Strong decays of charmed baryons in heavy hadron chiral perturbation theory: An update

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We first give a brief overview of the charmed baryon spectroscopy and discuss their possible structure and spin-parity assignments in the quark model. With the new Belle measurement of the widths of $\Sigma_c(2455)$ and $\Sigma_c(2520)$ and the recent CDF measurement of the strong decays of $\Lambda_c(2595)$ and $\Lambda_c(2625)$, we give updated coupling constants in heavy hadron chiral perturbation theory. We find $g_2 = 0.565^{+0.011}_{-0.024}$ for P-wave transitions between s-wave and s-wave baryons, and h_2, one of the couplings responsible for S-wave transitions between s-wave and p-wave baryons, is extracted from $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$ to be 0.63 ± 0.07. It is substantially enhanced compared to the old value of order 0.437. With the help from the quark model, two of the couplings h_{10} and h_{11} responsible for D-wave transitions between s-wave and p-wave baryons are determined from $\Sigma_c(2880)$ decays. There is a tension for the coupling h_2 as its value extracted from $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$ will imply $\Xi_c(2790)^0 \to \Xi_c'\pi$ and $\Xi_c(2815)^+ \to \Xi_c^*\pi$ rates slightly above the current limits. It is conceivable that SU(3) flavor symmetry breaking can help account for the discrepancy.

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I. INTRODUCTION

Many new excited charmed baryon states have been discovered by *BABAR*, Belle, CLEO, and LHCb in the past decade. A very rich source of charmed baryons comes both from B decays and from the $e^+e^- \to c\bar{c}$ continuum. Experimentally and theoretically, it is important to identify the quantum numbers of these new states and understand their properties. Since the pseudoscalar mesons involved in the strong decays of charmed baryons are soft, the charmed baryon system offers an excellent ground for testing the ideas and predictions of heavy quark symmetry of the heavy quark and chiral symmetry of the light quarks. The strong decays of charmed baryons are most conveniently described by the heavy hadron chiral perturbation theory (HHChPT) in which heavy quark symmetry and chiral symmetry are incorporated [1,2]. Heavy baryon chiral Lagrangians were first constructed in Ref. [1] for strong decays of s-wave charmed baryons and in Refs. [3,4] for p-wave ones. Previous phenomenological studies of the strong decays of p-wave charmed baryons based on HHChPT can be found in Refs. [3–7].

With the new Belle measurement of the $\Sigma_c(2455)$ and $\Sigma_c(2520)$ widths and the recent collider detector at Fermilab measurement of the strong decays of $\Lambda_c(2595)$ and $\Lambda_c(2625)$, we would like to update the coupling constants appearing in heavy hadron chiral perturbation theory. Indeed, this work is basically the update of Ref. [7]. We begin with the spectroscopy of charmed baryon states and discuss their possible spin-parity quantum numbers and inner structure in Sec. II. Then in Sec. III we consider the strong decays of s-wave and p-wave baryons within the framework of HHChPT and update the relevant coupling constants. Section IV comes to our conclusions.

II. SPECTROSCOPY

Charmed baryon spectroscopy provides an ideal place for studying the dynamics of the light quarks in the environment of a heavy quark. The singly charmed baryon is composed of a charmed quark and two light quarks, which we will often refer to as a diquark. Each light quark is a triplet of the flavor SU(3). Since $\mathbf{3} \times \mathbf{3} = \bar{\mathbf{3}} + \mathbf{6}$, there are two different SU(3) multiplets of charmed baryons: a symmetric sextet $\mathbf{6}$ and an antisymmetric antitriplet $\bar{\mathbf{3}}$. The Λ_c^+, Ξ_c^+, and Ξ_c^0 form a $\bar{\mathbf{3}}$ representation, and they all decay weakly. The Ω_c^0, $\Xi_c'^+$, $\Xi_c'^0$, and $\Sigma_c^{++,+,0}$ form a $\mathbf{6}$ representation; among them, only Ω_c^0 decays weakly. We have followed the Particle Data Group's convention [8] to use a prime to distinguish the Ξ_c in the $\mathbf{6}$ from the one in the $\bar{\mathbf{3}}$.

In the quark model, the orbital angular momentum of the light diquark can be decomposed into $\mathbf{L}_\ell = \mathbf{L}_\rho + \mathbf{L}_\lambda$ (not $L_\ell = L_\rho + L_\lambda$!), where $\mathbf{L}_\rho$ is the orbital angular momentum between the two light quarks and $\mathbf{L}_\lambda$ the orbital angular momentum between the diquark and the charmed quark. The lowest-lying orbitally excited baryon states are the p-wave charmed baryons. Denoting the quantum numbers L_ρ and L_λ as the eigenvalues of $\mathbf{L}_\rho^2$ and $\mathbf{L}_\lambda^2$, respectively, the p-wave heavy baryon can be either in the ($L_\rho = 0$, $L_\lambda = 1$) λ-state or the ($L_\rho = 1$, $L_\lambda = 0$) ρ-state. It is obvious that the orbital λ-state (ρ-state) is symmetric (antisymmetric) under the interchange of two light quarks q_1 and q_2. The total angular momentum of the diquark is $\mathbf{J}_\ell = \mathbf{S}_\ell + \mathbf{L}_\ell$, and

TABLE I. The p-wave charmed baryons denoted by $\mathcal{B}_{cJ_\ell}(J^P)$ and $\tilde{\mathcal{B}}_{cJ_\ell}(J^P)$ where J_ℓ is the total angular momentum of the two light quarks. In the quark model, the orbital λ-states with $L_\lambda = 1$ (ρ-states with $L_\rho = 1$) have even (odd) orbital wave functions under the permutation of the two light quarks. The ρ-states are denoted by a tilde. A prime is used to distinguish between the sextet and antitriplet SU(3) flavor states of the Ξ_c. The explicit quark model wave functions for p-wave charmed baryons can be found in Ref. [4].

State	SU(3)	S_ℓ	$L_\ell(L_\rho, L_\lambda)$	$J_\ell^{P_\ell}$	State	SU(3)	S_ℓ	$L_\ell(L_\rho, L_\lambda)$	$J_\ell^{P_\ell}$
$\Lambda_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\bar{\mathbf{3}}$	0	1 (0, 1)	1^-	$\Sigma_{c0}(\frac{1}{2}^-)$	$\mathbf{6}$	1	1 (0, 1)	0^-
$\tilde{\Lambda}_{c0}(\frac{1}{2}^-)$	$\bar{\mathbf{3}}$	1	1 (1, 0)	0^-	$\Sigma_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\mathbf{6}$	1	1 (0, 1)	1^-
$\tilde{\Lambda}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\bar{\mathbf{3}}$	1	1 (1, 0)	1^-	$\Sigma_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$	$\mathbf{6}$	1	1 (0, 1)	2^-
$\tilde{\Lambda}_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$	$\bar{\mathbf{3}}$	1	1 (1, 0)	2^-	$\tilde{\Sigma}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\mathbf{6}$	0	1 (1, 0)	1^-
$\Xi_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\bar{\mathbf{3}}$	0	1 (0, 1)	1^-	$\Xi'_{c0}(\frac{1}{2}^-)$	$\mathbf{6}$	1	1 (0, 1)	0^-
$\tilde{\Xi}_{c0}(\frac{1}{2}^-)$	$\bar{\mathbf{3}}$	1	1 (1, 0)	0^-	$\Xi'_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\mathbf{6}$	1	1 (0, 1)	1^-
$\tilde{\Xi}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\bar{\mathbf{3}}$	1	1 (1, 0)	1^-	$\Xi'_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$	$\mathbf{6}$	1	1 (0, 1)	2^-
$\tilde{\Xi}_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$	$\bar{\mathbf{3}}$	1	1 (1, 0)	2^-	$\tilde{\Xi}'_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$	$\mathbf{6}$	0	1 (1, 0)	1^-

the total angular momentum of the charmed baryon is $\mathbf{J} = \mathbf{S}_c + \mathbf{J}_\ell$. In the heavy quark limit, the spin of the charmed quark S_c and the total angular momentum of the two light quarks J_ℓ are separately conserved. In the following, we shall use the notation $\mathcal{B}_{cJ_\ell}(J^P)$ ($\tilde{\mathcal{B}}_{cJ_\ell}(J^P)$) to denote the states symmetric (antisymmetric) in the orbital wave functions under the exchange of two light quarks. The lowest-lying orbitally excited baryon states are the p-wave charmed baryons with their quantum numbers listed in Table I.

The next orbitally excited states are the positive-parity excitations with $L_\rho + L_\lambda = 2$. There are multiplets for the first positive-parity excited charmed baryons (e.g. Λ_{c2} and $\hat{\Lambda}_{c2}$) with the symmetric orbital wave function, corresponding to $L_\lambda = 2$, $L_\rho = 0$ and $L_\lambda = 0$, $L_\rho = 2$ (see Table II). They are distinguished by a hat.[1] For the case of $L_\lambda = L_\rho = 1$, the total orbital angular momentum L_ℓ of the diquark is 2, 1, or 0. Since the orbital states are antisymmetric under the interchange of two light quarks, we shall use a tilde to denote the $L_\lambda = L_\rho = 1$ states. Moreover, we shall use the notation $\tilde{\mathcal{B}}^{L_\ell}_{cJ_\ell}(J^P)$ for tilde states in the $\bar{\mathbf{3}}$ as the quantum number L_ℓ is needed to distinguish different states.[2]

The observed mass spectra and decay widths of charmed baryons are summarized in Table III. By now, the $J^P = \frac{1}{2}^+$ and $\frac{1}{2}^-$ $\bar{\mathbf{3}}$ states, (Λ_c^+, Ξ_c^+, Ξ_c^0), ($\Lambda_c(2595)^+$, $\Xi_c(2790)^+$, $\Xi_c(2790)^0$), and ($\Lambda_c(2625)^+$, $\Xi_c(2815)^+$, $\Xi_c(2815)^0$), respectively, and $J^P = \frac{1}{2}^+$ and $\frac{3}{2}^+$ $\mathbf{6}$ states, (Ω_c, Σ_c, Ξ'_c) and (Ω_c^*, Σ_c^*, $\Xi_c'^*$), respectively, are established. Notice that, except for the parity of the lightest Λ_c^+ and the heavier one $\Lambda_c(2880)^+$, none of the other J^P quantum numbers given in Table III has been measured. One has to rely on the quark model to determine the J^P assignments.

In the following we discuss some of the excited charmed baryon states.

