

Fracture Functions from Cut Vertices

M. Grazzini^a, L. Trentadue^a and G. Veneziano^b

^a *Dipartimento di Fisica, Università di Parma and
INFN Gruppo Collegato di Parma, 43100 Parma, Italy*

^b *Theory Division, CERN, CH-1211, Geneva, Switzerland*

Abstract

Using a generalized cut vertex expansion we introduce the concept of an extended fracture function for the description of semi-inclusive deep inelastic processes in the target fragmentation region. Extended fracture functions are shown to obey a standard homogeneous DGLAP-type equation which, upon integration over t , becomes the usual inhomogeneous evolution equation for ordinary fracture functions.

PACS 13.85.Ni

1 Introduction

Within QCD, Wilson's Operator Product Expansion (OPE) [1] finds its most straightforward application in lepton hadron Deep Inelastic Scattering (DIS). The cross section for this process is related to the time-ordered product of two currents, $T[J(x)J(0)]$, which, at short distances, is expanded as

$$T[J(x)J(0)] = \sum_k C_k(x) O_k(0). \quad (1.1)$$

The coefficient functions $C_k(x)$ embody the short-distance structure of the amplitude, while O_k are local operators whose matrix elements describe long-distance effects. For DIS the OPE leads to an expansion of the forward Compton amplitude which, by use of dispersion relations [2], gives predictions for the moments of the DIS inclusive cross section.

Unfortunately, for more general hard processes the straightforward use of the OPE technique fails. Two ways out of this impasse have been proposed. The first one [3] relies on a careful study of collinear and soft singularities order by order in perturbation theory and aims at proving that such singularities can all be lumped into some universal (i.e. hard-process independent) functions. This method has the advantage of being close to the physical parton picture for the hard process. As an alternative, an expansion in terms of *cut vertices* was proposed by Mueller [4] as a more direct extension of the OPE. Although more formal, the method of cut vertices has the advantage of giving directly an expansion for the absorptive part of the amplitude, i.e. for the cross section, in terms of the convolution of a cut vertex with a coefficient function. The cut vertex contains all the long-distance effects, and, as such, generalizes the matrix element of a local operator, while the coefficient function, as usual, is calculable in perturbation theory. The cut vertex method has been used in Ref.[5] to deal with a variety of hard processes in QCD.

Let us consider a deep inelastic semi-inclusive reaction in which a hadron in the final state is detected. When the transverse momentum of the hadron is of order Q the usual current fragmentation mechanism holds [6]. In the region in which the transverse momentum of the produced hadron is much smaller than Q , the so-called target fragmentation

region, the simple description based on parton densities and fragmentation functions fails. For this reason, in order to describe the target fragmentation region as well, a new formulation of semi-inclusive hard processes in terms of *fracture functions* was proposed [7] and later developed [8, 9]. A fracture function $M_{AA'}^i(x, z, Q^2)$ gives the probability of extracting a parton i with momentum fraction x from a hadron A while observing a final hadron A' in the target fragmentation region with longitudinal momentum fraction z . The possible relevance of such an object was already mentioned by Feynman [10] before the advent of QCD.

In this paper we will consider the case in which also the transverse momentum of the hadron A' (or equivalently the momentum transfer $t = -(p_A - p_{A'})^2$) is observed. We will argue that, in the region $t \ll Q^2$, an expansion in terms of cut vertices can be given and that a t dependent (or extended) fracture function¹ can be defined in terms of a new kind of cut vertex. We will show, within a $(\phi^3)_6$ toy-model field theory, that extended factorization is just a consequence of collinear power counting [12, 13].

The paper is organized as follows. In Sect. 2 we recall the definition of cut vertices and we obtain the expansion of the moments of the structure function in terms of matrix elements of local operators times coefficient functions by using cut vertices. In Sect. 3 the cut vertex approach is extended to deal with semi-inclusive processes and a generalized cut vertex expansion is given. In Sect. 4 we define extended fracture functions and their evolution equation is given and discussed. In Sect. 5 we discuss the results and draw our conclusions.

