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New Ω_c^0 baryons discovered by LHCb as the members of 1*P* and 2*S* states

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Inspired by the recently observed Ω_c^0 states at LHCb, we decode their properties by performing an analysis of their mass spectrum and decay behavior. Our studies show that the five narrow states, i.e., $\Omega_c(3000)^0$, $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$, could be grouped into the 1*P* states with negative parity. Among them, the $\Omega_c(3000)^0$ and $\Omega_c(3090)^0$ states could be the $J^P = 1/2^-$ candidates, while $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ are suggested as the $J^P = 3/2^-$ states. $\Omega_c(3066)^0$ could be regarded as a $J^P = 5/2^-$ state. Since the spin-parity, electromagnetic transitions, and possible hadronic decay channels $\Omega_c^{(*)}\pi$ have not been measured yet, other explanations are also probable for these narrow Ω_c^0 states. Additionally, we discuss the possibility of the broad structure $\Omega_c(3188)^0$ as a 2*S* state with $J^P = 1/2^+$ or $J^P = 3/2^+$. In our scheme, $\Omega_c(3119)^0$ cannot be a 2*S* candidate.

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I. INTRODUCTION

Establishing the higher radial and orbital excited states of heavy baryons has been an interesting and important research issue of hadron spectroscopy in recent years [1], by which we can gain more information about the nonperturbative behavior of quantum chromodynamics (QCD). With the efforts of experiments, significant progress has been made in searching for highly excited charmed baryons in the last few years. For example, the $\Xi_c(2790)$, $\Xi_c(2815)$, and $\Xi_c(2980)$ have been measured by Belle with greater precision [2]. A new decay mode, $D^+\Lambda$, was found for the $\Xi_c(3055)^+$ and $\Xi_c(3080)^+$ [3]. Recently, LHCb remeasured the resonance parameters of the $\Lambda_c(2880)^+$ and $\Lambda_c(2940)^+$ [4], and confirmed the previously measured masses, decay widths, and the spin of $\Lambda_c(2880)^+$. In addition, a constraint on the spin-parity of the $\Lambda_c(2940)^+$ was given for the first time. More importantly, LHCb also observed a new broad state, $\Lambda_c(2860)^+$, with $J^P = 3/2^+$ [4]. These new measurements have greatly enriched the information about charmed baryons with the cqq and csq ($q = u$ or d quark) configurations. It is obvious that the study of charmed baryons is still ongoing.

Very recently, the LHCb Collaboration again brought us a surprise due to the observation of five new narrow Ω_c^0 states plus a broad structure in the $\Xi_c^+ K^-$ invariant mass spectrum [5]. Here, the five narrow resonances are the $\Omega_c(3000)^0$, $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ states. Since this broad structure around 3188 MeV may be from a single resonance, a superposition of several states, or other mechanism [5], in this work we tentatively name this broad structure as the $\Omega_c(3188)^0$. The concrete experimental results of LHCb plus two established Ω_c^0 baryons, $\Omega_c(2700)$ and $\Omega_c(2770)$ [6], are collected in Table I for the reader's convenience. These newly detected Ω_c^0 states at LHCb not only make the Ω_c^0 charmed baryon family become abundant, but also let us face how to categorize them into the Ω_c^0 family.

These newly observed Ω_c^0 states have greatly aroused the theorists' interest [7–20]. In the following, we briefly review the research status of these new Ω_c^0 states. With the method of QCD sum rules the masses of 2*S* Ω_c^0 states were calculated [7], where the assignment of $\Omega_c(3066)^0$ and $\Omega_c(3119)^0$ as the 2*S* states was proposed. The strong decays of *P*-wave charmed baryons have been systematically investigated using the light-cone QCD sum rules [8], where the decay behaviors of the $\Omega_c(3000)^0$, $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ were discussed. Based on the analysis from the Regge trajectories and the mass calculation via the heavy quark-light diquark model, the $\Omega_c(3090)^0$ and $\Omega_c(3119)^0$ were suggested as the 2*S* states with $J^P = 1/2^+$ and $J^P = 3/2^+$, respectively [9]. Within a constituent quark model, the analysis of strong behavior also favors the 2*S* state assignment to the $\Omega_c(3119)^0$ [10]. This assignment to the $\Omega_c(3119)^0$ mentioned above is partly supported by the results in Refs. [21,22], where the masses of 2*S* Ω_c^0 states were predicted to be around 3100 MeV. Different explanations have been proposed in Refs. [11–14], where the $\Omega_c(3000)^0$, $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ were suggested as the 1*P* excited states with negative parity. In addition, some authors suggested that several newly observed narrow Ω_c^0 states may be the

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TABLE I. The experimental results, including masses, widths, and decay modes, for the observed Ω_c^0 states.

State [5,6]	Decay mode	Mass (MeV)	Width (MeV)
$\Omega_c(2700)^0$	week	2695.2 ± 1.7	
$\Omega_c(2770)^0$	$\Omega_c^0\gamma$	2765.9 ± 2.0	
$\Omega_c(3000)^0$	$\Xi_c K$	$3000.4 \pm 0.2 \pm 0.1^{+0.3}_{-0.5}$	$4.5 \pm 0.6 \pm 0.3$
$\Omega_c(3050)^0$	$\Xi_c K$	$3050.2 \pm 0.1 \pm 0.1^{+0.3}_{-0.5}$	$0.8 \pm 0.2 \pm 0.1$
$\Omega_c(3066)^0$	$\Xi_c K$, $\Xi_c' K$	$3065.6 \pm 0.1 \pm 0.3^{+0.3}_{-0.5}$	$3.5 \pm 0.4 \pm 0.2$
$\Omega_c(3090)^0$	$\Xi_c K$, $\Xi_c' K$	$3090.2 \pm 0.3 \pm 0.5^{+0.3}_{-0.5}$	$8.7 \pm 1.0 \pm 0.8$
$\Omega_c(3119)^0$	$\Xi_c K$, $\Xi_c' K$	$3119.1 \pm 0.3 \pm 0.9^{+0.3}_{-0.5}$	$1.1 \pm 0.8 \pm 0.4$
$\Omega_c(3188)^0$	$\Xi_c K$	$3188 \pm 5 \pm 13$	$60 \pm 15 \pm 11$

charmed exotic systems [16–20]. Obviously, different conclusions were made by different theoretical methods (see Table II in Ref. [10] for a concise review). When facing such a messy situation, more efforts are needed regarding these newly observed Ω_c^0 baryons.

In this work, we perform a systematic analysis of these new Ω_c^0 states by analyzing the mass spectra and calculating the strong decays. First, we adopt a simple quark potential model to calculate the mass spectrum of low excited Ω_c^0 states, where the heavy quark–light diquark picture is employed. Our study indicates that P-wave Ω_c baryons have masses around 3.05 GeV, which overlaps with experimental observation. Thus, identifying five observed Ω_c states as P-wave charmed baryons becomes possible. To further determine their spin-parity quantum numbers, we study their two-body Okubo-Zweig-Iizuka (OZI)-allowed strong decays using the quark pair creation (QPC) model, which may provide more information about their inner structure.

