

Θ^{++} from QCD sum rule

Tetsuo NISHIKAWA,* Yoshiko KANADA-EN'YO,† and Osamu MORIMATSU‡

*Institute of Particle and Nuclear Studies,
High Energy Accelerator Research Organization,
1-1, Ooho, Tsukuba, Ibaraki, 305-0801, Japan*

Yoshihiko KONDO§

Kokugakuin University, Higashi, Shibuya, Tokyo 150-8440, Japan

(Dated: October 31, 2004)

Abstract

We study the pentaquark with $J = 3/2$ and $I = 1$ (Θ^{++}) in the QCD sum rule approach. We derive the QCD sum rules for positive and negative parity states of the pentaquark. The QCD sum rule predicts that there exists Θ^{++} with negative parity and its mass is $1.5 \sim 1.6$ GeV. The negative parity Θ^{++} can be extremely narrow, since it lies much below the ΔK threshold and the decay into KN state is strongly suppressed due to the D -wave centrifugal barrier. Also, the possibility of the existence of the Θ^{++} with positive parity is not excluded. Although it nearly degenerates with the negative parity state, it may be broader than the negative parity state.

PACS numbers: 11.55.Fv, 11.55.Hx, 12.38.Aw, 12.39.Mk

Keywords: Pentaquark, QCD sum rules

*Electronic address: nishi@post.kek.jp

†Electronic address: yoshiko.enyo@kek.jp

‡Electronic address: osamu.morimatsu@kek.jp

§Electronic address: kondo@kokugakuin.ac.jp

An exotic baryon state with positive strangeness has recently been observed by LEPS collaboration in Spring-8 [1]. The mass is ~ 1540 MeV and the width is unusually small: $\Gamma < 25$ MeV. It was confirmed by subsequent experiments [2, 3, 4, 5, 6, 7, 8, 9, 10]. In all cases the mass is near 1540 MeV and the width is small. The spin and parity have not yet been experimentally determined. Concerning the isospin, there is an experimental indication that Θ^+ has $I = 0$ [3, 5]. This state cannot be a three-quark state since it has $S = +1$, and the minimal quark content is $(uudd\bar{s})$. It has come to be called ‘‘pentaquark Θ^+ ’’.

Besides the calculation in the chiral quark soliton model [11], which motivated the pentaquark search by LEPS collaboration, the discovery of Θ^+ has triggered an intense theoretical activity to clarify the quantum numbers and to understand the structure of the Θ^+ [12, 13, 14, 15, 16, 17, 18]. However, no consensus has yet been seen on the quantum numbers and the structure.

The discovery of Θ^+ leads us to think of the possibility of the existence of pentaquarks with various quantum numbers. Recently, it has been found from a quark model that the pentaquark with $J^P = 3/2^-$ and $I = 1$ can exist as a low lying state for $uudd\bar{s}$ system and nearly degenerate with $J^P = 1/2^+$ or $3/2^+$ and $I = 0$ pentaquark [19]. If such a state lies much below the ΔK threshold, it can be extremely narrow, since it decays only to D -wave KN states and the width is strongly suppressed due to the high centrifugal barrier [19]. In order to ascertain the existence of the narrow pentaquark with $J^P = 3/2^-$ and $I = 1$, it is crucial to estimate its absolute mass, since the width is sensitive to the energy difference from the ΔK threshold.

In this paper, we study $J = 3/2$ and $I = 1$ pentaquark (Θ^{++}) by using the method of QCD sum rule [22], which is closely related to the fundamental theory and able to evaluate the absolute masses of hadrons without any model assumptions. In QCD sum rule approach, a correlation function of an interpolating field is calculated by the use of the operator product expansion (OPE), and is compared with the spectral representation via dispersion relation. The sum rules relate hadron properties to the vacuum expectation values of QCD operators (condensates), such as $\langle 0|\bar{q}q|0\rangle$, $\langle 0|(\alpha_s/\pi)G^2|0\rangle$ and so on.