2 Cut vertices

In this section we will recall the definition of cut vertices in the simpler case of $(\phi^3)_6$ theory. This toy model, despite its simpler structure, shares several important properties with QCD. It has a dimensionless, asymptotically free coupling constant and the diagrams with leading mass singularities have the same topology as in (light cone gauge) QCD. For

¹Similar objects have been defined in the context of diffraction in Ref. [11].

these reasons $(\phi^3)_6$ is an excellent theoretical framework for the study of factorization properties [14].

Consider the inclusive deep inelastic process

$$p + J(q) \rightarrow X$$

off the current $J = \frac{1}{2}\phi^2$. We define as usual Q^2 and x as

$$Q^2 = -q^2 \quad x = \frac{Q^2}{2pq}. \quad (2.1)$$

Let us choose a frame in which $p = (p_+, p_-, \mathbf{0})$ with $p_+ \gg p_-$ and $pq \simeq p_+q_-$. Given a vector $k = (k_+, k_-, \mathbf{k})$ define $\hat{k} = (k_+, 0, \mathbf{0})$. The structure function, defined as

$$F(p, q) = \frac{Q^2}{2\pi} \int d^6y e^{iqy} \langle p | J(y) J(0) | p \rangle, \quad (2.2)$$

describes the interaction of the far off-shell current $J(q)$ with an elementary quantum of momentum p through the discontinuity of the forward scattering amplitude (see Fig. 1).

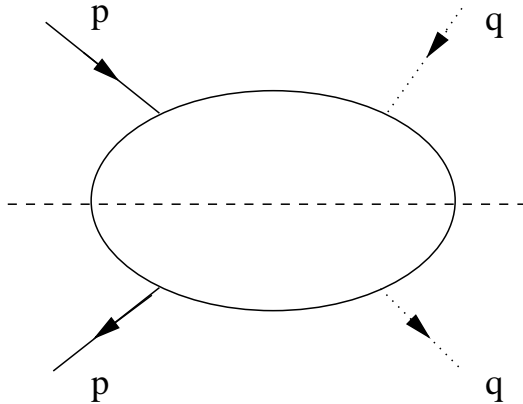


Figure 1: *Deep inelastic structure function in $(\phi^3)_6$*

The leading contribution to the structure function comes from the decomposition shown in Fig.2. Here, with the notations of Ref. [4], τ is the hard part of the diagram, i.e. the one in which the large momentum flows, while λ is the soft part. Decompositions with more than two legs connecting the hard to the soft part are suppressed by powers of $1/Q^2$.

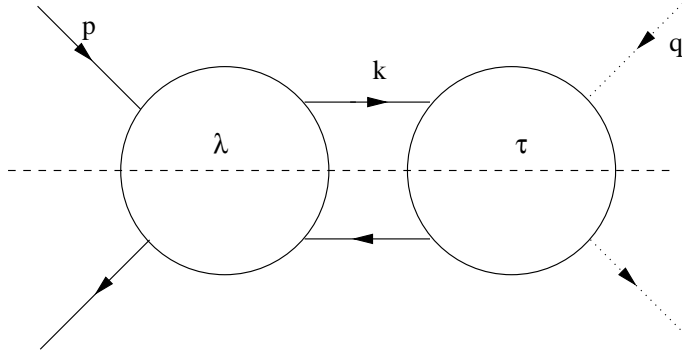


Figure 2: *Relevant decomposition for the deep inelastic structure function in $(\phi^3)_6$*

Such decomposition can be written in formulae as

$$F(p, q) = \sum_{\tau} \int V_{\lambda}(p, k) H_{\tau}(\hat{k}, q) \frac{d^6 k}{(2\pi)^6}, \quad (2.3)$$

where $V_{\lambda}(p, k)$ and $H_{\tau}(k, q)$ are the discontinuities of the long and short distance parts, respectively.

