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Pion valence-quark parton distribution function

Lei Chang *, Anthony W. Thomas

CSSM, School of Chemistry and Physics, University of Adelaide, Adelaide SA 5005, Australia

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ABSTRACT

Within the Dyson–Schwinger equation formulation of QCD, a rainbow ladder truncation is used to calculate the pion valence-quark distribution function (PDF). The gap equation is renormalized at a typical hadronic scale, of order 0.5 GeV, which is also set as the default initial scale for the pion PDF. We implement a corrected leading-order expression for the PDF which ensures that the valence-quarks carry all of the pion's light-front momentum at the initial scale. The scaling behavior of the pion PDF at a typical partonic scale of order 5.2 GeV is found to be $(1-x)^\nu$, with $\nu \simeq 1.6$, as x approaches one.

Given its dual roles as a conventional bound-state in quantum field theory and as the Goldstone mode associated with dynamical chiral symmetry breaking, the pion has been proven critical to explaining phenomena as diverse as the long-range nucleon–nucleon interactions and the flavor asymmetry observed in the quark sea of the nucleon [1]. The study of the pion structure function is of great interest as a fundamental test of our understanding of nonperturbative QCD. Experimental information on the parton distribution function (PDF) in the pion has primarily been inferred from the Drell–Yan reaction in pion–nucleon collisions [2–5].

Lattice QCD calculations [6–8] have traditionally been able to yield only the low-order moments of the PDFs. While there has been a recent suggestion of a very promising way [9] to directly compute the x-dependence in lattice QCD, it will take considerable effort to reliably extract the large-x behavior using this method. The calculation of PDFs within models is challenging and various models have given a diversity of results. Most models, including the QCD parton model [10], pQCD [11] and the Dyson–Schwinger equations (DSE) [12,13] indicate that at high-x the PDF should behave as $(1-x)^\alpha$, with $\alpha \simeq 2$. The Nambu–Jona–Lasinio (NJL) models [14] with translationally invariant regularization and Drell–Yan–West relation [15] favors a linear dependence on $1-x$.

The first DSE study of the pion PDF was based [12] upon an analysis that employed phenomenological parametrizations of both the Bethe–Salpeter amplitude and the dressed-quark propagators. A numerical solution of the DSE utilizing the rainbow-ladder (RL) truncation has been used to compute the pion and kaon PDFs following same line [13].

In this work we revisit the pion valence PDF within the DSE approach, with the following improvements: 1) the rainbow-ladder gap equation is renormalized at a typical hadron scale, ζ_H, that also serves as the initial scale for the PDF; 2) a corrected leading-order expression for the PDF is employed within the RL trunction; 3) the extraction of the PDF is based on its moments, a method that has been widely used in parton distribution amplitude calculations [16,18,19]. The large-x behavior is naturally reflected in the high moments. In the method used here, we can calculate any large moment and thus we have a reliable tool with which to analyze the large-x behavior.

In order to help explain the numerical results and place them in some perspective, we introduce several models which produce pointwise PDFs. Our suggestions cover a broad range of possibilities, against which the predictions of the present model may be compared, especially calculations that can be described within the amplitude language, such as the DSE and NJL models with various regularization frameworks.

In Ref. [17] a corrected, leading-order expression was given for the pion's valence-quark PDF. This expression produces the model-independent result that quarks dressed via the RL truncation carry all of the pion's light-front momentum at a characteristic hadronic scale, if the meson amplitude is momentum dependent. We quote the form of the quark distribution function in the RL truncation here:

* Corresponding author.
E-mail address: l.chang@adelaide.edu.au (L. Chang).