The correlation function from which we derive the QCD sum rule is

$$\Pi_{\mu\nu}(p) = -i \int dx^4 \exp(ipx) \langle 0|T [\eta_\mu(x)\bar{\eta}_\nu(0)] |0\rangle, \quad (1)$$

where η_μ is an interpolating field for the pentaquark with $J^P = 3/2^-$ and $I = 1$. We use

the following interpolating field,

$$\eta_\mu = \epsilon_{cfdg}(\epsilon_{abc}u_a^T C \gamma_5 d_b)(\epsilon_{def}u_d^T C \gamma_\mu u_e) C \bar{s}_g^T, \quad (2)$$

where u , d and s are up, down and strange quark fields, respectively, roman indices a, b, \dots are color, C denotes charge conjugation matrix, and T transpose. $\epsilon_{abc}u_a^T C \gamma_5 d_b$ is a color $\bar{3}$ scalar diquark operator with $I=0$. $\epsilon_{def}u_d^T C \gamma_\mu u_e$ is a color $\bar{3}$ axial-vector diquark operator with $I = 1$. Thus Eq.(2) is totally $I = 1$ and it contains the state with $J^P = 3/2^-$. The way of constructing the interpolating field, Eq.(2), is based on the picture of the pentaquark structure found from the quark model calculation mentioned above [19]. According to the Ref.[19], the pentaquark with $J^P = 3/2^-$ and $I = 1$ consists of two color $\bar{3}$ diquarks and an anti-strange quark. One of the diquarks has $S = 0$ and $I = 0$ and the other $S = 1$ and $I = 1$. Evidently, the interpolating field, Eq.(2), possesses the same diquark structure.

The correlation function, Eq.(1), has various tensor structures,

$$\Pi_{\mu\nu}(p) = g_{\mu\nu} \not{p} \Pi_1(p^2) + g_{\mu\nu} \Pi_2(p^2) + \gamma_\mu \gamma_\nu \Pi_3(p^2) + \dots \quad (3)$$

We are interested in the terms proportional to $g_{\mu\nu}$:

$$\Pi(p) \equiv \not{p} \Pi_1(p^2) + \Pi_2(p^2), \quad (4)$$

since these terms receive the contribution of pure $J = 3/2$ states. In the other terms, $J = 1/2$ states contribute as well as $J = 3/2$ states [23].

We can relate the correlation function with the spectral function via Lehman representation,

$$\Pi(p_0, \mathbf{p}) = \int_{-\infty}^{\infty} \frac{\rho(p'_0, \mathbf{p})}{p_0 - p'_0} dp'_0, \quad (5)$$

where $\rho(p_0, \mathbf{p})$ is the spectral function. On the other hand, in the deep Euclid region, $p_0^2 \rightarrow -\infty$, the correlation function can be evaluated by an operator product expansion. Then the correlation function is expressed as a sum of various vacuum condensates. Using the analyticity, we obtain a relation between the imaginary part of the correlation function evaluated by an OPE, ρ^{OPE} , and the spectral function as

$$\int_{-\infty}^{\infty} dp_0 \rho^{\text{OPE}}(p_0, \mathbf{p}) W(p_0) = \int_{-\infty}^{\infty} dp_0 \rho(p_0, \mathbf{p}) W(p_0), \quad (6)$$

where $W(p_0)$ is an analytic function of p_0 . Eq.(6) is a general form of the QCD sum rule. By properly parameterizing $\rho(p_0, \mathbf{p})$, we obtain QCD sum rules for physical quantities in $\rho(p_0, \mathbf{p})$.

Let us first consider the spectral function, $\rho(p_0, \mathbf{p})$. The interpolating field couples to the states whose parity is opposite to that of the interpolating field, as well as the states with the same parity of the interpolating field [24]. Therefore, in the zero-width approximation, Eq.(4) is expressed as

$$\Pi(p) = \sum_n \left[|\lambda_-^n|^2 \frac{\not{p} + m_-^n}{p^2 - m_-^n{}^2} + |\lambda_+^n|^2 \frac{\not{p} - m_+^n}{p^2 - m_+^n{}^2} \right], \quad (7)$$

where m_{\mp}^n are the masses of negative and positive parity states, λ_{\mp}^n are the coupling strengths of the interpolating field with negative and positive parity states, respectively. The spectral function in the rest frame, $\mathbf{p} = \mathbf{0}$, can be decomposed into two parts as follows,

$$\rho(p_0) = P_- \rho_-(p_0) + P_+ \rho_+(p_0), \quad (8)$$

where $P_{\mp} = (\gamma_0 \pm 1)/2$ and $\rho_{\mp}(p_0)$ are given by

$$\rho_-(p_0) = \sum_n \left[|\lambda_-^n|^2 \delta(p_0 - m_-^n) + |\lambda_+^n|^2 \delta(p_0 + m_+^n) \right], \quad (9)$$

$$\rho_+(p_0) = \sum_n \left[|\lambda_-^n|^2 \delta(p_0 + m_-^n) + |\lambda_+^n|^2 \delta(p_0 - m_+^n) \right]. \quad (10)$$

Next, we construct the sum rule for negative parity states and that for positive parity. We apply the projection operator P_{\mp} to Eq.(6) for $\mathbf{p} = \mathbf{0}$. Then we obtain

$$\int_{-\infty}^{\infty} dp_0 \rho_{\mp}^{\text{OPE}}(p_0) W(p_0) = \int_{-\infty}^{\infty} dp_0 \rho_{\mp}(p_0) W(p_0). \quad (11)$$

Note that in Eq.(11) the contribution from the positive and negative parity states are not decoupled, since, as can be seen from Eqs.(9) and (10), each of $\rho_-(p_0)$ and $\rho_+(p_0)$ contains the contribution from both of the parity states. What we want to do is to separate the negative or positive parity contribution from Eqs.(11). (The following procedure is essentially equivalent to that in Ref.[25].)

