1 Introduction

Recently, the LHCb collaboration reported a narrow peak in the $D^- K^+$ invariant mass spectrum in the decays $B^\pm \rightarrow D^+ D^- K^+$ with the statistical significance much greater than 5σ [1]. The peak has been parameterized in terms of two Breit-Wigner resonances:

\begin{align}
X_0(2900) & : \quad J^P = 0^+, \quad M_0 = 2866 \pm 7 \text{ MeV}, \quad \Gamma_0 = 57 \pm 13 \text{ MeV} ; \\
X_1(2900) & : \quad J^P = 1^-, \quad M_1 = 2904 \pm 5 \text{ MeV}, \quad \Gamma_1 = 110 \pm 12 \text{ MeV} .
\end{align}

This is the first exotic hadron with fully open flavor, the valence quarks or the constituent quarks are $cs\bar{u}d$ [1]. In Ref.[2], Karliner and Rosner assign the narrow peak to be the scalar-diquark-scalar-antidiquark ($SS$) type fully open flavor $cs\bar{u}d$ tetraquark states with the spin-parity $J^P = 0^+$ via the QCD sum rules. The predicted masses $M_{AA} = 2.91 \pm 0.12 \text{GeV}$ and $M_{SS} = 3.05 \pm 0.10 \text{GeV}$ support assigning the $X_0(2900)$ to be the $AA$-type scalar $cs\bar{u}d$ tetraquark state.

In 2019, the BESIII collaboration explored the process $J/\psi \rightarrow \phi \eta' \eta'$ and observed a structure $X$ in the $\phi \eta'$ mass spectrum [8]. The fitted mass and width are $M_X = (2002.1 \pm 27.5 \pm 15.0) \text{MeV}$ and $\Gamma_X = (129 \pm 17 \pm 7) \text{MeV}$ respectively with the assignment $J^P = 1^-$, while the fitted mass and width are $M_X = (2062.8 \pm 13.1 \pm 4.2) \text{MeV}$ and $\Gamma_X = (177 \pm 36 \pm 20) \text{MeV}$ respectively with the assignment $J^P = 1^+$. In Ref.[3], we study the axialvector-diquark-axialvector-antidiquark type scalar, axialvector, tensor and vector $ss\bar{s}s/qq\bar{q}q$ tetraquark states with the QCD sum rules in a systematic way. The predicted mass $M_X = 2.08 \pm 0.12 \text{GeV}$ for the axialvector tetraquark state supports assigning the new structure $X(2060)$ from the BESIII collaboration to be a $ss\bar{s}s$ tetraquark state with the spin-parity-charge-conjugation $J^{PC} = 1^{++}$. In Ref.[10], we construct various scalar, axialvector and tensor tetraquark currents to study the mass spectrum of the ground state hidden-charm tetraquark states with the QCD sum rules in a comprehensive way, and revisit the assignments of the $X$, $Y$, $Z$ states, such as $X(3860)$, $X(3872)$, $X(3915)$, $X(3940)$, $X(4160)$, $Z_c(3900)$, $Z_c(4020)$, $Z_c(4055)$, $Z_c(4100)$, $Z_c(4200)$, $Z_c(4250)$, $Z_c(4430)$, $Z_c(4600)$, etc in a consistent way. For the axialvector-diquark-axialvector-antidiquark type ($AA$-type) scalar tetraquark states, we obtain the masses [9][10],

\begin{align}
M_{qq\bar{q}\bar{q}} &= 1.86 \pm 0.11 \text{GeV} , \\
M_{ss\bar{s}s} &= 2.08 \pm 0.13 \text{GeV} , \\
M_{c\bar{q}c\bar{q}} &= 3.95 \pm 0.09 \text{GeV} .
\end{align}
Now we can estimate the mass of the $AA$-type $cs{\bar{u}}d$ tetraquark state crudely,

$$M_{cs{\bar{u}}d} = \frac{M_{qq{\bar{q}}} + M_{ss{\bar{s}}} + 2M_{cq{\bar{q}}}}{4} = 2.96 \pm 0.11 \text{ GeV},$$

which is consistent with the mass of the $X_0(2900)$ within uncertainties.

