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MULTIPLE-POMERON SPLITTING IN QCD – A NOVEL ANTI-SHADOWING EFFECT IN COHERENT DIJET PRODUCTION ON NUCLEI

N.N.Nikolaev[+*], *W.Schäfer*[*], *G.Schwiete*[*]

[+]*L.D.Landau Institute for Theoretical Physics RAS*
117940 Moscow, Russia

[*]*Institut für Kernphysik (Th), FZ Jülich*
D-52425 Jülich, Germany

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We discuss the salient features of Pomeron splitting mechanism for coherent diffraction of pions into hard dijets on nuclei. Our findings include anti-shadowing multiple-Pomeron splitting expansion for diffractive amplitudes, exact cancellation of nuclear attenuation and broadening/anti-shadowing effects to leading twist and parameter-free perturbative calculation of nuclear-rescattering driven higher twist correction. We comment on the pQCD interpretation of the E791 results on diffractive dijets.

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At QCD parton level diffraction dissociation of photons and hadrons is modeled by excitation of $q\bar{q}, q\bar{q}g, ...$ Fock states which are lifted on their mass shell through the t-channel exchange of a QCD Pomeron with the target hadron. The color singlet two-gluon structure of the Pomeron gives rise to the two distinct forward $q\bar{q}$ dijet production subprocesses: The first one, of Fig.1a, is a counterpart of the classic Landau, Pomeranchuk, Feinberg, and Glauber, [1 – 3] mechanism of diffraction dissociation of deuterons into the proton-

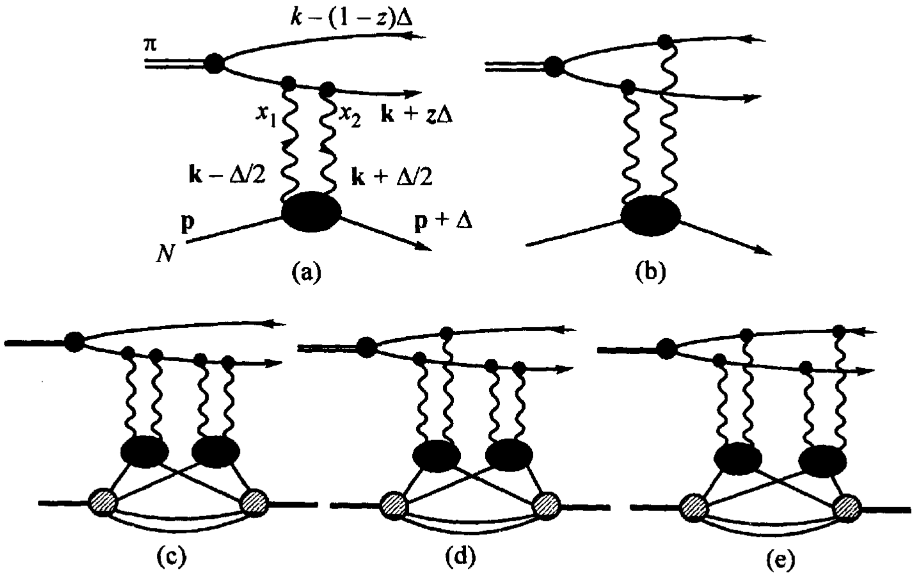

Fig.1. Sample Feynman diagrams for diffractive dijet excitation in πN collisions (diagrams 1a, 1b) and typical rescattering corrections to the nuclear coherent amplitude (diagrams 1c , 1d, 1e)

neutron continuum and can be dubbed the splitting of the beam particle into the dijet, because the transverse momentum $\boldsymbol{k}$ of jets comes from the intrinsic transverse momentum of quarks and antiquarks in the beam particle. Specific of QCD is the mechanism of Fig.1b where jets receive a transverse momentum from gluons in the Pomeron. In an extension of their earlier work [4], Nikolaev and Zakharov have shown in 1994 [5], that the second mechanism dominates at a sufficiently large $\boldsymbol{k}$ and in this regime diffractive amplitudes are proportional to the differential (unintegrated) gluon structure function $\mathcal{F}(x,\boldsymbol{k}^2) = \partial G(x,\boldsymbol{k}^2)/\partial \log \boldsymbol{k}^2$. Correspondingly, this mechanism has been dubbed splitting of the Pomeron into dijets. In diffractive DIS the Pomeron splitting dominates at $k \gg Q$, whereas the somewhat modified Landau et al. mechanism dominates at $k \lesssim Q$.

