Analysis of the $X(4350)$ as a Scalar $\bar{c}c$ and $D_s^*\bar{D}_s^*$ Mixing State with QCD Sum Rules

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Abstract

In this article, we assume that the narrow structure $X(4350)$ is a scalar $\bar{c}c$ – $D_s^*\bar{D}_s^*$ mixing state, and study its mass using the QCD sum rules. The numerical result $M_X = (4.37 \pm 0.15)$ GeV is in good agreement with the experimental data, the $X(4350)$ may be a scalar $\bar{c}c$ -- $D_s^*\bar{D}_s^*$ mixing state. Other possibility, such as a scalar (tensor) $cs\bar{c}\bar{s}$ tetraquark state is not excluded.

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1 Introduction

In 2009, the CDF collaboration observed a narrow structure (which is denoted as the $Y(4140)$ now) near the $J/\psi\phi$ threshold with statistical significance in excess of 3.8 standard deviations in exclusive decays $B^+ \rightarrow J/\psi\phi K^+$ produced in $\bar{p}p$ collisions at $\sqrt{s} = 1.96$ TeV [1]. The mass and width are $(4143.0 \pm 2.9 \pm 1.2)$ MeV and $(11.7^{+8.8}_{-5.0} \pm 3.7)$ MeV respectively. The narrow structure $Y(4140)$ is very similar to the charmonium-like state $Y(3930)$ near the $J/\psi\omega$ threshold [2, 3]. The mass and width of the $Y(3930)$ are $(3914.6^{+3.8}_{-3.4} \pm 2.0)$ MeV and $(34^{+12}_{-8} \pm 5)$ MeV respectively [3].

There have been several explanations for the nature of the narrow structure $Y(4140)$, such as a $D_s^*\bar{D}_s^*$ molecular state [4, 5, 6, 7, 8, 9, 10], an exotic ($J^{PC} = 1^{-+}$) hybrid charmonium [5], a $c\bar{c}s\bar{s}$ tetraquark state [11], the effect of the $J/\psi\phi$ threshold [12], or none a conventional charmonium state [13] nor a scalar $D_s^*\bar{D}_s^*$ molecular state [14, 15], etc. Assuming the $Y(4140)$ is a $D_s^*\bar{D}_s^*$ molecular state with $J^{PC} = 0^{++}$ or $2^{++}$, Branz et al predict its two-photon decay width is of order 1 KeV [6].

Recently, the Belle collaboration measured the process $\gamma\gamma \rightarrow \phi J/\psi$ for the $\phi J/\psi$ invariant mass distributions between the threshold and 5 GeV based on a data sample of 825 fb$^{-1}$, and observed a narrow peak of $8.8^{+4.2}_{-3.2}$ events with a significance of 3.2 standard deviations [16]. The mass and width of the structure (denoted as $X(4350)$) are $(4350.6^{+4.6}_{-5.1} \pm 0.7)$ MeV and $(13.3^{+12.9}_{-9.9} \pm 4.1)$ MeV respectively. No signal for the $Y(4140) \rightarrow \phi J/\psi$ structure was observed, this disfavors the scenario of the $Y(4140)$ as a $D_s^*\bar{D}_s^*$ molecular state.

The possible quantum numbers for a state $X$ decaying into $J/\psi\phi$ are $J^{PC} = 0^{-+}, 0^{++}, 1^{-+}, 2^{++}$, the corresponding strong interactions can be described by the

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following phenomenological Lagrangian,

\[ \mathcal{L}_{0^{-+}} = g \epsilon^{\mu \nu \alpha \beta} \partial_\mu \psi_\nu \partial_\alpha \phi_\beta X, \]
\[ \mathcal{L}_{1^{--}} = g \epsilon^{\mu \nu \alpha \beta} (\psi_\mu \partial_\nu \phi_\alpha - \phi_\mu \partial_\nu \psi_\alpha) X_\beta, \]
\[ \mathcal{L}_{0^{++}} = g \psi_\mu \phi^\mu X, \]
\[ \mathcal{L}_{2^{++}} = g X^{\mu \nu} \psi_\mu \phi_\nu, \]

(1)

the strong coupling constants \( g \) can be fitted phenomenologically or calculated by some theoretical approaches, for example, the QCD sum rules.