A. Λ_c states

There are seven lowest-lying p-wave Λ_c states arising from combining the charmed quark spin S_c with light constituents in the $J_\ell^{P_\ell} = 1^-$ state: $\Lambda_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$, $\tilde{\Lambda}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$, $\tilde{\Lambda}_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$, and one singlet $\tilde{\Lambda}_{c0}(\frac{1}{2}^-)$ (see Table I). $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ form a doublet $\Lambda_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$. The allowed strong decays are $\Lambda_{c1}(1/2^-) \to [\Sigma_c\pi]_S$, $[\Sigma_c^*\pi]_D$ and $\Lambda_{c1}(3/2^-) \to [\Sigma_c\pi]_D$, $[\Sigma_c^*\pi]_{S,D}$, $[\Lambda_c\pi\pi]_P$.

$\Lambda_c(2765)^+$ is a broad state ($\Gamma \approx 50$ MeV) first seen in $\Lambda_c^+\pi^+\pi^-$ by CLEO [12]. It appears to resonate through Σ_c and probably also Σ_c^*. It has a nickname $\Sigma_c(2765)^+$ because whether it is a Λ_c^+ or a Σ_c^+ and whether the width might be due to overlapping states are not known. It could be a first positive-parity excitation of Λ_c. It has also been proposed in the diquark model [13] to be either the first radial ($2S$) excitation of the Λ_c with $J^P = \frac{1}{2}^-$ containing the light scalar diquark or the first orbital excitation ($1P$) of the Σ_c with $J^P = \frac{3}{2}^-$ containing the light axial-vector diquark.

The state $\Lambda_c(2880)^+$ first observed by CLEO [12] in $\Lambda_c^+\pi^+\pi^-$ was also seen by *BABAR* in the D^0p spectrum [14]. Belle has studied the experimental constraint on its spin-parity quantum numbers [15] and found that $J^P = \frac{5}{2}^+$ is favored by the angular analysis of $\Lambda_c(2880)^+ \to \Sigma_c^{0,++}\pi^\pm$ together with the ratio of $\Sigma^*\pi/\Sigma\pi$ measured to be

$$R \equiv \frac{\Gamma(\Lambda_c(2880) \to \Sigma_c^*\pi^\pm)}{\Gamma(\Lambda_c(2880) \to \Sigma_c\pi^\pm)} = (24.1 \pm 6.4^{+1.1}_{-4.5})\%. \qquad (2.1)$$

In the quark model, the candidates for the parity-even spin-$\frac{5}{2}$ state are $\Lambda_{c2}(\frac{5}{2}^+)$, $\hat{\Lambda}_{c2}(\frac{5}{2}^+)$, $\tilde{\Lambda}^1_{c2}(\frac{5}{2}^+)$, $\tilde{\Lambda}^2_{c2}(\frac{5}{2}^+)$, and $\tilde{\Lambda}^2_{c3}(\frac{5}{2}^+)$

[1] In our original paper [7], we did not explicitly distinguish between $L_\lambda = 2$, $L_\rho = 0$ and $L_\lambda = 0$, $L_\rho = 2$ orbital states.

[2] In terms of the old notation in Ref. [7], $\tilde{\Xi}''_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ stands for $\tilde{\Xi}^1_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ and $\tilde{\Xi}'''_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ for $\tilde{\Xi}^2_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$, for example.

TABLE II. The first positive-parity excitations of charmed baryons denoted by $\mathcal{B}_{cJ_\ell}(J^P)$, $\hat{\mathcal{B}}_{cJ_\ell}(J^P)$, and $\tilde{\mathcal{B}}_{cJ_\ell}^{L_\ell}(J^P)$. States with antisymmetric orbital wave functions (i.e. $L_\rho = L_\lambda = 1$) under the interchange of two light quarks are denoted by a tilde. States with the symmetric orbital wave functions $L_\rho = 2$ and $L_\lambda = 0$ are denoted by a hat. A prime is used to distinguish between the sextet and antitriplet SU(3) flavor states of the Ξ_c. For convenience, we drop the superscript L_ℓ for tilde states in the sextet.

State	$\mathrm{SU}(3)_F$	S_ℓ	$L_\ell(L_\rho, L_\lambda)$	$J_\ell^{P_\ell}$	State	$\mathrm{SU}(3)_F$	S_ℓ	$L_\ell(L_\rho, L_\lambda)$	$J_\ell^{P_\ell}$
$\Lambda_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	0	2 (0, 2)	2^+	$\Sigma_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	**6**	1	2 (0, 2)	1^+
$\hat{\Lambda}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	0	2 (2, 0)	2^+	$\Sigma_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	**6**	1	2 (0, 2)	2^+
$\tilde{\Lambda}_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	$\bar{\mathbf{3}}$	1	0 (1, 1)	1^+	$\Sigma_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$	**6**	1	2 (0, 2)	3^+
$\tilde{\Lambda}_{c0}^1(\frac{1}{2}^+)$	$\bar{\mathbf{3}}$	1	1 (1, 1)	0^+	$\hat{\Sigma}_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	**6**	1	2 (2, 0)	1^+
$\tilde{\Lambda}_{c1}^1(\frac{1}{2}^+,\frac{3}{2}^+)$	$\bar{\mathbf{3}}$	1	1 (1, 1)	1^+	$\hat{\Sigma}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	**6**	1	2 (2, 0)	2^+
$\tilde{\Lambda}_{c2}^1(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	1	1 (1, 1)	2^+	$\hat{\Sigma}_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$	**6**	1	2 (2, 0)	3^+
$\tilde{\Lambda}_{c1}^2(\frac{1}{2}^+,\frac{3}{2}^+)$	$\bar{\mathbf{3}}$	1	2 (1, 1)	1^+	$\tilde{\Sigma}_{c0}(\frac{1}{2}^+)$	**6**	0	0 (1, 1)	0^+
$\tilde{\Lambda}_{c2}^2(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	1	2 (1, 1)	2^+	$\tilde{\Sigma}_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	**6**	0	1 (1, 1)	1^+
$\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+,\frac{7}{2}^+)$	$\bar{\mathbf{3}}$	1	2 (1, 1)	3^+	$\tilde{\Sigma}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	**6**	0	2 (1, 1)	2^+
$\Xi_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	0	2 (0, 2)	2^+	$\Xi'_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	**6**	1	2 (0, 2)	1^+
$\hat{\Xi}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	0	2 (2, 0)	2^+	$\Xi'_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	**6**	1	2 (0, 2)	2^+
$\tilde{\Xi}_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	$\bar{\mathbf{3}}$	1	0 (1, 1)	1^+	$\Xi'_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$	**6**	1	2 (0, 2)	3^+
$\tilde{\Xi}_{c0}^1(\frac{1}{2}^+)$	$\bar{\mathbf{3}}$	1	1 (1, 1)	0^+	$\hat{\Xi}'_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	**6**	1	2 (2, 0)	1^+
$\tilde{\Xi}_{c1}^1(\frac{1}{2}^+,\frac{3}{2}^+)$	$\bar{\mathbf{3}}$	1	1 (1, 1)	1^+	$\hat{\Xi}'_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	**6**	1	2 (2, 0)	2^+
$\tilde{\Xi}_{c2}^1(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	1	1 (1, 1)	2^+	$\hat{\Xi}'_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$	**6**	1	2 (2, 0)	3^+
$\tilde{\Xi}_{c1}^2(\frac{1}{2}^+,\frac{3}{2}^+)$	$\bar{\mathbf{3}}$	1	2 (1, 1)	1^+	$\tilde{\Xi}'_{c0}(\frac{1}{2}^+)$	**6**	0	0 (1, 1)	0^+
$\tilde{\Xi}_{c2}^2(\frac{3}{2}^+,\frac{5}{2}^+)$	$\bar{\mathbf{3}}$	1	2 (1, 1)	2^+	$\tilde{\Xi}'_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$	**6**	0	1 (1, 1)	1^+
$\tilde{\Xi}_{c3}^2(\frac{5}{2}^+,\frac{7}{2}^+)$	$\bar{\mathbf{3}}$	1	2 (1, 1)	3^+	$\tilde{\Xi}'_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$	**6**	0	2 (1, 1)	2^+

(see Table II), recalling that the superscript refers to the orbital angular momentum L_ℓ of the diquark. Based on heavy quark symmetry alone, one cannot predict the ratio R for these states except $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$. Its decays to $\Sigma_c^*\pi$, $\Sigma_c\pi$, and $\Lambda_c\pi$ all in F waves. It turns out that [7]

$$\frac{\Gamma(\tilde{\Lambda}_{c3}^2(5/2^+) \to [\Sigma_c^*\pi]_F)}{\Gamma(\tilde{\Lambda}_{c3}^2(5/2^+) \to [\Sigma_c\pi]_F)} = \frac{5}{4}\frac{p_\pi^7(\Lambda_c(2880) \to \Sigma_c^*\pi)}{p_\pi^7(\Lambda_c(2880) \to \Sigma_c\pi)} = \frac{5}{4} \times 0.29 = 0.36, \tag{2.2}$$

where the factor of $5/4$ follows from heavy quark symmetry. Although this deviates from the experimental measurement (2.1) by 1σ, it is a robust prediction.

It is worth mentioning that the Peking group [16] has studied the strong decays of charmed baryons based on the so-called $^3P_0$ recombination model. For the $\Lambda_c(2880)^+$, the Peking group found that all the symmetric states Λ_{c2} and $\hat{\Lambda}_{c2}$ are ruled out as they do not decay into D^0p according to the $^3P_0$ model. Moreover, the predicted ratio R is either too large or too small compared to experiment. Therefore, it appears that the $L = 2$ orbitally excited state $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ dictates the inner structure of $\Lambda_c(2880)^+$.

