Strong decays of charmed baryons in heavy hadron chiral perturbation theory

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Strong decays of charmed baryons are analyzed in the framework of heavy hadron chiral perturbation theory (HHChPT) in which heavy quark symmetry and chiral symmetry are synthesized. HHChPT works excellently for describing the strong decays of s-wave charmed baryons. For $L = 1$ orbitally excited states, two of the unknown couplings, namely, h_2 and h_{10}, are determined from the resonant $\Lambda_c^+\pi\pi$ mode produced in the $\Lambda_c(2593)$ decay and the width of $\Sigma_c(2800)$, respectively. Predictions for the strong decays of the p-wave charmed baryon states $\Lambda_c(2625)$, $\Xi_c(2790)$ and $\Xi_c(2815)$ are presented. Since the decay $\Lambda_c(2593)^+ \to \Lambda_c^+\pi\pi$ receives nonresonant contributions, our value for h_2 is smaller than the previous estimates. We also discuss the first positive-parity excited charmed baryons. We conjecture that the charmed baryon $\Lambda_c(2880)$ with $J^P = \frac{5}{2}^+$ is an admixture of $\Lambda_{c2}(\frac{5}{2}^+)$ with an $\tilde{\Lambda}''_{c3}(\frac{5}{2}^+)$; both are $L = 2$ orbitally excited states. The potential model suggests $J^P = \frac{5}{2}^-$ or $\frac{3}{2}^+$ for $\Lambda_c(2940)^+$. Measurements of the ratio of $\Sigma_c^*\pi/\Sigma_c\pi$ will enable us to discriminate the J^P assignments for $\Lambda_c(2940)$. We advocate that the J^P quantum numbers of $\Xi_c(2980)$ and $\Xi_c(3077)$ are $\frac{1}{2}^+$ and $\frac{5}{2}^+$, respectively. Under this J^P assignment, it is easy to understand why $\Xi_c(2980)$ is broader than $\Xi_c(3077)$.

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I. INTRODUCTION

In the past years many new excited charmed baryon states have been discovered by *BABAR*, Belle and CLEO. In particular, B factories have provided a very rich source of charmed baryons both from B decays and from the continuum $e^+e^- \to c\bar{c}$. A new era for the charmed baryon spectroscopy is opened by the rich mass spectrum and the relatively narrow widths of the excited states. 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 quarks and chiral symmetry of the light quarks.

The observed mass spectra and decay widths of charmed baryons are summarized in Tables I and II. Several new excited charmed baryon states such as $\Lambda_c(2765)^+$, $\Lambda_c(2880)^+$, $\Lambda_c(2940)^+$, $\Xi_c(2815)$, $\Xi_c(2980)$ and $\Xi_c(3077)$ have been measured recently and they are still not on the particle listings of 2006 Review of Particle Physics by the Particle Data Group [7]. By now, the $J^P = \frac{1}{2}^+$ and $\frac{1}{2}^-$ antitriplet states: $(\Lambda_c^+, \Xi_c^+, \Xi_c^0)$, $(\Lambda_c(2593)^+, \Xi_c(2790)^+, \Xi_c(2790)^0)$, and $J^P = \frac{1}{2}^+$ and $\frac{3}{2}^+$ sextet

TABLE I. Masses (first entry) and decay widths (second entry) in units of MeV for the excited charmed baryons $\Lambda_c(2880)$, $\Lambda_c(2940)$, $\Xi_c(2980)^{+,0}$, $\Xi_c(3077)^{+,0}$ and $\Omega_c(2768)$.

State	*BABAR* [1–3]	Belle [4,5]	CLEO [6]	Average
$\Lambda_c(2880)^+$	$2881.9 \pm 0.1 \pm 0.5$ $5.8 \pm 1.5 \pm 1.1$	$2881.2 \pm 0.2^{+0.4}_{-0.3}$ $5.5^{+0.7}_{-0.5} \pm 0.4$	2882.5 ± 2.2 <8	2881.5 ± 0.3 5.5 ± 0.6
$\Lambda_c(2940)^+$	$2939.8 \pm 1.3 \pm 1.0$ $17.5 \pm 5.2 \pm 5.9$	$2937.9 \pm 1.0^{+1.8}_{-0.4}$ $10 \pm 4 \pm 5$		2938.8 ± 1.1 13.0 ± 5.0
$\Xi_c(2980)^+$	$2967.1 \pm 1.9 \pm 1.0$ $23.6 \pm 2.8 \pm 1.3$	$2978.5 \pm 2.1 \pm 2.0$ $43.5 \pm 7.5 \pm 7.0$		2971.1 ± 1.7 25.2 ± 3.0
$\Xi_c(2980)^0$		$2977.1 \pm 8.8 \pm 3.5$ 43.5 (fixed)		2977.1 ± 9.5 43.5
$\Xi_c(3077)^+$	$3076.4 \pm 0.7 \pm 0.3$ $6.2 \pm 1.6 \pm 0.5$	$3076.7 \pm 0.9 \pm 0.5$ $6.2 \pm 1.2 \pm 0.8$		3076.5 ± 0.6 6.2 ± 1.1
$\Xi_c(3077)^0$		$3082.8 \pm 1.8 \pm 1.5$ $5.2 \pm 3.1 \pm 1.8$		3082.8 ± 2.3 5.2 ± 3.6
$\Omega_c(2768)^0$	2768.3 ± 3.0			2768.3 ± 3.0

TABLE II. Mass spectra and decay widths (in units of MeV) of charmed baryons. Experimental values are taken from the Particle Data Group [7] and Table I.

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		weak
$\Lambda_c(2593)^+$	$\frac{1}{2}^-$	0	1	1^-	2595.4 ± 0.6	$3.6^{+2.0}_{-1.3}$	$\Sigma_c\pi$, $\Lambda_c\pi\pi$
$\Lambda_c(2625)^+$	$\frac{3}{2}^-$	0	1	1^-	2628.1 ± 0.6	<1.9	$\Lambda_c\pi\pi$, $\Sigma_c\pi$
$\Lambda_c(2765)^+$	$?^?$	?	?	?	2766.6 ± 2.4	50	$\Sigma_c\pi$, $\Lambda_c\pi\pi$
$\Lambda_c(2880)^+$	$\frac{5}{2}^+$	?	?	?	2881.5 ± 0.3	5.5 ± 0.6	$\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$, D^0p
$\Lambda_c(2940)^+$	$?^?$	?	?	?	2938.8 ± 1.1	13.0 ± 5.0	$\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$, D^0p
$\Sigma_c(2455)^{++}$	$\frac{1}{2}^+$	1	0	1^+	2454.02 ± 0.18	2.23 ± 0.30	$\Lambda_c\pi$
$\Sigma_c(2455)^+$	$\frac{1}{2}^+$	1	0	1^+	2452.9 ± 0.4	<4.6	$\Lambda_c\pi$
$\Sigma_c(2455)^0$	$\frac{1}{2}^+$	1	0	1^+	2453.76 ± 0.18	2.2 ± 0.4	$\Lambda_c\pi$
$\Sigma_c(2520)^{++}$	$\frac{3}{2}^+$	1	0	1^+	2518.4 ± 0.6	14.9 ± 1.9	$\Lambda_c\pi$
$\Sigma_c(2520)^+$	$\frac{3}{2}^+$	1	0	1^+	2517.5 ± 2.3	<17	$\Lambda_c\pi$
$\Sigma_c(2520)^0$	$\frac{3}{2}^+$	1	0	1^+	2518.0 ± 0.5	16.1 ± 2.1	$\Lambda_c\pi$
$\Sigma_c(2800)^{++}$	$\frac{3}{2}^-?$	1	1	2^-	2801^{+4}_{-6}	75^{+22}_{-17}	$\Lambda_c\pi$, $\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$
$\Sigma_c(2800)^+$	$\frac{3}{2}^-?$	1	1	2^-	2792^{+14}_{-5}	62^{+60}_{-40}	$\Lambda_c\pi$, $\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$
$\Sigma_c(2800)^0$	$\frac{3}{2}^-?$	1	1	2^-	2802^{+4}_{-7}	61^{+28}_{-18}	$\Lambda_c\pi$, $\Sigma_c^{(*)}\pi$, $\Lambda_c\pi\pi$
Ξ_c^+	$\frac{1}{2}^+$	0	0	0^+	2467.9 ± 0.4		weak
Ξ_c^0	$\frac{1}{2}^+$	0	0	0^+	2471.0 ± 0.4		weak
$\Xi_c'^+$	$\frac{1}{2}^+$	1	0	1^+	2575.7 ± 3.1		$\Xi_c\gamma$
$\Xi_c'^0$	$\frac{1}{2}^+$	1	0	1^+	2578.0 ± 2.9		$\Xi_c\gamma$
$\Xi_c(2645)^+$	$\frac{3}{2}^+$	1	0	1^+	2646.6 ± 1.4	<3.1	$\Xi_c\pi$
$\Xi_c(2645)^0$	$\frac{3}{2}^+$	1	0	1^+	2646.1 ± 1.2	<5.5	$\Xi_c\pi$
$\Xi_c(2790)^+$	$\frac{1}{2}^-$	0	1	1^-	2789.2 ± 3.2	<15	$\Xi_c'\pi$
$\Xi_c(2790)^0$	$\frac{1}{2}^-$	0	1	1^-	2791.9 ± 3.3	<12	$\Xi_c'\pi$
$\Xi_c(2815)^+$	$\frac{3}{2}^-$	0	1	1^-	2816.5 ± 1.2	<3.5	$\Xi_c^*\pi$, $\Xi_c\pi\pi$, $\Xi_c'\pi$
$\Xi_c(2815)^0$	$\frac{3}{2}^-$	0	1	1^-	2818.2 ± 2.1	<6.5	$\Xi_c^*\pi$, $\Xi_c\pi\pi$, $\Xi_c'\pi$
$\Xi_c(2980)^+$	$?^?$	?	?	?	2971.1 ± 1.7	25.2 ± 3.0	see Table VII
$\Xi_c(2980)^0$	$?^?$	?	?	?	2977.1 ± 9.5	43.5	see Table VII
$\Xi_c(3077)^+$	$?^?$	?	?	?	3076.5 ± 0.6	6.2 ± 1.1	see Table VII
$\Xi_c(3077)^0$	$?^?$	?	?	?	3082.8 ± 2.3	5.2 ± 3.6	see Table VII
Ω_c^0	$\frac{1}{2}^+$	1	0	1^+	2697.5 ± 2.6		weak
$\Omega_c(2768)^0$	$\frac{3}{2}^+$	1	0	1^+	2768.3 ± 3.0		$\Omega_c\gamma$

states: (Ω_c, Σ_c, Ξ_c'), (Ω_c^*, Σ_c^*, $\Xi_c'^*$) are established. Notice that except for the parity of the lightest Λ_c^+ and the spin-parity of $\Lambda_c(2880)^+$, none of the other J^P quantum numbers given in Table II has been measured. One has to rely on the quark model to determine the J^P assignments.