If $\rho_{\mp}^{\text{OPE}}(p_0)$ are separable into $\rho_{\mp}^{\text{OPE}}(p_0 > 0)$ and $\rho_{\mp}^{\text{OPE}}(p_0 < 0)$, we can separate Eq.(11) into the contributions from $p_0 > 0$ and $p_0 < 0$. From the positive energy part, we obtain

$$\int_0^{\infty} dp_0 \rho_-^{\text{OPE}}(p_0) W(p_0) = \int_0^{\infty} dp_0 \rho_-(p_0) W(p_0), \quad (12)$$

$$\int_0^{\infty} dp_0 \rho_+^{\text{OPE}}(p_0) W(p_0) = \int_0^{\infty} dp_0 \rho_+(p_0) W(p_0). \quad (13)$$

Here we notice that only the negative (positive) parity states contribute to $\rho_-(p_0 > 0)$ ($\rho_+(p_0 > 0)$) (see Eqs.(9) and (10)). Eq.(12) is therefore the sum rule for the negative parity states and Eq.(13) is that for the positive parity states.

A comment is in order here. In order to separate the negative and positive parity states in the sum rule as Eqs.(12) and (13), it is necessary that $\rho_{\mp}^{\text{OPE}}(p_0)$ are separable into $\rho_{\mp}^{\text{OPE}}(p_0 > 0)$ and $\rho_{\mp}^{\text{OPE}}(p_0 < 0)$ as mentioned above. In general, $\rho_{\mp}^{\text{OPE}}(p_0)$ are not separable [26]. However, as will be seen below (Eqs.(18), (19) and (20)), $\rho_{\mp}^{\text{OPE}}(p_0)$ for pentaquark is separable as long as we truncate the OPE at certain order, since $\rho_{\mp}^{\text{OPE}}(p_0)$ up to dimension 6 operator have the p_0 dependence as $p_0^n [\theta(p_0) - \theta(-p_0)]$. We thus derive the sum rule for each parity state of the pentaquark as Eqs.(12) and (13).

We parameterize $\rho_{\mp}(p_0)$ with a pole plus continuum contribution,

$$\rho_{\mp}(p_0) = |\lambda_{\mp}|^2 \delta(p_0 - m_{\mp}) + |\lambda_{\pm}|^2 \delta(p_0 + m_{\pm}) + [\theta(p_0 - \omega_{\mp}) + \theta(-p_0 - \omega_{\pm})] \rho^{\text{OPE}}(p_0) \quad (14)$$

Substituting Eq.(14) into the right-hand sides of Eqs.(12) and (13), we obtain the following sum rules,

$$\int_0^{\omega_{\mp}} dp_0 \rho_{\mp}^{\text{OPE}}(p_0) p_0^n \exp(-\frac{p_0^2}{M^2}) = m_{\mp}^n |\lambda_{\mp}|^2 \exp(-\frac{m_{\mp}^2}{M^2}) \quad (15)$$

Here we have chosen the weight function as $W(p_0) = p_0^n \exp(-p_0^2/M^2)$. The parameter M is called Borel mass. From Eq.(15) for $n = 0$, we obtain the sum rule for the pole residues $|\lambda_{\mp}|^2$,

$$|\lambda_{\mp}|^2 \exp(-\frac{m_{\mp}^2}{M^2}) = \int_0^{\omega_{\mp}} dp_0 \rho_{\mp}^{\text{OPE}}(p_0) \exp(-\frac{p_0^2}{M^2}). \quad (16)$$

The ratio of Eq.(15) for $n = 0$ and $n = 2$ gives the sum rules for the masses,

$$m_{\mp}^2 = \frac{\int_0^{\omega_{\mp}} dp_0 \rho_{\mp}^{\text{OPE}}(p_0) p_0^2 \exp(-\frac{p_0^2}{M^2})}{\int_0^{\omega_{\mp}} dp_0 \rho_{\mp}^{\text{OPE}}(p_0) \exp(-\frac{p_0^2}{M^2})}. \quad (17)$$

