We report a study of the processes of coherent a pion and a photon diffraction dissociation into two jets. The structure of non-factorizable contributions in these reactions is discussed. We argue that production of hard dijets by real photons can provide direct evidence for chirality violation in hard processes and the first measurement of the magnetic susceptibility of the quark condensate.

1 Introduction

The process of a pion diffraction dissociation into a pair of jets on a nucleon target was suggested as a probe of the nuclear filtering of pion components with a small transverse size. The A-dependence of the coherent dijet cross section was first calculated in [1] and it was argued that the jet distribution with respect to the longitudinal momentum fraction has to follow the quark momentum distribution in the pion and hence provides a direct measurement of the pion distribution amplitude. Recent experimental data by the E791 collaboration [2] indeed confirm the strong A-dependence which is a signature for color transparency. Moreover, the jet longitudinal momentum fraction distribution turns out to be consistent with the \( \sim z^2(1-z)^2 \) shape corresponding to the asymptotic pion distribution amplitude. After these first successes, one naturally asks whether the QCD description of coherent dijet production can be made fully quantitative. First we discuss factorization and concentrate on a pion dissociation process, then we will consider production of dijets initiated by a real photon, a process which is sensitive to the chiral-odd properties of QCD vacuum. The results reported here have been obtained in collaboration with V. Braun, S. Gottwald, A. Schäfer and L. Szymanowski [3].
2 Pion dissociation

The kinematics of the process and the notations for momenta is shown in Fig. 1. The momenta of the incoming particles are $p_1$ and $p_2$, $z$ is the longitudinal momentum fraction ($\bar{z} \equiv 1 - z$) and $q_{\perp}$ the transverse momentum of the quark jet. We consider the forward limit, when transverse momenta of the jets compensate each other. In this kinematics the invariant mass of the produced $q\bar{q}$ pair is equal to $M^2 = q_{\perp}^2/z\bar{z}$, and the momentum of the outgoing nucleon $p'_2 = p_2(1 - \xi)/(1 + \xi)$, where $\xi = M^2/(2s - M^2) \simeq M^2/2s$, $s = (p_1 + p_2)^2$.

The possibility to constrain the pion distribution amplitude $\phi_\pi(u)$ in the dijet diffractive dissociation experiment assumes that the amplitude of this process can be calculated in the collinear approximation as suggested by Fig. 1:

$$M = \int_0^1 du \int_{-1}^1 dy \phi_\pi(u) T_{H}^{g}(u, y) \mathcal{H}_{g}(y, \xi).$$

Here $\mathcal{H}_{g}(y, \xi)$ is the generalized gluon distribution in the target nucleon, variable $-1 < y < 1$ parametrizes the momentum fractions of the emitted and the absorbed gluons. $T_{H}^{g}(u, y)$ is the hard scattering amplitude involving at least one hard gluon exchange.

There are two important regions in the integral (1), see [3] for more details. At $u \to 0, 1$

$$M|_{\text{end-points}} \sim i \bar{z}z \int_{u_{\text{min}}}^1 du \frac{\phi_\pi(u)}{u^2} \mathcal{H}_{g}(\xi, \xi).$$

Since $\phi_\pi(u) \sim u$ at $u \to 0$, the integral over $u$ diverges logarithmically. Remarkably, the integral containing the pion distribution amplitude does not involve any $z$-dependence. Therefore, the longitudinal momentum distribution of the jets in the nonfactorizable contribution is calculable and, as it turns out, has the shape of the asymptotic pion distribution amplitude $\phi_\pi^{as}(z) = 6z\bar{z}$.

