Investigating the excited Ω_c^0 states through $\Xi_c \bar{K}$ and $\Xi'_c \bar{K}$ decay channels

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Inspired by the five newly observed Ω_c^0 states by the LHCb detector, we study the Ω_c^0 states as the S -wave molecular pentaquarks with $I = 0$, $J^P = \frac{1}{2}^-$, $\frac{3}{2}^-$, and $\frac{5}{2}^-$ by solving the renormalization group method equation in the framework of the chiral quark model. Both the energies and the decay widths are obtained in this work. Our results suggest that $\Omega_c(3119)^0$ can be explained as an S -wave resonance state of ΞD with $J^P = \frac{1}{2}^-$, and the decay channels are the S -wave $\Xi_c \bar{K}$ and $\Xi'_c \bar{K}$. Other reported Ω_c^0 states cannot be obtained in our present calculation. Another Ω_c^0 state with much higher mass 3533 MeV with $J^P = \frac{5}{2}^-$ is also obtained. In addition, the calculation is extended to the Ω_b^0 states, and results similar to those of Ω_c^0 are obtained.

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I. INTRODUCTION

There has been important experimental progress in the sector of heavy baryons in the past decade. Many heavy baryons have been reported. For example, the triplet of excited Σ_c baryons, $\Sigma_c(2800)$, was observed by Belle [1] in 2005, and it tentatively identified the quantum numbers of these states as $J^P = \frac{3}{2}^-$. In 2008, the same neutral state Σ_c^0 was also observed by the *BABAR* Collaboration with the mean value of mass higher than that obtained by Belle [2]. The charmed baryons $\Lambda_c(2880)^+$ and $\Lambda_c(2940)^+$ were observed by both the *BABAR* and Belle Collaborations in 2007 [3,4]. The charm-strange baryons $\Xi_c(2980)^+$ and $\Xi_c(3077)^0$ were reported by Belle [5] and later confirmed by *BABAR* [6]. $\Xi_c(3055)^+$ and $\Xi_c(3123)^+$ were also investigated by *BABAR* [6]. Among the expected charmed baryons, the spectrum of the Ω_c^0 baryons, which have the quark content of ssc , is still unknown. Only two states, $\Omega_c(2695)^0$ and $\Omega_c(2768)^0$ with $J^P = 1/2^+$ and $J^P = 3/2^+$, respectively, have been observed before [7,8]. Very recently, the LHCb Collaboration reported five new narrow Ω_c^0 states in the $\Xi_c^+ K^-$ invariant mass spectrum. They are $\Omega_c(3000)^0$, $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ [9]. Moreover, the decay widths of these states

were also observed by the experiment, which are only a few MeV. However, the quantum numbers and the structures of these states are still unclear now.

All this experimental progress of heavy baryons has stimulated extensive interest in understanding the structures of the charmed baryons. A classical way to describe the charmed baryons is based on the assumption that they are conventional three-quark charmed baryons. And another way is treating them as candidates of molecular states. We take the states $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ as the example. Being considered as two traditional charmed baryons, the strong decays of these two states have been studied by using the heavy hadron chiral perturbation theory [10], the 3P_0 model [11], and the chiral quark model [12]. On the other hand, many works treat them as candidates of molecular states. J. R. Zhang found that $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ as the S -wave ND state and ND^* state, respectively, by means of QCD sum rules [13]. He and Liu explained the $\Lambda_c(2940)^+$ as an isoscalar S -wave or P -wave D^*N system within the one boson-exchange model [14]. For the newly observed Ω_c^0 states, some works treat them as traditional charmed baryons. Agaev *et al.* calculated the masses and the residues of these states with $J^P = 1/2^+$ and $J^P = 3/2^+$ in the framework of QCD two-point sum rules, and they were inclined to assign the $\Omega_c(2066)^0$ and $\Omega_c(3119)^0$ states as the first radially excited ($2S, 1/2^+$) and ($2S, 3/2^+$) charmed baryons [15]. Chen *et al.* studied the decay properties of the P -wave charmed baryons within the light-cone QCD sum rules, including some Λ_c , Σ_c , and Ξ_c states as well as these newly reported Ω_c^0 states [16]. They interpreted one of these Ω_c^0 states was a $J^P = 1/2^-$ state, two of them were a $J^P = 3/2^-$ state and a

