MAGNETIC MOMENT OF THE PENTAQUARK STATE $\Theta^+(1540)$

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Abstract

In this article, we study the magnetic moment of the pentaquark $\Theta^+(1540)$ as (scalar)diquark-(pseudoscalar)diquark-antiquark state with the QCD sum rules approach in the external electromagnetic field. Due to the special structure of the interpolating current, only the electromagnetic interactions of the $s$ quark with the external field have contributions to the magnetic moment with tensor structure $\{\sigma_{\mu\nu} \hat{p} + \hat{p}\sigma_{\mu\nu}\}$. The numerical results indicate the magnetic moment of the pentaquark state $\Theta^+(1540)$ is about $\mu_{\Theta^+} = (0.16 \pm 0.03)\mu_N$.

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1 Introduction

Intense theoretical investigations have been motivated to clarify the quantum numbers and to understand the under-structures of the pentaquark state $\Theta^+(1540)$ since its observation [1, 2, 3]. The extremely narrow width below 10 MeV puts forward a serious challenge to all theoretical models, in the conventional uncorrelated quark models the expected width is supposed to be of the order of several hundred MeV, since the strong decay $\Theta^+ \to K^+N$ is Okubo-Zweig-Iizuka (OZI) super-allowed. The zero of the third component of isospin $I_3 = 0$ and the absence of isospin partners suggest that the baryon $\Theta^+(1540)$ is an isosinglet, while the spin and parity have not been experimentally determined yet and no consensus has ever been reached on the theoretical side. Determining the parity of the pentaquark state $\Theta^+(1540)$ is of great importance in establishing its basic quantum numbers and in understanding the low energy QCD especially when multiquarks are involved. The experiments of photo- or electro-production and proton-proton collision can be used to determine the fundamental quantum numbers of the pentaquark state $\Theta^+(1540)$, such as spin and parity [4]. The magnetic moment of the $\Theta^+(1540)$ $\mu_{\Theta^+}$ is an important ingredient in studying the cross sections of the photo-production, and may be extracted from the experiments eventually in the future. In fact, the magnetic moments of the pentaquark states are fundamental parameters as their masses, which have copious information about the underlying quark structures, can be used to distinguish the preferred quark configurations from various theoretical models and deepen our understanding of the underlying dynamics.

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There have been several works on the magnetic moment $\mu_{\Theta^+}$ \[5, 6, 7, 8\], in this article, we take the point of view that the quantum numbers of the pentaquark state $\Theta^+(1540)$ are $J = \frac{1}{2}$, $I = 0$, $S = +1$, and study its magnetic moment $\mu_{\Theta^+}$ with the QCD sum rules approach \[9\].

The article is arranged as follows: we derive the QCD sum rules in the external electromagnetic field for the magnetic moment $\mu_{\Theta^+}$ in section II; in section III, numerical results; section IV is reserved for conclusion.

2 QCD Sum Rules in External Electromagnetic Field

In the following, we write down the two-point correlation function $\Pi_\Theta(p)$ in the presence of a weak external electromagnetic field $F_{\alpha\beta}$ \[10\],

$$\Pi_\Theta(p) = i \int d^4x e^{ip\cdot x} \langle 0 | T \{ J(x) J(0) \} | 0 \rangle F_{\alpha\beta},$$

$$= \Pi_0(p) + \Pi_{\mu\nu}(p) F^{\mu\nu} + \cdots,$$  \hspace{1cm} (1)

where the $\Pi_0(p)$ is the correlation function without the external field $F_{\alpha\beta}$ and the $\Pi_{\mu\nu}(p)$ is the linear response term. In this article, we take the (scalar)diquark-(pseudoscalar)diquark-antiquark type interpolating current $J(x)$ for the pentaquark state $\Theta^+(1540)$ \[11\],

$$J(x) = \epsilon^{abc} \epsilon^{def} \epsilon^{cfg} \{ u_a^T(x) C d_b(x) \} \{ u_d^T(x) C \gamma_5 d^e(x) \} C \bar{s}_g^T(x),$$  \hspace{1cm} (2)

