The $\Omega_{cc}$ resonances with negative parity in the chiral constituent quark model

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Spectrum of the low-lying $\Omega_{cc}$ resonances with negative parity, which are assumed to be dominated by $scc\bar{q}$ pentaquark components, is investigated using the chiral constituent quark model. Energies of the $\Omega_{cc}$ resonances are obtained by considering the hyperfine interaction between quarks by exchanging Goldstone boson. Possible $scc\bar{q}$ configurations with spin-parity $1/2^-$, $3/2^-$ and $5/2^-$ are taken into account. Numerical results show that the lowest $\Omega_{cc}$ resonances with negative parity may lie at $4050\pm100$ MeV. In addition, the transitions of the $\Omega_{cc}$ resonance to a pseudoscalar meson and a ground baryon state are also investigated within the chiral Lagrangian approach. We expect that these $\Omega_{cc}$ resonances could be observed in the $D\Xi_c$ channel by future experiments.

I. INTRODUCTION

Recently, five narrow $\Omega^{0}_{cc}$ resonances and a doubly charmed baryon $\Xi^{+}_{cc}$ were observed by LHCb collaboration Ref. [49]. These observations are of significant importance in hadron physics since the experimental data forbaryon resonances with one or more charm quarks are very poor before 2017 Ref. [4]. Accordingly, the theorists also made great efforts to describe the spectrum and decay behaviours of the observed heavy baryon resonances with charm quarks using various of approaches, such as the quark model within the three- or five-quark picture Ref. [5–13], QCD sum rules Ref. [14–18], the chiral perturbation approach Ref. [19–22], the lattice QCD Ref. [23–25], etc. Furthermore, the molecular nature of these heavy baryon resonances were also studied in Refs. [26–29], where they were dynamically generated from the meson–baryon interactions in the coupled channels. Besides, it is also very interesting that the observation of doubly heavy baryon is claimed to imply existence of heavy tetraquark mesons Ref. [30–31]. Thus, it is appropriate to perform further studies about these heavy baryons.

On the other hand, three hidden-charm $N_{cc}$ pentaquark states were observed by LHCb collaboration in recent years Ref. [32–34], where the experimental data are in agreement with the predictions made in Refs. Ref. [31–33]. Consequently, rather than the three-quark picture, it may be more interesting to study the properties of baryon resonances near or above 4 GeV within a five-quark picture, since the energy for pulling a light quark-antiquark pair to form a pentaquark configuration as the baryon excitation may be lower than that for the traditional orbital and radial excitations of a three-quark configuration Ref. [36]. Taking the five-quark picture, the $\Omega$ excited states with negative parity Ref. [37–44], nucleon excited states Ref. [45–51], and the newly observed $\Omega^+_1$ resonances Ref. [11–13] are investigated explicitly. Suggestions on how to observe the $\Omega(2012)$ state by looking at $\Lambda_c$ weak decay process have been made in Ref. [52]. It was found that the observed small energy splitting of the $\Omega^{++}$ resonances Ref. [13], the masses and decay behaviours of the observed $\Omega(2012)$ Ref. [41–44] can be well described by taking either the hadronic molecule picture or the compact pentaquark configuration, while it’s of course that one cannot rule out the three-quark components in the baryon resonances.

The observation of the $\Xi_{cc}$ states have brought new opportunities for us to study the doubly charmed baryons, since this finding suggests the potential of discovering more low-lying doubly charmed baryons in the near future, and thus one needs to have a solid theoretical calculations for the corresponding spectrum. In the present work, based on the chiral constituent quark model, we study the spectrum of the low-lying $\Omega_{cc}$ resonances with negative parity. And the transitions of these $\Omega_{cc}$ states to a pseudoscalar meson and a ground baryon state ($MB$) are also studied, employing the chiral Lagrangian approach which has been explicitly developed to study the strong decays of the $N_{s\bar{s}}$ nucleon resonances as in Ref. [49].

The present manuscript is organized as following: in section II we briefly present our theoretical formalism which includes the Hamiltonian and wave functions for the $\Omega_{cc}$ pentaquark system, and the chiral Lagrangian approach for strong decays of a five-quark system, we show our explicit numerical results in section III and section IV contains summary and conclusions.
II. THEORETICAL FRAME

We will briefly introduce the Hamiltonian and wave functions for the \( \Omega_{cc} \) resonances with negative parity as pentaquark states in Sec. II A and the chiral Lagrangian approach for strong decays of the \( \Omega_{cc} \) states in Sec. II B.