In Ref.\[9\], Zhang and Huang study the \( Q \bar{s} Q \) and \( Q \bar{s} Q' \) molecular states in a systematic way using the QCD sum rules before the Belle experiment, the mass of the \( D_s^*+ \bar{D}_{s0}^- \) molecular state is \((4.36 \pm 0.08) \text{ GeV}\), which is consistent with experimental data \[16\]. Such a state has \( J^P = 1^- \) and no definite charge conjugation. In Ref.\[17\], Albuquerque et al re-study the exotic \( D_s^*+ \bar{D}_{s0}^- \bar{D}_{s0}^- \bar{D}_{s0}^* \) molecular state with \( J^{PC} = 1^{--} \) by taking into account the contributions from the vacuum condensates up to dimension-8, the prediction \( M_{D_s^* \bar{D}_{s0}} = (5.05 \pm 0.19) \text{ GeV} \) disfavors the scenario of the \( X(4350) \) as a \( D_s^* \bar{D}_{s0} \) molecular state.

In Ref.\[18\], Drenska et al study the exotic tetraquark states of the kind \( c s \bar{c} \bar{s} \) by computing their spectrum and decay modes within a constituent diquark-antidiquark model, the predictions \( M_{0^{++}} = 4277, 4312 \text{ MeV} \) and \( M_{1^{--}} = 4321, 4356 \text{ MeV} \) are consistent with the experimental data \[16\]. On the other hand, the flux-tube model \[19, 20\] and the Lattice QCD \[21, 22, 23\] predict that the masses of the low lying hybrid charmonium states are about \((4.0 \sim 4.2) \text{ GeV} \) and \((4.0 \sim 4.4) \text{ GeV} \) respectively, which are also consistent with the experimental data \[16\]. However, the decay of a hybrid to two photons is generically forbidden \[24\].

In Ref.\[11\], Stancu study the mass spectrum of the \( c \bar{c} s \bar{s} \) tetraquarks using a simple quark model with chromo-magnetic interaction and observe that the \( Y(4140) \) may be the strange partner of the \( X(3872) \), the prediction for the mass of the \( 2^{++} \) tetraquark state is consistent with the \( X(4350) \). As noticed by the author, the amplitude of the singlet-singlet component seems too large comparing with the octet-octet component.

In Ref.\[25\], Liu et al discuss the possibility that the \( X(4350) \) is an excited \( P \)-wave charmonium state \( \chi''_{c_2} \) by studying the strong decays of the \( P \)-wave charmonium states with the \( 3P_0 \) model.

The CDF and Belle collaborations analyze the experimental data by assuming the vector mesons \( J/\psi \) and \( \phi \) have a relative \( S \)-wave \[1 \], so we will not focus on the scenarios of the \( X(4350) \) as the \( 0^{-+} \) and \( 1^{--} \) tetraquark state or hybrid charmonium.

In Refs.\[26, 27\], we study the mass spectrum of the scalar hidden charm and hidden bottom tetraquark states which consist of the scalar-scalar type, axial-vector-axial-vector type and vector-vector type diquark pairs in a systematic way using the QCD sum rules, the scalar-scalar type and axial-vector-axial-vector type scalar \( c \bar{c} s \bar{s} \) tetraquark states have masses about \((4.45 \pm 0.16) \text{ GeV} \), the lower bound of the
Table 1: The possible explanations for the nature of the $X(4350)$, where the RW stands for the relative wave of the final state mesons.

masses is consistent with the $X(4350)$, we cannot exclude that the $X(4350)$ is a scalar $c\bar{c}s\bar{s}$ tetraquark state. In Refs. [14, 15], we study the $D^*_s\bar{D}^*_s$, $B^*\bar{B}^*$ and $B^*_s\bar{B}^*_s$ molecular states in a systematic way using the QCD sum rules. The numerical result is inconsistent with the experimental data, the $D^*_s\bar{D}^*_s$ is probably a virtual state and not related to the meson $Y(4140)$. In this article, we study the $X(4350)$ as a linear superposition of a scalar charmonium state $c\bar{c}$ and a virtual state $D^*_s\bar{D}^*_s$ using the QCD sum rules [28, 29]. In Table 1, we present the possible explanations for the nature of the $X(4350)$.

The article is arranged as follows: we derive the QCD sum rules for the narrow structure $X(4350)$ in Sect. 2; in Sect. 3, we present the numerical results and discussions; and Sect. 4 is reserved for our conclusions.

