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Low-lying charmed and charmed-strange baryon states

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Abstract In this work, we systematically study the mass spectra and strong decays of $1P$ and $2S$ charmed and charmed-strange baryons in the framework of non-relativistic constituent quark models. With the light quark cluster–heavy quark picture, the masses are simply calculated by a potential model. The strong decays are studied by the Eichten–Hill–Quigg decay formula. Masses and decay properties of the well-established $1S$ and $1P$ states can be reproduced by our method. $\Sigma_c(2800)^{0,+,++}$ can be assigned as a $\Sigma_{c2}(3/2^-)$ or $\Sigma_{c2}(5/2^-)$ state. We prefer to interpret the signal $\Sigma_c(2850)^0$ as a $2S(1/2^+)$ state although at present we cannot thoroughly exclude the possibility that this is the same state as $\Sigma_c(2800)^0$. $\Lambda_c(2765)^+$ or $\Sigma_c(2765)^+$ could be explained as the $\Lambda_c^+(2S)$ state or $\Sigma_{c1}^+(1/2^-)$ state, respectively. We propose to measure the branching ratio of $\mathcal{B}(\Sigma_c(2455)\pi)/\mathcal{B}(\Sigma_c(2520)\pi)$ in the future, which may disentangle the puzzle of this state. Our results support $\Xi_c(2980)^{0,+}$ as the first radial excited state of $\Xi_c(2470)^{0,+}$ with $J^P = 1/2^+$. The assignment of $\Xi_c(2930)^0$ is analogous to $\Sigma_c(2800)^{0,+,++}$, i.e., a $\Xi'_{c2}(3/2^-)$ or $\Xi'_{c2}(5/2^-)$ state. In addition, we predict some typical ratios among partial decay widths, which are valuable for experimental search for these missing charmed and charmed-strange baryons.

1 Introduction

At present, the particle data group (PDG) lists nine charmed and ten charmed-strange baryons [1]. They are $\Lambda_c(2286)^+$, $\Lambda_c(2595)^+$, $\Lambda_c(2625)^+$, $\Lambda_c(2765)^+$ (or $\Sigma_c(2765)^+$), $\Lambda_c(2880)^+$, $\Lambda_c(2940)^+$, $\Sigma_c(2800)^{0,+,++}$, $\Sigma_c(2455)^{0,+,++}$,

$\Sigma_c(2520)^{0,+,++}$, $\Xi_c(2470)^{0,+}$, $\Xi'_c(2580)^{0,+}$, $\Xi'_c(2645)^{0,+}$, $\Xi_c(2790)^{0,+}$, $\Xi_c(2815)^{0,+}$, $\Xi_c^{(\prime)}(2930)^0$, $\Xi_c(2980)^{0,+}$, $\Xi_c(3055)^{0,+}$, $\Xi_c(3080)^{0,+}$, and $\Xi_c^{(\prime)}(3123)^+$. For brevity, we call these charmed and charmed-strange baryons just charmed baryons in the following. Among these observed states, some new measurements have been performed by experiments in the past several years. The masses and widths of $\Sigma_c(2455)^{0,+,++}$, $\Sigma_c(2520)^{0,+,++}$, $\Lambda_c(2595)^+$, and $\Lambda_c(2625)^+$ have been measured with significantly small uncertainties with the efforts of CDF [2] and Belle [3]. Very recently, the Belle Collaboration updated the measurements of $\Xi'_c(2580)^{0,+}$, $\Xi'_c(2645)^{0,+}$, $\Xi_c(2790)^{0,+}$, $\Xi_c(2815)^{0,+}$, and $\Xi_c(2980)^{0,+}$ [4]. On the other hand, new decay modes for the higher excited charmed baryon states have been found by experiment. For instance, the decay channel of ΛD^+ was first found for $\Xi_c(3055)^+$ and $\Xi_c(3080)^+$, and the following ratios of branching fractions were first reported by Belle several months ago [5]:

$$\frac{\mathcal{B}(\Xi_c(3055)^+ \to \Lambda D^+)}{\mathcal{B}(\Xi_c(3055)^+ \to \Sigma_c(2455)^{++} K^-)} = 5.09 \pm 1.01 \pm 0.76,$$

$$\frac{\mathcal{B}(\Xi_c(3080)^+ \to \Lambda D^+)}{\mathcal{B}(\Xi_c(3080)^+ \to \Sigma_c(2455)^{++} K^-)} = 1.29 \pm 0.30 \pm 0.15,$$

and

$$\frac{\mathcal{B}(\Xi_c(3080)^+ \to \Sigma_c(2520)^{++} K^-)}{\mathcal{B}(\Xi_c(3055)^+ \to \Sigma_c(2455)^{++} K^-)} = 1.07 \pm 0.27 \pm 0.01,$$

where the uncertainties are statistical and systematic. Obviously, these new measurements are very useful to understand the nature of these excited charmed baryon states.

Theoretically, the charmed baryons which contain one heavy quark and two light quarks occupy a particular position in the baryon physics. Since the chiral symmetry and heavy quark symmetry (HQS) can provide some qualitative insight into the dynamics of charmed baryons, the investiga-

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tion of charmed baryons should be more helpful for improving our understanding of the confinement mechanism. The spectroscopy of charmed baryons has been investigated in various models. So far, the several kinds of quark potential models [6–10], the relativistic flux tube (RFT) model [11,12], the coupled channel model [13], the QCD sum rule [14–16], and the Regge phenomenology [17] have been applied to study the mass spectra of excited charmed baryons, and so did lattice QCD [18,19]. The strong decay behaviors of charmed baryons have been studied by several methods, such as the heavy hadron chiral perturbation theory (HHChPT) [20,21], the chiral quark model [22,23], the $^3P_0$ model [24], and a non-relativistic quark model [25]. The decays of $1P$ Λ_c and Ξ_c baryons have also been investigated by a light front quark model [26,27], a relativistic three-quark model [28], and the QCD sum rule [29].

Although many experimental and theoretical efforts have been made for the research of charmed baryons, most of the $1P$ and $2S$ charmed baryons are not yet established. Several candidates, including $\Lambda_c(2765)^+$, $\Sigma_c(2800)^{0,+,++}$, $\Xi_c(2930)^0$, and $\Xi_c(2980)^{0,+}$ are still in controversy. $\Lambda_c(2765)^+$ was first observed by the CLEO Collaboration in the decay channel of $\Lambda_c(2765)^+ \to \Lambda_c^+ \pi^+ \pi^-$ [30], and confirmed by Belle in the $\Sigma_c(2455)\pi$ mode [31,32]. Because both Λ_c^+ and Σ_c^+ excitations can decay through $\Lambda_c^+ \pi^+ \pi^-$ and $\Sigma_c(2455)\pi$, we even do not know whether the observed charmed baryon signal around 2765 MeV is the Λ_c^+ or Σ_c^+ state, or their overlapping structure [33]. In the e^+e^- annihilation process, an isotriplet state, $\Sigma_c(2800)^{0,+,++}$, was observed by Belle in the channel of $\Lambda_c^+ \pi$, and it was tentatively identified as the Σ_{c2} state with $J^P = 3/2^-$ [34]. Interestingly, another neutral resonance was later found by BaBar in the process of $B^- \to \Sigma_c^{*0} \bar{p} \to \Lambda_c^+ \pi^- \bar{p}$ with the mass $2846 \pm 8 \pm 10$ MeV and decay width 86^{+33}_{-22} MeV [35]. The higher mass and the weak evidence of $J = 1/2$ indicate that the signal observed by BaBar might be different from the Belle observation. In this paper, we denote the signal discovered by BaBar by $\Sigma_c(2850)^0$. $\Xi_c(2930)^0$, which was only seen by BaBar in the decay mode $\Lambda_c^+ K^-$ [36], still needs more confirmation. $\Xi_c(2980)^{0,+}$ was first reported by Belle in the channels $\Lambda_c^+ K^- \pi^+$ and $\Lambda_c^+ K_S^0 \pi^-$ [37], and was later confirmed by Belle [4,38] and BaBar [39] in the channels $\Xi_c'(2580)\pi$, $\Xi_c(2645)\pi$ and $\Sigma_c(2455)K$, respectively. However, the decay widths reported by Refs. [4,37–39] were quite different from each other. More experimental information as regards the charmed baryons can be found in Refs. [40–43].

Obviously, a systematic study of masses and decays is required for these unestablished charmed baryons. More importantly, most of $2S$ and $1P$ charmed baryons have not yet been detected by any experiments. Such a research can also help the future experiments to find them. In the present work, we will focus on both the mass spectra and the strong decays

of low-lying $1P$ and $2S$ charmed baryons. We pay attention to only the charmed baryons inside of which degrees of freedom of two light quarks are frozen. It means that two light quarks are not considered to be excited, neither radially nor orbitally. As illustrated in Ref. [44], these kinds of charmed excitations carry lower-excited energies, which means these excited charmed baryons may more likely be detected by experiments. Fortunately, our results indicate that most of the observed charmed baryons can be accommodated in this way.

The paper is arranged as follows. In Sect. 2, the masses of low-excited charmed baryons are calculated by the non-relativistic quark potential model. In Sect. 3, the Eichten–Hill–Quigg (EHQ) decay formula, which is employed to study the strong decays of excited charmed baryons, is introduced. The properties of low-lying charmed baryon states are fully discussed in Sect. 4. Finally, the paper ends with a conclusion and an outlook. Some detailed calculations and definitions are collected in the appendices.

2 The deduction of mass spectra

2.1 Treating charmed baryon system as a two-body problem

To study the baryon dynamics, one crucial question which should be answered is "What are the relevant degrees of freedom in a baryon?" [45]. In constituent quark models, a baryon system consists of three confined quarks. Thus the dynamics of a baryon resonance is surely more complex than a meson. Due to the HQS, however, the dynamics of charmed baryons could be greatly simplified. The HQS suggests that the couplings between a c quark and two light quarks are weak [46]. Therefore, two light quarks in a charmed baryon could first couple with each other to form a light quark cluster.[1] Then the light quark cluster couples with a charm quark, and a charmed baryon resonance forms. With this assumption, two light quarks have the same status to a c quark, including the average distances to a c quark. In the light cluster–heavy quark picture, the dynamics of heavy baryon can be simplified. In the non-relativistic constituent quark model, the spin-independent part of the Hamiltonian is

$$H = \sum_{i=1}^{3} \left(\frac{p_i^2}{2m_i} + m_i \right) + \sum_{i<j} \left(-\frac{2\alpha}{3r_{ij}} + \frac{b}{2} r_{ij} \right), \quad (1)$$

where the Cornell potential [47] is used as a phenomenological confining term. The Jacobi coordinates are usually taken to deal with the three-body problem for baryons. ρ, λ, and R are related to the quark positions by

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[1] In some work, the light quark cluster may also be named a light diquark.

$$\vec{\rho} = \vec{r}_1 - \vec{r}_2,$$
$$\vec{\lambda} = \frac{m_1\vec{r}_1 + m_2\vec{r}_2}{m_1 + m_2} - \vec{r}_Q,$$
$$\vec{R} = \frac{m_1\vec{r}_1 + m_2\vec{r}_2 + m_Q\vec{r}_Q}{m_1 + m_2 + m_Q},$$

where the indices, 1, 2, and Q, are for two light quarks and a heavy quark, respectively. The momenta $\vec{p}_\rho$, $\vec{p}_\lambda$ and $\vec{p}_R$, which are conjugate to the Jacobi coordinates above, can be defined easily. Now the spin-independent Hamiltonian becomes

$$H = \frac{p_\rho^2}{2m_\rho} + \frac{p_\lambda^2}{2m_\lambda} + \frac{p_R^2}{2M} + M + \left(-\frac{2\alpha}{3r_{12}} + \frac{b}{2}r_{12}\right)$$
$$+ \left(-\frac{2\alpha}{3r_{1Q}} + \frac{b}{2}r_{1Q}\right) + \left(-\frac{2\alpha}{3r_{2Q}} + \frac{b}{2}r_{2Q}\right), \quad (2)$$

where $m_\rho = m_1 m_2/(m_1 + m_2)$, $m_\lambda = (m_1 + m_2)m_Q/M$, and $M = m_1 + m_2 + m_Q$. According to the definitions above, the relative motion between two light quarks is usually called the ρ-mode, while the one between the center of mass of the two light quarks and the heavy quark is called the λ-mode.