A. Hamiltonian and wave functions

In present work, the constituent quark model is employed to study the spectrum of \( \Omega_{cc} \) resonances, within which the Hamiltonian for a five-quark system can be written as

\[
H = \sum_{i<j} H_{hyp}^{ij} + \sum_{i=1,5} m_i + H_o ,
\]

where \( H_{hyp}^{ij} \) represents the hyperfine interaction between the \( i \)-th and \( j \)-th quarks in the five-quark system, \( m_i \) is the constituent mass of the \( i \)-th quark, and \( H_o \) is the Hamiltonian concerning orbital motions of the quarks, which should contain the kinetic term, the confinement potential of the quarks, and the flavor symmetry breaking term.

In general, the corresponding eigenvalue \( E_0 \) of \( H_o + \sum m_i \) in Eq. (1) should depend on the constituent masses of quarks and the model parameters in the quark confinement model, for instance, the confinement strength \( C \) and constant \( V_0 \) in the harmonic oscillator potential model \( \xi_m^2 \). In this work, we study the low-lying \( \Omega_{cc} \) resonances with negative parity as pentaquark states, which require all the quarks and antiquark to be in their ground states, accordingly, the eigenvalue \( E_o \) should be the same one for different pentaquark configurations.

The parameter \( E_0 \) has been taken to be 2127 MeV for investigations on the intrinsic sea flavor content of nucleon in Ref. 55, with which value the data for light sea quark asymmetry \( \bar{d} - \bar{u} \) in the proton can be well reproduced, while to fit the experimental data about \( \Omega^0_c \) resonances, we took \( E_0 = 3132 \) MeV in Ref. 13. As discussed in details in Ref. 13, the resulted different values of \( E_0 \) by fitting the experimental data should be consistent by the chiral constituent quark model if all the model parameters are taken to be the empirical values.

In this work, the value of \( E_0 \) should be \( \sim 1140 \) MeV higher than the one we took in Ref. 13, because of the different quark content in \( \Omega_{cc} \) and \( \Omega^0_c \) states, while the \( SU(4) \) flavor symmetry breaking effects caused by two charm quarks in present case will lower \( E_0 \) by \( \sim 170 \) MeV than those caused by one charm quark as in Ref. 13, if the Hamiltonian for symmetry breaking correction is taken to be the form similar as in Ref. 38. Consequently, hereafter we will take \( E_0 = 4102 \) MeV, based on the investigations on the intrinsic sea content of nucleon and spectrum of low-lying \( \Omega^0_c \) resonances, and the requirements of the chiral constituent quark model.

less, we will investigate the dependency of the results on \( E_0 \).

The hyperfine interaction between quarks is taken to be mediated by goldstone boson exchange and the corresponding \( H_{hyp}^{ij} \) is taken as following

\[
H_{hyp}^{ij} = -\vec{\sigma}_i \cdot \vec{\sigma}_j \left[ \sum_{a=1}^{3} V_\pi (r_{ij}) \lambda^a_i \lambda^a_j + \sum_{a=4}^{7} V_K (r_{ij}) \lambda^a_i \lambda^a_j + \sum_{a=9}^{12} V_D (r_{ij}) \lambda^a_i \lambda^a_j + \sum_{a=13}^{14} V_D (r_{ij}) \lambda^a_i \lambda^a_j + V_{\eta} (r_{ij}) \lambda^5_i \lambda^5_j \right] ,
\]

where \( V_M (r_{ij}) \) denotes the coupling strength for a meson \( M \) exchanged between the \( i \)-th and \( j \)-th quarks. In this work, the \( \pi, \eta, K, \eta_d \) and \( \eta \) mesons are taken into account.

For a five-quark system with the quark flavor as \( \Omega_{cc} \) resonances, namely, the \( sccq\bar{q} \) system, a general wave function can be written as

\[
\psi_{\bar{J}+J} = \sum_{a,b,c \{Y,T\},t_z,s_z} \sum_{s_\bar{c},s_c} C_{[4]}^{[1]}_{ia} C_{[31]}^{[3]}_{ia} [F^a]_{b,Y,T_z} [S^0]_{c,S_z} [211;C]_{\{Y,T,T_z,y,t_z\}0,0,0} \langle S,S_z,1/2,s_z,J,J_z|\bar{\xi}_{y,t_z}^c \varphi (\{\xi_j\}) \rangle .
\]

where \([F^a]_{b,Y,T_z} [S^0]_{c,S_z} \) and \([211;C]_{\{Y,T,T_z,y,t_z\}0,0,0}\) are the flavor, spin and color wave functions of the four-quark subsystem denoted by Young tableaux, the label \( i \) enumerates different pentaquark configurations. The \( \bar{\xi}_j \) is the Jacobi coordinates for a five quark system, which is defined as

\[
\bar{\xi}_j = \frac{1}{\sqrt{j+1}} \left( \sum_{i=1}^{j} r_i - j r_{j+1} \right) , j = 1, \ldots, 4 .
\]