The paper is organized as follows. We analyze the mass spectrum of low-lying Ω_c^0 excitations in the next section. In Sec. III, we present an introduction to the QPC model for calculating the Ω_c^0 baryon decays. In Sec. IV, we calculate the strong decays of the newly observed Ω_c^0 baryon states and discuss our results. The paper ends with a summary and outlook in Sec. IV.

II. AN ESTIMATE OF MASSES OF LOW-LYING Ω_c^0 STATES

The dynamics of a Ω_c^0 state could be separated into two parts, i.e., the degree of freedom between two s quarks is usually called the "ρ mode," while the degree of freedom between the center of mass of two s quarks and the c quark corresponds to the "λ mode." If we use an ordinary confining potential (like the linear or oscillator harmonic form) to depict the interaction, the mass of the ρ mode excitations is higher than that of the λ mode excitations [23,24]. So experiments may first detect the λ mode excited Ω_c^0 states. Just considering this point, we first discuss whether these newly observed Ω_c^0 states can be regarded as the λ mode excitations.[1]

Since the excitation between two s quarks is not considered here, we may treat two s quarks in these λ-mode excited Ω_c^0 baryons as a special part with the antitriplet color structure and a peculiar size, which is called a light-quark cluster or a light diquark. Under this heavy quark–light diquark picture, Isgur discussed the similarity between the dynamics of heavy baryons and heavy-light mesons [25]. Later, the heavy baryon masses were systematically calculated in such a diquark picture [21,26]. In our recent work [27], the properties of low-excited charmed and charmed-strange baryons were described well by this diquark picture. We need to emphasize that the heavy quark–light diquark picture for the charmed baryon system is partially supported by the recent Belle measurement of the production cross sections of charmed baryons [28], where a factor of 3 difference for Λ_c^+ states over Σ_c states was observed. This observation suggests a diquark structure in the ground state and low-lying excited Λ_c^+ baryons [28].

In studying the dynamics of these low-lying Ω_c^0 baryons, as described above, the light quark cluster (denoted as $\{ss\}$) could be an effective degree of freedom. Thus, a λ-mode excited Ω_c^0 state may be simplified as a quasi-two-body system, which is similar to a meson system composed of a quark and an antiquark. Here, the Cornell potential [29] is applied to phenomenologically describe the confining part of the interaction between a charm quark and an $\{ss\}$ cluster, i.e.,

$$H^{\rm conf} = -\frac{4}{3}\frac{\alpha_s}{r} + br + C. \tag{1}$$

The parameters α_s, b, and C stand for the strength of the color Coulomb potential, the strength of linear confinement, and a mass-renormalized constant, respectively. Furthermore, we expect that the following spin-dependent interaction terms which have been applied in studying the mass spectrum of different flavor mesons [30] could also be suitable for the λ-mode excited Ω_c^0 baryons. The first one is the magnetic dipole–magnetic dipole color hyperfine interaction

$$H_{\rm hyp} = \frac{4}{3}\frac{\alpha_s}{m_l m_c}\left(\frac{8\pi}{3}\delta^3(\vec{r})\vec{s}_l\cdot\vec{s}_c + \frac{1}{r^3}\hat{S}_{lc}\right), \tag{2}$$

where the tensor operator is defined as

$$\hat{S}_{lc} = \frac{3\vec{s}_l\cdot\vec{r}\,\vec{s}_c\cdot\vec{r}}{r^2} - \vec{s}_l\cdot\vec{s}_c.$$

[1] In principle, a heavy-light baryon state should contain both a ρ mode and a λ mode. However, the investigation by Yoshida *et al.*. indicated that the ρ and λ modes are well separated for the charmed and bottomed baryons [23].

Here, $\vec{s}_l$ and $\vec{s}_c$ represent the spins of the $\{ss\}$ diquark and c quark, respectively. The second one corresponds to the spin-orbit interactions which contains two contributions. The first piece is the color magnetic interaction due to one-gluon exchange,

$$H_{\rm so}^{\rm cm} = \frac{4\alpha_s}{3r^3}\left(\frac{1}{m_l}+\frac{1}{m_c}\right)\left(\frac{\vec{s}_l}{m_l}+\frac{\vec{s}_c}{m_c}\right)\cdot\vec{L}. \quad (3)$$

The second piece is the Thomas-precession term,

$$H_{\rm so}^{\rm tp} = -\frac{1}{2r}\frac{\partial H^{\rm conf}}{\partial r}\left(\frac{\vec{s}_l}{m_l^2}+\frac{\vec{s}_c}{m_c^2}\right)\cdot\vec{L}. \quad (4)$$

With the expression for $H^{\rm conf}$ in Eq. (1), the total spin-orbit interaction is given as

$$\begin{aligned} H_{\rm so} = &\left[\left(\frac{2\alpha}{3r^3}-\frac{b}{2r}\right)\frac{1}{m_l^2}+\frac{4\alpha}{3r^3}\frac{1}{m_l m_c}\right]\vec{s}_l\cdot\vec{L} \\ &+\left[\left(\frac{2\alpha}{3r^3}-\frac{b}{2r}\right)\frac{1}{m_c^2}+\frac{4\alpha}{3r^3}\frac{1}{m_l m_c}\right]\vec{s}_c\cdot\vec{L}, \end{aligned} \quad (5)$$

Due to the Pauli exclusion principle, the total wave function of the diquark $\{ss\}$ should be antisymmetric in the exchange of two strange quarks. Since the spatial and flavor parts of this light-quark cluster are always symmetric and the color part is antisymmetric, the spin wave function should be symmetric, i.e., $s_l = 1$ for the Ω_c^0 state. In addition, L denotes the orbital quantum number between the light $\{ss\}$ cluster and the charm quark. By solving the Schrödinger equation, the masses of low-lying Ω_c^0 states can be obtained directly. In our calculations, all spin-dependent interactions are treated as the leading-order perturbations for the orbital Ω_c^0 excitations.

In our scheme, the mass matrix is calculated in the jj coupling scheme. To this end, we adopt the basis $|s_l, L, j_l, s_c, J\rangle$, where $\vec{s}_l+\vec{L}=\vec{j}_l$ and $\vec{j}_l+\vec{s}_c=\vec{J}$. The five eigenvectors for the $1P$ λ-mode excited Ω_c^0 states include $|1,1,0,1/2,1/2\rangle \equiv |0,1/2^-\rangle$, $|1,1,1,1/2,1/2\rangle \equiv |1,1/2^-\rangle$, $|1,1,1,1/2,3/2\rangle \equiv |1,3/2^-\rangle$, $|1,1,2,1/2,3/2\rangle \equiv |2,3/2^-\rangle$, and $|1,1,2,1/2,5/2\rangle \equiv |2,5/2^-\rangle$, where the notation $|s_l, L, j_l, s_Q, J\rangle$ is abbreviated as $|j_l, J^P\rangle$ and the superscript P denotes parity.