As emphasized above, the average distances to c quark should be equal for two light quarks in a cluster, i.e., $\tau \equiv r_{1Q} = r_{2Q}$ (see Fig. 1). In practice, we should solve the following Schrödinger equation for the mass of a heavy baryon:

$$\left[-\frac{\nabla_\rho^2}{2m_\rho} - \frac{\nabla_\lambda^2}{2m_\lambda} + \left(\frac{b\rho}{2} - \frac{2\alpha}{3\rho}\right) + \left(b\tau - \frac{4\alpha}{3\tau}\right) + C\right]$$
$$\times \psi = E\psi. \quad (3)$$

Since the excited mode between two light quarks is not considered in this paper, the light quark cluster can be treated as a block with the antitriplet color structure and peculiar size. Specifically, a color singlet baryon system should be formed as $3_{q_1} \otimes 3_{q_2} \otimes 3_Q \ni \bar{3}_{cl.} \otimes 3_Q \ni \bar{1}_{Q-cl.}$. In this way, a heavy baryon could be treated as a quasi-two-body system. In the light quark cluster–heavy quark picture, Isgur has discussed the similarity of dynamics between heavy baryons and heavy–light mesons [48]. It should be stressed that the scenario of a light cluster–heavy quark picture is not contradictory to the one-gluon exchange interaction. Since the
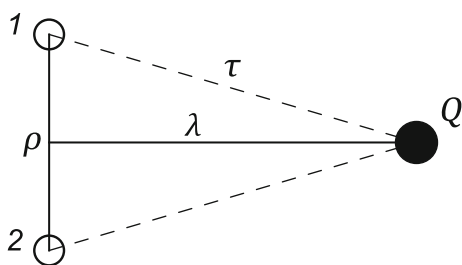

Fig. 1 A single heavy baryon in the light cluster–heavy quark picture. With the SU(3) flavor symmetry, the relation of $\tau = \sqrt{\lambda^2 + \rho^2/4}$ is obvious

color–spin interaction is proportional to the inverse of the quark masses, two light quarks in the heavy-baryon system are expected to strongly couple to each other [49]. Thus, they may develop into a quark cluster. In fact, the existence of a light quark cluster correlations was partly confirmed by the lattice QCD [50] and the Bethe–Salpeter equation [51].

Thus, we expect the light quark cluster to be an effective degree of freedom for a charmed baryon. Since the ρ mode of a cluster is not considered here, the Schrödinger equation (3) is simplified as

$$\left(-\frac{\nabla_\lambda^2}{2m_\lambda} - \frac{4\alpha}{3\tau} + b\tau + C\right)\psi = E\psi. \quad (4)$$

2.2 Adopted effective potentials

As a whole, the light quark cluster which occupies an antitriplet color structure interacts with the c quark. Then we would like to substitute λ (distance between light cluster and c quark) for τ (distance between light quark and c quark). To this end, the effective potential [52]

$$H_{Q-cl.}^{conf}(\lambda) = -\frac{4}{3}\frac{\alpha_s}{\lambda} + b\lambda^\nu - C_{Qqq'} \quad (5)$$

describes the interaction between the cluster and c quark, where ν is an adjustable parameter. This approximation can greatly decrease the computational complexity. As shown in Tables 2 and 3, the mass spectra given in this way are reasonable for the low-lying excited charmed baryons.

As a two-body problem, we treat the masses of different kinds of clusters as parameters and first fix them before calculating the masses of the low-lying charmed baryons. According to the flavor and spin, the light cluster can be classified into two kinds: one is the "scalar" cluster, and another is the "vector" cluster. Constrained by the Pauli exclusion principle, the total wave function of the light quark cluster should be antisymmetric in the exchange of two quarks. Because the spatial and color parts of this light quark cluster are always symmetric and antisymmetric, respectively, the function $|\text{flavor}\rangle \times |\text{spin}\rangle$ should be symmetric. Therefore, the scalar light quark cluster $[qq]$ ($S = 0$) is always flavor antisymmetric, and the axial-vector light quark cluster $\{qq\}$ ($S = 1$) is flavor symmetric. In Jaffe's terminology [49], the "scalar" and "vector" quark clusters are named the "good" and "bad" quark clusters, respectively. The masses of the "good" light quark cluster $[qq]$ and $[qs]$ are taken from our previous work where $m_{[qq]}$ and $m_{[qs]}$ were fixed as 450 and 630 MeV by the RFT model [11], respectively. Following Jaffe's method [49], the bad light quark cluster masses can be evaluated by the following relationships:

$$\frac{4 \times \Sigma_c(2520) + 2 \times \Sigma_c(2455)}{6} - \Lambda_c(2286) \approx 210 \text{ MeV},$$

$$\frac{4 \times \Xi_c^*(2645) + 2 \times \Xi_c'(2580)}{6} - \Xi_c(2470) \approx 150 \text{ MeV}.$$

Evidently, the masses of $\{qq\}$ and $\{qs\}$ are about 660 MeV and 780 MeV, respectively. Henceforth, we will call the Λ_c^+ and $\Xi_c^{0,+}$ baryons the $\mathcal{G}$-type baryons, and $\Sigma_c^{0,+,++}$ and $\Xi_c'^{0,+}$ the $\mathcal{B}$-type baryons for convenience.

Due to an antitriplet color structure, the spin-dependent interactions between light cluster and c quark are expected to be the same as the meson systems. In a constituent quark model [53], the spin-dependent interaction is written as

$$H_S = H_{\text{Q-cl.}}^{\text{cont}} + H_{\text{Q-cl.}}^{\text{ten}} + H_{\text{Q-cl.}}^{SO}. \tag{6}$$

The color contact interaction $H_{\text{Q-cl.}}^{\text{cont}}$ is usually given by the following form:

$$H_{\text{Q-cl.}}^{\text{cont}} = \frac{32\pi}{9} \frac{\alpha_s}{m_Q m_{\text{cl.}}} \tilde{\delta}_\sigma(\lambda) \vec{S}_Q \cdot \vec{S}_{\text{cl.}}, \tag{7}$$

where $\vec{S}_Q$ and $\vec{S}_{\text{cl.}}$ refer to the heavy quark and light cluster spins. A Gaussian-smeared function $(\sigma/\sqrt{\pi})^3 e^{-\sigma^2\lambda^2}$ is normally used for $\tilde{\delta}_\sigma(\lambda)$ [54]. If the SU(3) flavor symmetry is kept for the charmed baryons, we may modify the color contact interaction as

$$H_{\text{Q-cl.}}^{\text{cont}} = \frac{32}{9\sqrt{\pi}} \frac{\alpha_s \sigma^3}{m_Q} e^{-\sigma^2\lambda^2} \vec{S}_Q \cdot \vec{S}_{\text{cl.}}, \tag{8}$$

where the mass of a light quark cluster is just replaced by a unit. This assumption is supported by the mass differences of the $1S$ $\mathcal{B}$-type charmed baryons,

$$\Sigma_c(2520)^{++} - \Sigma_c(2455)^{++} = 64.44_{-0.24}^{+0.25} \text{ MeV},$$
$$\Xi_c(2645)^+ - \Xi_c'(2580)^+ = 70.2 \pm 3.0 \text{ MeV},$$
$$\Omega_c(2770)^0 - \Omega_c(2695)^0 = 70.7 \pm 2.6 \text{ MeV}.$$

The mass differences shown above are mainly due to the color contact interaction in the quark potential model. Clearly, these values are almost independent of the light quark cluster masses. The color tensor interaction in Eq. (6) is

$$H_{\text{Q-cl.}}^{\text{ten}} = \frac{4}{3} \frac{\alpha_s}{m_Q m_{\text{cl.}}} \frac{1}{\lambda^3} \left(\frac{3 \left(\vec{S}_Q \cdot \vec{\lambda} \right) \left(\vec{S}_{\text{cl.}} \cdot \vec{\lambda} \right)}{\lambda^2} - \vec{S}_Q \cdot \vec{S}_{\text{cl.}} \right). \tag{9}$$

Finally, $H_{\text{Q-cl.}}^{SO}$ denotes the spin–orbit interaction which contains two terms. One is the color magnetic interaction which arises from one-gluon exchange,

$$H_{\text{Q-cl.}}^{SO(cm)} = \frac{4}{3} \frac{\alpha_s}{\lambda^3} \left(\frac{\vec{S} \cdot \vec{L}}{m_Q m_{\text{cl.}}} + \frac{\vec{S}_Q \cdot \vec{L}}{m_Q^2} + \frac{\vec{S}_{\text{cl.}} \cdot \vec{L}}{m_{\text{cl.}}^2} \right), \tag{10}$$

where $\vec{S}$ denotes the spin of a baryon, $\vec{S} = \vec{S}_Q + \vec{S}_{\text{cl.}}$. Another spin–orbit interaction is the Thomas-precession term,

$$H_{\text{Q-cl.}}^{SO(tp)} = -\frac{1}{2\lambda} \frac{\partial H_{\text{Q-cl.}}^{\text{conf}}}{\partial \lambda} \left(\frac{\vec{S}_Q \cdot \vec{L}}{m_Q^2} + \frac{\vec{S}_{\text{cl.}} \cdot \vec{L}}{m_{\text{cl.}}^2} \right). \tag{11}$$

To reflect the importance of the heavy quark symmetry, we rewrite the spin-dependent interactions as

$$H_S = V_{ss} \vec{S}_Q \cdot \vec{S}_{\text{cl.}} + V_1 \vec{S}_{\text{cl.}} \cdot \vec{L} + V_2 \vec{S}_Q \cdot \vec{j}_{\text{cl.}} + V_t \hat{S}_{12}. \tag{12}$$

The degrees of freedom of the light quark cluster are characterized by its total angular momentum $\vec{j}_{\text{cl.}}$, i.e., $\vec{j}_{\text{cl.}} = \vec{S}_{\text{cl.}} + \vec{L}$. Obviously, the orbital angular momentum $\vec{L}$ of a charmed baryon in the present picture is defined by the angular momentum between light quark cluster and c quark, i.e., $\vec{L} = \vec{L}_\lambda$. The tensor operator is defined as $\hat{S}_{12} = 3 \left(\vec{S}_Q \cdot \vec{\lambda} \right) \left(\vec{S}_{\text{cl.}} \cdot \vec{\lambda} \right)/\lambda^2 - \vec{S}_Q \cdot \vec{S}_{\text{cl.}}$.

With the confining term of Eq. (5), the coefficients V_{ss}, V_1, V_2, and V_t in Eq. (12) are defined by

$$V_{ss} = \frac{1}{m_Q} \left[\frac{32\alpha_s}{9\sqrt{\pi}} \sigma^3 e^{-\sigma^2\lambda^2} - \frac{1}{m_Q} \left(\frac{2\alpha_s}{3\lambda^3} - \frac{bv}{2}\lambda^{v-2} \right) \right.$$
$$\left. - \frac{4\alpha_s}{3\lambda^3} \frac{1}{m_{\text{cl.}}} \right],$$
$$V_1 = \frac{1}{m_{\text{cl.}}} \left[\frac{1}{m_{\text{cl.}}} \left(\frac{2\alpha_s}{3\lambda^3} - \frac{bv}{2}\lambda^{v-2} \right) + \frac{4\alpha_s}{3\lambda^3} \frac{1}{m_Q} \right],$$
$$V_2 = \frac{1}{m_Q} \left[\frac{1}{m_Q} \left(\frac{2\alpha_s}{3\lambda^3} - \frac{bv}{2}\lambda^{v-2} \right) + \frac{4\alpha_s}{3\lambda^3} \frac{1}{m_{\text{cl.}}} \right],$$
$$V_t = \frac{4\alpha_s}{3\lambda^3} \frac{1}{m_Q m_{\text{cl.}}}. \tag{13}$$

2.3 Getting masses of charmed baryons

In our calculation, the following Schrödinger equation is solved for the nS state:

$$\left(-\frac{\nabla_\lambda^2}{2m_\lambda} + H_{\text{Q-cl.}}^{\text{conf}} + H_{\text{Q-cl.}}^{\text{cont}} \right) \Psi = E\Psi. \tag{14}$$

The confining and contact terms have been given by Eqs. (5) and (8). For the orbital excitations, all spin-dependent interactions are treated as the leading-order perturbations. Our calculation indicates that the color contact interaction can be ignored for the orbital excitations.

