MAGNETIC MOMENT OF THE PENTAQUARK $\Theta^+(1540)$ WITH LIGHT-CONE QCD SUM RULES

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Abstract

In this article, we study the magnetic moment of the pentaquark state $\Theta^+(1540)$ as diquark-diquark-antiquark ($[ud][ud]\bar{s}$) state in the framework of the light-cone QCD sum rules approach. The numerical results indicate the magnetic moment of the pentaquark state $\Theta^+(1540)$ is about $\mu_{\Theta^+} = -(0.49 \pm 0.06)\mu_N$.

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Key Words: Light-cone QCD Sum Rules, Magnetic moment, Pentaquark

1 Introduction

The observation of the new baryon $\Theta^+(1540)$ with positive strangeness and minimal quark content $udud\bar{s}$ [1] has motivated intense theoretical investigations to clarify its quantum numbers and to understand the under-structures [3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 15, 16]. Through the pentaquark state $\Theta^+(1540)$ can be signed to the top of the antidecuplet $\bar{10}$ with isospin $I = 0$, the spin and parity have not been experimentally determined yet and no consensus has ever been reached on the theoretical side [2, 3, 4, 5, 6, 7, 12, 13, 14, 17, 18, 19, 20, 21, 22, 23]. The discovery has opened a new field of strong interaction and provides a new opportunity for a deeper understanding of low energy QCD especially when multiquarks are involved. The magnetic moments of the pentaquark states are fundamental parameters as their masses, which have copious information about the underlying quark structures i.e. different substructures result in different predictions, can be used to distinguish the preferred configurations from various theoretical models and deepen our understanding of the underlying dynamics. Furthermore, the magnetic moment of the pentaquark state $\Theta^+(1540)$ is an important ingredient in studying the cross sections of the photo- or electro-production, which can be used to determine the fundamental quantum number of the pentaquark state $\Theta^+(1540)$, such as spin and parity [24, 25], and may be extracted from experiments eventually in the near future.

There have been several works on the magnetic moments of the pentaquark state $\Theta^+(1540)$ [25, 26, 27, 28, 29, 30, 31, 32, 33, 34, 35, 36], in this article, we take the point of view that the $\Theta^+(1540)$ baryon is diquark-diquark-antiquark ($[ud][ud]\bar{s}$)

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state with the quantum numbers $J = \frac{1}{2}, I = 0, S = +1$, and study its magnetic moment in the framework of the light-cone QCD sum rules approach, which carries out the operator product expansion near the light-cone $x^2 \approx 0$ instead of the short distance $x \approx 0$ while the nonperturbative matrix elements are parameterized by the light-cone distribution amplitudes which classified according to their twists instead of the vacuum condensates \[37\] \[38\].

The article is arranged as follows: we derive the light-cone QCD sum rules for the magnetic moment of the pentaquark state $\Theta^+(1540)$ in section II; in section III, numerical results; section IV is reserved for conclusion.

2 Light-cone QCD Sum Rules for the Magnetic Moment

We can study the magnetic moments of the baryons using two-point correlation functions in an external electromagnetic field, with vacuum susceptibilities introduced as parameters for nonperturbative propagation in the external field, i.e. the QCD sum rules in the external field \[40\] \[41\]. As nonperturbative vacuum properties, the susceptibilities can be introduced for both small and large momentum transfers in the external fields. The alternative way is the light-cone QCD sum rules approach, which was firstly used to calculate the magnetic moments of nucleons in Ref.\[38\]. For more discussions about the magnetic moments of the baryons in the framework of the light-cone QCD sum rules approach, one can consult Ref. \[42\].

In the following, we write down the two-point correlation function $\Pi_{\eta}(p)$ in the framework of the light-cone QCD sum rules approach \[6\] \[29\].

\[
\Pi_{\eta}(p, q) = i \int d^4x \, e^{ipx} \langle \gamma(q)|T\{\eta(x)\bar{\eta}(0)\}|0\rangle,
\]

\[
\eta(x) = \{t\eta_1(x) + \eta_2(x)\},
\]

\[
\eta_1(x) = \frac{1}{\sqrt{2}} \epsilon^{abc} \{[u^T_{a}(x)C\gamma_5d_b(x)] [u^T_{c}(x)C\gamma_5d_e(x)] C\bar{s}^T_{e}(x) - (u \leftrightarrow d)\},
\]

\[
\eta_2(x) = \frac{1}{\sqrt{2}} \epsilon^{abc} \{[u^T_{a}(x)Cd_b(x)] [u^T_{c}(x)Cd_e(x)] C\bar{s}^T_{e}(x) - (u \leftrightarrow d)\}.
\]