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$$q(x) = N_c \mathrm{tr} \int_{dk}^{\Lambda} \delta(n \cdot k - xn \cdot P)\, \partial_k \left[\Gamma(k - \frac{P}{2}; -P) S(k) \right]$$
$$\Gamma(k - \frac{P}{2}; P) S(k - P) \,. \qquad (1)$$

In the infinite momentum frame, $q(x)$ is the number density for a single parton of flavor q to carry the momentum fraction $x = n \cdot k/n \cdot P$, which is positive definite over the physical region $0 < x < 1$. Here, n is a light-like four-vector, $n^2 = 0$; P is the pion's four-momentum, $P^2 = -m_\pi^2$, with m_π the pion mass; $\int_{dk}^{\Lambda}$ is a Poincaré-invariant regularization of the four-dimensional momentum integral (over k), with Λ the ultraviolet regularization mass-scale. In addition, S and Γ are the quark propagator and pion Bethe–Salpeter amplitude, respectively. In the present work the ultraviolet behavior of S and Γ is controlled by the one-gluon exchange interaction. In this case the above integral is ultraviolet divergence free and Λ can be set to infinity safely.

As the derivative in Eq. (1) acts on the full expression within the brackets it naturally yields two terms. The term related to the derivative of the quark propagator yields the so-called impulse-approximation. That corresponds to the textbook "handbag" contribution to virtual Compton scattering. The second term, arising from the action of the derivative on the amplitude originates in the initial/final state interactions. This expression is the minimal expression that retains the contribution to the quark distribution function from the gluons which bind dressed-quarks into the meson. This contribution may be thought of as a natural consequence of the nonlocal properties of the pion wave function. That is, it expresses the process where a photon is absorbed by a dressed quark, which then proceeds to become part of the pion bound-state before re-emitting the photon. It is easy to prove that the distribution function is symmetric, $q(x) = q(1 - x)$, under isospin symmetry and the valence quarks carry all of the momentum of the meson.

We describe pion as bound state using the Bethe–Salpeter equation. This takes the abbreviated form:

$$\Gamma_\pi(k; P) = \int_{dq}^{\Lambda} K(q, k; P) \chi_\pi(q; P) \qquad (2)$$

where q and k are the relative momenta between the quark–antiquark pair, P is the pion's four momentum and

$$\chi_\pi(q; P) = S(q_+) \Gamma_\pi(q; P) S(q_-) \qquad (3)$$

is the pion's Poincaré-covariant Bethe–Salpeter wave-function, with Γ_π the Bethe–Salpeter amplitude. Using isospin symmetry we label the dressed quark propagators $S(q_\pm)$, where $q_\pm = q \pm \frac{P}{2}$, without loss of generality. Explicitly, these take the form:

$$S^{-1}(k) = i\gamma \cdot k A(k^2) + B(k^2) \qquad (4)$$

where the scalar functions A, B depend on both momentum and the choice of renormalization point.

In this work, we perform the ladder truncation for the quark–antiquark scattering kernel, $K(q, k; P)$. This approximation has been widely used to compute the spectrum of meson bound states and related properties. In this framework the quark–gluon vertex is bare and a judicious choice of effective gluon propagator provides a connection between the infrared and ultraviolet scales. We use the interaction provided in Ref. [20], which contains two different parts. Its ultraviolet composition preserves the one-loop renormalization group behavior of QCD so that, as we shall see, the leading Bethe–Salpeter amplitude takes the well known, model independent ultraviolet behavior. The parameters of the infrared interaction, $D\omega$ and ω, manifest the strength and width of the interaction, respectively. It is chosen deliberately to be consistent with that determined in modern studies of the gauge sector of QCD.

The rainbow ladder truncation of the DSEs preserves the chiral symmetry of QCD. The renormalization constants for the wave function and mass function $Z_{2,4}(\zeta, \Lambda)$ must be included to regularize the logarithmic ultraviolet divergences. In the present calculation we follow the current quark mass independent renormalization approach introduced in Ref. [21]. In practice, the renormalization is defined by the conditions $A(k = \zeta) = 1$ and $\frac{\partial B(k=\zeta)}{\partial m_\zeta} = 1$ in the chiral limit. It should be noted that the renormalization point can be chosen in either the ultraviolet or infrared region and the quark mass function is independent of this choice. The renormalization constants $Z_{2,4}(\zeta, \Lambda)$ decrease as the scale decreases, reflecting the increase in the coupling strength in the infrared region. Here we choose $\zeta = 0.5$ GeV.