Two bases are employed to extract the mass matrix elements. One is defined by the eigenstates $|S_{\text{cl.}}, L, j_{\text{cl.}}, S_Q, J\rangle$ (jj coupling scheme) and the other is by $|S_{\text{cl.}}, S_Q, S, L, J\rangle$ (LS coupling scheme). The relation between these two bases is
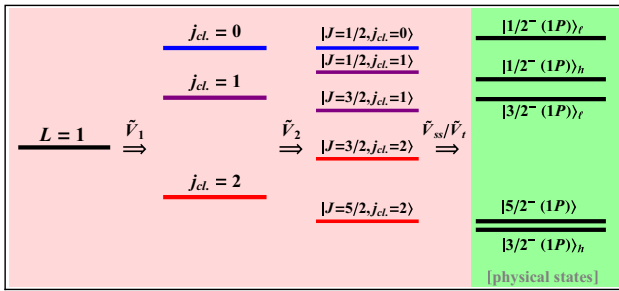
Fig. 2 A schematic diagram for the splittings of p-wave Σ_c and Ξ'_c. Here $\vec{j}_l = \vec{L} + \vec{j}_{cl.}$ and subindices ℓ and h of the last column mean low and high states in mass after including $\tilde{V}_{ss}$ and $\tilde{V}_t$ interactions

$$|[S_{cl.}, L]_{j_{cl.}}, S_Q\rangle_J = \sum_S (-1)^{S_{cl.}+S_Q+L+J}$$
$$\times \sqrt{(2J_{cl.} + 1)(2S + 1)} \qquad (15)$$
$$\times \begin{Bmatrix} s_{cl.} & L & j_{cl.} \\ J & S_Q & S \end{Bmatrix} |[S_{cl.}, S_Q]_S, L\rangle_J.$$

Due to $S_{cl.} = 0$, only $V_2 \, \vec{S}_Q \cdot \vec{j}_{cl.}$ contributes to the masses of $\mathcal{G}$-type charmed baryons. With a bad light quark cluster, however, $\mathcal{B}$-type charmed baryons have more complicated splitting structures. Within the framework of the heavy quark effective theory, the spin of an axial-vector light quark cluster, $S_{cl.}$, first couples with the orbital angular momentum L. As illustrated in Fig. 2, in the heavy quark limit $m_c \to \infty$, there are only three states which are characterized by $j_{cl.}$ for $1P$ charmed baryons. When the heavy quark spin S_Q couples with $j_{cl.}$, the degeneracy is resolved and the five states appear. These are two $J^P = 1/2^-$, two $J^P = 3/2^-$, and one $J^P = 5/2^-$ states. Lastly, the states with the same J^P mix with each other by the interactions of $V_{ss} \, \vec{S}_Q \cdot \vec{S}_{cl.}$ and $V_t \, \hat{S}_{12}$, and physical states are formed.

We now turn to a calculation of the mass matrix in the jj coupling scheme. For $1P$ states with $J^P = 1/2^-$, the mass matrix is given by

$$\langle \Phi_{1/2} \mid H_S \mid \Phi_{1/2} \rangle = \begin{pmatrix} -2V_1 - 4V_t & \frac{V_{ss}-4V_t}{\sqrt{2}} \\ \frac{V_{ss}-4V_t}{\sqrt{2}} & -V_1 - V_2 - \frac{V_{ss}}{2} - 2V_t \end{pmatrix}.$$

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

$$\langle \Phi_{3/2} \mid H_S \mid \Phi_{3/2} \rangle$$
$$= \begin{pmatrix} -V_1 + \frac{V_2}{2} + \frac{V_{ss}}{4} + 4V_t & \frac{5V_{ss}+16V_t}{4\sqrt{5}} \\ \frac{5V_{ss}+16V_t}{4\sqrt{5}} & V_1 - \frac{3V_2}{2} - \frac{3V_{ss}}{4} + \frac{4V_t}{5} \end{pmatrix}.$$

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

$$\langle j_l = 2, J^P = 5/2^- \mid H_S \mid j_l = 2, J^P = 5/2^- \rangle$$
$$= V_1 + V_2 + \frac{1}{2}V_{ss} - \frac{6}{5}V_t.$$

In the following, we denote $|S_{cl.}, L, j_{cl.}, S_Q, J\rangle$ as $|j_{cl.}, J^P\rangle$ for brevity. Then the notations $| \Phi_{1/2}\rangle$ and $| \Phi_{3/2}\rangle$ appearing above are defined by

Table 1 Values of the parameters of the non-relativistic quark potential model. The unit of b is $\mathrm{GeV}^{\nu+1}$ which varies depending on each value of ν

m_c	1.68 GeV	b	0.145	C_{Λ_C}	0.233 GeV
$m_{[qq]}$	0.45 GeV	α	0.45	C_{Σ_C}	0.100 GeV
$m_{[qs]}$	0.63 GeV	$\nu_{[\Lambda_c, \Xi_c]}$	0.84	C_{Ξ_C}	0.156 GeV
$m_{\{qq\}}$	0.66 GeV	$\nu_{[\Sigma_c, \Xi'_c]}$	0.70	$C_{\Xi'_C}$	0.060 GeV
$m_{\{qs\}}$	0.78 GeV	σ	1.00 GeV		

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

The mass matrix of $1D$ states can also be obtained by the similar procedure. As shown above, there are seven parameters in the non-relativistic quark potential model, which are $m_Q, m_{cl.}, b, \alpha, \gamma, \nu,$ and $C_{Qqq'}$. All values of parameters are listed in Table 1. If the SU(3) flavor symmetry is taken into account for the charmed and charmed-strange baryons, the dynamics of Λ_c^+ states should be like Ξ_c. The case of Σ_c and Ξ'_c is alike. Accordingly, the same value of γ is selected for the $\mathcal{G}$-type charmed baryons, as well as the case of $\mathcal{B}$-type.

We have adopted the typical values in the quark potential models for $m_c, b, \alpha,$ and ν (see Table 1). It is an effective method to investigate charmed baryons in the heavy quark–light quark cluster picture. We do not expect the values of ν to be the same for $\mathcal{G}$-type and $\mathcal{B}$-type baryons. Here, ν of Λ_c^+ / Ξ_c is slightly larger than Σ_c / Ξ'_c. The predicted masses of the low-excited charmed baryons are collected in Tables 2 and 3.

As mentioned earlier, the nonzero off-diagonal elements in mass matrices of $\langle \Phi_{1/2} \mid H_S \mid \Phi_{1/2}\rangle$ and $\langle \Phi_{3/2} \mid H_S \mid \Phi_{3/2}\rangle$ cause the mixing between two states with the same J^P but different $j_{cl.}$. However, the mechanism of mixing effects in hadron physics is still unclear. In principle, a physical hadron state with a specific J^P comprises all possible Fock states with the same total spin and parity. As the most famous member of the XYZ family, $X(3872)$ may be explained as a mixture of charmonium and the molecular state with $J^{PC} = 1^{++}$ [55]. Here we take the $|j_{cl.}, J^P\rangle$ basis to describe the mixing for the $\mathcal{B}$-type baryons. Then two physical states characterized by different masses can be denoted

$$\begin{pmatrix} |\text{High}, J^P\rangle \\ |\text{Low}, J^P\rangle \end{pmatrix} = \begin{pmatrix} \cos\phi & \sin\phi \\ -\sin\phi & \cos\phi \end{pmatrix} \begin{pmatrix} |J - 1/2, J^P\rangle \\ |J + 1/2, J^P\rangle \end{pmatrix}. \qquad (16)$$

For example, two $1P$ Σ_c states with $J^P = 1/2^-$ can be represented as

$$\begin{pmatrix} \Sigma_c(2765) \\ \Sigma_c(2702) \end{pmatrix} = \begin{pmatrix} \cos 125.4° & \sin 125.4° \\ -\sin 125.4° & \cos 125.4° \end{pmatrix}$$
$$\times \begin{pmatrix} |0, 1/2^-\rangle \\ |1, 1/2^-\rangle \end{pmatrix}. \qquad (17)$$

Table 2 Predicted masses for Λ_c^+ and Ξ_c states of ours and other approaches in Refs. [10,11,56,57] compared to experimental data [1] (in MeV)

States	Λ_c^+ baryons					Ξ_c baryons				
	PDG [1]	Prediction	Ref. [10]	Ref. [11]	Refs. [56]	PDG [1]	Prediction	Ref. [10]	Ref. [11]	Ref. [57]
$\lvert 1S, 1/2^+\rangle$	2286.86	2286	2286	2286	2265	2470.88	2470	2476	2467	2466
$\lvert 2S, 1/2^+\rangle$	2766.6	2772	2769	2766	2775	2968.0	2940	2959	2959	2924
$\lvert 3S, 1/2^+\rangle$		3116	3130	3112	3170		3265	3323	3325	
$\lvert 1P, 1/2^-\rangle$	2592.3	2614	2598	2591	2630	2791.8	2793	2792	2779	2773
$\lvert 1P, 3/2^-\rangle$	2628.1	2639	2627	2629	2640	2819.6	2820	2819	2814	2783
$\lvert 1D, 3/2^+\rangle$		2843	2874	2857	2910	3054.2	3033	3059	3055	3012
$\lvert 1D, 5/2^+\rangle$	2881.53	2851	2880	2879	2910	3079.9	3040	3076	3076	3004
$\lvert 2P, 1/2^-\rangle$	2939.3	2980	2983	2989	3030	3122.9	3140	3179	3195	
$\lvert 2P, 3/2^-\rangle$		3004	3005	3000	3035		3164	3201	3204	

Table 3 Predicted masses for Σ_c and Ξ'_c states of ours and other approaches in Refs. [9,10,56,58] compared to experimental data [1] (in MeV)

States	Σ_c baryons						Ξ'_c baryons			
	PDG [1]	Prediction	Ref. [9]	Ref. [10]	Refs. [56]	Ref. [58]	PDG [1]	Prediction	Ref. [9]	Ref. [10]
$\lvert 1S, 1/2^+\rangle$	2452.9	2456	2439	2443	2440	2452	2575.6	2579	2579	2579
$\lvert 1S, 3/2^+\rangle$	2517.5	2515	2518	2519	2495	2501	2645.9	2649	2654	2649
$\lvert 2S, 1/2^+\rangle$	2846[a]	2850	2864	2901	2890	2961		2977	2984	2983
$\lvert 2S, 3/2^+\rangle$		2876	2912	2936	2985	2996		3007	3035	3026
$\lvert 3S, 1/2^+\rangle$		3091		3271	3035	3381		3215		3323
$\lvert 3S, 3/2^+\rangle$		3109		3293	3200	3403		3236		3396
$\lvert 1P, 1/2^-\rangle_l$		2702	2795	2713	2765	2832		2839	2928	2854
$\lvert 1P, 1/2^-\rangle_h$	2766.6	2765	2805	2799	2770	2841		2900	2934	2936
$\lvert 1P, 3/2^-\rangle_l$		2785	2761	2773	2770	2812	2931	2921	2900	2912
$\lvert 1P, 3/2^-\rangle_h$	2801	2798	2798	2798	2805	2822		2932	2931	2935
$\lvert 1P, 5/2^-\rangle$		2790	2799	2789	2815	2796		2927	2921	2929
$\lvert 1D, 1/2^+\rangle$		2949	3014	3041	3005			3075	3132	3163
$\lvert 1D, 3/2^+\rangle_l$		2952	3005	3040	3060			3089	3127	3160
$\lvert 1D, 3/2^+\rangle_h$		2964	3010	3043	3065			3081	3131	3167
$\lvert 1D, 5/2^+\rangle_l$		2942	2960	3023	3065			3091	3087	3153
$\lvert 1D, 5/2^+\rangle_h$		2962	3001	3038	3080			3077	3123	3166
$\lvert 1D, 7/2^+\rangle$		2943	3015	3013	3090			3078	3136	3147
$\lvert 2P, 1/2^-\rangle_l$		2971	3176	3125	3185	3245		3094	3294	3267
$\lvert 2P, 1/2^-\rangle_h$		3018	3186	3172	3195	3256		3144	3300	3313
$\lvert 2P, 3/2^-\rangle_l$		3036	3147	3151	3195	3223		3172	3269	3293
$\lvert 2P, 3/2^-\rangle_h$		3044	3180	3172	3210	3233		3165	3296	3311
$\lvert 2P, 5/2^-\rangle$		3040	3167	3161	3220	3203		3170	3282	3303

[a] The mass value for the $\lvert 2S, 1/2^+\rangle$ state is taken from the measurement of BaBar [35]

Here we have denoted the physical states by their masses (see Table 3). The mixing angles for other states in Table 3 with the same J^P are listed in Table 4.

