X(1835) as a Baryonium State with QCD Sum Rules

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Abstract

In this article, we take the point of view that the X(1835) be a baryonium state and calculate its mass within the framework of the QCD sum rules approach. The numerical value of the mass of the X(1835) is consistent with the experimental data. There may be some baryonium component in the X(1835) state.

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1 Introduction

In 2003, the BES collaboration observed a significant narrow near-threshold enhancement in the proton-antiproton (p\bar{p}) invariant mass spectrum from the radiative decay J/\psi \rightarrow \gamma p\bar{p} \cite{1}. The enhancement can be fitted with either an S- or P-wave Breit-Wigner resonance function. In the case of the S-wave fitted form, the Breit-Wigner mass \( M = 1859^{+3}_{-10} \pm 25 \) MeV and the width \( \Gamma < 30 \) MeV. Recently the BES collaboration observed a resonance state X(1835) in the \( \eta' \pi^+ \pi^- \) invariant mass spectrum in the process J/\psi \rightarrow \gamma \pi^+ \pi^- with the Breit-Wigner mass \( M = (1833.7 \pm 6.2 \pm 2.7) \) MeV and the width \( \Gamma = (67.7 \pm 20.3 \pm 7.7) \) MeV \cite{2}. Many theoretical works were stimulated to interpret the nature and underlying structures of the new particle, there exist many possibilities, for example, the p\bar{p} bound state, the pseudoscalar glueball, non-exotic state, etc \cite{3, 4, 5, 6}.

The X(1835) may be pseudoscalar glueball \cite{3}, which can take into account the observation of the X(1835) in the \( \eta' \pi^+ \pi^- \) channel not in the \( \pi^0 \pi^+ \pi^- \) channel, while the strong coupling to the p\bar{p} state can be related to the large contribution of the gluon axial anomaly to the proton spin. The calculations with lattice QCD and QCD sum rules indicate that the pure scalar glueballs lie around \( (1.5 - 1.7) GeV \) and the pure pseudoscalar glueballs lie around \( 2.6 GeV \) \cite{7, 8}, we have to resort to special mechanism to pull the mass down to \( 1.835 GeV \).

The X(1835) may also not be exotic state and be the candidate for the second radial excitation of the \( \eta' \) meson, which fulfil the pseudoscalar nonet \( \pi(1800), K(1830), \eta(1760), X(1835) \) \cite{5}, such assignment can explain the mass, total decay width, production rate and decay pattern phenomenologically. The decay X(1835) \rightarrow \eta' \pi^+ \pi^-

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takes place through the emission of a pair of S-wave \( \pi \) mesons, while the decay \( X(1835) \rightarrow \eta \pi \pi \) has not been observed experimentally yet. Whether or not there exists this decay mode is of great importance, further experiments are needed to prove or exclude the possibility.

In this article, we take the point of view that the \( X(1835) \) be a baryonium with the quantum numbers \( J^P C = 0^- + \) [4], and calculate its mass in the framework of the QCD sum rules approach [9]. The radiative decay of the \( J/\psi \) is generally believed to be glue-rich, which can explain the branching ratio of the decay \( J/\psi \rightarrow \gamma \eta' \) (through \( J/\psi \rightarrow \gamma G \tilde{G} \rightarrow \gamma \eta' \) ) is large while the decay ratio of the \( J/\psi \rightarrow \gamma \eta \) (through the \( \eta' - \eta \) mixing) is small, about \( (9.8 \pm 1.0) \times 10^{-4} \) [10]. The observation of the \( X(1835) \) in the \( \eta' \) channel not in the \( \eta \) channel may be due to the intermediate virtual gluons are flavor-neutral and the \( \eta' \) meson is mainly a \( SU(3) \) flavor singlet. The threshold \( 2m_p = 1876 > 1835 \) and the width \( \Gamma = 68 \), the decay \( X(1835) \rightarrow p\bar{p} \) takes place through the fall apart mechanism with re-arrangement in the color space, while suppressed kinematically and the decay occurs only through the tail of the mass distribution.

On the other hand, whether or not there exist some quark configurations which can result in the baryonium state is of great importance itself, we explore this possibility and propose a special quark configuration, later experimental data can confirm or reject this assumption.