This work is organized as follows. In Sec. II, the experimental status of the charmed baryon spectroscopy is reviewed. The p-wave charmed baryons and the first positive-parity excitations are discussed. In Sec. III we first present the relevant chiral Lagrangians which combine heavy quark and chiral symmetries. Then we proceed to the phenomenological implications to the strong decays of s-wave and p-wave charmed baryons as well as first positive-parity excited charmed baryon states. Conclusions are presented in Sec. IV.

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 charmed baryon of interest contains 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 spin-flavor-space wave functions of baryons are totally symmetric since the color wave function is totally antisymmetric. For the ground-state s-wave baryons in the quark model, the symmetries in the flavor and spin of the diquarks are thus correlated. Consequently, the diquark in the

TABLE III. The p-wave charmed baryons and their quantum numbers, where S_ℓ (J_ℓ) is the total spin (angular momentum) of the two light quarks. The quantum number in the subscript labels J_ℓ, while the quantum number in parentheses is referred to the spin of the baryon. In the quark model, the upper (lower) eight multiplets have even (odd) orbital wave functions under the permutation of the two light quarks. That is, L_ℓ for the former is referred to the orbital angular momentum between the diquark and the charmed quark, while L_ℓ for the latter is the orbital angular momentum between the two light quarks. The explicit quark model wave functions for p-wave charmed baryons can be found in [8]. The states antisymmetric in orbital wave functions are denoted by a tilde, while the superscript prime is reserved for the Ξ_c charmed baryons to distinguish between the sextet and antitriplet SU(3) flavor states.

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

flavor-symmetric sextet has spin 1, while the diquark in the flavor-antisymmetric antitriplet has spin 0. When the diquark combines with the charmed quark, the sextet contains both spin-$\frac{1}{2}$ and spin-$\frac{3}{2}$ charmed baryons. However, the antitriplet contains only spin-$\frac{1}{2}$ ones. More specifically, 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. Note that we follow the Particle Data Group (PDG) [7] to use a prime to distinguish the Ξ_c in the $\mathbf{6}$ from the one in the $\bar{\mathbf{3}}$.

The lowest-lying orbitally excited baryon states are the p-wave charmed baryons with their quantum numbers listed in Table III. Although the separate spin angular momentum S_ℓ and orbital angular momentum L_ℓ of the light degrees of freedom are not well defined, they are included for guidance from the quark model. In the heavy quark limit, the spin of the charmed quark S_c and the total angular momentum of the two light quarks $J_\ell = S_\ell + L_\ell$ are separately conserved. It is convenient to use them to enumerate the spectrum of states. There are two types of $L_\ell = 1$ orbital excited charmed baryon states: states with the unit of orbital angular momentum between the diquark and the charmed quark, and states with the unit of orbital angular momentum between the two light quarks. The orbital wave function of the former (latter) is symmetric (antisymmetric) under the exchange of two light quarks. To see this, one can define two independent relative momenta $\mathbf{k} = \frac{1}{2}(\mathbf{p}_1 - \mathbf{p}_2)$ and $\mathbf{K} = \frac{1}{2}(\mathbf{p}_1 + \mathbf{p}_2 - 2\mathbf{p}_c)$ from the two light quark momenta $\mathbf{p}_1$, $\mathbf{p}_2$ and the heavy quark momentum $\mathbf{p}_c$. (In the heavy quark limit, $\mathbf{p}_c$ can be set to zero.) Denoting the quantum numbers L_k and L_K as the eigenvalues of $\mathbf{L}_k^2$ and $\mathbf{L}_K^2$,[1] the k-orbital momentum L_k describes relative orbital excitations of the two light quarks, and the K-orbital momentum L_K describes orbital excitations of the center of the mass of the two light quarks relative to the heavy quark [10]. The p-wave heavy baryon can be either in the ($L_k = 0, L_K = 1$) K-state or the ($L_k = 1, L_K = 0$) k-state. It is obvious that the orbital K-state (k-state) is symmetric (antisymmetric) under the interchange of $\mathbf{p}_1$ and $\mathbf{p}_2$.

There are seven lowest-lying p-wave Λ_c arising from combining the charmed quark spin S_c with light constituents in $J_\ell^{P_\ell} = 1^-$ state: three $J^P = \frac{1}{2}^-$ states, three $J^P = \frac{3}{2}^-$ states and one $J^P = \frac{5}{2}^-$ state. They form three doublets $\Lambda_{c1}(\frac{1}{2},\frac{3}{2})$, $\tilde{\Lambda}_{c1}(\frac{1}{2},\frac{3}{2})$, $\tilde{\Lambda}_{c2}(\frac{3}{2},\frac{5}{2})$ and one singlet $\tilde{\Lambda}_{c0}(\frac{1}{2})$, where we have used a tilde to denote the multiplets antisymmetric in the orbital wave functions under the exchange of two light quarks. In terms of K- and k-states introduced above, the doublets Λ_{c1}, $\tilde{\Lambda}_{c1}$, $\tilde{\Lambda}_{c2}$ and the singlet $\tilde{\Lambda}_{c0}$ are sometimes denoted by Λ_{cK1}, Λ_{ck1}, Λ_{ck2} and Λ_{ck0}, respectively, in the literature [10]. Quark models [9] indicate that the untilde states for Λ- and Σ-type charmed baryons with symmetric orbital wave functions lie about 150 MeV below the tilde ones. The two states in each doublet with $J = J_\ell \pm \frac{1}{2}$ are nearly degenerate; their masses split only by a chromomagnetic interaction.

Since the spin-parity of the newly measured $\Lambda_c(2880)^+$ was recently pinned down to be $\frac{5}{2}^+$ by Belle [4], we shall

[1] In the notation of [9], L_k and L_K correspond to ℓ_ρ and ℓ_λ, respectively.

TABLE IV. The first positive-parity excitations of charmed baryons and their quantum numbers. States with antisymmetric orbital wave functions (i.e. $L_K = L_k = 1$) under the interchange of two light quarks are denoted by a tilde. A prime is used to distinguish between the sextet and antitriplet SU(3) flavor states of the excited Ξ_c.

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

briefly discuss the positive-parity excitations of charmed baryons. Referring to the orbital angular momentum quantum numbers L_K and L_k, the first positive-parity excitations are those states with $L_K + L_k = 2$. For $L_K = 2$, $L_k = 0$, $L = 2$ or $L_K = 0$, $L_k = 2$, $L = 2$, there is one multiplet for positive-parity excited Λ_c and three multiplets for Σ_c as tabulated in Table IV.[2] The orbital states of these multiplets are symmetric under the interchange of the two light quarks. For the case of $L_K = L_k = 1$, 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 $L_K = L_k = 1$ states. The Fermi-Dirac statistics for baryons yields seven more multiplets for positive-parity excited Λ_c states and three more multiplets for Σ_c baryons. The reader is referred to [11] for more details.

In the following we discuss some of the new excited charmed baryon states:

A. Λ_c

$\Lambda_c(2593)^+$ and $\Lambda_c(2625)^+$ form a doublet $\Lambda_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ [12]. The dominant decay mode is $\Sigma_c\pi$ in an S wave for $\Lambda_{c1}(\frac{1}{2}^-)$ and $\Lambda_c\pi\pi$ in a P wave for $\Lambda_{c1}(\frac{3}{2}^-)$. (The two-body mode $\Sigma_c\pi$ is a D-wave in $\Lambda_c(\frac{3}{2}^-)$ decay.) This explains why the width of $\Lambda_c(2625)^+$ is narrower than that of $\Lambda_c(2593)^+$.

$\Lambda_c(2765)^+$ is a broad state ($\Gamma \approx 50$ MeV) first seen in $\Lambda_c^+\pi^+\pi^-$ by CLEO [6]. It appears to resonate through Σ_c and probably also Σ_c^*. However, whether it is a Λ_c^+ or a Σ_c^+ or whether the width might be due to overlapping states are not known. The Skyrme model [13] and the quark model [9] suggest a $J^P = \frac{1}{2}^+$ Λ_c state with a mass 2742 and 2775 MeV, respectively. Therefore, $\Lambda_c(2765)^+$ could be a first positive-parity excitation of Λ_c.

The state $\Lambda_c(2880)^+$ first observed by CLEO [6] in $\Lambda_c^+\pi^+\pi^-$ was also seen by *BABAR* in the D^0p spectrum [1]. It was originally conjectured that, based on its narrow width, $\Lambda_c(2880)^+$ might be a $\tilde{\Lambda}_{c0}^+(\frac{1}{2}^-)$ state [6]. Recently, Belle has studied the experimental constraint on the J^P quantum numbers of $\Lambda_c(2880)^+$ [4]. The angular analysis of $\Lambda_c(2880)^+ \to \Sigma_c^{0,++}\pi^\pm$ indicates that $J = \frac{5}{2}$ is favored over $J = \frac{1}{2}$ or $\frac{3}{2}$, while the study of the resonant structure of $\Lambda_c(2880)^+ \to \Lambda_c^+\pi^+\pi^-$ implies the existence of the $\Sigma_c^*\pi$ intermediate states and $\Gamma(\Sigma_c^*\pi^\pm)/\Gamma(\Sigma_c\pi^\pm) = (24.1 \pm 6.4^{+1.1}_{-4.5})\%$. This value is in agreement with heavy quark symmetry predictions [14] and favors the $\frac{5}{2}^+$ over the $\frac{5}{2}^-$ assignment. We shall return back to this point in Sec. III C. It is interesting to notice that, based on the diquark idea, the quantum numbers $J^P = \frac{5}{2}^+$ have already been predicted in [15] for the $\Lambda_c(2880)$ before the Belle experiment.