Moreover, in order to pick up the leading contribution in eq. (2.3), the momentum k which enters τ is taken to be collinear to the external momentum p . Neglecting renormalization, let us define for a given decomposition into a λ and a τ subdiagram

$$v_{\lambda}(p^2, x) = \int V_{\lambda}(p, k) x \delta\left(x - \frac{k_+}{p_+}\right) \frac{d^6 k}{(2\pi)^6} \quad (2.4)$$

and

$$C_{\tau}(x, Q^2) = H_{\tau}(k^2 = 0, x, q^2). \quad (2.5)$$

Here $v_{\lambda}(p^2, x)$ represents the contribution of λ when the hard part is contracted to a point, while $C_{\tau}(x, Q^2)$ is the hard part in which one neglects the virtuality of the incoming momentum with respect to Q^2 . Since

$$x \simeq \frac{Q^2}{2p_+q_-} \quad (2.6)$$

and

$$H_{\tau}(\hat{k}, q) = H_{\tau}\left(0, \frac{Q^2}{2k_+q_-}, q^2\right), \quad (2.7)$$

using the definition of \hat{k} and eqs. (2.3)-(2.7) we can write

$$\begin{aligned}
F(p, q) &= \sum_{\tau} \int V_{\lambda}(p, k) H_{\tau}(\hat{k}, q) \frac{d^6 k}{(2\pi)^6} \\
&= \sum_{\tau} \int V_{\lambda}(p, k) \delta\left(u - \frac{k_{+}}{p_{+}}\right) du C_{\tau}(x/u, Q^2) \frac{d^6 k}{(2\pi)^6} \\
&= \sum_{\tau} \int v_{\lambda}(p^2, u) C_{\tau}(x/u, Q^2) \frac{du}{u} \equiv \int v(p^2, u) C(x/u, Q^2) \frac{du}{u}.
\end{aligned} \tag{2.8}$$

The last integral defines the spacelike *cut vertex* $v(p^2, x)$ and the corresponding coefficient function $C(x, Q^2)$. As usual a simpler factorized expression for the structure function is obtained by taking moments with respect to x . Defining the Mellin transform as

$$f_{\sigma} = \int_0^1 dx x^{\sigma-1} f(x), \tag{2.9}$$

we find immediately

$$F_{\sigma}(p^2, Q^2) = v_{\sigma}(p^2) C_{\sigma}(Q^2). \tag{2.10}$$

It was shown in Refs.[15, 16] that the cut vertex represents the analytic continuation in the spin variable of a matrix element of operators of minimal twist. This correspondence has been confirmed up to two loops by direct calculation of the anomalous dimensions of cut vertices and leading twist operators [15]. Hence, in the case of DIS, the factorized expression (2.10) can be identified with the one given by OPE

$$F_n(p^2, Q^2) = A_n(p^2) C_n(Q^2), \tag{2.11}$$

where $A_n(p^2)$ are now matrix elements of local operators. Thus, for integer values of σ , the coefficient function which appears in (2.10) is the same as in (2.11). This fact will be used in the next section where the evolution of the extended fracture function will be shown to be driven by the anomalous dimension of the same set of local operators.

3 A cut vertex approach to semi-inclusive processes

Let us consider now, still within $(\phi^3)_6$, a deep inelastic reaction in which a particle with momentum p' is inclusively observed in the final state, i.e. the process

$$p + J(q) \rightarrow p' + X.$$

By using the same line of reasoning as for the inclusive case we may define a semi-inclusive structure function as (see Fig. 3)

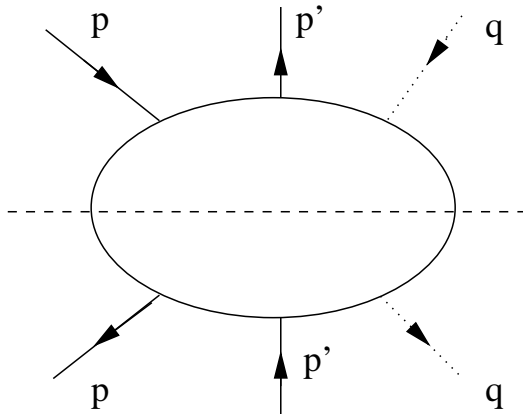


Figure 3: *Deep inelastic semi-inclusive structure function in $(\phi^3)_6$*

$$W(p, p', q) = \frac{Q^2}{2\pi} \sum_x \int d^6x e^{iqx} \langle p | J(x) | p' X \rangle \langle X p' | J(0) | p \rangle \quad (3.1)$$

in terms of matrix elements of the current operator between the incoming hadron with momentum p and the outgoing hadron with momentum p' plus anything.