Let us now turn to the OPE. We have taken into account the terms up to dimension 6 operator. We show the result of the OPE,

$$\rho^{\text{OPE}}(p_0) = \gamma_0 A(p_0) + B(p_0), \quad (18)$$

where $A(p_0)$ and $B(p_0)$ are given by

$$A(p_0) = \left[\frac{1}{5^2 \cdot 3^2 \cdot 2^{18} \pi^8} (p_0)^{11} + \frac{-1}{5 \cdot 3^4 \cdot 2^{17} \pi^6} \langle 0 | \frac{\alpha_s}{\pi} G^{a\mu\nu} G_{\mu\nu}^a | 0 \rangle (p_0)^7 \right]$$

$$\begin{aligned}
& + \frac{1}{3^3 \cdot 2^{12} \pi^6} m_s \langle 0 | \bar{s} s | 0 \rangle (p_0)^7 + \frac{-1}{5 \cdot 3^2 \cdot 2^{10} \pi^6} m_s g \langle 0 | \bar{s} \sigma^{\mu\nu} (\lambda^a / 2) G_{\mu\nu}^a s | 0 \rangle (p_0)^5 \\
& + \frac{1}{5 \cdot 2^8 \pi^4} \langle 0 | \bar{q} q | 0 \rangle^2 (p_0)^5 \Big] \times [\theta(p_0) - \theta(-p_0)], \tag{19}
\end{aligned}$$

$$\begin{aligned}
B(p_0) = & \left[\frac{1}{7 \cdot 5^2 \cdot 3 \cdot 2^{14} \pi^8} m_s (p_0)^{10} + \frac{-1}{3^3 \cdot 2^{13} \pi^6} \langle 0 | \bar{s} s | 0 \rangle (p_0)^8 \right. \\
& \left. + \frac{1}{5 \cdot 3^2 \cdot 2^{10} \pi^6} g \langle 0 | \bar{s} \sigma^{\mu\nu} (\lambda^a / 2) G_{\mu\nu}^a s | 0 \rangle (p_0)^6 \right] \times [\theta(p_0) - \theta(-p_0)], \tag{20}
\end{aligned}$$

In Eqs.(19) and (20), $q = u, d$, m_s is the strange quark mass, and $\langle 0 | \mathcal{O} | 0 \rangle$ denotes the vacuum expectation value of the operator \mathcal{O} .

Here, before deriving the QCD sum rules for Θ^{++} , we comment on the contribution of the continuum states. In Ref.[27], it was pointed out that pentaquark correlation functions receive contribution of two-hadron-reducible (2HR) diagrams, which represent baryons and mesons propagating independently without interacting with each other. The 2HR diagrams are related only with the background (continuum states). We can make the background contribution in the sum rules as small as possible by subtracting the 2HR diagrams. It is better to subtract them especially in the sum rules for $J^P = 1/2^\pm$ pentaquark, where the NK continuum contribution should be significant. However, in the sum rules for $J^P = 3/2^-$ pentaquark, the NK continuum contribution itself is expected to be small, since the N and K are relatively D -wave in this channel. Hence, in this paper, we consider the correlation function without subtracting the 2HR parts.

We substitute $\rho_{\mp}^{\text{OPE}}(p_0) = A(p_0) \pm B(p_0)$ with Eqs.(19) and (20) into the right hand sides of Eqs.(16) and (17). Then we obtain the QCD sum rules for Θ^{++} .

We plotted in Fig.1 the right-hand side of Eq.(16) for the negative parity state as a function of the Borel mass, M . Here and hereafter we use the standard values of the QCD parameters, $\langle 0 | \bar{q} q | 0 \rangle = (-0.23 \text{ GeV})^3$, $m_s = 0.12 \text{ GeV}$, $\langle 0 | \bar{s} s | 0 \rangle = 0.8 \langle 0 | \bar{q} q | 0 \rangle$, $g \langle 0 | \bar{s} \sigma^{\mu\nu} (\lambda^a / 2) G_{\mu\nu}^a s | 0 \rangle = (0.8 \text{ GeV}^2) \langle 0 | \bar{s} s | 0 \rangle$, $\langle 0 | \frac{\alpha_s}{\pi} G^{a\mu\nu} G_{\mu\nu}^a | 0 \rangle = (0.33 \text{ GeV})^4$. As can be seen, the right-hand side of Eq.(16) is positive. This implies that the present QCD sum rule does not exclude the possibility of the existence of the Θ^{++} with negative parity since the left-hand side of Eq.(16) must be positive.