Due to the spin-orbit interaction $\vec{s}_c\cdot\vec{L}$ and the tensor interaction, two physical $1P$ states with $J^P = 1/2^-$ should be mixtures of $|0,1/2^-\rangle$ and $|1,1/2^-\rangle$, which satisfy

$$\begin{pmatrix} |1P,1/2^-\rangle_L \\ |1P,1/2^-\rangle_H \end{pmatrix} = \begin{pmatrix} \cos\theta_1 & -\sin\theta_1 \\ \sin\theta_1 & \cos\theta_1 \end{pmatrix}\begin{pmatrix} |0,1/2^-\rangle \\ |1,1/2^-\rangle \end{pmatrix}. \quad (6)$$

For the case of $J^P = 3/2^-$ states, there exists

$$\begin{pmatrix} |1P,3/2^-\rangle_H \\ |1P,3/2^-\rangle_L \end{pmatrix} = \begin{pmatrix} \cos\theta_2 & -\sin\theta_2 \\ \sin\theta_2 & \cos\theta_2 \end{pmatrix}\begin{pmatrix} |1,3/2^-\rangle \\ |2,3/2^-\rangle \end{pmatrix}. \quad (7)$$

The physical states with the same J^P are distinguished by their different masses and widths. Here, the states with the lower and higher masses are denoted by the subscripts "L" and "H," respectively. Since the contribution from the contact hyperfine interaction [the first term of Eq. (2)] is small for orbitally excited states, the spin-dependent interactions for the P-wave Ω_c^0 baryons can be further simplified as

$$H_S = V_l\vec{s}_l\cdot\vec{L} + V_c\vec{s}_c\cdot\vec{L} + V_t\hat{S}_{lc}. \quad (8)$$

By comparing Eqs. (2) and (5), we may define the expressions for V_l, V_c, and V_t, i.e.,

$$\begin{aligned} V_l &= \left(\frac{2\alpha}{3r^3}-\frac{b}{2r}\right)\frac{1}{m_l^2}+\frac{4\alpha}{3r^3}\frac{1}{m_l m_c}, \\ V_c &= \left(\frac{2\alpha}{3r^3}-\frac{b}{2r}\right)\frac{1}{m_c^2}+\frac{4\alpha}{3r^3}\frac{1}{m_l m_c}, \\ V_t &= \frac{4}{3}\frac{\alpha_s}{m_l m_c}\frac{1}{r^3}. \end{aligned}$$

For P-wave states with $J^P = 1/2^-$, the mass matrix is

$$\langle\Phi_{1/2}|H_S|\Phi_{1/2}\rangle = \begin{pmatrix} -2V_l - \frac{4}{3}V_t & -\frac{3V_c+V_t}{3\sqrt{2}} \\ -\frac{3V_c+V_t}{3\sqrt{2}} & -V_l - \frac{1}{2}V_c + \frac{1}{3}V_t \end{pmatrix}.$$

Similarly, for two states with $J^P = 3/2^-$, we have

$$\langle\Phi_{3/2}|H_S|\Phi_{3/2}\rangle = \begin{pmatrix} \frac{1}{4}V_c - V_l + \frac{5}{6}V_t & \frac{4V_t-15V_c}{12\sqrt{5}} \\ \frac{4V_t-15V_c}{12\sqrt{5}} & V_l - \frac{3}{4}V_c - \frac{1}{30}V_t \end{pmatrix}.$$

For the $J^P = 5/2^-$ state, we get

$$\langle 2,5/2^-|H_S|2,5/2^-\rangle = V_l + \frac{1}{2}V_c - \frac{1}{5}V_t.$$

The notations $|\Phi_{1/2}\rangle$ and $|\Phi_{3/2}\rangle$ in the above expressions have the definition

$$|\Phi_J\rangle = \begin{pmatrix} |j_l = J-1/2, J^P\rangle \\ |j_l = J+1/2, J^P\rangle \end{pmatrix}.$$

There are five parameters, m_c, m_l, b, α_s, σ, and C, in this nonrelativistic quark potential model. The c quark mass is taken to be $m_c = 1.68$ GeV from our previous work [27]. Based on the SU(3) flavor symmetry, there is a similarity between the dynamics for Σ_c and Ω_c baryon families. The averaged mass of the ground Σ_c states [$\Sigma_c(2455)$ and $\Sigma_c(2520)$] is about 2496 MeV, while the averaged mass of the ground Ω_c states [$\Omega_c(2700)$ and $\Omega_c(2770)$] is about 2743 MeV. Thus, the splitting of the ground Σ_c and Ω_c states is given by

$$\bar{M}_{[\Sigma_C(2455),\Sigma_C(2520)]} - \bar{M}_{[\Omega_C(2700),\Omega_C(2770)]} \approx 247 \text{ MeV}, \qquad (9)$$

which could be regarded as the mass difference between the light-quark clusters of $\{uu\}$ and $\{ss\}$ [31]. Then, we evaluate the mass of the axial-vector light-quark cluster m_l as 0.91 GeV since the mass of the $\{uu\}$ cluster in Σ_c baryons has been fixed as 660 MeV by a relativistic flux tube model in our previous work [26]. The strength of the color Coulomb potential α_s and the strength of linear confinement b are taken as 0.34 and 0.120 GeV^{-1}, respectively, which can fairly reproduce the averaged masses of the observed $1S$, $1P$, and $2S$ Ω_c^0 candidates. To reproduce the mass splitting of the $\Omega_c(2700)^0$ and $\Omega_c(2770)^0$, the parameter σ is fixed as 1.00. Finally, the constant C is determined as 0.16 GeV.

With the parameters above, the masses of the two ground ($1S$) Ω_c states are predicted as 2698 and 2765 MeV, which are in good agreement with the experimental results (see Table I). For a clear comparison, we collect the predicted $1P$ and $2S$ Ω_c^0 masses and the newly observed states by LHCb together in Fig. 1. According to the predicted masses, we find the following:

(1) Among the five $1P$ states, a $J^P = 1/2^-$ state should have the lowest mass. This conclusion is supported by most theoretical works [11–14,23,32,33]. Then, we may conjecture that $\Omega_c(3000)^0$ is a good candidate for the $|1P, 1/2^-\rangle_L$ state if we try to assign these five narrow states to the $1P$ Ω_c^0 family.

(2) The masses of the $|1P, 1/2^-\rangle_H$ state, two $3/2^-$ states, and one $5/2^-$ state are predicted in the range 3068–3092 MeV. Considering the intrinsic uncertainties of the quark potential models, we cannot determine the quantum numbers for the $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ states using only the analysis of the mass spectrum mentioned above.

(3) The broad structure—the $\Omega_c(3188)^0$—may be regarded as a $2S$ state with $J^P = 1/2^+$ or $3/2^+$, or their overlapping structure.

