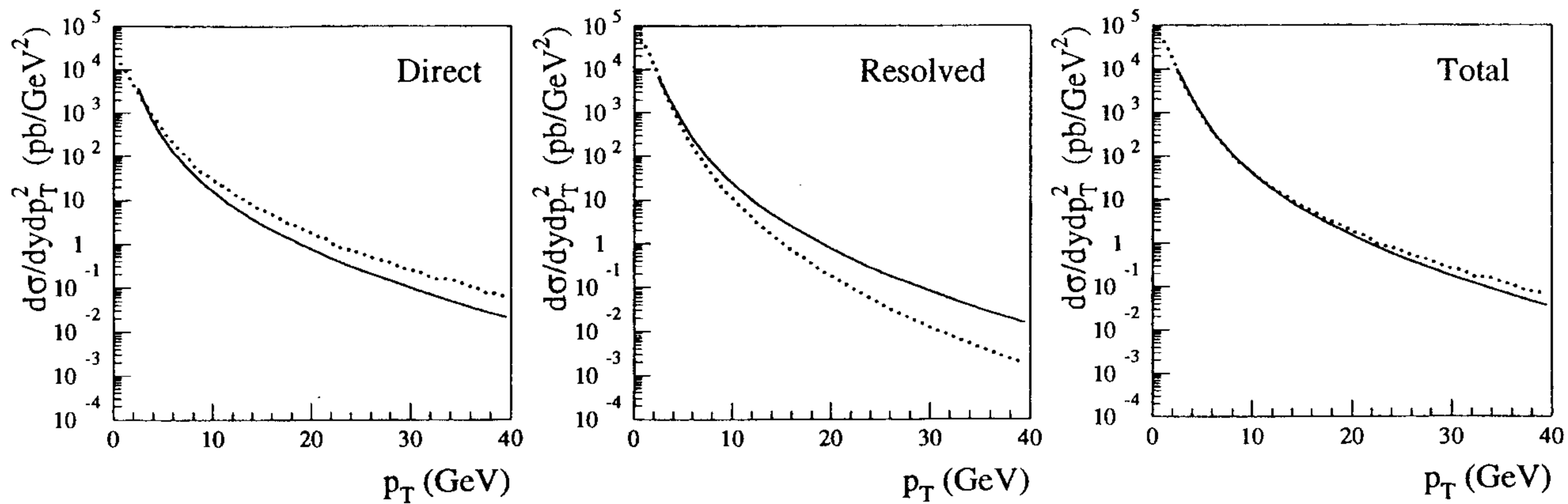Let us consider the inclusive photoproduction of a c quark at present HERA energy.

Fig. 1. Comparison between the fixed order (FMNR, dotted line) and the perturbative fragmentation function (PFF, full line) prediction of the charm photoproduction p_T spectrum. ACFGP-mc $\otimes$MRSA, $\sqrt{s} = 296$ GeV (26.7$\otimes$820), $y_{lab} = 1$, $\xi = 1$

For the sake of simplicity we'll take a photon of fixed energy $E_\gamma = 26.7$ GeV which scatters against a $E_p = 820$ GeV proton. This amounts to a center of mass energy of 296 GeV. According to the usual conventions, we'll take positive rapidities in the direction of the incoming proton. Then the rapidities in the laboratory frame and in the center of mass frame are connected by the usual formula

$$y_{lab} = y_{cm} + \frac{1}{2}\log\frac{E_p}{E_\gamma}. \tag{14}$$

We use the MRSA [16] structure function set for the proton and the ACFGP-mc [17] one for the photon. The $\Lambda_5^{\overline{MS}}$ is fixed at the central value in the MRSA set, namely 151 MeV and the mass of the charm quark is assumed 1.5 GeV.

Then Fig. 1 shows how the fixed order calculation (from here on referred to as FMNR) and the PFF one do distribute differently the overall cross section between the direct and the resolved contributions. The bidifferential cross section $d\sigma/dp_T^2 dy$ is considered at the rapidity[2] value $y_{lab} = 1$. In these plots the r/f scale is taken equal to the reference value defined in (12). Namely, $\mu = \xi\mu_{ref}$ with $\xi = 1$.