here the $a, b, c \cdots$ are color indexes. There have been several works on the magnetic moment of the $\Theta^+(1540)$ using the diquark-triquark type current $J_1(x)$ and diquark-diquark-antiquark type current $J_2(x)$ within the framework of the QCD sum rules approach \[6, 7, 8\],

$$J_1(x) = \frac{1}{\sqrt{2}} \epsilon^{abc} \{ u_a^T(x) C \gamma_5 d_b(x) \} \{ u_c(x) \bar{s}_e(x) i \gamma_5 d_e(x) - u_c(x) \bar{s}_e(x) i \gamma_5 u_e(x) \},$$  \hspace{1cm} (3)

$$J_2(x) = \{ t \eta_1(x) + \eta_2(x) \},$$  \hspace{1cm} (4)

$$\eta_1(x) = \frac{1}{\sqrt{2}} \epsilon^{abc} \left[ u_a^T(x) C \gamma_5 d_b(x) \right] \left[ u_c^T(x) C \gamma_5 d_e(x) \right] C \bar{s}_g^T(x) - (u \leftrightarrow d),$$

$$\eta_2(x) = \frac{1}{\sqrt{2}} \epsilon^{abc} \left[ u_a^T(x) C d_b(x) \right] \left[ u_c^T(x) C d_e(x) \right] C \bar{s}_g^T(x) - (u \leftrightarrow d).$$

There exist a great number of possible quark configurations satisfy the Fermi statistics and the color singlet condition for the under-structures of the pentaquark state $\Theta^+(1540)$, if we release stringent dynamical constraints. Different quark configurations can be implemented by different interpolating currents, which are always

\[^2\text{For technical details, one can consult Refs.\[7, 8\].}\]
lead to substantially different magnetic moments. Although the three interpolating baryon currents $J$, $J_1(x)$, $J_2(x)$ can all give satisfactory masses about 1.5$GeV$ for the $\Theta^+(1540)$, the resulting magnetic moments are substantially different. Without additionally powerful constraints (such as the magnetic moment), we cannot select the preferred quark configurations from various theoretical quark models.

The Fierz re-ordering of the interpolating current $J(x)$ can lead to the following sub-structures,

$$J(x) = \frac{1}{4} e^{abc}(u^T_a C d_b) \left\{ -d_c(\bar{s}\gamma_5 u) - \gamma^\mu d_c(\bar{s}\gamma_5 \gamma_\mu u) + \frac{1}{2} \sigma^{\mu\nu} d_c(\bar{s}\sigma_{\mu\nu} \gamma_5 u) + \gamma^\mu \gamma_5 d_c(\bar{s}\gamma_\mu u) - \gamma_5 d_c(\bar{s}u) - (u \leftrightarrow d) \right\}.$$  \hspace{1cm} (5)

A naive result of the Fierz re-ordering may be the appearance of the reducible contributions with the sub-structure of $udd - u\bar{s}$ (i.e. $N - K$) clusters in the two-point correlation function $\Pi_0(p)$ \cite{12}, however, in our calculations with the interpolating current $J(x)$ in Eq.\,(2), no such factorable $udd - u\bar{s}$ terms appear, so there are no reducible $N - K$ contributions to the correlation function $\Pi_0(p)$. The re-ordering in Dirac spin space is always accompanied with color re-arrangement, which involves the underlying dynamics. The appearance of the $N - K$ component in the Fierz re-ordering maybe manifest the possibility (not the probability) of the evolution from the $\Theta^+(1540)$ to the $NK$ final state without net quark-antiquark pairs creation (maybe the quark-antiquark pairs created and annihilated subsequently), which is significantly in contrast to the conventional baryons, however, we have no knowledge about the detailed process of the evolution. If there are really some $N - K$ components in the interpolating current $J(x)$, they should be factorized out, the remainder can not have the correct quantum numbers to interpolate the $\Theta^+(1540)$.