[2]Strictly spaking, there are two multiplets for positive-parity excited Λ_c and six multiplets for Σ_c coming from two different orbital states $L_K = 2$, $L_k = 0$ and $L_K = 0$, $L_k = 2$. For simplicity, here we will not distinguish between them.

The highest $\Lambda_c(2940)^+$ was first discovered by *BABAR* in the $D^0 p$ decay mode [1] and confirmed by Belle in the decays $\Sigma_c^0\pi^+$, $\Sigma_c^{++}\pi^-$ which subsequently decay into $\Lambda_c^+\pi^+\pi^-$ [4,16]. Since the mass of $\Lambda_c(2940)^+$ is barely below the threshold of $D^{*0}p$, this observation has motivated the authors of [17] to suggest an exotic molecular state of D^{*0} and p with a binding energy of order 6 MeV for $\Lambda_c(2940)^+$. Its quantum numbers J^P could be $\frac{3}{2}^+$ or $\frac{5}{2}^-$ as suggested by the quark model calculation [9].

B. Σ_c

The highest isotriplet charmed baryons $\Sigma_c(2800)^{++,+,0}$ decaying to $\Lambda_c^+\pi$ were first measured by Belle [18]. They are most likely to be the $J^P = \frac{3}{2}^-$ Σ_{c2} states because the $\Sigma_{c2}(\frac{3}{2}^-)$ baryon decays principally into the $\Lambda_c\pi$ system in a D-wave, while $\Sigma_{c1}(\frac{3}{2}^-)$ decays mainly to the two-pion system $\Lambda_c\pi\pi$ in a P-wave. The state $\Sigma_{c0}(\frac{1}{2}^-)$ can decay into $\Lambda_c\pi$ in an S-wave, but it is very broad with width of order 406 MeV (see Sec. III C).

C. Ξ_c

The states $\Xi_c(2790)$ and $\Xi_c(2815)$ form a doublet $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$. Since the diquark transition $1^- \to 0^+ + \pi$ is prohibited, $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ cannot decay to $\Xi_c\pi$. The dominant decay mode is $[\Xi_c'\pi]_S$ for $\Xi_{c1}(\frac{1}{2}^-)$ and $[\Xi_c^*\pi]_S$ for $\Xi_{c1}(\frac{3}{2}^-)$ where Ξ_c^* stands for $\Xi_c(2645)$.

The new charmed strange baryons $\Xi_c(2980)^+$ and $\Xi_c(3077)^+$ that decay into $\Lambda_c^+K^-\pi^+$ were first observed by Belle [5] and confirmed by *BABAR* [2]. In the recent *BABAR* measurement [2], the $\Xi_c(2980)^+$ is found to decay resonantly through the intermediate state $\Sigma_c(2455)^{++}K^-$ with 4.9σ significance and nonresonantly to $\Lambda_c^+K^-\pi^+$ with 4.1σ significance. With 5.8σ significance, the $\Xi_c(3077)^+$ is found to decay resonantly through $\Sigma_c(2455)^{++}K^-$, and with 4.6σ significance, it is found to decay through $\Sigma_c(2520)^{++}K^-$. The significance of the signal for the nonresonant decay $\Xi_c(3077)^+ \to \Lambda_c^+K^-\pi^+$ is 1.4σ.

D. Ω_c

At last, the $J^P = \frac{3}{2}^+$ $\Omega_c(2768)$ charmed baryon was recently observed by *BABAR* in the decay $\Omega_c(2768)^0 \to \Omega_c^0\gamma$ [3]. With this new observation, the $\frac{3}{2}^+$ sextet is finally completed. However, it will be very difficult to measure the electromagnetic decay rate because the width of Ω_c^*, which is predicted to be of order 0.9 keV [19], is too narrow to be experimentally resolvable.

III. STRONG DECAYS

Because of the rich mass spectrum and the relatively narrow widths of the excited states, the charmed baryon system offers an excellent ground for testing the ideas and predictions of heavy quark symmetry and light flavor SU(3) symmetry. The pseudoscalar mesons involved in the strong decays of charmed baryons such as $\Sigma_c \to \Lambda_c\pi$ are soft. Therefore, heavy quark symmetry of the heavy quark and chiral symmetry of the light quarks will have interesting implications for the low-energy dynamics of heavy baryons interacting with the Goldstone bosons.

The strong decays of charmed baryons are most conveniently described by the heavy hadron chiral Lagrangians in which heavy quark symmetry and chiral symmetry are incorporated [20,21]. The Lagrangian involves two coupling constants g_1 and g_2 for P-wave transitions between s-wave and s-wave baryons [20], 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 [8].

Since the general chiral Lagrangian for heavy baryons coupling to the pseudoscalar mesons can be expressed compactly in terms of superfields, we first introduce the superfields for s-wave baryons given by

$$\begin{aligned}
\mathcal{S}_\mu^{ij} &= \frac{1+\not v}{2}B_{6\mu}^{*ij} + \frac{1}{\sqrt{3}}(\gamma_\mu + v_\mu)\gamma_5\frac{1+\not v}{2}B_6^{ij},\\
\bar{\mathcal{S}}_\mu^{ij} &= \bar{B}_{6\mu}^{*ij}\frac{1+\not v}{2} - \frac{1}{\sqrt{3}}\bar{B}_6^{ij}\frac{1+\not v}{2}\gamma_5(\gamma_\mu + v_\mu),\\
\mathcal{T}_i &= \frac{1+\not v}{2}(\Xi_c^0, -\Xi_c^+, \Lambda_c^+)_i = \frac{1}{2}\epsilon_{ijk}\frac{1+\not v}{2}(B_{\bar{3}})_{jk}.
\end{aligned} \tag{3.1}$$

where the matrices B_6 and $B_{6\mu}^*$ are defined in [20]

$$\begin{aligned}
(B_6)_{ij} &= \begin{pmatrix} \Sigma_c^{++} & \frac{1}{\sqrt{2}}\Sigma_c^+ & \frac{1}{\sqrt{2}}\Xi_c'^+ \\ \frac{1}{\sqrt{2}}\Sigma_c^+ & \Sigma_c^0 & \frac{1}{\sqrt{2}}\Xi_c'^0 \\ \frac{1}{\sqrt{2}}\Xi_c^{+'} & \frac{1}{\sqrt{2}}\Xi_c'^0 & \Omega_c^0 \end{pmatrix}_{ij},\\
(B_{\bar{3}})_{ij} &= \begin{pmatrix} 0 & \Lambda_c^+ & \Xi_c^+ \\ -\Lambda_c^+ & 0 & \Xi_c^0 \\ -\Xi_c^+ & -\Xi_c^0 & 0 \end{pmatrix}_{ij}.
\end{aligned} \tag{3.2}$$

The superfield for p-wave $\bar{\mathbf{3}}$ multiplets symmetric in orbital wave functions such as $\Lambda_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ and $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ is given by

$$\mathcal{R}^i_\mu = \frac{1}{\sqrt{3}}(\gamma_\mu + v_\mu)\gamma_5 R^i + R_\mu^{*i}, \tag{3.3}$$

with

$$\begin{aligned}
R_i &= \frac{1+\not v}{2}(\Xi_{c_1}^0, -\Xi_{c_1}^+, \Lambda_{c_1}^+)_i,\\
R_\mu^{*i} &= \frac{1+\not v}{2}(\Xi_{c_1}^{*0}, -\Xi_{c_1}^{*+}, \Lambda_{c_1}^{*+})_{\mu i},
\end{aligned} \tag{3.4}$$

where $\mathcal{B}_c^*$ denotes a spin-$\frac{3}{2}$ charmed baryon. Note that $v \cdot \mathcal{S} = v \cdot \mathcal{R} = 0$.

There are three other p-wave sextet multiplets with antisymmetric orbital states and with quantum numbers $J_\ell^{P_\ell} = 0^-, 1^-, 2^-$. Their $I = 1$ members are $\Sigma_{c0}(\frac{1}{2}^-)$,

$\Sigma_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ and $\Sigma_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$, while the corresponding $I = \frac{1}{2}$ members are $\Xi'_{c0}(\frac{1}{2}^-)$, $\Xi'_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ and $\Xi'_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ (see Table III), The $J_\ell^{P_\ell} = 0^-$ multiplet will be represented as a symmetric matrix $(U)_{ij}$ defined in the same manner as $(B_6)_{ij}$

$$U_{ij} = \begin{pmatrix} \Sigma_{c0}^{++} & \frac{1}{\sqrt{2}}\Sigma_{c0}^{+} & \frac{1}{\sqrt{2}}\Xi_{c0}^{\prime +} \\ \frac{1}{\sqrt{2}}\Sigma_{c0}^{+} & \Sigma_{c0}^{0} & \frac{1}{\sqrt{2}}\Xi_{c0}^{\prime 0} \\ \frac{1}{\sqrt{2}}\Xi_{c0}^{\prime +} & \frac{1}{\sqrt{2}}\Xi_{c0}^{\prime 0} & \Omega_{c0}^{0} \end{pmatrix}_{ij}, \qquad (3.5)$$

The $J_\ell^{P_\ell} = 1^-$ multiplet will be represented as a superfield similar to (3.3) but with a symmetric matrix $\mathcal{V}_\mu^{ij}$

$$\mathcal{V}_\mu^{ij} = \frac{1}{\sqrt{3}}(\gamma_\mu + v_\mu)\gamma_5 V^{ij} + V_\mu^{*ij}, \qquad (3.6)$$

where V_{ij} has the same expression as U_{ij} except for the replacement of the superscript "$c0$" by "$c1$". The superfield corresponding to the $J_\ell^{P_\ell} = 2^-$ baryons is constructed as [22]

$$\mathcal{X}_{\mu\nu}^{ij} = X_{\mu\nu}^{*ij} + \frac{1}{\sqrt{10}}\{(\gamma_\mu + v_\mu)\gamma_5 g_\nu^\alpha + (\gamma_\nu + v_\nu)\gamma_5 g_\mu^\alpha\}X_\alpha^{ij}, \qquad (3.7)$$

with $X_{\mu\nu}^{*ij}$ a spin-$\frac{5}{2}$ Rarita-Schwinger field and X_α^{ij} its spin-$\frac{3}{2}$ heavy quark symmetry partner.