By looking at the plots we immediately appreciate that the direct and resolved contributions separately do behave differently in the two approaches. The FMNR direct part tends to be higher than the PFF one. The opposite happens in the resolved part, the PFF result being markedly higher, especially in the large p_T region. The total cross section however, where the two components have been added together, displays a good similarity between the FMNR and the PFF results. The two curves are practically coincident up to p_T values around 10–15 GeV, and above these values the PFF one is only slightly lower. This deviation is due to the multiple gluon emission from the final state charm, which becomes relevant for $p_T \gg m$.

This effect is of course not described by the fixed order NLO calculation, which only deals with one gluon emission and therefore predicts an overall harder charm fragmentation function.

The different descriptions that FMNR and PFF give of the direct and resolved part separately are due to the different treatment of the diagrams in Fig. 3, where the photon splits into a $c\bar{c}$ pair. In the FMNR calculation the diagram a) is entirely included in the NLO direct part, the finite mass of the charm quark preventing the splitting from being singular. Consistently the diagram b), i.e a scattering initiated by the heavy quark, does not appear in the resolved part (we would have a double counting otherwise). In the PFF approach on the other hand the charm is massless, and the process of photon splitting has to be separated into two. Above the factorization scale μ_γ the contribution is assigned to the direct component, while below this scale the splitting process is instead considered non perturbative: it is taken into account via the photon structure function F_γ^c (Fig. 3b) and assigned therefore to the resolved part. This different separation explains why in the PFF approach the direct component is smaller and the resolved one bigger than in the FMNR one.

The direct and resolved components of the cross section are shown in Fig. 2, where the rapidity distribution at $p_T = 10$ and 20 GeV is plotted. Once again, the PFF direct term falls below the FMNR one, due to the subtraction of part of the photon splitting to $c\bar{c}$. By considering the resolved part, on the other hand, we see that FMNR displays the typical symmetrical bell shape usually observed in hadroproduction. The PFF shows instead a peak in the low rapidity region, similarly to what happens in the direct component. Still, after adding the two contributions the two approaches show a remarkable agreement, except at the very border of phase space where PFF generally produces a lower cross section. This discrepancy is due to the fragmentation functions being here probed exclusively at values of z near one. In this region they display an unphysical singularity, and resummation of Sudakov terms is needed to ensure a proper behaviour [11].