Motivated by the recent data from the E791 Fermilab experiment [6], in this communication we discuss peculiarities of the Pomeron splitting contribution to diffractive excitation of pions into dijets on free nucleon and nuclear targets. First, because of the non-pointlike nature of the $\pi \to q\bar{q}$ vertex, it is precisely the Pomeron splitting mechanism which dominates at large $\boldsymbol{k}$. Second, in the Pomeron splitting regime diffraction amplitudes are found to be proportional to the much discussed pion distribution amplitude [7]. Third, we show that mutiple-Pomeron splitting emerges as the sole mechanism of nuclear broadening of the jet momentum distribution. Fourth,we derive an antishadowing property of a multiple-Pomeron splitting expansion of nuclear diffraction amplitudes. Fifth, we establish a model-independence of the higher twist effect from nuclear rescatterings.

We first turn to the description of the nucleon amplitude $\pi p \to (q\bar{q})p$ which is a building block of the nuclear multiple scattering series. In a slight adaption of the results of [4, 5] one should replace the pointlike $\gamma^* q\bar{q}$ vertex $eA_\mu \bar{\Psi}\gamma_\mu \Psi$ by the non–pointlike $\pi q\bar{q}$ vertex $i\Gamma(M^2)\bar{\Psi}\gamma_5\Psi$. Here $M^2 = (\boldsymbol{k}^2 + m_f^2)/z(1-z)$ is the invariant mass squared of the dijet system, the pion momentum is shared by the jets in the partitioning $z, 1-z$. The explicit form of the spinor vertex reads

$$\overline{\Psi}_\lambda(\boldsymbol{k})\gamma_5\Psi_{\bar{\lambda}}(-\boldsymbol{k}) = \frac{\lambda}{\sqrt{z(1-z)}}[m_f\delta_{\lambda-\bar{\lambda}} - \sqrt{2}\boldsymbol{k}\cdot\boldsymbol{e}_{-\lambda}\delta_{\lambda\bar{\lambda}}], \tag{1}$$

where m_f is the quark mass, λ and $\bar{\lambda}$ are the quark and antiquark helicities and $\boldsymbol{e}_\lambda = -(\lambda \boldsymbol{e}_x + i\boldsymbol{e}_y)/\sqrt{2}$. The two helicity amplitudes $\Phi_0(z,\boldsymbol{k},\boldsymbol{\Delta})$ for $\lambda+\bar{\lambda}=0$ and $\Phi_1(z,\boldsymbol{k},\boldsymbol{\Delta})$ for $\lambda+\bar{\lambda}=\pm 1$, can be cast in the form (we concentrate on the forward limit, $\boldsymbol{\Delta}=0$)

$$\Phi_0(z,\boldsymbol{k}) = \alpha_S(\boldsymbol{k}^2)\sigma_0\left[\int d^2\boldsymbol{\kappa}\, m_f\psi_\pi(z,\boldsymbol{k})f^{(1)}(\boldsymbol{\kappa}) - \int d^2\boldsymbol{\kappa}\, m_f\psi_\pi(z,\boldsymbol{\kappa})f^{(1)}(\boldsymbol{k}-\boldsymbol{\kappa})\right], \tag{2}$$

and Φ_1 is obtained from the substitution $m_f\psi(z,\boldsymbol{k}) \to \boldsymbol{k}\psi(z,\boldsymbol{k})$. We introduced

$$\sigma_0 = \frac{4\pi}{3}\int d^2\boldsymbol{k}\frac{\mathcal{F}(\frac{1}{2}x_{\mathbb{P}},\boldsymbol{k}^2)}{\boldsymbol{k}^4}, \; f^{(1)}(\boldsymbol{k}) = \frac{4\pi}{3\sigma_0}\frac{\mathcal{F}(\frac{1}{2}x_{\mathbb{P}},\boldsymbol{k}^2)}{\boldsymbol{k}^4}, \tag{3}$$