With the mass matrices above, the mixing angles in Eqs. (6) and (7) are obtained as $\theta_1 = 158°$ and $\theta_2 = 159°$,

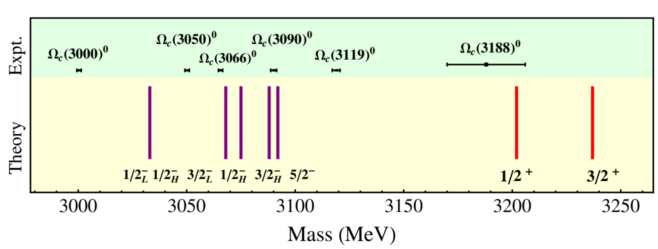

FIG. 1. A comparison of the predicted $1P$ and $2S$ Ω_c^0 masses with the newly observed states by LHCb.

respectively. However, the spin-dependent interactions for the excited Ω_c^0 states have never been understood well in the three-body picture. We should treat these obtained mixing angles with some care. Although the spin-parity quantum numbers of $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ cannot be fixed under the simplified quark potential model in the diquark picture, we may conclude that the $1P$ states of the Ω_c baryon family have a mass around 3.05 GeV which overlaps with the masses of these observed states. Thus, assigning $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ states as $1P$ states of the Ω_c family is still possible. Our tentative conclusion is also supported by the result given recently using lattice QCD [12].

In the following section, we will employ the QPC model to give further constraints for the assignments of these narrow Ω_c^0 states by analyzing their strong decays.

III. THE Ω_c^0 BARYON DECAYS IN THE QPC MODEL

The idea of the QPC strong decay model was introduced by Micu [34] and Carlitz and Kislinger [35], and later formulated by the Orsay group [36,37]. For an OZI-allowed decay process of a hadron system, the QPC model suggests that a quark-antiquark pair is created from the vacuum and then regroups into two outgoing hadrons by a quark rearrangement process. Thus, the $q\bar{q}$ pair will carry the quantum number of 0^{++}, suggesting that they are in a $^3P_0$ state. So the QPC model is also named "the $^3P_0$ model." This model has been successfully applied to study different kinds of hadron systems for their OZI-allowed strong decays. Here we just quote some works that were devoted to the strong decay behaviors of excited heavy baryons. The strong decays of the S-wave, P-wave, D-wave, and radially excited charmed baryons were studied using the QPC model [38]. The QPC model was used to study the decay processes of ground and excited bottom baryons [39]. Recently, the decays of P-wave excitations of charmed strange baryons were studied systematically using the QPC model [40].

According to the excited energy, there are two kinds of decay processes for an excited charmed baryon state (see Fig. 2). When the mass of a Ω_c^0 excitation is higher than about 2.97 GeV (the threshold of the $\Xi_c K$ channel), only the left process (labeled by [a]) is allowed. When the mass is higher than the threshold of the $D\Xi$ channel, two kinds of decay processes (depicted in Fig. 2) are possible. We now take the left process, $\Omega_c(A) \to \Xi_c(B) + K(C)$, as an example to show how to obtain the partial wave amplitudes. To describe the decay process, the transition operator $\hat{\mathcal{T}}$ of the $^3P_0$ model is given by

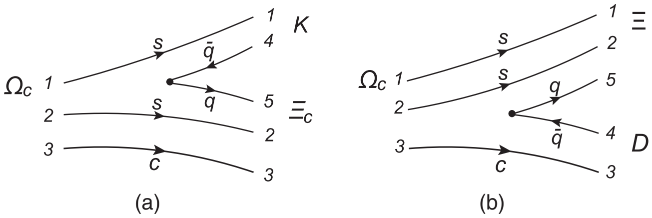

FIG. 2. A diagram for the decay processes $\Omega_c^0 \to \Xi_c K$ and $\Omega_c^0 \to \Xi D$, where the quarks have been numbered.

$$\hat{\mathcal{T}} = -3\gamma \sum_m \langle 1, m; 1, -m | 0, 0\rangle \iint d^3\vec{k}_4 d^3\vec{k}_5 \delta^3(\vec{k}_4 + \vec{k}_5) \times \mathcal{Y}_1^m\left(\frac{\vec{k}_4 - \vec{k}_5}{2}\right)\omega^{(4,5)}\varphi_0^{(4,5)}\chi_{1,-m}^{(4,5)} d_4^\dagger(\vec{k}_4) d_5^\dagger(\vec{k}_5) \quad (10)$$

in the nonrelativistic limit. Here, $\omega_0^{(4,5)}$ and $\varphi_0^{(4,5)}$ are the color and flavor wave functions of the $\bar{q}_4 q_5$ pair created from the vacuum. Therefore, $\omega^{(4,5)} = (R\bar{R} + G\bar{G} + B\bar{B})/\sqrt{3}$ and $\varphi_0^{(4,5)} = (u\bar{u} + d\bar{d} + s\bar{s})/\sqrt{3}$ are color and flavor singlets. $\chi_{1,-m}^{(4,5)}$ represents the pair production in a spin triplet state. The solid harmonic polynomial $\mathcal{Y}_1^m(\vec{k}) \equiv |\vec{k}|\mathcal{Y}_1^m(\theta_k, \phi_k)$ reflects the momentum-space distribution of the $\bar{q}_4 q_5$. The dimensionless parameter γ describes the strength of the quark-antiquark pair created from the vacuum. The value of γ is usually fixed as a constant by fitting the well-measured partial decay widths.

When the mock state [41] is adopted to describe the spatial wave function of a hadron state, the helicity amplitude can be easily constructed in the JJ basis. The mock state for the initial state A is given by

$$|A(n_A{}^{2S_A+1}L_A^{J_A j_A}(\vec{P}_A))\rangle \equiv \omega_A^{123}\phi_A^{123}\prod_A \int d^3\vec{k}_1 d^3\vec{k}_2 d^3\vec{k}_3 \delta^3(\vec{k}_1 + \vec{k}_2 + \vec{k}_3 - \vec{P}_A) \times \Psi_{n_A}^{L_A l_A}(\vec{k}_1, \vec{k}_2, \vec{k}_3)|q_1(\vec{k}_1)q_2(\vec{k}_2)q_3(\vec{k}_3)\rangle, \quad (11)$$

where ω_A^{123} and ϕ_A^{123} are the color and flavor wave functions of baryon A. The wave function of a Ξ_c baryon in Fig. 3 can be constructed in the same way. The wave function of the K meson is

$$|C(n_C{}^{2S_C+1}L_C^{J_C j_C}(\vec{P}_C))\rangle \equiv \omega_C^{14}\phi_C^{14}\prod_C \int d^3\vec{k}_1 d^3\vec{k}_4 \times \delta^3(\vec{k}_1 + \vec{k}_4 - \vec{P}_C)\psi_{n_C}^{L_C l_C}(\vec{k}_1, \vec{k}_4)|q_1(\vec{k}_1)\bar{q}_4(\vec{k}_4)\rangle. \quad (12)$$

Here, the symbols $\prod_i$ ($i = A$, B, and C) represent the Clebsch-Gordan coefficients for the initial and final