Here the $\gamma(q)$ represents the external electromagnetic field with the vector potential $A_\mu(x) = \varepsilon_\mu e^{iq \cdot x}$, the $\varepsilon_\mu$ is the photon polarization vector and the field strength $F_{\mu\nu}(x) = i(\varepsilon_\nu q_\mu - \varepsilon_\mu q_\nu) e^{iq \cdot x}$. The $a, b, c$ and $e$ are color indexes, the $C = -C^T$ is the charge conjugation operator, and the $t$ is an arbitrary parameter. The constituents $\epsilon^{abc} u^T_{a}(x)C\gamma_5d_b(x)$ and $\epsilon^{abc} u^T_{b}(x)Cd_e(x)$ represent the scalar $J^P = 0^+$ and the pseudoscalar $J^P = 0^-$ $ud$ diquarks respectively. They both belong to the antitriplet $3$ representation of the color $SU(3)$ group and can cluster together with diquark-diquark-antiquark type structure to give the correct spin and parity for the pentaquark $\Theta^+(1540) J^P = \frac{1}{2}^+$. The scalar diquarks correspond to the $^1S_0$ states.
of \( ud \) quark system. The one gluon exchange force and the instanton induced force can lead to significant attractions between the quarks in the \( 0^+ \) channels. The pseudoscalar diquarks do not have nonrelativistic limit, can be taken as the \( ^3P_0 \) states.

According to the basic assumption of current-hadron duality in the QCD sum rules approach, we insert a complete series of intermediate states satisfying the unitarity principle with the same quantum numbers as the current operator \( \eta(x) \) into the correlation function in Eq.(1) to obtain the hadronic representation. After isolating the pole terms of the lowest pentaquark states, we get the following result,

\[
\Pi_\eta(p,q) = -f_0^2 \epsilon^\mu \frac{\not{p} + m_{\Theta^+}}{p^2 - m_{\Theta^+}^2} \left[ F_1(q^2) \gamma_\mu + \frac{i \sigma_{\mu\nu} q^\nu}{2m_{\Theta^+}} F_2(q^2) \right] \frac{\not{p} + \not{q} + m_{\Theta^+}}{(p + q)^2 - m_{\Theta^+}^2} + \cdots
\]

\[
= -\frac{f_0^2}{(p^2 - m_{\Theta^+}^2)((p + q)^2 - m_{\Theta^+}^2)} \not{p} \not{q} + \cdots
\]

\[
= \Pi(p,q) i \epsilon_{\mu\nu\alpha\beta} \gamma^\mu \gamma^\nu \epsilon^{\alpha} p^\beta + \cdots, \tag{3}
\]

where

\[
\langle 0 | \eta(0) | \Theta^+(p) \rangle = f_0 u(p), \tag{4}
\]

and the \( F_1(q^2) \) and \( F_2(q^2) \) are electromagnetic form factors. Here we write down only the double-pole term which corresponds to the magnetic moment of the pentaquark state \( \Theta^+(1540) \) explicitly, and choose the tensor structure \( \epsilon^\mu \epsilon^\nu \epsilon^\alpha p^\beta \). The contributions concerning the excited and continuum states are suppressed after double Borel transformation, and not shown explicitly for simplicity. From the electromagnetic form factors \( F_1(q^2) \) and \( F_2(q^2) \), we can obtain the magnetic moment of the pentaquark state \( \Theta^+(1540) \),

\[
\mu_{\Theta^+} = \left\{ F_1(0) + F_2(0) \right\} \frac{e_{\Theta^+}}{2m_{\Theta^+}}. \tag{5}
\]

The calculation of the operator product expansion near the light-cone \( x^2 \approx 0 \) at the level of quark and gluon degrees of freedom is straightforward and tedious, here technical details are neglected for simplicity. The photon can couple to the quark lines perturbatively and nonperturbatively, which results in two classes of diagrams. In the first class of diagrams, the photon couples to the quark lines perturbatively through the standard QED,

\[
\langle 0 | T \left\{ q^a(x) q^b(0) \right\} | 0 \rangle_{F_{\mu\nu}} = \frac{\delta^{ab} e_q}{16\pi^2 x^2} \int_0^1 du \left\{ 2(1 - 2u)x_\mu \gamma_\nu + i e_{\mu\rho\sigma} \gamma_5 \gamma_\rho x^\sigma \right\} F^{\mu\nu}(ux). \tag{6}
\]