The p-wave states with antisymmetric orbital wave functions can be constructed in complete analogy to the symmetric ones. Following [8], we use the superfield $\tilde{\mathcal{R}}_\mu^{ij}$ constructed in analogy to $\mathcal{S}_\mu^{ij}$ to represent the two sextets $\tilde{\Sigma}_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ and $\tilde{\Xi}'_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$. Likewise, we use the superfields $\tilde{\mathcal{U}}_\mu^i$, $\tilde{\mathcal{V}}_\mu^i$, $\tilde{\mathcal{X}}_{\mu\nu}^i$ to denote the antitriplets: $\tilde{\Lambda}_{c0}^+$, $\tilde{\Lambda}_{c1}^+$, $\tilde{\Lambda}_{c2}^+$ in $I = 0$ and $\tilde{\Xi}_{c0}$, $\tilde{\Xi}_{c1}$, $\tilde{\Xi}_{c2}$ in $I = \frac{1}{2}$.

The leading Lagrangian terms describing P-wave couplings among the s-wave baryons and S-wave couplings between the s-wave and p-wave baryons are

$$\mathcal{L}_P = \frac{3}{2} i g_1 \epsilon_{\mu\nu\sigma\lambda} \mathrm{Tr}(\bar{\mathcal{S}}^\mu v^\nu A^\sigma \mathcal{S}^\lambda) - \sqrt{3} g_2 \mathrm{Tr}(\bar{B}_{\bar{3}} A^\mu \mathcal{S}_\mu + \bar{\mathcal{S}}^\mu A_\mu B_{\bar{3}}), \qquad (3.8)$$

and[3]

$$\begin{aligned}\mathcal{L}_S = {} & h_2\{\epsilon_{ijk}\bar{\mathcal{R}}_i^\mu v_\nu A_{jl}^\nu \mathcal{S}_\mu^{kl} + \epsilon_{ijk}\bar{\mathcal{S}}_\mu^{kl} v_\nu A_{lj}^\nu \mathcal{R}_i^\mu\} \\ & + h_3 \mathrm{Tr}(\bar{B}_{\bar{3}} v_\mu A^\mu U + \bar{U} v^\mu A_\mu B_{\bar{3}}) \\ & + h_4 \mathrm{Tr}\{\bar{\mathcal{V}}_\mu v_\nu A^\nu \mathcal{S}^\mu + \bar{\mathcal{S}}_\mu v_\nu A^\nu \mathcal{V}^\mu\} \\ & + h_5 \mathrm{Tr}(\bar{\tilde{\mathcal{R}}}_\mu v_\nu A^\nu \mathcal{S}^\mu + \bar{\mathcal{S}}^\mu v_\nu A^\nu \tilde{\mathcal{R}}_\mu) \\ & + h_6(\bar{\mathcal{T}}_i v_\nu A_{ji}^\nu \tilde{\mathcal{U}}_j + \bar{\tilde{\mathcal{U}}}_i v_\nu A_{ji}^\nu \mathcal{T}_j) \\ & + h_7\{\epsilon_{ijk}\bar{\tilde{\mathcal{V}}}_\mu^i v_\nu A_{jl}^\nu \mathcal{S}_{kl}^\mu + \epsilon_{ijk}\bar{\mathcal{S}}_{kl}^\mu v_\nu A_{lj}^\nu \tilde{\mathcal{V}}_\mu^i\}, \end{aligned} \qquad (3.9)$$

respectively. The Goldstone bosons couple to the matter fields through the nonlinear axial-vector field A_μ defined as

$$A_\mu = \frac{i}{2}(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger) = -\frac{1}{f_\pi}\partial_\mu \phi + \frac{1}{6f_\pi^3}[\phi, [\phi, \partial_\mu \phi]] + \cdots, \qquad (3.10)$$

with $\xi = \exp(i\phi/f_\pi)$, $\phi \equiv \frac{1}{\sqrt{2}}\pi^a \lambda^a$ and $f_\pi = 132$ MeV.

The D-wave couplings of the p-wave baryons to s-wave baryons are described by dimension-5 terms in the effective Lagrangian [8]

$$\begin{aligned}\mathcal{L}_D = {} & i h_8 \epsilon_{ijk} \bar{\mathcal{S}}_\mu^{kl}\Big(\mathcal{D}^\mu A^\nu + \mathcal{D}^\nu A^\mu + \frac{2}{3} g^{\mu\nu}(v\cdot\mathcal{D})(v\cdot A)\Big)_{lj}\mathcal{R}_\nu^i \\ & + i h_9 \mathrm{Tr}\Big\{\bar{\mathcal{S}}_\mu\Big(\mathcal{D}^\mu A^\nu + \mathcal{D}^\nu A^\mu + \frac{2}{3} g^{\mu\nu}(v\cdot\mathcal{D})(v\cdot A)\Big)\mathcal{V}_\nu\Big\} + i h_{10}\epsilon_{ijk}\bar{\mathcal{T}}_i(\mathcal{D}_\mu A_\nu + \mathcal{D}_\nu A_\mu)_{jl}\mathcal{X}_{kl}^{\mu\nu} \\ & + h_{11}\epsilon_{\mu\nu\sigma\lambda}\mathrm{Tr}\{\bar{\mathcal{S}}^\mu(\mathcal{D}^\nu A_\alpha + \mathcal{D}_\alpha A^\nu)\mathcal{X}^{\alpha\sigma}\}v^\lambda + i h_{12}\mathrm{Tr}\Big\{\bar{\mathcal{S}}_\mu\Big(\mathcal{D}^\mu A^\nu + \mathcal{D}^\nu A^\mu + \frac{2}{3} g^{\mu\nu}(v\cdot\mathcal{D})(v\cdot A)\Big)\tilde{\mathcal{R}}_\nu\Big\} \\ & + i h_{13}\epsilon_{ijk}\bar{\mathcal{S}}_\mu^{kl}\Big(\mathcal{D}^\mu A^\nu + \mathcal{D}^\nu A^\mu + \frac{2}{3} g^{\mu\nu}(v\cdot\mathcal{D})(v\cdot A)\Big)_{lj}\tilde{\mathcal{V}}_\nu^i + i h_{14}\bar{\mathcal{T}}_i(\mathcal{D}^\mu A^\nu + \mathcal{D}^\nu A^\mu)_{ji}\tilde{\mathcal{X}}_{\mu\nu}^j \\ & + h_{15}\epsilon_{\mu\nu\sigma\lambda}\epsilon_{ijk}\bar{\mathcal{S}}_{kl}^\mu(\mathcal{D}^\nu A_\alpha + \mathcal{D}_\alpha A^\nu)_{lj}\tilde{\mathcal{X}}_i^{\alpha\sigma}v^\lambda, \end{aligned} \qquad (3.11)$$

where the covariant derivative of the axial-vector field A_μ is defined as $\mathcal{D}_\mu A_\nu = \partial_\mu A_\nu + [V_\mu, A_\nu]$ with

$$V_\mu = \frac{1}{2}(\xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger) = \frac{1}{2f_\pi^2}[\phi, \partial_\mu \phi] + \cdots, \qquad (3.12)$$

and satisfies the relation $\mathcal{D}_\mu A_\nu - \mathcal{D}_\nu A_\mu = 0$. Note that a pure D-wave is described by the configuration

[3] The original h_1 term defined in [12] is now the g_2 term in Eq. (3.8) where we have followed [20] for the definition of g_1 and g_2 couplings. The $h_3, \cdots, h_7$ terms were first introduced in [8].

$$D_{\mu\nu} = (\partial_\mu - v_\mu v\cdot\partial)(\partial_\nu - v_\nu v\cdot\partial) - \frac{1}{3}(g_{\mu\nu} - v_\mu v_\nu)(\partial - v v\cdot\partial)^2, \tag{3.13}$$

satisfying $v^\mu D_{\mu\nu} = 0$, $D_\mu{}^\mu = 0$ and $D_{\mu\nu} = D_{\nu\mu}$. It is straightforward to show that the structure $\mathcal{D}^\mu A^\nu + \mathcal{D}^\nu A^\mu + \frac{2}{3}g^{\mu\nu}(v\cdot\mathcal{D})(v\cdot A)$ appearing in Eq. (3.11) indeed projects out a pure D-wave.

Some of the partial widths derived from the Lagrangians (3.8) and (3.9) are [8]:

$$\begin{aligned}
\Gamma(\Sigma_c^*\to\Sigma_c\pi) &= \frac{g_1^2}{2\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Sigma_c^*}}p_\pi^3,\\
\Gamma(\Sigma_c\to\Lambda_c\pi) &= \frac{g_2^2}{2\pi f_\pi^2}\frac{m_{\Lambda_c}}{m_{\Sigma_c}}p_\pi^3,\\
\Gamma(\Lambda_{c1}(1/2^-)\to\Sigma_c\pi) &= \frac{h_2^2}{2\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Lambda_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\Sigma_{c0}(1/2^-)\to\Lambda_c\pi) &= \frac{h_3^2}{2\pi f_\pi^2}\frac{m_{\Lambda_c}}{m_{\Sigma_{c0}}}E_\pi^2 p_\pi,\\
\Gamma(\Sigma_{c1}(1/2^-)\to\Sigma_c\pi) &= \frac{h_4^2}{4\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Sigma_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\tilde\Sigma_{c1}(1/2^-)\to\Sigma_c\pi) &= \frac{h_5^2}{4\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde\Sigma_{c1}}}E_\pi^2 p_\pi,\\
\Gamma(\tilde\Xi_{c0}(1/2^-)\to\Xi_c\pi) &= \frac{h_6^2}{2\pi f_\pi^2}\frac{m_{\Xi_c}}{m_{\tilde\Xi_{c0}}}E_\pi^2 p_\pi,\\
\Gamma(\tilde\Lambda_{c1}(1/2^-)\to\Sigma_c\pi) &= \frac{h_7^2}{2\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\tilde\Lambda_{c1}}}E_\pi^2 p_\pi,
\end{aligned} \tag{3.14}$$

where p_π is the c.m. momentum of the pion and $f_\pi = 132$ MeV. Unfortunately, the decay $\Sigma_c^*\to\Sigma_c\pi$ is kinematically prohibited since the mass difference between Σ_c^* and Σ_c is only of order 65 MeV. Consequently, the coupling g_1 cannot be extracted directly from the strong decays of heavy baryons. Note that since the charge of the final state is not specified in Eq. (3.14), care must be taken for the neutral pion state. For example, an additional factor of $\frac{1}{2}$ should be taken for $\Sigma_c^{*+}\to\Sigma_c^+\pi^0$ and $\tilde\Xi_{c0}^+\to\Xi_c^+\pi^0$, but not for $\Sigma_c^+\to\Lambda_c^+\pi^0$.