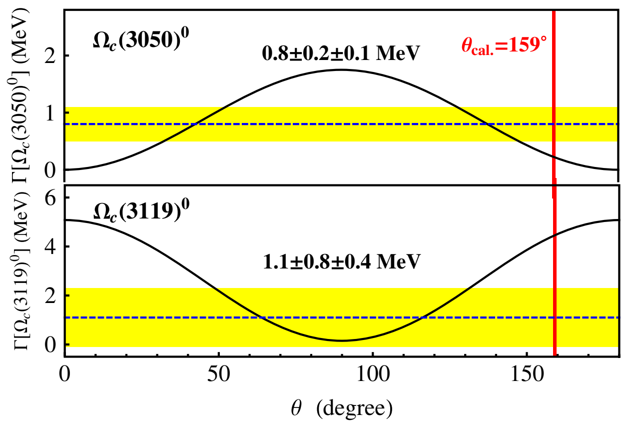

FIG. 3. The dependence of the total decay widths of the $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ states on the mixing angle θ_2 in Eq. (7). Here, the $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ are treated as $|1P, 3/2^-\rangle_L$ and $|1P, 3/2^-\rangle_H$ states, respectively. The experimental widths are also presented for comparison.

hadrons, which arise from the couplings among the orbital, spin, and total angular momentum and their projection of l_z and s_z to j_z. More specifically, $\prod_i$ ($i = A$, B, and C) are

$$\langle s_1 m_1, s_2 m_2 | s_{12} m_{12}\rangle\langle s_{12} m_{12}, L_A l_A | J_A^l j_A^l\rangle\langle J_A^l j_A^l, s_3 m_3 | J_A j_A\rangle,$$
$$\langle s_2 m_2, s_5 m_5 | s_{25} m_{25}\rangle\langle s_{25} m_{25}, L_B l_B | J_B^l j_B^l\rangle\langle J_B^l j_B^l, s_3 m_3 | J_B j_B\rangle,$$
$$\langle s_1 m_1, s_4 m_4 | S_C s_C\rangle\langle L_C l_C, S_C s_C | J_C j_C\rangle,$$

respectively. J_α^l ($\alpha = A$ or B) refers to the total angular momentum of the light diquark in the heavy baryons. To obtain the analytical amplitudes, the following simple harmonic oscillator (SHO) wave function is usually employed to construct the spatial wave function of the hadron state:

$$\psi_{Lm}^n(\mathbf{p}) = \frac{(-1)^n}{\beta^{3/2}}\sqrt{\frac{2(2n-1)!}{\Gamma(n + L + \frac{1}{2})}}\left(\frac{p}{\beta}\right)^L L_{n-1}^{L+1/2}\left(\frac{p^2}{\beta^2}\right) \times e^{-\frac{p^2}{2\beta^2}}\mathcal{Y}_{Lm}(\mathbf{p}), \quad (13)$$

where $L_{n-1}^{L+1/2}(p^2/\beta^2)$ is an associated Laguerre polynomial. The helicity amplitude $\mathcal{M}^{j_A, j_B, j_C}(q)$ is defined by

$$\langle BC|\hat{\mathcal{T}}|A\rangle = \delta^3(\vec{P}_A - \vec{P}_B - \vec{P}_C)\mathcal{M}^{j_A, j_B, j_C}(q). \quad (14)$$

Here, q represents the momentum of an outgoing meson in the rest frame of a meson A, which is given by

$$q = \frac{\sqrt{[M_A^2 - (M_B + M_C)^2][M_A^2 - (M_B - M_C)^2]}}{2M_A}. \quad (15)$$

For comparison with experiments, one needs to obtain the partial wave amplitudes $\mathcal{M}_{LS}(q)$ via the Jacob-Wick formula [42]

$$\mathcal{M}_{LS}(q) = \frac{\sqrt{2L+1}}{2J_A+1}\sum_{J_B,J_C}\langle L0Jj_A|J_Aj_A\rangle \times \langle J_Bj_B,J_Cj_C|Jj_A\rangle \mathcal{M}^{j_A,j_B,j_C}(q). \quad (16)$$

Finally, the decay width $\Gamma(A\to BC)$ is derived analytically in terms of the partial wave amplitudes in the A rest frame,

$$\Gamma(A\to BC) = 2\pi\frac{E_BE_C}{M_A}q\sum_{L,S}|\mathcal{M}_{LS}(q)|^2. \quad (17)$$

With Eqs. (10), (11), (12), (14), and (16), the full expression for $\mathcal{M}_{LS}(q)$ in the rest frame of the baryon A is

$$\mathcal{M}_{LS}(q) = -3\gamma\xi\frac{\sqrt{2L+1}}{2J_A+1}\sum_{l_i,m_j}\langle L0;Jj|J_Aj_A\rangle\langle J_Bj_B;J_Cj_C|Jj\rangle\langle s_1m_1;s_2m_2|s_{dA}m_{12}\rangle\langle s_{dA}m_{12};L_Al_A|J_A^lj_A^l\rangle\langle J_A^lj_A^l;s_3m_3|J_Aj_A\rangle$$
$$\times\langle s_2m_2;s_5m_5|s_{dB}m_{25}\rangle\langle s_{dB}m_{25};L_Bl_B|J_B^lj_B^l\rangle\langle J_B^lj_B^l;s_3m_3|J_Bj_B\rangle\langle s_1m_1;s_4m_4|S_Cs_C\rangle\langle S_Cs_C;L_Cl_C|J_Cj_C\rangle$$
$$\times\langle s_4m_4;s_5m_5|1-m\rangle\langle 1,m;1,-m|0,0\rangle\langle\omega_B^{235}\omega_C^{14}|\omega_0^{45}\omega_A^{123}\rangle\int\cdots\int d^3\vec{k}_1\cdots d^3\vec{k}_5\delta^3(\vec{k}_1+\vec{k}_2+\vec{k}_3)$$
$$\times\delta^3(\vec{q}-\vec{k}_1-\vec{k}_4)\delta^3(\vec{k}_4+\vec{k}_5)\delta^3(\vec{q}+\vec{k}_2+\vec{k}_3+\vec{k}_5)\Psi_A(\vec{k}_1,\vec{k}_2,\vec{k}_3)\Psi_B^*(\vec{k}_1,\vec{k}_2,\vec{k}_4)\psi_C^*(\vec{k}_3,\vec{k}_5)\mathcal{Y}_1^m\left(\frac{\vec{k}_4-\vec{k}_5}{2}\right), \quad (18)$$

where, $i=A,B,C$ and $j=1,2,\ldots,5$. The color matrix element $\langle\omega_B^{235}\omega_C^{14}|\omega_0^{45}\omega_A^{123}\rangle$ is a constant which can be absorbed into the parameter γ. More details for deducing these flavor matrix elements, $\xi=\langle\varphi_B^{235}\varphi_C^{14}|\varphi_0^{45}\varphi_A^{123}\rangle$, will be presented in Appendix A. In addition, we collect these obtained partial wave amplitudes in Appendix B.