For detailed discussions about this subject, one can see Ref.\,[13].

The interpolating current $J(x)$ in Eq.\,(2) can couple to the pentaquark states with both negative and positive parity, and picks out only the state with the lowest mass without knowledge about its parity, for example, we use the Ioffe current $J_\rho(x) = e^{abc}(u^T_a(x)C\gamma_\mu u_b(x))\gamma_5\gamma^\mu d_c(x)$ to interpolate the proton and pick out only the lowest energy state \cite{14}, which happen to have positive parity, while the possible negative parity states are included in the high resonances and continuum states. It is not necessary to include both the negative and positive states in the phenomenological spectral density. As the electromagnetic vertex of the pentaquark states with either negative or positive parity is the same, we can extract the absolute value of the magnetic moment of the lowest pentaquark state.

The linear response term $\Pi_{\mu\nu}(p)$ in the weak external electromagnetic field $F_{\alpha\beta}$ has three different Dirac tensor structures,

$$\Pi_{\mu\nu}(p) = \Pi(p) \{ \sigma_{\mu\nu} \hat{\rho} + \hat{\rho} \sigma_{\mu\nu} \} + \Pi_1(p)i \{ p_\mu \gamma_\nu - p_\nu \gamma_\mu \} \hat{\rho} + \Pi_2(p) \sigma_{\mu\nu}.$$  \hspace{1cm} (6)

The first structure has an odd number of $\gamma$-matrix and conserves chirality, the second and third have even number of $\gamma$-matrices and violate chirality. In the original QCD
sum rules analysis of the nucleon magnetic moments \[10\], the interval of dimensions (of the condensates) for the odd structure is larger than the interval of dimensions for the even structures, one may expect a better accuracy of the results obtained from the sum rules at the odd structure. In this article, the spin of the pentaquark state $\Theta^+(1540)$ is supposed to be $\frac{1}{2}$, just like the nucleon, we can choose the first Dirac tensor structure $\{\sigma_{\mu\nu}\vec{p} + \vec{p}\sigma_{\mu\nu}\}$. The phenomenological spectral density can be written as

$$\frac{\text{Im}\Pi(s)}{\pi} = \frac{1}{4}\{F_1(0) + F_2(0)\} f_0^2 \delta'(s - m_{\Theta^+}^2) + C_{\text{subtract}} \delta(s - m_{\Theta^+}^2) + \cdots,$$

where the first term corresponds to the magnetic moment $\mu_{\Theta^+}$, and is of double-pole. The second term comes from the electromagnetic transitions between the pentaquark state $\Theta^+(1540)$ and the excited states, and is of single-pole. Here we introduce the quantity $C_{\text{subtract}}$ to parameterize the electromagnetic transitions between the ground pentaquark state and the high resonances, it may have complex dependence on the energy $s$ and high resonance masses. However, we have no knowledge about the high resonances, even the existence of the ground pentaquark state $\Theta^+(1540)$ is not firmly established, which is in contrast to the conventional baryons, in those channels we can use the experimental data as guides in constructing the phenomenological spectral densities. In practical manipulations, we can take the $C_{\text{subtract}}$ as an unknown constant, and fitted to reproduce reliable values for the form factors $F_1(0) + F_2(0)$. The higher resonances and continuum states in Eq.(7) are neglected for simplicity. From the electromagnetic form factors $F_1(0)$ and $F_2(0)$, we can obtain the magnetic moment $\mu_{\Theta^+}$,

$$\mu_{\Theta^+} = \{F_1(0) + F_2(0)\} \frac{e_{\Theta^+}}{2m_{\Theta^+}}.$$

After performing the operator product expansion in the deep Euclidean space-time region, we can express the correlation functions at the level of quark-gluon degrees of freedom into the following form through dispersion relation,

$$\Pi(P^2) = \frac{e_s}{\pi} \int_{m_0^2}^{s_0} ds \frac{\text{Im}[A(s)]}{s + P^2} + \cdots,$$

where

$$\frac{\text{Im}[A(s)]}{\pi} = \frac{s^4}{21254!\pi^8} - \frac{m_s\chi\langle\bar{s}s\rangle s^3}{2754!\pi^6} + \frac{s^2}{2134!\pi^6} \left(\frac{\alpha_sGG}{\pi}\right).$$