When the observed particle has transverse momentum p_{\perp}^2 of order Q^2 the cross section is dominated by the current fragmentation mechanism and can be written in the usual factorized way [6]

$$W(p, p', q) = \int \frac{dx'}{x'} \frac{dz'}{z'} f_A(x', Q^2) \hat{\sigma}(x/x', z', Q^2) D_{A'}(z/z', Q^2), \quad (3.2)$$

where

$$z = \frac{pp'}{pq} \simeq \frac{p'_-}{q_-}. \quad (3.3)$$

In the language of cut vertices, eq.(3.2) is a convolution of a spacelike and a timelike cut vertex through a coefficient function [5]

$$W(p, p', q) = \int \frac{dx'}{x'} \frac{dz'}{z'} v(p^2, x') C(x/x', z', Q^2) v'(p'^2, z/z'). \quad (3.4)$$

By contrast, the limit $t = -(p - p')^2 \ll Q^2$ is dominated by the target fragmentation mechanism and has not been considered in either approach². Following the same steps as in the previous section, one can argue that, in the region $t \ll Q^2$, the leading contribution to the semi-inclusive cross section is given by the decomposition shown in Fig.4.

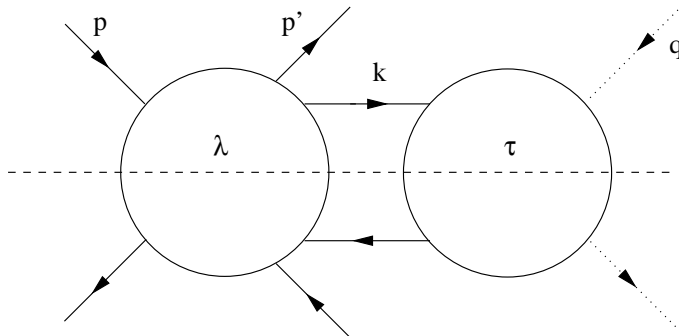


Figure 4: *Relevant decomposition for the semi-inclusive structure function in $(\phi^3)_6$ in the limit $t \ll Q^2$*

Such a decomposition implies that an expansion similar to (2.8) holds, in terms of a new function $v(p, p', \bar{x})$ and a coefficient function $C(\bar{x}, Q^2)$

$$W(p, p', q) = \int v(p, p', u) C(\bar{x}/u, Q^2) \frac{du}{u} \quad (3.5)$$

where we have defined a new variable z as

$$z = \frac{p'q}{pq} \simeq \frac{p'_+}{p_+} \quad (3.6)$$

and a rescaled variable $\bar{x} = x/(1 - z)$. The new function $v(p, p', \bar{x})$ is given by

$$v(p, p', \bar{x}) = \int T(p, p', k) \bar{x} \delta\left(\bar{x} - \frac{k_+}{p_+ - p'_+}\right) \frac{d^6k}{(2\pi)^6}, \quad (3.7)$$

where $T(p, p', k)$ is the discontinuity of a six-point amplitude in the channel $(p - p' - k)^2$. The function $v(p, p', \bar{x})$ is a new object that we will call a *generalized cut vertex*, which depends both on p and p' and embodies all the leading mass singularities of the cross

²In Ref. [8] it has been shown at one loop that in the limit $t \rightarrow 0$ a new collinear singularity appears in the semi-inclusive cross section which cannot be absorbed into parton densities and fragmentation functions, and so must be lumped into a new phenomenological distribution, i.e. the fracture function.

section. By taking moments with respect to \bar{x} as in eq.(2.9), eq. (3.5) becomes

$$W_\sigma(p, p', q) = v_\sigma(p, p') C_\sigma(Q^2) \quad (3.8)$$

that is a completely factorized expression analogous to (2.10).

We are now going to show that this expansion holds up to corrections suppressed by powers of $1/Q^2$. In order to do so we will use the method of infra-red power counting [12, 13] applied to our process.