B. Σ_c states

The highest isotriplet charmed baryons $\Sigma_c(2800)^{++,+,0}$ decaying to $\Lambda_c^+\pi$ were first measured by Belle [17]. The measured widths of order 70 MeV are shown in Table III. The possible quark states are $\Sigma_{c0}(\frac{1}{2}^-)$, $\Sigma_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$, $\tilde{\Sigma}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-)$, and $\Sigma_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$ (see Table I). Obviously, the mass analysis alone is not adequate to fix the quantum numbers J^P of $\Sigma_c(2800)$, and the study of its strong decays is necessary. The states Σ_{c1} and $\tilde{\Sigma}_{c1}$ are ruled out because their decays to $\Lambda_c^+\pi$ are excluded in the heavy quark limit. They decay mainly to the two pion system $\Lambda_c\pi\pi$ in a P-wave. Now the $\Sigma_{c2}(\frac{3}{2}^-,\frac{5}{2}^-)$ baryon decays principally into the $\Lambda_c\pi$ system in a D-wave, while $\Sigma_{c0}(\frac{1}{2}^-)$ decays into $\Lambda_c\pi$ in an S-wave. Since HHChPT implies a very broad Σ_{c0} with width of order 885 MeV (see Sec. III. B below), this p-wave state is also excluded. Therefore, $\Sigma_c(2800)^{++,+,0}$ are likely to be either $\Sigma_{c2}(\frac{3}{2}^-)$ or $\Sigma_{c2}(\frac{5}{2}^-)$ or their mixing. In the quark-diquark model [18],

TABLE III. Mass spectra and widths (in units of MeV) of charmed baryons. Experimental values are taken from the Particle Data Group [8]. For the widths of the $\Sigma_c(2455)^{0/++}$ and $\Sigma_c(2520)^{0/++}$ baryons, we have taken into account the recent Belle measurement [9] for the average. The width of $\Xi_c(2645)^+$ is taken from Ref. [10]. For $\Xi_c(3055)^0$, we quote the preliminary result from Belle [11].

State	J^P	S_ℓ	L_ℓ	$J_\ell^{P_\ell}$	Mass	Width	Decay modes
Λ_c^+	$\frac{1}{2}^+$	0	0	0^+	2286.46 ± 0.14	$\cdots$	weak
$\Lambda_c(2595)^+$	$\frac{1}{2}^-$	0	1	1^-	2592.25 ± 0.28	2.59 ± 0.56	$\Lambda_c\pi\pi$, $\Sigma_c\pi$
$\Lambda_c(2625)^+$	$\frac{3}{2}^-$	0	1	1^-	2628.11 ± 0.19	< 0.97	$\Lambda_c\pi\pi$, $\Sigma_c\pi$
$\Lambda_c(2765)^+$	$?^?$	?	?	?	2766.6 ± 2.4	50	$\Sigma_c\pi$, $\Lambda_c\pi\pi$
$\Lambda_c(2880)^+$	$\frac{5}{2}^+$	?	?	?	2881.53 ± 0.35	5.8 ± 1.1	$\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$, D^0p
$\Lambda_c(2940)^+$	$?^?$	?	?	?	$2939.3^{+1.4}_{-1.5}$	17^{+8}_{-6}	$\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$, D^0p
$\Sigma_c(2455)^{++}$	$\frac{1}{2}^+$	1	0	1^+	2453.98 ± 0.16	$1.94^{+0.08}_{-0.16}$	$\Lambda_c\pi$
$\Sigma_c(2455)^+$	$\frac{1}{2}^+$	1	0	1^+	2452.9 ± 0.4	< 4.6	$\Lambda_c\pi$
$\Sigma_c(2455)^0$	$\frac{1}{2}^+$	1	0	1^+	2453.74 ± 0.16	$1.87^{+0.09}_{-0.17}$	$\Lambda_c\pi$
$\Sigma_c(2520)^{++}$	$\frac{3}{2}^+$	1	0	1^+	2517.9 ± 0.6	$14.8^{+0.3}_{-0.4}$	$\Lambda_c\pi$
$\Sigma_c(2520)^+$	$\frac{3}{2}^+$	1	0	1^+	2517.5 ± 2.3	< 17	$\Lambda_c\pi$
$\Sigma_c(2520)^0$	$\frac{3}{2}^+$	1	0	1^+	2518.8 ± 0.6	$15.3^{+0.3}_{-0.4}$	$\Lambda_c\pi$
$\Sigma_c(2800)^{++}$	$?^?$	?	?	?	2801^{+4}_{-6}	75^{+22}_{-17}	$\Lambda_c\pi$, $\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$
$\Sigma_c(2800)^+$	$?^?$	?	?	?	2792^{+14}_{-5}	62^{+60}_{-40}	$\Lambda_c\pi$, $\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$
$\Sigma_c(2800)^0$	$?^?$	?	?	?	2806^{+5}_{-7}	72^{+22}_{-15}	$\Lambda_c\pi$, $\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$
Ξ_c^+	$\frac{1}{2}^+$	0	0	0^+	$2467.8^{+0.4}_{-0.6}$	$\cdots$	weak
Ξ_c^0	$\frac{1}{2}^+$	0	0	0^+	$2470.88^{+0.34}_{-0.80}$	$\cdots$	weak
$\Xi_c^{\prime +}$	$\frac{1}{2}^+$	1	0	1^+	2575.6 ± 3.1	$\cdots$	$\Xi_c\gamma$
$\Xi_c^{\prime 0}$	$\frac{1}{2}^+$	1	0	1^+	2577.9 ± 2.9	$\cdots$	$\Xi_c\gamma$
$\Xi_c(2645)^+$	$\frac{3}{2}^+$	1	0	1^+	$2645.9^{+0.5}_{-0.6}$	2.6 ± 0.5	$\Xi_c\pi$
$\Xi_c(2645)^0$	$\frac{3}{2}^+$	1	0	1^+	2645.9 ± 0.9	< 5.5	$\Xi_c\pi$
$\Xi_c(2790)^+$	$\frac{1}{2}^-$	0	1	1^-	2789.9 ± 3.2	< 15	$\Xi_c'\pi$
$\Xi_c(2790)^0$	$\frac{1}{2}^-$	0	1	1^-	2791.8 ± 3.3	< 12	$\Xi_c'\pi$
$\Xi_c(2815)^+$	$\frac{3}{2}^-$	0	1	1^-	2816.6 ± 0.9	< 3.5	$\Xi_c^*\pi$, $\Xi_c\pi\pi$, $\Xi_c'\pi$
$\Xi_c(2815)^0$	$\frac{3}{2}^-$	0	1	1^-	2819.6 ± 1.2	< 6.5	$\Xi_c^*\pi$, $\Xi_c\pi\pi$, $\Xi_c'\pi$
$\Xi_c(2930)^0$	$?^?$	?	?	?	2931 ± 6	36 ± 13	$\Lambda_c\bar{K}$
$\Xi_c(2980)^+$	$?^?$	?	?	?	2971.4 ± 3.3	26 ± 7	$\Sigma_c\bar{K}$, $\Lambda_c\bar{K}\pi$, $\Xi_c\pi\pi$
$\Xi_c(2980)^0$	$?^?$	?	?	?	2968.0 ± 2.6	20 ± 7	$\Sigma_c\bar{K}$, $\Lambda_c\bar{K}\pi$, $\Xi_c\pi\pi$
$\Xi_c(3055)^+$	$?^?$	?	?	?	3054.2 ± 1.3	17 ± 13	$\Sigma_c\bar{K}$, $\Lambda_c\bar{K}\pi$, $D\Lambda$
$\Xi_c(3055)^0$	$?^?$	?	?	?	3059.7 ± 0.8	7.4 ± 3.9	$\Sigma_c\bar{K}$, $\Lambda_c\bar{K}\pi$, $D\Lambda$
$\Xi_c(3080)^+$	$?^?$	?	?	?	3077.0 ± 0.4	5.8 ± 1.0	$\Sigma_c\bar{K}$, $\Lambda_c\bar{K}\pi$, $D\Lambda$
$\Xi_c(3080)^0$	$?^?$	?	?	?	3079.9 ± 1.4	5.6 ± 2.2	$\Sigma_c\bar{K}$, $\Lambda_c\bar{K}\pi$, $D\Lambda$
$\Xi_c(3123)^+$	$?^?$	?	?	?	3122.9 ± 1.3	4.4 ± 3.8	$\Sigma_c^*\bar{K}$, $\Lambda_c\bar{K}\pi$
Ω_c^0	$\frac{1}{2}^+$	1	0	1^+	2695.2 ± 1.7	$\cdots$	weak
$\Omega_c(2770)^0$	$\frac{3}{2}^+$	1	0	1^+	2765.9 ± 2.0	$\cdots$	$\Omega_c\gamma$

both of them have very close masses compatible with experiment. However, if we consider the Regge trajectory in the (J, M^2) plane, $\Sigma_c(2800)$ with $J^P = 3/2^-$ fits nicely to the parent Σ_c trajectory (see Fig. 2(a) of Ref. [18]). Hence, we will advocate a $\Sigma_{c2}(3/2^-)$ quark state for $\Sigma_c(2800)$. It is worth mentioning that for light strange baryons the first orbital excitation of the Σ also has the quantum numbers $J^P = 3/2^-$ [8].

TABLE IV. Antitriplet and sextet states of charmed baryons. The spin-parity quantum numbers of $\Xi_c(3080)$ are not yet established. Mass differences $\Delta m_{\Xi_c\Lambda_c} \equiv m_{\Xi_c} - m_{\Lambda_c}$, $\Delta m_{\Xi'_c\Lambda_c} \equiv m_{\Xi'_c} - m_{\Lambda_c}$, and $\Delta m_{\Omega_c\Xi'_c} \equiv m_{\Omega_c} - m_{\Xi'_c}$ are in units of MeV.

	$\mathcal{B}_{cJ_\ell}(J^P)$	States	Mass difference
$\bar{\mathbf{3}}$	$\mathcal{B}_{c0}(\frac{1}{2}^+)$	$\Lambda_c(2287)^+$, $\Xi_c(2470)^+$, $\Xi_c(2470)^0$	$\Delta m_{\Xi_c\Lambda_c} = 183$
	$\mathcal{B}_{c1}(\frac{1}{2}^-)$	$\Lambda_c(2595)^+$, $\Xi_c(2790)^+$, $\Xi_c(2790)^0$	$\Delta m_{\Xi_c\Lambda_c} = 198$
	$\mathcal{B}_{c1}(\frac{3}{2}^-)$	$\Lambda_c(2625)^+$, $\Xi_c(2815)^+$, $\Xi_c(2815)^0$	$\Delta m_{\Xi_c\Lambda_c} = 190$
	$\tilde{\mathcal{B}}^2_{c3}(\frac{5}{2}^+)$	$\Lambda_c(2880)^+$, $\Xi_c(3080)^+$, $\Xi_c(3080)^0$	$\Delta m_{\Xi_c\Lambda_c} = 196$
$\mathbf{6}$	$\mathcal{B}_{c1}(\frac{1}{2}^+)$	$\Omega_c(2695)^0$, $\Xi'_c(2575)^{+,0}$, $\Sigma_c(2455)^{++,+,0}$	$\Delta m_{\Xi'_c\Sigma_c} = 124$, $\Delta m_{\Omega_c\Xi'_c} = 119$
	$\mathcal{B}_{c1}(\frac{3}{2}^+)$	$\Omega_c(2770)^0$, $\Xi'_c(2645)^{+,0}$, $\Sigma_c(2520)^{++,+,0}$	$\Delta m_{\Xi'_c\Sigma_c} = 128$, $\Delta m_{\Omega_c\Xi'_c} = 120$

C. Ξ_c states

There are seven lowest-lying p-wave Ξ_c states in the $\bar{\mathbf{3}}$, $\tilde{\Xi}_{c0}(\frac{1}{2}^-)$, $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, $\tilde{\Xi}_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, and $\tilde{\Xi}_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$, and seven states in the $\mathbf{6}$, $\Xi'_{c0}(\frac{1}{2}^-)$, $\Xi'_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, $\tilde{\Xi}'_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, and $\Xi'_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$. The states $\Xi_c(2790)$ and $\Xi_c(2815)$ form a doublet $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$. Their strong decays are $\Xi_{c1}(1/2^-) \to [\Xi'_c\pi]_S$ and $\Xi_{c1}(3/2^-) \to [\Xi^*_c\pi]_S$ where Ξ^*_c stands for $\Xi_c(2645)$.