The appearance of the end point divergence is due to pinching of the $y$ contour at the point $y = \xi$ in case that the variable $u$ is close to the end–points. One can trace [3] that this pinching occurs between soft gluon interactions in the initial and in the final state, and is related with the existence of the unitarity cuts of the amplitude in different, $s$ and $M^2$, channels. The other important integration region in Eq. (1) is the one when the longitudinal momentum fraction carried by the quark is close (for high energies) to that of the quark jet in the final state

$$M|_{\xi \ll |u - z| \ll 1} \sim 4i\phi_\pi(z) \int_{\xi} dy \frac{\mathcal{H}_{g}(y, \xi)}{y + \xi}.$$
the high transverse momentum can be emitted in a broad rapidity interval and is not constrained to the pion fragmentation region. The integral on the r.h.s. of Eq. 3 can be identified with the unintegrated generalized gluon distribution. Therefore, in this region hard gluon exchange can be viewed as a large transverse momentum part of the gluon distribution in the proton, cf. 6. This contribution is proportional to the pion distribution amplitude \( \phi_\pi(z) \) whereas the end-point contribution 2 imitates the shape of \( \phi_\pi^s(z) \). This implies, in a contradiction to 1, that jets longitudinal momentum distribution does not proportional to \( |\phi_\pi(z)|^2 \).

3 Photon dissociation

The wave function of a real photon contains both the perturbative chiral-even (CE) contribution of the quark-antiquark pair with opposite helicities, and the nonperturbative chiral-odd (CO) contribution with quarks having the same helicity and which is due to the chiral symmetry breaking. It is proportional to fundamental parameters of QCD vacuum, quark condensate \( \langle \bar{q}q \rangle \) and magnetic susceptibility \( \chi \). The perturbative CE contribution is singular \( \sim 1/|r| \) at small transverse distances \( r \). The nonperturbative CO contribution is regular at small transverse separations and can be parametrized by the photon distribution amplitude \( \phi_\gamma(u, \mu) \).

\[
\langle 0|\bar{q}(0)\sigma_{\alpha\beta}q(x)|\gamma^{(\lambda)}(q)\rangle = i e_q \chi \langle \bar{q}q \rangle \left( e^{(\lambda)}_{\alpha}q_{\beta} - e^{(\lambda)}_{\beta}q_{\alpha} \right) \int_0^1 du e^{-iu(xz)} \phi_\gamma(u, \mu). \tag{4}
\]

\( \phi_\gamma(u, \mu \geq 1 \text{ GeV}) \) is believed to be not far from the asymptotic form \( \phi_\gamma^s(u) = 6u(1 - u) \). \( \chi \) was estimated using the vector dominance approximation and QCD sum rules 38. \( \chi \langle \bar{q}q \rangle \simeq 40 - 70 \text{ MeV} \). However, any direct experimental evidence on both \( \chi \) and \( \phi_\gamma(u) \) is absent. This structure can be studied in experiments similar to the studies of coherent dijets in pion dissociation by the E791 collaboration 2.

Since the CE and CO contributions lead to final states with different helicity, they do not interfere and the dijet cross section is given by the incoherent sum, for the linearly polarized photon

\[
\left. \frac{d\sigma_{\gamma \rightarrow 2 \text{jets}}}{d\phi dq^2_{1} dddz} \right|_{t=0} = \sum_q e_q^2 \alpha_E M q^2 (1 + \xi^2)^2 \left( 1 - 4z\bar{z}\cos^2\phi \right) |J_{\text{CE}}|^2 + \frac{\pi^2 \alpha_s^2 \chi \langle \bar{q}q \rangle^2}{N_c^2 q^2} |J_{\text{CO}}|^2, \tag{5}
\]

where \( \phi \) is the azimuthal angle between the jet direction and the photon polarization \( (e^{(\lambda)} \cdot q_\perp) \sim \cos \phi \). \( J_{\text{CE}} \) and \( J_{\text{CO}} \) are the CE and CO amplitudes respectively. Note that the CE contribution is \( \sim 1/q^2_\perp \) and the CO contribution is suppressed by one extra power of \( q^2_\perp \) which follows from twist counting. The different \( \phi \) dependence can be traced to the fact that the \( \bar{q}q \) pair is produced in a state with orbital angular momentum \( L_z = 0 \) and \( L_z = \pm 1 \) for the CO and CE contributions, respectively. The CE contribution originates from the region of large momenta flowing through the photon vertex. To leading order (LO) in the strong coupling \( \alpha_s = \alpha_s(q_\perp) \) the amplitude is given by the sum of Feynman diagrams of the type shown in Fig. 1, a.