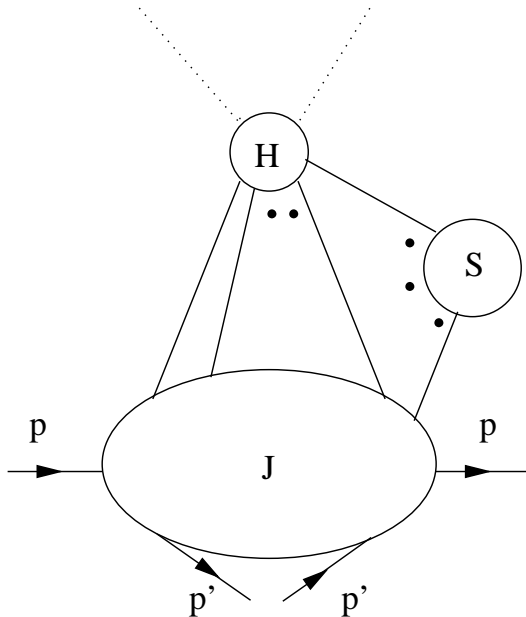


Figure 5: *General reduced graphs which contribute to the semi-inclusive structure function in $(\phi^3)_6$*

In order to get insight into the large Q^2 limit of the semi-inclusive cross section, let us look at the singularities in the limit $p^2, p'^2, t \rightarrow 0$. The infra-red power counting technique can predict the strength of such singularities. Starting from a given diagram, its *reduced* form in the large Q^2 limit is constructed by simply contracting to a point all the lines whose momenta are not on shell. The general reduced diagrams in the large Q^2 limit for the process under study involve a jet subdiagram J , composed by on-shell lines collinear to the incoming particle, from which the detected particle emerges in the forward direction³, a hard subgraph H in which momenta of order Q circulate, which is

³In the large Q limit p and p' can be taken as parallel.

connected to the jet by an arbitrary number of collinear lines. Soft connections between J and H can be possibly collected into a soft blob S which is connected to the rest of the diagram by an arbitrary number of lines (see Fig.5). In $(\phi^3)_6$, by using power counting [13], we find⁴ that the leading contributions come from graphs with no soft lines and the minimum number of collinear lines connecting the hard to the jet subdiagram, as in Fig.6.

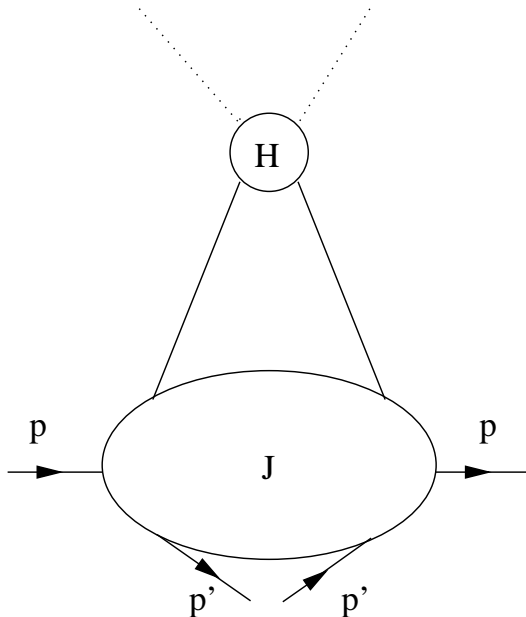


Figure 6: *Leading contributions to the semi-inclusive structure function in $(\phi^3)_6$*

Any other diagram containing additional collinear lines between J and H is suppressed by powers of $1/Q^2$. It follows that $W(p, p', q)$ is of the following form

$$W(p, p', q) = \int \frac{d^6 k}{(2\pi)^6} T(p, p', k) H(\hat{k}, q) + \mathcal{O}(1/Q^2). \quad (3.9)$$

It is now straightforward to show that eq. (3.9) is equivalent to (3.5) with the substitution $H(0, x, q^2) = C(x, Q^2)$. Thus the expansion (3.5) corresponds to taking the leading part of the semi-inclusive cross section.

⁴This fact has been verified by an explicit one-loop calculation in Ref. [17].

4 Extended fracture functions

In the previous section we have given arguments for the validity of a generalized cut vertex expansion for the process $p + J(q) \rightarrow p' + X$ in the region $t \ll Q^2$. We now want to investigate the consequences of such a result.

The coefficient function which appears in (3.5) is the same as that of (2.8) since it comes from the hard part of the graphs which is exactly the same as in DIS. So we can draw the important conclusion that the evolution of the coefficient function appearing in (3.5) is directly related to the anomalous dimension of the leading twist local operator which drives the evolution of the DIS coefficient function.