According to the \( SU(2) \) symmetry, the spin wave function of a four-quark system may be \([4]_S, [31]_S \) or \([22]_S \), the corresponding spin quantum numbers are 2, 1 and 0, respectively. While the coupling between spin of the four-quark subsystem and the antiquark leads to \( J = 1/2, 3/2 \) or 5/2. Given that all the quarks and antiquark are in their ground states, namely, the orbital wave function of the four-quark system is \([4]_Y \), then the flavor wave function of the \( sccq\bar{q} \) subsystem can be \([4]_{[2]} , [31]_{[2]} \), \([31]_{[2]} \), \([22]_{[2]} \) and \([211]_{[2]} \). Finally, the possible pentaquark configurations denoted by \([i]\) for spin-parity quantum number \( J^P = 1/2^- \) are

\[
\begin{align*}
|1\rangle & : sccq[4]_{[2]} [211]_{[2]} [22]_{[2]} [211]_{c} \bar{q} ,
|2\rangle & : sccq[4]_{[2]} [31]_{[2]} [22]_{[2]} [211]_{c} \bar{q} ,
|3\rangle & : sccq[4]_{[2]} [31]_{[2]} [22]_{[2]} [211]_{c} \bar{q} ,
\end{align*}
\]
for \( J^P = 1/2^- \) or \( 3/2^- \) are

\[
\begin{align*}
|4\rangle & : \text{sc}c\text{c}q[4]^A_4[x_{[211]},r_{[31]}]s_{[211]}c^q, \\
|5\rangle & : \text{sc}c\text{c}q[4]^A_4[x_{[22]},r_{[31]}]s_{[211]}c^q, \\
|6\rangle & : \text{sc}c\text{c}q[4]^A_4[x_{[31]},r_{[31]}]s_{[211]}c^q, \\
|7\rangle & : \text{sc}q[4]^A_4[x_{[31]}]_3^p_{[31]}s_{[211]}c^q, \\
|8\rangle & : \text{sc}q[4]^A_4[x_{[4]}]_p_{[31]}s_{[211]}c^q,
\end{align*}
\]

(6)

and for \( J^P = 3/2^- \) or \( 5/2^- \) are

\[
\begin{align*}
|9\rangle & : \text{sc}c\text{c}q[4]^A_4[x_{[31]}]_3^p_{[4]}s_{[211]}c^q, \\
|10\rangle & : \text{sc}c\text{c}q[4]^A_4[x_{[31]}]_3^p_{[31]}s_{[211]}c^q,
\end{align*}
\]

(7)

respectively.

One should note that the different spin symmetries of the four-quark subsystem result in vanishing coupling between different five-quark configurations, this is another reason for us to categorize the states in three groups by the four-quark spin wave functions.

**B. The chiral Lagrangian approach**

We consider the decays \( \text{sc}c\text{c}q \rightarrow MB \), which mainly proceed through the process of \( q \bar{q} \rightarrow M \), where the final baryon and meson are assumed to be composed of three-quark and a quark-antiquark pair, respectively. We name this kind of decays as the annihilation transitions. The \( \text{sc}c\text{c}q \rightarrow \bar{K}\Xi \) and \( D\Xi \) transitions are shown in Fig. 1.

To compute the transitions of the \( \Omega_{cc} \rightarrow \bar{K}\Xi \) and \( \Omega_{cc} \rightarrow D\Xi \) shown in Fig. 1 we use the chiral Lagrangian approach. Within this approach, the quark pseudoscalar (P) and vector (V) meson couplings are

\[
H_{eff}^{Pqq} = \sum_j \bar{\psi}_j (a \gamma_\mu + \frac{i b \sigma_{\mu\nu} k^\nu}{2m_j}) \phi^\mu \psi_j,
\]

\[
H_{eff}^{Vqq} = \sum_j \bar{\psi}_j (a \gamma_\mu + \frac{i b \sigma_{\mu\nu} k^\nu}{2m_j}) \phi^\mu \psi_j,
\]

(8)

(9)

respectively, where the summation on \( j \) runs over the quark in the initial hadron. \( \psi_j \) represents the quark field, and \( \phi^\mu \) are the pseudoscalar and vector meson fields. \( m_j \) is the constituent mass of the \( j \)th quark, while \( k^\mu_M \) denotes the four-momentum of the vector meson. \( a \) and \( b \) are the vector and tensor coupling constants, respectively.