$x_{\mathbb{P}} = M^2/W^2$, W is the total πp cms-energy, herebelow the argument $x = x_{\mathbb{P}}/2$ of the unintegrated gluon structure function accounts for the skewedness of the t–channel gluons [8]. The distribution $f^{(1)}(\boldsymbol{k})$ is normalized to unity, $\int d^2\boldsymbol{k}\, f^{(1)}(\boldsymbol{k}) = 1$, and σ_0 is the soft gluon exchange dominated nonperturbative parameter.The radial wave function of the $q\bar{q}$ Fock state of the pion is related to the $\pi q\bar{q}$ vertex function as

$$\psi_\pi(z,\boldsymbol{k}) = \frac{N_c}{4\pi^3 z(1-z)}\frac{\Gamma_\pi(M^2)}{(M^2-m_\pi^2)} \tag{4}$$

and is normalized to the $\pi \to \mu\nu$ decay constant $F_\pi = 131$ MeV through

$$F_\pi = \int d^2k dz m_f \psi_\pi(z,\boldsymbol{k}) = F_\pi \int_0^1 dz \phi_\pi(z)\,, \qquad (5)$$

where $\phi_\pi(z)$ is the pion distribution amplitude [7]. Finally, our normalization of helicity amplitudes is such, that the differential cross section of forward dijet production equals

$$\left.\frac{d\sigma_D}{dz dk^2 d\Delta^2}\right|_{\Delta=0} = \frac{\pi^3}{24}\left\{|\Phi_0|^2 + |\Phi_1|^2\right\}. \qquad (6)$$

We turn to the discussion of the asymptotics for large jet momenta $\boldsymbol{k}$. The first term in eq.(2) comes from the Landau et al. pion splitting mechanism of Fig.1a, whereas the second one is the contribution from the Pomeron splitting of Fig.1b. Because it is a convolution $\psi_\pi \otimes f^{(1)}$, it is necessarily a broader function than $\psi_\pi(z,\boldsymbol{k})$ and thus will take over if only $\boldsymbol{k}$ is large enough. For the quantitative estimate, we remind the reader, that for $x \sim 10^{-2}$ relevant to the kinematics of E791, the large-$\boldsymbol{k}$ behaviour of $f^{(1)}(\boldsymbol{k})$ is well described by the inverse power law $f^{(1)}(\boldsymbol{k}) \propto \boldsymbol{k}^{-2\delta}$ with an exponent $\delta \sim 2.15$ [9]. Clearly, $f^{(1)}(\boldsymbol{k})$ decreases much slower than $\psi(z,\boldsymbol{k})$, and hence the asymptotics of the convolution integral is controlled by the asymptotics of $f^{(1)}(\boldsymbol{k})$:

$$\int d^2\kappa m_f \psi_\pi(z,\boldsymbol{\kappa}) f^{(1)}(\boldsymbol{k}-\boldsymbol{\kappa}) \approx f^{(1)}(\boldsymbol{k}) \int d^2\kappa m_f \psi_\pi(z,\boldsymbol{\kappa}) = f^{(1)}(\boldsymbol{k})\,\phi_\pi(z)\,F_\pi\,, \qquad (7)$$

which shows that in this regime the dijet momentum comes entirely from the momentum of gluons in the Pomeron. Furthermore, in the same regime diffraction into dijets probes the pion distribution amplitude $\phi_\pi(z)$. We note in passing, that the explicit proportionality to the unintegrated gluon structure function $\mathcal{F}(x,\boldsymbol{k}^2)$ is in marked contrast to the erroneous claims in [10], where the diffractive amplitude was alleged to be proportional to the *integrated* gluon structure function $G(x_{\mathbb{P}},\boldsymbol{k}^2)$.