In the quark model, various couplings in Eqs. (3.9) and (3.11) are related to each other. The S-wave couplings between the s-wave and the p-wave baryons are related by [8]

$$\frac{|h_3|}{|h_4|} = \frac{\sqrt{3}}{2}, \qquad \frac{|h_2|}{|h_4|} = \frac{1}{2}, \qquad \frac{|h_5|}{|h_6|} = \frac{2}{\sqrt{3}}, \qquad \frac{|h_5|}{|h_7|} = 1. \tag{3.15}$$

The D-wave couplings satisfy the relations [8]

$$|h_8| = |h_9| = |h_{10}|, \qquad \frac{|h_{11}|}{|h_{10}|} = \frac{|h_{15}|}{|h_{14}|} = \sqrt{2}, \qquad \frac{|h_{12}|}{|h_{13}|} = 2, \qquad \frac{|h_{14}|}{|h_{13}|} = 1. \tag{3.16}$$

From the dimensional analysis, it is expected that the dimensional D-wave couplings $h_{8,\cdots,14}$ are of order $1/\Lambda_\chi \sim (1.0-1.2)\times 10^{-3}$ MeV^{-1}, where $\Lambda_\chi = (0.83\sim 1)$ GeV is the chiral symmetry breaking scale.

As discussed in Sec. II, there exist first positive-parity excited charmed baryon states. Their orbital angular momentum L is 2, 1 or 0. The transitions of these excited baryons to the s-wave baryons involve S-wave, P-wave and F-wave couplings. However, given the complications with the even-parity excitations, here we will not generalize HHChPT to include them.

In terms of the Wigner 6-j symbol, the decay rate for the baryon decay process $J\to J'+\pi$ after spin-averaging over the initial spin and summing over final spins has the expression [23,24]

$$\Gamma(J\to J'+\pi) = (2J_\ell+1)(2J'+1)\left|\begin{Bmatrix} L_\pi & J'_\ell & J_\ell \\ S_Q & J & J' \end{Bmatrix}\right|^2 \times p_\pi^{2L_\pi+1}|M_{L_\pi}|^2, \tag{3.17}$$

where L_π is the orbital angular momentum of the pion, $S_Q = \frac{1}{2}$ is the heavy quark spin and M_{L_π} is the reduced matrix element which is independent of J and J'. This relation is very useful in relating strong decays into different multiplets for a given partial wave.

A. Strong decays of s-wave charmed baryons

In the framework of heavy hadron chiral pertrubation theory (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. We shall use the measured rates of $\Sigma_c^{++}\to\Lambda_c^+\pi^+$, $\Sigma_c^{*++}\to\Lambda_c^+\pi^+$ and $\Sigma_c^{*0}\to\Lambda_c^+\pi^-$ as inputs to obtain

$$|g_2| = 0.605^{+0.039}_{-0.043}, \qquad 0.57\pm 0.04, \qquad 0.60\pm 0.04, \tag{3.18}$$

respectively, where we have neglected the tiny contributions from electromagnetic decays. Note that $|g_2|$ obtained from $\Sigma_c^0\to\Lambda_c^+\pi^-$ has the same central value as the first one in Eq. (3.18) except that the errors are slightly large.[4] Hence, the averaged g_2 is[5]

$$|g_2| = 0.591\pm 0.023. \tag{3.19}$$

[4]Historically, based on the nonrelativistic quark model, the prediction $\Gamma(\Sigma_c^0\to\Lambda_c^+\pi^-) = 2.45$ MeV was made long before experiment [20].

[5]For previous efforts of extracting g_2 from experiment using HHChPT, see [8,25].

As pointed out in [20], 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

$$g_1 = \frac{4}{3} g_A^q, \qquad g_2 = \sqrt{\frac{2}{3}} g_A^q. \tag{3.20}$$

Using $g_A^q = 0.75$ which is required to reproduce the correct value of the nucleon axial coupling $g_A^N = 1.25$, we obtain

$$g_1 = 1, \qquad g_2 = 0.61. \tag{3.21}$$

Hence, the quark model prediction is in good agreement with experiment, while the large-N_c prediction $|g_2| = g_A^N/\sqrt{2} = 0.88$ [26] deviates from the data by 2σ. Applying (3.18) leads to (see also Table V)

$$\begin{aligned} \Gamma(\Xi_c^{*+}) &= \Gamma(\Xi_c^{*+} \to \Xi_c^+ \pi^0, \Xi_c^0 \pi^+) \\ &= \frac{g_2^2}{4\pi f_\pi^2}\left(\frac{1}{2}\frac{m_{\Xi_c^+}}{m_{\Xi_c^{*+}}} p_\pi^3 + \frac{m_{\Xi_c^0}}{m_{\Xi_c^{*+}}} p_\pi^3\right) \\ &= (2.7 \pm 0.2)\ \mathrm{MeV}, \\ \Gamma(\Xi_c^{*0}) &= \Gamma(\Xi_c^{*0} \to \Xi_c^+ \pi^-, \Xi_c^0 \pi^0) \\ &= \frac{g_2^2}{4\pi f_\pi^2}\left(\frac{m_{\Xi_c^+}}{m_{\Xi_c^{*0}}} p_\pi^3 + \frac{1}{2}\frac{m_{\Xi_c^0}}{m_{\Xi_c^{*0}}} p_\pi^3\right) \\ &= (2.8 \pm 0.2)\ \mathrm{MeV}. \end{aligned} \tag{3.22}$$

Note that we have neglected the effect of $\Xi_c - \Xi_c'$ mixing in calculations (for recent considerations, see [31,32]). Therefore, the predicted total width of Ξ_c^{*+} is in the vicinity of the current limit $\Gamma(\Xi_c^{*+}) < 3.1$ MeV [33].

It is clear from Table V that the strong decay width of Σ_c is smaller than that of Σ_c^* by a factor of $\sim$7, although they will become the same in the limit of heavy quark symmetry. This is ascribed to the fact that the c.m. momentum of the pion is around 90 MeV in the decay $\Sigma_c \to \Lambda_c \pi$ while it is 2 times bigger in $\Sigma_c^* \to \Lambda_c \pi$. Since Σ_c states are significantly narrower than their spin-$\frac{3}{2}$ counterparts, this explains why the measurement of their widths came out much later. Instead of using the data to fix the coupling constants in a model-independent manner, there exist some calculations of couplings in various models such as the relativistic light-front model [27], the relativistic three-quark model [28] and light-cone sum rules [29,34]. The calculated results are summarized in Table V.

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 $\Xi_c^{*0} \to \Xi_c^0 \gamma$, which is prohibited at tree level by SU(3) symmetry but can be induced by chiral loops. A measurement of $\Gamma(\Xi_c^{*0} \to \Xi_c^0 \gamma)$ will yield two possible solutions for g_1. Assuming the validity of the quark model relations among different coupling constants, the experimental value of g_2 implies $|g_1| = 0.93 \pm 0.16$ [25].

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

Some of the *S*-wave and *D*-wave couplings of *p*-wave baryons to *s*-wave baryons can be determined. In principle, the coupling h_2 is readily extracted from $\Lambda_c(2593)^+ \to \Sigma_c^0 \pi^+$ with $\Lambda_c(2593)$ being identified as $\Lambda_{c1}(\frac{1}{2}^-)$. However, since $\Lambda_c(2593)^+ \to \Sigma_c \pi$ is kinematically barely allowed, the finite width effects of the intermediate resonant states could become important [35]. If these effects are neglected, then from Eq. (3.14) and the measured decay rates of $\Lambda_c(2593)^+ \to \Sigma_c^0 \pi^+$ and $\Lambda_c(2593)^+ \to \Sigma_c^{++} \pi^-$, we find

$$|h_2| = 0.41 \pm 0.11. \tag{3.23}$$

Before proceeding to a more precise determination of h_2, we make several remarks on the partial widths of $\Lambda_c(2593)^+$ decays. (i) PDG [7] has assumed the isospin relation, namely, $\Gamma(\Lambda_c^+ \pi^+ \pi^-) = 2\Gamma(\Lambda_c^+ \pi^0 \pi^0)$ to extract the branching ratios for $\Sigma_c \pi$ modes. However, the decay $\Lambda_c(2593) \to \Lambda_c \pi\pi$ occurs very close to the threshold as $m_{\Lambda_c(2593)} - m_{\Lambda_c} = 308.9 \pm 0.6$ MeV. Hence, the phase space is very sensitive to the small isospin-violating mass differences between members of pions and charmed Sigma

TABLE V. Decay widths (in units of MeV) of *s*-wave charmed baryons. Theoretical predictions of [27] are taken from Table IV of [28].