2 QCD sum rules for the $X(4350)$ as a mixing state

In the following, we write down the two-point correlation function $\Pi(p)$ in the QCD sum rules,

$$\Pi(p) = i \int d^4 x e^{ipx} \langle 0 | T \{ J(x) J^\dagger(0) \} | 0 \rangle,$$

$$J(x) = \frac{J_1(x) + J_2(x)}{\sqrt{2}},$$

$$J_1(x) = \bar{c}(x) \gamma_\mu s(x) \bar{s}(x) \gamma^\mu c(x),$$

$$J_2(x) = -\frac{\langle \bar{s}s \rangle}{3} \bar{c}(x) c(x),$$

where the $J_2(x)$ is the normalized two-quark current [30].

We can insert a complete set of intermediate hadronic states with the same quantum numbers as the current operator $J(x)$ into the correlation function $\Pi(p)$
to obtain the hadronic representation \cite{28, 29}. After isolating the ground state contribution from the pole term of the lowest state \(X\), we get the following result,

\[
\Pi(p) = \frac{\lambda_X^2}{M_X^2 - p^2} + \cdots ,
\]

(4)

where the pole residue (or coupling) \(\lambda_X\) is defined by

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

(5)

The two-quark current \(\bar{c}(x)c(x)\) has non-vanishing coupling with the charmonia \(\chi_{c0}(1P), \chi_{c0}(2P), \chi_{c0}(3P), \cdots\); while the molecule type current \(\bar{c}(x)\gamma_\mu s(x)\bar{s}(x)\gamma^\mu c(x)\) has non-vanishing coupling with the molecular states \(\bar{D}_s^* \bar{D}_s^*, \bar{D}_s^* \bar{D}_s^*, \bar{D}_s^* \bar{D}_s^*, \cdots\) and the scattering states \(\bar{D}_s^* - \bar{D}_s^*, \bar{D}_s^* - \bar{D}_s^*, \bar{D}_s^* - \bar{D}_s^*, \cdots\) \cite{31}. We cannot distinguish those contributions and study them exclusively. In this article, we take the assumption that the interpolating current \(J(x)\) couples to a particular resonance, which is a special superposition of the scalar charmonia \(\chi_{c0}(1P), \cdots\) and the virtual molecular states \(\bar{D}_s^* \bar{D}_s^*, \cdots\). In other words, we take a single pole approximation, the pole embodies the net effects.

We carry out the operator product expansion for the correlation function \(\Pi(p)\) at the large space-like momentum region \(p^2 \ll 0\),

\[
\Pi(p) = \frac{1}{2} \Pi_{11}(p) + \frac{\langle \bar{s}s \rangle^2}{6} \Pi_{22}(p) ,
\]

(6)

where

\[
\Pi_{11}(p) = i \int d^4xe^{ip\cdot x} \langle 0|T \left\{ J_1(x) J_1^\dagger(0) \right\} |0\rangle = \int_{\Delta^2}^{s_0} ds \frac{\rho_{11}(s)}{s - p^2} + \cdots ,
\]

\[
\Pi_{22}(p) = i \int d^4xe^{ip\cdot x} \langle 0|T \left\{ J_2(x) J_2^\dagger(0) \right\} |0\rangle = \int_{\Delta^2}^{s_0} ds \frac{\rho_{22}(s)}{s - p^2} + \cdots ,
\]

\[
\rho_{22}(s) = \frac{9}{4\pi^2} \int_{x_i}^{x_f} dx (1 - x)(s - \tilde{m}_c^2)
\]

\[
+ \frac{1}{8} \frac{\alpha_s GG}{\pi} \int_0^1 dx \left[ 1 - \frac{(x^2 - x + 1)\tilde{m}_c^2}{x(1-x)M^2} \right] \delta(s - \tilde{m}_c^2) ,
\]

(7)

the explicit expression of the spectral density \(\rho_{11}(s)\) can be found in Refs.\cite{14, 15}. \(\Delta^2 = 4(m_c + m_s)^2, \tilde{m}_c^2 = \frac{m_c^2}{x(1-x)}\), \(x_f = \left(1 + \sqrt{1 - \frac{4m_c^2}{s}}\right)/2\), \(x_i = \left(1 - \sqrt{1 - \frac{4m_c^2}{s}}\right)/2\).