In the non-relativistic approximation, Eqs. (8) and (9) lead to the operators for process involving \( q \rightarrow q'M \) transitions as following

\[
\begin{align*}
T_{d,P}^{Pqq} &= \sum_j \left( \frac{\omega_M}{E_f + M_f} \sigma_j \cdot \vec{P}_f + \frac{\omega_M}{E_i + M_i} \sigma_j \cdot \vec{P}_i - \sigma_j \cdot \vec{k}_M + \frac{\omega_M}{2 \mu_q} \sigma_j \cdot \vec{p}_j \right) X_p^\mu \exp(-i \vec{k}_M \cdot \vec{r}_j), \\
T_{d,T}^{Vqq} &= \sum_j \left\{ \frac{i b'}{2m_j} \sigma_j \cdot (\vec{k}_M \times \vec{e}) + \frac{a}{2 \mu_q} \vec{p}_j \cdot \vec{e} \right\} X_p^\mu \exp(-i \vec{k}_M \cdot \vec{r}_j), \\
T_{d,L}^{Vqq} &= \sum_j \frac{a M_V}{|k_M|} X_p^\mu \exp(-i \vec{k}_M \cdot \vec{r}_j),
\end{align*}
\]

(10)

(11)

(12)

where \( \omega_M \) and \( \vec{k}_M \) are the energy and three-momentum of the final meson, \( E_i(f), M_i(f) \) and \( \vec{P}_i(f) \) are the energy, mass and three-momentum of the initial (final) baryon, while \( \vec{p}_j, \vec{r}_j \) and \( m_j \) are the momentum, coordinate and constituent mass of the quark which emits a meson. The \( \mu_q \) is the reduced mass of the \( j \)th quark before and after emitting the meson. For the vector meson emission, in Eqs. (11) and (12), the transition operators are denoted as \( T_{d,T}^{Vqq} \) and \( T_{d,L}^{Vqq} \) for the meson being transversely and longitudinally polarized, respectively. The \( b' \) in Eq. (11) is defined by \( b' = b - a \). \( M_V \) is the mass of the vector meson, and the polarization vectors of the final vector meson are taken to be

\[
\begin{align*}
\epsilon^L_\mu &= \frac{1}{M_V} \left( \frac{\vec{k}_M}{E_V} \frac{\vec{k}_M}{|k_M|} \right) , \\
\epsilon^T_\mu &= \left( 0 \ \vec{e} \right),
\end{align*}
\]

(13)

with

\[
\vec{e}(\pm) = \frac{1}{\sqrt{2}}(\mp 1, -i, 0)^T,
\]

(14)

and \( E_V \) is the energy of the final vector meson.

Finally, \( X_p^\mu \) and \( X_p^\mu \) are the operators in flavor space for a pseudoscalar and vector meson emission, which only depends on the quark-antiquark content of the emitted meson. For instance, \( X_p^\mu \) for a light pseudoscalar meson
emission in Eq. (10) can be defined as

\[
X^j_{π^±} = \mp \frac{1}{\sqrt{2}} (\lambda^j_1 \mp i \lambda^j_2), \\
X^j_{π^0} = \lambda^j_3, \\
X^j_{K^±} = \mp \frac{1}{\sqrt{2}} (\lambda^j_4 \mp i \lambda^j_5), \\
X^j_{K^0,K^±} = \mp \frac{1}{\sqrt{2}} (\lambda^j_6 \mp i \lambda^j_7), \\
X^j_η = \cos θ λ^j_8 - \sin θ \sqrt{\frac{2}{3}} I^j, \\
X^j_η' = \sin θ λ^j_8 + \cos θ \sqrt{\frac{2}{3}} I^j,
\]

with \(λ^j_i\) and \(I\) the Gell-Mann matrix and unit matrix in flavor space. \(θ\) denotes the mixing angle for the mixing between \(η_1\) and \(η_8\), leading to the physical states \(η\) and \(η'\)

\[
\begin{align*}
η &= η_8 \cos θ - η_1 \sin θ \\
η' &= η_8 \sin θ + η_1 \cos θ,
\end{align*}
\]

where the empirical value for the mixing angle is \(θ = -23°\). The flavor operators for other pseudoscalar mesons or the vector mesons can be obtained straightforward.