The charmed strange baryons $\Xi_c(2980)$ and $\Xi_c(3080)$ that decay into $\Lambda_c^+ K^- \pi^+$ and $\Lambda_c^+ K_S^0 \pi^-$ were first observed by Belle [19] and confirmed by *BABAR* [20]. In the same paper, *BABAR* also claimed evidence of two new resonances $\Xi_c(3055)^+$ and $\Xi_c(3123)^+$. The former was confirmed by Belle, while no signature of the latter was seen [10]. The neutral $\Xi_c(3055)^0$ was observed recently by Belle in ΛD^0 decays [11]. Another state $\Xi_c(2930)^0$ omitted from the Particle Data Group (PDG) summary table has been only seen by *BABAR* in the $\Lambda_c^+ K^-$ mass projection of $B^- \to \Lambda_c^+ \bar{\Lambda}_c^- K^-$ [21].

The charmed baryons $\Xi_c(2980)$, $\Xi_c(3055)$, $\Xi_c(3080)$, and $\Xi_c(3123)$ could be the first positive-parity excitations of the Ξ_c. The study of the Regge phenomenology is very useful for the J^P assignment of charmed baryons [18,22]. Just as the two states $\Lambda_c(2880)(5/2^+)$ and $\Lambda_c(2625)(3/2^-)$ fit nicely the parent Λ_c Regge trajectory in the (J, M^2) plane, $\Xi_c(3080)$ and $\Xi_c(2815)(3/2^-)$ fall into the parent Ξ_c Regge trajectory (see Fig. 3(a) of Ref. [18]). Hence, this suggests that $\Xi_c(3080)$ has $J^P = 5/2^+$. Likewise, $\Xi_c(3055)$ with $3/2^+$ fits to the parent $\Xi_c(2790)(1/2^-)$ Regge trajectory (see Fig. 3(b) of Ref. [18]).

Since the mass difference between the antitriplets Λ_c and Ξ_c for $J^P = \frac{1}{2}^+, \frac{1}{2}^-, \frac{3}{2}^-$ is of order $180 \sim 200$ MeV, it is conceivable that $\Xi_c(2980)$ and $\Xi_c(3080)$ are the counterparts of $\Lambda_c(2765)$ and $\Lambda_c(2880)$, respectively, in the strange charmed baryon sector. As noted in passing, the state $\Lambda_c(2765)^+$ could be a radial excitation $(2S)$ of Λ_c^+, and $\Lambda_c(2880)$ has the quantum numbers $J^P = \frac{5}{2}^+$; it is thus tempting to assign $J^P = \frac{1}{2}^+$ to $\Xi_c(2980)$ with first radial excitation and $\frac{5}{2}^+$ to $\Xi_c(3080)$. From Table IV we see that $\Xi_c(3080)$ and $\Lambda_c(2880)$ form nicely a $J^P = 5/2^+$ antitriplet as the mass difference between $\Xi_c(3080)$ and $\Lambda_c(2880)$ is consistent with that observed in other antitriplets.

In the relativistic quark-diquark model [18], $\Xi_c(2980)$ is a sextet $J^P = \frac{1}{2}^+$ state. According to Table II, possible sextet candidates are $\Xi'_{c1}(\frac{1}{2}^+)$, $\hat{\Xi}'_{c1}(\frac{1}{2}^+)$, $\tilde{\Xi}'_{c0}(\frac{1}{2}^+)$, and $\tilde{\Xi}'_{c1}(\frac{1}{2}^+)$, recalling that a tilde is to denote states with antisymmetric orbital wave functions (i.e. $L_\rho = L_\lambda = 1$) under the interchange of two light quarks and a hat for $L_\rho = 2$ and $L_\lambda = 0$ states. Strong decays of these four states have been studied in Ref. [16] using the $^3P_0$ model. It turns out that $\Gamma(\tilde{\Xi}'_{c0}(\frac{1}{2}^+)) \approx 2.0$ MeV is too small compared to the experimental value of order 25 MeV (see Table III), while $\hat{\Xi}'_{c1}(\frac{1}{2}^+)$ yields too large $\Lambda_c^+ \bar{K}$ and $\Xi_c\pi$ rates. In the $^3P_0$ model, the strong decay of $\Xi'_{c1}(\frac{1}{2}^+)$ to $\Sigma_c\bar{K}$ is largely suppressed relative to $\Lambda_c^+\bar{K}$. This is not favored by experiment as the decay modes $\Lambda_c^+\bar{K}\pi$, $\Sigma_c\bar{K}$, $\Xi_c\pi\pi$, and $\Xi_c(2645)\pi$ of $\Xi_c(2980)$ have been seen, but not $\Lambda_c^+\bar{K}$. $\tilde{\Xi}'_{c1}(\frac{1}{2}^+)$ does not decay to $\Xi_c\pi$ and $\Lambda_c\bar{K}$ and has a width of 28 MeV consistent with experiment. Therefore, the favored candidate for $\Xi_c(2980)$ is $\tilde{\Xi}'_{c1}(\frac{1}{2}^+)$ which has $J_\ell = L_\ell = 1$.

Just as $\Lambda_c(2880)$, $\Xi_c(3080)$ is mostly likely an antitriplet $J^P = \frac{5}{2}^+$ state as noted in passing. The possible quark states are $\Xi_{c2}(\frac{5}{2}^+)$, $\hat{\Xi}_{c2}(\frac{5}{2}^+)$, $\tilde{\Xi}^1_{c2}(\frac{5}{2}^+)$, $\tilde{\Xi}^2_{c2}(\frac{5}{2}^+)$, and $\tilde{\Xi}^2_{c3}(\frac{5}{2}^+)$ (see Table II). Since $\Xi_c(3080)$ is above the $D\Lambda$ threshold, the two-body mode $D\Lambda$ should exist though it has not been searched for in the $D\Lambda$ spectrum. Recall that the neutral $\Xi_c(3055)^0$ was observed recently by Belle in the $D^0\Lambda$ spectrum [11]. According to the $^3P_0$ model, the first four quark states are excluded as they do not decay into $D\Lambda$ [16]. The only possibility left is $\tilde{\Xi}^2_{c3}(\frac{5}{2}^+)$. Although it can decay into $D\Lambda$, the identification of $\tilde{\Xi}^2_{c3}(\frac{5}{2}^+)$ with $\Xi_c(3080)$ encounters two potential difficulties: (i) its width is dominated by $\Xi_c\pi$ and $\Lambda_c^+\bar{K}$ modes which have not been seen experimentally, and (ii) the predicted width of order 47 MeV [16] is too large compared to the measured one of order 5.7 MeV, even though one may argue that the $^3P_0$ model's prediction can be easily off by a factor of $2 \sim 3$ from the experimental measurement due to its inherent uncertainties [16].

D. Ω_c states

Only two ground states have been observed thus far: $1/2^+$ Ω_c^0 and $3/2^+$ $\Omega_c(2770)^0$. The latter was seen by *BABAR* in the electromagnetic decay $\Omega_c(2770) \to \Omega_c\gamma$ [23].

Charmed baryon spectroscopy has been studied extensively in various models. The interested readers are referred to Refs. [24–26] for further references. It appears that the spectroscopy is well described by the model based on the relativitsic heavy quark-light diquark model by Ebert, Faustov, and Galkin (EFG) [18] (see also Ref. [27]). Indeed, the quantum numbers $J^P = \frac{5}{2}^+$ of $\Lambda_c(2880)$ have been correctly predicted in the model based on the diquark idea before the Belle experiment [28]. Moreover, EFG have shown that all available experimental data on heavy baryons fit nicely to the linear Regge trajectories, namely, the trajectories in the (J, M^2) and (n_r, M^2) planes for orbitally and radially excited heavy baryons, respectively,

$$J = \alpha M^2 + \alpha_0, \qquad n_r = \beta M^2 + \beta_0, \tag{2.3}$$

where n_r is the radial excitation quantum number; α, β are the slopes; and α_0, β_0 are intercepts. The Regge trajectories can be plotted for charmed baryons with natural $(P = (-1)^{J-1/2})$ and unnatural $(P = (-1)^{J+1/2})$ parties. The linearity, parallelism, and equidistance of the Regge trajectories were verified. The predictions of the spin-parity quantum numbers of charmed baryons and their masses in Ref. [18] can be regarded as a theoretical benchmark. Specifically, the J^P assignments are given by $\Lambda_c(2765)$: $1/2^+(2S)$; $\Sigma_c(2800)$: $3/2^-(1P)$; $\Xi'_c(2930)$: $1/2^-$, $3/2^-$, $5/2^-(1P)$; $\Xi'_c(2980)$: $1/2^+(2S)$; $\Xi_c(3055)$: $3/2^+(1D)$; $\Xi'_c(3080)$: $5/2^+(1D)$; $\Xi_c(3123)$: $7/2^+(1D)$.