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$J^P = 5/2^-$ state, and another two may have been with $J^P = 1/2^+$ and $3/2^+$. Karliner and Rosner explained these Ω_c^0 baryons as bound states of a P -wave ss diquark and a c quark [17], and they predicted two of spin 1/2, two of spin 3/2, and one of spin 5/2, all with negative parity. Wang *et al.* investigated the strong and radiative decay properties of the low-lying Ω_c^0 states in a constituent quark model [18]. Their results show that the $\Omega_c(3000)^0$ and $\Omega_c(3090)^0$ can be assigned to have $J^P = 1/2^-$, the $\Omega_c(3050)^0$ can be assigned to have $J^P = 3/2^-$, the $\Omega_c(3066)^0$ can be assigned to have $J^P = 5/2^-$, and the $\Omega_c(3119)^0$ might be one of the two $2S$ states of the first radial excitations. Another way to describe these states is by assuming they are pentaquark states. Kim *et al.* employed the chiral quark-soliton model to describe excited baryons with one heavy quark and found two out of five narrow Ω_c^0 states may correspond to the exotic pentaquarks [19]. An *et al.* investigated the spectrum of several low-lying $sscq\bar{q}$ pentaquark configurations by using the constituent quark model and showed that the $sscq\bar{q}$ pentaquark configurations may form sizable components of the observed Ω_c^0 states [20]. Besides, both Refs. [21,22] interpreted some Ω_c^0 excited baryons as meson-baryon molecular states. To study the structure of the observed Ω_c^0 , G. Yang *et al.* did a dynamical calculation of the Ω_c^0 within the chiral quark model by taking the advantage of Gaussian expansion method [23]. Both the $3q$ configurations and five-quark configurations were studied in this work, and the results suggested that two of the observed states, $\Omega_c(3090)^0$ and $\Omega_c(3119)^0$, may be explained as the $2S$ states of the ssc system and the $\Xi\bar{D}$, $\Xi_c\bar{K}$ and $\Xi_c^*\bar{K}$ are possibly the candidates of some newly reported Ω_c^0 states by LHCb.

Actually, the hadron-hadron scattering is one of the important ways to generate and identify multi-quark states. Therefore, to provide the necessary information for experiment to search for the multi-quark states, we should not only calculate the mass spectrum but also study the corresponding scattering process. The multi-quark state will appear as a resonance in the scattering phase shifts. By using the constituent quark models and the resonating group method (RGM) [24], we have obtained the d^* resonance during the NN scattering process, and we found that the energy and decay width of the partial wave of NN are consistent with the experiment data [25]. Extending to the pentaquark system, we investigated the $N\phi$ state in the different scattering channels: $N\eta'$, ΛK , and ΣK [26]. Both the resonance mass and decay width were obtained, which provided the necessary information for experimental searching at Jefferson Lab. Therefore, it is interesting to extend such a study to the newly observed Ω_c^0 states. In this work, we assume Ω_c^0 states as pentaquark states, calculate both the masses and decay widths of these states, and analyze if there are some Ω_c^0 states which can be explained as pentaquarks by comparing with the LHCb data. Finally,

we will also extend the study to the Ω_b^0 states because of the heavy flavor symmetry.

The structure of this paper is as follows. A brief introduction of a constituent quark model and the calculation method is given in Sec. II. Section III is devoted to the numerical results and discussions. The summary is shown in the last section.

II. QUARK MODEL AND CALCULATION METHOD

Here, we use the chiral quark model to study the Ω_c^0 states. The Salamanca version was chosen as the representative of the chiral quark models, because the Salamanca group's work covers the hadron spectra, nucleon-nucleon interaction, and multi-quark states. We also have used this model to study the nucleon-nucleon interaction, dibaryon resonance states, such as the d^* state [25], $N\Omega$ [27,28], and so on. In this model, the constituent quarks interact with each other through the one-gluon exchange and the Goldstone-boson exchange in addition to the color confinement. For the system with strangeness, the model is extended to include full SU(3) scalar nonet meson exchange instead of one chiral partner σ meson [29,30]. These scalar potentials have the same functional form as the one of SU(2) chiral quark model (ChQM) but a different SU(3) operator dependence [29]. The model details can be found in Ref. [31]. Here, we only give the Hamiltonian,