Accordingly, the transition operators for a pseudoscalar meson emission \(T^p_{aqq}\), a transversely polarized vector meson emission \(T^v_{a'T}\) and a longitudinally polarized vector meson emission can be obtained as

\[
\begin{align*}
T^p_{aqq} &= \sum_j (m_j + m_b)C^j_{X_FSC}\chi^j_{1} I^j_2 \chi^j_{2} X^j \exp\{-i\vec{k}_M \cdot (\vec{r}_j + \vec{r}_q)/2\}, \\
T^v_{a'T} &= \sum_j \left\{ a - \frac{m_j + m_b}{2m_j} \right\} \bar{σ} \cdot \bar{c} X^j V \exp\{-i\vec{k}_M \cdot (\vec{r}_j + \vec{r}_q)/2\}, \\
T^v_{a'L} &= \sum_j \left\{ a - \frac{m_j + m_b}{2m_j} \right\} \frac{E_V \bar{σ} \cdot \bar{k}_M}{M\sqrt{|\vec{k}_M|}} X^j V \exp\{-i\vec{k}_M \cdot (\vec{r}_j + \vec{r}_q)/2\},
\end{align*}
\]

where \(m_j\) and \(m_q\) are the constituent masses of the \(j\)th quark and the antiquark, respectively. \(C^j_{X_FSC}\) denotes the operator to calculate the orbital, flavor, spin and color overlap factor between the residual wave function of the pentaquark configuration after the quark-antiquark annihilation and the wave function of the final baryon. \(\chi^j_{1} I^j_2 \chi^j_{2}\) is the spin operator for the quark-antiquark annihilation.

III. NUMERICAL RESULTS AND DISCUSSIONS

In this section, we present our theoretical results for the mass spectrum of the low-lying \(scqq\) states with \(J^P = \frac{3}{2}^-\), \(\frac{5}{2}^-\) and \(\frac{7}{2}^-\), and the decay behaviours of the obtained \(Ω_{1cc}\) pentaquark states.
A. The mass spectrum of the low-lying \( scc\bar{q}\bar{q} \) states

With the Hamiltonian in Eq. (1) and wave function in Eq. (2), one can obtain the following nonzero \( H_{ij} = \langle i|H|j \rangle \) matrix elements

\[
\begin{align*}
H_{11} &= E_0 - 7.50C_D - 7.50C_{D_s}, \\
H_{21} &= E_{21} = 4.90C_D - 4.90C_{D_s}, \\
H_{31} &= H_{31} = -4.33C_K + 0.87C_D + 0.87C_{D_s} + 2.60C_{c\bar{c}}, \\
H_{32} &= H_{32} = -5.66C_D + 5.66C_{D_s}, \\
H_{33} &= E_0 - 5.0C_D - 2.50C_{D_s} + 3.0C_{c\bar{c}},
\end{align*}
\]

between the pentaquark configurations in Eqs. (3), and

\[
\begin{align*}
H_{44} &= E_0 - 2.5C_K - 4.50C_D - 4.50C_{D_s} - 1.5C_{c\bar{c}}, \\
H_{45} &= H_{54} = -6.12C_D + 6.12C_{D_s}, \\
H_{46} &= H_{64} = -3.54C_D + 3.54C_{D_s}, \\
H_{47} &= H_{74} = 5.00C_K - 2.50C_D - 2.50C_{D_s}, \\
H_{55} &= E_0 - 0.50C_K - 4.00C_D - 4.00C_{D_s} - 0.5C_{c\bar{c}}, \\
H_{56} &= H_{65} = 1.73C_K - 1.73C_{c\bar{c}}, \\
H_{57} &= H_{75} = -1.22C_D + 1.22C_{D_s}, \\
H_{58} &= H_{85} = -1.41C_K + 1.41C_D + 1.41C_{D_s} - 1.41C_{c\bar{c}}, \\
H_{66} &= E_0 + 1.50C_K - 4.0C_D - 4.0C_{D_s} + 1.50C_{c\bar{c}}, \\
H_{67} &= H_{76} = -0.71C_D + 0.71C_{D_s}, \\
H_{68} &= H_{86} = -2.45C_K + 2.45C_{c\bar{c}}, \\
H_{77} &= E_0 - 2.5C_K - 0.50C_D - 0.50C_{D_s} - 1.5C_{c\bar{c}}, \\
H_{78} &= H_{87} = -3.46C_D + 3.46C_{D_s}, \\
H_{88} &= E_0 + 0.50C_K + 1.0C_D + 1.0C_{D_s} + 0.50C_{c\bar{c}},
\end{align*}
\]