The second class of diagrams involve the nonperturbative interactions of photons with the quark lines, which are parameterized by the photon light-cone distribution amplitudes in stead of the vacuum susceptibilities. In this article, the following
two-particle photon light-cone distribution amplitudes are useful \[\text{II}\],

\[
\langle \gamma(q)|\bar{q}(x)\sigma_{\alpha\beta}q(0)|0\rangle = i e_q \langle \bar{q} q \rangle \int_0^1 du e^{iux} \left\{ (\varepsilon_{\alpha} q_{\beta} - \varepsilon_{\beta} q_{\alpha}) \chi \varphi(u) + x^2 [(g_1(u) - g_2(u))] + \{ q x (\varepsilon_{\alpha} x_{\beta} - \varepsilon_{\beta} x_{\alpha}) + \varepsilon x (\alpha_{\alpha} x_{\beta} - \alpha_{\beta} x_{\alpha}) \} g_2(u) \right\},
\]

(7)

\[
\langle \gamma(q)|\bar{q}(x)\gamma_5 q(0)|0\rangle = f^e_q \varepsilon_{\mu\rho\sigma} \varepsilon^{\nu} q^\rho x^\sigma \int_0^1 du e^{iux} \psi(u),
\]

(8)

\[
\langle \gamma(q)|\bar{q}(x)\gamma_\mu q(0)|0\rangle = f^{(V)} e_q \varepsilon_{\mu} \int_0^1 du e^{iux} \psi^{(V)} (u),
\]

(9)

where the \(\chi\) is the magnetic susceptibility of the quark condensate, its values with different theoretical approaches are different from each other, for a short review, one can see Ref.\[46\]. The \(e_q\) is the quark charge, \(\varphi(u)\) and \(\psi(u)\) are the twist–2 photon light-cone distribution amplitudes, while \(g_1(u)\) and \(g_2(u)\) are the twist–4 light-cone distribution amplitudes. The twist–3 photon light-cone distribution amplitudes are neglected due to their small contributions.

Once the analytical results are obtained, then we can express the correlation functions at the level of quark-gluon degrees of freedom into the following form through dispersion relation,

\[
\Pi(p, q) = e_s \int_0^1 du \left\{ \frac{1}{\pi} \int_0^{s_0} ds \frac{\text{Im}[A(-s)]}{s - up^2 - (1 - u)(p + q)^2} + B(p, q) \right\} + \cdots,
\]

(10)

where

\[
\frac{\text{Im}[A(-s)]}{\pi} = \frac{(5t^2 + 2t + 5)s^4}{2^{12}5!4!\pi^8} - \frac{(5t^2 + 2t + 5)s^3 f \psi(u)}{2^{13}5!\pi^6} + \frac{(7t^2 - 2t - 5)s \langle \bar{q} q \rangle^2}{2^{7}3^2 \pi^4} - \frac{(7t^2 - 2t - 5) \langle \bar{q} q \rangle^2 f \psi(u)}{2^{6}3^2 \pi^4} + \frac{(5t^2 + 2t + 5)s \langle \bar{q} q \rangle^4}{2^{11}3^4 \pi},
\]

\[
B(p, q) = -\frac{(5t^2 + 2t + 5) \langle \bar{q} q \rangle^4}{2^{3}3^3 (-up^2 - (1 - u)(p + q)^2)^2}.
\]

From Eq.(10), we can see that due to the special interpolating current (Eq.(2)), only the \(s\) quark has contributions to the magnetic moment of the pentaquark \(\Theta^+(1540)\), which is significantly different from the results obtained in Refs.\[29, 35\], where all the \(u\), \(d\) and \(s\) quarks have contributions. In Refs.\[29, 35\], the interpolating current \(J(x)\) is used,

\[
J(x) = \frac{1}{\sqrt{2}} \langle \sum_{a} u^T_a(x) C \gamma_5 d_b(x) \rangle \{ u_e(x) \bar{s}_e(x) i \gamma_5 d_e(x) - d_e(x) \bar{s}_e(x) i \gamma_5 u_e(x) \}.
\]

(11)