In this article, we construct the scalar-diquark-scalar-antidiquark type ($SS$-type) and axialvector-diquark-antidiquark type ($AA$-type) scalar currents to study the masses of the $cs{\bar{u}}d$ tetraquark states with the QCD sum rules in details and explore the possible assignment of the $X_0(2900)$ as the scalar tetraquark state.

The article is arranged as follows: we obtain the QCD sum rules for the masses and pole residues of the scalar tetraquark states in Sect.2; in Sect.3, we present the numerical results and discussions; and Sect.4 is reserved for our conclusion.

## 2 QCD sum rules for the scalar tetraquark states

Firstly, we write down the two-point correlation functions $\Pi(p^2)$ in the QCD sum rules,

$$\Pi(p^2) = i \int d^4x e^{ipx} \langle 0| T \{ J(x)\bar{J}(0) \} | 0 \rangle,$$

where $J(x) = J_{AA}(x), J_{SS}(x),$$

$$J_{AA}(x) = \epsilon^{ijk}\epsilon^{imn} s_j^T(x)C\gamma_\alpha c_k(x)\bar{u}_m(x)\gamma^\alpha \bar{C}d_n^T(x),$$

$$J_{SS}(x) = \epsilon^{ijk}\epsilon^{imn} s_j^T(x)C\gamma_\alpha c_k(x)\bar{u}_m(x)\gamma_5 C\bar{d}_n^T(x),$$

the $i, j, k, m$ and $n$ are color indexes, the $C$ is the charge conjugation matrix. The attractive interactions of one-gluon exchange favor formation of the diquarks in color antitriplet [11]. The QCD sum rules calculations indicate that the favored quark-quark configurations are the scalar and axialvector diquark states [12] [13] [14].

At the hadron side, we insert a complete set of scalar tetraquark states with the same quantum numbers as the current operators $J(x)$ into the correlation functions $\Pi(p^2)$ to obtain the hadronic representation [13] [16]. After isolating the pole terms of the lowest $cs{\bar{u}}d$ tetraquark states $X_0$, we obtain the result:

$$\Pi(p^2) = \frac{\lambda_X^3}{M_X^2 - p^2} + \cdots,$$

where the pole residues $\lambda_X$ are defined by $\langle 0| J(0)| X(p) \rangle = \lambda_X$.

Now, we briefly outline the operator product expansion for the correlation functions $\Pi(p^2)$ in perturbative QCD. Firstly, we contract the $u, d, s$ and $c$ quark fields in the correlation functions $\Pi(p^2)$ with Wick theorem, and obtain the result:

$$\Pi_{AA}(p^2) = i \epsilon^{ijk}\epsilon^{imn} \epsilon^{i'j'k'} \epsilon^{i'm'n'} \int d^4x e^{ipx}$$

$$\text{Tr} \left[ \gamma_\mu C_{k'k}(x)\gamma_\nu C S_{j'j}(x)C \right] \text{Tr} \left[ \gamma^\nu U_{m'm}(x)\gamma^\mu C D_{n'n}(x) \right],$$

$$\Pi_{SS}(p^2) = i \epsilon^{ijk}\epsilon^{imn} \epsilon^{i'j'k'} \epsilon^{i'm'n'} \int d^4x e^{ipx}$$

$$\text{Tr} \left[ \gamma_5 C_{k'k}(x)\gamma_5 C S_{j'j}(x)C \right] \text{Tr} \left[ \gamma_5 U_{m'm}(x)\gamma_5 C D_{n'n}(x) \right],$$

where the $U_{ij}(x), D_{ij}(x), S_{ij}$ and $C_{ij}(x)$ are the full $u, d, s$ and $c$ quark propagators, respectively,

$$\frac{U/D_{ij}(x)}{2\pi^2 x^4} = \frac{\delta_{ij} \delta(q\pi)}{12} - \frac{\delta_{ij} x^2 (q\pi, \sigma Gq)}{192} - \frac{ig_\sigma G_{\alpha\beta} (x) \alpha^\beta (x) x^2}{32\pi^2 x^2}$$