IV. RESULTS AND DISCUSSIONS

As stressed here, an important motivation of this work is to test whether the five narrow Ω_c^0 states can be grouped into the $1P$ family. The primary challenge to this question is to explain their narrow widths together. According to our result given by the potential model, the $\Omega_c(3000)^0$ state could be regarded as a $|1P,1/2^-\rangle_L$ candidate with the predominant component $|0,1/2^-\rangle$. If this is true, however, $\Omega_c(3000)^0$ might be a broad state since it can decay into Ξ_cK through the S-wave channel with a large phase space. The results given in Refs. [9,10,15] confirmed this point. Thus $\Omega_c(3000)^0$ should have the predominant component $|1,1/2^-\rangle$. If this is true, the narrow width of $\Omega_c(3000)^0$ can be understood since the decay process of Ξ_cK is forbidden for the component of $|0,1/2^-\rangle$. Then the decay of $\Omega_c(3000)^0$ is strongly suppressed due to the mixing effect. According to Eq. (6), we have

$$\begin{pmatrix}|\Omega_c(3000)^0\rangle\\|\Omega_c(X)^0\rangle\end{pmatrix} = \begin{pmatrix}\cos\theta_1 & -\sin\theta_1\\ \sin\theta_1 & \cos\theta_1\end{pmatrix}\begin{pmatrix}|0,1/2^-\rangle\\|1,1/2^-\rangle\end{pmatrix}. \quad (19)$$

Then the question arises of which state among $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ can be the mixture partner of $\Omega_(3000)^0$. To solve this problem, we calculate the ratio of their decay widths with respect to the width of $\Omega_c(3000)^0$. The advantage of this method is that one can avoid the uncertainty of the phenomenological parameter γ in the QPC model. With the decay amplitudes listed in Appendix B, we first fix the mixing angle θ_1 directly. Furthermore, we obtain the value of γ by reproducing their measured decay widths. The results for θ_1 and γ are collected in Table II.

If $\Omega_c(3050)^0$ is the partner of $\Omega_c(3000)^0$, the fixed mixing angle θ_1 is about 151.8° which is comparable to the value obtained using quark potential model. But the value of γ is only about 0.208 which is too small to reproduce the widths of the other three states. The case of $\Omega_c(3066)^0$ is similar. We find that the $\Omega_c(3119)^0$ cannot be regarded as the partner of $\Omega_c(3000)^0$ in our scheme. Indeed only $\Omega_c(3090)^0$ could be a candidate for the partner of $\Omega_c(3000)^0$. The angle $\theta_1=116.3°$ fixed by $\Omega_c(3000)^0$ and $\Omega_c(3090)^0$ indicates that the quark potential model overestimates the mixing of two $1/2^-$ Ω_c^0 states. As shown

TABLE II. As the mixture partner of $\Omega_c(3000)^0$ [Eq. (19)], the values of γ and θ_1 obtained by the ratios of the decay widths of $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$ to $\Omega_c(3000)^0$.

$\Omega_c(X)^0$	$\Omega_c(3050)^0$	$\Omega_c(3066)^0$	$\Omega_c(3090)^0$	$\Omega_c(3119)^0$
γ	0.208	0.336	0.469	$\cdots$
θ_1	151.8°	128.2°	116.3°	$\cdots$

TABLE III. The partial widths of $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, and $\Omega_c(3119)^0$ with the $|1,3/2^-\rangle$, $|2,3/2^-\rangle$, and $|2,5/2^-\rangle$ assignments. The measured widths in MeV in square brackets are listed for comparison.

Assignment	$\Omega_c(3050)^0$ [0.8 ± 0.3]		$\Omega_c(3066)^0$ [3.5 ± 0.6]		$\Omega_c(3119)^0$ [1.1 ± 1.2]	
$\|1,3/2^-\rangle$	$\Xi_c K$	×	$\Xi_c K$	×	$\Xi_c K$	×
					$\Xi_c' K$	0.2
$\|2,3/2^-\rangle$	$\Xi_c K$	1.8	$\Xi_c K$	2.4	$\Xi_c K$	4.8
					$\Xi_c' K$	0.3
$\|2,5/2^-\rangle$	$\Xi_c K$	1.8	$\Xi_c K$	2.4	$\Xi_c K$	4.8
					$\Xi_c' K$	0.1

TABLE IV. The partial and total decay widths in MeV, and branching fractions in %, of the $2S$ Ω_c^0 states. The partial widths of $D\Xi$ are not listed since they are no more than 1 MeV for the $2S$ Ω_c^0 states.

Decay modes	$\Omega_c^0[2S(1/2^+)]$ Γ_i	$\Omega_c^0[2S(1/2^+)]$ $\mathcal{B}_i$	$\Omega_c^0[2S(3/2^+)]$ Γ_i	$\Omega_c^0[2S(3/2^+)]$ $\mathcal{B}_i$
$\Xi_c(2470)K$	8.4	60.9%	9.0	68.7%
$\Xi_c'(2570)K$	4.7	34.0%	0.7	5.3%
$\Xi_c^*(2645)K$	0.7	5.1%	3.4	26.0%
Theory	13.8	100%	13.1	100%
Expt. [5]	60 ± 15 ± 11			

later, the value of $\gamma = 0.469$ can also naturally reproduce the widths of $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, and $\Omega_c(3119)^0$.

In the following, we further study the $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, and $\Omega_c(3119)^0$, and try to find which one is the most likely $5/2^-$ state. To this end, we calculate their partial widths with the $|1,3/2^-\rangle$, $|2,3/2^-\rangle$, and $|2,5/2^-\rangle$ assignments. All results are listed in Table III where $\gamma = 0.469$ is used. As a $|2,5/2^-\rangle$ state, we find that the predicted widths of $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ are about 2–4 times larger than the center values of the experimental widths, which implies that the possibility of $\Omega_c(3050)^0$ or $\Omega_c(3119)^0$ as the $|2,5/2^-\rangle$ state could be preliminarily excluded. The predicted width of $\Omega_c(3066)^0$ is about 2.4 MeV which is compatible with experimental results (see Table III). So $\Omega_c(3066)^0$ is the most likely of the $J^P = 5/2^-$ candidates. We may stress that the mass of the $1P$ $5/2^-$ Ω_c^0 state was also predicted around 3.05 GeV in Refs. [21,23,32]. So the assignment of $\Omega_c(3066)^0$ as a $5/2^-$ state is probable.

To investigate the possibility of $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ as two mixtures of $|1,3/2^-\rangle$ and $|2,3/2^-\rangle$ [see Eq. (2)], we study their decay behaviors in Fig. 3 where the dependence of their total decay widths on the mixing angle θ_2 is illustrated. The mixing angle θ_2 given by the quark potential model also seems to be overestimated since the predicted width of $\Omega_c(3119)^0$ is still larger than experimental values. If we choose a suitable mixing angle, e.g., $\theta_2 \approx 130°$, the calculated widths are about 1.0 and 2.2 MeV for $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$, respectively. Obviously, the measured decay widths of $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ could be understood in the mixing scenario. In a word, the $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ are suggested to be the $|1P,3/2^-\rangle_L$ and $|1P,3/2^-\rangle_H$ candidates in our scheme.