In Fig.2, we plotted the mass of Θ^{++} with negative parity against the Borel mass which is obtained from Eq.(17). We see that the dependence on the Borel mass is weak. This implies that the sum rule works well. However, the result depends on the choice of the continuum threshold ω_- . The continuum mainly comes from the S -wave ΔK scattering states, whose

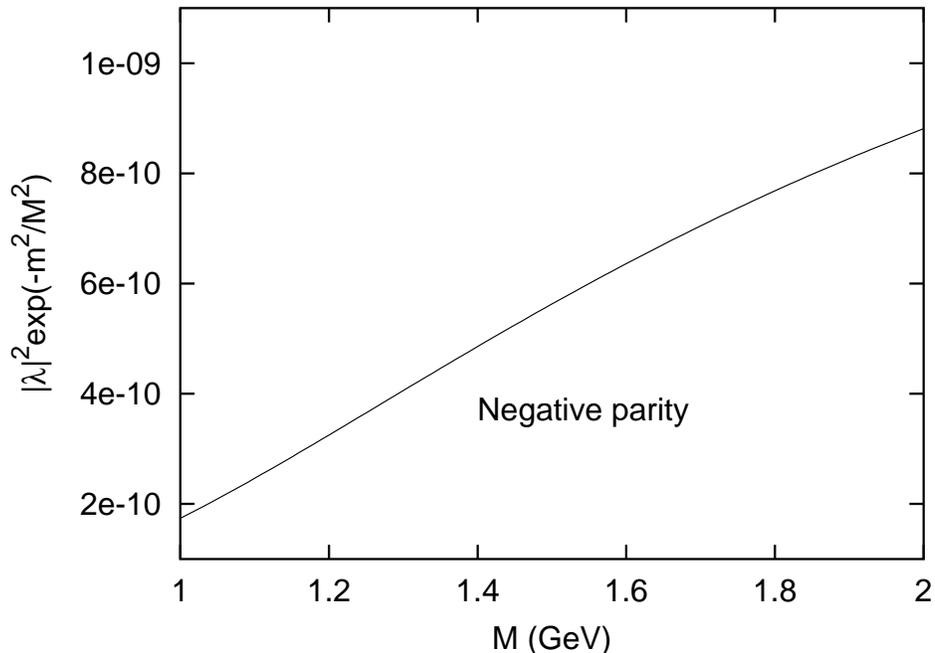


FIG. 1: $|\lambda_-|^2 \exp(-\frac{m_\pi^2}{M^2})$ as functions of Borel mass, M with the continuum threshold parameter $\omega_- = 1.8 \text{ GeV}$.

threshold is 1.73 GeV , since NK scattering state must be D -wave in this channel and it starts up very gradually. Thus we choose $\omega_- = 1.73, 1.8, 1.9 \text{ GeV}$. From the stabilized region of the curve, we predict the mass to be $1.5 \sim 1.6 \text{ GeV}$, which is close to the observed Θ^+ mass. The mass is much below the ΔK threshold. This implies that the Θ^{++} with $J^P = 3/2^-$ can be extremely narrow since it is allowed to decay only to D -wave KN state and the width is strongly suppressed due to the large centrifugal barrier.

It is remarkable that such a high spin and isovector state can be a low lying state, which is not the case for ordinary baryons. The possibility of $J^P = 3/2^-, I = 1$ pentaquark being low lying state has been suggested by the previous calculation from a quark model [19]. In Ref.[19], a simple quark model in which constituent quarks interact via one-gluon exchange force at short distances and confining (or string) potential at long distances was considered. A $q^4\bar{q}$ system has a connected string configuration corresponding to a confined state, in addition to an ordinary meson-baryon like configuration. A variational method called antisymmetrized molecular dynamics (AMD) [20, 21] was applied to the confined $uudd\bar{s}$ system and all the possible spin parity states were calculated. The narrow and low lying states they have found are $J^P = 1/2^+$ or $3/2^+$ with $I = 0$ and $J^P = 3/2^-$ with $I = 1$

