PHYSICAL REVIEW D, VOLUME 59, 091502

# Two-pion light-cone distribution amplitudes from the instanton vacuum

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(Received 8 July 1998; published 18 March 1999)

We calculate the two-pion light-cone distribution amplitudes in the effective low-energy theory based on the instanton vacuum. These generalized distribution amplitudes describe the soft (non-perturbative) part of the process $\gamma^*\gamma\to\pi\pi$ in the region where the c.m. energy is much smaller than the photon virtuality. They can also be used in the analysis of exclusive processes such as $\gamma^* p\to p+2\pi,3\pi$, etc. [S0556-2821(99)50207-5]

PACS number(s): 12.38.Lg, 13.60.Fz, 13.60.Le

Hadron production in photon-photon collisions at low invariant masses has become a subject of great interest recently [1]. Under certain conditions such processes are amenable to a partonic description, if the virtuality of the photon, $Q^2$, is much larger than the squared c.m. energy, $W^2$. The simplest such process, $\gamma^*\gamma\to\pi^0$, provides a unique framework for studying the pion light-cone distribution amplitude [2]. Recently, Diehl *et al.* have argued that also production of a pair of hadrons, $\gamma^*\gamma\to h\bar{h}$, can be described in factorized form [3]. The amplitudes for these processes contain unknown non-perturbative functions describing the exclusive fragmentation of a quark-antiquark pair into two hadrons.[1] In addition to the usual parton momentum fraction with respect to the total momentum of the hadronic final state, these functions depend on the distribution of longitudinal momentum between the two hadrons, as well as on the invariant mass of the produced system, $W^2$. In particular, the authors of Ref. [3] consider the production of two pions, the soft part of which is contained in a generalized two-pion distribution amplitude. Quantitative estimates of these functions are urgently needed for the computation of the cross section for such processes.

In this Rapid Communication we calculate the two-pion distribution amplitudes at a low normalization point ($\sim$600 MeV) in the effective low-energy theory based on the instanton vacuum [5]. This approach has recently been used to study the pion distribution amplitude [6,7], which was found to be close to the asymptotic form, consistent with the recent CLEO measurements [8]. Our aim is to discuss qualitative features of the two-pion amplitudes, and to calculate the convolution with the tree-level hard scattering amplitude. An account of a more extensive investigation will be published elsewhere [9].

By crossing symmetry, the process $\gamma^*\gamma\to h\bar{h}$ is related to virtual Compton scattering, $\gamma^* h\to\gamma h$, which can be factorized into a hard photon-parton scattering amplitude and an off-forward parton distribution (OFPD) [10]. In addition to the light-cone momentum fraction, $x$, the latter depends on the longitudinal component of the momentum transfer to the hadron, $\xi$. The off-forward isosinglet quark-antiquark distributions in the nucleon have been computed in the effective low-energy theory in Ref. [11]. It was found that at $x=\pm\xi/2$, the OFPD exhibits characteristic discontinuities (see also Ref. [12]). Actually, these would-be discontinuities are reduced to sharp, but continuous crossovers when one takes into account the momentum dependence of the dynamical quark mass. We shall find similar behavior in the two-pion amplitudes considered here.

The generalized two-pion distribution amplitudes (GDA's) are defined as [3]

$$\Phi(z,\zeta,W^2)=\frac{1}{4\pi}\int dx^- e^{-(i/2)zP^+x^-} \times\langle\pi^a(p_1)\pi^b(p_2)|\overline{\psi}(x) \times n^\mu\gamma_\mu T\psi(0)|0\rangle|_{x^+=0,x_\perp=0}. \tag{1}$$

Here, $n_\mu=(1,0,0,1)$ is a light-like vector ($n^2=0$), and for any vector, $V$, the ''plus'' and ''minus'' light-cone components are defined as $V^+\equiv(n\cdot V)=V^0+V^3$, $V^-=V^0-V^3$. The outgoing pions have momenta $p_1,p_2$, and $P\equiv p_1+p_2$ is the total momentum of the final state. $T$ is a flavor matrix; $T=1$ for the isosinglet, $T=\tau^3$ for the isovector GDA.