Besides these five narrow Ω_c^0 states, LHCb also reported a broad structure around 3188 MeV, denoted as the $\Omega_c(3188)^0$, in the $\Xi_c^+ K^-$ invariant mass distribution [5]. As shown in Fig. 1, the $\Omega_c(3188)^0$ can be grouped into the $2S$ family with $J^P = 1/2^+$ or $3/2^+$. Comparable predicted masses for the $2S$ Ω_c^0 states were also obtained in Refs. [33,43–45]. To further check this possibility, we calculated the partial and total decay widths of two $2S$ Ω_c^0 states using the predicted masses. As shown in Table IV, the largest decay channel for the $2S$ Ω_c^0 states is $\Xi_c K$, which may explain why the $\Omega_c(3188)^0$ was first observed in this channel. Nevertheless, we notice that theoretical total decay widths of $2S$ Ω_c^0 states with $J^P = 1/2^+$ and $J^P = 3/2^+$ are only 13.8 and 13.1 MeV, respectively, which are much smaller than measurements even though the experimental errors are considered. We notice that LHCb only used one Breit-Wigner distribution to depict the broad structure around 3188 MeV, which is still very rough. According to our result, two Ω_c^0 states with $J^P = 1/2^+$ and $J^P = 3/2^+$ have a mass around 3.2 GeV, and mainly decay into $\Xi_c(2470)K$. This means that two $2S$ states should appear in the $\Xi_c(2470)K$ invariant mass spectrum simultaneously. Thus, we speculate that such a broad structure around 3188 MeV [5] may contain at least two resonance structures corresponding to two $2S$ Ω_c^0 states. If this is true, we may partly understand why the theoretical total decay widths of the $2S$ Ω_c^0 state with $J^P = 1/2^+$ and $J^P = 3/2^+$ are smaller than the experimental width of the $\Omega_c(3188)^0$. This speculation should be further tested with more precise experimental data.

Here we present some typical ratios of partial decay widths for the $2S$ Ω_c^0 states, i.e.,

$$\frac{\Gamma(\Omega_c^0(1/2^+) \to \Xi_c'(2570)K)}{\Gamma(\Omega_c^0(1/2^+) \to \Xi_c(2470)K)} = 0.56, \tag{20}$$

$$\frac{\Gamma(\Omega_c^0(1/2^+) \to \Xi_c^*(2645)K)}{\Gamma(\Omega_c^0(1/2^+) \to \Xi_c(2470)K)} = 0.08, \tag{21}$$

$$\frac{\Gamma(\Omega_c^0(3/2^+) \to \Xi_c'(2570)K)}{\Gamma(\Omega_c^0(3/2^+) \to \Xi_c(2470)K)} = 0.08, \tag{22}$$

$$\frac{\Gamma(\Omega_c^0(3/2^+) \to \Xi_c^*(2645)K)}{\Gamma(\Omega_c^0(3/2^+) \to \Xi_c(2470)K)} = 0.38, \tag{23}$$

which can be tested by future experiments. Finally, we should point out that $\Omega_c(3119)^0$ could not be a $2S$ candidate in our scheme. As a $2S$ state with $J^P = 1/2^+$ or $3/2^+$, the total widths of $\Omega_c(3119)^0$ are predicted to be about 10.2 and 8.3 MeV, respectively, which are much larger than the experimental result (<2.6 MeV, 95% C.L. [5]).

V. SUMMARY AND OUTLOOK

With the observation of five narrow Ω_c^0 states at LHCb [5], the study of higher orbital and radial excitations of charmed baryons is becoming a hot issue. These new Ω_c^0 states also stimulated our interest in revealing their inner structures. By performing an analysis of their masses and strong decays, we found that five narrow Ω_c^0 states, i.e., $\Omega_c(3000)^0$, $\Omega_c(3050)^0$, $\Omega_c(3066)^0$, $\Omega_c(3090)^0$, and $\Omega_c(3119)^0$, could be grouped into the P-wave charmed baryon family with a css configuration. Specifically, both $\Omega_c(3000)^0$ and $\Omega_c(3090)^0$ are the $J^P = 1/2^-$ states with mixtures of the $|0, 1/2^-\rangle$ and $|1, 1/2^-\rangle$ components, while the $\Omega_c(3050)^0$ and $\Omega_c(3119)^0$ are the $J^P = 3/2^-$ states with mixtures of the $|1, 3/2^-\rangle$ and $|2, 3/2^-\rangle$ components. Our results indicate that the mixing angles [see Eqs. (6) and (7)] determined by the simple quark potential model are overestimated. In our scheme, the $\Omega_c(3066)^0$ is most like a $|2, 5/2^-\rangle$ state. We also studied the possibility of the broad structure, $\Omega_c(3188)^0$ [5], as a $2S$ Ω_c^0 state with $J^P = 1/2^+$ or $J^P = 3/2^+$. Our results suggest that such a broad structure around 3188 MeV [5] may contain at least two resonance structures corresponding to the $2S$ Ω_c^0 states. Due to its very narrow decay width, in our scheme, $\Omega_c(3119)^0$ cannot be regarded as a $2S$ candidate.

Since some important properties—such as the spin-parity quantum numbers, the electromagnetic (EM) transitions and the hadronic modes $\Omega_c^{(*)0}\pi$—have not been measured yet, other possible assignments may also exist for these five narrow Ω_c^0 states. For example, the measurements of these Ω_c^0 states in the decay channels of $\Omega_c^{(*)}\pi$ may help us to test the pentaquark scenario [20]. The EM transitions are also useful for providing important information about their internal structures [10]. Even in our scheme, we can also explain the $\Omega_c(3050)^0$ and $\Omega_c(3090)^0$ states as the $1/2^-$ mixtures. Here the $\Omega_c(3050)^0$ state has the predominant component $|1, 1/2^-\rangle$ while $\Omega_c(3090)^0$ has the predominant component $|0, 1/2^-\rangle$. Then the exotic assignment should be considered for the $\Omega_c(3000)^0$ state since a $1/2^-$ Ω_c^0 state in the $1P$ Ω_c^0 family is suggested to have the lowest energy in most works (see the discussion in Sec. II). So an important task for future experiments like LHCb and the forthcoming Belle II is to carry out the measurement of the quantum numbers, the EM transitions, and other possible decay modes for these narrow Ω_c^0 excitations.

Although the masses and widths of these five narrow Ω_c^0 excitations could be explained under the P-wave assignment, at least two questions were not solved in the present work. The first one is why the value of γ fixed in the Ω_c^0 baryon family is so small. In our previous work [27], a larger value of γ in the QPC model was fixed by the process of $\Sigma_c(2520) \to \Lambda_c + \pi$. Although the study of meson decays indicates that the phenomenological parameter γ in the QPC model may have a complicated structure [46,47], this parameter has never been systematically investigated in the baryon sector. The second question is why the mixing angles determined by the simple quark potential model are overestimated. To answer this question, more works are needed to investigate the spin-dependent interactions of these excited Ω_c^0 states.