\[
J_{\text{CE}} = i\xi \mathcal{H}'_g(\xi, \xi) + \frac{i\alpha_s N_c}{\pi} \int_1^1 dy \frac{dy}{y + \xi} \mathcal{H}_g(y, \xi), \tag{6}
\]

where \( \mathcal{H}'_g(\xi, \xi) = d\mathcal{H}_g(y, \xi)/dy|_{y=\xi} \). The second term in Eq. 6 originates from the diagrams with additional gluon exchange between the t-channel gluons, see Fig. 1b. It corresponds to the leading at large energies (enhanced by \( \log \xi \) ) NLO contribution. Since \( \mathcal{H}_g(y, \xi) \sim G(y) \) at \( y \gg \xi \) and as the factor \( \alpha_s N_c/(\pi y) \) is nothing but the low-\( y \) limit of the DGLAP gluon splitting function, the integral in Eq. 6 can be identified to logarithmic accuracy with the unintegrated gluon distribution \( f(\xi, q^2) = \partial G(\xi, q^2)/\partial \ln q^2 \).
For the nonperturbative CO contribution the large momenta are not allowed in the photon vertex and the factorization formula contains a convolution with the photon distribution amplitude. In this case an additional hard gluon exchange is mandatory and the diagram in Fig. 1b presents one example of the existing 31 LO contributions. Similar to the pion case we found that the result for the amplitude may be approximated well by the sum of two contributions in analogy to Eqs. (2) and (3). The origin of the end-point divergence is the same as in a pion dissociation. Assuming that the photon distribution amplitude is close to the asymptotic form, we obtain $J_{CO} \sim z(1-z)$ for both integration regions, up to small corrections. The presence of nonfactorizable contribution does not have, therefore, any significant effect on the jet distribution but mainly influences the normalization.

In the numerical calculation performed for HERA kinematics we have introduced an infrared cutoff $u_{\text{min}} = \mu_{\text{IR}}^2/q_\perp^2$ to regularize the nonfactorizable contribution, $\mu_{\text{IR}} = 500 \text{ MeV}$. We found that the nonperturbative CO contribution integrated over $\phi$, $z$ and $t$ is of the order of $d\sigma_{\text{CO}}/d\sigma_{\text{CE}} \sim (7 \pm 2 \text{ GeV})^2 \cdot \frac{\alpha_s(q_\perp)^2}{q_\perp^2} \left( \frac{\chi(\bar{q}q)}{50 \text{ MeV}} \right)^2$. (7)

For $q_\perp > 4 \text{ GeV}$ the cross section is dominated by the perturbative CE contribution, for smaller transverse momenta the dijet cross section is saturated by the CO contribution. The transition between the two different regimes is seen very clearly from the dependence of the cross section on the dijet longitudinal momentum fraction $t$ and the azimuthal angle. At $q_\perp > 4 \text{ GeV}$ the $z$-distribution is almost flat, while the $\phi$ distribution is almost purely $\sim 1 - \cos^2 \phi$. In contrast to this at $q_\perp < 4 \text{ GeV}$ the $z$-distribution becomes comparable with $\sim z^2(1-z)^2$ while the $\phi$-distribution becomes flat. Our main result is that the nonperturbative CO contribution is large in the region of intermediate $q_\perp \sim 2 - 4 \text{ GeV}$ and can be clearly separated from the perturbative contribution by a different $z$- and $\phi$-dependence. Observation of the CO contribution would be the first clear evidence for the chirality violation in hard processes and also provide the first direct measurement of the magnetic susceptibility of the quark condensate.

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