$$- \frac{1}{8} (q\pi, \sigma^{\mu\nu} q_\nu) \sigma_{\mu\nu} + \cdots.$$
\[
S_{ij}(x) = \frac{i\delta_{ij} x}{2\pi^2 x^4} - \frac{\delta_{ij} m_s}{4\pi^2 x^2} - \frac{\delta_{ij} \langle \bar{s}s \rangle}{12} + \frac{i\delta_{ij} \not{\!m_s} \langle \bar{s}s \rangle}{48} - \frac{\delta_{ij} x^2 \langle \bar{s}g_s \sigma Gs \rangle}{192} + \frac{i\delta_{ij} x^2 \not{\!m_s} \langle \bar{s}g_s \sigma Gs \rangle}{1152} - \frac{ig_s G_{\alpha \beta} G_{\alpha \beta}^i \hat{t}^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} s^n (\tilde{f} x) + \cdots ,
\]

\[
C_{ij}(x) = i \frac{(2\pi)^4}{(2\pi)^4} \int d^4k e^{-ik \cdot x} \left\{ \frac{\delta_{ij} m_s}{k - m_c} - \frac{g_s G_{\alpha \beta} G_{\mu \nu}^i \sigma^{\alpha \beta} (k + m_c)}{(k^2 - m_c^2)^2} \left( \frac{g_s G_{\alpha \beta} G_{\mu \nu}^i \sigma^{\alpha \beta} (k + m_c)}{(k^2 - m_c^2)^2} \right) + \cdots \right\} ;
\]

\[
f^{\alpha \beta \mu \nu} = (k + m_c) \gamma^\alpha (k + m_c) \gamma^\beta (k + m_c) \gamma^\mu (k + m_c) \gamma^\nu (k + m_c) .
\]

and \( t^n = \lambda^n \), the \( \lambda^n \) is the Gell-Mann matrix [16] [17] [18]. We retain the terms \( \langle \bar{q}_j \sigma_{\mu \nu} q_i \rangle \) and \( \langle \bar{s}_j \sigma_{\mu \nu} s_i \rangle \) come from Fierz re-ordering of the \( \langle \bar{q}_j \hat{q}_j \rangle \) and \( \langle \bar{s}_j \hat{s}_j \rangle \) to absorb the gluons emitted from other quark lines to extract the mixed condensate \( \langle \bar{q} g_s \sigma Gq \rangle \) and \( \langle \bar{s} g_s \sigma Gs \rangle \), respectively [18]. Then we compute the integrals both in the coordinate space and momentum space to obtain the correlation functions \( \Pi(p^2) \). Finally, we obtain the QCD spectral densities \( \rho(s) \) at the quark level through dispersion relation,

\[
\rho(s) = \lim_{\epsilon \to 0} \frac{\text{Im} \Pi(s + i\epsilon)}{\pi} .
\]

In this article, we carry out the operator product expansion up to the vacuum condensates of dimension-11, and assume vacuum saturation for the higher dimensional vacuum condensates. There are three light quark propagators and one heavy quark propagator in the correlation functions \( \Pi(p^2) \), if the heavy quark line emits a gluon and each light quark line contributes a quark-antiquark pair, we obtain a quark-gluon operator \( g_s G_{\mu \nu} \bar{q} q \bar{q} q \bar{q} s s \), which is of dimension 11, and can lead to the vacuum condensates \( \langle \bar{q} q \rangle^2 \langle \bar{s} g_s \sigma Gs \rangle \) and \( \langle \bar{q} q \rangle \langle \bar{s} s \rangle \langle \bar{q} g_s \sigma Gq \rangle \), which we should take into account the vacuum condensate up to dimension 11 in a consistent way. As the vacuum condensates are the vacuum expectations of the quark-gluon operators, we take into account the quark-gluon operators of the orders \( O(a_s^k) \) with \( k \leq 1 \) [18] [19].

Now we can take the quark-hadron duality below the continuum thresholds \( s_0 \) and perform the Borel transform to obtain the QCD sum rules:

\[
\lambda_X \exp \left( \frac{-M_X^2}{T^2} \right) = \int_{m_s^2}^{s_0} ds \rho(s) \exp \left( \frac{-s}{T^2} \right) ,
\]

where the \( T^2 \) is the Borel parameter, \( \rho(s) = \rho_{AA}(s), \rho_{SS}(s) \), we neglect the explicit expressions for simplicity.