The GDA, Eq. (1), depends on the quark momentum fraction with respect to the total momentum of the two-pion state, $z$; the variable $\zeta\equiv p_1^+/P^+$ characterizing the distribution of longitudinal momentum between the two pions, and the invariant c.m. energy, $W^2=P^2$. In what follows we shall work in the reference frame where $P_\perp=0$. In this frame

$$P^-=\frac{W^2}{P^+},$$

$$p_1^-=\frac{W^2(1-\zeta)}{P^+}, \tag{2}$$

$$p_\perp^2\equiv p_{1\perp}^2=p_{2\perp}^2=W^2\zeta(1-\zeta)-m_\pi^2.$$

From these relations, one obtains the kinematical constraint

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[1]Similar functions have previously been introduced to describe multi-hadron production at large invariant masses [4].

 

$$\zeta(1-\zeta)\geqslant\frac{m_\pi^2}{W^2}. \tag{3}$$

We compute the GDA, Eq. (1), within the effective theory of pions interacting with quarks with a dynamical mass, which has been derived from the instanton model of the QCD vacuum [5]. The coupling of the pion field to the quarks, whose non-linear form is dictated by chiral invariance, is described by the action

$$S_{\rm int}=\int d^4x\,\overline{\psi}(x)\sqrt{M(\partial^2)}e^{i\gamma_5\tau^a\pi^a(x)/F_\pi}\sqrt{M(\partial^2)}\,\psi(x), \tag{4}$$

where $F_\pi=93$ MeV is the weak pion decay constant.

The momentum-dependent dynamical quark mass, $M(-p^2)$, plays the role of an UV regulator. Its form for Euclidean momenta was derived in Ref. [5]. This form factor cuts loop integrals at momenta of order of the inverse average instanton size, $\overline{\rho}^{-1}\approx 600$ MeV. One should note that the value of the mass at zero momentum, $M(0)$, is parametrically small compared to $\overline{\rho}^{-1}$; the product $M(0)\overline{\rho}$ is proportional to the small packing fraction of the instanton medium. Numerically, a value $M(0)=345$ MeV was obtained in Ref. [5].

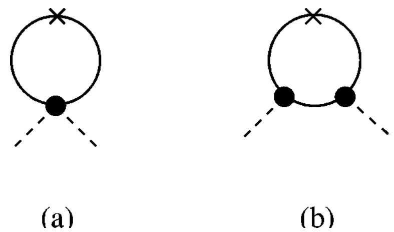

FIG. 1. The Feynman diagrams defining the two contributions to the two-pion distribution amplitude in the effective low-energy theory, $\Phi^{(1)}$ and $\Phi^{(2)}$. The solid line denotes the massive quark propagator, the solid points the vertices obtained by expanding the non-linear quark-pion coupling of Eq. (4) in powers of the pion field, involving the form factors $\sqrt{M(-p^2)}$.

In the leading order of the $1/N_c$-expansion, the matrix element in Eq. (1) is given by the sum of the two Feynman graphs shown in Fig. 1. The first graph, Fig. 1(a), contributes only to the isosinglet part of the GDA ($T=1$), whereas the second graph, Fig. 1(b), contributes to both the isosinglet and isovector part ($T=\tau^3$). We first consider the contribution from Fig. 1(a). By straightforward calculation, we obtain

$$\Phi^{(1)}(z,\zeta,W^2)=\frac{4iN_cP^+}{F_\pi^2}\,\delta^{ab}{\rm Tr}(T)\int\frac{d^4k}{(2\pi)^4}\,\delta(k^+-zP^+)\sqrt{M(k)M(k-P)}\frac{zM(k-P)-(1-z)M(k)}{[k^2-M(k)^2+i0][(k-P)^2-M(k-P)^2+i0]}. \tag{5}$$

This contribution is independent of $\zeta$, which is a consequence of the contact nature of the two-pion-quark vertex in this model, Eq. (4)—the quark loop in Fig. 1(a) is independent of the relative momentum of the pions. Also, $\Phi^{(1)}(z,W^2)=-\Phi^{(1)}(1-z,W^2)$, a property which actually follows from $C$-invariance [3].

Let us first evaluate the integral in Eq. (5), neglecting for a moment the momentum dependence of the constituent quark mass. In this case the result takes a simple form

$$\Phi^{(1)}(z,W^2)=-\frac{N_cM_0^2}{\pi F_\pi^2}\,\delta^{ab}{\rm Tr}(T)\,\theta[z(1-z)](2z-1)\times\int\frac{d^2k_\perp}{(2\pi)^2}\frac{1}{k_\perp^2+M_0^2-W^2z(1-z)}, \tag{6}$$

where $M_0\equiv M(0)$ and $\theta$ denotes the step function. This expression is nonzero at the endpoints, $z\to 0$ and 1. Such behavior of the distribution amplitude would violate the factorization theorem for the process $\gamma^*\gamma\to\pi\pi$, because the amplitude contains the integral