Since the $J^P = 1/2^-$ and $3/2^-$ antitriplets are well established (see Table IV), one may wonder what the counterparts are in the **6**. It turns out that there is no $J^P = \frac{1}{2}^-$ sextet as the $\Sigma_c(2800)$ cannot be assigned with such spin-parity quantum numbers. This should not be a surprise given that the light Σ baryon with $J^P = 1/2^-$ also has not been seen [8]. The next possible sextet is for $J^P = 3/2^-$: $(\Omega_c(3050)^0,\ \Xi'_c(2930)^{+,0},\ \Sigma_c(2800)^{++,+,0})$ where the $\Omega_c(3/2^-)$ is predicted to have a mass 3050 MeV by the quark-diquark model [18]. The mass differences in this sextet, $\Delta m_{\Xi'_c\Sigma_c} = 131$ MeV and $\Delta m_{\Omega_c\Xi'_c} = 119$ MeV, are consistent with that measured in $J^P = 1/2^+$ and $3/2^+$ sextets (cf. Table IV).

On the basis of QCD sum rules, many charmed baryon multiplets classified according to $[\mathbf{6}_F$ (or $\bar{\mathbf{3}}_F$), J_ℓ, S_ℓ, $\rho/\lambda)]$ were studied in Ref. [24] with a focus on the physics of ρ- and λ-mode excitations. Three sextets were proposed in this work: $(\Omega_c(3250),\ \Xi'_c(2980),\ \Sigma_c(2800))$ for $J^P = 1/2^-$, $3/2^-$ and $(\Omega_c(3320), \Xi'_c(3080), \Sigma_c(2890))$ for $J^P = 5/2^-$. Notice that $\Xi'_c(2980)$ and $\Xi'_c(3080)$ were treated as p-wave baryons rather than first positive-parity excitations. The results on the multiplet $[\mathbf{6}_F, 1, 0, \rho]$ led the authors of Ref. [24] to suggest that there are two $\Sigma_c(2800)$, $\Xi'_c(2980)$, and $\Omega_c(3250)$ states with $J^P = 1/2^-$ and $J^P = 3/2^-$. The mass splittings are 14 ± 7, 12 ± 7, and 10 ± 6 MeV, respectively. The predicted mass of $\Omega_c(1/2^-, 3/2^-)$ around 3250 ± 200 MeV is to be compared with 3050 MeV calculated in the quark-diquark model. Using the central value of the predicted masses to label the states in the multiplet $[\mathbf{6}_F, 1, 0, \rho]$ (see Table I of Ref. [24]), one will have

$$\begin{aligned}
J^P = 1/2^-\colon\ &(\Omega_c(3250), \Xi'_c(2960), \Sigma_c(2730)),\\
&\Delta m_{\Xi'_c\Sigma_c} = 230 \pm 234,\\
&\Delta m_{\Omega_c\Xi'_c} = 290 \pm 250,\\
J^P = 3/2^-\colon\ &(\Omega_c(3260), \Xi'_c(2980), \Sigma_c(2750)),\\
&\Delta m_{\Xi'_c\Sigma_c} = 230 \pm 234,\\
&\Delta m_{\Omega_c\Xi'_c} = 280 \pm 242
\end{aligned} \tag{2.4}$$

in units of MeV. Because of the large theoretical uncertainties in masses, it is not clear if the QCD sum-rule calculations are compatible with the mass differences measured in $J^P = 1/2^+$ and $3/2^+$ sextets, namely, $\Delta m_{\Xi'_c\Sigma_c} \approx 125$ MeV and $\Delta m_{\Omega_c\Xi'_c} \approx 120$ MeV. At any rate, it will be interesting to test these two different model predictions for $J^P = 3/2^-$ and $1/2^-$ sextets in the future.

Finally, we would like to remark that in recent years there have been intensive lattice studies of singly, doubly, and triply charmed baryon spectra by many different groups; see e.g. Refs. [29,30] and references therein. However, the current lattice QCD calculations on singly charmed baryons focus mostly on the low-lying $1/2^+$ and $3/2^+$ states. There exist some preliminary lattice results on excited charmed baryon spectroscopy, but the identification with observed charmed baryon states has not been made [30,31]. It will be very interesting if the lattice studies in the future can provide us information on the spin-parity quantum numbers of p-wave and d-wave excited states such as $\Lambda_c(2765), \Sigma_c(2800), \Xi_c(2980), \Xi_c(3055), \ldots$, etc.

III. STRONG DECAYS

As stated in the Introduction, strong decays of charmed baryons involving soft pseudoscalar mesons are most conveniently described by HHChPT. The chiral Lagrangian involves two coupling constants g_1 and g_2 for P-wave transitions between s-wave and s-wave baryons [1], six couplings h_2–h_7 for the S-wave transitions between s-wave and p-wave baryons, and eight couplings h_8–h_{15} for the D-wave transitions between s-wave and p-wave baryons [4]. The general chiral Lagrangian for heavy baryons coupling to the pseudoscalar mesons can be expressed compactly in terms of superfields. We will not write down the relevant Lagrangians here; instead the reader is referred to Eqs. (3.1) and (3.3) of Ref. [4]. Nevertheless, we list some of the partial widths derived from the Lagrangian [4],

$$\begin{aligned}
\Gamma(\Sigma_c^* \to \Sigma_c\pi) &= \frac{g_1^2}{2\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Sigma_c^*}}p_\pi^3,\\
\Gamma(\Sigma_c \to \Lambda_c\pi) &= \frac{g_2^2}{2\pi f_\pi^2}\frac{m_{\Lambda_c}}{m_{\Sigma_c}}p_\pi^3,\\
\Gamma(\Lambda_{c1}(1/2^-) \to \Sigma_c\pi) &= \frac{h_2^2}{2\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Lambda_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\Sigma_{c0}(1/2^-) \to \Lambda_c\pi) &= \frac{h_3^2}{2\pi f_\pi^2}\frac{m_{\Lambda_c}}{m_{\Sigma_{c0}}}E_\pi^2 p_\pi,\\
\Gamma(\Sigma_{c1}(1/2^-) \to \Sigma_c\pi) &= \frac{h_4^2}{4\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Sigma_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\tilde{\Sigma}_{c1}(1/2^-) \to \Sigma_c\pi) &= \frac{h_5^2}{4\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde{\Sigma}_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\tilde{\Xi}_{c0}(1/2^-) \to \Xi_c\pi) &= \frac{h_6^2}{2\pi f_\pi^2}\frac{m_{\Xi_c}}{m_{\tilde{\Xi}_{c0}}}E_\pi^2 p_\pi,\\
\Gamma(\tilde{\Lambda}_{c1}(1/2^-) \to \Sigma_c\pi) &= \frac{h_7^2}{2\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde{\Lambda}_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\Lambda_{c1}(3/2^-) \to \Sigma_c\pi) &= \frac{2h_8^2}{9\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Lambda_{c1}(3/2)}}p_\pi^5,\\
\Gamma(\Sigma_{c1}(3/2^-) \to \Sigma_c^{(*)}\pi) &= \frac{h_9^2}{9\pi f_\pi^2}\frac{m_{\Sigma_c^{(*)}}}{m_{\Sigma_{c1}(3/2)}}p_\pi^5,\\
\Gamma(\Sigma_{c2}(3/2^-) \to \Lambda_c\pi) &= \frac{4h_{10}^2}{15\pi f_\pi^2}\frac{m_{\Lambda_c}}{m_{\Sigma_{c2}}}p_\pi^5,\\
\Gamma(\Sigma_{c2}(3/2^-) \to \Sigma_c^{(*)}\pi) &= \frac{h_{11}^2}{10\pi f_\pi^2}\frac{m_{\Sigma_c^{(*)}}}{m_{\Sigma_{c2}}}p_\pi^5,\\
\Gamma(\Sigma_{c2}(5/2^-) \to \Sigma_c\pi) &= \frac{2h_{11}^2}{45\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Sigma_{c2}}}p_\pi^5,\\
\Gamma(\Sigma_{c2}(5/2^-) \to \Sigma_c^*\pi) &= \frac{7h_{11}^2}{45\pi f_\pi^2}\frac{m_{\Sigma_c^*}}{m_{\Sigma_{c2}}}p_\pi^5,\\
\Gamma(\tilde{\Sigma}_{c1}(3/2^-) \to \Sigma_c\pi) &= \frac{h_{12}^2}{9\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde{\Sigma}_{c1}}}p_\pi^5,\\
\Gamma(\tilde{\Lambda}_{c1}(3/2^-) \to \Sigma_c\pi) &= \frac{4h_{13}^2}{9\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde{\Lambda}_{c1}}}p_\pi^5,\\
\Gamma(\tilde{\Xi}_{c2}(3/2^-) \to \Xi_c\pi) &= \frac{4h_{14}^2}{15\pi f_\pi^2}\frac{m_{\Xi_c}}{m_{\tilde{\Xi}_{c2}}}p_\pi^5,\\
\Gamma(\tilde{\Lambda}_{c2}(3/2^-) \to \Sigma_c\pi) &= \frac{h_{15}^2}{5\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde{\Lambda}_{c2}}}p_\pi^5,
\end{aligned} \tag{3.1}$$

where p_π is the pion's momentum and $f_\pi = 132$ MeV. The dependence on the pion momentum is proportional to p_π, p_π^3, and p_π^5 for S-wave, P-wave and D-wave transitions, respectively. It is obvious that the couplings $g_1, g_2, h_2, \ldots, h_7$ are dimensionless, while $h_8, \ldots, h_{15}$ have canonical dimension E^{-1}.

A. Strong decays of *s*-wave charmed baryons

Since the strong decay $\Sigma_c^* \to \Sigma_c\pi$ is kinematically prohibited, the coupling g_1 cannot be extracted directly from the strong decays of heavy baryons. In the framework of HHChPT, one can use some measurements as input to fix the coupling g_2 which, in turn, can be used to predict the rates of other strong decays. Among the strong decays $\Sigma_c^{(*)} \to \Lambda_c\pi$, $\Sigma_c^{++} \to \Lambda_c^+\pi^+$ is the most well measured. Hence, we shall use this mode to extract the coupling g_2.