Before discussing any numerical results, it is interesting to recall some general features of the shape of the PDF and the Bethe–Salpeter amplitude for a meson. It has been shown that the significant features of $q(x)$, in Eq. (1), can be illustrated algebraically with some simple models. To take a close look at the relation between $q(x)$ and the pion Bethe–Salpeter amplitude, we consider an algebraic model where the quark propagator and the meson amplitude take of constituent mass form [16]

$$S^{-1}(k) = i\gamma \cdot k + M \qquad (5)$$

and

$$\Gamma_\pi(k; P) = i\gamma_5 \frac{12}{5} \frac{M}{f_\pi} \int_{-1}^{1} dz \rho(z) \frac{M^2}{k^2 + zk \cdot P + M^2} \,, \qquad (6)$$

where M is a dressed-quark mass, f_π is the pion decay constant and we focus on the case of a massless pion. The factor 12/5 is the normalization constant needed to ensure the charge conservation. k is the relative momentum of the quarks in the pion and we choose the amplitude to behave like $1/k^2$ asymptotically, as this is the leading order result if one takes a one-gluon exchange interaction between the quark and antiquark. We introduce the Nakanishi-like representation [22] with weight function $\rho(z)$, which takes a form different from that considered in Ref. [16]. We will see that different choices of ρ lead to different behaviors of $q(x)$. The present algebraic model makes it possible for us to determine the x-dependence of the PDF.

In Ref. [16] it has been shown that $\rho(z) = \frac{1}{2}(\delta(1-z) + \delta(1+z))$ describes a bound state with point-particle-like characteristics. It should be noted that it also gives a constant PDF, if one calculates the PDF exactly, even though the Bethe–Salpeter amplitude is momentum dependent. We infer that such behavior corresponds to the NJL prediction if one performs a Pauli–Villas regularization.

The QCD conformal limit can be reproduced with the weight function $\rho(z) = \frac{3}{4}(1 - z^2)$. Ref. [17] deduced a PDF which can be approximately expressed by $30x^2(1 - x)^2$. Following this line, we extend the model of spectral density as $\rho(z) = \frac{3}{4}(1 - z^2)\big(1 + 6a_2 C_2^{3/2}(z)\big)$, where a second Gegenbauer polynomial has been introduced and a_2 is a parameter. The corresponding PDA has the form $\varphi(x) = 6x(1 - x)\big(1 + a_2 C_2^{3/2}(2x - 1)\big)$. Obviously this form reproduces the Chernyak–Zhitnitsky (CZ) form, with $a_2 = 2/3$ [23]. Following the method in Ref. [17], the PDF related to the CZ PDA can be computed consistently. The result is depicted in Fig. 1. Near $x = 1$ this model has the power-law behavior, $\beta(1 - x)^2$, predicted by the QCD parton model. Here a_2 only affects the coefficient β, not the power. However, the PDF shows oscillatory behavior that

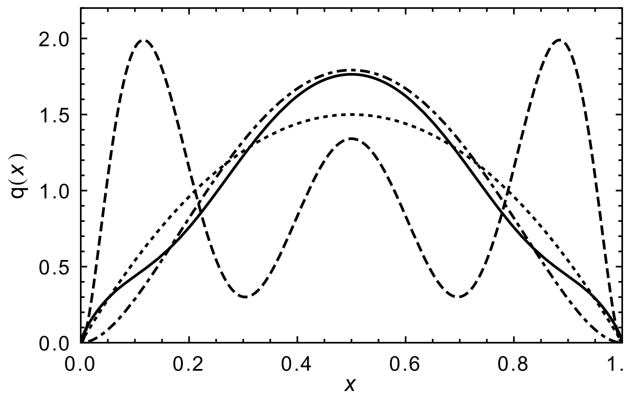

Fig. 1. Parton distribution function at ζ_H. Curves: *solid*, rainbow-ladder computation herein corresponding to Eq. (11); *dashed*, $\rho(z)=\frac{3}{4}(1-z^2)\left(1+6a_2C_2^{3/2}(z)\right)$ with $a_2=2/3$ (*dotdashed* $a_2=0$); *dotted*, $\rho(z)=\frac{1}{\pi}(1-z^2)^{-\frac{1}{2}}$.

is difficult to reconcile with the physical meaning of the parton distribution function. Of course, it is known that the CZ-like PDA is also double-humped and it can be argued that this is possible because it only has the interpretation of an amplitude. In our consistent calculation we have show that the PDF is triple-humped and so we treat it with some caution.