states. The former has just the same structure as that conjectured by Jaffe and Wilczek [12]. We represent it as $[ud]_{S=0,I=0}[ud]_{S=0,I=0}[\bar{s}]$, where $[ud]_{S,I}$ denotes a color $\bar{\mathbf{3}}$ ud -diquark with spin S and isospin I . Both of the two diquarks gain color magnetic interaction since they have $S = 0$. However, this state loses the kinetic and string energy, since the two diquarks, which are to be antisymmetric in color, are identical and must be relatively P -wave. In Ref.[19], another energetically favorable state has been predicted, which consists of an $S = 0$ diquark and an $S = 1$ diquark: $[ud]_{S=0,I=0}[ud]_{S=1,I=1}[\bar{s}]$. This state is totally $J^P = 3/2^-$ and $I = 1$. It loses color magnetic interaction since one of the diquarks has $S = 1$. However, it gains kinetic and string energy, since the two diquarks are no longer identical and they can be relatively S -wave. Owing to the balance between the energy gain and loss, $J^P = 3/2^-$, $I = 1$ state degenerate with $J^P = 1/2^+$ or $3/2^+$, $I = 0$ state. Within the quark model employed in Ref.[19], however, one cannot predict the absolute masses but only the level structure of the pentaquarks, because this quark model relies on the zero-point energy of the confining potential. In Ref.[19], it was adjusted to reproduce the observed mass of Θ^+ . Whereas, the QCD sum rule is able to estimate the absolute mass. We confirmed from the QCD sum rule that the $J^P = 3/2^-$, $I = 1$ state actually can be a low lying state, using the interpolating field, Eq.(2), which has the same structure as that suggested by the quark model.

The pentaquark with $J^P = 3/2^-$ and $I = 1$ has also been found from the chiral unitary approach, as a resonance state in the ΔK channel [29]. This state is generated due to an attractive interaction in that channel existing in the lowest order chiral Lagrangian. The attractive interaction leads to a pole of the complex energy plane and manifests itself in a large strength of the ΔK scattering amplitude with $L = 0$ and $I = 1$. We note that the interpolating field, Eq.(2), can also couple with such a ΔK resonance states because it contains the ΔK component as is shown by Fierz transformation.

Let us turn to the sum rule for the positive parity state. We plotted in Fig.3 the right-hand side of Eq.(16) for the positive parity state as a function of the Borel mass, M . The right-hand side of Eq.(16) is positive, which implies that the existence of the positive parity state is not excluded. The mass against the Borel mass is shown in Fig.4. The continuum in this channel mainly comes from the P -wave NK scattering states. We choose $\omega_+ = 1.7, 1.8, 1.9$ GeV. Although the curve depends on the choice of the continuum threshold parameter, we can say that the positive parity state nearly degenerate with the negative parity state.

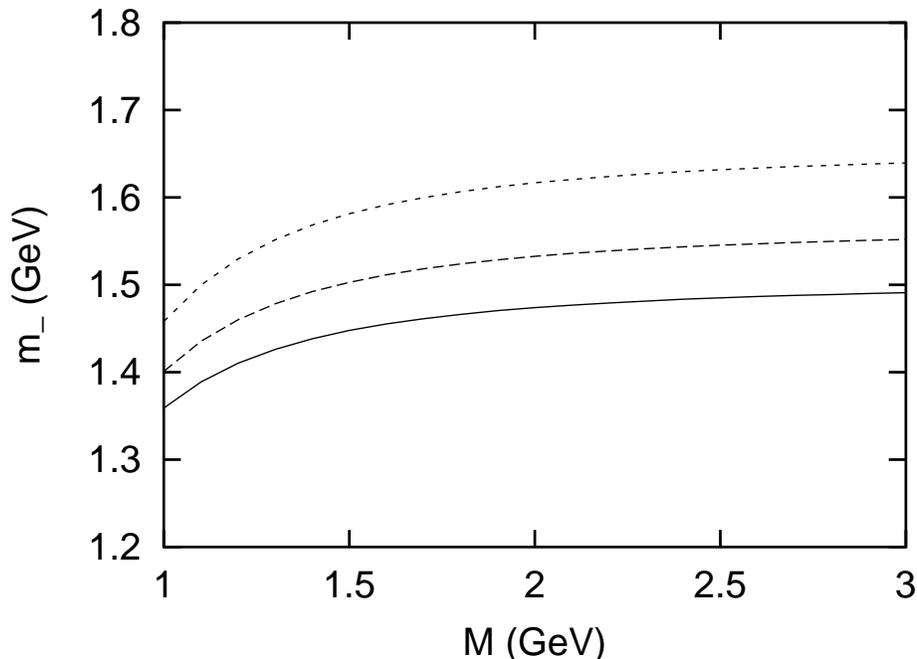


FIG. 2: Mass of $J^P = 3/2^-, I = 1$ pentaquark as a function of Borel mass, M , with the continuum threshold parameters $\omega_- = 1.73$ GeV (solid line), 1.8 GeV (dashed), 1.9 GeV (dotted).

However, the positive parity state is expected to be broader than the negative parity state, since the former can decay into P -wave NK states while the latter only to D -wave NK states. The present result is consistent with a recent calculation by Skyrme model [30]. The authors in Ref.[30] predicted that there exists a new isotriplet of Θ -baryons with $J^P = 3/2^+$ and $I = 1$. Its mass is 1595 MeV and the width is large: $\Gamma \sim 80$ MeV.