---

[2]The PFF approach actually deals with the pseudorapidity η. In the large p_T limit however, where $p_T \gg m$, the two quantities are coincident
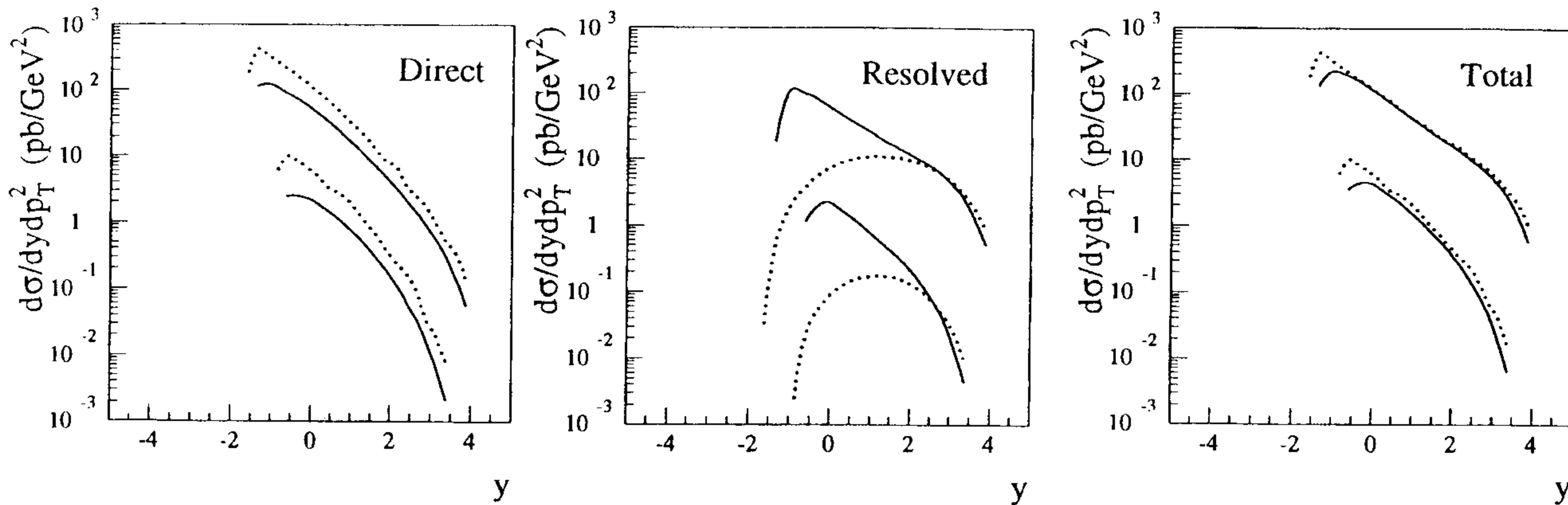
Fig. 2. Comparison between the FMNR (dotted line) and the PFF (full line) prediction for the rapidity distribution (in the lab frame) in charm photoproduction. The upper curves refer to $p_T = 10$ GeV, the lower ones to $p_T = 20$ GeV. ACFGP-mc$\otimes$MRSA, $\sqrt{s} = 296$ GeV (26.7$\otimes$820), $\xi = 1$
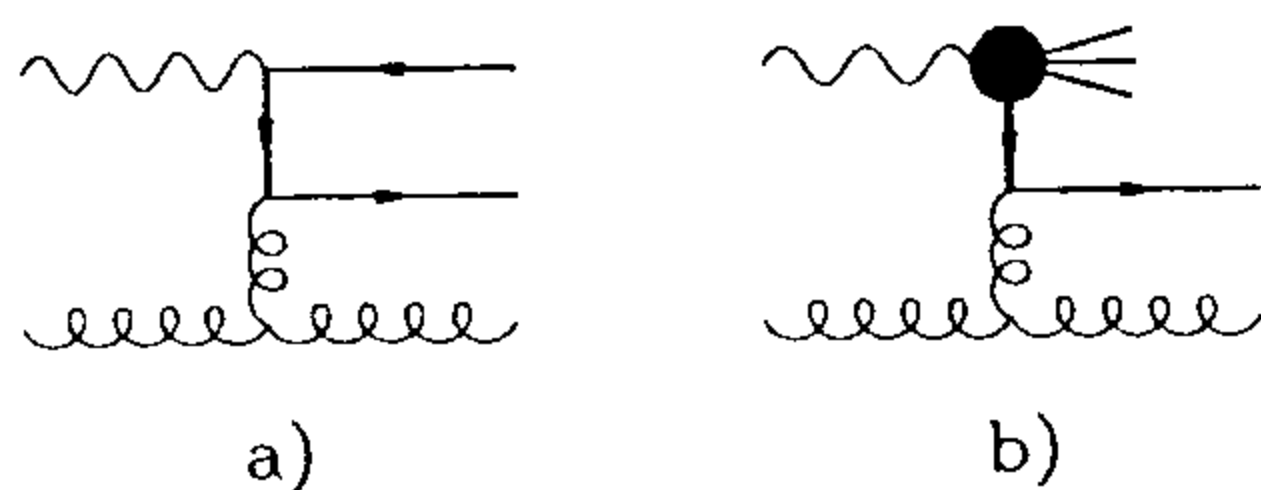

Fig. 3. The two different ways a charm can come from a photon: via pointlike coupling (a) or via the photon structure function (b)

Fig. 4. Scale sensitivity of the FMNR and PFF prediction for the direct contribution to charm photoproduction

Next we consider the scale sensitivity of the two contributions in the two approaches. To this aim we recall that in the case of hadroproduction the PFF approach leads to a reduced sensitivity to the r/f scales, improving considerably the theoretical uncertainty at large p_T. Fig. 4 shows the uncertainty band gi.en by varying the r/f scale between 0.25 μ_{ref} and 2 μ_{ref} in the direct component. In both approaches $\xi = 2$ gives the lower curve. This is coherent with the definition of "direct contribution" as what is above the factorization scale. Then raising this scale, less room is left of course for this process to take place. From the same figure we also notice a much larger sensitivity in the PFF approach. This is due to the fact that in this case a change of the factorization scale also affects the charm content of the photon. In FMNR, on the other hand, this process is not factorized and therefore its large contribution to the cross section is not affected by a variation of μ_γ.