We differentiate Eq. (14) with respect to \( \tau = \frac{1}{T^2} \), then eliminate the pole residues \( \lambda_X \) and obtain the QCD sum rules for the masses of the scalar \( c \bar{s} u d \) tetraquark states,

\[
M_X^2 = \frac{\frac{d}{d\tau} \int_{m_s^2}^{s_0} ds \rho(s) \exp \left( -s\tau \right)}{\int_{m_s^2}^{s_0} ds \rho(s) \exp \left( -s\tau \right) } .
\]

3 numerical results and discussions

We take the standard values of the vacuum condensates \( \langle \bar{q} q \rangle = -(0.24 \pm 0.01 \text{ GeV})^3 \), \( \langle \bar{s} s \rangle = (0.8 \pm 0.1) \langle \bar{q} q \rangle \), \( \langle \bar{q} g_s \sigma Gq \rangle = m_0^2 \langle \bar{q} q \rangle \), \( \langle \bar{s} g_s \sigma Gs \rangle = m_0^2 \langle \bar{s} s \rangle \), \( m_0^2 = (0.8 \pm 0.1) \text{ GeV}^2 \), \( g_s G_{\mu \nu} \langle \bar{q} q \rangle = (0.33 \pm 0.05) \text{ GeV}^4 \) at the energy scale \( \mu = 1 \text{ GeV} \) [15] [17] [20], and take the \( \overline{\text{MS}} \) masses \( m_{c_i} = (1.275 \pm 0.025) \text{ GeV} \) and \( m_s(\mu = 2 \text{ GeV}) = (0.95 \pm 0.005) \text{ GeV} \) from the Particle Data Group [21]. Furthermore, we
take into account the energy-scale dependence of the quark condensates, mixed quark condensates and $\overline{M}\overline{S}$ masses according to the renormalization group equation \cite{22},

$$
\langle \bar{q}q \rangle (\mu) = \langle \bar{q}q \rangle (1\text{GeV}) \left[ \frac{\alpha_s(1\text{GeV})}{\alpha_s(\mu)} \right]^{\frac{12}{\pi^2}},
$$

$$
\langle \bar{s}s \rangle (\mu) = \langle \bar{s}s \rangle (1\text{GeV}) \left[ \frac{\alpha_s(1\text{GeV})}{\alpha_s(\mu)} \right]^{\frac{12}{\pi^2}},
$$

$$
\langle \bar{q}g_\alpha \sigma \bar{g}q_\beta \rangle (\mu) = \langle \bar{q}g_\alpha \sigma \bar{g}q_\beta \rangle (1\text{GeV}) \left[ \frac{\alpha_s(1\text{GeV})}{\alpha_s(\mu)} \right]^{\frac{2}{\pi^2}},
$$

$$
\langle \bar{s}g_\alpha \sigma Gs \rangle (\mu) = \langle \bar{s}g_\alpha \sigma Gs \rangle (1\text{GeV}) \left[ \frac{\alpha_s(1\text{GeV})}{\alpha_s(\mu)} \right]^{\frac{4}{\pi^2}},
$$

$$
m_c(\mu) = m_c(m_c) \left[ \frac{\alpha_s(\mu)}{\alpha_s(m_c)} \right]^{\frac{12}{\pi^2}},
$$

$$
m_s(\mu) = m_s(2\text{GeV}) \left[ \frac{\alpha_s(\mu)}{\alpha_s(2\text{GeV})} \right]^{\frac{12}{\pi^2}},
$$

$$
\alpha_s(\mu) = \frac{1}{b_0 t} \left[ 1 - \frac{b_1 \log t}{b_0^2} + \frac{b_2 (\log^2 t - \log t - 1)}{b_0^2 t^2} \right], \quad (16)
$$

where $t = \log \frac{\mu^2}{\Lambda^2}$, $b_0 = \frac{33 - 2n_f}{12\pi}$, $b_1 = \frac{153 - 19n_f}{24\pi}$, $b_2 = \frac{2387 - 2038n_f + 225n_f^2}{128\pi^2}$, $\Lambda = 213 \text{ MeV}$, $296 \text{ MeV}$ and 339 MeV for the flavors $n_f = 5$, 4 and 3, respectively. For the fully open flavor $cs\bar{u}\bar{d}$ tetraquark states, we choose the flavor numbers $n_f = 4$, and the typical energy scale $\mu = 1 \text{ GeV}$.