Then we make double Borel transformation with respect to the variables \(p^2\) and \((p + q)^2\) in Eq.(10),

\[
\mathcal{B}^{M^2}_{(p + q)^2} \mathcal{B}^{M^2}_p \frac{\Gamma(n)}{[m^2 - (1 - u)(p + q)^2 - up^2]^n} = \langle M^2 \rangle^{2-n} e^{-\frac{m^2}{M^2}} \delta(u - u_0),
\]

(12)
here $M^2 = \frac{M_1^2 M_2^2}{M_1^2 + M_2^2}$ is the Borel parameter and $u_0 \equiv \frac{M_1^2}{M_1^2 + M_2^2}, 1 - u_0 \equiv \frac{M_2^2}{M_1^2 + M_2^2}$. In this way the single-pole terms can be eliminated. For more discussions about double Borel transformation, one can consult Ref. [47]. Finally we obtain the sum rule,

$$\{F_1(0) + F_2(0)\} f_0^2 e^{-\frac{m_0^2}{M^2}} = -e_s AA,$$

where $AA = \frac{(5t^2 + 2t + 5)M^{12}E_4}{2^{12}5!\pi^8} - \frac{(5t^2 + 2t + 5)M^{10}E_3 f \psi(u_0)}{2^{15}5!\pi^6} + \frac{\langle \frac{d^2E}{dt^2} \rangle \langle \frac{u}{m} \rangle}{2^{17}3!\pi^4} - \frac{m}{2^{14}4!\pi^6} \langle \frac{\alpha_s GG}{\pi} \rangle - \frac{(5t^2 + 2t + 5)M^6 E_1}{2^{14}3!\pi^4} \langle \frac{\alpha_s GG}{\pi} \rangle - \frac{(5t^2 + 2t + 5)\langle \bar{q}q \rangle^4}{2^{3}3!}.$$

$E_n = 1 - \exp \left( -\frac{s_0}{M^2} \right) \sum_{k=0}^{n} \left( \frac{s_0}{M^2} \right)^k \frac{1}{k!},$

here we have used the functions $E_n$ to subtract contributions come from the excited states and continuum states. As the initial and final states are the same pentaquark state $\Theta^+(1540)$. It is natural to take the values $M_1^2 = M_2^2 = 2M^2$ and $u_0 = \frac{1}{2}$.

If we replace the final $\gamma(q)$ state with the vacuum state in Eq.(1), we can obtain the sum rules for the coupling constant $f_0$ [5],

$$f_0^2 e^{-\frac{m_0^2}{M^2}} = BB,$$

$$BB = \frac{3(5t^2 + 2t + 5)M^{12}E_5}{2^{11}7!\pi^8} + \frac{(5t^2 + 2t + 5)m_s \langle \bar{s}s \rangle M^8 E_3}{2^{10}5!\pi^6} + \frac{(1 - t)^2 M^8 E_3}{2^{13}5!\pi^6} \langle \frac{\alpha_s GG}{\pi} \rangle + \frac{(7t^2 - 2t - 5)\langle \bar{q}q \rangle^2 M^6 E_2}{2^{9}3!\pi^4} - \frac{(5t^2 + 2t + 5)m_s \langle \bar{s}g_s \cdot Gs \rangle M^6 E_2}{2^{14}3!\pi^6} + \frac{(7t^2 - 2t - 5)m_s \langle \bar{s}s \rangle \langle \bar{q}q \rangle^2 M^2 E_0}{2^{6}3!\pi^2} + \frac{(5t^2 + 2t + 5)\langle \bar{q}q \rangle^4}{6^3}.$$

From above equations, we can obtain

$$\{F_1(0) + F_2(0)\} = -e_s \frac{AA}{BB}.$$

### 3 Numerical Results

In this article, we take the values of the parameter $t$ for the interpolating current in Eq.(2) to be $t = -1$, which can give stable mass i.e. $m_{\Theta^+} \approx 1540$ MeV with respect to the variations of the Borel mass $M^2$ in the considered interval $M^2 =$
The other parameters are taken as \( \langle \bar{s}s \rangle = 0.8 \langle \bar{u}u \rangle, \langle \bar{d}d \rangle = (\text{-}219\text{MeV})^2, \langle \bar{s}g_s \cdot Gs \rangle = 0.8 \langle \bar{s}s \rangle, \langle \bar{s}G \rangle = (0.33\text{GeV})^4, m_u = m_d = 0, m_s = 150\text{MeV}, \psi(u) = 1 \text{ and } f = 0.028\text{GeV}, \text{[44-45].} \) Here we have neglected the uncertainties about the vacuum condensates, small variations of those condensates will not lead to larger changes about the numerical values. The threshold parameter \( s_0 \) is chosen to vary between \((3.6 - 3.8)\text{GeV}^2\) to avoid possible contaminations from higher resonances and continuum states. For \( s_0 = (3.6 - 3.8)\text{GeV}^2 \), we obtain the values