$$H = \sum_{i=1}^5 \left(m_i + \frac{p_i^2}{2m_i} \right) - T_{CM} + \sum_{j>i=1}^5 (V_{ij}^C + V_{ij}^G + V_{ij}^\chi + V_{ij}^\sigma), \quad (1)$$

$$V_{ij}^C = -a_c \lambda_i^c \cdot \lambda_j^c (r_{ij}^2 + v_0), \quad (2)$$

$$V_{ij}^G = \frac{1}{4} \alpha_s \lambda_i^c \cdot \lambda_j^c \left[\frac{1}{r_{ij}} - \frac{\pi}{2} \delta(r_{ij}) \left(\frac{1}{m_i^2} + \frac{1}{m_j^2} + \frac{4\sigma_i \cdot \sigma_j}{3m_i m_j} \right) - \frac{3}{4m_i m_j r_{ij}^3} S_{ij} \right] \quad (3)$$

$$V_{ij}^\chi = V_\pi(\mathbf{r}_{ij}) \sum_{a=1}^3 \lambda_i^a \cdot \lambda_j^a + V_K(\mathbf{r}_{ij}) \sum_{a=4}^7 \lambda_i^a \cdot \lambda_j^a + V_\eta(\mathbf{r}_{ij}) [(\lambda_i^8 \cdot \lambda_j^8) \cos \theta_P - (\lambda_i^0 \cdot \lambda_j^0) \sin \theta_P] \quad (4)$$

$$V_\chi(\mathbf{r}_{ij}) = \frac{g_{ch}^2}{4\pi} \frac{m_\chi^2}{12m_i m_j} \frac{\Lambda_\chi^2}{\Lambda_\chi^2 - m_\chi^2} m_\chi \times \left\{ (\sigma_i \cdot \sigma_j) \left[Y(m_\chi r_{ij}) - \frac{\Lambda_\chi^3}{m_\chi^3} Y(\Lambda_\chi r_{ij}) \right] + \left[H(m_\chi r_{ij}) - \frac{\Lambda_\chi^3}{m_\chi^3} H(\Lambda_\chi r_{ij}) \right] S_{ij} \right\}, \quad (5)$$

$\chi = \pi, K, \eta,$

TABLE I. Model parameters: $m_\pi = 0.7 \text{ fm}^{-1}$, $m_k = 2.51 \text{ fm}^{-1}$, $m_\eta = 2.77 \text{ fm}^{-1}$, $m_\sigma = 3.42 \text{ fm}^{-1}$, $m_{a_0} = m_\kappa = m_{f_0} = 4.97 \text{ fm}^{-1}$, $\Lambda_\pi = \Lambda_\sigma = 4.2 \text{ fm}^{-1}$, $\Lambda_K = \Lambda_\eta = \Lambda_{a_0} = \Lambda_\kappa = \Lambda_{f_0} = 5.2 \text{ fm}^{-1}$, and $g_{ch}^2/(4\pi) = 0.54$, $\theta_p = -15^\circ$.

b (fm)	m_u (MeV)	m_d (MeV)	m_s (MeV)	m_c (MeV)	m_b (MeV)
0.518	313	313	450	1635	4988
a_c (MeV fm $^{-2}$)	V_0 (MeV)	α_{su}	α_{us}	α_{sc}	α_{sb}
48.59	-0.961	0.84	0.82	0.35	0.27
α_{ss}	α_{sc}	α_{sb}	α_{cc}	α_{bb}	
0.30	0.25	0.20	0.20	0.15	

$$V_{ij}^{\sigma_a} = V_{a_0}(\mathbf{r}_{ij}) \sum_{a=1}^3 \lambda_i^a \cdot \lambda_j^a + V_\kappa(\mathbf{r}_{ij}) \sum_{a=4}^7 \lambda_i^a \cdot \lambda_j^a + V_{f_0}(\mathbf{r}_{ij}) \lambda_i^8 \cdot \lambda_j^8 + V_\sigma(\mathbf{r}_{ij}) \lambda_i^0 \cdot \lambda_j^0 \quad (6)$$