Our results of the mixing angles in Table 4 indicate that the heavier $1/2^-$ state, $\Sigma_c(2765)$, is dominated by a $\lvert 1, 1/2^-\rangle$ component, while $\Sigma_c(2702)$ is dominated by a $\lvert 0, 1/2^-\rangle$ component. For two $3/2^-$ states, the light $\Sigma_c(2785)$ is dom-inated by $\lvert 2, 1/2^-\rangle$, and the heavy $\Sigma_c(2798)$ by $\lvert 1, 1/2^-\rangle$. The mixing of $2P$ states is similar to the $1P$ states. For the $1D$ states, one notices that both $3/2^+$ and $5/2^+$, with heavier masses, are dominated by smaller $j_{cl.}$ components.

The uncertainty may exist in the mixing angles. Firstly, the loop corrections to the spin-dependent one-gluon-exchange potential may be important for the heavy–light hadrons. As

Table 4 The mixing angles for the $1P$, $2P$, and $1D$ Σ_c/Ξ_c' states

	$1P(1/2^-)$	$1P(3/2^-)$	$2P(1/2^-)$	$2P(3/2^-)$	$1D(3/2^+)$	$1D(5/2^+)$
Σ_c	125.4°	−156.8°	124.8°	−151.4°	172.2°	−175.6°
Ξ_c'	125.0°	−153.6°	124.3°	−145.1°	168.9°	−173.8°

an example, the lower mass of $D_s(2317)^\pm$ compared with the old calculations [53] can be well explained by the corrected spin-dependent potential [59,60]. If we use this type of potential in our calculation, of course, the mixing angle will change. Secondly, the mixing angles depend on the parameters. Thirdly, there are other mechanisms, e.g., hadron loop effects [61], which may contribute to the mixing phenomenon in hadron physics. Anyway, we expect that the mixing angles in Table 4 reflect the main features of the mixing states. Due to the uncertainties of the mixing angles, however, we ignore the mixing effects as the first step to the study of the decays of charmed excitations in the next subsection. Obviously, it is a good approximation only when the mixing effects are not large. Fortunately, this crude procedure is partially supported by the former analysis of charmed mesons [62–64]. If the decay properties obtained in this way describe the principal characteristics of the mixing states, the angles obtained by the potential model may be overestimated.

2.4 Simple harmonic oscillator (SHO) β values

In the next section, the Okubo–Zweig–Iizuka (OZI) allowed decays will be calculated for the $1P$ and $2S$ charmed baryons where the SHO wave functions are used to evaluate the transition factors via the $^3P_0$ model. We will also discuss the mixing effects for the decays of the relevant states. Following the method of Ref. [65], all values of the SHO wave function scale, denoted β in the following, are calculated (see Table 5). The values of β reflect the distances between the light quark cluster and the c quark.

In our calculation of strong decays, we will consider the structures of light diquarks. What is more important is that the possible final states of an excited charmed baryon may

Table 5 The meson effective β values in GeV

States	Λ_c^+	Ξ_c	Σ_c	Ξ_c'
1S				
$1/2^+$	0.291	0.331	0.335	0.362
$3/2^+$			0.296	0.315
2S	0.145	0.162	0.144	0.152
3S	0.102	0.113	0.098	0.103
1P	0.184	0.205	0.182	0.192
2P	0.117	0.130	0.112	0.118
1D	0.142	0.156	0.136	0.143

contain a light flavor meson, a charmed meson, a light flavor baryon, e.g., π, K, D, p, and Λ. For the β of these hadrons, the following potential will be used:

$$V(r) = \mathbf{F}_{q_1} \cdot \mathbf{F}_{q_2} \left(\frac{\alpha_s}{r} - \frac{3}{4}br + \frac{3}{4}C + \frac{32\alpha_s\sigma^3 e^{-\sigma^2 r^2}}{9\sqrt{\pi}m_q m_q} \vec{S}_{q_1} \cdot \vec{S}_{q_2} \right),$$
$$(18)$$

where

$$\langle \mathbf{F}_{q_1} \cdot \mathbf{F}_{q_2} \rangle = \begin{cases} -\dfrac{4}{3} & \text{for } q_1\bar{q}_2 \quad (19a) \\ -\dfrac{2}{3} & \text{for } q_1 q_2 \quad (19b) \end{cases}$$

Here, the parameters α_s and b are taken as 0.45 and 0.145 GeV$^{\nu+1}$ as in Table 1, respectively. To reproduce the masses of the light quark clusters in Table 1, the masses of u/d, s are fixed as 0.195 and 0.380 GeV. While σ and C are treated as adjustable parameters, the masses of the π/ρ, K/K^*, D/D^*, p/Δ, and Λ families are fitted with experimental data. Meanwhile, the values of β for the corresponding states are also obtained; they are collected in Table 6.

Before ending this section, we briefly summarize the complicated deduction presented here. Firstly, the dynamics of heavy baryon is simplified as a two-body system when the symmetric configuration is considered. Secondly, the mass matrices were calculated in the jj coupling scheme. By solving the Schrödinger equation, we obtain the mass spectra and mixing angles for the relevant states. For estimating the two-body strong decays in the next section, finally, we also present the values of the SHO wave function scale for all initial and final states.

3 Strong decays

In this section, we will use the formula provided by Eichten, Hill, and Quigg (EHQ) [66] to extract the decay widths of excited charmed baryons. Since the dynamical behavior of the heavy–light hadrons is governed by the light degrees of freedom in the limit of heavy quark symmetry, a doublet formed by two states with the same $j_{\text{cl.}}$ but different J shall have similar decay properties. More specifically, the transitions between two doublets should be determined by a single amplitude which is proportional to the products of four Clebsch–Gordan coefficients [46]. Some typical ratios of excited charmed baryons with negative parity were pre-

Table 6 The effective β values in GeV for the light quark cluster and various hadrons (the second row). The values of σ and C are given in the square brackets for various hadron structures (the third row)

$[qq]$	$\{qq\}$	$[qs]$	$\{qs\}$	π	ρ	K	K^*	D	D^*	p	Λ
0.201	0.143	0.207	0.159	0.298	0.179	0.291	0.201	0.250	0.230	0.189	0.226
[1.17, 0.39]		[1.57, 0.38]		[0.73, 0.63]		[0.83, 0.48]		[1.20, 0.63]		[−, 0.38]	[−, 0.26]

dicted by this law [46]. Later, a more concise formula (the EHQ formula) was proposed for the widths of heavy–light mesons [66]. The EHQ formula has been applied systematically to the decays of excited open-charm mesons [62–64]. Recently, the EHQ formula has been extended to the study of the decay properties of $1D$ Λ_c and Ξ_c states [11].

For the charmed baryons, the EHQ formula can be written as

$$\Gamma^{A \to BC}_{jC,\ell} = \xi \left(\mathcal{C}^{s_Q,j_B,J_B}_{jC,j_A,J_A} \right)^2 \left| \mathcal{M}^{j_A,j_B}_{jC,\ell}(q) \right|^2 q^{2\ell+1} e^{-q^2/\tilde{\beta}^2}, \tag{20}$$

where ξ is the flavor factor given in Table 13 in Appendix B. $q = |\vec{q}|$ denotes the three-momentum of a final state in the rest frame of an initial state. A and B represent the initial and final heavy–light hadrons, respectively. C denotes the light flavor hadron (see Fig. 3). The explicit expression of $\tilde{\beta}$ is given in Eq. (A.11) in Appendix A. In addition, $\mathcal{C}^{s_Q,j_B,J_B}_{jC,j_A,J_A}$ is a normalized coefficient given by the following equation:

$$\mathcal{C}^{s_Q,j_B,J_B}_{jC,j_A,J_A} = (-1)^{J_A+j_B+jc+s_Q} \sqrt{(2j_A+1)(2J_B+1)} \\ \times \begin{Bmatrix} s_Q & j_B & J_B \\ jC & J_A & j_A \end{Bmatrix}, \tag{21}$$

where $\vec{j}_C \equiv \vec{s}_C + \vec{\ell}$. The symbols s_C and ℓ represent the spin of the light hadron C and the orbital angular momentum relative to B, respectively. The transition factors $\mathcal{M}^{j_A,j_B}_{jC,\ell}(q)$ involved in the concrete dynamics can only be calculated by various phenomenological models. For the decays of heavy–light mesons, transition factors have been calculated by the relativistic chiral quark model [67] and the $^3P_0$ model [62, 64,68]. In our work, we will employ the $^3P_0$ model [69–71] to obtain the transition factors. More details of an estimate of the transition factors are given in Appendix A.
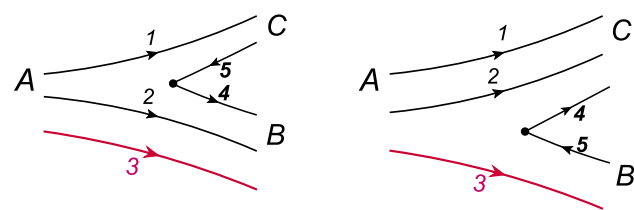

Fig. 3 The two topological diagrams for an excited charmed baryon A decaying into the final states B and C. The *brown line* 3 denotes a charm quark

4 Discussion

4.1 Experimentally well-established $1S$ and $1P$ states

At present, all the ground states and $1P$ $\mathcal{G}$-type charmed states have been experimentally established [1]. These states have been observed, at least, by two different collaborations, and their properties including masses and decays have been well determined. With good precision, the strong decays of these states provide a crucial test of our method.

Among the $1S$ charmed baryons, the measurements of $\Sigma_c(2455)$ and $\Sigma_c(2520)$ have been largely improved [2,3] (see Table 7). In our calculation, the mass and decay width of $\Sigma_c(2520)^{++}$ measured by CDF will be taken as input data to fix the constant γ peculiar to the $^3P_0$ model. With the transition factor for the process $\Sigma_c(2520) \to \Lambda_c(2286) + \pi$ [see Eq. (A.12) in Appendix A], the value of γ is fixed as 1.296.[2]

As shown in Tables 8 and 9, the predicted widths of other ground and $1P$ $\mathcal{G}$-type charmed baryons are well consistent with experiments. Our results of mass spectra and decay widths indicate that $\Lambda_c(2595)^+$, $\Lambda_c(2625)^+$, $\Xi_c(2790)^{0,+}$, and $\Xi_c(2815)^{0,+}$ can be accommodated with the $1P$ $\mathcal{G}$-type charmed baryons. $\Lambda_c(2595)^+$ and $\Xi_c(2790)^{0,+}$ can be classified into the $1/2^-$ states while $\Lambda_c(2625)^+$ and $\Xi_c(2815)^{0,+}$ into the $3/2^-$ states. The predicted mass splittings between the $1P$ $1/2^-$ and $3/2^-$ states are 25 MeV and 27 MeV for the Λ_c and Ξ_c baryons, respectively, which are also consistent with the experiments. The assignments of $\Lambda_c(2595)^+$, $\Lambda_c(2625)^+$, $\Xi_c(2790)^{0,+}$, and $\Xi_c(2815)^{0,+}$ are also supported by other work [9–12] in which the light quark cluster scenario was also employed. In addition, the mass spectra obtained by different types of the quark potential models in the three-body picture also support these assignments [7,8,56–58]. However, the investigations by QCD sum rules indicate that these $1P$ candidates may have more complicated structures [14–16]. Especially, the work by Chen et al. suggested that $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ form the heavy doublet $\tilde{\Lambda}_{c1}(1/2^-, 3/2^-)$ (the same assignments as the case of $\Xi_c(2790)^{0,+}$ and $\Xi_c(2815)^{0,+})$ [16], which is

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[2] For the different conventions to extract the color and flavor factors, the value of γ here is different from those in Refs. [24,73,74]. The deviation, of course, does not affect the predictions since γ is regarded as an adjustable parameter in the $^3P_0$ model.