The presence of the external electromagnetic field $F_{\mu\nu}$ induces three new vacuum condensates i.e. the vacuum susceptibilities $\chi$, $\kappa$ and $\xi$ in the QCD vacuum \[10\]. The values with different theoretical approaches are different from each other, for a short review, one can see Ref.\[15\]. Here we shall adopt the values $\chi = -4.4 \text{ GeV}^{-2}$, $\kappa = 0.4$ and $\xi = -0.8$ \[10,16\]. From Eqs.(9-10), we can see that due to the special structure of the diquark-diquark-antiquark type interpolating current $J(x)$ (also
the $J_2(x)$, the $u$ and $d$ quarks which constitute the diquarks have no contributions to the magnetic moment though they have electromagnetic interactions with the external field, the net contributions to the magnetic moment come from the $s$ quark only, which is different significantly from the results obtained in Refs.\[8, 9\] with the diquark-triquark type interpolating current $J_1(x)$. Although the diquark-diquark-antiquark type and diquark-triquark type configurations implemented by the interpolating currents $J(x)$ (or $J_1(x)$) and $J_2(x)$ respectively can give satisfactory masses for the pentaquark state $\Theta^+(1540)$, the resulting magnetic moments are substantially different, once the magnetic moment can be extracted from the electro- or photo-production experiments, we can select the preferred configuration.

Finally we obtain the sum rules for the form factors $F_1(0)$ and $F_2(0)$,

$$-\frac{1}{4} \{ F_1(0) + F_2(0) \} \frac{1 + CM^2}{M^4} f_0^2 e^{-\frac{m_s^2}{M^2}} = \frac{1}{M^2} \int_{m_s^2}^{s_0} ds \frac{\text{Im}[A(s)]}{\pi} e^{-\frac{s}{M^2}},$$

(11)

where the definition $\langle 0 | J(0) | \Theta^+(p) \rangle = f_0 u(p)$ has been used, the $m_s$ is the strange quark mass and $s_0$ is the threshold parameter used to subtract the contributions from the higher resonances and continuum states. The Borel transform can not eliminate the contaminations from the single-pole terms, we introduce the parameter $C$ which proportional to the $C_{\text{subtract}}$ in Eq.(7) to the subtract the contaminations. We have no knowledge about the electromagnetic transitions between the pentaquark state $\Theta^+(1540)$ and the excited states (or high resonances), the $C$ can be taken as a free parameter, we choose the suitable values for $C$ to eliminate the contaminations from the single-pole terms to obtain the reliable sum rules. The contributions from the single-pole terms may as large as or larger than the double-pole term, in practical calculations, the $C$ can be fitted to give stable sum rules with respect to variations of the Borel parameter $M^2$ in a suitable interval. Taking the $C$ as an unknown constant has smeared the complex energy $s$ and high resonances masses dependence, which will certainly impair the prediction power. As there really exists a platform with the variations of the Borel parameter $M^2$, the predictions still make sense. Furthermore, from the correlation function $\Pi_0(p)$ in Eq.(1), we can obtain the sum rules for the coupling constant $f_0$ \[11\],

$$f_0^2 e^{-\frac{m_s^2}{M^2}} = \int_{m_s^2}^{s_0} ds e^{-\frac{s}{M^2}} \rho_0(s),$$

(12)

where

$$\rho_0 = \frac{s^5}{2^{10}5!17\pi^8} + \frac{m_s \langle \bar{s}s \rangle s^3}{2^{8}5!3!\pi^6} - \frac{m_s \langle \bar{s}g_sGs \rangle s^2}{2^{9}4!3!\pi^6} + \frac{s^3}{2^{10}5!3!\pi^6} \frac{\langle \alpha_sGG \rangle}{\pi}.$$