\[
F_1(0) + F_2(0) = -(0.80 \pm 0.10), \\
\mu_{\Theta^+} = -{(0.80 \pm 0.10)e_{\Theta^+}}{2m_{\Theta^+}}, \\
= -(0.49 \pm 0.06)\mu_N,
\]

where the \( \mu_N \) is the nucleon magneton. From the Table 1, we can see that although the numerical values for the magnetic moment of the pentaquark state \( \Theta^+(1540) \) vary with theoretical approaches, they are small in general; our numerical results are consistent with most of the existing values of theoretical estimations in magnitude, however, with negative sign. In previously work \([35, 36]\), we observe that the sum rules in the external electromagnetic field with different interpolating currents can lead to very different predictions for the magnetic moment, although they can both give satisfactory masses for the pentaquark state \( \Theta^+(1540) \). The present results show that the same interpolating current with different approaches i.e. the QCD sum rules in the external fields and the light-cone QCD sum rules, can result in different predictions in magnitude with the same sign \([29, 35, 36]\). When the experimental measurement of the magnetic moment of the pentaquark state \( \Theta^+(1540) \) is possible in the near future, we might be able to test the theoretical predictions of the magnetic moment and select the preferred quark configurations. In the region \( M^2 = (2 - 3)\text{GeV}^2 \), the sum rules for \( F_1(0) + F_2(0) \) as functions of the Borel parameter \( M^2 \) is plotted in Fig.1 for \( s_0 = 3.7\text{GeV}^2 \) as an example.

4 Conclusion

In summary, we have calculated the magnetic moment of the pentaquark state \( \Theta^+(1540) \) as diquark-diquark-antiquark \((\bar{u}d)[\bar{u}d][\bar{s}]\) state in the framework of the light-cone QCD sum rules approach. The numerical results are consistent with most of the existing values of theoretical estimations in magnitude, however, with negative sign, \( \mu_{\Theta^+} = -(0.80 \pm 0.10)e_{\Theta^+}{2m_{\Theta^+}} = -(0.49 \pm 0.06)\mu_N. \) In previously work \([35, 36]\), we observe that the sum rules in the external electromagnetic field with different interpolating currents can lead to very different predictions for the magnetic moment, although they can both give satisfactory masses for the pentaquark state \( \Theta^+(1540) \). The present results show that the same interpolating current with
Table 1: The values of $\mu_{\Theta^+}$ (in unit of $\mu_N$)

| Reference | $\mu_{\Theta^+}$ ($\mu_N$) |
|-----------|--------------------------|
| 25        | 0.12 ± 0.06              |
| 26        | 0.08 ~ 0.6               |
| 27        | 0.2~0.3                  |
| 28        | 0.2~0.5                  |
| 29        | 0.08 or 0.23 or 0.19 or 0.37 |
| 30        | 0.4                      |
| 31        | 0.38                     |
| 32        | -1.19 or -0.33           |
| 33        | 0.71 or 0.56             |
| 34        | 0.362                    |
| 35        | 0.24±0.02                |
| 36        | -(0.134± 0.006)          |
| This Work | -(0.49± 0.06)            |

different approaches i.e. the QCD sum rules in the external fields and the light-cone QCD sum rules, can result in different predictions in magnitude with the same sign [29, 35, 36]. The magnetic moments of the baryons are fundamental parameters as their masses, which have copious information about the underlying quark structures, different substructures can lead to very different results. The width of the pentaquark state $\Theta^+(1540)$ is so narrow, the small magnetic moment may be extracted from photo-production experiments eventually in the near future, which may be used to distinguish the preferred configurations from various theoretical models, obtain more insight into the relevant degrees of freedom and deepen our understanding about the underlying dynamics that determines the properties of the exotic pentaquark states.