Table 7 The masses and widths (in units of MeV) of $\Sigma_c(2455)^{++}$ and $\Sigma_c(2520)^{++}$ measured by CDF [2] and Belle [3]

$\Sigma_c(2455)^{++}$	$2453.90 \pm 0.13 \pm 0.14$	2.34 ± 0.47	CDF
	$2453.97 \pm 0.01 \pm 0.02 \pm 0.14$	$1.84 \pm 0.04^{+0.07}_{-0.20}$	Belle
$\Sigma_c(2520)^{++}$	$2517.19 \pm 0.46 \pm 0.14$	15.03 ± 2.52	CDF
	$2518.45 \pm 0.10 \pm 0.02 \pm 0.14$	$14.77 \pm 0.25^{+0.18}_{-0.30}$	Belle

Table 8 Open-flavor strong decay widths of $1S$ Σ_c and Ξ'_c in MeV

$1S$ Σ_c and Ξ'_c			
$1/2^+$		$3/2^+$	
$\Sigma_c(2455)^{++}$	$\Xi'_c(2580)^+$	$\Sigma_c(2520)^{++}$	$\Xi'_c(2645)^+$
$\Lambda_c^+\pi^+$ 1.53		$\Lambda_c^+\pi^+$ Input	$\Xi_c^0\pi^+$ 1.54
			$\Xi_c^+\pi^0$ 1.01
1.53	–	Input	2.55
$1.89^{+0.09}_{-0.18}$ [1]	–	14.9 ± 1.5 [1]	2.6 ± 0.6 [72]

Table 9 Open-flavor strong decay widths of $1P$ Λ_c and Ξ_c in MeV

$1P$ Λ_c and Ξ_c			
$1/2^-$		$3/2^-$	
$\Lambda_c(2595)^+$	$\Xi_c(2790)^+$	$\Lambda_c(2625)^+$	$\Xi_c(2815)^+$
$\Sigma_c\pi$ 2.78	$\Xi'_c\pi$ 6.01	$\Sigma_c\pi$ 0.04	$\Xi'_c\pi$ 0.15
			$\Xi_c^*\pi$ 4.09
2.78	6.01	0.04	4.24
2.6 ± 0.6 [1]	8.9 ± 1.4 [4]	<0.97 [1]	2.43 ± 0.37 [4]

different from our conclusion. Since the quantum numbers of J^P have not yet been determined for these $1P$ charmed states, more experiments are required in the future.

4.2 $1P$ $\Sigma_c^{0,+,++}$ states

As shown in Tables 2 and 3, the masses of $1P$ Σ_c states are predicted in the range of 2700–2800 MeV. Then $\Sigma_c(2765)^+$ and $\Sigma_c(2800)^{0,+,++}$ can be grouped into the candidates of $1P$ Σ_c family. The predicted mass of the $|1P, 1/2^-\rangle_h$ state is about 2765 MeV, which is in good agreement with the measured mass of $\Sigma_c(2765)^+$. In addition, the theoretical result for the decay width of the $\Sigma_{c1}(1/2^-)$ state in Table 10 is about 63.52 MeV, which is also in agreement with the measurements [30–32]. Furthermore, the signal of $\Sigma_c(2765)^+$ has been observed in the $\Sigma_c(2455)\pi$ intermediate state, while there is no clear evidence for the decay of $\Sigma_c(2765)^+$ through $\Sigma_c(2520)\pi$ [31,32]. This is also consistent with our results of the $|1P, 1/2^-\rangle_h$. Based on the combined analysis of the mass spectrum and strong decays, we therefore conclude that $\Sigma_c(2765)^+$ could be regarded as a good candidate of $\Sigma_{c1}(1/2^-)$. Considering the uncertainties of the quark potential models, the masses obtained by Refs. [9,10,56] are not contradictory to our assignment to $\Sigma_c(2765)^+$.

According to the predicted masses in Table 3, $\Sigma_c(2800)^{0,+,++}$ could be assigned to either the $|1P, 3/2^-\rangle_l$, or the $|1P, 3/2^-\rangle_h$, or the $|1P, 5/2^-\rangle$ states. When we consider the decay properties of these three states (see Table 10), the possibility of assignment to the $|1P, 3/2^-\rangle_l$ state can be excluded since the Belle Collaboration observed this state in the $\Lambda_c^+\pi$ mode.[3] At present, the Belle Collaboration tentatively identified $\Sigma_c(2800)^{0,+,++}$ as members of the $\Sigma_{c2}(3/2^-)$ isospin triplet, which agrees with our results of both mass spectrum and strong decays. When the measured mass of $\Sigma_c(2800)^0$ (2806 MeV) is used for the $\Sigma_{c2}(3/2^-)$ state, the predicted width is about 40.1 MeV, which is comparable with the experiment [34]. However, we notice that the quantum number J^P of $\Sigma_c(2800)^{0,+,++}$ has not yet been measured. Then the possibility of this state as the $\Sigma_{c2}(5/2^-)$ candidate cannot be excluded by our results since the decay mode of $\Lambda_c^+\pi$ is dominant for this state. In addition, the predicted mass and total width of the $\Sigma_{c2}(5/2^-)$ state are also compatible with experimental data of the $\Sigma_c(2800)^{0,+,++}$ baryon. Therefore, we would like to point out that the signal of $\Sigma_c(2800)^{0,+,++}$ found by Belle might be their overlapping structure. We hope the future experiments measure the following branching ratios to disentangle this state.

For the $\Sigma_{c2}(3/2^-)$ state,

$$\frac{\mathcal{B}(\Sigma_{c2}(3/2^-) \rightarrow \Sigma_c(2455)\,\pi)}{\mathcal{B}(\Sigma_{c2}(3/2^-) \rightarrow \Sigma_c(2520)\,\pi)} = 1.90; \qquad (22)$$

$$\frac{\mathcal{B}(\Sigma_{c2}(3/2^-) \rightarrow \Lambda_c(2286)\,\pi)}{\mathcal{B}(\Sigma_{c2}(3/2^-) \rightarrow \Sigma_c(2455)\,\pi)} = 4.07. \qquad (23)$$

For the $\Sigma_{c2}(5/2^-)$ state,

$$\frac{\mathcal{B}(\Sigma_{c2}(5/2^-) \rightarrow \Sigma_c(2455)\,\pi)}{\mathcal{B}(\Sigma_{c2}(5/2^-) \rightarrow \Sigma_c(2520)\,\pi)} = 0.58; \qquad (24)$$

$$\frac{\mathcal{B}(\Sigma_{c2}(5/2^-) \rightarrow \Lambda_c(2286)\,\pi)}{\mathcal{B}(\Sigma_{c2}(5/2^-) \rightarrow \Sigma_c(2455)\,\pi)} = 9.85. \qquad (25)$$

As mentioned earlier, the signal $\Sigma_c(2850)^0$ discovered by the BaBar Collaboration may be a $J = 1/2$ state. If $\Sigma_c(2850)^0$ is the $1/2^+(2S)$ state, the corresponding ratios (see Sect. 4.5) are different from Eqs. (22)–(25). So the measurements of

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[3] Even if the possible mixing between $\Sigma_{c1}(3/2^-)$ and $\Sigma_{c2}(3/2^-)$ is considered, the partial width of $\Lambda_c^+\pi$ is only 3.87 MeV for the $|1P, 3/2^-\rangle_l$ state where the mixing angle obtained in Table 4 has been used.

Table 10 The partial and total decay widths of $1P$ Σ_c states in MeV

Decay modes	$1/2^-$ ($1P$)		$3/2^-$ ($1P$)		$5/2^-$ ($1P$)
	$\Sigma_{c0}(2702)$	$\Sigma_{c1}(2765)$	$\Sigma_{c1}(2798)$	$\Sigma_{c2}(2785)$	$\Sigma_{c2}(2790)$
$\Lambda_c\pi$	3.64	$\times$	$\times$	24.06	24.63
$\Sigma_c(2455)\pi$	$\times$	58.94	3.48	5.22	2.50
$\Sigma_c(2520)\pi$	$\times$	1.70	63.72	2.47	4.34
$\Lambda_c(2595)\pi$		2.88	2.31	1.93	0.03
$\Lambda_c(2625)\pi$			3.12	0.07	0.63
Theory	3.64	63.52	72.63	33.75	32.13
Expt. [1]		≈ 50		72^{+22}_{-15}	

these ratios of branching fractions can help us understand the nature of $\Sigma_c(2800)^{0,+,++}$ and $\Sigma_c(2850)^0$.

Although, at present, the $\Sigma_{c0}(1/2^-)$ and $\Sigma_{c1}(3/2^-)$ states are still missing in the experiments, our results indicate that the $\Sigma_{c0}(1/2^-)$ state may be a narrow resonance and its predominant decay channel $\Lambda_c^+\pi$ is only about 3.64 MeV (see Table 10). Since the decay mode of $\Sigma_c(2520)\pi$ is the largest for the $\Sigma_{c1}(3/2^-)$ state, we suggest to search this channel for this state in future experiments. In the heavy quark limit, the branching ratio for the $\Sigma_{c1}(3/2^-)$ state,

$$\frac{\mathcal{B}(\Sigma_{c1}(3/2^-) \to \Sigma_c(2455)\,\pi)}{\mathcal{B}(\Sigma_{c1}(3/2^-) \to \Sigma_c(2520)\,\pi)} = 0.04, \qquad (26)$$

is much smaller than $\Sigma_{c2}(3/2^-)$ (Eq. 22).

4.3 $1P$ $\Xi_c'^{\,0,+}$ states

As shown in Table 3, the predicted mass of $1P$ Ξ_c' is in the range from 2840 to 2930 MeV. Then the resonance structure observed by BaBar [36] in the decay channel $B^- \to \Xi_c'(2930)^0\bar{\Lambda}_c^- \to \Lambda_c^+K^-\bar{\Lambda}_c^-$ with an invariant mass of 2.93 GeV could be a good candidate of $1P$ Ξ_c' members. The results of the decays in Table 11 favor $\Xi_c'(2930)^0$ as the $\Xi_{c2}'(3/2^-)$ or $\Xi_{c2}'(5/2^-)$ state. Then $\Xi_c'(2930)^0$ might be regarded as the strange partner of $\Sigma_c(2800)^{0,+,++}$ by our results. Interestingly, the mass difference between $\Xi_c'(2930)^0$

and $\Sigma_c(2800)^{0,+,++}$ is about 130 MeV, which is comparable with the mass differences among sextet states of the ground charmed baryons [21]. With a chiral quark model, Liu et al. also analyzed the $\Xi_c'(2930)^0$ by the two-body strong decays [23]. Their results support $\Xi_c'(2930)^0$ as the $|\Xi_c'^2 P_\lambda,\ 1/2^-\rangle$ or $|\Xi_c'^4 P_\lambda,\ 1/2^-\rangle$ state. Since the heavy quark symmetry was not considered in Ref. [23], the notations of charm-strange baryons in Ref [23] are different from our $\Xi_{c0}'(1/2^-)$ and $\Xi_{c1}'(1/2^-)$. Although the results in Table 11 indicate that the Λ_c^+K decay mode dominates the decay of $\Xi_{c0}'(1/2^-)$ state, the mass of this state is predicted to be about 2840 MeV, which is much smaller than $\Xi_c'^0(2930)$. In addition, the Λ_c^+K decay mode is forbidden for the $\Xi_{c1}'(1/2^-)$ state. Thus, according to our results, $\Xi_c'^0(2930)$ is unlikely to be a $1P$ state with $J^P = 1/2^-$.

Another charm-strange baryon, $\Xi_c(2980)^{0,+}$, is slightly higher than the predicted mass range of $1P$ Ξ_c' states. This state has been observed in $\Sigma_c(2455)K$, $\Xi_c'(2580)\pi$, $\Xi_c'(2645)\pi$, and nonresonant $\Lambda_c^+\bar{K}\pi$ decay channels. However, it was not seen in the decay modes of $\Lambda_c^+\bar{K}$ and $\Xi_c\pi$ [37–39]. Comparing the mass and decay properties of $\Xi_c(2980)^{0,+}$ with our results, the possibility of a $1P$ Ξ_c' state might be excluded. As shown in the next subsection, $\Xi_c(2980)^{0,+}$ could be a good $2S$ Ξ_c candidate. Based on our results on strong decays, we find that the $\Xi_{c1}'(1/2^-)$ and $\Xi_{c1}'(3/2^-)$ are quite narrow (see Table 11).