between the pentaquark configurations in Eqs. (4), and

\[
\begin{align*}
H_{99} &= E_0 - 1.50C_K + 1.0C_D + 1.0C_{D_s} - 1.50C_{c\bar{c}}, \\
H_{910} &= H_{109} = 2.83C_D - 2.83C_{D_s}, \\
H_{1010} &= E_0 + 2.50C_K - C_D - C_{D_s} - 1.50C_{c\bar{c}}.
\end{align*}
\]

between the pentaquark configurations in Eq. (4).

In the above equations, \( C_M \) are the corresponding matrix elements of the hyperfine interaction coupling strength \( V_M(r_{ij}) \) between the S-wave orbital wave functions of the quarks in \( scc\bar{q}\bar{q} \) system, namely

\[
C_M = \langle \phi(\{\bar{q}_j\})|V_M(r_{ij})|\phi(\{\bar{q}_j\}) \rangle.
\]

\( C_{c\bar{c}} \) is obtained from the last term in Eq. (2), and it contains the exchanges of the \( uu, dd, ss \) and \( cc \) pairs. The coupling strength constants \( C_M \) are taken to be the empirical values [53] as shown in Table I.

With the above values for the model parameters and the diagonalization of the matrices obtained by Eqs. (24), one can get the physical states which are shown in Fig. 2 while the explicit probability amplitudes are shown in Table I. For instances, Eq. (25) leads to the following energy matrix

\[
\begin{pmatrix}
C_\pi & C_K & C_{s\bar{s}} & C_D & C_{D_s} & C_{c\bar{c}} \\
21.0 & 15.5 & 11.5 & 6.5 & 6.5 & 0
\end{pmatrix}
\]

Then one can directly obtain the eigenvalues and eigenvectors of matrix in Eq. (29). The three obtained eigenvalues are the energies of the physical states \( |i'\rangle \) with \( i = 1, 2, 3 \), respectively, and a obtained eigenvector just show the coefficients for the decoupling of a corresponding physical state \( |i\rangle \) to the configurations \( |i'\rangle \) listed in Eq. (29).

The energies for the obtained \( \Omega_{cc} \) states in present work are at 4050 ± 100 MeV. Similar as the results for \( \Omega_c^0 \) obtained in Ref. [12], mixing between the pentaquark configurations \( |i\rangle \) caused by the goldstone boson exchange is strong, while the mass splitting for the obtained states \( |i'\rangle \) is not very large. On the other hand, the spectrum of the ten obtained states is not sensitive to the values of the coupling strength \( C_M \).

Up to now, there are no solid experimental data for the \( \Omega_{cc} \) resonances, while theoretical investigations on the doubly heavy baryon resonances have been intensively taken using various of approaches, such as the constituent quark model [5,10,60], QCD sum rules [17], chiral perturbation theory [20,22], the unitarized coupled channel approach [61], and the lattice QCD calculations [23], etc. The corresponding obtained energies for the \( P \)-wave \( \Omega_{cc} \) in a three-quark picture are around 4000 – 4200 MeV in most of the literatures, and one may note that in Ref. [61], the S-wave interactions between pseudo-Nambu-Goldstone bosons (\( \pi, K \) and \( \eta \)) and the
$J^P = 1/2^+$ ground state doubly charmed baryons in the energy region around the corresponding thresholds are investigated, two quasistable narrow $J^P = 1/2^- \Omega_{cc}$ are predicted to lie at the energy below 4200 MeV, and their strong decay mode is predicted to be only the $\Omega_{cc}\pi^0$, which is isospin breaking channel. Therefore, the two obtained $\Omega_{cc}$ resonances to be dominated by three-quark configurations:

$$|\text{scc}, 1\rangle = [21]_X [21]_F [21]_S [21]_F [21]_S [13]_C,$$

$$|\text{scc}, 2\rangle = [21]_X [21]_F [21]_S [3]_F [21]_S [13]_C,$$

$$|\text{scc}, 3\rangle = [21]_X [21]_F [21]_S [21]_F [3]_S [13]_C,$$

whose spin-parity quantum number $J^P$ may be $1/2^-$ or $3/2^-$ for the first two configurations, and $1/2^-, 3/2^-$ or $5/2^-$ for the latest one. Direct calculations employing the chiral constituent quark model as in Ref. [54] lead to the following values for the energies of the three $\Omega_{cc}$ states,

$$E_1 = 4219 \text{ MeV}, \quad E_2 = 4246 \text{ MeV}, \quad E_3 = 4257 \text{ MeV},$$

respectively. Consequently, the energies of low-lying $\Omega_{cc}$ states in the five-quark picture are lower than those in the three-quark picture, this conclusion is the same as that for the $\Omega^*$ resonances [37].