The article is arranged as follows: we derive the QCD sum rules for the mass \( M_X \) of the \( X(1835) \) in section II; in section III, numerical results and discussions; section VI is reserved for conclusion.

2 QCD sum rules for the \( X(1835) \)

In the following, we write down the two-point correlation function \( \Pi(q^2) \) in the framework of the QCD sum rules approach,

\[
\Pi(q^2) = i \int d^4x e^{iq \cdot x} \langle 0 | T \{ J_5(x) J_5^+(0) \} | 0 \rangle ,
\]

\[
J_5(x) = \bar{J}_{kl}(x) i \gamma_5 J_{kl}(x),
\]

\[
J_{kl}(x) = \epsilon_{kmn} u_m^T(x) C \gamma^\alpha u_n(x) \gamma_5 \gamma_\alpha d_l(x),
\]

\[
\lambda_X = \langle 0 | J_5(0) | X \rangle.
\]

Here the \( k, l, m, n \) are color indexes, the \( C \) is charge conjunction matrix, the \( \alpha \) is Lorentz index. The hexaquark states can be classified as "baryonia" if they can be described as a single multiquark cluster, or "molecules" if they consist of weakly bound \( NN \) pairs. We take the pseudoscalar proton-antiproton type interpolating current \( J_5(x) \) to represent the \( X(1835) \), if we smear the color indexes, the colored constituent \( J_{kl}(x) \) has the same structure as the Ioffe current \( \eta(x) \) which interpolates the proton,

\[
\eta(x) = \epsilon_{kmn} u_m^T(x) C \gamma^\alpha u_n(x) \gamma_5 \gamma_\alpha d_k(x),
\]
the color indexes \( k \) and \( l \) in the \( J_{kl}(x) \) (in other words, the strong color interactions) bind the two constituents as a single multiquark cluster \( uud\bar{u}\bar{d} \). If we take the color singlet operator \( \eta(x) \) as the basic constituent and choose \( \eta_5(x) = \bar{\eta}(x)i\gamma_5\eta(x) \), the current \( \eta_5(x) \) can interpolate a hexaquark state, whether the compact state or the loosely deuteron-like \( pp \) bound state, it is difficult to separate the contributions of the bound state from the scattering \( pp \) state. In this article, we take the \( X(1835) \) as a baryonium state and choose the current \( J_5(x) \), although the \( \eta_5(x) \) has non-vanishing coupling with the \( X(1835) \).

According to the basic assumption of current-hadron duality in the QCD sum rules approach \[9\], we insert a complete series of intermediate states satisfying the unitarity principle with the same quantum numbers as the current operator \( J_5(x) \) into the correlation function in Eq.(1) to obtain the hadronic representation. After isolating the pole term of the lowest \( X(1835) \) state, we obtain the following result,

\[
\Pi(q^2) = \frac{\lambda_\eta^2}{M^2_X - q^2} + \cdots .
\] (5)

In the following, we briefly outline the calculations of the operator product expansion in the deep Euclidean space. In order to evaluate the correlation function \( \Pi(q^2) \) at the level of quark-gluon degrees of freedom, we determine the quark propagator in the presence of the quark and gluon condensates firstly\[^2\].

\[
S_{ab}(x) \equiv \langle 0|T\{q_a(x)\bar{q}_b(0)\}|0\rangle = \frac{i\delta_{ab}}{2\pi^2 x^4} - \frac{ig_s G_{\mu\nu}^{ab}}{32\pi^2 x^2} (\sigma_{\mu\nu}\hat{x} + \hat{x}\sigma_{\mu\nu}) - \frac{\delta_{ab}\langle \bar{q}q \rangle}{12} - \frac{\delta_{ab}\langle g_s \bar{q}\sigma Gq \rangle}{192} x^2 + \cdots ,
\] (6)

where the small masses of the \( u \) and \( d \) quarks are neglected. Then we substitute the quark propagator into the following correlation function to obtain the spectral density with the vacuum condensates adding up to dimension-12,

\[
\Pi(q^2) = 4i \int d^4xe^{iq\cdot x}\epsilon_{kij}\epsilon_{kmn}\epsilon_{k'i'j'}\epsilon_{k'm'n'}\mathrm{Tr} \left[ \gamma_\alpha \gamma_\beta S_{ll'}(x)\gamma_\alpha' \gamma_\beta' S_{ll'}(-x) \right] \mathrm{Tr} \left[ C\gamma^\beta S_{n'i'}(-x)\gamma^\alpha C S^T_{m'j'}(-x) \right] \mathrm{Tr} \left[ C\gamma^\beta S_{n'i'}(x)\gamma^\alpha' C S^T_{m'j'}(x) \right] .
\] (7)