3 Numerical Results

The input parameters are taken as $\chi = -(4.4 \pm 0.4)GeV^{-2}$, $\langle \bar{s}s \rangle = (0.8 \pm 0.1)\langle \bar{q}q \rangle$, $\langle \bar{s}g_sGs \rangle = m_0^2 \langle \bar{s}s \rangle$, $m_0^2 = (0.8 \pm 0.1)GeV^2$, $\langle \bar{q}q \rangle = -(0.24 \pm 0.01GeV)^3$, $\langle \alpha_sGG \rangle = $
(0.33GeV)^4, \, m_u = m_d = 0 \text{ and } m_s = (140 \pm 10)MeV. \text{ Here we use the standard values \([9]\), small variations of those condensates will not lead to large changes about the numerical values. The threshold parameter } \sqrt{s_0} \text{ is chosen to vary between } (1.7 - 1.9)GeV \text{ to avoid possible contaminations from higher resonances and continuum states. In the region } M^2 = (1.5 - 3.5)GeV^2, \text{ the sum rules for } F_1(0) + F_2(0) \text{ are almost independent of the Borel parameter } M^2, \text{ which are shown in Fig.1 for } \chi = -4.4GeV^{-2}, \langle \bar{s}s \rangle = 0.8\langle \bar{q}q \rangle, \, m_0^2 = 0.8GeV^2, \, \langle \bar{q}q \rangle = (-0.24GeV)^3, \, m_s = 140MeV. \text{ For } \sqrt{s_0} = (1.7 - 1.9)GeV, \text{ we obtain the values}

\[ F_1(0) + F_2(0) = 0.27 \pm 0.05, \]

\[ \mu_{\Theta^+} = (0.27 \pm 0.05)\frac{e_{\Theta^+}}{2m_{\Theta^+}}, \]

\[ = (0.16 \pm 0.03)\mu_N, \]

(13)

where the } \mu_N \text{ is the nucleon magneton. Although the numerical values for the magnetic moment } \mu_{\Theta^+} \text{ vary with theoretical approaches (i.e. } \mu_{\Theta^+} \approx (0.1 - 0.7)\mu_N \text{ or } \mu_{\Theta^+} \approx -(0.1 - 1.2)\mu_N), \text{ they are small in general; our numerical results are consistent with most of the existing values of theoretical estimations \[5, 14, 15, 16\]. For a short review of the existing calculations of the magnetic moments of the } \Theta^+(1540), \text{ one can consult Ref.\[7\]. The main contributions to the magnetic moments } \mu_{\Theta^+} \text{ come from the perturbative term in Eq.}(10), \text{ about } 70\%, \text{ the contributions from terms of the quark condensates and gluons condensates have opposite sign, and the resulting net contributions are about } 30\%, \text{ the high dimensional condensates are neglected as they are suppressed by large denominators. For the conventional ground state mesons and baryons, due to the resonance dominates over the QCD continuum contributions, the good convergence of the operator product expansion, and the useful experimental guidance on the threshold parameter } s_0, \text{ we can obtain the}

\[ \text{Figure 1: } |F_1(0) + F_2(0)| \text{ with the Borel Parameter } M^2. \]
fiducial Borel mass region. However, in the QCD sum rules for the pentaquark states, the spectral density $\rho(s) \sim s^m$ with $m$ larger than the corresponding ones in the sum rules for the conventional baryons, larger $m$ means stronger dependence on the continuum or the threshold parameter $s_0$ \cite{17}. In Eq.\(12\), due to the large continuum contributions, the threshold parameter $s_0$ has to be fixed ad hoc or intuitively. In this article, the threshold parameter $s_0$ are taken to be $\sqrt{s_0} = (1.7 - 1.9) GeV$, the mass $m_{\Theta^+} = 1540 MeV$ and the width $\Gamma_{\Theta^+} < 10 MeV$, the contributions from the lowest pentaquark state can be successfully included in. Although the uncertainties of the condensates, the neglect of the higher dimension condensates, the lack of perturbative QCD corrections, etc, will result in errors, we have stable sum rules, which are shown in Fig.\(1\), the predictions still make sense, or qualitative at least.