The generalization to the nuclear target is most conveniently done in the $q\bar{q}$ color dipole representation [11], for brevity we focus on Φ_0, which takes the form

$$\Phi_0(z,\boldsymbol{k}) = \int d^2\boldsymbol{r}\, e^{-i\boldsymbol{k}\boldsymbol{r}}\, \sigma(x,\boldsymbol{r})\, m_f \Psi_\pi(z,\boldsymbol{r})\,, \qquad (8)$$

Here the color dipole cross section $\sigma(x,\boldsymbol{r})$ for the interaction of a color dipole of transverse size $\boldsymbol{r}$ and the color dipole distribution amplitude $\Psi_\pi(z,\boldsymbol{r})$ in the pion in the impact parameter plane are defined through

$$\sigma(x,\boldsymbol{r}) = \alpha_S(\boldsymbol{k}^2)\sigma_0 \int d^2\kappa f^{(1)}(\boldsymbol{\kappa})\left[1 - e^{i\boldsymbol{\kappa}\boldsymbol{r}}\right]\,,\ \Psi_\pi(z,\boldsymbol{r}) = \int \frac{d^2\boldsymbol{k}}{(2\pi)^2}\psi_\pi(z,\boldsymbol{k})e^{i\boldsymbol{k}\boldsymbol{r}}\,. \qquad (9)$$

The Glauber – Gribov representation [12, 13] of the nuclear amplitude is readily obtained [11] by substituting in eq.(8)

$$\sigma(x,\boldsymbol{r}) \to \sigma_A(x,\boldsymbol{r}) = 2\int d^2\boldsymbol{b}(1 - \exp(-\frac{1}{2}\sigma(x,\boldsymbol{r})T_A(\boldsymbol{b})),$$

where $\boldsymbol{b}$ is the pion-nucleus impact parameter, $T_A(\boldsymbol{b})$ is the nuclear optical thickness, $\int d^2\boldsymbol{b} T_A(\boldsymbol{b}) = A$. Typical nuclear double scattering diagrams of Figs.1c–e can conveniently be classified as nuclear shadowing of the pion splitting (Fig.1c), nuclear shadowing of single Pomeron splitting (Fig.1d) and double Pomeron splitting (fig.1e) contributions. The j Pomeron splitting is due to diagrams, in which j of the Pomerons exchanged between the color dipole and the nucleus couple with one gluon to the quark and with one gluon to the antiquark and involve the j-fold convolutions of the unintegrated gluon distribution

$$f^{(j)}(\boldsymbol{k}) = \int d^2\boldsymbol{\kappa}_1 ... d^2\boldsymbol{\kappa}_j f^{(1)}(\boldsymbol{\kappa}_1) f^{(1)}(\boldsymbol{\kappa}_2) ... f^{(1)}(\boldsymbol{\kappa}_j) \delta(\boldsymbol{k} - \sum_{i=1}^{j} \boldsymbol{\kappa}_i) . \tag{10}$$

After some algebra we obtain the central result of this paper – the multiple-Pomeron splitting expansion for nuclear diffractive amplitude

$$\Phi_0^{(A)}(z, \boldsymbol{k}, \boldsymbol{\Delta}) = 2m_f \int d^2\boldsymbol{b} e^{-i\boldsymbol{b}\boldsymbol{\Delta}} \left\{ \Psi_\pi(z, \boldsymbol{k}) \left[1 - \exp\left(- \frac{\sigma_{eff}(\boldsymbol{k}^2)}{2} T_A(\boldsymbol{b}) \right) \right] - \right.$$

$$\left. - \sum_{j \geq 1} \int d^2\boldsymbol{\kappa} \Psi_\pi(z, \boldsymbol{\kappa}) f^{(j)}(\boldsymbol{k} - \boldsymbol{\kappa}) \frac{1}{j!} \left[\frac{\sigma_{eff}(\boldsymbol{k}^2) T_A(\boldsymbol{b})}{2} \right]^j \exp\left[- \frac{\sigma_{eff}(\boldsymbol{k}^2)}{2} T_A(\boldsymbol{b}) \right] \right\} . \tag{11}$$