In Ref. [54], a relativistic quark model was applied to study the spectrum of doubly heavy baryons. Considering the $\Omega_{cc}$ resonances to be dominated by three-quark components, it was obtained that the low-lying $\Omega_{cc}$ resonances with negative parity fall in the range of $4200 - 4300 \text{ MeV}$, which are consistent with the results obtained in Ref. [10] by employing a three-quark model.

While in Ref. [5], a three-quark model was employed to investigate the spectrum of the doubly heavy baryons, in which model the two heavy quarks were treated as a diquark, and the resulting energies of the low-lying $\Omega_{cc}$ were in the range of $4050 - 4150 \text{ MeV}$. Those results are about 100 MeV lower than the present rough estimation using a three-quark model, and the results in [10, 56]. So one may expect that the diquark assumption for the two heavy quark in $\Omega_{cc}$ resonances may reduce the energies.

In any case, we can conclude that the $\Omega_{cc}$ resonances should lie at a energy below 4200 MeV in both the compact five-quark model (present) and the meson-baryon model [61].

Finally, we show the dependency of presently obtained spectrum on the model parameter $E_0$. By taking $E_0 = 3132 \text{ MeV}$ as given in Ref. [13], one can get

$$E_{i'} \simeq 3075 \pm 100 \text{ MeV}. \quad (31)$$

Obviously, the obtained energies are much lower than those predicted by using other approaches. Namely, the value $E_0 = 4102 \text{ MeV}$ employed in our calculations should be reasonable. We also present the numerical results with $E_0$ changed by 2%.

$$E_{i'} \simeq 4125 \pm 100 \text{ MeV}. \quad (32)$$

In fact, change of $E_0$ should lead to almost the same change for energy of each physical state $|i\rangle$. In addition, the coefficients for decompositions of the physical states $|i\rangle$ are not sensitive to $E_0$.

### B. S-wave coupling of the $sccq\bar{q}$ to pseudoscalar meson and ground baryon states

From Fig. [2] and Table [11] one can find that most of the obtained physical $sccq\bar{q}$ states are above the threshold of...
the $SU(2)$ isospin breaking $\pi\Omega_{cc}$ channel, but, below the
thresholds of the other pseudoscalar meson and ground state baryons channels. It is expected that the decay
widths of the presently obtained $\Omega_{cc}$ should not be very large. This is in consistent with these findings in Ref. [61]. Therefore, in present work, we only try to estimate the partial
widths of the obtained $\Omega_{cc}$ resonances to $MB$ channels. 

Using the transition operator in Eq. (22), and the wave
functions obtained in Sec. IIIA one can calculate the
transition matrix elements of the obtained $\Omega_{cc}$ resonances
to $\bar{K}\Xi_{cc}$, $K\Xi_{cc}$, $D\Xi_{cc}$, and $D\Xi_{cc}'$ channels, respectively. It
is found that all the transition amplitudes of $\Omega_{cc} \rightarrow MB$
processes share a common factor which involves the overlap
between the orbital wave functions of the pentaquark
configurations and the final meson-baryon, namely,

$$\mathcal{F}(\vec{k}_M^2) = \langle \phi_B(\{\vec{\xi}_i\})|\exp\{-i\vec{k}_M \cdot (\vec{r}_j + \vec{r}_q)/2\}|\varphi(\{\vec{\xi}_i\})\rangle,$$  (33)

which is shown in Table III.

Compared to the overlap factor [49, 50] for the
strangeness five-quark configurations $|uuds\bar{s}\rangle$ to $np$ channels that is about ~ 0.75, which may account for the
strong coupling between $S_{11}(1535)$ and strangeness
canals, one can expect that the presently obtained physical
states $|i'\rangle$ may couple strongly to the $MB$ channels for
which the overlap factors shown in Table IV are larger
than 0.8.