In Ref.\[8\], the authors take the diquark-diquark-antiquark current $J_2(x)$ which is linear superposition of both S-type and P-type baryon currents ( $\eta_1(x)$ and $\eta_2(x)$) to calculate the magnetic moment $\mu_{\Theta^+}$ in two approaches (i.e. the QCD sum rules in the external field and the light-cone QCD sum rules), and obtain $\mu_{\Theta^+} = -(0.11 \pm 0.02) \mu_N$ and $\mu_{\Theta^+} = -(0.1 \sim 0.5) \mu_N$, respectively. As the values obtained from the QCD sum rules in the external field are more stable than the corresponding ones from the light-cone QCD sum rules, $\mu_{\Theta^+} = -(0.11 \pm 0.02) \mu_N$ is more reliable. For the diquark-diquark-antiquark type interpolating currents $J(x)$ and $J_2(x)$, only the electromagnetic interactions of the $s$ quark with the external field have contributions to magnetic moment with the tensor structure $\{\sigma_{\mu\nu} \tilde{p} \sigma_{\mu\nu} \}$, the resulting magnetic moments are substantially different. For the diquark-triquark type interpolating currents $J_1(x)$, the electromagnetic interactions of all the $u$, $d$ and $s$ quarks with the external field have contributions to the magnetic moment, and $\mu_{\Theta^+} = (0.24 \pm 0.02) \mu_N$ \[7\]. The interpolating currents $J(x)$, $J_1(x)$ and $J_2(x)$ with different quark configurations can all give satisfactory masses for the $\Theta^+ (1540)$, without additional powerful dynamical constraints, we can not pick out the preferred configurations from various theoretical models ( constituent quark models or cluster quark models ) . The magnetic moment plays an important role in understanding the understructures of the pentaquark state.

4 Conclusion

In summary, we have calculated the magnetic moment of the pentaquark $\Theta^+ (1540)$ as (scalar)diquark-(pseudoscalar)diquark-antiquark state with the QCD sum rules approach in the weak external electromagnetic field. The numerical results are consistent with most of the existing values of theoretical estimations, $\mu_{\Theta^+} = (0.27 \pm 0.05)\frac{\mu_N}{2m_{\Theta^+}} = (0.16 \pm 0.03) \mu_N$. 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 small magnetic moment $\mu_{\Theta^+}$ may be extracted from the electro- or photo-production experiments eventually in the future, which may be used to distinguish the preferred
quark configurations and QCD sum rules 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. Rev. 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] For example, D. Diakonov, V. Petrov and M. V. Polyakov, Z. Phys. A359 (1997) 305; R. Jaffe and F. Wilczek, Phys. Rev. Lett. 91 (2003) 232003; M. Karliner and H. J. Lipkin, Phys. Lett. B575 (2003) 249; S. L. Zhu, Phys. Rev. Lett. 91 (2003) 232002; R. D. Matheus, F. S. Navarra, M. Nielsen, R. Rodrigues da Silva and S. H. Lee, Phys. Lett. B578 (2004) 323; M. Eidemüller, Phys. Lett. B597 (2004) 314; T. W. Chiu and T. H. Hsieh, hep-ph/0403020; C. E. Carlson C. D. Carone, H. J. Kwee and V. Nazaryan, Phys. Rev. D70 (2004) 037501; E. Shuryak and I. Zahed, Phys. Lett. B589 (2004) 21; F. Huang, Z. Y. Zhang, Y. W. Yu and B. S. Zou, Phys. Lett. B586 (2004) 69; C. E. Carlson, C. D. Carone, H. J. Kwee and V. Nazaryan, Phys. Lett. B579 (2004) 52; Y. S. Oh, H. C. Kim and S. H. Lee, Phys. Rev. D69 (2004) 014009; J. Haidenbauer and G. Krein, Phys. Rev. C68 (2003) 052201; F. Csikor, Z. Fodor, S. D. Katz and T. G. Kovacs, JHEP 0311 (2003) 070; T. D. Cohen, Phys. Lett. B581 (2004) 175; L. Ya. Glozman, Phys. Lett. B575 (2003) 18; 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; F.
Stancu and D. O. Riska, Phys. Lett. B575 (2003) 242; C. E. Carlson C. D. Carone, H. J. Kwee and V. Nazaryan, Phys. Lett. B573 (2003) 101.