Here we reinstated a dependence on the momentum transfer $\boldsymbol{\Delta}$, which in the final result can be related to the nuclear charge form factor $G_{em}(\boldsymbol{\Delta}^2)$. Those Pomeron exchanges, in which both gluons couple to the same constituent in the color dipole give rise to nuclear attenuation with an effective cross section $\sigma_{eff} = \alpha_S(\boldsymbol{k}^2)\, \sigma_0$, the appearance of the strong coupling at the hard scale $\boldsymbol{k}^2$ being a reminder of color transparency. Remarkably, all the multiple Pomeron splittings enter with the same sign and are an anti-shadowing, i.e., enhancement, contribution to the impulse approximation, i.e., single-pomeron splitting, term $\propto f^{(1)}(\boldsymbol{k})$. Again, one may obtain $\Phi_1^{(A)}$ by substituting $m_f \psi_\pi(z, \boldsymbol{k}) \to \boldsymbol{k} \psi_\pi(z, \boldsymbol{k})$. Obviously, similar formulas apply for dijet photoproduction as well as deep inelastic diffractive scattering. As we argued above, for the pion beam the multiple Pomeron splitting terms will rather quickly take over.

We estimate the large-$\boldsymbol{k}$ behaviour of $f^{(j)}(\boldsymbol{k})$ and of their convolution with the pion wave function by making use of the small-$\boldsymbol{\kappa}$ expansion of $f^{(1)}(\boldsymbol{k} - \boldsymbol{\kappa})$. To the next-to-leading twist one can readily derive, that the power law asymptotics $f^{(1)}(\boldsymbol{k}) \propto k^{-2\delta}$ entails the large-$\boldsymbol{k}$ behaviour of $f^{(j)} \otimes \psi_\pi$ given by

$$\int d^2\boldsymbol{\kappa} m_f \psi_\pi(z, \boldsymbol{\kappa}) f^{(j)}(\boldsymbol{k} - \boldsymbol{\kappa}) \simeq$$

$$\simeq F_\pi \phi_\pi(z) j f^{(1)}(\boldsymbol{k}) \left[1 + \frac{\delta^2}{\boldsymbol{k}^2} \left(\langle \kappa_\pi^2(z) \rangle + \frac{4\pi^2}{3\sigma_0} (j-1) G(x, \boldsymbol{k}^2) \right) \right] . \tag{12}$$

Here the proportionality to j of the leading twist reflects the broadening properties of the multiple convolution. Note that due to the factor $j - 1$ this broadening is still stronger for nuclear higher twist. The higher twist induced by intrinsic momentum of constituents in the pions is controlled by the nonperturbative parameter $\langle \kappa_\pi^2(z) \rangle = m_f \int d^2\boldsymbol{\kappa}\, \boldsymbol{\kappa}^2 \psi_\pi(z, \boldsymbol{\kappa}) / \phi_\pi(z) F_\pi$. The realistic soft wave functions of the pion give $\langle \kappa_\pi^2(z) \rangle \sim 0.15 \div 0.2\,\mathrm{GeV}^2$ [14].

Making use of the explicit j dependence of the large-$\boldsymbol{k}$ asymptotics (12) in the multiple-Pomeron splitting expansion (11) one readily finds the most intriguing result - an *exact* *cancellation* of the broadening and anti-shadowing and the soft shadowing effects to the both leading and next-to-leading order:

$$\Phi_0^{(A)}(z,\boldsymbol{k},\boldsymbol{\Delta}=0) \simeq -A\frac{4\pi F_\pi \phi_\pi(z)\alpha_S(\boldsymbol{k}^2)\mathcal{F}(x,\boldsymbol{k}^2)}{3\boldsymbol{k}^4}\times$$
$$\times\left\{1+\frac{\delta^2}{\boldsymbol{k}^2}\left[\langle\kappa_\pi^2(z)\rangle+\frac{\pi C_A A\alpha_S(\boldsymbol{k}^2)G(\frac{1}{2}x_{\mathbf{P}},\boldsymbol{k}^2)}{2\langle R_{ch}^2\rangle}\right]\right\}. \quad (13)$$