Decay	Expt. [7]	This work HHChPT	Tawfiq *et al.* [27]	Ivanov *et al.* [28]	Huang *et al.* [29]	Albertus *et al.* [30]
$\Sigma_c^{++} \to \Lambda_c^+ \pi^+$	2.23 ± 0.30	input	1.51 ± 0.17	2.85 ± 0.19	2.5	2.41 ± 0.07
$\Sigma_c^{+} \to \Lambda_c^+ \pi^0$	<4.6	2.5 ± 0.2	1.56 ± 0.17	3.63 ± 0.27	3.2	2.79 ± 0.08
$\Sigma_c^{0} \to \Lambda_c^+ \pi^-$	2.2 ± 0.4	input	1.44 ± 0.16	2.65 ± 0.19	2.4	2.37 ± 0.07
$\Sigma_c(2520)^{++} \to \Lambda_c^+ \pi^+$	14.9 ± 1.9	input	11.77 ± 1.27	21.99 ± 0.87	8.2	17.52 ± 0.75
$\Sigma_c(2520)^{+} \to \Lambda_c^+ \pi^0$	<17	16.6 ± 1.3			8.6	17.31 ± 0.74
$\Sigma_c(2520)^{0} \to \Lambda_c^+ \pi^-$	16.1 ± 2.1	input	11.37 ± 1.22	21.21 ± 0.81	8.2	16.90 ± 0.72
$\Xi_c(2645)^{+} \to \Xi_c^{0,+} \pi^{+,0}$	<3.1	2.7 ± 0.2	1.76 ± 0.14	3.04 ± 0.37		3.18 ± 0.10
$\Xi_c(2645)^{0} \to \Xi_c^{+,0} \pi^{-,0}$	<5.5	2.8 ± 0.2	1.83 ± 0.06	3.12 ± 0.33		3.03 ± 0.10

baryon multiplets. Since the neutral pion is slightly lighter than the charged one, it turns out that both $\Lambda_c^+\pi^+\pi^-$ and $\Lambda_c^+\pi^0\pi^0$ have very similar rates [see Eq. (3.35) below]. (ii) Taking $\mathcal{B}(\Lambda_c(2593)^+ \to \Lambda_c^+\pi^+\pi^-) \approx 0.5$ and using the measured ratios [7]

$$\frac{\Gamma(\Lambda_c(2593)^+ \to \Sigma_c^{++}\pi^-)}{\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi^+\pi^-)} = 0.36 \pm 0.10, \quad \frac{\Gamma(\Lambda_c(2593)^+ \to \Sigma_c^{0}\pi^+)}{\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi^+\pi^-)} = 0.37 \pm 0.10, \tag{3.24}$$

we obtain

$$\mathcal{B}(\Lambda_c(2593)^+ \to \Sigma_c^{++}\pi^-) = 0.18 \pm 0.05, \quad \mathcal{B}(\Lambda_c(2593)^+ \to \Sigma_c^{0}\pi^+) = 0.19 \pm 0.05, \tag{3.25}$$

and

$$\Gamma(\Lambda_c(2593)^+ \to \Sigma_c^{++}\pi^-) = 0.65^{+0.41}_{-0.31}\ \text{MeV}, \quad \Gamma(\Lambda_c(2593)^+ \to \Sigma_c^{0}\pi^+) = 0.67^{+0.41}_{-0.31}\ \text{MeV}. \tag{3.26}$$

(iii) The nonresonant or direct three-body decay mode $\Lambda_c^+\pi^+\pi^-$ has a branching ratio of 0.14 ± 0.08 [7]. Assuming the same for $\Lambda_c^+\pi^0\pi^0$, then the fractions for resonant and nonresonant $\Lambda_c^+\pi\pi$ are 0.73 ± 0.15 and 0.27 ± 0.15, respectively, where $\Lambda_c^+\pi\pi = \Lambda_c^+\pi^+\pi^- + \Lambda_c^+\pi^0\pi^0$. From the measured total width of $\Lambda_c(2593)^+$, $3.6^{+2.0}_{-1.3}$ MeV [7], we are led to

$$\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi\pi)_{\rm R} = (2.63^{+1.56}_{-1.09})\ \text{MeV}, \quad \Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi\pi)_{\rm NR} = (0.97^{+0.76}_{-0.64})\ \text{MeV}. \tag{3.27}$$

Pole contributions to the decays $\Lambda_c(2593)^+$, $\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi$ have been considered in [8,12,29] with the finite width effects included. The intermediate states of interest are Σ_c and Σ_c^* poles. The resonant contribution arises from the Σ_c pole, while the nonresonant term receives a contribution from the Σ_c^* pole. (Since $\Lambda_c(2593)^+$, $\Lambda_c(2625)^+ \to \Sigma_c^*\pi$ are not kinematically allowed, the Σ_c^* pole is not a resonant contribution.) The decay rates thus depend on two coupling constants h_2 and h_8. The decay rate for the process $\Lambda_{c_1}^+(2593) \to \Lambda_c^+\pi^+\pi^-$ can be calculated in the framework of heavy hadron chiral perturbation theory [8]

$$\frac{d^2\Gamma(\Lambda_{c1}^+(2593) \to \Lambda_c^+\pi^+(E_1)\pi^-(E_2))}{dE_1 dE_2} = \frac{g_2^2}{16\pi^3 f_\pi^4} m_{\Lambda_c^+}\{\mathbf{p}_2^2|A|^2 + \mathbf{p}_1^2|B|^2 + 2\mathbf{p}_1\cdot\mathbf{p}_2\,\mathrm{Re}(AB^*)\}, \tag{3.28}$$

with

$$A(E_1, E_2) = \frac{h_2 E_1}{\Delta_R - \Delta_{\Sigma_c^0} - E_1 + i\Gamma_{\Sigma_c^0}/2} - \frac{\frac{2}{3}h_8\mathbf{p}_2^2}{\Delta_R - \Delta_{\Sigma_c^{*0}} - E_1 + i\Gamma_{\Sigma_c^{*0}}/2} + \frac{2h_8\mathbf{p}_1\cdot\mathbf{p}_2}{\Delta_R - \Delta_{\Sigma_c^{*++}} - E_2 + i\Gamma_{\Sigma_c^{*++}}/2}, \tag{3.29}$$

$$B(E_1, E_2; \Delta_{\Sigma_c^{(*)0}}, \Delta_{\Sigma_c^{(*)++}}) = A(E_2, E_1; \Delta_{\Sigma_c^{(*)++}}, \Delta_{\Sigma_c^{(*)0}}), \tag{3.30}$$

where $\Delta_R = m_{\Lambda_c(2593)} - m_{\Lambda_c}$ and $\Delta_{\Sigma_c^{(*)}} = m_{\Sigma_c^{(*)}} - m_{\Lambda_c}$.

For the spin-$\frac{3}{2}$ state $\Lambda_c(2625)$, its decay is dominated by the three-body channel $\Lambda_c^+\pi\pi$ as the major two-body decay $\Sigma_c\pi$ is a D-wave one. As for the decay $\Lambda_{c1}^+(2625) \to \Lambda_c^+\pi^+\pi^-$, its rate is given by [8]

$$\begin{aligned}\frac{d^2\Gamma(\Lambda_{c1}^+(2625) \to \Lambda_c^+\pi^+(E_1)\pi^-(E_2))}{dE_1 dE_2} &= \frac{g_2^2}{16\pi^3 f_\pi^4} M_{\Lambda_c^+}\{\mathbf{p}_1^2|C|^2 + \mathbf{p}_2^2|E|^2 \\ &\quad + 2\mathbf{p}_1\cdot\mathbf{p}_2\,\mathrm{Re}(CE^*) + [\mathbf{p}_1^2\mathbf{p}_2^2 - (\mathbf{p}_1\cdot\mathbf{p}_2)^2] \\ &\quad \times [\mathbf{p}_1^2|D|^2 + \mathbf{p}_2^2|F|^2 - \mathrm{Re}(CF^*) + \mathrm{Re}(DE^*) \\ &\quad + 2\mathbf{p}_1\cdot\mathbf{p}_2\,\mathrm{Re}(DF^*)]\},\end{aligned} \tag{3.31}$$

with

$$\begin{aligned} C(E_1, E_2) &= \left(h_2 E_2 - \frac{2}{3}h_8\mathbf{p}_2^2\right)\frac{1}{\Delta_{R^*} - \Delta_{\Sigma_c^{*++}} - E_2 + i\Gamma_{\Sigma_c^{*++}}/2} \\ &\quad + \frac{2}{3}h_8\mathbf{p}_1\cdot\mathbf{p}_2\left(\frac{1}{\Delta_{R^*} - \Delta_{\Sigma_c^0} - E_1 + i\Gamma_{\Sigma_c^0}/2} + \frac{2}{\Delta_{R^*} - \Delta_{\Sigma_c^{*0}} - E_1 + i\Gamma_{\Sigma_c^{*0}}/2}\right), \\ D(E_1, E_2) &= \frac{2}{3}h_8\left(-\frac{1}{\Delta_{R^*} - \Delta_{\Sigma_c^0} - E_1 + i\Gamma_{\Sigma_c^0}/2} + \frac{1}{\Delta_{R^*} - \Delta_{\Sigma_c^{*0}} - E_1 + i\Gamma_{\Sigma_c^{*0}}/2}\right), \\ E(E_1, E_2; \Delta_{\Sigma_c^{(*)0}}, \Delta_{\Sigma_c^{(*)++}}) &= C(E_2, E_1; \Delta_{\Sigma_c^{(*)++}}, \Delta_{\Sigma_c^{(*)0}}), \\ F(E_1, E_2; \Delta_{\Sigma_c^{(*)0}}, \Delta_{\Sigma_c^{(*)++}}) &= -D(E_2, E_1; \Delta_{\Sigma_c^{(*)++}}, \Delta_{\Sigma_c^{(*)0}}), \end{aligned} \tag{3.32}$$

where $\Delta_{R^*} = m_{\Lambda_c(2625)} - m_{\Lambda_c}$.

The total widths of the $\Lambda_c(2593)$ and $\Lambda_c(2625)$ states obtained after integrating out the variables E_1 and E_2 and including the $\pi^0\pi^0$ channel are

TABLE VI. Same as Table V except for p-wave charmed baryons.