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References

[1] LEPS collaboration, T. Nakano et al., Phys. Rev. Lett. 91 (2003) 012002; DIANA Collaboration, V. V. Barmin et al., Phys. Atom. Nucl. 66 (2003) 1715; CLAS Collaboration, S. Stepanyan et al., Phys. Rev. Lett. 91 (2003) 252001; CLAS Collaboration, V. Kubarovsky et al., Phys. Rev. Lett. 92 (2004) 032001; SAPHIR Collaboration, J. Barth et al., Phys. Lett. B572 (2003) 127; A. E. Asratyan, A. G. Dolgolenko and M. A. Kubantsev, Phys. Atom. Nucl. 67 (2004) 682; HERMES Collaboration, A. Airapetian et al., Phys. Lett. B585 (2004) 213; SVD Collaboration, A. Aleev et al., hep-ex/0401024; ZEUS Collaboration, S. V. Chekanov, hep-ex/0404007; CLAS Collaboration, H. G. Juengst, nucl-ex/0312019.

[2] D. Diakonov, V. Petrov and M. V. Polyakov, Z. Phys. A359 (1997) 305.

[3] R. Jaffe and F. Wilczek, Phys. Rev. Lett. 91 (2003) 232003.

[4] M. Karliner and H. J. Lipkin, Phys. Lett. B575 (2003) 249.

[5] S. L. Zhu, Phys. Rev. Lett. 91 (2003) 232002.

[6] R. D. Matheus, F. S. Navarra, M. Nielsen, R. Rodrigues da Silva and S. H. Lee, Phys. Lett. B578 (2004) 323.

[7] J. Sugiyama, T. Doi and M. Oka, Phys. Lett. B581 (2004) 167.

[8] M. Eidemüller, Phys. Lett. B597 (2004) 314.

[9] R. D. Matheus and S. Narison, hep-ph/0412063.

[10] Y. Kanada-En’yo, O. Morimatsu and T. Nishikawa, hep-ph/0404144.

[11] S. Takeuchi and K. Shimizu, hep-ph/0410286.

[12] T. W. Chiu and T. H. Hsieh, hep-ph/0403020; T. W. Chiu and T. H. Hsieh, hep-ph/0501227.

[13] S. Sasaki, Phys. Rev. Lett. 93 (2004) 152001; F. Csikor, Z. Fodor, S. D. Katz and T. G. Kovacs, JHEP 0311 (2003) 070.

[14] N. Ishii, T. Doi, H. Iida, M.Oka, F. Okiharu, H. Suganuma, Phys. Rev. D71 (2005) 034001; N. Mathur, F.X. Lee, A. Alexandru, C. Bennhold, Y. Chen, S. J. Dong, T. Draper, I. Horvath, K. F. Liu, S. Tamhankar and J.B. Zhang, Phys. Rev. D70 (2004) 074508.

[15] E. Shuryak and I. Zahed, Phys. Lett. B589 (2004) 21.

[16] B. K. Jennings and K. Maltman, Phys. Rev. D69 (2004) 094020.
[17] F. Stancu and D. O. Riska, Phys. Lett. **B575** (2003) 242.

[18] A. Hosaka, Phys. Lett. **B571** (2003) 55.

[19] C. E. Carlson, C. D. Carone, H. J. Kwee and V. Nazaryan, Phys. Rev. **D70** (2004) 037501.

[20] F. Huang, Z. Y. Zhang, Y. W. Yu and B. S. Zou, Phys. Lett. **B586** (2004) 69.

[21] C. E. Carlson C. D. Carone, H. J. Kwee and V. Nazaryan, Phys. Lett. **B573** (2003) 101.

[22] B. Wu and B. Q. Ma, Phys. Rev. **D70** (2004) 097503.

[23] Z. G. Wang, W. M. Yang and S. L. Wan, [hep-ph/0501015](http://arxiv.org/abs/hep-ph/0501015).

[24] A. W. Thomas, K. Hicks and A. Hosaka, Prog. Theor. Phys. **111** (2004) 291.

[25] P. Z. Huang, W. Z. Deng, X. L. Chen and S. L. Zhu, Phys. Rev. **D69** (2004) 074004.

[26] Q. Zhao, Phys. Rev. **D69** (2004) 053009; Erratum-ibid. **D70** (2004) 039901.

[27] H. C. Kim and M. Praszalowicz, Phys. Lett. **B585** (2004) 99.

[28] S. I. Nam, A. Hosaka and H. C. Kim, Phys. Lett. **B579** (2004) 43.