Namely, in contrast to the nonperturbative higher twist $\propto \langle\kappa_\pi^2(z)\rangle$ the arguably dominant, nuclear rescattering driven, higher twist is parameter free and perturbatively calculable. It is controlled by the scaling violations in the gluon structure function $G(x,Q^2)$ and rises for heavy nuclei $\propto A^{1/3}$. Although to both the leading and next-to-leading twist any dependence on the soft rescattering parameter $\sigma_{eff}(\boldsymbol{k}^2)$ disappears that doesn't imply that higher order rescatterings are unimportant: the rescatterings up to j =9–10 do contribute and $\langle j\rangle = 4.2$ for the ^{192}Pt target and $\langle j\rangle = 2.05$ for the ^{12}C target. In contrast to $\Phi_0(z,\boldsymbol{k},\boldsymbol{\Delta})$ the large-$\boldsymbol{k}$ asymptotics of $\Phi_1(z,\boldsymbol{k})$ is of pure higher twist:

$$\Phi_1^{(A)}(z,\boldsymbol{k},\boldsymbol{\Delta}=0) \simeq -A\frac{4\pi F_\pi \phi_\pi(z)\alpha_S(\boldsymbol{k}^2)\mathcal{F}(x,\boldsymbol{k}^2)}{3m_f\boldsymbol{k}^4}\frac{\delta\langle\kappa_\pi^2(z)\rangle}{\boldsymbol{k}^2}. \quad (14)$$

For large-$\boldsymbol{k}$, where the terms $\propto \Psi_\pi(z,\boldsymbol{k})$ die out, using eq.(12) we obtain upon integration over the transverse momentum transfer $\boldsymbol{\Delta}$ the following simple form for the large-$\boldsymbol{k}$ asymptotics of the diffractive dijet cross section:

$$\frac{d\sigma_D}{dzd\boldsymbol{k}^2} = \frac{2\pi^5}{27}F_\pi^2\phi_\pi^2(z)G_{em}^2(x_{\mathbf{P}}^2m_N^2)\alpha_S^2(\boldsymbol{k}^2)\left[\frac{\mathcal{F}(x,\boldsymbol{k}^2)}{\boldsymbol{k}^4}\right]^2\frac{3A^2}{\langle R_{ch}^2\rangle}\times$$
$$\times\left\{1+\frac{\delta^2}{\boldsymbol{k}^2}\left[2\langle\kappa_\pi^2(z)\rangle+\frac{\langle\kappa_\pi^2(z)\rangle^2}{m_f^2}+\frac{2\pi C_A A\alpha_S(\boldsymbol{k}^2)}{3\langle R_{ch}^2\rangle}G(x,\boldsymbol{k}^2)\right]\right\}. \quad (15)$$

Here the parameter $C_A \approx 1$ depends slightly on the shape of the nuclear matter distribution, R_{ch} is the nuclear charge radius, the factor $G_{em}^2(x_{\mathbf{P}}^2m_N^2)$ accounts for the finite longitudinal momentum transfer $x_{\mathbf{P}}m_N$ to the target nucleus. The second term in curly braces in (15) is the higher twist contribution $\propto \Phi_1^2$. Notice that the behaviour $\propto k^{-8}$ expected from naive dimensional counting receives potentially large corrections from the scaling violations in $\mathcal{F}(x,\boldsymbol{k}^2)$. Remarkably, the leading twist term does not contain any free parameters and thus is the perturbatively calculable quantity. The z-dependence of the nonperturbative parameter $\langle\kappa_\pi^2(z)\rangle$ is sufficiently weak and does not preclude the determination of the pion distribution amplitude from the z-distribution of dijets.