In Eq.(7), we have taken the assumption of the vacuum saturation for the condensates, the high dimension vacuum condensates are always factorized to lower condensates with the vacuum saturation in the QCD sum rules, the factorization works well in the large \( N_c \) limit. It is obvious that such an assumption can not take into account

\[^2\]For the \( u, d \) and \( s \) quarks, the current masses are small, it is convenient to work in the \( x \)-representation and adopt the external field method, we follow the routine presented in page-28 and page-36 in the (review) article of Ref.\[9\] to carry out the operator product expansion. For the technical details, one can consult the excellent review "Hadron Properties from QCD Sum Rules", L. J. Reinders, H. Rubinstein and S. Yazaki, Phys. Rept. 127 (1985) 1.
some information in the parameter space, the straight forward calculations with the standard operator product expansions can lead to a more general expression for the vacuum condensates. We take a simple routine in Eqs.(6-7) to simplify the calculation and obtain the result in a special case, it is the common approach to deal with the multiquark states with the QCD sum rules, for example, the tetraquark states in Ref.[11]. In this article, we take into account the contributions from the quark condensates $\langle \bar{q}q \rangle$, gluon condensates $\langle \alpha_s G G \rangle$, mixed condensates $\langle \bar{q}g_8 \sigma G q \rangle$, and neglect the contributions from other high dimension condensates which are suppressed by large denominators and would not play significant roles. Once the analytical results are obtained, then we can take the current-hadron duality below the threshold $s_0$ and perform the Borel transformation with respect to the variable $Q^2 = -q^2$, finally we obtain the following sum rule,

$$\lambda_X^2 e^{-\frac{M_X^2}{M_B^2}} = \int_0^{s_0} ds e^{-\frac{s}{M_B^2}} \frac{\text{Im}\Pi(s)}{\pi},$$

$$\frac{\text{Im}\Pi(s)}{\pi} = \frac{9s^7}{2^9 \cdot 7! \cdot \pi^{10}} + \frac{s^4 \langle \bar{q}q \rangle^2}{2^4 \cdot 4! \cdot \pi^6} + \frac{5s^3 \langle \bar{q}q \rangle \langle \bar{q}g_8 \sigma Gq \rangle}{2^2 \cdot 2^2 \cdot \pi^2} - \frac{5s^3 \langle \bar{q}g_8 \sigma Gq \rangle^2}{2^{10} \cdot 3 \cdot \pi^6} - \frac{s^2 \langle \bar{q}q \rangle \langle \alpha_s G G \rangle}{2^6 \cdot 3^3 \cdot \pi^4} - \frac{s \langle \bar{q}g_8 \sigma Gq \rangle}{2^8 \cdot 3^2 \cdot \pi^4} \langle \alpha_s G G \rangle.$$

Differentiate the above sum rule with respect to the variable $\frac{1}{M_B^2}$, then eliminate the quantity $\lambda_X$, we obtain the QCD sum rule for the mass,

$$M_X^2 = \int_0^{s_0} ds e^{-\frac{s}{M_B^2}} \frac{\text{Im}\Pi(s)}{\pi} / \int_0^{s_0} ds e^{-\frac{s}{M_B^2}} \frac{\text{Im}\Pi(s)}{\pi}.$$  

It is easy to integrate over the variable $s$, we prefer this formulation for simplicity. If we replace the $e^{-\frac{s}{M_B^2}}$ with $s^n$, $n = 0, 1, 2, 3, \ldots$, we obtain the finite energy sum rule (FESR) [12],

$$\lambda_X^2 M_X^{2n} = \int_0^{s_0} ds s^n \frac{\text{Im}\Pi(s)}{\pi},$$

$$M_X^2 = \int_0^{s_0} ds s^{n+1} \frac{\text{Im}\Pi(s)}{\pi} / \int_0^{s_0} ds s^n \frac{\text{Im}\Pi(s)}{\pi}.$$  