In earlier work on the pion PDA we found that it is concave and broader than the asymptotic form at a typical hadronic scale. To capture this characteristic we suggest another model for the weight function, $\rho(z)=\frac{1}{\pi}(1-z^2)^{-\frac{1}{2}}$, which is divergent at $z=\pm 1$ but nevertheless integrable, $\int_{-1}^{1}\rho(z)=1$. The corresponding PDA and PDF are $\varphi(x)=\frac{8}{\pi}\sqrt{x(1-x)}$ and $6x(1-x)$, respectively. This model for the PDA follows from the precise mapping of string amplitudes in Anti-deSitter space to the light-front wavefunction of a hadron in physical space-time using holographic methods [24]. It should be emphasized that the related PDF has the power-law behavior, $1-x$, near $x=1$. That is in contrast to the QCD parton prediction.

Although the ultraviolet k^2 dependence is the same for the three different models described earlier, the parton distribution function is very different. Ezawa [25] predicted that the pion PDF would behave as $(1-x)^{2\alpha}$ if the pion amplitude behaved as $\frac{1}{(k_1^2)^{\alpha}}$, where k_1 is the struck quark momentum.

Let us write a general form of the amplitude as $\int_{-1}^{1}dz(1-z^2)^{\nu}\frac{1}{(k^2+zk\cdot P+M^2)^{\beta}}$ with the relative momentum k. If we fix one quark momentum and set the other to infinity we will find that the leading order amplitude is $\frac{1}{(k_1^2)^{1+\nu}}$ for $-1<\nu<0$, whatever the value of β. Based on Ezawa's work one readily finds a PDF which behaves as $(1-x)^{2(1+\nu)}$, which is consistent with our results. If one works with a free gluon propagator that produces an amplitude with $\nu=1$; $\beta=1$, then this yields the well known large-x behavior, $(1-x)^2$. At the present time we can only find a model independent formula if the relative momentum tends to infinity. However, the infrared interaction does effect the asymptotic form if one quark momentum goes to infinity with the other fixed.

With the normalization constant set by hand, we can input the light quark mass at $\zeta=0.5$ GeV and arrange that it yields the correct pion mass. With an input current quark mass of 18.6 MeV we obtain $m_\pi=0.14$ GeV and $f_\pi=0.092$ GeV. In the present work, the computation of the moments of the PDF is relatively straightforward because we employ algebraic parameterizations of the of quark propagator and the Bethe–Salpeter amplitude. The dressed-quark propagators are represented as [26]

Table 1
Fit parameters for the pseudoscalar meson Bethe–Salpeter amplitudes.

	c^{i}	c^{u}	ν^{i}	ν^{u}	a	Λ^{i}	Λ^{u}
E_π	5.8577	0.195	−0.656	1.08	2.682	1.247	1
F_π	2.9870	0.021	1.82	1.08	2.655	1.027	1

$$S(p)=\sum_{j=1}^{j_m}\left[\frac{z_j}{i\gamma\cdot p+m_j}+\frac{z^*}{i\gamma\cdot p+m_j^*}\right], \tag{7}$$

with $\Im m_j\neq 0\ \forall j$, so that $\sigma_{V,S}$ are meromorphic functions with no poles on the real p^2-axis, a feature consistent with confinement [27]. The pseudoscalar Bethe–Salpeter amplitude has the form

$$\Gamma(q;P)=\gamma_5\big[iE(q;P)+\gamma\cdot PF(q;P)+\gamma\cdot q\,G(q;P)+\sigma_{\mu\nu}q_\mu P_\nu H(q;P)\big]. \tag{8}$$