Using the 2006 data of $\Gamma(\Sigma_c^{++}) = \Gamma(\Sigma_c^{++} \to \Lambda_c^+\pi^+) = 2.23 \pm 0.30$ MeV [32], we obtain the coupling g_2 to be

$$|g_2|_{2006} = 0.605^{+0.039}_{-0.043}, \tag{3.2}$$

where we have neglected the tiny contributions from electromagnetic decays. The predicted rates of other modes are shown in Table V, for example,

$$\begin{aligned}
\Gamma(\Xi_c^{\prime *+}) &= \Gamma(\Xi_c^{\prime *+} \to \Xi_c^+\pi^0, \Xi_c^0\pi^+)\\
&= \frac{g_2^2}{4\pi f_\pi^2}\left(\frac{1}{2}\frac{m_{\Xi_c^+}}{m_{\Xi_c^{\prime +}}}p_\pi^3 + \frac{m_{\Xi_c^0}}{m_{\Xi_c^{\prime +}}}p_\pi^3\right)\\
&= (2.8 \pm 0.4)\ \text{MeV},\\
\Gamma(\Xi_c^{\prime *0}) &= \Gamma(\Xi_c^{\prime *0} \to \Xi_c^+\pi^-, \Xi_c^0\pi^0)\\
&= \frac{g_2^2}{4\pi f_\pi^2}\left(\frac{m_{\Xi_c^+}}{m_{\Xi_c^{\prime 0}}}p_\pi^3 + \frac{1}{2}\frac{m_{\Xi_c^0}}{m_{\Xi_c^{\prime 0}}}p_\pi^3\right)\\
&= (2.9 \pm 0.4)\ \text{MeV}.
\end{aligned} \tag{3.3}$$

Note that we have neglected the effect of $\Xi_c - \Xi_c'$ mixing in calculations (for recent considerations, see Refs. [33,34]). It is clear from Table V that the agreement between theory and experiment is excellent except the predicted width for $\Sigma_c^{*++} \to \Lambda_c^+\pi^+$ is a bit too large.

Using the new data from 2014 Particle Data Group [8] in conjunction with the new measurements of Σ_c and Σ_c^* widths by Belle [9], we have $\Gamma(\Sigma_c^{++} \to \Lambda_c^+\pi^+) = 1.94^{+0.08}_{-0.16}$ MeV. Therefore, the coupling g_2 is reduced to

$$|g_2|_{2015} = 0.565^{+0.011}_{-0.024}. \tag{3.4}$$

From Table V we see that the agreement between theory and experiment is further improved: The predicted $\Xi_c(2645)^+$ width is consistent with the first new measurement by Belle [10], and the new calculated width for $\Sigma_c^{*++} \to \Lambda_c^+\pi^+$ is now in agreement with

TABLE V. Decay widths (in units of MeV) of s-wave charmed baryons. Data under the label Expt.(2015) are taken from 2014 PDG [8] together with the new measurements of Σ_c, Σ_c^* [9] and $\Xi_c(2645)^+$ widths [10].

Decay	Expt.(2006)	HHChPT(2006)	Expt.(2015)	HHChPT(2015)
$\Sigma_c^{++} \to \Lambda_c^+ \pi^+$	2.23 ± 0.30	input	$1.94^{+0.08}_{-0.16}$	input
$\Sigma_c^{+} \to \Lambda_c^+ \pi^0$	< 4.6	2.6 ± 0.4	< 4.6	$2.3^{+0.1}_{-0.2}$
$\Sigma_c^{0} \to \Lambda_c^+ \pi^-$	2.2 ± 0.4	2.2 ± 0.3	$1.9^{+0.1}_{-0.2}$	$1.9^{+0.1}_{-0.2}$
$\Sigma_c(2520)^{++} \to \Lambda_c^+ \pi^+$	14.9 ± 1.9	16.7 ± 2.3	$14.8^{+0.3}_{-0.4}$	$14.5^{+0.5}_{-0.8}$
$\Sigma_c(2520)^{+} \to \Lambda_c^+ \pi^0$	< 17	17.4 ± 2.3	< 17	$15.2^{+0.6}_{-1.3}$
$\Sigma_c(2520)^{0} \to \Lambda_c^+ \pi^-$	16.1 ± 2.1	16.6 ± 2.2	$15.3^{+0.4}_{-0.5}$	$14.7^{+0.6}_{-1.2}$
$\Xi_c(2645)^{+} \to \Xi_c^{0,+} \pi^{+,0}$	< 3.1	2.8 ± 0.4	2.6 ± 0.5	$2.4^{+0.1}_{-0.2}$
$\Xi_c(2645)^{0} \to \Xi_c^{+,0} \pi^{-,0}$	< 5.5	2.9 ± 0.4	< 5.5	$2.5^{+0.1}_{-0.2}$

experiment. It is also clear that the Σ_c width is smaller than that of Σ_c^* by a factor of ~ 7, although they will become the same in the limit of heavy quark symmetry. This is ascribed to the fact that the pion's momentum is around 90 MeV in the decay $\Sigma_c \to \Lambda_c \pi$ while it is two times bigger in $\Sigma_c^* \to \Lambda_c \pi$.

It is worth remarking that, although the coupling g_1 cannot be determined directly from the strong decay such as $\Sigma_c^* \to \Sigma_c \pi$, some information of g_1 can be learned from the radiative decay $\Xi_c^{\prime *0} \to \Xi_c^0 \gamma$, which is prohibited at tree level by SU(3) symmetry but can be induced by chiral loops. A measurement of $\Gamma(\Xi_c^{\prime *0} \to \Xi_c^0 \gamma)$ will yield two possible solutions for g_1. As pointed out in Ref. [1], within the framework of the nonrelativistic quark model, the couplings g_1 and g_2 can be related to g_A^q, the axial-vector coupling in a single quark transition of $u \to d$, via

$$g_1 = \frac{4}{3} g_A^q, \qquad g_2 = \sqrt{\frac{2}{3}} g_A^q. \tag{3.5}$$

Assuming the validity of the quark model relations among different coupling constants, the experimental value of g_2 implies $|g_1| = 0.93 \pm 0.16$ [35].

The couplings g_1 and g_2 have been evaluated using lattice QCD with the results [36][3]

$$g_1 = 0.56 \pm 0.13, \qquad g_2 = 0.41 \pm 0.08. \tag{3.6}$$

Hence, the quark model values of g_1 and g_2 are significantly larger than the above lattice QCD results. This is ascribed to the fact that $1/m_Q$ corrections to strong decays have been taken into account in lattice calculations [36]. For example, the $1/m_c$ effect on the amplitude of $\Sigma_c^{(*)} \to \Lambda_c \pi$ is about 40%. As a consequence, the lattice values of g_1 and g_2 are significantly smaller than the quark model results.

[3]Our definitions of g_1 and g_2 are related to that of Detmold, Lin, and Meinel [36] by the relations $g_1 = (2/3) g_2^{\rm DLM}$ and $g_2 = g_3^{\rm DLM}/\sqrt{3}$.

B. Strong decays of *p*-wave charmed baryons

As noted in passing, six couplings h_2–h_7 are needed to describe the S-wave transitions between s-wave and p-wave baryons, and eight couplings h_8–h_{15} are needed for the D-wave transitions between s-wave and p-wave baryons [4]. Since $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ form a doublet $\Lambda_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, the couplings h_2 and h_8 in principle can be extracted from $\Lambda_c(2595) \to \Sigma_c \pi$ and $\Lambda_c(2625) \to \Sigma_c \pi$, respectively. However, this method is not practical as only the decay $\Lambda_c(2595)^+ \to \Sigma^+ \pi^0$ is kinematically (barely) allowed (see the discussions below), while the $\Lambda_c(2625)$ decay to $\Sigma_c \pi$ via a D-wave transition is kinematically suppressed.

Likewise, the information on the couplings h_{10} and h_{11} can be inferred from the strong decays of $\Sigma_c(2800)$ identified with $\Sigma_{c2}(3/2^-)$. Couplings other than h_2, h_8, and h_{10} can be related to each other via the quark model. The S-wave couplings between the s-wave and the p-wave baryons are related by [4]

$$\frac{|h_3|}{|h_4|} = \frac{\sqrt{3}}{2}, \qquad \frac{|h_2|}{|h_4|} = \frac{1}{2}, \qquad \frac{|h_5|}{|h_6|} = \frac{2}{\sqrt{3}}, \qquad \frac{|h_5|}{|h_7|} = 1. \tag{3.7}$$

The D-wave couplings satisfy the relations

$$|h_8| = |h_9| = |h_{10}|, \qquad \frac{|h_{11}|}{|h_{10}|} = \frac{|h_{15}|}{|h_{14}|} = \sqrt{2},$$
$$\frac{|h_{12}|}{|h_{13}|} = 2, \qquad \frac{|h_{14}|}{|h_{13}|} = 1. \tag{3.8}$$

The reader is referred to Ref. [4] for further details.

TABLE VI. Same as Table V except for p-wave charmed baryons. The results under the label HHChPT(2015) are obtained using $g_2 = 0.565$, $h_2 = 0.63 \pm 0.07$ and $h_8 = (0.85^{+0.11}_{-0.08}) \times 10^{-3}\ \mathrm{MeV}^{-1}$.

Decay	Expt.(2006)	HHChPT(2006)	Expt.(2015)	HHChPT(2015)
$\Lambda_c(2595)^+ \to (\Lambda_c^+\pi\pi)_R$	$2.63^{+1.56}_{-1.09}$	Input	2.59 ± 0.56	Input
$\Lambda_c(2595)^+ \to \Sigma_c^{++}\pi^-$	$0.65^{+0.41}_{-0.31}$	$0.72^{+0.43}_{-0.30}$	$\cdots$	$\cdots$
$\Lambda_c(2595)^+ \to \Sigma_c^0\pi^+$	$0.67^{+0.41}_{-0.31}$	$0.77^{+0.46}_{-0.32}$	$\cdots$	$\cdots$
$\Lambda_c(2595)^+ \to \Sigma_c^+\pi^0$	$\cdots$	$1.57^{+0.93}_{-0.65}$	$\cdots$	$2.38^{+0.56}_{-0.50}$
$\Lambda_c(2625)^+ \to \Sigma_c^{++}\pi^-$	< 0.58	0.029	< 0.30	0.028
$\Lambda_c(2625)^+ \to \Sigma_c^0\pi^+$	< 0.60	0.029	< 0.30	0.029
$\Lambda_c(2625)^+ \to \Sigma_c^+\pi^0$	$\cdots$	0.041	$\cdots$	0.040
$\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi$	< 1.9	0.21	< 0.97	0.32
$\Sigma_c(2800)^{++} \to \Lambda_c\pi,\ \Sigma_c^{(*)}\pi$	75^{+22}_{-17}	Input	75^{+22}_{-17}	Input
$\Sigma_c(2800)^{+} \to \Lambda_c\pi,\ \Sigma_c^{(*)}\pi$	62^{+60}_{-40}	Input	62^{+60}_{-40}	Input
$\Sigma_c(2800)^{0} \to \Lambda_c\pi,\ \Sigma_c^{(*)}\pi$	61^{+28}_{-18}	Input	72^{+22}_{-15}	Input
$\Xi_c(2790)^+ \to \Xi_c'^{0,+}\pi^{+,0}$	< 15	$8.0^{+4.7}_{-3.3}$	< 15	$16.7^{+3.6}_{-3.6}$
$\Xi_c(2790)^0 \to \Xi_c'^{+,0}\pi^{-,0}$	< 12	$8.5^{+5.0}_{-3.5}$	< 12	$17.7^{+2.9}_{-3.8}$
$\Xi_c(2815)^+ \to \Xi_c^{*+,0}\pi^{0,+}$	< 3.5	$3.4^{+2.0}_{-1.4}$	< 3.5	$7.1^{+1.5}_{-1.5}$
$\Xi_c(2815)^0 \to \Xi_c^{*+,0}\pi^{-,0}$	< 6.5	$3.6^{+2.1}_{-1.5}$	< 6.5	$7.7^{+1.7}_{-1.7}$