Decay	Expt. [7]	This work HHChPT	Tawfiq *et al.* [27]	Ivanov *et al.* [28]	Huang *et al.* [29]	Zhu [34]
$\Lambda_c(2593)^+ \to (\Lambda_c^+\pi\pi)_R$	$2.63^{+1.56}_{-1.09}$	input			2.5	
$\Lambda_c(2593)^+ \to \Sigma_c^{++}\pi^-$	$0.65^{+0.41}_{-0.31}$	$0.72^{+0.43}_{-0.30}$	1.47 ± 0.57	0.79 ± 0.09	$0.55^{+1.3}_{-0.55}$	0.64
$\Lambda_c(2593)^+ \to \Sigma_c^0\pi^+$	$0.67^{+0.41}_{-0.31}$	$0.77^{+0.46}_{-0.32}$	1.78 ± 0.70	0.83 ± 0.09	0.89 ± 0.86	0.86
$\Lambda_c(2593)^+ \to \Sigma_c^+\pi^0$		$1.57^{+0.93}_{-0.65}$	1.18 ± 0.46	0.98 ± 0.12	1.7 ± 0.49	1.2
$\Lambda_c(2625)^+ \to \Sigma_c^{++}\pi^-$	<0.10	$\lesssim 0.029$	0.44 ± 0.23	0.076 ± 0.009	0.013	0.011
$\Lambda_c(2625)^+ \to \Sigma_c^0\pi^+$	<0.09	$\lesssim 0.029$	0.47 ± 0.25	0.080 ± 0.009	0.013	0.011
$\Lambda_c(2625)^+ \to \Sigma_c^+\pi^0$		$\lesssim 0.041$	0.42 ± 0.22	0.095 ± 0.012	0.013	0.011
$\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi$	<1.9	$\lesssim 0.21$			0.11	
$\Sigma_c(2800)^{++} \to \Lambda_c\pi, \Sigma_c^{(*)}\pi$	75^{+22}_{-17}	input				
$\Sigma_c(2800)^{+} \to \Lambda_c\pi, \Sigma_c^{(*)}\pi$	62^{+60}_{-40}	input				
$\Sigma_c(2800)^{0} \to \Lambda_c\pi, \Sigma_c^{(*)}\pi$	61^{+28}_{-18}	input				
$\Xi_c(2790)^+ \to \Xi_c^{\prime 0,+}\pi^{+,0}$	<15	$8.0^{+4.7}_{-3.3}$				
$\Xi_c(2790)^0 \to \Xi_c^{\prime +,0}\pi^{-,0}$	<12	$8.5^{+5.0}_{-3.5}$				
$\Xi_c(2815)^+ \to \Xi_c^{*+,0}\pi^{0,+}$	<3.5	$3.4^{+2.0}_{-1.4}$	2.35 ± 0.93	0.70 ± 0.04		
$\Xi_c(2815)^0 \to \Xi_c^{*+,0}\pi^{-,0}$	<6.5	$3.6^{+2.1}_{-1.5}$				

$$\begin{aligned}\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi\pi) &= 13.82h_2^2 + 26.28h_8^2 - 2.97h_2h_8,\\ \Gamma(\Lambda_c(2625)^+ \to \Lambda_c^+\pi\pi) &= 0.617h_2^2 + 0.136\times 10^6 h_8^2 - 27h_2h_8,\end{aligned} \tag{3.33}$$

where use of (3.18) for g_2 has been made. It is clear that the experimental 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 $\Lambda_c(2593)$ is entirely governed by the coupling h_2. This indicates that the direct nonresonant $\Lambda_c^+\pi\pi$ contribution cannot be described by the Σ_c^* pole alone. Some other mechanisms are needed to account for the nonresonant contributions. Identifying the calculated $\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi\pi)$ with the resonant one, we find

$$|h_2| = 0.437^{+0.114}_{-0.102}, \qquad |h_8| < 3.65\times 10^{-3}. \tag{3.34}$$

Comparing (3.34) with (3.23) we see that the magnitude of h_2 is enhanced slightly by finite width effects.

Assuming that the total width of $\Lambda_c(2593)^+$ is saturated by the resonant $\Lambda_c^+\pi\pi$ 3-body decays, Pirjol and Yan obtained $|h_2| = 0.572^{+0.322}_{-0.197}$ and $|h_8| \le (3.50 - 3.68)\times 10^{-3}$ MeV^{-1} [8]. Using the updated hadron masses and $\Gamma(\Lambda_c(2593))$,[6] we find $|h_2| = 0.499^{+0.134}_{-0.100}$. Taking into account the fact that the Σ_c and Σ_c^* poles only describe the resonant contributions to the total width of $\Lambda_c(2593)$, we finally reach at the value of h_2 given by (3.34). Using this result for h_2, the two-body $\Lambda_c(2593) \to \Sigma_c\pi$ rates are shown in Table VI. The three-body partial rates are found to be

[6]The CLEO result $\Gamma(\Lambda_c(2593)) = 3.9^{+2.4}_{-1.6}$ MeV [36] is used in [8] to fix h_2.

$$\begin{aligned}\Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi^+\pi^-) &= 1.29\ \mathrm{MeV},\\ \Gamma(\Lambda_c(2593)^+ \to \Lambda_c^+\pi^0\pi^0) &= 1.34\ \mathrm{MeV}.\end{aligned} \tag{3.35}$$

Therefore, isospin violation is manifested in the relations $\Gamma(\Sigma_c^+\pi^0) \approx 2\Gamma(\Sigma_c^{++}\pi^-) \sim 2\Gamma(\Sigma_c^0\pi^+)$ and $\Gamma(\Lambda_c^+\pi^0\pi^0) \approx \Gamma(\Lambda_c^+\pi^+\pi^-)$ in $\Lambda_c(2593)$ decays.

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 (3.34) and the experimental observation that the $\Xi_c\pi\pi$ mode in $\Xi_c(2815)$ decays is consistent with being entirely via $\Xi_c^*\pi$ [37], the predicted $\Xi_c(2790)$ and $\Xi_c(2815)$ widths are shown in Table VI, where uses have been made of

$$\begin{aligned}\Gamma(\Xi_{c1}(1/2)^+) &\approx \Gamma(\Xi_{c1}(1/2)^+ \to \Xi_c^{\prime +}\pi^0, \Xi_c^{\prime 0}\pi^+)\\ &= \frac{h_2^2}{4\pi f_\pi^2}\bigg(\frac{1}{2}\frac{m_{\Xi_c^{\prime +}}}{m_{\Xi_{c1}(1/2)}}E_\pi^2 p_\pi\\ &\quad + \frac{m_{\Xi_c^{\prime 0}}}{m_{\Xi_{c1}(1/2)}}E_\pi^2 p_\pi\bigg),\\ \Gamma(\Xi_{c1}(3/2)^+) &\approx \Gamma(\Xi_{c1}(3/2)^+ \to \Xi_c^{*+}\pi^0, \Xi_c^{*0}\pi^+)\\ &= \frac{h_2^2}{4\pi f_\pi^2}\bigg(\frac{1}{2}\frac{m_{\Xi_c^{*+}}}{m_{\Xi_{c1}(3/2)}}E_\pi^2 p_\pi\\ &\quad + \frac{m_{\Xi_c^{*0}}}{m_{\Xi_{c1}(3/2)}}E_\pi^2 p_\pi\bigg),\end{aligned} \tag{3.36}$$

and similar expressions for the neutral $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ states.

The predictions are consistent with the current experimental limits.

Some information on the coupling h_{10} can be inferred from the strong decays of $\Sigma_c(2800)$. As noticed in passing, the states $\Sigma_c(2800)^{++,+,0}$ which are observed in the $\Lambda_c^+\pi$ spectrum are most likely to be $\Sigma_{c2}(\frac{3}{2}^-)$. From Table III we see that there are three low-lying p-wave Σ_c multiplets: Σ_{c0}, Σ_{c1} and Σ_{c2}. Both Σ_{c0} and Σ_{c2} decay to $\Lambda_c\pi$ in an S-wave and a D-wave, respectively, while Σ_{c1} decays mainly to the two-pion system $\Lambda_c\pi\pi$ in a P-wave. From Eqs. (3.14), (3.15), and (3.34) we find $\Gamma(\Sigma_{c0}\to\Lambda_c\pi)\approx$ 406 MeV. Hence, it is too broad to be observable. Therefore, $\Sigma_c(2800)^{++,+,0}$ are likely to be $\Sigma_{c2}(\frac{3}{2}^-)$. Assuming their widths are dominated by the two-body D-wave modes $\Lambda_c\pi$, $\Sigma_c\pi$ and $\Sigma_c^*\pi$, we have [8]

$$\begin{aligned}\Gamma(\Sigma_{c2}(3/2)^{++}) \approx\ & \Gamma(\Sigma_{c2}(3/2)^{++}\to\Lambda_c^+\pi^+)\\ &+\Gamma(\Sigma_{c2}(3/2)^{++}\to\Sigma_c^+\pi^+)\\ &+\Gamma(\Sigma_{c2}(3/2)^{++}\to\Sigma_c^{*+}\pi^+)\\ &+\Gamma(\Sigma_{c2}(3/2)^{++}\to\Sigma_c^{++}\pi^0)\\ &+\Gamma(\Sigma_{c2}(3/2)^{++}\to\Sigma_c^{*++}\pi^0),\end{aligned} \tag{3.37}$$

and similar expressions for $\Sigma_{c2}(\frac{3}{2})^+$ and $\Sigma_{c2}(\frac{3}{2})^0$. Using[7]

$$\begin{aligned}\Gamma(\Sigma_{c2}(3/2^-)\to\Lambda_c\pi) &= \frac{4h_{10}^2}{15\pi f_\pi^2}\frac{m_{\Lambda_c}}{m_{\Sigma_{c2}}}p_\pi^5,\\ \Gamma(\Sigma_{c2}(3/2^-)\to\Sigma_c^{(*)}\pi) &= \frac{h_{11}^2}{10\pi f_\pi^2}\frac{m_{\Sigma_c^{(*)}}}{m_{\Sigma_{c2}}}p_\pi^5,\end{aligned} \tag{3.38}$$

and the quark model relation $h_{11}^2=2h_{10}^2$ [cf. Equation (3.16)] and the measured widths of $\Sigma_c(2800)^{++,+,0}$ (Table II), we obtain

$$|h_{10}| = (0.86^{+0.08}_{-0.10})\times 10^{-3}\ \mathrm{MeV}^{-1}. \tag{3.39}$$

This is consistent with the naive expectation that $h_{10}\sim 1/\Lambda_\chi$. Since the state $\Sigma_{c1}(\frac{3}{2}^-)$ is broader, even a small mixing of $\Sigma_{c2}(\frac{3}{2}^-)$ with $\Sigma_{c1}(\frac{3}{2}^-)$ could enhance the decay width of the former [8]. Moreover, the nonresonant three-body mode $\Lambda_c^+\pi\pi$ may have contributions to the width of Σ_{c2}, the above value for h_{10} should be regarded as an upper limit of $|h_{10}|$. Using the quark model relation $|h_8|=|h_{10}|$ [Eq. (3.16)], we then have

$$|h_8| \lesssim (0.86^{+0.08}_{-0.10})\times 10^{-3}\ \mathrm{MeV}^{-1}, \tag{3.40}$$

which improves the previous limit (3.34) by a factor of 4.