Let us choose the continuum threshold parameters as $\sqrt{s_0} = M_X + (0.5 \sim 0.7) \text{ GeV}$ tentatively according to the mass gap $m_{\psi'} - m_{J/\psi} = 0.59 \text{ GeV}$ \cite{21}, and vary the parameters $\sqrt{s_0}$ to obtain the best Borel parameters $T^2$ to satisfy pole dominance at the hadron side and convergence of the operator product expansion at the QCD side.

After trial and error, we obtain the ideal Borel parameters or Borel windows $T^2$ and continuum threshold parameters $s_0$, therefore the pole contributions of the ground state scalar $cs\bar{u}\bar{d}$ tetraquark states and the convergent behaviors of the operator product expansion, see Table I. In the Borel windows, the pole contributions are about $(38 - 67)\%$, while the central values exceed 52%, the pole dominance is well satisfied. The absolute values of the contributions of the highest dimensional vacuum condensates $|D(11)|$ are about $(2 - 4)\%$ and $(0 - 1)\%$ for the $AA$-type and $SS$-type tetraquark states, respectively. The operator product expansion is well convergent.

At the beginning, we assume the ground states of the scalar tetraquark states $cs\bar{u}\bar{d}$ have the masses about 2.9 GeV, just like the $X_{b}(2900)$, and choose the continuum threshold parameters $\sqrt{s_0} = 2.9 + (0.5 \sim 0.7) \text{ GeV}$ tentatively to search for the optimal values via trial and error to satisfy the constraint $\sqrt{s_0} = M_X + (0.5 \sim 0.7) \text{ GeV}$ besides the two basic criteria of the QCD sum rules. From Table I, we can see that for the $AA$-type scalar tetraquark state, the continuum threshold parameter $\sqrt{s_0} = 2.9 + (0.5 \sim 0.7) \text{ GeV}$ happens to coincide with the optimal value $3.5 \pm 0.1 \text{ GeV}$, while for the $SS$-type scalar tetraquark state, the continuum threshold parameter $\sqrt{s_0} = 2.9 + (0.5 \sim 0.7) \text{ GeV}$ is slightly smaller than the optimal value $3.6 \pm 0.1 \text{ GeV}$. In fact, we can choose other values of the continuum threshold parameters $\sqrt{s_0}$, for example, $\sqrt{s_0} = 2.5 + (0.5 \sim 0.7) \text{ GeV}$ as the initial point, and obtain the optimal values shown Table I.

Now we take into account all uncertainties of the input parameters, and obtain the values of the masses and pole residues of the fully open flavor $cs\bar{u}\bar{d}$ tetraquark states, which are shown explicitly in Table II and Fig II. In Fig II, we plot the masses of the $AA$-type and $SS$-type scalar $cs\bar{u}\bar{d}$ tetraquark states with variations of the Borel parameters $T^2$ in much larger ranges than the Borel windows. From the figure, we can see that there appear platforms in the Borel windows, it is reliable to extract the tetraquark masses.
The predicted mass $M_{AA} = 2.91 \pm 0.12$ GeV is consistent with the experimental value 2866 ± 7 MeV from the LHCb collaboration [1], and supports assigning the $X_0(2900)$ to be the $AA$-type $cs\bar{u}d$ tetraquark state with the spin-parity $J^P = 0^+$. While the predicted mass $M_{SS} = 3.05 \pm 0.10$ GeV lies above the experimental value 2866 ± 7 MeV from the LHCb collaboration [1].

The two-body strong decays $X_0(2900) \to D\bar{K}$ can take place with the fall-apart mechanism and are kinematically allowed, therefore it is Okubo-Zweig-Iizuka super-allowed. The current $J_{AA}(x)$ also couples potentially to the two-meson scattering states $D\bar{X}_{cs}$, where energies $\Sigma_D$.

Then the hadron sides of the QCD sum rules in Eqs.(14)-(15) undergo the replacements, of the $X$ part to modify the dispersion relation, with the central value of the continuum threshold parameter $\lambda_X(3900)$ plays an un-substitutable role, we can saturate the QCD sum rules with or without the two-particle scattering state contributions. The conclusion is applicable in the present case.