$$I(\zeta,W^2)=\int_0^1 dz\,\frac{2z-1}{z(1-z)}\Phi(z,\zeta,W^2), \tag{7}$$

which would be divergent if $\Phi$ were nonzero at the endpoints. However, such conclusion would be premature, as it is easily seen that the momentum dependence of the dynamical quark mass becomes crucial at the end points (see the discussion in Refs. [7,11]). This means that when computing the isosinglet GDA, one cannot neglect the momentum dependence of the constituent quark mass. For the numerical estimates, we employ a simple numerical fit to the momentum-dependent mass obtained from the instanton vacuum:

$$M(-p^2)=\frac{M_0}{(1+0.5p^2\overline{\rho}^2)^2}. \tag{8}$$

The result for the $\zeta$-independent contribution $\Phi^{(1)}(z,W^2)$ to the isosinglet GDA is shown in Fig. 2 for a value of $W^2=0.25$ GeV$^2$.

It is useful to expand the isosinglet GDA in Gegenbauer polynomials of index 3/2, which are the eigenfunctions of the evolution kernel [13],

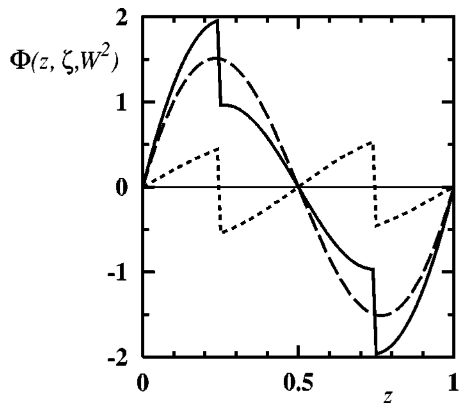


FIG. 2. The contributions to the isosinglet two-pion distribution amplitude, $\Phi(z,\zeta,W^2)$, as a function of $z$, for $\zeta=0.25$ and $W^2=0.25\ \mathrm{GeV}^2$. *Dashed line:* contribution $\Phi^{(1)}$, Eq. (5). *Dotted line:* contribution $\Phi^{(2)}$, Eqs. (13,14,15). *Solid line:* total, $\Phi^{(1)}+\Phi^{(2)}$.

$$\Phi^{(1)}(z,W^2)=\delta^{ab}\mathrm{Tr}(T)6z(1-z)\sum_{\substack{n\\ \mathrm{odd}}} a_n(W^2)C_n^{3/2}(2z-1). \tag{9}$$

Due to the antisymmetry with respect to $z\to 1-z$, only polynomials of odd degree appear in the isosinglet part. The numerical results for the first few coefficients in the range $0\leqslant W^2\leqslant 4M_0^2$ can be approximated by[2]

$$a_1(W^2)=-0.5556(1-1.33W^2),$$

$$a_3(W^2)=-0.036(1-6.2W^2)(1-1.22W^2)^{-2}, \tag{10}$$

$$a_5(W^2)=-0.0036(1-2.0W^2)^{-3/2},$$

where $W^2$ is in $\mathrm{GeV}^2$. One sees that the coefficients decrease rapidly with increasing order. We also give the result for the contribution of $\Phi^{(1)}$ to the isosinglet part of the integral of Eq. (7). With good accuracy the numerical results can be fitted by

$$I^{(1)}(W^2)=-\delta^{ab}\mathrm{Tr}(T)\times 3.58(1-1.0W^2). \tag{11}$$

We now turn to the contribution to the GDA given by the graph of Fig. 1(b). One can easily convince oneself that the neglect of the momentum dependence of the dynamical quark mass in this case does not lead to any violation of factorization. To make the discussion more transparent, we shall evaluate this contribution with a constant quark mass; the logarithmic divergence of the loop integrals can be absorbed in the pion decay constant,

$$F_\pi^2=4N_c(-i)\int\frac{d^4k}{(2\pi)^4}\frac{M_0^2}{(k^2-M_0^2+i0)^2}. \tag{12}$$