[29] Y. R. Liu, P. Z. Huang, W. Z. Deng, X. L. Chen and S. L. Zhu, Phys. Rev. **C69** (2004) 035205.

[30] T. Inoue, V. E. Lyubovitskij, T. Gutsche and A. Faessler, [hep-ph/0408057](http://arxiv.org/abs/hep-ph/0408057).

[31] R. Bijker, M. M. Giannini and E. Santopinto, Phys. Lett. **B595** (2004) 260.

[32] K. Goeke, H. C. Kim, M. Praszalowicz and G. S. Yang, [hep-ph/0411195](http://arxiv.org/abs/hep-ph/0411195).

[33] D. K. Hong, Y. J. Sohn and I. Zahed, Phys. Lett. **B596** (2004) 191.

[34] P. Jimenez Delgado, [hep-ph/0409128](http://arxiv.org/abs/hep-ph/0409128).

[35] Z. G. Wang, W. M. Yang and S. L. Wan, [hep-ph/0501278](http://arxiv.org/abs/hep-ph/0501278).

[36] Z. G. Wang, S. L. Wan and W. M. Yang, [hep-ph/0503007](http://arxiv.org/abs/hep-ph/0503007).

[37] I. I. Balitsky, V. M. Braun and A. V. Kolesnichenko, Sov. J. Nucl. Phys. **44** (1986) 1028; Nucl. Phys. **B312** (1989) 509; V. L. Chernyak and I. R. Zhitnitsky, Nucl. Phys. **B345** (1990) 137; V. L. Chernyak and A. R. Zhitnitsky, Phys. Rep. **112** (1984) 173.

[38] V. M. Braun and I. E. Filyanov, Z. Phys. **C44** (1989) 157.
[39] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. B147 (1979) 385, 448.

[40] B. L. Ioffe and A. V. Smilga Nucl. Phys. B232 (1984) 109.

[41] I. I. Balitsky and A. V. Yung, Phys. Lett. B129 (1983) 328.

[42] T.M. Aliev, I. Kanik and M. Savci, Phys. Rev. D68 (2003) 056002; T.M. Aliev, A. Ozpineci and M. Savci. Phys. Rev. D66 (2002) 016002, Erratum-ibid. D67 (2003) 039901; Phys. Rev. D65 (2002) 096004; Phys. Rev. D65 (2002) 056008; Phys. Lett. B516 (2001) 299; Nucl. Phys. A678 (2000) 443; T. M. Aliev and A. Ozpineci. Phys. Rev. D62 (2000) 053012.

[43] A. De Rujula, H. Georgi and S. L. Glashow, Phys. Rev. D12 (1975) 147; T. DeGrand, R. L. Jaffe, K. Johnson and J. E. Kiskis, Phys. Rev. D12 (1975) 2060; G. ’t Hooft, Phys. Rev. D14 (1976) 3432 [Erratum-ibid. Phys. Rev. D18 (1978) 2199]; E. V. Shuryak, Nucl. Phys. B203 (1982) 93; T. Schafer and E. V. Shuryak, Rev. Mod. Phys. 70 (1998) 323; E. Shuryak and I. Zahed, Phys. Lett. B 589 (2004) 21.

[44] I. I. Balitsky, V. M. Braun, and A. V. Kolesnichenko, Nucl. Phys. B312 (1989) 509.

[45] V. M. Braun and I. E. Filyanov, Z. Phys. C48 (1990) 239; A. Ali and V. M. Braun, Phys. Lett. B359 (1995) 223; G. Eilam, I. Halperin and R. R. Mendel, Phys. Lett. B361 (1995) 137; P. Ball, V. M. Braun and N. Kivel, Nucl. Phys. B649 (2003) 263.

[46] Z. G. Wang, J. Phys. G28 (2002) 3007.

[47] V. M. Belyaev, V. M. Braun, A. Khodjamirian and R. Rückl, Phys. Rev. D51 (1995) 6177; H. Kim, S. H. Lee and M. Oka, Prog. Theor. Phys. 109 (2003) 371; Z. G. Wang, W. M. Yang and S. L. Wan, Eur. Phys. J. C37 (2004) 223.
Fig. 1  \(- (F_1(0) + F_2(0))\) as function of \(M^2\)

\(- (F_1(0) + F_2(0))\) as function of \(M^2\)

\(s_0 = 3.7 \text{GeV}^2\)