The contributions of the intermediate two-meson scattering states $D\bar{K}$, $D^*\bar{K}^*$, etc besides the scalar tetraquark candidate $X_0(2900)$ can be written as,

$$\Pi_{AA}(p^2) = -\frac{\lambda_X^2}{p^2 - M_X^2 + \Sigma_{D\bar{K}}(p^2) + \Sigma_{D^*\bar{K}^*}(p^2) + \cdots}$$ (17)

We choose the bare mass and pole residue $\tilde{M}_X$ and $\tilde{\lambda}_X$ to absorb the divergences in the self-energies $\Sigma_{D\bar{K}}(p^2)$, $\Sigma_{D^*\bar{K}^*}(p^2)$, etc. The renormalized self-energies contribute a finite imaginary part to modify the dispersion relation,

$$\Pi_{AA}(p^2) = -\frac{\lambda_X^2}{p^2 - M_X^2 + i\sqrt{p^2}\Gamma_X(p^2) + \cdots},$$ (18)

with the (central value of) physical width $\Gamma_X(M_X^2) = 57$ MeV from the LHCb collaboration [1].

We can take into account the finite width with the simple replacement of the hadronic spectral density,

$$\lambda_X^2 \delta(s - M_X^2) \rightarrow \lambda_X^2 \int^{\infty} ds \frac{M_X\Gamma_X(s)}{\pi(s - M_X^2)^2 + M_X^2\Gamma_X^2(s)} = \lambda_X^2 \int^{\infty} ds \frac{M_X\Gamma_X(s)}{\pi(s - M_X^2)^2 + M_X^2\Gamma_X^2(s)},$$ (19)

where

$$\Gamma_X(s) = \Gamma_X \frac{M_X^2}{s} \sqrt{\frac{s - (M_D + M_K)^2}{M_X^2 - (M_D + M_K)^2}}.$$ (20)

Then the hadron sides of the QCD sum rules in Eqs.(14)-(15) undergo the replacements,

$$\lambda_X^2 \exp\left(-\frac{M_X^2}{T^2}\right) \rightarrow \lambda_X^2 \int^{\infty}_{(M_D + M_K)^2} ds \frac{1}{\pi(s - M_X^2)^2 + M_X^2\Gamma_X^2(s)} \exp\left(-\frac{s}{T^2}\right),$$

$$= (0.97 \sim 0.98) \lambda_X^2 \exp\left(-\frac{M_X^2}{T^2}\right),$$ (21)

$$\lambda_X^2 M_X^2 \exp\left(-\frac{M_X^2}{T^2}\right) \rightarrow \lambda_X^2 \int^{\infty}_{(M_D + M_K)^2} ds \frac{1}{\pi(s - M_X^2)^2 + M_X^2\Gamma_X^2(s)} \exp\left(-\frac{s}{T^2}\right),$$

$$= (0.96 \sim 0.96) \lambda_X^2 M_X^2 \exp\left(-\frac{M_X^2}{T^2}\right),$$ (22)

with the central value of the continuum threshold parameter $\sqrt{s_0} = 3.50$ GeV. We can absorb the numerical factors $0.97 \sim 0.98$ and $0.96 \sim 0.96$ into the pole residue with the simple replacement $\lambda_X \rightarrow (0.98 \sim 0.99)\lambda_X$ safely. It is indeed that the intermediate meson-loops cannot affect the mass $M_X$ and pole residue $\lambda_X$ remarkably.
\( T^2(\text{GeV}^2) \quad \sqrt{s_0}(\text{GeV}) \quad \text{pole} \quad |D(11)| \quad M(\text{GeV}) \quad \lambda(\text{GeV}^2) \)

\begin{tabular}{|c|c|c|c|c|c|}
\hline
\( |c|s_A|\bar{u}d|_A \) & 1.9 - 2.3 & 3.5 \pm 0.1 & (38 - 67)\% & (2 - 4)\% & 2.91 \pm 0.12 & (1.60 \pm 0.33) \times 10^{-2} \\
\( |c|s|\bar{u}d|_S \) & 2.1 - 2.5 & 3.6 \pm 0.1 & (39 - 66)\% & (0 - 1)\% & 3.05 \pm 0.10 & (1.20 \pm 0.21) \times 10^{-2} \\
\hline
\end{tabular}

Table 1: The Borel windows, continuum threshold parameters, pole contributions, contributions of the vacuum condensates of dimension 11, masses and pole residues for the scalar tetraquark states.