Figure 5 shows the analogous band for the resolved component. For FMNR the $\xi = 2$ choice still gives the lower curve. This is consistent with what was observed when studing the hadroproduction case [10]. For PFF, on the contrary, $\xi = 2$ now gives the higher curve, at least in the large p_T region. This is related to the fact that the photon splitting into $c\bar{c}$ is assigned to the resolved component below the factorization scale.

Finally, Fig. 6 shows the scale sensitivity of the overall cross section. Due to the opposite behaviours of the two partial contributions previously considered the two approaches can be seen to display similar sensitivity.
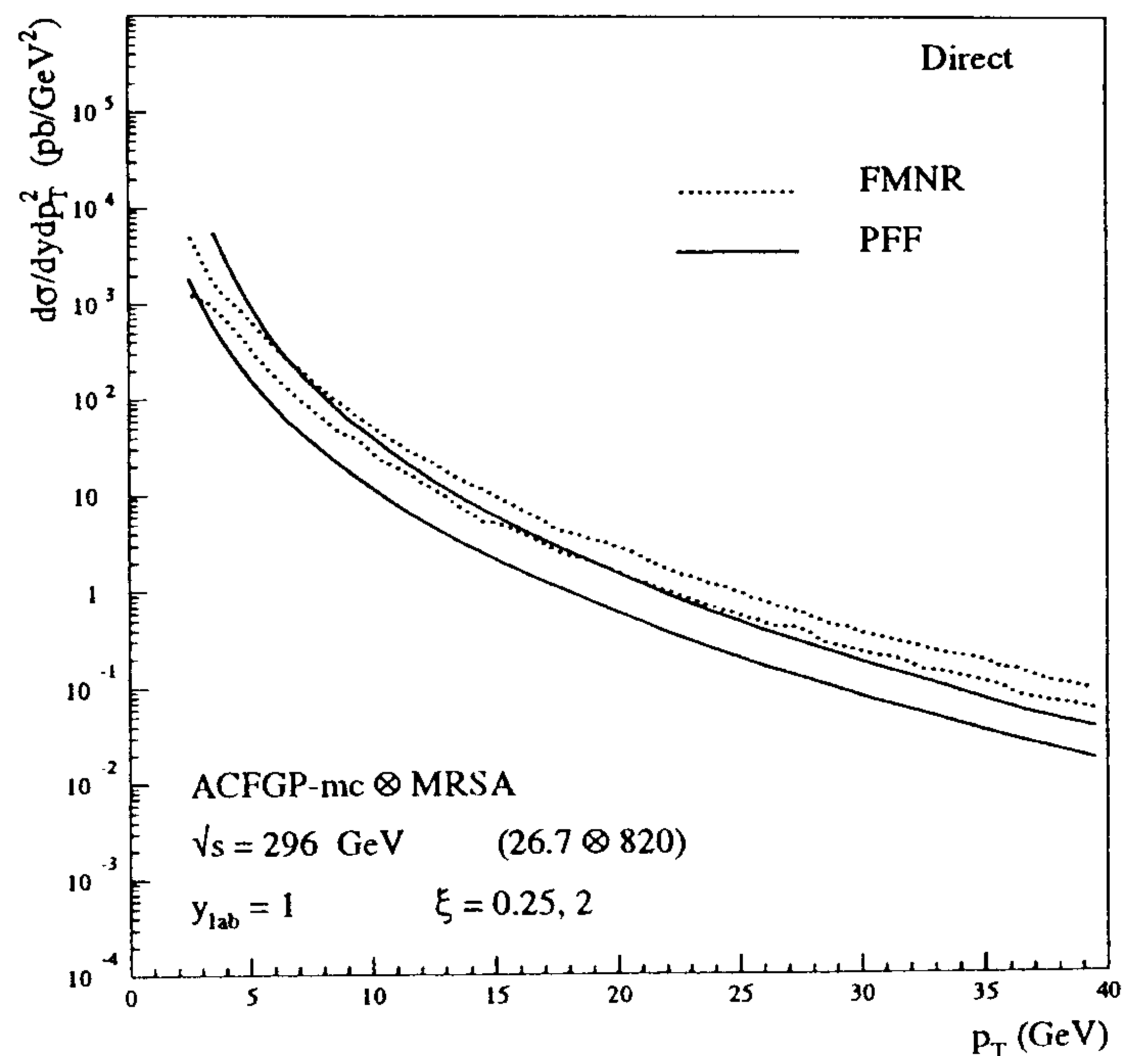
A more detailed analysis of the scale dependence is shown in Fig. 7, where the cross section is plotted for three different p_T values (10.5, 20.5 and 30.5 GeV) as a function of the r/f scale in the range $\xi = 0.25$–2. As previously noted, the dependencies are pretty similar, the PFF approach being slightly less sensitive. For comparison a PFF result obtained using leading order kernel cross sections and α_s and fragmentation functions evolutions has been plotted. It predicts a lower cross section and displays, as expected, a larger scale sensitivity.