In calculation, we use the Fierz re-ordering in the color space and Dirac spin space to express the correlation functions \(\Pi_{12}(p)\) and \(\Pi_{21}(p)\) in terms of the \(\Pi_{22}(p)\). In this article, we carry out the operator product expansion to the vacuum condensates adding up to dimension-10 and take the assumption of vacuum saturation for the high dimensional vacuum condensates, they are always factorized to lower condensates with vacuum saturation in the QCD sum rules, and factorization works well in large \(N_c\) limit.
Once analytical result is obtained, then we can take the quark-hadron duality and perform the Borel transform with respect to the variable \( P^2 = -p^2 \), finally we obtain the following sum rule:

\[
\lambda_X^2 e^{\frac{M_X^2}{M^2}} = \int_{\Delta^2}^{s_0} ds \left[ \frac{1}{2} \rho_{11}(s) + \frac{(\bar{s}s)^2}{6} \rho_{22}(s) \right] e^{-\frac{s}{M^2}}. \quad (8)
\]

Differentiating Eq.(8) with respect to \( \frac{1}{M^2} \), then eliminate the pole residue \( \lambda_X \), we can obtain the sum rule for the mass,

\[
M_X^2 = \frac{\int_{\Delta^2}^{s_0} ds \frac{d}{d(-1/M^2)} [3\rho_{11}(s) + (\bar{s}s)^2 \rho_{22}(s)] e^{-\frac{s}{M^2}}}{\int_{\Delta^2}^{s_0} ds [3\rho_{11}(s) + (\bar{s}s)^2 \rho_{22}(s)] e^{-\frac{s}{M^2}}}. \quad (9)
\]

3 Numerical results and discussions

The input parameters are taken to be the standard values \( \langle q\bar{q} \rangle = -(0.24 \pm 0.01 \text{ GeV})^3 \), \( \langle \bar{s}s \rangle = (0.8 \pm 0.2) \langle q\bar{q} \rangle \), \( \langle \bar{s}q_i \sigma Gs \rangle = m_\alpha^2 \langle \bar{s}s \rangle \), \( m_\alpha^2 = (0.8 \pm 0.2) \text{ GeV}^2 \), \( \langle \frac{a_{GG}}{\pi} \rangle = (0.33 \text{ GeV})^4 \), \( m_s = (0.14 \pm 0.01) \text{ GeV} \) and \( m_c = (1.35 \pm 0.10) \text{ GeV} \) at the energy scale \( \mu = 1 \text{ GeV} \). In Refs.[28, 29], there are two criteria (pole dominance and convergence of the operator product expansion) for choosing the Borel parameter \( M^2 \) and threshold parameter \( s_0 \). In Refs.[14, 15], we study the \( D^* \bar{D}^*, D_s^* \bar{D}_s^*, B^* \bar{B}^* \) and \( B_s^* \bar{B}_s^* \) molecular states in a systematic way, the threshold parameters are \( s_0 = (24 \pm 1) \text{ GeV}^2 \), \( (25 \pm 1) \text{ GeV}^2 \), \( (138 \pm 2) \text{ GeV}^2 \) and \( (140 \pm 2) \text{ GeV}^2 \) in the \( c\gamma_{\mu}ud\gamma_{\mu}c, c\gamma_{\mu}\bar{s}s\gamma_{\mu}c, b\gamma_{\mu}ud\gamma_{\mu}b \) and \( b\gamma_{\mu}\bar{s}s\gamma_{\mu}b \) channels respectively; the Borel parameters are \( M^2 = (2.6 - 3.0) \text{ GeV}^2 \) and \( (7.0 - 8.0) \text{ GeV}^2 \) in the hidden charm and hidden bottom channels respectively. In those regions, the two criteria of the QCD sum rules are satisfied. In this article, we choose the interpolating current \( J(x) \), which is a special superposition of the scalar currents \( \bar{c}(x)c(x) \) and \( \bar{c}(x)\gamma_{\mu}s(x)\bar{s}(x)\gamma_{\mu}c(x) \). So we can take the same threshold parameter and Borel parameter as in the channel \( \bar{c}\gamma_{\mu}\bar{s}s\gamma_{\mu}c \), i.e. \( s_0 = (25 \pm 1) \text{ GeV}^2 \) and \( M^2 = (2.6 - 3.0) \text{ GeV}^2 \).

The contributions from the different terms in the operator product expansion are shown in Fig.1, from the figure, we can see that the dominant contribution comes from the perturbative term and the operator product expansion is well convergent. In Fig.2, we show the contribution from the pole term with variation of the Borel parameter and the threshold parameter. The pole contribution is larger than 50%, the pole dominant condition is also satisfied.