The detailed numerical analysis for the kinematics of the E791 experiment shall be reported elsewhere. In Fig.2 we only show the numerical result of the $\boldsymbol{k}$ dependence of the dijet cross section with the (unnormalized) preliminary data from E791. We note in passing that the region of jet momenta $k \lesssim 1.5$ GeV is contaminated by diffractive excitation of heavy mesons a_1, π', etc. and in this region the use of plane wave parton model formulas is not warranted. Our calculations for smaller k only serve to give an idea on how the pion splitting dominates at small k and how the pomeron splitting mechanism takes over beyond the dip in Φ_0^2 and Φ_1^2. In the pomeron splitting dominance region of $k > 1.5$ Gev we find good agreement with experiment. We wish to warn the reader

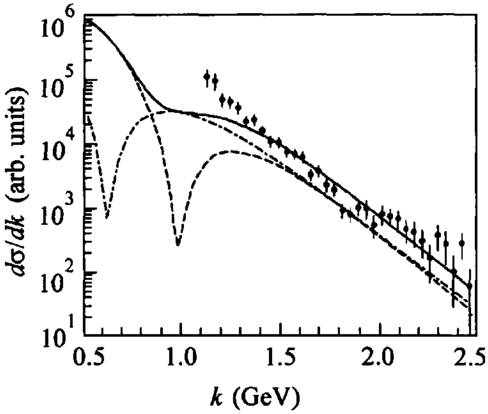

Fig.2. The E791 data [6] for the differential diffractive dijet cross section $d\sigma/dk$ for the ^{196}Pt target with the theoretical calculations. The data are not normalized. The dash–dotted line shows the contribution of the helicity amplitude $\Phi_0^{(A)}$, the dashed line is the contribution from $\Phi_1^{(A)}$. The solid line is the total result

however, that for the kinematics of E791 the dijet cross section receives huge contributions from higher twists, which is exemplified by the large contribution from higher twist amplitude (14), i.e., the second term in the curly braces in the expansion (15).

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1. L.D.Landau and I.Ya.Pomeranchuk, ZhETF **24**, 505 (1953).
2. I.Ya.Pomeranchuk and E.L. Feinberg, Doklady Akademii Nauk SSSR **93**, 439 (1953); E.L.Feinberg, ZhETF **28**, 242 (1955); E.L.Feinberg and I.Ya.Pomeranchuk, Nuovo Cim. (Suppl.) **4**, 652 (1956); A.I.Akhiezer and A.G.Sitenko, ZhETF **32**, 794 (1957); **33**, 1040 (1957).
3. R.J.Glauber, Phys. Rev. **99**, 1515 (1955).
4. N.N.Nikolaev and B.G.Zakharov, Z. Phys. **C53**, 331 (1992).
5. N.N.Nikolaev and B.G.Zakharov, Phys. Lett. **B332**, 177 (1994).
6. D.Ashery, *Invited talk at X International Light-Cone Meeting*, Heidelberg, June 2000, hep-ex/0008036.
7. V.L.Chernyak and A.R.Zhitnitsky, Phys. Rept. **112**, 173 (1984); P.Kroll, πN Newslett. **15**, 205 (1999); R.Jakob and P.Kroll, Phys. Lett. **B315**, 463 (1993), Erratum-ibid. **B319**, 545 (1993); G.P.Lepage and S.J.Brodsky, Phys. Rev. **D22**, 2157 (1980); S.J.Brodsky, H.-C.Pauli, and S.S.Pinsky, Phys. Rept. **301**, 299 (1998).
8. A.V.Radyushkin, Phys. Lett. **B385**, 333 (1996); K.Golec-Biernat, J.Kwiecinski, and A.D.Martin, Phys. Rev. **D58**, 094001 (1998).
9. I.P.Ivanov and N.N.Nikolaev, hep-ph/0004206.
10. L.L.Frankfurt, G.A.Miller, and M.I.Strikman, Phys. Lett. **B304**, 1 (1993).
11. N.N.Nikolaev and B.G.Zakharov, Z. Phys. **C49**, 607 (1991).
12. R.J.Glauber, in: *Lectures in Theoretical Physics*, Eds. W.Brittain and L.G.Dunham, Interscience, New York, 1959, Vol.1; R.J.Glauber and G.Matthiae, Nucl. Phys. **B21**, 135 (1970).
13. V.N.Gribov, Sov. Phys. JETP **29**, 483 (1969); ZhETF **56**, 892 (1969).
14. W.Jaus, Phys. Rev. **D44**, 2851 (1991); G.Schwiete, Diploma Thesis, University of Bonn, forthcoming.