To conclude, we have considered the open charm production in γp collisions to next-to-leading order within the perturbative fragmentation functions formalism, to allow the resummation of $\alpha_s \log(p_T^2/m^2)$ terms. Both direct
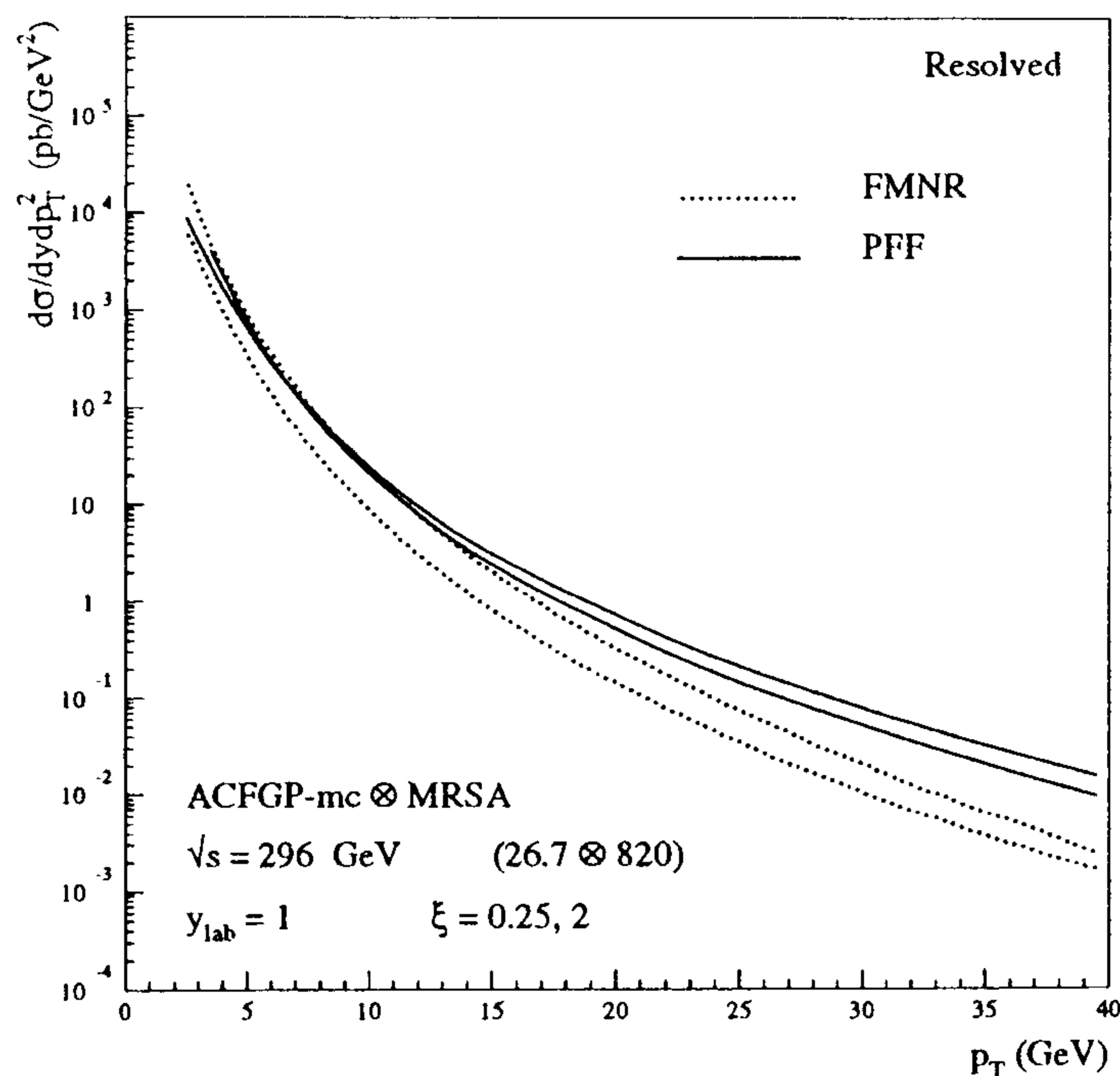
Fig. 5. Scale sensitivity of the FMNR and PFF prediction for the resolved contribution to charm photoproduction
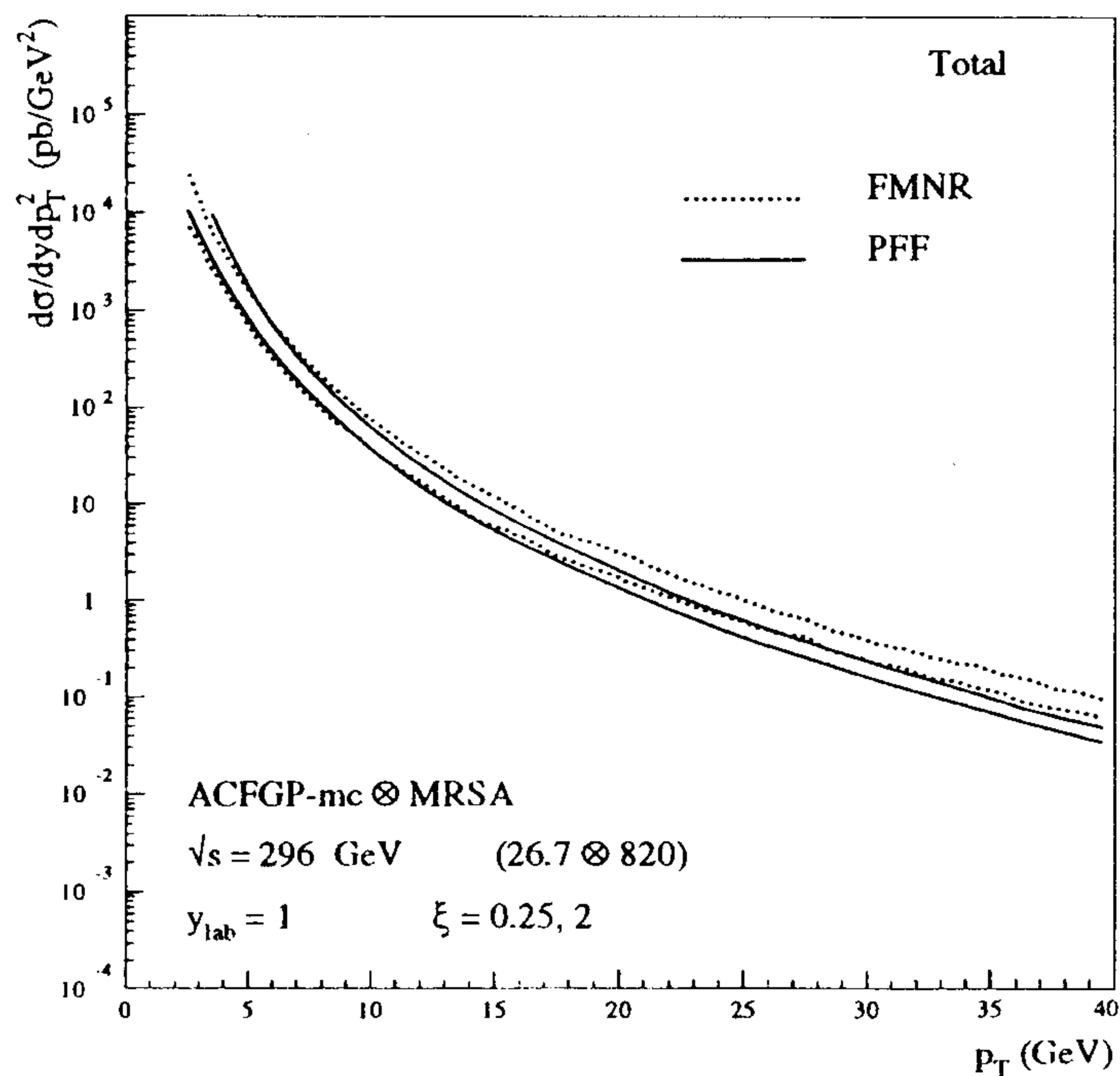