Figure 1: The masses of the AA-type and SS-type tetraquark states with variations of the Borel parameters \( T^2 \).

4 Conclusion

In this article, we construct the axialvector-diquark-axialvector-antidiquark type and scalar-diquark-scalar-antidiquark type currents to study the fully open flavor \( cs\bar{u}\bar{d} \) tetraquark states with the spin-parity \( J^P = 0^+ \) via the QCD sum rules by carrying out the operator product expansion up to the vacuum condensates of dimension 11 in a consistent way. We obtain the predictions \( M_{AA} = 2.91 \pm 0.12 \text{ GeV} \) and \( M_{SS} = 3.05 \pm 0.10 \text{ GeV} \), the predicted mass for the axialvector-diquark-axialvector-antidiquark type scalar tetraquark state is consistent with the experimental value 2866 \pm 7 \text{ MeV} \) from the LHCb collaboration, and supports assigning the \( X_{0}(2900) \) to be the axialvector-diquark-axialvector-antidiquark type \( cs\bar{u}\bar{d} \) tetraquark state with the spin-parity \( J^P = 0^+ \).

Acknowledgements

This work is supported by National Natural Science Foundation, Grant Number 11775079.

References

[1] R. Aaij et al, arXiv:2009.00025; R. Aaij et al, arXiv:2009.00026.
[2] M. Karliner and J. L. Rosner, arXiv:2008.05993.
[3] M. W. Hu, X. Y. Lao, P. Ling and Q. Wang, arXiv:2008.06894; M. Z. Liu, J. J. Xie and L. S. Geng, arXiv:2008.07389; H. X. Chen, W. Chen, R. R. Dong and N. Su, arXiv:2008.07516.
[4] X. G. He, W. Wang and R. L. Zhu, arXiv:2008.07145.
[5] X. H. Liu, M. J. Yan, H. W. Ke, G. Li and J. J. Xie, arXiv:2008.07190.

[6] J. R. Zhang, arXiv:2008.07295.

[7] Q. F. Lu, D. Y. Chen and Y. B. Dong, arXiv:2008.07340.

[8] M. Ablikim et al, Phys. Rev. D99 (2019) 112008.

[9] Z. G. Wang, Adv. High Energy Phys. 2020 (2020) 6438730.

[10] Z. G. Wang, Phys. Rev. D102 (2020) 014018.

[11] 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.

[12] H. G. Dosch, M. Jamin and B. Stech, Z. Phys. C42 (1989) 167; M. Jamin and M. Neubert, Phys. Lett. B238 (1990) 387.

[13] Z. G. Wang, Eur. Phys. J. C71 (2011) 1524; R. T. Kleiv, T. G. Steele and A. Zhang, Phys. Rev. D87 (2013) 125018.

[14] Z. G. Wang, Commun. Theor. Phys. 59 (2013) 451.

[15] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. B147 (1979) 385; Nucl. Phys. B147 (1979) 448.

[16] L. J. Reinders, H. Rubinstein and S. Yazaki, Phys. Rept. 127 (1985) 1.

[17] P. Pascual and R. Tarrach, “QCD: Renormalization for the practitioner”, Springer Berlin Heidelberg (1984).

[18] Z. G. Wang and T. Huang, Phys. Rev. D89 (2014) 054019.

[19] Z. G. Wang and T. Huang, Eur. Phys. J. C74 (2014) 2891; Z. G. Wang, Eur. Phys. J. C74 (2014) 2874.

[20] P. Colangelo and A. Khodjamirian, hep-ph/0010175.

[21] C. Patrignani et al, Chin. Phys. C40 (2016) 100001.

[22] S. Narison and R. Tarrach, Phys. Lett. 125 B (1983) 217.

[23] S. S. Agaev, K. Azizi and H. Sundu, Phys. Rev. D93 (2016) 094006.

[24] Z. G. Wang, Int. J. Mod. Phys. A35 (2020) 2050138.