[3] M. Oka, Prog. Theor. Phys. 112 (2004) 1; S. L. Zhu, Int. J. Mod. Phys. A19 (2004) 3439; S. L. Zhu, hep-ph/0410002; F. E. Close, hep-ph/0311087 B. K. Jennings and K. Maltman, Phys. Rev. D69 (2004) 094020; and references therein.

[4] A. W. Thomas, K. Hicks and A. Hosaka, Prog. Theor. Phys. 111 (2004) 291; S. I. Nam, A. Hosaka and H. C. Kim, Phys. Lett. B579 (2004) 43.

[5] Q. Zhao, Phys. Rev. D69 (2004) 053009; Erratum-ibid. D70 (2004) 039901; H. C. Kim and M. Praszalowicz, Phys. Lett. B585 (2004) 99; K. Goeke, H. C. Kim, M. Praszalowicz and G. S. Yang, hep-ph/0411195; G. S. Yang, H. C. Kim, M. Praszalowicz and K. Goeke, Phys. Rev. D70 (2004) 114002; Y. R. Liu, P. Z. Huang, W. Z. Deng, X. L. Chen and S. L. Zhu, Phys. Rev. C69 (2004) 035205; R. Bijker, M. M. Giannini and E. Santopinto, Phys. Lett. B595 (2004) 260; T. Inoue, V. E. Lyubovitskij, T. Gutsche and A. Faessler, hep-ph/0408057; D. K. Hong, Y. J. Sohn and I. Zahed, Phys. Lett. B596 (2004) 191; P. Jimenez Delgado, hep-ph/0409128.

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

[7] Z. G. Wang, W. M. Yang and S. L. Wan, J. Phys. G31 (2005) 703.

[8] Z. G. Wang, S. L. Wan and W. M. Yang, hep-ph/0503007.

[9] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. B147 (1979) 385, 448; L. J. Reinders, H. Rubinstein and S. Yazaki, Phys. Rept. 127 (1985) 1; S. Narison, World Sci. Lect. Notes Phys. 26 (1989) 1.

[10] B. L. Ioffe and A. V. Smilga, Nucl. Phys. B232 (1984) 109; I. I. Balitsky and A. V. Yung, Phys. Lett. B129 (1983) 328.

[11] J. Sugiyama, T. Doi and M. Oka, Phys. Lett. B581 (2004) 167; S. Sasaki, Phys. Rev. Lett. 93 (2004) 152001; Z. G. Wang, W. M. Yang and S. L. Wan, hep-ph/0501015.

[12] Y. Kondo, O. Morimatsu, T. Nishikawa, Phys. Lett. B611 (2005) 93; S. H. Lee, H. Kim, Y. Kwon, Phys. Lett. B609 (2005) 252; Y. Kwon, A. Hosaka, S. H. Lee, hep-ph/0505040.

[13] Z. G. Wang, W. M. Yang, S. L. Wan, hep-ph/0504151.

[14] B. L. Ioffe, Nucl. Phys. B188 (1981) 317, Erratum-ibid. B191 (1981) 591.
[15] Z. G. Wang, J. Phys. **G28** (2002) 3007.

[16] V. M. Belyaev and Y. I. Kogan, Yad. Fiz. **40** (1984) 1035; I. I. Balitsky, V. M. Braun and A. V. Kolesnichenko, Nucl. Phys. **B312** (1989) 509.

[17] R. D. Matheus, S. Narison, [hep-ph/0412063](http://arxiv.org/abs/hep-ph/0412063).