Although the coupling h_2 can be inferred from the two-body decay $\Lambda_c(2595) \to \Sigma_c\pi$, this method is less accurate because the decay is kinematically barely allowed or even prohibited depending on the mass of $\Lambda_c(2595)^+$. Since $m(\Sigma_c^{++}) + m(\pi^-) = 2593.55$ MeV, $m(\Sigma_c^+) + m(\pi^0) = 2587.88$ MeV, and $m(\Sigma_c^0) + m(\pi^+) = 2593.31$ MeV, it is obvious that the decays $\Lambda_c(2595)^+ \to \Sigma_c^{++}\pi^-, \Sigma_c^0\pi^+$ and $\Lambda_c(2595)^+ \to \Sigma^+\pi^0$ are kinematically barely allowed for $m(\Lambda_c(2595)) = 2595.4$ MeV, while only the last mode is allowed for $m(\Lambda_c(2595)) = 2592.25$ MeV. Moreover, the finite width effect of the intermediate resonant states could become important [6].

We next turn to the three-body decays $\Lambda_c^+\pi\pi$ of $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ to extract h_2 and h_8. Since the three-body decay of the latter proceeds in a P-wave, it is expected to be suppressed. Using the measured ratios of $\Gamma(\Lambda_c(2595)^+ \to \Sigma_c^{++}\pi^-)$ and $\Gamma(\Lambda_c(2595)^+ \to \Sigma_c^0\pi^+)$ relative to $\Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi^+\pi^-)$, assuming $\Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi^0\pi^0) \approx \Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi^+\pi^-)$ (see the discussion below for the justification) and using the 2006 data from PDG [32] for $\Gamma(\Lambda_c(2595)) = 3.6^{+2.0}_{-1.3}$ MeV and the mass $m(\Lambda_c(2595)) = 2595.4 \pm 0.6$ MeV, we have obtained the experimental resonant rate [7]

$$\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi\pi)_R = (2.63^{+1.56}_{-1.09})\ \mathrm{MeV} \tag{3.9}$$

as shown in Table VI.

Assuming the pole contributions to $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$ due to the intermediate states Σ_c and Σ_c^*, the resonant rate for the process $\Lambda_{c_1}^+(2595) \to \Lambda_c^+\pi^+\pi^-$ can be calculated in the framework of HHChPT to be [4]

$$\frac{d^2\Gamma(\Lambda_{c1}^+(2595) \to \Lambda_c^+\pi^+(E_1)\pi^-(E_2))}{dE_1 dE_2} = \frac{g_2^2}{16\pi^3 f_\pi^4} m_{\Lambda_c^+}\{\mathbf{p}_2^2|A|^2 + \mathbf{p}_1^2|B|^2 + 2\mathbf{p}_1\cdot\mathbf{p}_2\mathrm{Re}(AB^*)\}, \tag{3.10}$$

with

$$A(E_1,E_2) = \frac{h_2 E_1}{\Delta_R - \Delta_{\Sigma_c^0} - E_1 + i\Gamma_{\Sigma_c^0}/2} - \frac{\frac{2}{3}h_8\mathbf{p}_2^2}{\Delta_R - \Delta_{\Sigma_c^{*0}} - E_1 + i\Gamma_{\Sigma_c^{*0}}/2} + \frac{2h_8\mathbf{p}_1\cdot\mathbf{p}_2}{\Delta_R - \Delta_{\Sigma_c^{*++}} - E_2 + i\Gamma_{\Sigma_c^{*++}}/2}, \tag{3.11}$$

$$B(E_1,E_2;\Delta_{\Sigma_c^{(*)0}},\Delta_{\Sigma_c^{(*)++}}) = A(E_2,E_1;\Delta_{\Sigma_c^{(*)++}},\Delta_{\Sigma_c^{(*)0}}), \tag{3.12}$$

where $\Delta_R = m_{\Lambda_c(2593)} - m_{\Lambda_c}$ and $\Delta_{\Sigma_c^{(*)}} = m_{\Sigma_c^{(*)}} - m_{\Lambda_c}$. Likewise, a similar relation can be derived for $\Lambda_{c_1}^+(2625) \to \Lambda_c^+\pi^+\pi^-$ (see Ref. [4]). Numerically, we found [7]

$$\Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi)_R = 13.82h_2^2 + 26.28h_8^2 - 2.97h_2h_8,$$
$$\Gamma(\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi)_R = 0.617h_2^2 + 0.136\times 10^6 h_8^2 - 27h_2h_8, \qquad (3.13)$$

where $\Lambda_c^+\pi\pi = \Lambda_c^+\pi^+\pi^- + \Lambda_c^+\pi^0\pi^0$. It is clear that the limit on $\Gamma(\Lambda_c(2625))$ gives an upper bound on h_8 of order 10^{-3} (in units of MeV^{-1}), whereas the decay width of $\Lambda_c(2595)$ is entirely governed by the coupling h_2

$$|h_2|_{2006} = 0.437^{+0.114}_{-0.102}, \qquad |h_8|_{2006} < 3.65\times 10^{-3}\ \mathrm{MeV}^{-1}. \qquad (3.14)$$

It was pointed out in Ref. [6] that the proximity of the $\Lambda_c(2595)^+$ mass to the sum of the masses of its decay products will lead to an important threshold effect which will lower the $\Lambda_c(2595)^+$ mass by 2–3 MeV than the one observed. A more sophisticated treatment of the mass line shape of $\Lambda_c(2595)^+ \to \Lambda_c^+\pi^+\pi^-$ by CDF with data sample 25 times larger than previous measurements yields $m(\Lambda_c(2595)) = 2592.25 \pm 0.28$ MeV [37], which is 3.1 MeV smaller than the 2006 world average. Therefore, strong decays of $\Lambda_c(2595)$ into $\Lambda_c\pi\pi$ are very close to the threshold as $m_{\Lambda_c(2595)} - m_{\Lambda_c} = 305.79 \pm 0.24$ MeV [8]. Hence, its phase space is very sensitive to the small isospin-violating mass differences between members of pions and charmed Sigma baryon multiplets.

For $m(\Lambda_c(2595)) = 2592.25 \pm 0.28$ MeV [37], we obtain (in units of MeV)

$$\Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi)_R = g_2^2(20.45h_2^2 + 43.92h_8^2 - 8.95h_2h_8),$$
$$\Gamma(\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi)_R = g_2^2(1.78h_2^2 + 4.557\times 10^6 h_8^2 - 79.75h_2h_8). \qquad (3.15)$$

By performing a fit to the measured $M(pK^-\pi^+\pi^+) - M(pK^-\pi^+)$ and $M(pK^-\pi^+\pi^+\pi^-) - M(pK^-\pi^+)$ mass difference distributions and using $g_2^2 = 0.365$, CDF found $h_2^2 = 0.36 \pm 0.08$ or $|h_2| = 0.60 \pm 0.07$ [37]. This corresponds to a decay width $\Gamma(\Lambda_c(2595)^+) = 2.59 \pm 0.30 \pm 0.47$ MeV [37].[4] Note that the decay width of $\Lambda_c(2595)^+$ measured by CDF is the quantity $\Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi)_R$ instead of the natural width associated with a Breit–Wigner curve. For the width of $\Lambda_c(2625)^+$, CDF observed a value consistent with zero and therefore calculated an upper limit 0.97 MeV using a Bayesian approach. According to the PDG, $\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi$ is dominated by direct nonresonant contributions [8].

From the CDF measurements $\Gamma(\Lambda_c(2595)^+) = 2.59 \pm 0.56$ MeV and $\Gamma(\Lambda_c(2625)^+) < 0.97$ MeV, we obtain

$$|h_2|_{2015} = 0.63 \pm 0.07, \qquad |h_8|_{2015} < 2.32\times 10^{-3}\ \mathrm{MeV}^{-1}. \qquad (3.16)$$

Hence, the magnitude of the coupling h_2 is greatly enhanced from 0.437 to 0.63. Our h_2 is slightly different from the value of 0.60 obtained by CDF. This is because CDF used $|g_2| = 0.604$ to calculate the mass dependence of $\Gamma(\Lambda_c^+\pi\pi)$, while we used $|g_2| = 0.565$. Since $\Gamma(\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi)$ is basically proportional to $g_2^2h_2^2$, a smaller g_2 will lead to a larger h_2.