It should be very interesting to compare the decay be-
behaviours of the presently obtained $\Omega_{cc}$ resonances with
those in a three-quark model. In Ref. [61], five $\Omega_{cc}$
resonances lying at 4208 – 4303 MeV were obtained, and the
decay widths of these resonances to $\bar{K}\Xi_{cc}$ or $K\Xi_{cc}$
channels were estimated explicitly. It was found that some of the obtained decay widths should be larger
than 100 MeV. This is very different from the conclusion that
most of the obtained $\Omega_{cc}$ resonances using a pentaquark
picture can only decay to the isospin breaking channel
$\pi\Omega_{cc}$.

In a three-quark picture, one can also estimate the
flavor-spin-color overlap factor of the $\Omega_{cc}$ resonances and the $MB$ channels using Eq. (10). For instance, a straight-
ward calculation on the overlap factors of the three-
quark states given in Eq. (30), shows that coupling for
$\Omega_{cc} \rightarrow K\Xi_{cc}$ may be comparable to that for $\Omega_{cc} \rightarrow D\Xi_{cc}$,
since the obtained flavor-spin-color overlap factor for a
given three-quark $\Omega_{cc}$ resonance to the $K\Xi_{cc}$ channel is
$\sqrt{2}$ times of that for the $D\Xi_{cc}$ channel, this is determined
by the flavor-spin structure of the $\Omega_{cc}$ resonances and
the effective chiral Lagrangian. However, as we can see in
Table IV, the presently obtained numerical results for
several states are very different from the three-quark re-

IV. SUMMARY

In present work, we investigate the spectrum of low-
lyliyng $\Omega_{cc}$ resonances with negative parity as pentaquark
states, using the chiral constituent quark model within
a five-quark picture. We obtain ten pentaquark states
TABLE IV: Color-flavor-spin overlap factors for the transitions \( J^P = 1/2^- \) (shown in the 2nd to 4th rows) and \( 3/2^- \) (shown in the last two rows) \( \Xi_{cc} \) physical states \( |i'\rangle \) to \( MB \) channels.

| \( i' \) | \( J^P \) | \( |1\rangle \) | \( |2\rangle \) | \( |3\rangle \) | \( |4\rangle \) | \( |5\rangle \) | \( |6\rangle \) | \( |7\rangle \) | \( |8\rangle \) | \( |9\rangle \) | \( |10\rangle \) |
|---|---|---|---|---|---|---|---|---|---|---|---|
| \( K\Xi_{cc} \) | \( -1 \) | \( \sqrt{6}/3 \) | \( \sqrt{3}/3 \) | \( \sqrt{3}/2 \) | \( \sqrt{6}/3 \) | \( \sqrt{3}/3 \) | \( 0 \) | \( - \) | \( - \) | \( - \) | \( - \) |
| \( D\Xi_{c} \) | \( 0 \) | \( \sqrt{12}/3 \) | \( 0 \) | \( 0 \) | \( -2 \) | \( \sqrt{12}/3 \) | \( 0 \) | \( 0 \) | \( - \) | \( - \) | \( - \) |
| \( D\Xi_{c}^* \) | \( \sqrt{5}/3 \) | \( 0 \) | \( \sqrt{2} \) | \( -\sqrt{2} \) | \( 0 \) | \( 0 \) | \( \sqrt{2} \) | \( 0 \) | \( - \) | \( - \) | \( - \) |
| \( K\Xi_{cc}^* \) | \( - \) | \( - \) | \( - \) | \( - \) | \( 0 \) | \( 0 \) | \( \sqrt{12}/3 \) | \( 0 \) | \( \sqrt{2} \) | \( -\sqrt{30}/3 \) | \( 0 \) |
| \( D\Xi_{c}^* \) | \( - \) | \( - \) | \( - \) | \( - \) | \( 0 \) | \( 0 \) | \( \sqrt{12}/3 \) | \( 0 \) | \( \sqrt{2} \) | \( -\sqrt{30}/3 \) | \( 0 \) |

with spin-parity \( J^P = 1/2^- \), \( 3/2^- \), \( 5/2^- \), which lie at 4050 ± 100 MeV. Most of the obtained states are above the isospin breaking decay channel \( \pi \Omega_{cc} \), but below the other meson-baryon channels. So we just try to calculate the flavor, spin, orbital, and color overlap factor for the final \( MB \) states and the residual three-quark-meson configurations of the \( \Omega_{cc} \) states after the annihilation of the quark-antiquark \( q\bar{q} \rightarrow M \). It is found that several ones of the presently obtained \( \Omega_{cc} \) may couple strongly to \( D\Xi_{c} \) or \( K\Xi_{cc} \) channels. One may expect that these calculations here could be compared with the future experimental measurements which are likely to be done by Belle II and/or LHCb.

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