We retain all four terms in the pseudoscalar meson Bethe–Salpeter amplitude in the numerical calculation of the BS equation. In the computation of the PDF we find that the first two terms in the BS amplitude dominate the parton structure. For this reason, for the main part of the present work we retain only the first two terms and leave the full calculation for future work. We fit the associated scalar functions E and F via

$$\mathcal{F}(q;P)=\mathcal{F}^{\mathrm{i}}(q;P)+\mathcal{F}^{\mathrm{u}}(q;P)\,, \tag{9a}$$

$$\mathcal{F}^{\mathrm{i}}(q;P)=c^{\mathrm{i}}\int_{-1}^{1}dz\,\rho_{\nu^{\mathrm{i}}}(z)\left[a\hat{\Delta}^4_{\Lambda^{\mathrm{i}}}(q_z^2)+(1-a)\hat{\Delta}^5_{\Lambda^{\mathrm{i}}}(q_z^2)\right],$$

$$\mathcal{F}^{\mathrm{u}}(q;P)=c^{\mathrm{u}}\int_{-1}^{1}dz\,\rho_{\nu^{\mathrm{u}}}(z)\hat{\Delta}_{\Lambda^{\mathrm{u}}}(q_z^2)\,,$$

where $\rho_\nu(z)=\frac{\Gamma(\nu+\frac{3}{2})}{\sqrt{\pi}\Gamma(\nu+1)}(1-z^2)^\nu$ and $\Delta_\Lambda(q_z^2)=\frac{\Lambda^2}{q^2+zq\cdot P+\Lambda^2}$. This choice of denominator makes the fit easier when the meson mass is not zero. The resulting parameter values are listed in Table 1.

In the actual calculations we use a small power to respect the possible anomalous dimension [28] in the BS amplitude and the leading order behavior of the amplitude in the ultraviolet region is $1/q^{2+2\alpha}$, which is the natural outcome of using a one-gluon exchange interaction. The leading-order weight function, $\rho(z)$, in the ultraviolet region can be obtained exactly if the free gluon propagator is considered [10]. We choose the simple form to respect the infrared behavior. The use of denominator powers 4 and 5 to fit the infrared behavior of the pion's E amplitude is consistent with the chiral condition which makes it identical with chiral propagator $B(p^2)$ function. In particular, this presentation preserves the infrared inflection point which is an indicator of quark confinement [29]. For the light quark meson the infrared power $\nu^{\mathrm{i}}<0$ in the weight function, yields an integrable singularity at $z=\pm 1$. We have shown that this is true for ρ and K mesons.

In practice we usually fit ν^{i} using the limit number of Chebyshev polynomials expansion of amplitude [30]. In another alternative approach one can get the best fit by using a CZ-like weight function with a positive value of a_2 if one only has limited knowledge of $q\cdot P$ dependence of amplitude $\Gamma_\pi(q;P)$. We have shown that this choice would produce an unlikely humped parton distribution that is not well understood. An improvement might be considered with the complete picture of amplitude and then the ambiguity in the weight function would be controlled.

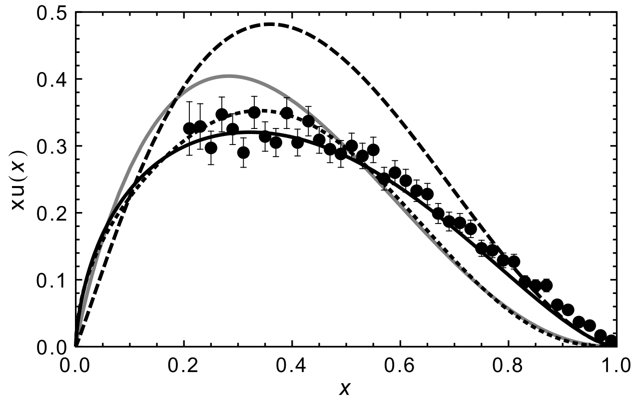

Fig. 2. Comparison of the pion parton distribution function calculated using the rainbow-ladder truncation computation described herein at $\zeta_5 = 5.2$ GeV with experimental data extracted from the π N Drell–Yan reaction [4]. Curves: *dashed*, result obtained with $2\langle x\rangle = 1$, illustrating the dressed quarks carry all of the pion momentum and *solid*, result obtained with $2\langle x\rangle = 0.75$, illustrating the effect of shifting 25% of the dressed quark's momentum into sea and glue; *dotted*, next-leading order evolution of Fit-3 in Table-I in [31] by considering soft gluon resummation; *gray solid*, result obtained with $2\langle x\rangle = 1$ at $\zeta_H = 0.36$ GeV.