Fig. 7. Scale sensitivity of the charm photoproduction cross section as a function of the renormalization/factorization scale. "PFF LO" has been obtained with leading order cross section kernels and α_s and fragmentation functions evolutions

and the relative contribution of the direct aₓ ' resolved terms, but essentially agree on their sum. The resummation is found to lead to a slightly reduced sensitivity to the choice of the renormalization/factorization scale.

Note added: After this paper was written it came to our attention the work [18], where a similar analysis was performed.

Fig. 6. Scale sensitivity of the FMNR and PFF prediction for charm photoproduction. Both the direct and resolved component at the NLO are here taken into account

and resolved components have been considered in detail. In the large p_T region ($p_T > m$) the results are found to be consistent with the fixed order NLO calculations, small discrepancies being found for very large p_T and at the edge of phase space. The two approaches differ in the definition
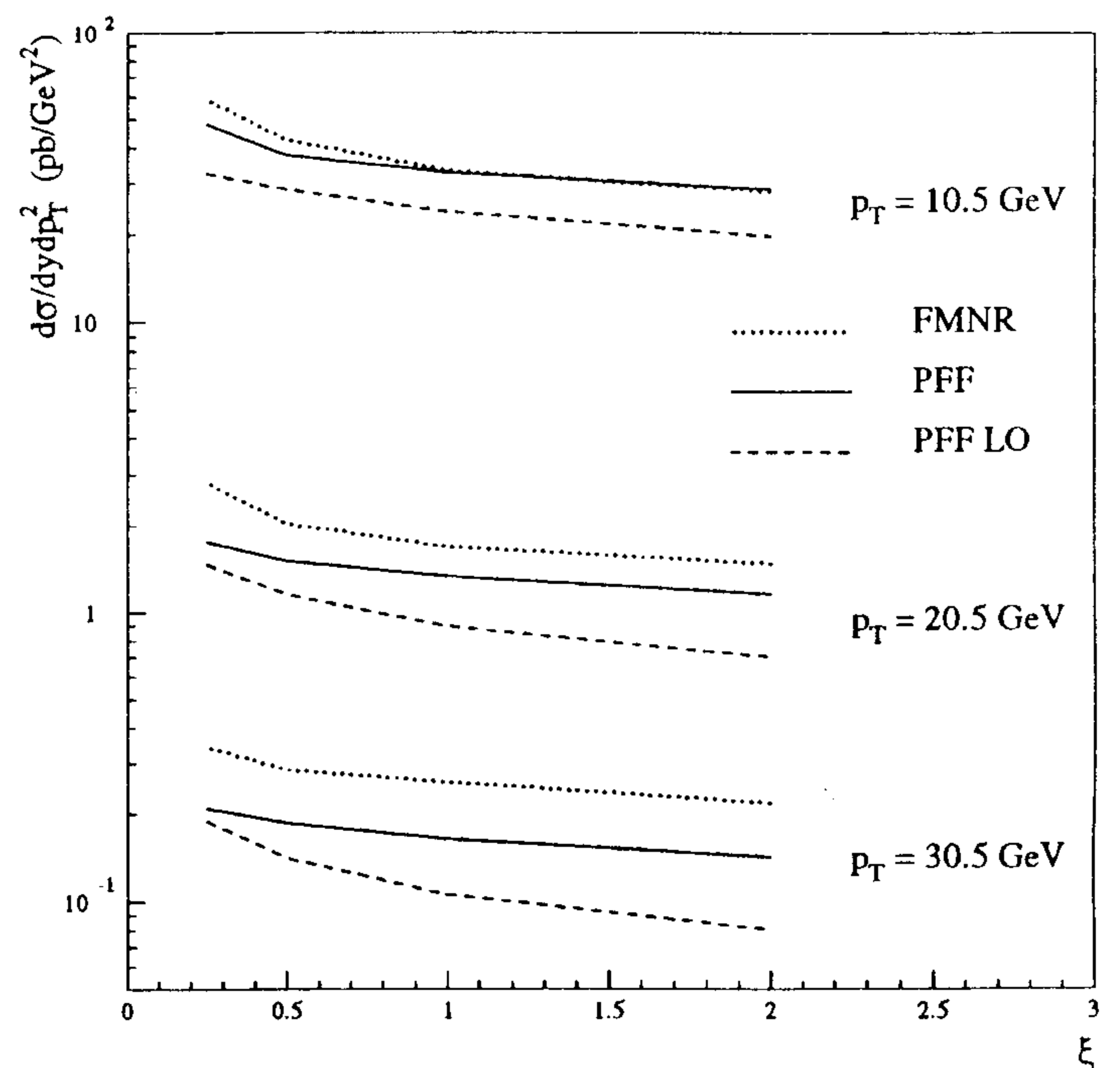
References