The threshold parameter $s_0$ is determined by the condition,

$$\frac{d}{ds_0} M_X = 0.$$  

The FESRs correlate the ground state mass with the continuum threshold $s_0$, and separate the ground state from the continuum contributions at the very beginning, for some pentaquark currents, there happen exist reasonable stability regions $s_0$ [13].
The weight function $s^n$ enhances the continuum or the high mass resonances rather than the lowest ground state, we must make sure that only the lowest pole terms contribute to the FESR below the $s_0$, in some case, a naive stability region $s_0$ can not guarantee a physically reasonable value of the $s_0$ [13]. For the hexaquark state, the situation is much worse, the stability condition in Eq.(13) can not be satisfied, we discard the FESR in this article.

3 Numerical results and discussions

The input parameters are taken to be the standard values $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle = \langle \bar{q}q \rangle = -(0.24\pm0.01GeV)^3$, $\langle \bar{q}g_s\sigma Gq \rangle = m_0^2\langle \bar{q}q \rangle$, $m_0^2 = (0.8\pm0.1)GeV^2$, $\langle \frac{\alpha_sGG}{\pi} \rangle = (0.33GeV)^4$ and $m_u = m_d = 0$. In numerical calculation, we observe that the contributions from the terms with the gluon condensate $\langle \frac{\alpha_sGG}{\pi} \rangle$ are very small, and neglect the uncertainty of the gluon condensate. The main contributions to the correlation function in Eq.(8) come from the terms with $\langle \bar{q}q \rangle^2$ and $\langle \bar{q}q \rangle\langle \bar{q}g_s\sigma Gq \rangle$, about 85%; the contributions from the terms with $\langle \bar{q}q \rangle^4$ and $\langle \bar{q}g_s\sigma Gq \rangle^2$ are about 35% and 20% respectively, and cancel out with each other, the resulting net contributions are less than 15%; the contribution comes from the perturbative term is very small, about 1%. We neglect the contributions from other high dimension condensates which are suppressed by large denominators. In the QCD sum rules with the interpolating currents constructed from the multiquark configurations, the main contributions come from the terms with the condensates $\langle \bar{q}q \rangle$ and $\langle \bar{s}s \rangle$ [14], sometimes the mixed condensates $\langle \bar{q}g_s\sigma Gq \rangle$ and $\langle \bar{s}g_s\sigma Gs \rangle$ also play important roles [15].

In the following, we discuss the criterion for selecting the threshold parameter $s_0$ and Borel parameter $M_B$ in the QCD sum rules dealing with the multiquark states. For the conventional (two-quark) mesons and (three-quark) baryons, the hadronic spectral densities are experimentally well known, the separations between the ground state and excited states are large enough, the ”single-pole + continuum states” model works well in representing the phenomenological spectral densities. The continuum states can be approximated by the contributions from the asymptotic quarks and gluons, and the single-pole dominance condition can be well satisfied,

$$\int_{s_0}^{\infty} \rho_{pert} e^{\frac{s}{M_B}} ds < \int_{0}^{s_0} (\rho_{pert} + \rho_{nomp}) e^{\frac{s}{M_B}} ds , \quad (14)$$