Despite the fact that the theoretical framework in which we have been working is the model field theory $(\phi^3)_6$, we expect that the main consequences expressed in (3.5) will remain valid also in a gauge theory such as QCD. The only further complications which are expected to arise are due to soft gluon lines connecting the hard to the jet subdiagrams. Unlike in $(\phi^3)_6$, in QCD these diagrams are not suppressed by power counting; hence the only way to get rid of such contributions is to show that they cancel out. This is the issue of a complete factorization proof, which is beyond the aim of this paper. As already argued in Ref.[11] we do not expect that this complication will destroy factorization. In QCD, by using renormalization group, we have

$$C_n^i(Q^2) \equiv C_n^i(Q^2/Q_0^2, \alpha_S) = \left[e^{\int_{\alpha_S}^{\alpha_S(Q^2)} d\alpha \frac{\gamma^{(n)}(\alpha)}{\beta(\alpha)}} \right]_{ij} C_n^j(1, \alpha_S(Q^2)), \quad (4.1)$$

where Q_0 is the renormalization scale, $\alpha_S \equiv \alpha_S(Q_0^2)$, $\gamma^{(n)}$ is the anomalous dimension matrix of the relevant operators, and an ordered exponential is to be understood. Thus we can write the analogue of eq.(3.8) in QCD as

$$W_n(z, t, Q^2) = \sum_i \mathcal{M}_n^i(z, t, Q^2) C_n^i(1, \alpha_S(Q^2)) \quad (4.2)$$

where, by following Ref. [15], we have defined a t -dependent fracture function

$$\mathcal{M}_n^j(z, t, Q^2) \equiv V_n^i(z, t, Q_0^2) \left[e^{\int_{\alpha_S}^{\alpha_S(Q^2)} d\alpha \frac{\gamma^{(n)}(\alpha)}{\beta(\alpha)}} \right]_{ij} \quad (4.3)$$

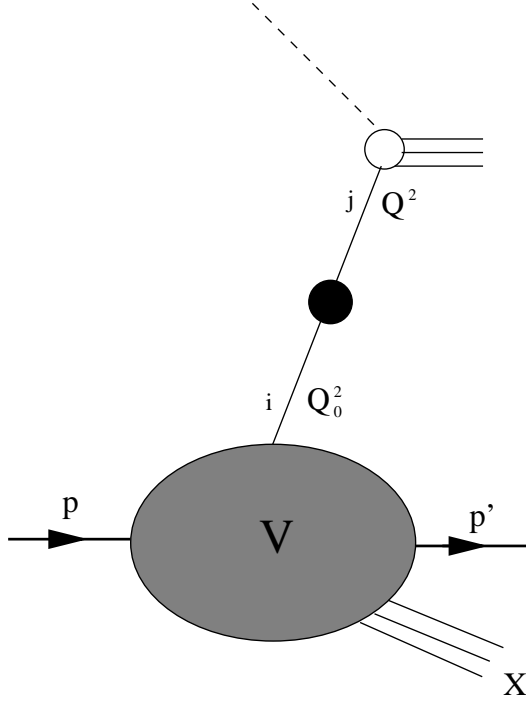


Figure 7: *Extended fracture function*

just in terms of a cut vertex $V_n^i(z, t, Q_0^2)$ (see Fig. 7).

Inverting the moments and expressing the extended fracture function in terms of the usual Bjorken variable x , one finds that $\mathcal{M}_{A,A'}^i(x, z, t, Q^2)$ obeys the simple homogeneous evolution equation

$$Q^2 \frac{\partial}{\partial Q^2} \mathcal{M}_{A,A'}^i(x, z, t, Q^2) = \sum_j \int_{\frac{x}{1-z}}^1 \frac{du}{u} K_{ij}(u, \alpha_S(Q^2)) \mathcal{M}_{A,A'}^j(x/u, z, t, Q^2) \quad (4.4)$$

where $K_{ij}(u, \alpha)$, defined as

$$K_{ij}(u, \alpha) \equiv \frac{1}{2\pi i} \int_{\frac{1}{2}-i\infty}^{\frac{1}{2}+i\infty} dn \gamma_{ij}^{(n)}(\alpha) u^{-n}, \quad (4.5)$$