All formulas below can easily be generalized to the case of a momentum-dependent mass. Computing the Feynman integral corresponding to Fig. 1(b), we thus obtain

$$\begin{aligned}\Phi^{(2)}(z,\zeta,W^2)=&\tfrac{1}{2}\mathrm{Tr}(T[\tau^a,\tau^b])[\phi^{(2)}(z,\zeta,W^2)\\&-\phi^{(2)}(z,1-\zeta,W^2)]+\mathrm{Tr}(T)\delta^{ab}\\&\times[\phi^{(2)}(z,\zeta,W^2)+\phi^{(2)}(z,1-\zeta,W^2)],\end{aligned} \tag{13}$$

where

$$\phi^{(2)}(z,\zeta,W^2)\overset{z<\zeta}{=}\frac{M_0^2N_c}{\pi F_\pi^2}\theta[z(1-z)]z\int\frac{d^2k_\perp}{(2\pi)^2}\frac{1}{k_\perp^2+M_0^2-W^2z(1-z)}\left(1+\frac{\vec{k}_\perp\cdot\vec{p}_\perp+m_\pi^2z-W^2z(1-\zeta)}{(k_\perp^2+M_0^2)\zeta-2\vec{k}_\perp\cdot\vec{p}_\perp z-m_\pi^2z+W^2(1-\zeta)z^2}\right), \tag{14}$$

$$\phi^{(2)}(z,\zeta,W^2)\overset{z>\zeta}{=}-\{z\to 1-z,\zeta\to 1-\zeta\}. \tag{15}$$

Note the different symmetry of the isosinglet and isovector part with respect to $z\to 1-z$.

Some comments are in order here. The expressions for the GDA given by Eqs. (6) and (13), (14), (15) are valid in the parametrically wide region of invariant c.m. energies $4m_\pi^2<W^2<4M_0^2$. For $W^2>4M_0^2$ the graphs of Fig. 1 develop an imaginary part due to the presence of a two-quark threshold in this model. This singularity is not physical because for momenta of order $\overline{\rho}^{-1}\sim 600$ MeV degrees of freedom not accounted for in the effective theory (e.g., heavy resonances) start to play a significant role. These contributions can in principle be estimated phenomenologically by adding the coupling of the resonances to constituent quarks to the effective action, Eq. (4). Note also that the GDA's computed in this model are real at $W^2<4M_0^2$ only in the leading order of the $1/N_c$-expansion; an imaginary part appears in the next-to-leading order due to pion loop contributions.

The contribution to the GDA given by Eqs. (13),(14), (15) exhibits discontinuities at $z=\zeta$ and $z=1-\zeta$. The nature of these discontinuities is the same as those found in the non-diagonal parton distributions computed in this model [11]. As was shown in Ref. [11], the inclusion of the momentum dependence of the constituent quark mass smears this discontinuity. However, in contrast to the case of non-diagonal parton distributions, the discontinuities in the two-pion GDA

[2]The functions given here should be considered as purely numerical fits in the given range of $W^2$. In particular, the different powers in the denominators have no physical meaning.

do not lead to violation of factorization, so we can be negligent and evaluate Eqs. (13),(14) with a constant quark mass. The result is shown in Fig. 2, for a value of $\zeta=0.25$ and $W^2=0.25\ \mathrm{GeV}^2$, where we also give the total result for the isosinglet GDA.

Let us take in Eq. (14) the chiral limit, $m_\pi=0$, and expand this contribution in powers of $W^2$. In this way one obtains an analytic expression for the GDA,

$$\phi^{(2)}(z,\zeta,W^2) \overset{z<\zeta}{=} \theta[z(1-z)]z\left\{1+\frac{N_c W^2}{8\pi^2F_\pi^2}\times\frac{z[2\zeta(1-z)-(1-\zeta)]}{\zeta}+\cdots\right\}, \quad (16)$$

and for $z>\zeta$ as above, Eq. (15). From this expression we find for the lowest two moments

$$\int_0^1 dz\Phi^{(2)}(z,\zeta,W^2)=\frac{1}{2}(2\zeta-1)\mathrm{Tr}(T[\tau^a,\tau^b])\times\left\{1+\frac{N_c}{24\pi^2F_\pi^2}W^2+\cdots\right\}, \quad (17)$$

$$\int_0^1 dz(2z-1)\Phi^{(2)}(z,\zeta,W^2)=\delta^{ab}\mathrm{Tr}(T)\times\left\{\frac{1}{3}-2\zeta(1-\zeta)+\frac{N_cW^2}{120\pi^2F_\pi^2}\times[1-5\zeta(1-\zeta)]+\cdots\right\}. \quad (18)$$

The first moment, Eq. (17), corresponds to the expansion of the pion electromagnetic form factor for small time-like momenta. The pion charge radius read off from Eq. (17) is $\langle r^2\rangle_{\mathrm{e.m.}}=N_c/(4\pi^2F_\pi^2)$, which coincides with the result of Ref. [5]. Only the isovector part of the GDA contributes to the first moment, as it should be on grounds of C-parity.