Based on the method we developed for calculating the PDA in Ref. [16], it is straightforward to compute the PDF with the quark propagator and meson amplitude in hand. The first step is to compute the moments $\langle x^m\rangle = \int_0^1 dx x^m q(x)$. The algebraic form of the input makes it possible to compute arbitrarily many moments. Considering the fact that PDF is an even function under $x \leftrightarrow (1-x)$ and vanishes at the endpoints unless the underlying interaction is momentum-independent, we can reconstruct the PDF by expanding in Gegenbauer polynomials of order α, that are a complete set with respect to the measure $(x(1-x))^{\alpha-1/2}$. Therefore, with complete generality, the PDF for π may accurately be approximated as follows:

$$q(x) \approx q_m(x) = N_\alpha [x\bar{x}]^{\alpha-1/2}\left[1 + \sum_{j=2,4,\ldots}^{j_{\max}} a_j^\alpha C_j^\alpha(2x-1)\right], \qquad (10)$$

where $\bar{x} = 1-x$, $N_\alpha = \Gamma(2\alpha+1)/[\Gamma(\alpha+1/2)]^2$. The parameters α, a_j can be fitted by the moments. In practice, this procedure converged very rapidly: $j_m = 8$ was sufficient for the pion PDF. Our results for the PDFs are depicted in Fig. 1 with the functions defined in Eqs. (10) and

$$\begin{array}{cccccc} & \alpha & a_2 & a_4 & a_6 & a_8 \\ \pi & 1.158 & -0.175 & 0.1 & -0.019 & -0.015 \end{array}. \qquad (11)$$

The parton distribution function at $\zeta = 0.5$ GeV cannot be simply expressed by a one parameter representation like $x^\alpha(1-x)^\alpha$. There is a point of inflection around $x = 0.85$, which can be thought of as the transition from soft to hard scales. The PDF behaves as $(1-x)^\nu$ for $x > 0.85$, with $\nu \simeq 2(1+\nu_E^i)$, consistent with our model analysis. The F part of amplitude exhibits an infrared weight function with positive power that would contribute to the parton distribution as higher-twists. Including F does not effect the region $x > 0.85$ but does make the PDF more broad in the infrared region.

In Fig. 2 we show the result of evolving the PDF, using the leading order DGLAP equations, from $\zeta_H \to \zeta_5$ with $\zeta_5 = 5.2$ GeV. For the dashed line we suppose the valence quarks carry all the pion momentum at ζ_H which produce the momentum fraction carried by valence quarks as $2\langle x\rangle = 0.55$ at ζ_5. The rainbow-ladder truncation just describes pion solely from a dressed-quark and dressed-antiquark which include no mechanism that can shift momentum from the dressed-quarks into sea-quarks and gluons. A phenomenological method has been introduced in Ref. [17] to address this effect. Following this perspective we shift 25% dressed quark momentum to sea and glue at ζ_H as the solid line show. Including sea and glue suppress the momentum distribution at mediate region of x. The both numerical results favor a power-law in the valence region of the form $(1-x)^{1.6}$. The experimental data has been revisited due to soft gluon resummation [31] (dotted line in Fig. 2) which suggest an approximate power-law form $(1-x)^\nu$ with $\nu > 2$. In any model calculation the value of ν at ζ_5 depends on the initial scale ζ_H. As shown by the gray line in Fig. 2 we plot the momentum distribution evolved from $\zeta_H = 0.36$ GeV and suppose the quarks carry all the momentum at this scale. The momentum distribution matches the dotted line well at $x > 0.5$.

To summarize, we have presented the pion PDF within a rainbow-ladder truncation of the DSE approach. By employing a Nakanishi representation of Bethe–Salpeter amplitude [22] and calculating the moments of the PDF to arbitrarily large values we have been able to calculate the x-dependence of the PDF at a typical hadronic scale. We analyze the relation between the power-law behavior and the infrared interaction that binds the meson. The present DSEs favor a power $(1-x)^{1.6}$, at a typical experimental scale, $\zeta = 5.2$ GeV after leading order QCD evolution.