Table 11 The partial and total decay widths of $1P$ Ξ_c' states in MeV

Decay modes	$1/2^-$ ($1P$)		$3/2^-$ ($1P$)		$5/2^-$ ($1P$)
	$\Xi_{c0}'(2839)$	$\Xi_{c1}'(2900)$	$\Xi_{c1}'(2932)$	$\Xi_{c2}'(2921)$	$\Xi_{c2}'(2927)$
$\Lambda_c K$	46.59	$\times$	$\times$	11.59	12.43
$\Xi_c\pi$	4.39	$\times$	$\times$	7.42	7.75
$\Xi_c'(2580)\pi$	$\times$	9.44	0.76	1.20	0.57
$\Xi_c'(2645)\pi$	$\times$	0.52	3.23	0.75	1.31
$\Xi_c(2790)\pi$			0.01		
Theory	50.98	9.96	4.00	20.96	22.06
Expt.				$36 \pm 7 \pm 11$ [36]	

Table 12 The partial and total decay widths of $2S$ Λ_c^+ and $\Xi_c^{+,0}$ states in MeV

$1/2^+$ $(2S)$					$1/2^+$ $(2S)'$				$3/2^+$ $(2S)'$			
$\Lambda_c(2765)^+$		$\Xi_c(2980)$			$\Sigma_c(2850)^0$		$\Xi_c'(3000)$		$\Sigma_c(2880)^0$		$\Xi_c'(3030)$	
$\Sigma_c(2455)\pi$	26.23	$\Sigma_c(2455)K$	3.14		$\Lambda_c^+\pi$	35.11	Λ_c^+K	17.42	$\Lambda_c^+\pi$	34.96	Λ_c^+K	18.37
$\Sigma_c(2520)\pi$	19.28	$\Xi_c'(2580)\pi$	11.47		$\Sigma_c(2455)\pi$	57.16	$\Xi_c'(2580)\pi$	12.56	$\Sigma_c(2455)\pi$	15.98	$\Xi_c'(2580)\pi$	3.50
		$\Xi_c'(2645)\pi$	12.83		$\Sigma_c(2520)\pi$	17.54	$\Xi_c'(2645)\pi$	4.13	$\Sigma_c(2520)\pi$	54.52	$\Xi_c'(2645)\pi$	12.92
					$\Lambda_c(2595)\pi$	6.92	$\Xi_c(2790)\pi$	5.89	$\Lambda_c(2595)\pi$	1.07	$\Xi_c(2790)\pi$	0.40
					$\Lambda_c(2625)\pi$	1.57	$\Xi_c(2815)\pi$	0.13	$\Lambda_c(2625)\pi$	7.62	$\Xi_c(2815)\pi$	6.22
					D^0n	0.03	$\Sigma_c(2455)K$	15.34	D^0n	3.03	$\Sigma_c(2455)K$	6.49
											$D^0\Lambda$	0.01
45.51		27.44			118.33		55.47		117.18		47.91	
≈ 50 [1]		$28.1\pm2.4^{+1.0}_{-5.0}$ [4]			86^{+33}_{-22} [35]							

4.4 $2S$ Λ_c^+ and $\Xi_c^{0,+}$ states

According to the mass spectrum (see Table 2), $\Lambda_c/\Sigma_c(2765)^+$ can also be regarded as the first radial $(2S)$ excitation of the $\Lambda_c(2286)^+$ with $J^P = 1/2^+$. Interestingly, the results of strong decays in Table 12 do not contradict this assignment. Our calculation indicates that the decay channel $\Sigma_c(2455)\pi$ is a dominant decay channel for the $\Lambda_c^+(2S)$ state. This is in line with the observations by Belle [31,32]. At present, both $1/2^+(2S)$ Λ_c^+ and $1/2^-(1P)$ Σ_c^+ are possible for the assignment of $\Lambda_c/\Sigma_c(2765)^+$. However, there is the very important feature for experiments to distinguish these two assignments in future. Specifically, we suggest to search $\Lambda_c/\Sigma_c(2765)^+$ in the channel of $\Sigma_c(2520)\pi$. As shown in Table 12, the channel $\Sigma_c(2520)\pi$ is large enough to find the $\Lambda_c^+(2S)$ state. On the other hand, this mode seems to be too small to be detected for the $\Sigma_{c1}(1/2^-)$ (see Table 10). Explaining the criteria concretely, we give the following branching ratios for these two states.

For the $\Lambda_c(2S)$ state,

$$\frac{\mathcal{B}(\Lambda_c(2765) \to \Sigma_c(2520)\,\pi)}{\mathcal{B}(\Lambda_c(2765) \to \Sigma_c(2455)\,\pi)} = 0.74. \tag{27}$$

For the $\Sigma_{c1}(1/2^-)$ state,

$$\frac{\mathcal{B}(\Sigma_c(2765) \to \Sigma_c(2520)\,\pi)}{\mathcal{B}(\Sigma_c(2765) \to \Sigma_c(2455)\,\pi)} = 0.03. \tag{28}$$

The branching ratio of $\mathcal{B}(\Sigma_c(2520)\pi)/\mathcal{B}(\Sigma_c(2455)\pi)$ for the $\Sigma_{c1}(1/2^-)$ state is roughly an order of magnitude smaller than $\Lambda_c(2S)$. If $\Lambda_c(2765)$ is the $2S$ excitation, $\Xi_c(2980)$ could be a good candidate as its charm-strange analog [21] as seen in Table 12. The mass difference between $\Lambda_c(2765)$ and $\Xi_c(2980)$ is about 200 MeV, which nearly equals the mass difference between $\Lambda_c(2287)$ and $\Xi_c(2470)$. The predicted width of $\Xi_c(2980)$ is 27.44 MeV which is in good agreement with the experimental data [4,39]. As the $2S$ excitation of

$\Xi_c(2470)$, the branching ratio

$$\frac{\mathcal{B}(\Xi_c(2980) \to \Xi_c'(2580)\,\pi)}{\mathcal{B}(\Xi_c(2980) \to \Xi_c(2645)\,\pi)} = 0.89 \tag{29}$$

is predicted for $\Xi_c(2980)$, which can be tested by future experiments. Recently, the ratio of branching fractions

$$\frac{\mathcal{B}(\Xi_c(2980)^+ \to \Xi_c'(2580)^0\pi^+)}{\mathcal{B}(\Xi_c(2815)^+ \to \Xi_c(2645)^0\pi^+ \to \Xi_c^+\pi^-\pi^+)} \approx 75\% \tag{30}$$

has been estimated by the Belle Collaboration [4]. Combining this with the predicted partial widths of $\Xi_c(2815)$ and $\Xi_c(2645)$ in Tables 8 and 9, the branching fraction $\mathcal{B}(\Xi_c(2980)^+ \to \Xi_c'(2580)^0\pi^+)$ is evaluated to be about 40%, which is consistent with our direct result of 41.8%.

4.5 $2S$ $\Sigma_c^{0,+,++}$ and $\Xi_c'^{0,+}$ states

In Table 3, masses of the $2S$ $\Sigma_c(1/2^+, 3/2^+)$ states are predicted as 2850 and 2876 MeV, respectively. The neutral $\Sigma_c(2850)^0$ found by the BaBar Collaboration in the decay channel $B^- \to \Sigma_c(2850)^0\bar{p} \to \Lambda_c^+\pi^-\bar{p}$ [35] can be regarded as the $2S$ Σ_c state with $J^P = 1/2^+$. The mass and width of the neutral $\Sigma_c(2800)^0$ and $\Sigma_c(2850)^0$ are collected here:

$\Sigma_c(2800)^0$: $m = 2806^{+5}_{-7}$ MeV, $\Gamma = 72^{+22}_{-15}$ MeV;
$\Sigma_c(2850)^0$: $m = 2846 \pm 8 \pm 10$ MeV, $\Gamma = 86^{+33}_{-22}$ MeV.

For lack of experimental information, at present, PDG treated $\Sigma_c(2850)^0$ and $\Sigma_c(2800)^{0,+,++}$ as the same state [1]. As pointed out by the BaBar Collaboration [35], however, there are indications that these two signals detected by Belle [34] and BaBar [35] are two different Σ_c^* states. The main reasons are listed as follows:

1. Although the widths of $\Sigma_c(2800)^{0,+,++}$ and $\Sigma_c(2850)^0$ are consistent with each other, their masses are 3σ apart.

2. The Belle Collaboration tentatively identified the $\Sigma_c(2800)^{0,+,++}$ as the $J = 3/2$ isospin triple, while the BaBar Collaboration found weak evidence of $\Sigma_c(2850)^0$ as a $J = 1/2$ state.

Our results also indicate that $\Sigma_c(2800)^{0,+,++}$ and $\Sigma_c(2850)^0$ are the different Σ_c excited states. One notices that the predicted mass of $1/2^+(2S)$ Σ_c state in this work and in Ref. [9] are around 2850 MeV. Even the results in Refs. [10,56] are only about 50 MeV larger than the measurements. Due to the intrinsic uncertainties of the quark potential model, it is appropriate to assign $\Sigma_c(2850)^0$ as a $2S$ $1/2^+$ state. More importantly, the predicted decay width of $\Sigma_c(1/2^+, 2S)$ state is 118.36 MeV, which is comparable with the measurement by BaBar [35]. The partial width of $\Lambda_c\pi$ is 35.11 MeV, which can explain why $\Sigma_c(2850)^0$ was first found in this channel. We find that the decay modes of $\Sigma_c(2455)\pi$ and $\Sigma_c(2520)\pi$ are also large. Finally, we give the following branching ratios:

$$\frac{\mathcal{B}(\Sigma_c(2850) \to \Sigma_c(2455)\,\pi)}{\mathcal{B}(\Sigma_c(2850) \to \Sigma_c(2520)\,\pi)} = 3.26, \tag{31}$$

$$\frac{\mathcal{B}(\Sigma_c(2850) \to \Sigma_c(2286)\,\pi)}{\mathcal{B}(\Sigma_c(2850) \to \Sigma_c(2455)\,\pi)} = 0.61, \tag{32}$$

which can be tested by future experiments. If $\Sigma_c(2850)^0$ is the $1/2^+(2S)$ state, the mass of its doublet partner in the heavy quark effective theory is predicted as 2876 MeV (denoted $\Sigma_c(2880)$). According to the predicted decay widths in Table 12, this state might also be broad. $\Lambda_c^+\pi$, $\Sigma_c(2455)\pi$, and $\Sigma_c(2520)\pi$ are also dominant for the decay of $\Sigma_c(2880)$. The ratio of $\Gamma(\Sigma_c(2455)\pi)/\Gamma(\Sigma_c(2520)\pi)$ for $\Sigma_c(2880)$ is different from $\Sigma_c(2850)$, whose numerical value is given by

$$\frac{\mathcal{B}(\Sigma_c(2880) \to \Sigma_c(2455)\,\pi)}{\mathcal{B}(\Sigma_c(2880) \to \Sigma_c(2520)\,\pi)} = 0.29. \tag{33}$$

Even though the strange partners of $\Sigma_c(2850)$ and $\Sigma_c(2880)$ have not been found by any experiments, their decay properties are calculated and presented in Table 12. Our results indicate that $\Lambda_c^+ K$, $\Xi_c'(2580)\pi$, and $\Sigma_c(2455)K$ are the dominant decay modes of the $\Xi_c'(3000)$ state with $J^P = 1/2^+$, while $\Lambda_c^+ K$ and $\Xi_c'(2645)\pi$ are those of the $\Xi_c'(3030)$.[4] Besides the masses and decay widths, the following branching ratios may also be valuable for future experiments:

$$\frac{\mathcal{B}(\Xi_c'(3000) \to \Xi_c'(2580)\,\pi)}{\mathcal{B}(\Xi_c'(3000) \to \Xi_c'(2645)\,\pi)} = 3.04, \tag{34}$$

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[4] If $\Xi_c(2980)$ is the first radial excited state of $\Xi_c(2470)$. Then our predicted masses for $2S$ charm-strange baryons may be about 20–30 MeV lower than the experiments. To compensate for this difference, we increase by about 25 MeV for the $2S$ Ξ_c' states in this case.

$$\frac{\mathcal{B}(\Xi_c'(3030) \to \Xi_c'(2580)\,\pi)}{\mathcal{B}(\Xi_c'(3030) \to \Xi_c'(2645)\,\pi)} = 0.27. \tag{35}$$

5 Summary and outlook

In principle, both ρ and λ modes can be excited in a baryon system. For charmed baryons, the excitation energies of the ρ and λ modes are different due to the heavier mass of a c quark. For the ordinary confining potential, such as the linear or harmonic form, the excited energy of the ρ mode is larger than the λ mode [44]. Hence the low-excited charmed baryons may be dominated by the λ mode excitations. Recently, the investigation by Yoshida et al. confirmed this point [75]. Furthermore, they find that the ρ and λ modes are well separated for the charmed and bottom baryons, which means the component of the ρ mode can be ignored for the low-excited charmed baryons. Interestingly, Refs. [9–11] have also shown that the masses of existing charmed baryons can be explained by the λ mode. Hence, our study of strong decays of the low-excited charmed baryons is an important complement to this work [9–11,75].