Of $\Phi^{(2)}$ only the isosinglet part contributes to the integral Eq. (7). It gives

$$I^{(2)}(\zeta,W^2)=4\delta^{ab}\mathrm{Tr}(T)\left[-\left(1+\frac{1}{2}\log[\zeta(1-\zeta)]\right)+\frac{N_cW^2}{48\pi^2F_\pi^2}\{4\zeta(1-\zeta)+3\zeta^2\log[\zeta(1-\zeta)]+3(1-2\zeta)\log(1-\zeta)\}\right]. \quad (19)$$

Note that this contribution is symmetric with respect to $\zeta\to1-\zeta$.

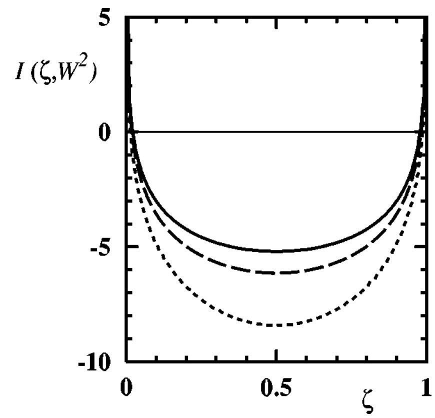


FIG. 3. The integral $I(\zeta,W^2)$ determining the $\gamma^*\gamma\to\pi\pi$ amplitude, Eq. (7), given by the sum of Eqs. (11) and (19), as a function of $\zeta$ for various values of $W^2$. *Solid line:* $W=0.3$ GeV. *Dashed line:* $W=0.5$ GeV. *Dotted line:* $W=2M_0=0.69$ GeV.

The total result for the integral Eq. (7) is given by the sum of Eqs. (11) and (19). In Fig. 3 we plot the total result as a function of $\zeta$ at several values of $W^2$. We see that the absolute value of this integral for intermediate values of $\zeta$ increases with $W^2$, and that the dependence on $\zeta$ is strong only for small values of $\zeta(1-\zeta)$, a region which is difficult to access because of the kinematical constraint, Eq. (3).

The scale dependence of the GDA is governed by the usual Efremov-Radyushkin-Brodsky-Lepage evolution [3,13]. This is obvious when one notes that the hard-scattering kernel of the amplitude for $\gamma^*\gamma\to\pi\pi$ is the same as that for a process with only one meson in the final state (with appropriate quantum numbers). The evolution equations can be solved, in analogy to those for the usual pion wave function [14], by expanding in eigenfunctions of the evolution kernel. In the isosinglet part of the GDA, one has to take into account the mixing with gluon operators [15]; this part will be considered elsewhere [9]. However, one can easily establish the scale dependence of the isovector part of the GDA, which does not mix with gluons. This function does not enter in the amplitude for $\gamma^*\gamma\to\pi\pi$, but it can appear in exclusive processes of the type $\gamma^*p\to Xp$, where $X$ is a two- (or more) pion state (see below). The asymptotic behavior of the isovector GDA is given by

$$\Phi(z,\zeta,W^2)=3\ \mathrm{Tr}(T[\tau^a,\tau^b])z(1-z)(2\zeta-1)F_\pi(W^2), \quad (20)$$

where $F_\pi(W^2)$ is the pion electromagnetic form factor in the time-like domain. Note that, contrary to the one-pion distribution amplitude [6,7], the isovector two-pion GDA calculated in the low-energy effective theory, Eq. (13), is far from the asymptotic form.

In this note we have given a first model estimate of the two-pion GDA's at a low scale. A number of qualitative features have emerged, which are independent of the details of the quantitative approximations made. First, the GDA's generally exhibit discontinuities at $z=\zeta$ and $1-\zeta$, which are in complete analogy to those found in the off-forward parton distributions[3] [11]. As in the OFPD's, the momentum dependence of the dynamical quark mass derived from the instanton vacuum turns these discontinuities into continuous, but sharp crossovers. However, since the discontinuities of the GDA's are not at odds with the factorization of the amplitude for $\gamma^*\gamma \to h\bar{h}$, we are able to estimate the amplitude using the discontinuous GDA's. Second, the isosinglet and -nonsinglet GDA's for the two-pion final state computed at the low scale exhibit different symmetry with respect to $z \to 1-z$.