where the $\rho_{pert}$ and $\rho_{nomp}$ stand for the contributions from the perturbative and non-perturbative part of the spectral density respectively. From the condition in Eq.(14), we can obtain the maximal value of the Borel parameter $M_B^{max}$, exceed this value, the single-pole dominance will be spoiled. On the other hand, the Borel parameter must be chosen large enough to warrant the convergence of the operator product expansion and the contributions from the high dimension vacuum condensates which are known poorly are of minor importance, the minimal value of the Borel parameter $M_B^{min}$ can be determined.
For the conventional mesons and baryons, the Borel window $M_B^{\text{max}} - M_B^{\text{min}}$ is rather large and the reliable QCD sum rules can be obtained. However, for the multiquark states i.e. tetraquark states, pentaquark states, hexaquark states, etc, the spectral densities $\rho \sim s^n$ with $n$ is larger than the ones for the conventional hadrons, the integral $\int_0^\infty s^n e^{-\frac{s}{M_B}} ds$ converges more slowly [13]. If one do not want to release the condition in Eq.(14), we have to either postpone the threshold parameter $s_0$ to very large values or choose very small values of the Borel parameter $M_B^{\text{max}}$. With large values of the threshold parameter $s_0$, for example, $s_0 \gg M_{\text{gr}}^2$, here the $\text{gr}$ stands for the ground state, the contributions from the high resonance states and continuum states are included in, we can not use the single-pole (or ground state) approximation for the spectral densities; on the other hand, with very small values of the Borel parameter $M_B^{\text{max}}$, the operator product expansion is broken down, and the Borel window $M_B^{\text{max}} - M_B^{\text{min}}$ shrinks to zero or negative values. We should resort to the "multi-pole + continuum states" to approximate the phenomenological spectral densities. The onset of the continuum states is not abrupt, the ground state, the first excited state, the second excited state, etc, the continuum states appear sequentially; the excited states may be loose bound states and have large widths. The threshold parameter $s_0$ is postponed to large value, at that energy scale, the spectral densities can be well approximated by the contributions from the asymptotic quarks and gluons, and of minor importance for the sum rules.

The present experimental knowledge about the phenomenological hadronic spectral densities of the multiquark states is rather vague, even the existence of the multiquark states is not confirmed with confidence, and no knowledge about either there are high resonances or not.

In this article, the following criteria are taken. We choose the suitable values of the Borel parameter $M_B$, on the one hand the minimal values $M_B^{\text{min}}$ are large enough to warrant the convergence of the operator product expansion, for $M_B^{\text{min}} > \sqrt{3.5}\text{GeV}$, the dominating contributions come from the terms with $\langle \bar{q}q \rangle^2$ and $\langle \bar{q}q \rangle \langle \bar{q}g_s \sigma Gq \rangle$, about 85%; on the other hand the maximal values $M_B^{\text{max}}$ are small enough to suppress the contributions from the high excited states and continuum states, we choose the naive analysis $e^{-s_0/(M_B^{\text{max}})^2} < e^{-1}$. For the hadronic spectral density, the more phenomenological analysis is preferred, we approximate the spectral density with the contribution from the single-pole term, the threshold parameter $s_0$ is taken slightly above the ground state mass ($\sqrt{s_0} > M_{\text{gr}} + \frac{\Gamma_{\text{gr}}}{2}$) to subtract the contributions from the excited states and continuum states. In this article, the threshold parameter $s_0$ is taken to be $\sqrt{s_0} = (2.1 - 2.3)\text{GeV} > 1.9\text{GeV}$. It is reasonable for the Breit-Wigner mass $M = (1833.7 \pm 6.2 \pm 2.7)\text{MeV}$ and width $\Gamma = (67.7 \pm 20.3 \pm 7.7)\text{MeV}$. The values of the $\lambda_X^2$ from the sum rules in Eq.(8) increase quickly with $\sqrt{s_0} > 2.3\text{GeV}$, which are shown in Fig.1, it may serve as indication of the onset of the high resonances and continuum states. From the Fig.1, we can see that the Borel parameter can be chosen to be $M_B^2 = (3.5 - 5.5)\text{GeV}^2$, $M_B^{\text{max}} \leq \sqrt{s_0} \leq 2.3\text{GeV}$. 


Finally, we obtain the value of the mass of the $X(1835)$,

$$M_X = (1.9 \pm 0.1)\,GeV.$$  \hspace{1cm} (15)

The numerical result is compatible with experimental data, one may reject taking the result from the more phenomenological analysis as quantitatively reliable, the result is qualitative at least. The systematic studies with the random instanton liquid model indicate that the masses of the diquarks are $m_S = (420 \pm 30)\,MeV$, $m_V = m_A = (940 \pm 20)\,MeV$, $m_T = (570 \pm 20)\,MeV$ \textsuperscript{[16]}, here the $S$, $V$, $A$ and $T$ stand for the scalar, vector, axial-vector and tensor diquarks respectively. In this article, the chosen quark configuration has two diquark constituents, a diquark $(\epsilon_{abc}u_b^T \gamma_\mu a_c)$ and an antidiquark $(\epsilon_{abc}\bar{u}_b^T \gamma_\mu \bar{u}_c^T)$, it is not surprise that the energy scale set by the diquarks is about $1.9\,GeV$. Different quark configurations can result in different masses for the hadrons, for example, the meson-meson type interpolating currents indicate the masses of the tetraquark states are less than $1\,GeV$ \textsuperscript{[17]}.