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APPENDIX A: DEDUCTION OF THE FLAVOR FACTORS

We take the decay process $\Omega_c^0 \to \Xi_c^+ + K^-$ as an example to show the evaluation of the flavor factors ξ in Eq. (18). The flavor wave functions of initial and final states are given as

$$\varphi_{\Omega_c^0} = ssc, \qquad \varphi_{\Xi_c^+} = (us - su)c/\sqrt{2}, \qquad \varphi_{K^-} = s\bar{u}.$$

Then the flavor factors ξ can be deduced in the following way (see Fig. 2):

$$\begin{aligned} \xi &= \langle \varphi_{\Xi_c^+}\varphi_{K^-} | \varphi_{\Omega_c^0}\varphi_0 \rangle \\ &= \left\langle \frac{(u_2 s_5 - s_2 u_5)c_3}{\sqrt{2}} \times s_1 \bar{u}_4 | s_1 s_2 c_3 \times \frac{u_5\bar{u}_4 + d_5\bar{d}_4 + s_5\bar{s}_4}{\sqrt{3}} \right\rangle \\ &= 1/\sqrt{6}. \end{aligned} \tag{A1}$$

Since the s quarks in the initial excited Ω_c^0 state have two possible ways to recombine in final states, we should multiply the value of the flavor factor in Eq. (A1) by a statistical factor of $\sqrt{2}$. The flavor factors ξ of other decay processes related to this work can be calculated in the same way.

APPENDIX B: THE PARTIAL WAVE AMPLITUDES FOR THE DECAYS OF 1*P* and 2*S* Ω_c^0 STATES

All partial wave amplitudes for the $1P$ and $2S$ Ω_c^0 decays can be written as

$$\mathcal{M}_{LS}(q) = \mathcal{P}_{(\beta_i,m_j)} q^L e^{-\frac{4fh-g^2}{4f}q^2}, \tag{B1}$$

where the $\mathcal{P}_{(\beta_i,m_j)}$ are listed in Table V. Here we have defined

$$f_s = \frac{6f\mu-(2fg\nu-g^2\mu)q^2}{16\sqrt{3}\pi^{5/4}f^{7/2}\lambda^{5/2}\beta_A^{5/2}\beta_{dA}^{3/2}\beta_B^{3/2}\beta_{dB}^{3/2}\beta_C^{3/2}},$$
$$f_p = \frac{2f[12fg\lambda(1-\lambda\beta_A^2)+\mu(5g\mu-4f\nu)]+g(g\mu-2f\nu)^2p^2}{192\pi^{5/4}f^{9/2}\lambda^{7/2}\beta_A^{7/2}\beta_{dA}^{3/2}\beta_B^{3/2}\beta_{dB}^{3/2}\beta_C^{3/2}},$$
$$f_d = \frac{g(g\mu-2f\nu)}{32\sqrt{5}\pi^{5/4}f^{7/2}\lambda^{5/2}\beta_A^{5/2}\beta_{dA}^{3/2}\beta_B^{3/2}\beta_{dB}^{3/2}\beta_C^{3/2}}, \tag{B2}$$

for the s-, p-, and d-wave decays. μ, ν, λ, f, g, and h in Eqs. (B1) and (B2) are given as

$$f = \frac{1}{2\beta_{dA}^2}+\frac{1}{2\beta_{dB}^2}+\frac{1}{2\beta_C^2}-\frac{\mu^2}{4\lambda},$$
$$g = \frac{1}{\beta_{dA}^2}+\frac{\varepsilon_3}{\beta_{dB}^2}+\frac{\varepsilon_4}{\beta_C^2}-\frac{\mu\nu}{2\lambda},$$
$$h = \frac{1}{2\beta_{dA}^2}+\frac{\varepsilon_2^2}{2\beta_B^2}+\frac{\varepsilon_3^2}{2\beta_{dB}^2}+\frac{\varepsilon_4^2}{2\beta_C^2}-\frac{\nu^2}{4\lambda},$$
$$\lambda = \frac{1}{2\beta_A^2}+\frac{1}{2\beta_B^2}+\frac{\varepsilon_1^2}{2\beta_{dA}^2}+\frac{\varepsilon_3^2}{2\beta_{dB}^2},$$
$$\mu = \frac{\varepsilon_1}{\beta_{dA}^2}+\frac{\varepsilon_3}{\beta_{dB}^2}; \qquad \nu = \frac{\varepsilon_1}{\beta_{dA}^2}+\frac{\varepsilon_2}{\beta_B^2}+\frac{\varepsilon_3^2}{\beta_{dB}^2}.$$

TABLE V. The $\mathcal{P}_{(\beta_i,m_j)}$ for different decay modes of $1P$ and $2S$ Ω_c^0 states.

$\vert j_l J^P\rangle$	$\Xi_c(2470)K$	$\Xi_c'(2570)K$	$\Xi_c^*(2645)K$
$\vert 0,1/2^-\rangle$	$-\sqrt{1/2}f_S$	×	⋯
$\vert 1,1/2^-\rangle$	×	$\sqrt{1/3}f_S$	⋯
$\vert 1,3/2^-\rangle$	×	$-\sqrt{5/9}f_D$	⋯
$\vert 2,3/2^-\rangle$	$\sqrt{4/3}f_D$	f_D	⋯
$\vert 2,5/2^-\rangle$	$\sqrt{4/3}f_D$	$-2/3f_D$	⋯
$\vert 1,1/2^+\rangle$	$\sqrt{1/2}f_P$	$\sqrt{2/3}f_P$	$\sqrt{1/3}f_P$
$\vert 1,3/2^+\rangle$	$\sqrt{1/2}f_P$	$-\sqrt{1/6}f_P$	$\sqrt{5/6}f_P$

$\varepsilon_1, \ldots, \varepsilon_4$ above are defined as

$$\varepsilon_1 = \frac{m_1}{m_1+m_2}, \qquad \varepsilon_2 = \frac{m_3}{m_1+m_3+m_5},$$
$$\varepsilon_3 = \frac{m_5}{m_2+m_5}, \qquad \varepsilon_4 = \frac{m_4}{m_1+m_4}.$$

The SHO wave function scale parameters $\beta_{d\alpha}$ ($\alpha = A$ or B) reflect the sizes of the diquarks in the initial and finial baryons. β_α reflect the distance between the light cluster and c quark. β_C reflects the size of the K meson. For the $1P$ and $2S$ Ω_c^0 states, the averaged values of the SHO wave function scale, denoted by βs, are obtained by the quark potential model, i.e., $\beta_{1P} = 0.236$ GeV and $\beta_{2S} = 0.191$ GeV. The SHO wave function scale of the light diquark, $\{ss\}$, is given as $\beta_{\{ss\}} = 0.184$ GeV. The SHO wave function scales of other hadrons related to this work can be found in our previous work [27]. m_j ($j = 1, \ldots, 5$) denotes the quark masses in Fig. 2. In the calculation of decays, the masses of u/d, s, and c quarks are taken as 0.195, 0.380, and 1.680 GeV, respectively [27].

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