1. R.K. Ellis, P. Nason: Nucl. Phys. B312 (1989) 551
2. J Smith, W.L. van Neerven: Nucl. Phys. B374 (1992) 36
3. P. Nason, S. Dawson, R.K. Ellis: Nucl. Phys. B303 (1988) 607 P. Nason, S. Dawson, R.K. Ellis: Nucl. Phys. B327 (1989) 49
4. W. Beenakker, H. Kuijf, W.L. van Neerven, J. Smith: Phys. Rev. D40 (1989) 54
   W. Beenakker, W.L. van Neerven, R. Meng, G.A. Schuler, J. Smith: Nucl. Phys. B351 (1991) 507
5. M.L. Mangano, P. Nason, G. Ridolfi: Nucl. Phys. B373 (1992) 295
6. S. Frixione, M.L. Mangano, P. Nason, G. Ridolfi: Nucl. Phys. B412 (1994) 225
7. M. Krämer, J. Zunft, J. Steegborn, P.M. Zerwas: Phys. Lett. B348 (1995) 657
8. J. Collins, F. Wilczek, A. Zee: Phys. Rev. D18 (1978) 242
   W.J. Marciano: Phys. Rev. D29 (1984) 580
9. M.A.G. Aivazis, F.I. Olness, Wu-Ki Tung: Phys. Rev. D50 (1994) 3085
   M.A.G. Aivazis, J.C. Collins, F.I. Olness, Wu-Ki Tung: Phys. Rev. D50 (1994) 3102
   F.I. Olness, S.T. Riemersma: Phys. Rev. D51 (1995) 4746

10. M. Cacciari, M. Greco: Nucl. Phys. B421 (1994) 530
11. B. Mele, P. Nason: Nucl. Phys. B361 (1991) 626
12. M. Greco, A. Vicini: Nucl. Phys. B415 (1994) 386
13. R.K. Ellis, M.A. Furman, H.E. Haber, I. Hinchliffe: Nucl. Phys. B173 (1980) 397
14. F. Aversa, P. Chiappetta, M. Greco, J.Ph. Guillet: Nucl. Phys. B327 (1989) 105
15. P. Aurenche, R. Baier, A. Douiri, M. Fontannaz, D. Schiff: Nucl. Phys. B286 (1987) 553
16. A.D. Martin, R.G. Roberts, W.J. Stirling: Phys. Rev. D50 (1994) 6734
17. P. Aurenche, P. Chiappetta, M. Fontannaz, J.Ph. Guillet, E. Pilon: Z. Phys. C56 (1992) 589
18. B.A. Kniehl, M. Krämer, G. Kramer, M. Spira: DESY 95-098