is the same DGLAP kernel which controls the evolution of the ordinary parton distribution functions. This result looks particularly appealing since it means that the evolution of the extended fracture function follows the usual perturbative behaviour. One may ask at this point how this result matches with the peculiar equation which drives the evolution

of ordinary fracture functions [7]:

$$Q^2 \frac{\partial}{\partial Q^2} M_{A,A'}^j(x, z, Q^2) = \frac{\alpha_S(Q^2)}{2\pi} \int_{\frac{x}{1-z}}^1 \frac{du}{u} P_i^j(u) M_{A,A'}^i(x/u, z, Q^2) + \frac{\alpha_S(Q^2)}{2\pi} \int_x^{\frac{x}{x+z}} \frac{du}{x(1-u)} F_A^i(x/u, Q^2) \hat{P}_i^{jl}(u) D_{l,A'} \left(\frac{zu}{x(1-u)}, Q^2 \right). \quad (4.6)$$

The evolution equation for $M_{A,A'}^j(x, z, Q^2)$ contains two terms: a homogeneous term describing the non-perturbative production of a hadron coming from target fragmentation and an inhomogeneous term whose origin is the perturbative fragmentation due to initial state bremsstrahlung. As discussed in Ref. [7], the separation between perturbative and non-perturbative fragmentation introduces an arbitrary scale but the fracture function itself $M_{A,A'}^j(x, z, Q^2)$ does not depend on it.

It can be shown [18] that the evolution equation (4.6) can be derived by an explicit calculation using eq. (4.4) together with Jet Calculus rules [19]. By defining in fact the ordinary fracture function as an integral over t up to a cut-off of order Q^2 , e.g. ϵQ^2 with $\epsilon < 1$:

$$M_{A,A'}^j(x, z, Q^2) = \int^{\epsilon Q^2} dt \mathcal{M}_{A,A'}^j(x, z, t, Q^2), \quad (4.7)$$

the inhomogeneous term in the evolution equation is obtained by taking into account the Q^2 dependence of the integration cut-off.

Moreover, the interplay between the scales Q^2 and t has a sizeable effect in terms of a new class of perturbative corrections of the form $\log Q^2/t$. Such corrections are large and potentially dangerous in the region $t \ll Q^2$ since they can ruin a reliable perturbative expansion. Those terms are naturally resummed into eq. (4.3). For the extended fracture function these corrections do play an important role for understanding the dynamics of semi-inclusive processes in the kinematic region we have been considering here [18].

5 Conclusions

In this paper we have presented an extension of the cut vertex formalism to deal with semi-inclusive processes. We have shown that an extended fracture function can be defined in

terms of a new kind of cut vertex. Such an unintegrated, t -dependent fracture function obeys a DGLAP evolution equation with the same kernel as the one appearing in ordinary parton densities. This result is particularly evident in the cut vertex approach and agrees with the analysis made in a different context in Ref.[11]. As a non-trivial check, the more complicated evolution equation for ordinary fracture functions [7] can be obtained from the simpler one obeyed by extended fracture functions.

Although our results have been obtained within the $(\phi^3)_6$ model, where a treatment based on mass singularities and infra-red power counting takes a much simpler form, we believe the extension to QCD to be harmless. Indeed, the picture that emerges for semi-inclusive hadronic processes in the target fragmentation region appears to be fully consistent with the intuitive QCD-corrected parton picture.

In addition, a new class of logarithmic contributions appears at $\Lambda^2 \ll t \ll Q^2$. Extended fracture functions could give an additional tool to handle this interesting regime which certainly deserves further study [18].

We finally note that the relation between ordinary and extended fracture functions is quite analogous to the one between the experimentally defined/measured diffractive structure functions $F_2^{D(3)}(x_{\mathcal{P}}, \beta, Q^2)$ and $F_2^{D(4)}(x_{\mathcal{P}}, \beta, Q^2, t)$ [20] which depend upon observing large rapidity gaps. Relating such quantities to fracture functions is another interesting issue for future investigation.

Acknowledgements

We would like to thank J.C. Collins, J. Kodaira and G.M. Shore for helpful discussions.

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