Up to now, several candidates of the $1P$ and $2S$ charm and charm-strange baryons have been found by experiments, and some of them are still open to debate. To better understand these low-excited charmed baryons, in this paper, we carry out a systematical study of the mass spectra and strong decays for the $1P$ and $2S$ charmed baryon states in the framework of the non-relativistic constituent quark model. The masses have been calculated in the potential model where the charmed baryons are simply treated as a quasi-two-body system in a light quark cluster picture. The strong decays are computed by the EHQ decay formula where the transition factors are determined by the $^3P_0$ model. When calculating the decays, the inner structure of a light quark cluster has also been considered. Except for the unique parameter γ of the QPC model, the parameters in the potential model and in the EHQ decay formula have the same values.

The well-established ground and $1P$ $\mathcal{G}$-type charmed baryons provide a good test to our method. The experimental properties including both masses and widths for these states can be well explained by our results. This success has made us more confident of our predictions for other $1P$ and $2S$ states. Our main conclusions are as follows.

The broad state $\Lambda_c(2765)^+$ (or $\Sigma_c(2765)^+$) which is still ambiguous could be assigned to the $1/2^+(2S)$ Λ_c^+, or the $1/2^-(1P)$ Σ_{c1}^+ state. The branching ratio $\mathcal{B}(\Sigma_c(2455)\pi)/\mathcal{B}(\Sigma_c(2520)\pi)$ is found to be different for these two assignments, which may help us understand the nature of this state.

$\Sigma_c(2800)^{0,+,++}$ observed by the Belle Collaboration in e^+e^- annihilation processes [34] can be regarded as a negative-parity state with $J^P = 3/2^-$, or $5/2^-$, or their overlapping structure. We suggest to measure the

$\mathcal{B}(\Sigma_c(2455)\pi)/\mathcal{B}(\Sigma_c(2520)\pi)$ in the future. Another neutral state, $\Sigma_c(2850)^0$, which was found in the B^- meson decay [35] could be a good candidate for the first radial excited state of $\Sigma_c(2455)$. With the above assignments, the ratios of $\mathcal{B}(\Lambda_c(2287)\pi)/\mathcal{B}(\Sigma_c(2455)\pi)$ shall be very different for $\Sigma_c(2800)^{0,+,++}$ and $\Sigma_c(2850)^0$, i.e., 4.07 for $\Sigma_c(2800)^{0,+,++}$ and 0.61 for $\Sigma_c(2850)^0$. The puzzle of $\Sigma_c(2800)^{0,+,++}$ and $\Sigma_c(2850)^0$ may be disentangled if these branching ratios are measured in the future. In addition, the ratio of branching fractions $\mathcal{B}(\Sigma_c(2455)\pi)/\mathcal{B}(\Sigma_c(2520)\pi)$ for $\Sigma_c(2850)^0$ is predicted to be 3.26.

The analysis of the mass and decay properties supports that $\Xi_c(2980)^{0,+}$ is the $2S$ excitation (the first radial excited state of $\Xi_c(2470)$). The existence of $\Xi_c(2930)^0$ is still in dispute. If it exists, the assignments of $\Xi'_{c2}(3/2^-)$ and $\Xi'_{c2}(5/2^-)$ are possible. In other words, it could be regarded as a strange partner of $\Sigma_c(2800)^{0,+,++}$. Some useful ratios of partial decay widths are also presented for $\Xi_c(2980)^{0,+}$ and $\Xi_c(2930)^0$.

Although both the masses and the strong decays have been explained in the heavy quark–light quark cluster picture for the observed $2S$ and $1P$ candidates, it is not the end of the story of the study of the excited charmed baryon states. Investigation of the ρ mode excited states with higher energies are also important to identify the effective degrees of freedom of charmed baryons. However, this topic needs much laborious work and is beyond the scope of the present work. In addition, the quark model employed here neglects the effect of virtual hadronic loops. In the future, a more reasonable scheme for studying the properties of heavy baryons will be obtained by the unquenched quark model. Another topic which is left as a future task is to calculate the sum rules among the branching fractions of charmed baryons by applying the technique found in Ref. [76].

Appendix A: Transition factor $\mathcal{M}^{j_A,j_B}_{j_C,\ell}(q)$ in the QPC model

In the following, we will show how to obtain the partial wave amplitudes by the $^3P_0$ strong decay model for the decays of excited charmed baryons. As an example, the process $\Sigma_c(2520) \rightarrow \Lambda_c(2280)\pi$ will be constructed and the transition factor for the EHQ formula will be extracted.

As pointed out in Sect. 3, there are two possible decay processes for an excited charmed baryon state (see Fig. 3). The final states of the left figure contain a charmed baryon and a light meson. The right one contains a charmed meson and a light baryon. If a baryon decays via the so-called $^3P_0$ mechanism, a quark–antiquark pair is created from the vacuum and then regroups two outgoing hadrons by a quark rearrangement process. In the non-relativistic limit, the transition operator $\hat{\mathcal{T}}$ of the $^3P_0$ model is given by

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

where $\omega^{(4,5)}$ and $\varphi_0^{(4,5)}$ are the color and flavor wave functions of the $q_4\bar{q}_5$ pair created from the vacuum. Thus, $\omega^{(4,5)} = (R\bar{R} + G\bar{G} + B\bar{B})/\sqrt{3}$ and $\varphi_0^{(4,5)} = (u\bar{u} + d\bar{d} + s\bar{s})/\sqrt{3}$ are color and flavor singlets. The pair is also assumed to carry the quantum number of 0^{++}, suggesting that they are in a $^3P_0$ state. The $\chi_{1,-m}^{(4,5)}$ 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 $q_4\bar{q}_5$. γ is a dimensionless constant which expresses the strength of the quark–antiquark pair created from the vacuum. The value of γ is usually fixed by fitting the well-measured partial decay widths.

When the mock state [77] is adopted to describe the spatial wave function of a meson, the helicity amplitude $\mathcal{M}^{j_A,j_B,j_C}(q)$ can easily be constructed in the LS basis [71]. The mock state for an A meson is

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

As for the left decay process in Fig. 3, the wave function of a B baryon can be constructed in the same way. The wave function of a C meson is

$$\left| C\left(n_C\,^{2S_C+1}L_C^{J_C jc}\left(\vec{P}_C\right)\right)\right\rangle$$
$$\equiv \omega_C^{15}\phi_C^{15} \prod_C \int d^3\vec{k}_1 d^3\vec{k}_5 \delta^3\left(\vec{k}_1 + \vec{k}_5 - \vec{P}_C\right)$$
$$\times \Psi_{n_C}^{L_C lc}\left(\vec{k}_1, \vec{k}_5\right)\left|q_1\left(\vec{k}_1\right) \bar{q}_5\left(\vec{k}_5\right)\right\rangle. \quad \text{(A.3)}$$

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

$$\prod_A = \langle s_1 m_1, s_2 m_2 | s_{12} m_{12} \rangle \langle s_{12} m_{12}, s_3 m_3 | S_A s_A \rangle$$
$$\times \langle L_A l_A, S_A s_A | J_A j_A \rangle,$$
$$\prod_B = \langle s_2 m_2, s_5 m_5 | s_{25} m_{25} \rangle \langle s_{25} m_{25}, s_3 m_3 | S_B s_B \rangle$$
$$\times \langle L_B l_B, S_B s_B | J_B j_B \rangle,$$
$$\prod_C = \langle s_1 m_1, s_4 m_4 | S_C s_C \rangle \langle L_C l_C, S_C s_C | J_C j_C \rangle,$$

respectively.

The helicity amplitude $\mathcal{M}^{j_A, j_B, j_C}(q)$ is defined by

$$\langle BC | \hat{T} | A \rangle = \delta^3 (\vec{P}_A - \vec{P}_B - \vec{P}_C) \mathcal{M}^{j_A, j_B, j_C}(q), \quad (A.4)$$

where q represents the momentum of an outgoing meson in the rest frame of a meson A. For comparison with experiments, one obtains the partial wave amplitudes $\mathcal{M}_{LS}(q)$ via the Jacob–Wick formula [78],

$$\mathcal{M}_{LS}(q) = \frac{\sqrt{2L+1}}{2J_A+1} \sum_{j_B, j_C} \langle L0 J j_A | J_A j_A \rangle$$
$$\times \langle J_B j_B, J_C j_C | J j_A \rangle \mathcal{M}^{j_A, j_B, j_C}(q). \quad (A.5)$$

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

$$\Gamma(A \rightarrow BC) = 2\pi \frac{E_B E_C}{M_A} q \sum_{L,S} |\mathcal{M}_{LS}(q)|^2. \quad (A.6)$$

Finally, the full expression of $\mathcal{M}_{LS}(q)$ in the rest frame of the baryon A is

$$\mathcal{M}_{LS}(q) = -3\gamma \sum_{l_i, m_j} \langle L0; Jj | J_A j_A \rangle \langle J_B j_B; J_C j_C | Jj \rangle$$
$$\times \langle s_1 m_1; s_2 m_2 | s_{dA} m_{12} \rangle \langle s_{dA} m_{12}; s_3 m_3 | S_A s_A \rangle$$
$$\times \langle S_A s_A; L_A l_A | J_A j_A \rangle \langle s_2 m_2; s_5 m_5 | s_{dB} m_{25} \rangle$$
$$\times \langle s_{dB} m_{25}; s_3 m_3 | S_B s_B \rangle \langle S_B s_B; L_B l_B | J_B j_B \rangle$$
$$\times \langle s_1 m_1; s_4 m_4 | S_C s_C \rangle$$
$$\times \langle S_C s_C; L_C l_C | J_C j_C \rangle \langle s_4 m_4; s_5 m_5 | 1 - m \rangle$$
$$\times \langle 1, m; 1, -m | 0, 0 \rangle$$
$$\times \langle \varphi_B^{235} \varphi_C^{14} | \varphi_0^{45} \varphi_A^{123} \rangle \langle \omega_B^{235} \omega_C^{14} | \omega_0^{45} \omega_A^{123} \rangle \int \cdots \int d^3 \vec{k}_1 \cdots$$
$$\times d^3 \vec{k}_5 \delta^3(\vec{k}_1 + \vec{k}_2 + \vec{k}_3) \delta^3(\vec{q} - \vec{k}_1 - \vec{k}_4) \delta^3(\vec{q} + \vec{k}_2 + \vec{k}_3 + \vec{k}_5)$$
$$\times \delta^3(\vec{k}_4 + \vec{k}_5) \Psi_A(\vec{k}_1, \vec{k}_2, \vec{k}_3) \Psi_B^*(\vec{k}_1, \vec{k}_2, \vec{k}_4)$$
$$\times \psi_C^*(\vec{k}_3, \vec{k}_5) \mathcal{Y}_1^m \left(\frac{\vec{k}_3 - \vec{k}_4}{2} \right), \quad (A.7)$$