Our results concerning the $W^2$-dependence of the GDA's and the total amplitude should be regarded as crude estimates. In particular, the imaginary part of the functions at $W^2 > 4M_0^2$ should not be regarded as physical, since it occurs in a region where the effective theory is no longer applicable. A more conservative approach would be to calculate the GDA's in an expansion in the external pion momenta and the pion mass, adopting, so to speak, the philosophy of the chiral Lagrangian. The physical imaginary part, which is due to final state interactions, appears in the next order of the $1/N_c$ expansion and can be estimated using the methods of chiral perturbation theory [9].

The studies reported in this note can easily be generalized to distribution amplitudes for $n > 2$ pions in the final state. Such multi-pion GDA's play a role in the analysis of exclusive processes of the type $\gamma^* p \to Xp$, where $X = 2\pi, 3\pi$ etc.[4] [17]. For instance, the two-pion isovector GDA computed here can be used to analyze the $2\pi$ background in the diffractive leptoproduction of rho mesons.

We acknowledge helpful suggestions by M. Diehl, as well as interesting discussions with L. Frankfurt, V. Petrov, and M. Strikman. It is a pleasure to thank K. Goeke for encouragement and support.

[3]The OFPD's in the pion show qualitatively similar behavior to those in the nucleon [16].

[4]We are grateful to L. Frankfurt and M. Strikman for discussion of this point.

[1] See, e.g., S. J. Brodsky, Talk given at the International Conference on the Structure and the Interactions of the Photon (Photon 97) including the 11th International Workshop on Photon-Photon Collisions, Egmond aan Zee, Netherlands, 1997, Report No. SLAC-PUB-7604, hep-ph/9708345.

[2] P. Kroll and M. Raulfs, Phys. Lett. B **387**, 848 (1996); A. V. Radyushkin and R. T. Ruskov, Nucl. Phys. **B481**, 625 (1996); I. V. Musatov and A. V. Radyushkin, Phys. Rev. D **56**, 2713 (1997).

[3] M. Diehl, T. Gousset, B. Pire, and O. Teryaev, Phys. Rev. Lett. **81**, 1782 (1998).

[4] V. N. Baier and A. G. Grozin, Sov. J. Nucl. Phys. **35**, 899 (1982).

[5] D. Diakonov and V. Petrov, Nucl. Phys. **B272**, 457 (1986); LNPI Report No. LNPI-1153, 1986, published (in Russian) in: Hadron matter under extreme conditions (Naukova Dumka, Kiev, 1986), p. 192.

[6] V. Yu. Petrov and P. V. Pobylitsa, hep-ph/9712203.

[7] V. Yu. Petrov, M. V. Polyakov, R. Ruskov, C. Weiss, and K. Goeke, Bochum Report No. RUB-TPII-8/98, hep-ph/9807229.

[8] CLEO Collaboration, J. Groenberg *et al.*, Phys. Rev. D **57**, 33 (1998).

[9] M. V. Polyakov, Bochum Report No. RUB-TPII-14/98, hep-ph/9809483; M. V. Polyakov and C. Weiss (in preparation).

[10] X. D. Ji, Phys. Rev. Lett. **78**, 610 (1997); Phys. Rev. D **55**, 7114 (1997); A. V. Radyushkin, Phys. Lett. B **285**, 333 (1996); **380**, 417 (1996).

[11] V. Yu. Petrov, P. V. Pobylitsa, M. V. Polyakov, I. Börnig, K. Goeke, and C. Weiss, Phys. Rev. D **57**, 4325 (1998).

[12] A. V. Radyushkin, Phys. Rev. D **59**, 014030 (1999).

[13] G. P. Lepage and S. J. Brodsky, Phys. Lett. **87B**, 359 (1979); A. V. Efremov and A. V. Radyushkin, *ibid.* **94B**, 245 (1980).

[14] For a review, see V. L. Chernyak and A. R. Zhitnitsky, Phys. Rep. **112**, 173 (1984).

[15] M. K. Chase, Nucl. Phys. **B174**, 109 (1980); V. N. Baier and A. G. Grozin, *ibid.* **B192**, 476 (1981).

[16] M. V. Polyakov and C. Weiss, Bochum Report No. RUB-TPII-1-99, hep-ph/9902451.

[17] S. J. Brodsky, L. Frankfurt, J. F. Gunion, A. H. Mueller, and M. Strikman, Phys. Rev. D **50**, 3134 (1994).