$$V_k(\mathbf{r}_{ij}) = -\frac{g_{ch}^2}{4\pi} \frac{\Lambda_k^2 m_k}{\Lambda_k^2 - m_k^2} \left[Y(m_k r_{ij}) - \frac{\Lambda_k}{m_k} Y(\Lambda_k r_{ij}) \right], \quad (7)$$

$$k = a_0, \kappa, f_0, \sigma,$$

$$S_{ij} = \left\{ 3 \frac{(\boldsymbol{\sigma}_i \cdot \mathbf{r}_{ij})(\boldsymbol{\sigma}_j \cdot \mathbf{r}_{ij})}{r_{ij}^2} - \boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j \right\}, \quad (8)$$

$$H(x) = (1 + 3/x + 3/x^2)Y(x), \quad Y(x) = e^{-x}/x, \quad (9)$$

where S_{ij} is the quark tensor operator; $Y(x)$ and $H(x)$ are standard Yukawa functions; T_c is the kinetic energy of the center of mass; α_s is the quark-gluon coupling constant; and g_{ch} is the coupling constant for chiral field, which is determined from the $NN\pi$ coupling constant through

$$\frac{g_{ch}^2}{4\pi} = \left(\frac{3}{5}\right)^2 \frac{g_{\pi NN}^2}{4\pi} \frac{m_{u,d}^2}{m_N^2}. \quad (10)$$

The other symbols in the above expressions have their usual meanings. Although the same quark model is used in both Ref. [23] and this work, there are some differences in the Hamiltonian. The first one is the confinement interaction. In this work, we use the quadratic confinement potential from Ref. [31], while in Ref. [23], a screened potential which was from Ref. [32] was used. We have found that the difference is small by using these two kinds of confinement potential when dealing with low-excited states, and the discrepancy will come up in the high-excited states. The second one is the $\delta(\mathbf{r}_{ij})$ function in one gluon-exchange potential, which has been regularized in another way in Ref. [23]. This helps one to avoid obtaining an unbound baryon spectrum when solving the Schrödinger equation [32,33]. The last one is the tensor interaction, which is included in this work but not considered in Ref. [23]. The tensor force is not responding to the

S -wave state, but it is necessary to couple the S - and D -wave channels.

Generally, we use the parameters from our former work of dibaryons [28]; only the mass of charm quark is adjusted to fit the charmed mesons and baryons used in this work. However, the former parameters can describe the ground baryons well but cannot fit the ground mesons, especially the K meson, the mass of which is much higher than the experimental value. This situation will lead to a consequence that some bound states cannot decay to the open channel $\Xi_c K$, because of the much larger mass of K . To solve this problem, we adjust the parameters which are related to s and c quarks in this work and keep the parameters which are related to u and d quarks unchanged. By doing this, the parameters can describe the nucleon-nucleon interaction well, and at the same time, it will lower the mass of K . The model parameters are fixed by fitting the spectrum of baryons and mesons we used in this work. All the parameters of the Hamiltonian are given in Table I. The calculated masses of baryons and mesons in comparison with experimental values are shown in Table II. One can see the parameters are also different from those used in Ref. [23]. It is not only because some different potentials are used in these two works, but also different calculation methods, which are introduced below, are used.

Here, we investigate the Ω_c^0 states as the molecular pentaquarks by using the RGM [24], a well-established method for studying a bound-state problem or a scattering

TABLE II. The masses of baryons and mesons used in this work (in MeV).

	Ξ	Ξ^*	Ξ_c	Ξ'_c	Ξ_c^*	Ω_c	Ω_c^*
Experimental value	1318	1533	2469	2577	2646	2695	2766
ChQM	1225	1359	2448	2527	2543	2662	2672
	K	K^*	η	ω	D	D^*	
Experimental value	495	892	548	783	1864	2007	
ChQM	593	830	523	702	1980	2007	
	Ξ_b	Ξ'_b	Ξ_b^*	Ω_b	Ω_b^*	B	B^*
Experimental value	5795	5935	5949	6046	?	5279	5325
ChQM	5794	5880	5884	6008	6011	5351	5358

one. The wave function of the baryon-meson system is of the form

$$\Psi = \mathcal{A}[\hat{\psi}_A(\xi_1, \xi_2)\hat{\psi}_B(\xi_3)\chi_L(\mathbf{R}_{AB})], \quad (11)$$