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References

[1] A.W. Thomas, Phys. Lett. B 126 (1983) 97.
[2] J. Badier, et al., NA3 Collaboration, Z. Phys. C 18 (1983) 281.
[3] B. Betev, et al., NA10 Collaboration, Z. Phys. C 28 (1985) 15.
[4] J.S. Conway, et al., Phys. Rev. D 39 (1989) 92.
[5] K. Wijesooriya, P.E. Reimer, R.J. Holt, Phys. Rev. C 72 (2005) 065203.
[6] D. Brommel, et al., in: PoS, LAT2007, 2007, p. 140.
[7] C. Best, et al., Phys. Rev. D 56 (1997) 2743.
[8] W. Detmold, W. Melnitchouk, A.W. Thomas, Phys. Rev. D 68 (2003) 034025.
[9] Xiangdong Ji, Phys. Rev. Lett. 110 (2013) 262002.
[10] G.R. Farrar, D.R. Jackson, Phys. Rev. Lett. 35 (1975) 1416.
[11] Xi. Ji, J. Ma, F. Yuan, Phys. Lett. B 610 (2005) 247;
S.J. Brodsky, M. Burkardt, I. Schmidt, Nucl. Phys. B 441 (1995) 197.
[12] M.B. Hecht, C.D. Roberts, S.M. Schmidt, Phys. Rev. C 63 (2001) 025213.
[13] T. Nguyen, A. Bashir, C.D. Roberts, P.C. Tandy, Phys. Rev. C 83 (2011) 062201(R).
[14] W. Bentz, T. Hama, T. Matsuki, K. Yazaki, Nucl. Phys. A 651 (1999) 143;
E. Ruiz Arriola, W. Broniowski, Phys. Rev. D 66 (2002) 094016.
[15] S.D. Drell, T.-M. Yan, Phys. Rev. Lett. 24 (1970) 181;
G.B. West, Phys. Rev. Lett. 24 (1970) 1206.
[16] L. Chang, et al., Phys. Rev. Lett. 110 (2013) 132001.
[17] L. Chang, et al., Phys. Lett. B 737 (2014) 23.
[18] Chao Shi, et al., arXiv:1406.3353.
[19] Fei Gao, et al., Phys. Rev. D 90 (2014) 014011.
[20] Si-Xue Qin, L. Chang, Y. -x Liu, C.D. Roberts, D.J. Wilson, Phys. Rev. C 84 (2011) 042202.
[21] S. Weinberg, Phys. Rev. D 8 (1973) 3497;
H. Pagels, Phys. Rev. D 19 (1979) 3080.
[22] N. Nakanishi, Phys. Rev. 130 (1963) 1230.
[23] V.L. Chernyak, A.R. Zhitnitsky, Phys. Rep. 112 (1984) 173.
[24] S.J. Brodsky, G.F. de Teramond, Phys. Rev. Lett. 96 (2006) 201601;
S.J. Brodsky, G.F. de Teramond, Phys. Rev. D 77 (2008) 056007.
[25] Z.F. Ezawa, Il Nuovo Cimento 23 (1974) 271.
[26] M. Bhagwat, M. Pichowsky, P.C. Tandy, Phys. Rev. D 67 (2003) 054019.
[27] V.N. Gribov, Eur. Phys. J. C 10 (1999) 91–105;
G. Krein, C.D. Roberts, A.G. Williams, Int. J. Mod. Phys. A 7 (1992) 5607–5624.
[28] L. Chang, et al., Phys. Rev. Lett. 111 (2013) 141802.
[29] K. Langfeld, H. Reinhardt, J. Gattnar, Nucl. Phys. B 621 (2002) 131.
[30] P. Jain, H.J. Munczek, Phys. Rev. D 48 (1993) 5403.
[31] M. Aicher, A. Schafer, W. Vogelsang, Phys. Rev. Lett. 105 (2010) 252003.