In summary, we have studied $J = 3/2, I = 1$ pentaquark, Θ^{++} , using the method of QCD sum rule. We used the interpolating field constructed from a color anti-triplet scalar isoscalar diquark, a color anti-triplet axial-vector isovector diquark and an anti-strange quark. We have derived the QCD sum rules for the negative and positive parity states. QCD sum rule predicts a narrow Θ^{++} ($J^P = 3/2^-$). Its mass is predicted to be $1.5 \sim 1.6$ GeV, which is much below the ΔK threshold. Since only the D -wave decay to NK channel is allowed, it should be an extremely narrow state. QCD sum rule also shows the possibility of the existence of the $J^P = 3/2^+$ state. It nearly degenerates with the negative parity state. It may be broader than the negative parity state, since it is allowed to decay into P -wave NK state. It is worth mentioning that this is the first QCD sum rule analysis of high spin states of the pentaquark. Most of the works using QCD sum rules and Lattice QCD concentrate

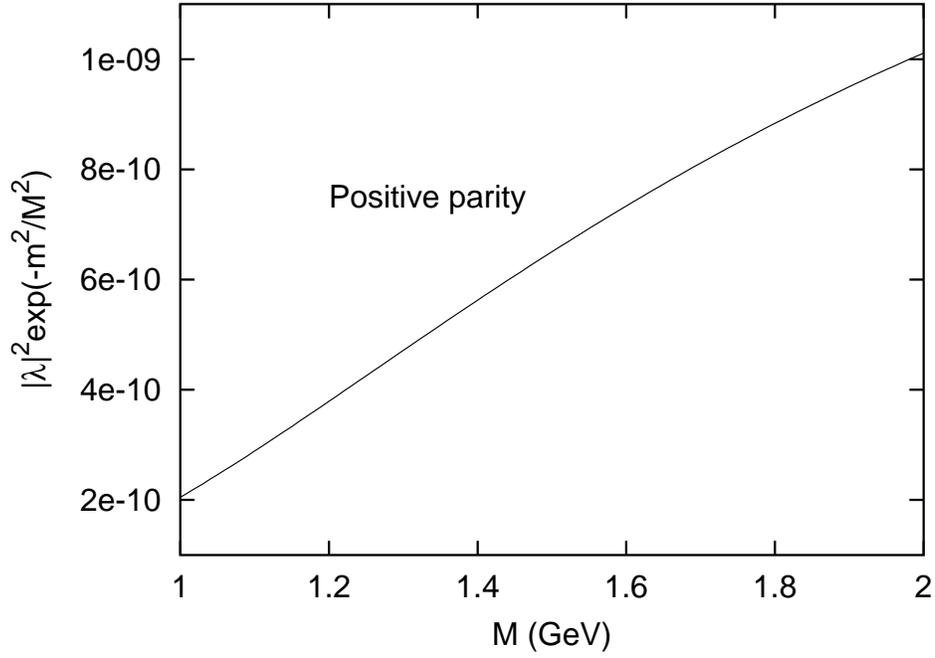


FIG. 3: $|\lambda_+|^2 \exp(-\frac{m_\mp^2}{M^2})$ as functions of Borel mass, M with the continuum threshold parameter $\omega_+ = 1.8$ GeV.