where ξ_1 and ξ_2 are the internal coordinates for the baryon cluster A and ξ_3 is the internal coordinate for the meson cluster B. $\mathbf{R}_{AB} = \mathbf{R}_A - \mathbf{R}_B$ is the relative coordinate between the two clusters A and B. The $\hat{\psi}_A$ and $\hat{\psi}_B$ are the antisymmetrized internal cluster wave functions of the baryon A and meson B, and $\chi_L(\mathbf{R}_{AB})$ is the relative motion wave function between two clusters. Here, we study only the molecular baryon-meson systems of the S wave, so the orbital angular momentum $L = 0$. The symbol $\mathcal{A}$ is the antisymmetrization operator defined as

$$\mathcal{A} = 1 - P_{14} - P_{24} - P_{34}, \quad (12)$$

where 1, 2, and 3 stand for the quarks in the baryon cluster and 4 stands for the quark in the meson cluster. We expand this relative motion wave function by Gaussian bases,

$$\begin{aligned} \chi_L(\mathbf{R}_{AB}) = & \frac{1}{\sqrt{4\pi}} \left(\frac{6}{5\pi b^2} \right) \sum_{i=1}^n C_i \\ & \times \int \exp \left[-\frac{3}{5b^2} (\mathbf{R}_{AB} - \mathbf{S}_i)^2 \right] Y_{LM}(\hat{\mathbf{S}}_i) d\hat{\mathbf{S}}_i, \end{aligned} \quad (13)$$

where $\mathbf{S}_i$ is called the generate coordinate and n is the number of the Gaussian bases, which is determined by the stability of the results. By doing this, the integro-differential equation of the RGM can be reduced to the algebraic equation, generalized eigenequation. Then, we can obtain the energy of the system by solving this generalized eigenequation. The details of solving the RGM equation can be found in Ref. [24]. In our calculation, the distance between the baryon-meson clusters is identified by the stability of the results. The results are stable when the distance between the baryon-meson clusters is larger than 6 fm. To keep the dimensions of the matrix manageably small, the baryon-meson separation is taken to be less than 6 fm. Note that the wave function used here is different from that in Ref. [23]. The purpose of this work is to study the mass and decay width of the molecular pentaquarks and analyze if there are some reported Ω_c^0 states which can be explained as pentaquarks. So, the internal motions of clusters are frozen, and the relative motion wave function between two clusters is expanded by Gaussian bases. This indicates that we consider only one relative motion of this pentaquark system. The binding energy and scattering phase shifts can be obtained conveniently by doing this. In Ref. [23], the authors aimed to study the structure of the Ω_c^0 states, so they used four relative motions for the five-body system. The Gaussian

expansion method [34] is employed to find the orbital wave functions, and every relative motion wave function is expended by a Gaussian with various sizes. Through this approach, the distances between quark pairs can be obtained, which suggests the structure of the pentaquarks.

III. RESULTS AND DISCUSSIONS

In this work, we investigate the S -wave Ω_c^0 states with $I = 0$, $J^P = \frac{1}{2}^-$, $\frac{3}{2}^-$, and $\frac{5}{2}^-$. Three structures are considered here, and they are structures 1) $uss - c\bar{u}$, 2) $usc - s\bar{u}$, and 3) $ssc - u\bar{u}$. All the channels involved are listed in Table III.

A. Bound-state calculation

For the first step, we do a dynamic calculation to check whether or not there is any bound state. The single-channel calculation shows that the energy of each channel is located above the threshold of the corresponding channel, which means that there is no bound state in this case. By coupling all the channels with different structures, there exist some bound states. The binding energies and the masses of the bound states as well as the percentages of each channel in the eigenstates are listed in Table IV. Before discussing the features of the states, we should mention how we obtain the mass of these states. First, we define the binding energy $B = M^{\text{the}} - M_B^{\text{the}} - M_M^{\text{the}}$, where M^{the} , M_B^{the} , and M_M^{the} stand for the theoretical mass of the molecular state, a baryon, and a meson, respectively. To minimize the theoretical errors and to compare calculated results to the experimental data, then we shift the mass of a molecular state to $M = M_B^{\text{exp}} + M_M^{\text{exp}} + B$, where the experimental values of a baryon and a meson are used. Taking the state $J^P = \frac{1}{2}^- \Xi D$ as an example, the calculated mass of this state is 3169 MeV; then, the binding energy B is obtained by subtracting the theoretical masses of Ξ and D , $3169 - 1225 - 1980 = -36$ (MeV). Adding the experimental masses of the hadrons, the mass of this state $M = 1318 + 1864 + (-36) = 3146$ (MeV) is reached.