Taking into account all uncertainties of the relevant parameters, finally we obtain the values of the mass and pole residue of the narrow structure \( X(4350) \), which are shown in Fig.3,

\[
M_X = (4.37 \pm 0.15) \text{ GeV}, \quad \lambda_X = (4.1 \pm 0.8) \times 10^{-2} \text{ GeV}^5. \quad (10)
\]
The prediction is in good agreement with the experimental data $M_X = (4350.6^{+4.6}_{-5.1} \pm 0.7)$ MeV [16], the $X(4350)$ may be a scalar $\bar{c}c - D_s^*\bar{D}_s^*$ mixing state. Other possibility, such as a scalar (tensor) $c\bar{c}s\bar{s}$ tetraquark state is not excluded.

The nominal thresholds of the $D_s - \bar{D}_s$ and $D_s^* - \bar{D}_s^*$ are $M_{D_s\bar{D}_s} = 3.937$ GeV and $M_{D_s^*\bar{D}_s^*} = 4.225$ GeV respectively [31], the strong decays $X(4350) \rightarrow D_s\bar{D}_s, D_s^*\bar{D}_s^*$ can take place, so we can search for the $X(4350)$ in the $D_s\bar{D}_s$ and $D_s^*\bar{D}_s^*$ invariant mass distributions. The decay channel $X(4350) \rightarrow D_s^*\bar{D}_s^*$ has much smaller phase space comparing with the decay channel $X(4350) \rightarrow D_s\bar{D}_s$, the strong decay $X(4350) \rightarrow D_s\bar{D}_s$ is of great importance. By measuring the relative angular distributions of the pseudoscalar mesons $D_s$ and $\bar{D}_s$, we can determine the spin of the $X(4350)$.

In Ref. [33], the light nonet scalar mesons are taken as tetraquark states, and the strong coupling constants among the light scalar mesons and pseudoscalar mesons are calculated with the QCD sum rules. The numerical results indicate that the values of the strong coupling constants for the tetraquark states are always smaller than the corresponding ones for the $q\bar{q}$ states [34, 35]. In Ref. [36], Maiani et al take the diquarks as the basic constituents, examine the rich spectrum of the diquark-antidiquark states with the constituent diquark masses and the spin-spin interactions, and try to accommodate some of the newly observed charmonium-like resonances not fitting a pure $c\bar{c}$ assignment. The predictions (also the Ref. [18]) depend on the assumption that the light scalar mesons $a_0(980)$ and $f_0(980)$ are tetraquark states, the basic parameters (constituent diquark masses) are estimated thereafter. If the scenarios of the light nonet scalar mesons as the tetraquark states are robust, the scalar (tensor) $c\bar{c}s\bar{s}$ tetraquark state will have smaller $D_s\bar{D}_s$ decay width than the corresponding ones of the $\bar{c}c - D_s^*\bar{D}_s^*$ mixing state.

4 Conclusion

In this article, we assume that the $X(4350)$ is a scalar $\bar{c}c - D_s^*\bar{D}_s^*$ mixing state, and study its mass using the QCD sum rules. Our prediction depends heavily on the two criteria (pole dominance and convergence of the operator product expansion) of the QCD sum rules. The numerical result is in good agreement with the experimental data, the $X(4350)$ may be a scalar $\bar{c}c - D_s^*\bar{D}_s^*$ mixing state. Other possibility, such as a scalar (tensor) $c\bar{c}s\bar{s}$ tetraquark state is not excluded. We can search for $X(4350)$ in the $D_s\bar{D}_s$ and $D_s^*\bar{D}_s^*$ invariant mass distributions, especially the $D_s\bar{D}_s$. By measuring the relative angular distributions of the pseudoscalar mesons $D_s$ and $\bar{D}_s$, we can determine the spin of the $X(4350)$.

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Figure 1: The contributions from the different terms with variation of the Borel parameter $M^2$ for $s_0 = 25\text{ GeV}^2$ in the operator product expansion. The $A$, $B$, $C$, $D$, $E$ and $F$ correspond to the contributions from the perturbative term, $\langle ss\rangle + \langle s g_s \sigma G s \rangle$ term, $\langle \alpha_s \frac{G G}{\pi} \rangle$ term, $\langle \frac{\alpha_s G G}{\pi} \rangle + \langle \frac{\alpha_s G G}{\pi} \rangle \left[ \langle ss \rangle + \langle sg_s \sigma G s \rangle + \langle ss \rangle^2 \right]$ term, $\langle ss \rangle^2 + \langle ss \rangle \langle sg_s \sigma G s \rangle$ term and $\langle sg_s \sigma G s \rangle^2$ term, respectively.

Figure 2: The contribution of the pole term with variation of the Borel parameter $M^2$. 

Figure 3: The mass and pole residue of the $X(4350)$ with variation of the Borel parameter $M^2$.

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