The reader may wonder why the coupling h_2 obtained in 2006 and 2015 is so different even though the resonant rate of $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$ used in 2006 and 2015 is very similar in its central value. This is ascribed to the fact that the mass of $\Lambda_c(2595)^+$ is 3.1 MeV lower than the previous world average due to the threshold effect. To illustrate this, following Ref. [37] we consider the dependence of $\Gamma(\Lambda_c^+\pi^+\pi^-)/h_2^2$ and $\Gamma(\Lambda_c^+\pi^0\pi^0)/h_2^2$ on $\Delta M(\Lambda_c(2595)) \equiv M(\Lambda_c(2595)^+) - M(\Lambda_c^+)$ as depicted in Fig. 1. For $\Delta M(\Lambda_c(2595)) = 308.9$ MeV, we see that $\Gamma(\Lambda_c^+\pi^0\pi^0) \approx \Gamma(\Lambda_c^+\pi^+\pi^-)$, while $\Gamma(\Lambda_c^+\pi^0\pi^0) \approx 4.5\Gamma(\Lambda_c^+\pi^+\pi^-)$ for $\Delta M(\Lambda_c(2595)) = 305.79$ MeV. Due to the threshold effect, the isospin relation $\Gamma(\Lambda_c^+\pi^0\pi^0) = \frac{1}{2}\Gamma(\Lambda_c^+\pi^+\pi^-)$ as advocated by the PDG is strongly violated in $\Lambda_c(2595) \to \Lambda_c\pi\pi$ decays, though it is still valid in $\Lambda_c(2625) \to \Lambda_c\pi\pi$ decays. It is evident from Fig. 1 that $\Gamma(\Lambda_c^+\pi\pi)/h_2^2$ at

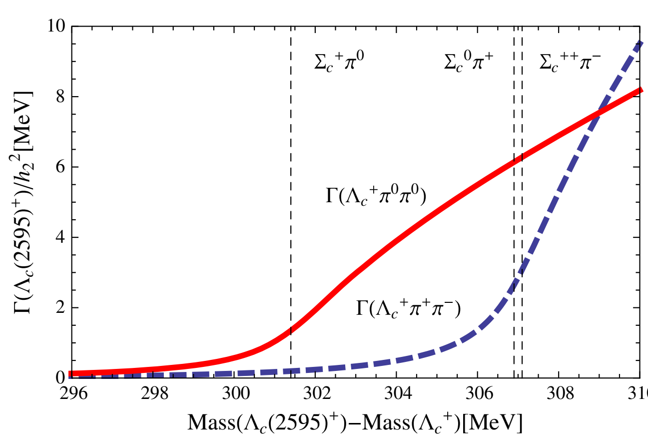

FIG. 1 (color online). Calculated dependence of $\Gamma(\Lambda_c^+\pi^0\pi^0)/h_2^2$ (full curve) and $\Gamma(\Lambda_c^+\pi^+\pi^-)/h_2^2$ (dashed curve) on $m(\Lambda_c(2595)^+) - m(\Lambda_c^+)$, where we have used the parameters $g_2 = 0.565$, $h_2 = 0.63$, and $h_8 = 0.85\times 10^{-3}\ \mathrm{MeV}^{-1}$.

[4]Note that the contributions from the h_8 terms to $\Gamma(\{\Lambda_c(2595)^+, \Lambda_c(2625)^+\} \to \Lambda_c^+\pi\pi)$ have been ignored in the CDF fit to the data [37]. Using $g_2^2 = 0.365$ and $h_2^2 = 0.36 \pm 0.08$, we obtain $\Gamma(\Lambda_c(2595)^+) \approx 2.68$ MeV from Eq. (3.15), which is slightly larger than the CDF value of 2.59 ± 0.56 MeV. This is mainly because we have used the updated widths for Σ_c and Σ_c^*.

$\Delta M(\Lambda_c(2595)) = 305.79$ MeV is smaller than that at 308.9 MeV. This explains why h_2 should become larger when $\Delta M(\Lambda_c(2595))$ becomes smaller.

The $\Xi_c(2790)$ and $\Xi_c(2815)$ baryons form a doublet $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$. $\Xi_c(2790)$ decays to $\Xi'_c\pi$, while $\Xi_c(2815)$ decays to $\Xi_c\pi\pi$, resonating through Ξ_c^*, i.e. $\Xi_c(2645)$. Using the coupling h_2 obtained from (3.16) and assuming SU(3) flavor symmetry, the predicted $\Xi_c(2790)$ and $\Xi_c(2815)$ widths are shown in Table VI, where uses have been made of

$$\begin{aligned}
&\Gamma(\Xi_{c1}(1/2^-)^+) \\
&\quad\approx \Gamma(\Xi_{c1}(1/2^-)^+ \to \Xi_c'^+\pi^0, \Xi_c'^0\pi^+) \\
&\quad= \frac{h_2^2}{4\pi f_\pi^2}\left(\frac{1}{2}\frac{m_{\Xi_c'^+}}{m_{\Xi_{c1}(1/2)}}E_\pi^2 p_\pi + \frac{m_{\Xi_c'^0}}{m_{\Xi_{c1}(1/2)}}E_\pi^2 p_\pi\right), \\
&\Gamma(\Xi_{c1}(3/2^-)^+) \\
&\quad\approx \Gamma(\Xi_{c1}(3/2^-)^+ \to \Xi_c^{*+}\pi^0, \Xi_c^{*0}\pi^+) \\
&\quad= \frac{h_2^2}{4\pi f_\pi^2}\left(\frac{1}{2}\frac{m_{\Xi_c^{*+}}}{m_{\Xi_{c1}(3/2)}}E_\pi^2 p_\pi + \frac{m_{\Xi_c^{*0}}}{m_{\Xi_{c1}(3/2)}}E_\pi^2 p_\pi\right),
\end{aligned} \tag{3.17}$$

and similar expressions for the neutral $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ states based on the experimental observation that the $\Xi_c\pi\pi$ mode in $\Xi_c(2815)$ decays is consistent with being entirely via $\Xi_c^*\pi$ [38]. It is evident that the predicted two-body decay rates of $\Xi_c(2790)^0$ and $\Xi_c(2815)^+$ exceed the current experimental limits because of the enhancement of h_2 (see Table VI). Hence, there is a tension for the coupling h_2 as its value extracted from $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$ will imply $\Xi_c(2790)^0 \to \Xi'_c\pi$ and $\Xi_c(2815)^+ \to \Xi_c^*\pi$ rates slightly above current limits. It is conceivable that SU(3) flavor symmetry breaking can help account for the discrepancy. For example, if h_2 is allowed to have 25% uncertainties due to SU(3) breaking between Ξ_{c1} and Λ_{c1} decays, one will have $\Gamma(\Xi_c(2790)^0) \approx 9.9$ MeV and $\Gamma(\Xi_c(2815)^+) \approx 4.0$ MeV for $h_2 = 0.47$. It is clear that the former is consistent with the measured limit, while the discrepancy between theory and experiment for the latter is much improved. In HHChPT, SU(3) breaking effects arise from chiral loops due to the light quark masses. Applications to the strong decays of heavy baryons have been considered in Ref. [39]. We plan to pursue this issue in the future.

Some information on the coupling h_{10} can be inferred from the strong decays of $\Sigma_c(2800)$. As noticed in passing, the states $\Sigma_c(2800)^{++,+,0}$ are most likely to be $\Sigma_{c2}(\frac{3}{2}^-)$. From Eq. (3.1) and the quark model relation $|h_3| = \sqrt{3}|h_2|$ from Eq. (3.7), we obtain, for example, $\Gamma(\Sigma_{c0}^{++} \to \Lambda_c^+\pi^+) \approx 885$ MeV. Hence, $\Sigma_c(2800)$ cannot be identified with $\Sigma_{c0}(1/2^-)$. Assuming their widths are dominated by the two-body D-wave modes $\Lambda_c\pi$, $\Sigma_c\pi$, and $\Sigma_c^*\pi$, we have [4]

$$\begin{aligned}
&\Gamma(\Sigma_{c2}(3/2)^{++}) \\
&\approx \Gamma(\Sigma_{c2}(3/2)^{++} \to \Lambda_c^+\pi^+) \\
&\quad+ \Gamma(\Sigma_{c2}(3/2)^{++} \to \Sigma_c^+\pi^+) + \Gamma(\Sigma_{c2}(3/2)^{++} \to \Sigma_c^{*+}\pi^+) \\
&\quad+ \Gamma(\Sigma_{c2}(3/2)^{++} \to \Sigma_c^{++}\pi^0) + \Gamma(\Sigma_{c2}(3/2)^{++} \to \Sigma_c^{*++}\pi^0)
\end{aligned} \tag{3.18}$$

and similar expressions for $\Sigma_{c2}(\frac{3}{2})^+$ and $\Sigma_{c2}(\frac{3}{2})^0$. The first term is governed by the h_{10}^2 coupling and the rest by h_{11}^2. Using Eq. (3.1), the quark model relation $h_{11}^2 = 2h_{10}^2$ [cf. Eq. (3.8)] and the measured widths of $\Sigma_c(2800)^{++,+,0}$ (Table III), we obtain

$$|h_{10}| = (0.85^{+0.11}_{-0.08}) \times 10^{-3}\ \mathrm{MeV}^{-1}. \tag{3.19}$$

Roughly speaking, the $\Sigma_c(2800)$ widths are about 55 and 15 MeV due to $\Lambda_c\pi$ and $\Sigma_c^{(*)}\pi$, respectively. Hence, the strong decays of $\Sigma_c(2800)$ are indeed dominated by the $\Lambda_c\pi$ mode.

The quark model relation $|h_8| = |h_{10}|$ then leads to

$$|h_8| \approx (0.85^{+0.11}_{-0.08}) \times 10^{-3}\ \mathrm{MeV}^{-1}, \tag{3.20}$$

which improves the previous limit (3.16) by a factor of 3. The calculated partial widths of $\Lambda_c(2625)^+$ shown in Table VI are consistent with experimental limits.

IV. CONCLUSIONS

We began with a brief overview of the charmed baryon spectroscopy and discussed their possible structure and J^P assignment in the quark model. We have assigned $\Sigma_{c2}(\frac{3}{2}^-)$ to $\Sigma_c(2800)$. As for first positive-parity excitations, with the help of the relativistic quark-diquark model and the $^3P_0$ model, we have identified $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ with $\Lambda_c(2800)$, $\tilde{\Xi}_c{}'(\frac{1}{2}^+)$ with $\Xi_c(2980)$, and $\tilde{\Xi}_{c3}^2(\frac{5}{2}^+)$ with $\Xi_c(3080)$, though the last assignment may encounter some potential problems.

With the new Belle measurement of the $\Sigma_c(2455)$ and $\Sigma_c(2520)$ widths and the recent CDF measurement of the strong decays of $\Lambda_c(2595)$ and $\Lambda_c(2625)$, we have updated coupling constants in heavy hadron chiral perturbation theory. We found $g_2 = 0.565^{+0.011}_{-0.024}$ for the P-wave transition between s-wave and s-wave baryons, and $h_2 = 0.63 \pm 0.07$ extracted from $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$. It is substantially enhanced compared to the old value of order 0.437. With the help from the quark model, two of the couplings h_{10} and h_{11} responsible for D-wave transitions between s-wave and p-wave baryons are determined from $\Sigma_c(2880)$ decays. There is a tension for the coupling h_2 as its value extracted from $\Lambda_c(2595)^+ \to \Lambda_c^+\pi\pi$ will imply $\Xi_c(2790)^0 \to \Xi'_c\pi$ and $\Xi_c(2815)^+ \to \Xi_c^*\pi$ rates slightly above the current limits. It is conceivable that SU(3) flavor symmetry breaking can help account for the discrepancy.

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