For the $J^P = \frac{1}{2}^-$ system, the channel coupling of each structure cannot make any state bound. Then, we do a

TABLE III. The channels calculated in this work.

J^P	Structure	Channels
$\frac{1}{2}^-$	1. $uss - c\bar{u}$	$\Xi D, \Xi D^*, \Xi^* D^*$
	2. $usc - s\bar{u}$	$\Xi_c \bar{K}, \Xi'_c \bar{K}, \Xi_c \bar{K}^*, \Xi'_c \bar{K}^*, \Xi_c^* \bar{K}^*$
	3. $ssc - u\bar{u}$	$\Omega_c \eta, \Omega_c \omega, \Omega_c^* \omega$
$\frac{3}{2}^-$	1. $uss - c\bar{u}$	$\Xi D^*, \Xi^* D, \Xi^* D^*$
	2. $usc - s\bar{u}$	$\Xi_c \bar{K}^*, \Xi'_c \bar{K}^*, \Xi_c^* \bar{K}, \Xi_c^* \bar{K}^*$
	3. $ssc - u\bar{u}$	$\Omega_c \omega, \Omega_c^* \eta, \Omega_c^* \omega$
$\frac{5}{2}^-$	1. $uss - c\bar{u}$	$\Xi^* D^*$
	2. $usc - s\bar{u}$	$\Xi_c^* \bar{K}^*$
	3. $ssc - u\bar{u}$	$\Omega_c^* \omega$

TABLE IV. The binding energy and masses (in MeV) of the molecular pentaquarks with channel coupling and the percentages of each channel in the eigenstates.

$J^P = \frac{1}{2}^-$		$J^P = \frac{3}{2}^-$		$J^P = \frac{5}{2}^-$	
E_B	-36	E_B	-1	E_B	-7
M_{cc}	3146	M_{cc}	3324	M_{cc}	3533
ΞD	58.1%	ΞD^*	63.6%	$\Xi^* D^*$	50.3%
ΞD^*	3.3%	$\Xi^* D$	3.6%	$\Xi_c^* \bar{K}^*$	13.4%
$\Xi^* D^*$	1.9%	$\Xi^* D^*$	0.2%	$\Omega_c^* \omega$	36.3%
$\Omega_c \eta$	31.1%	$\Omega_c \omega$	1.2%		
$\Omega_c \omega$	1.9%	$\Omega_c^* \eta$	31.3%		
$\Omega_c^* \omega$	3.7%	$\Omega_c^* \omega$	0.1%		

channel coupling among different structures. We find that there is no bound state by coupling channels of structures 1 and 2, or structures 2 and 3. However, by coupling the channels of structures 1 and 3, we obtain a stable state, the mass of which is 3146 MeV, -36 MeV lower than the threshold of ΞD , and the main component of this state is ΞD , as shown in Table IV. While coupling these channels to the channels of structure 2, in which there are open channels $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$, the thresholds of which are lower than the mass of ΞD , the lowest eigenenergy is very close to but still higher than the threshold of $\Xi_c \bar{K}$, the channel with the lowest energy among three structures. So, there is no bound state by all channels coupling. However, we also obtain a quasistable state, the mass of which is smaller than the threshold of ΞD , but it is higher than the threshold of $\Xi_c \bar{K}$, and it fluctuates around 3146 MeV with about 1 MeV with the variation of the baryon-meson separation. To confirm whether or not the state ΞD can survive as a resonance state after coupling to the open channels, the study of the scattering process of the open channels is needed, which is discussed in Sec. III. B.

For the $J^P = \frac{3}{2}^-$ system, we find a bound state ΞD^* with a binding energy of only -1 MeV by coupling channels of structures 1 and 3, as shown in Table IV. There is neither any bound state nor any quasistable state by the full channels coupling. However, we still need to calculate the scattering process of the open channel $\Xi_c^* \bar{K}$ to check if the state ΞD^* is a resonance or not. The result is also shown in Sec. III. B.