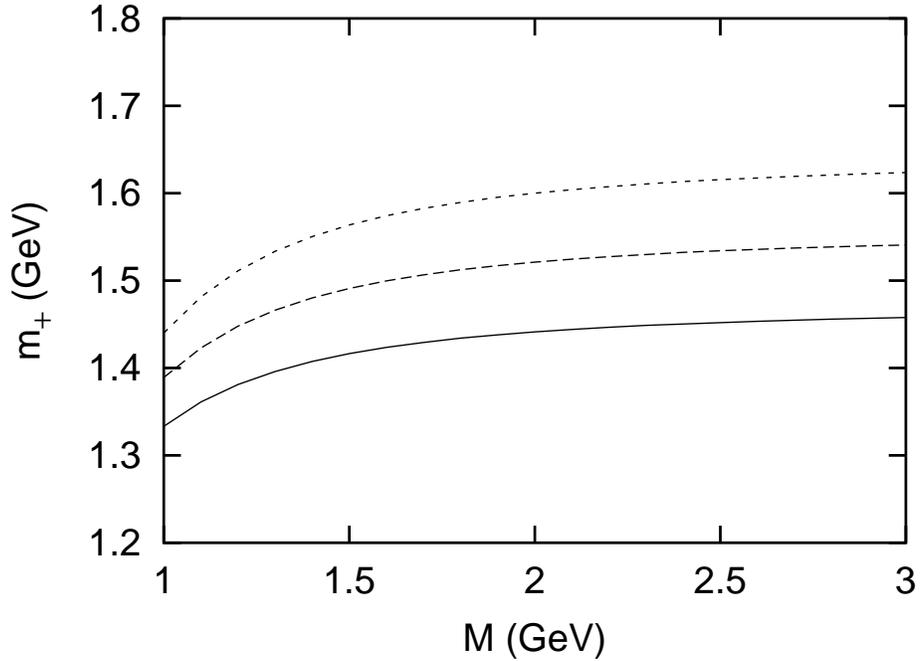


FIG. 4: Mass of $J^P = 3/2^+$, $I = 1$ pentaquark as a function of Borel mass, M with the continuum threshold parameters $\omega_+ = 1.7$ GeV (solid line), 1.8 GeV (dashed), 1.9 GeV (dotted).

on $J = 1/2$ and $I = 0$ pentaquark states. It would be interesting to see if lattice calculation could confirm these findings.

-
- [1] LEPS collaboration, T. Nakano *et al.* , Phys. Rev. Lett. **91**, 012002 (2003).
 - [2] DIANA Collaboration, V. V. Barmin *et al.* , Phys. Atom. Nucl. **66**, 1715 (2003).
 - [3] CLAS Collaboration, S. Stepanyan *et al.* , Phys. Rev. Lett. **91**, 252001 (2003).
 - [4] CLAS Collaboration, V. Kubarovsky *et al.* , Phys. Rev. Lett. **92**, 032001 (2004)
 - [5] SAPHIR Collaboration, J. Barth *et al.* , Phys. Lett. **B572**, 127 (2003)
 - [6] A. E. Asratyan, A. G. Dolgolenko and M. A. Kubantsev, hep-ex/0309042.
 - [7] HERMES Collaboration, A. Airapetian *et al.* , hep-ex/0312044.
 - [8] SVD Collaboration, A. Aleev *et al.* , hep-ex/0401024.
 - [9] ZEUS Collaboration, S. V. Chekanov, hep-ex/0404007.
 - [10] CLAS Collaboration, H. G. Juengst, nucl-ex/0312019.
 - [11] D. Diakonov, V. Petrov and M. V. Polyakov, Z. Phys. **A359**, 305 (1997).
 - [12] R. Jaffe and F. Wilczek, Phys. Rev. Lett. **91**, 232003 (2003).
 - [13] S. Capstick, P. R. Page and W. Roberts, Phys. Lett. **B570**, 185 (2003).
 - [14] M. Karliner and H. J. Lipkin, Phys. Lett. **B575**, 249 (2003).
 - [15] S. Sasaki, hep-lat/0310014.
 - [16] Shi-Lin Zhu, Phys. Rev. Lett. **91**, 232002 (2003).
 - [17] R. D. Matheus, F. S. Navarra, M. Nielsen, R. Rodrigues da Silva and S. H. Lee, Phys. Lett. **B578**, 323 (2004)
 - [18] J. Sugiyama, T. Doi and M. Oka, Phys. Lett. **B581**, 167 (2004).
 - [19] Y. Kanada-En'yo, O. Morimatsu and T. Nishikawa, hep-ph/0404144.
 - [20] Y. Kanada-En'yo, H. Horiuchi and A. Ono, Phys. Rev. **C52**, 628 (1995); Y. Kanada-En'yo and H. Horiuchi, Phys. Rev. **C52**, 647 (1995).
 - [21] Y. Kanada-En'yo, M. Kimura and H. Horiuchi, Comptes Rendus Physique Vol.4, 497(2003).
 - [22] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. **B147**, 385 (1979); Nucl. Phys. **B147**, 448 (1979).
 - [23] B. L. Ioffe, Nucl. Phys. **B188**, 317 (1981).
 - [24] Y. Chung, H. G. Dosch, M. Kremer and D. Schall, Nucl. Phys. **B197**, 55 (1982).

- [25] D. Jido, N. Kodama and M. Oka, *Phys. Rev.* **D54**, 4532 (1996).
- [26] Y. Kondo, O. Morimatsu and T. Nishikawa, in preparation
- [27] Y. Kondo, O. Morimatsu and T. Nishikawa, [hep-ph/0404285](#).
- [28] Talks by S.H. Lee and M. Oka in the international workshop “PENTAQUARK04” held on July 20-23, 2004 at SPring-8, Japan.
- [29] S. Sarkar, E. Oset and M.J. V. Vacas, [nucl-th/0404023](#); [nucl-th/0407025](#).
- [30] D. Borisyuk, M. Faber and A. Kobushkin, [hep-ph/0307370](#).