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. The flavor matrix element $\xi = \langle \varphi_B^{235} \varphi_C^{14} | \varphi_0^{45} \varphi_A^{123} \rangle$ will be presented in the next subsection. To obtain the analytical amplitudes, the SHO wave

functions are employed to describe the spatial wave function of a hadron. In the momentum space, the SHO radial wave function, $\psi_{Lm}^n(\mathbf{q})$, is given by

$$\psi_{Lm}^n(\mathbf{q})$$
$$= \frac{(-1)^n}{\beta^{3/2}} \sqrt{\frac{2(2n-1)!}{\Gamma(n+L+\frac{1}{2})}} \left(\frac{q}{\beta} \right)^L e^{-\frac{q^2}{2\beta^2}} L_{n-1}^{L+1/2} \left(\frac{q^2}{\beta^2} \right) \mathcal{Y}_{Lm}(\mathbf{q}), \quad (A.8)$$

with $\mathbf{q} = (m_i \vec{k}_j - m_j \vec{k}_i)/(m_i + m_j)$ and $\mathcal{Y}_{Lm}(\mathbf{q}) = |\mathbf{q}|^L Y_{Lm}(\Omega_{\mathbf{q}})$. $L_{n-1}^{L+1/2}(q^2/\beta^2)$ is an associated Laguerre polynomial. The values of the SHO wave function scale parameter β have been given in Tables 5 and 6. In the light quark cluster picture, the wave function of a charmed baryon can easily be constructed. Taking the A baryon as an example, the wave functions corresponding to the $1S$, $2S$, and $1P$ states are given as follows, respectively:

$$\Psi_{00}^0 = \frac{3^{3/4}}{\pi^{3/2} \beta_{dA}^{3/2} \beta_A^{3/2}} e^{-\frac{1}{2\beta_{dA}^2} \left(\frac{m_1 \vec{k}_2 - m_2 \vec{k}_1}{m_1 + m_2} \right)^2 - \frac{1}{2\beta_A^2} \left[\frac{(m_1+m_2)\vec{k}_3 - m_Q(\vec{k}_1 + \vec{k}_2)}{m_1 + m_2 + m_Q} \right]^2};$$

$$\Psi_{00}^1 = -\frac{3^{3/4}}{\sqrt{6}\pi^{3/2} \beta_{dA}^{3/2} \beta_A^{3/2}} e^{-\frac{1}{2\beta_{dA}^2} \left(\frac{m_1 \vec{k}_2 - m_2 \vec{k}_1}{m_1 + m_2} \right)^2 - \frac{1}{2\beta_A^2} \left[\frac{(m_1+m_2)\vec{k}_3 - m_Q(\vec{k}_1 + \vec{k}_2)}{m_1 + m_2 + m_Q} \right]^2}$$
$$\times \left\{ 3 - \frac{2}{\beta_A^2} \left[\frac{(m_1+m_2)\vec{k}_3 - m_Q(\vec{k}_1 + \vec{k}_2)}{m_1 + m_2 + m_Q} \right]^2 \right\};$$

$$\Psi_{1m}^0 = \frac{3^{3/4} \times 2\sqrt{2/3}}{\pi \beta_{dA}^{3/2} \beta_A^{5/2}} e^{-\frac{1}{2\beta_{dA}^2} \left(\frac{m_1 \vec{k}_2 - m_2 \vec{k}_1}{m_1 + m_2} \right)^2 - \frac{1}{2\beta_A^2} \left[\frac{(m_1+m_2)\vec{p}_3 - m_Q(\vec{k}_1 + \vec{k}_2)}{m_1 + m_2 + m_Q} \right]^2}$$
$$\times \mathcal{Y}_{1m} \left(\frac{(m_1+m_2)\vec{k}_3 - m_Q(\vec{k}_1 + \vec{k}_2)}{m_1 + m_2 + m_Q} \right).$$

With the help of Eq. (A.7), the transition amplitude can be obtained. In the following, we take the process $\Sigma_c(2520) \rightarrow \Lambda_c(2280)^+ \pi$ as an example. The wave functions of initial and final states are

$$\Psi_A = \frac{3^{3/4}}{\pi^{3/2} \beta_{dA}^{3/2} \beta_A^{3/2}} e^{-\frac{1}{2\beta_{dA}^2} \left(\frac{m_1 \vec{k}_2 - m_2 \vec{k}_1}{m_1 + m_2} \right)^2 - \frac{1}{2\beta_A^2} \left[\frac{(m_1+m_2)\vec{k}_3 - m_3(\vec{k}_1 + \vec{k}_2)}{m_1 + m_2 + m_3} \right]^2};$$

$$\Psi_B = \frac{3^{3/4}}{\pi^{3/2} \beta_{dB}^{3/2} \beta_B^{3/2}} e^{-\frac{1}{2\beta_{dB}^2} \left(\frac{m_5 \vec{k}_2 - m_2 \vec{k}_5}{m_2 + m_5} \right)^2 - \frac{1}{2\beta_B^2} \left[\frac{(m_2+m_5)\vec{k}_3 - m_Q(\vec{p}_2 + \vec{k}_5)}{m_2 + m_3 + m_5} \right]^2};$$

$$\psi_C = -\frac{1}{\pi^{3/4} \beta_C^{3/2}} e^{-\frac{1}{2\beta_C^2} \left(\frac{m_1 \vec{k}_4 - m_4 \vec{k}_1}{m_1 + m_2} \right)^2}.$$

Based on Eq. (A.7), we obtain the amplitude:

$$\mathcal{M}_{1\frac{1}{2}}(q) = -\frac{3g}{8\pi^{5/4} f^{5/2} \lambda^{3/2} \beta_A^{3/2} \beta_{dA}^{3/2} \beta_B^{3/2} \beta_{dB}^{3/2} \beta_C^{3/2}} p e^{-\frac{4fg-g^2}{4f} q^2}, \quad (A.9)$$

where

$$f = \frac{1}{2\beta_{dA}^2} + \frac{1}{2\beta_{dB}^2} + \frac{1}{2\beta_M^2} - \frac{\mu^2}{4\lambda};$$

$$g = \frac{1}{\beta_{dA}^2} + \frac{\varepsilon_3}{\beta_{dB}^2} + \frac{\varepsilon_4}{\beta_M^2} - \frac{\mu\nu}{2\lambda};$$

$$h = \frac{1}{2\beta_{dA}^2} + \frac{\varepsilon_2^2}{2\beta_B^2} + \frac{\varepsilon_3^2}{2\beta_{dB}^2} + \frac{\varepsilon_4^2}{2\beta_M^2} - \frac{\nu^2}{4\lambda};$$

$$\lambda = \frac{1}{2\beta_A^2} + \frac{1}{2\beta_B^2} + \frac{\varepsilon_1^2}{2\beta_{dA}^2} + \frac{\varepsilon_3^2}{2\beta_{dB}^2};$$

$$\mu = \frac{\varepsilon_1}{\beta_{dA}^2} + \frac{\varepsilon_3}{\beta_{dB}^2}; \quad \nu = \frac{\varepsilon_1}{\beta_{dA}^2} + \frac{\varepsilon_3}{\beta_B^2} + \frac{\varepsilon_3^2}{\beta_{dB}^2};$$

$$\varepsilon_1 = \frac{m_1}{m_1 + m_2}; \quad \varepsilon_2 = \frac{m_3}{m_1 + m_3 + m_5};$$

$$\varepsilon_3 = \frac{m_5}{m_2 + m_5}; \quad \varepsilon_4 = \frac{m_4}{m_1 + m_4},$$

and

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

(A.10)

Here, M_A, M_B, and M_C are the masses of hadrons A, B, and C, respectively. Then the $\tilde{\beta}$ in Eq. (20) is given by

$$\tilde{\beta} = 2\sqrt{\frac{f}{4fg - g^2}},$$

(A.11)

where f and g have been defined above. For the decay channel of $\Sigma_c(2520) \rightarrow \Lambda_c(2280)\pi$, the value of $C_{jC,j_A,J_A}^{s_Q,j_B,J_B}$ is -1. Therefore, we obtain

$$\mathcal{M}_{11}^{1,0}(q) = \frac{3g}{8\pi^{5/4} f^{5/2} \lambda^{3/2} \beta_A^{3/2} \beta_{dA}^{3/2} \beta_B^{3/2} \beta_{dB}^{3/2} \beta_C^{3/2}},$$

(A.12)

where a phase space factor $(2\pi E_B E_C / M_A)^{1/2}\gamma$ is omitted. One notices that the unitary rotation between the LS coupling and jj coupling (Eq. (15)) should be performed to reduce the transition factors of $1P$ state with the same J^P. More details for calculating the decay amplitudes of an excited baryon in the $^3P_0$ model can be found in Refs. [24,79].

Appendix B: Flavor factors

Based on the light SU(3) flavor symmetry, the flavor wave functions of charmed and charmed-strange baryons are given by [57]

$$\Lambda_c^+ = \frac{1}{\sqrt{2}}(ud - du)c; \quad \Sigma_c^{++} = uuc;$$

$$\Xi_c^+ = \frac{1}{\sqrt{2}}(us - su)c; \quad \Sigma_c^+ = \frac{1}{\sqrt{2}}(ud + du)c;$$

$$\Xi_c^0 = \frac{1}{\sqrt{2}}(ds - sd)c; \quad \Sigma_c^0 = ddc;$$

$$\Xi_c'^+ = \frac{1}{\sqrt{2}}(us + su)c; \quad \Xi_c'^0 = \frac{1}{\sqrt{2}}(ds + sd)c.$$

As shown in Fig. 3, the final states of an excited charmed baryons may contain a light meson and a low energy charmed baryon or a light baryon and a charmed meson. The flavor wave functions for the final states are collected in the following:

Table 13 The flavor matrix element ξ for different decay channels of the charmed baryons

Initial state	Final states				
Λ_c^+	$\Sigma_c^{++,+,0}\pi^{-,0,+}$	$D^+ n / D^0 p$			
	$\sqrt{1/3}$	$\sqrt{1/3}$			
Σ_c^{++}	$\Sigma_c^{++,+}\pi^{0,+}$	$\Lambda_c^+\pi^+$	$D^+ p$		
	$\sqrt{1/3}$	$\sqrt{1/3}$	$\sqrt{1/6}$		
Σ_c^+	$\Sigma_c^{++,0}\pi^{-,+}$	$\Lambda_c^+\pi^0$	$D^+ n / D^0 p$		
	$\sqrt{1/3}$	$-\sqrt{1/3}$	$\sqrt{1/12}$		
Σ_c^0	$\Sigma_c^+\pi^- / \Sigma_c^0\pi^0$	$\Lambda_c^+\pi^-$	$D^0 n$		
	$\sqrt{1/3}$	$\sqrt{1/3}$	$\sqrt{1/6}$		
$\Xi_c^{(\prime)+}$	$\Lambda_c^+ K^0$	$\Xi_c^{(\prime)0}\pi^+$	$\Xi_c^{(\prime)+}\pi^0$	$\Sigma_c^{++} K^-$	$\Sigma_c^+ K^0$
	$\sqrt{1/6}$	$\sqrt{1/6}$	$\sqrt{1/12}$	$\sqrt{1/3}$	$\sqrt{1/6}$
$\Xi_c^{(\prime)0}$	$\Lambda_c^+ K^-$	$\Xi_c^{(\prime)+}\pi^-$	$\Xi_c^{(\prime)0}\pi^0$	$\Sigma_c^+ K^-$	$\Sigma_c^0 K^0$
	$\sqrt{1/6}$	$\sqrt{1/6}$	$\sqrt{1/12}$	$\sqrt{1/6}$	$\sqrt{1/3}$

$$\pi^+ = u\bar{d}; \quad \pi^- = d\bar{u}; \quad \pi^0 = (u\bar{u} - d\bar{d})/\sqrt{2};$$

$$K^- = \bar{u}s; \quad \bar{K}^0 = \bar{d}s; \quad D^+ = \bar{d}c; \quad D^0 = \bar{u}c;$$

$$p = \frac{1}{\sqrt{2}}(du - ud)u; \quad n = \frac{1}{\sqrt{2}}(du - ud)d;$$

$$\Lambda^0 = \frac{1}{\sqrt{2}}(du - ud)s.$$

With the above flavor wave functions, the flavor matrix elements ξ for different decay processes are presented in Table 13.

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