For the $J^P = \frac{5}{2}^-$ system, it includes three channels, and each structure contributes one channel. Although it is not bound for each structure, there exists a bound state by three channels coupling. The binding energy and the mass of this system, as well as the percentages of each channel in the eigenstate, are shown in Table IV, from which we can see that the mass of the $J^P = \frac{5}{2}^-$ system is 3533 MeV. Moreover, this state can also decay to some open channels, for example, ΞD and $\Xi_c \bar{K}$, but they must be in the D wave, because of the total angular momentum $J = \frac{5}{2}$. The decay can only proceed through the tensor force, which is always

very weak in our quark model calculation [26]. So, we can estimate the effect of this kind of coupling is small here. We will do this S - and D -wave channel coupling in the future.

B. Resonance states and decay widths

Resonances are unstable particles usually observed as bell-shaped structures in scattering cross sections of their open channels. For a simple narrow resonance, its fundamental properties correspond to the visible cross section features: mass M is at the peak position, and decay width Γ is the half-width of the bell shape. To find the resonance mass and decay width of the quasistable states discussed in Sec. III. A, we calculate the cross section of the corresponding open channels in this section. The cross section can be obtained from the scattering phase shifts by the formula

$$\sigma = \frac{4\pi}{k^2} \cdot (2l + 1) \cdot \sin^2 \delta, \quad (14)$$

where $k = \sqrt{2\mu E_{cm}}/\hbar$; μ is the reduced mass of two hadrons of the open channel; E_{cm} is the incident energy; and δ is the scattering phase shift of the open channel, which can be obtained by the well-developed RGM [24].

For the $J^P = \frac{1}{2}^-$ system, we have found that ΞD is a quasistable state by coupling with the open channels $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$. Here, to confirm whether or not the state ΞD is a resonance state, we study the scattering process of two open channels $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$. The phase shifts of the S -wave $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$ are shown in Figs. 1(a) and 1(b), respectively, where $cc1$ is the phase shift of the open channel by coupling all channels of structures 1 and 3; $cc2$ is the phase shift of the open channel by coupling all channels of structures 1 and 3, except for the ΞD channel; $cc3$ is the phase shift of the open

 FIG. 1. The phase shifts of the open channels $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$ for the $J^P = \frac{1}{2}^-$ system.

FIG. 2. The cross section of the open channels $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$ for the $J^P = \frac{1}{2}^-$ system.

channel by subtracting the phase shift of $cc2$ from that of $cc1$. Coupling to the open channel $\Xi_c \bar{K}$ causes the bound state ΞD to change into an elastic resonance where the phase shift of $\Xi_c \bar{K}$ (see $cc1$) shows a sharp increase around the resonance mass. If we take out of this resonance state ΞD , there is no rapid increase in the phase shifts (see $cc2$), which means that there is no other resonance state here and the phase shift of $cc2$ is the one with the effect of nonresonance states. So, if we subtract the phase shift of $cc2$ from that of $cc1$, we can obtain the phase shift with the effect of the resonance state, shown as $cc3$, where the phase shift rises through $\pi/2$ at a resonance mass. Then, substituting the phase shift of $cc3$ into Eq. (14), the cross section of open channel $\Xi_c \bar{K}$ is achieved, shown in Fig. 2, from which we can obtain the resonance mass at the peak position and decay width at the half-width of the bell shape. We find the resonance mass of the ΞD state is 3144.6 MeV, and the decay width of the ΞD is only 0.2 MeV. By coupling to another open channel $\Xi_c' \bar{K}$, similar results are obtained. The ΞD bound state becomes a resonance state with the same resonance mass and decay width. Therefore, we can obtain a resonance state ΞD with $J^P = \frac{1}{2}^-$ in the decay channel $\Xi_c \bar{K}$ or $\Xi_c' \bar{K}$, with the resonance mass 3144.6 MeV and decay width 0.2 MeV, which is consistent with the newly reported $\Omega_c(3119)^0$, the decay width of which is $1.1 \pm 0.8 \pm 0.4$ MeV. What is more, this $\Omega_c(3119)^0$ state was observed in both $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$ channels in the LHCb experiment [9]. So, in our quark model calculation, we can explain the $\Omega_c(3119)^0$ as a resonance state ΞD with $J^P = \frac{1}{2}^-$.