## 4 Conclusion

In this article, we take the point of view that the $X(1835)$ be a baryonium state and calculate its mass within the framework of the QCD sum rule approach. The numerical value of the mass of the $X(1835)$ is consistent with the experimental data. There may be some baryonium component in the $X(1835)$ state.
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References

[1] J. Z. Bai et al, Phys. Rev. Lett. 91 (2003) 022001.

[2] M. Ablikim et al, Phys. Rev. Lett. 95 (2005) 262001.

[3] N. Kochelev and D. P. Min, Phys. Rev. D72 (2005) 097502; G. Hao, C. F. Qiao and A. L. Zhang, [hep-ph/0512214]; X. G. He, X. Q. Li, X. Liu and J. P. Ma, [hep-ph/0509140]; N. Kochelev and D. P. Min, Phys. Lett. B633 (2006) 283.

[4] A. Datta and P. J. O'Donnell, Phys. Lett. B567 (2003) 273; S. L. Zhu and C. S. Gao, [hep-ph/0507050]; M. L. Yan, S. Li, B. Wu and B. Q. Ma, Phys. Rev. D72 (2005) 034027; G. J. Ding and M. L. Yan, [hep-ph/0511186]; G. J. Ding, J. l. Ping, M. L. Yan, [hep-ph/0510013]; G. J. Ding and M. L. Yan, Phys. Rev. C72 (2005) 015208.

[5] T. Huang and S. L. Zhu, Phys. Rev. D73 (2006) 014023.

[6] J. L. Rosner, Phys. Rev. D68 (2003) 014004; D. V. Bugg, Phys. Lett. B598 (2004) 8; B. S. Zou and H. C. Chiang, Phys. Rev. D69 (2004) 034004; B. Kerbikov, A. Stavinsky and V. Fedotov, Phys. Rev. C69 (2004) 055205; X. A. Liu, X. Q. Zeng, Y. B. Ding, X. Q. Li, H. Shen and P. N. Shen, [hep-ph/0406118]; B. Loiseau and S. Wycech, Int. J. Mod. Phys. A20 (2005) 1990; B. Loiseau and S. Wycech, Phys. Rev. C72 (2005) 011001; X. G. He, X. Q. Li and J. P. Ma, Phys. Rev. D71 (2005) 014031.

[7] C. Michael, [hep-lat/0302001] and references therein.

[8] K. Senba, M. Tanimoto, Phys. Lett. B105 (1981) 297; S. Narison, Nucl. Phys. B509 (1998) 312; A. l. Zhang, T. G. Steele, Nucl. Phys. A728 (2003) 165.

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

[10] W. M. Yao et al, J. Phys. G33 (2006) 1.

[11] H. X. Chen, A. Hosaka, S. L. Zhu, Phys. Rev. D74 (2006) 054001.
[12] S. Narison, QCD Spectral Sum Rules, World Scientific Lecture Notes in Physics 26; and references there in.

[13] R. D. Matheus and S. Narison, hep-ph/0412063; W. Wei, P. Z. Huang, H. X. Chen, S. L. Zhu, JHEP 0507 (2005) 015; Z. G. Wang, S. L. Wan and W. M. Yang, Eur. Phys. J. C45 (2006) 201.

[14] N. Kodama, M. Oka and T. Hatsuda, Nucl. Phys. A580 (1994) 445; Z. G. Wang, W. M. Yang and S. L. Wan, J. Phys. G31 (2005) 971; Z. G. Wang and W. M. Yang, Eur. Phys. J. C42 (2005) 89; Z. G. Wang, S. L. Wan, Nucl. Phys. A778 (2006) 22.

[15] For example, B. L. Ioffe, A. G. Oganesian, JETP Lett. 80 (2004) 386; H. J. Lee, N. I. Kochelev, V. Vento, Phys. Rev. D73 (2006) 014010.

[16] T. Schafer, E. V. Shuryak, J. J. M. Verbaarschot, Nucl. Phys. B412 (1994) 143.

[17] A. I. Zhang, Phys. Rev. D61 (2000) 114021.