For the $J^P = \frac{3}{2}^-$ system, we study the scattering process of the open channel S -wave $\Xi_c^* \bar{K}$, and we do not find any resonance state. This is reasonable because the binding energy of the ΞD^* is only -1 MeV, which is too small. The channel coupling to the $\Xi_c^* \bar{K}$ pushes the energy of ΞD^*

FIG. 3. The phase shifts of the open channels $\Xi_b \bar{K}$ and $\Xi_b' \bar{K}$ for the $J^P = \frac{1}{2}^-$ system.

FIG. 4. The cross section of the open channels $\Xi_b \bar{K}$ and $\Xi_b' \bar{K}$ for the $J^P = \frac{1}{2}^-$ system.

above its threshold. So, there is no resonance state with $J^P = \frac{3}{2}^-$ in our calculation.

In addition, we also extend the study to the Ω_b^0 system. Because of the heavy flavor symmetry, the results are similar to that of the Ω_c^0 system. We obtain a resonance state ΞB with $J^P = \frac{1}{2}^-$ in the decay channels $\Xi_b \bar{K}$ and $\Xi_b' \bar{K}$, with the same resonance mass of 6560 MeV and decay width 0.2 MeV (see Figs. 3 and 4). Besides, there is a bound state $\Xi^* B^*$ with $J^P = \frac{5}{2}^-$, and the energy of it is 6856 MeV.

IV. SUMMARY

In summary, we investigate the excited Ω_c^0 states as the S -wave molecular pentaquarks with $I = 0$, $J^P = \frac{1}{2}^-, \frac{3}{2}^-$, and $\frac{5}{2}^-$ by solving the RGM equation in the framework of the chiral quark model. In this work, we study not only the

energies of the Ω_c^0 states but also the decay widths of them. Our results show that the $\Omega_c(3119)^0$ can be explained as an S -wave resonance state ΞD with $J^P = \frac{1}{2}^-$, and the decay channels are the S -wave $\Xi_c \bar{K}$ and $\Xi_c' \bar{K}$. Other newly reported Ω_c^0 states cannot be obtained in our present calculation. They maybe the conventional charmed baryons with P -wave or radial excitation states. It is also possible that these states are mixing states with q^3 and $q^4 \bar{q}$. The unquenched quark model, which takes into account the higher Fock components, is feasible to do this mixing. Therefore, some of those new Ω_c^0 states could be studied in the future in the unquenched quark model for heavy baryons as has already been done in the unquenched quark model for light baryons [35–40] and light mesons [41] and heavy ones [42–51]. We also obtain another Ω_c^0 state with much higher mass, which is 3533 MeV with $J^P = \frac{5}{2}^-$. Besides, the calculation is extended to the Ω_b^0 states, and results similar to that of Ω_c^0 are obtained.

Our calculation also shows that the coupling of the structure $uss - c\bar{u}$ and $ssc - u\bar{u}$ is important to make the ΞD bound, and the coupling with structure $usc - s\bar{u}$ is very weak, which can be used to explain why the reported $\Omega_c(3119)^0$ has a narrow decay width. It also gives us some information that the decay width of these Ω_c^0 states is related to the structure of these states. Reference [23] studied the structure of the observed Ω_c^0 , and suggested that $\Xi \bar{D}$, $\Xi_c \bar{K}$, and $\Xi_c^* \bar{K}$ are possible the candidates of

some newly announced Ω_c^0 states with the energy below 3.2 GeV. The 100% color-singlet component of $\Xi \bar{D}$ may infer the molecular structure of the state. In this work, we only study the Ω_c^0 states as meson-baryon molecular states. Comparing these two works indicates that the channel coupling with more structures may lower the energy of the pentaquark system. Meanwhile, Ref. [23] also studied the $3q$ configurations and found two first radial excitation states ($2S$), the masses of which are close to 3090 and 3119 MeV, respectively. However, the authors did not study the decay width of the Ω_c^0 states. Actually, it is still difficult for many theoretical works to describe the five observed Ω_c^0 states well. The structures and decay width may be the significant nature of the multiquark states. So, we should develop our model and calculation method to investigate the nature of the multiquark states and give more information on the experimental search for multiquark states.

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