Using the above value of h_8, the rates of $\Lambda_c(2625)$ decays to $\Lambda_c\pi\pi$ and $\Sigma_c\pi$ are presented in Table VI, where we have used Eq. (3.33) and

$$\Gamma(\Lambda_{c1}(3/2^-)\to\Sigma_c\pi) = \frac{2h_8^2}{9\pi f_\pi^2}\frac{m_{\Sigma_c}}{m_{\Lambda_{c1}(3/2)}}p_\pi^5. \tag{3.41}$$

C. Strong decays of first positive-parity excited charmed baryons

Besides the p-wave charmed baryons discussed in the previous subsection, some of the higher orbitally excited charmed baryons listed in Table I are likely to be the first positive-parity excitations. For example, the recent Belle studies favor the J^P quantum numbers of $\Lambda_c(2880)^+$ to be $\frac{5}{2}^+$ [4]. The quantum numbers for the first positive-parity excited charmed baryons are listed in Table IV. Those states have $L_K+L_k=2$ and hence the orbital angular momentum L_ℓ can be 2, 1 or 0. Besides $\Lambda_c(2880)^+$, the states $\Lambda_c(2765)^+$, $\Lambda_c(2940)^+$, $\Xi_c(2980)^{+,0}$ and $\Xi_c(3077)^{+,0}$ are also likely to be the first positive-parity excitations of charmed baryons as we are going to discuss.

As noticed in Sec. II, Belle has studied the experimental constraint on the J^P quantum numbers of $\Lambda_c(2880)^+$ and found that the assignment of $J=\frac{5}{2}$ is favored over $J=\frac{1}{2}$ or $\frac{3}{2}$ by the angular analysis of $\Lambda_c(2880)^+\to\Sigma_c^{0,++}\pi^\pm$ [4]. The measurement of the ratio of $\Lambda_c(2880)$ partial widths [4]

$$R=\frac{\Gamma(\Lambda_c(2880)\to\Sigma_c^*\pi^\pm)}{\Gamma(\Lambda_c(2880)\to\Sigma_c\pi^\pm)} = (24.1\pm 6.4^{+1.1}_{-4.5})\% \tag{3.42}$$

can be used to determine the parity assignment. From Tables III and IV we see that the candidates for the spin-$\frac{5}{2}$ state are $\tilde\Lambda_{c2}(\frac{5}{2}^-)$, $\Lambda_{c2}(\frac{5}{2}^+)$, $\tilde\Lambda'_{c2}(\frac{5}{2}^+)$, $\tilde\Lambda''_{c2}(\frac{5}{2}^+)$ and $\tilde\Lambda''_{c3}(\frac{5}{2}^+)$. For $J^P=\frac{5}{2}^-$, $\Lambda_c(2880)$ decays to $\Sigma_c^*\pi$ and $\Sigma_c\pi$ in a D wave. From Eq. (3.17) we obtain

$$\begin{aligned}\frac{\Gamma(\tilde\Lambda_{c2}(5/2^-)\to[\Sigma_c^*\pi]_D)}{\Gamma(\tilde\Lambda_{c2}(5/2^-)\to[\Sigma_c\pi]_D)} &= \frac{7}{2}\frac{p_\pi^5(\Lambda_c(2880)\to\Sigma_c^*\pi)}{p_\pi^5(\Lambda_c(2880)\to\Sigma_c\pi)}\\ &=\frac{7}{2}\times 0.42 = 1.45.\end{aligned} \tag{3.43}$$

Hence, the assignment of $J^P=\frac{5}{2}^-$ for $\Lambda_c(2880)$ is disfavored. For $J^P=\frac{5}{2}^+$, Λ_{c2}, $\tilde\Lambda'_{c2}$ and $\tilde\Lambda''_{c2}$ with $J_\ell=2$ decay to $\Sigma_c\pi$ in a F wave and $\Sigma_c^*\pi$ in F and P waves. Neglecting the P-wave contribution for the moment,

$$\begin{aligned}\frac{\Gamma(\Lambda_{c2}(5/2^+)\to[\Sigma_c^*\pi]_F)}{\Gamma(\Lambda_{c2}(5/2^+)\to[\Sigma_c\pi]_F)} &= \frac{4}{5}\frac{p_\pi^7(\Lambda_c(2880)\to\Sigma_c^*\pi)}{p_\pi^7(\Lambda_c(2880)\to\Sigma_c\pi)}\\ &=\frac{4}{5}\times 0.29 = 0.23.\end{aligned} \tag{3.44}$$

At first glance, it appears that this is in good agreement with experiment. However, the $\Sigma_c^*\pi$ channel is available via a P-wave and is enhanced by a factor of $1/p_\pi^4$ (or more precisely, $(\Lambda_\chi/p_\pi)^4$) relative to the F-wave one. Unfortunately, we cannot apply Eq. (3.17) to calculate the contribution of the $[\Sigma_c^*\pi]_F$ channel to the ratio R as the reduced matrix elements are different for P-wave and

[7]It is useful to apply Eq. (3.17) to check the consistency of the partial decay rate formulas.

F-wave modes. In any event, the $\Sigma_c^*\pi$ mode produced in $\Lambda_c(2880)$ is *a priori* not necessarily suppressed relative to $[\Sigma_c\pi]_F$. Therefore, if $\Lambda_c(2880)^+$ is one of the states Λ_{c2}, $\tilde{\Lambda}'_{c2}$ and $\tilde{\Lambda}''_{c2}$, the prediction $R = 0.23$ is not robust as it can be easily upset by the contribution from the P-wave $\Sigma_c^*\pi$.

As for $\tilde{\Lambda}''_{c3}(\frac{5}{2}^+)$, it decays to $\Sigma_c^*\pi$, $\Sigma_c\pi$ and $\Lambda_c\pi$ all in F waves. Since $J_\ell = 3$, $L_\ell = 2$, it turns out that

$$\frac{\Gamma(\Lambda''_{c3}(5/2^+) \to [\Sigma_c^*\pi]_F)}{\Gamma(\Lambda''_{c3}(5/2^+) \to [\Sigma_c\pi]_F)} = \frac{5}{4}\frac{p_\pi^7(\Lambda_c(2880) \to \Sigma_c^*\pi)}{p_\pi^7(\Lambda_c(2880) \to \Sigma_c\pi)} = \frac{5}{4} \times 0.29 = 0.36. \tag{3.45}$$

Although this deviates from the experimental measurement (3.42) by 1σ, it is a robust prediction. However, there are two issues with this assignment. First, $\Lambda''_{c3}(\frac{5}{2}^+)$ can decay to a F-wave $\Lambda_c\pi$ and this has not been seen by *BABAR* and Belle. Second, the quark model indicates a $\Lambda_{c2}(\frac{5}{2}^+)$ state around 2910 MeV which is close to the mass of $\Lambda_c(2880)$, while the mass of $\Lambda''_{c3}(\frac{5}{2}^+)$ is higher [9]. Therefore, we conjecture that the first positive-parity excited charmed baryon $\Lambda_c(2880)^+$ could be an admixture of $\Lambda_{c2}(\frac{5}{2}^+)$ and $\Lambda''_{c3}(\frac{5}{2}^+)$.

The quark potential model predicts a $\frac{5}{2}^-$ Λ_c state at 2900 MeV and a $\frac{3}{2}^+$ Λ_c state at 2910 MeV [9]. Given the uncertainty of order 50 MeV for the quark model calculation, this suggests that the possible allowed J^P numbers of the highest $\Lambda_c(2940)^+$ are $\frac{5}{2}^-$ and $\frac{3}{2}^+$. Hence, the potential candidates are $\tilde{\Lambda}_{c2}(\frac{5}{2}^-)$, $\Lambda_{c2}(\frac{3}{2}^+)$, $\tilde{\Lambda}'_{c1}(\frac{3}{2}^+)$,

TABLE VII. Possible strong decays of the first positive-parity excitations of the Ξ_c, where L denotes the orbital angular momentum of the light meson(s). The final state Σ_c^*K is kinematically allowed for $\Xi_c(3077)$.

J^P	Diquark transition	L	Decay channel	Final states
$\frac{1}{2}^+$	$0^+ \to 1^+ + 0^-$	1	$\frac{1}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^-$	$\Xi'_c\pi$, $\Xi_c^*\pi$, $\Sigma_c^{(*)}K$
	$0^+ \to 0^+ + 0^- + 0^-$	0	$\frac{1}{2}^+ \to \frac{1}{2}^+ + 0^- + 0^-$	$\Xi_c\pi\pi$, $\Lambda_cK\pi$
	$1^+ \to 0^+ + 0^-$	1	$\frac{1}{2}^+ \to \frac{1}{2}^+ + 0^-$	$\Xi_c\pi$, Λ_cK, $D\Lambda$
	$1^+ \to 1^+ + 0^-$	1	$\frac{1}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^-$	$\Xi'_c\pi$, $\Xi_c^*\pi$, $\Sigma_c^{(*)}K$
	$1^+ \to 1^- + 0^-$	0	$\frac{1}{2}^- \to \frac{1}{2}^+ + 0^-$	$\Xi_c(2790)\pi$, $\Lambda_c(2593)K$
	$1^+ \to 1^- + 0^-$	2	$\frac{1}{2}^- \to \frac{3}{2}^+ + 0^-$	$\Xi_c(2815)\pi$, $\Lambda_c(2625)K$
	$1^+ \to 0^+ + 0^- + 0^-$	2	$\frac{1}{2}^+ \to \frac{1}{2}^+ + 0^- + 0^-$	$\Xi_c\pi\pi$, $\Lambda_cK\pi$
	$1^+ \to 1^+ + 0^- + 0^-$	0	$\frac{1}{2}^+ \to \frac{1}{2}^+ + 0^- + 0^-$	$\Xi'_c\pi\pi$
	$1^+ \to 1^+ + 0^- + 0^-$	2	$\frac{1}{2}^+ \to \frac{3}{2}^+ + 0^- + 0^-$	$\Xi_c^*\pi\pi$
$\frac{3}{2}^+$	$1^+ \to 0^+ + 0^-$	1	$\frac{3}{2}^+ \to \frac{1}{2}^+ + 0^-$	$\Xi_c\pi$, Λ_cK, $D\Lambda$
	$1^+ \to 1^+ + 0^-$	1	$\frac{3}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^-$	$\Xi'_c\pi$, $\Xi_c^*\pi$, $\Sigma_c^{(*)}K$
	$1^+ \to 1^- + 0^-$	0	$\frac{3}{2}^- \to \frac{3}{2}^+ + 0^-$	$\Xi_c(2815)\pi$, $\Lambda_c(2625)K$
	$1^+ \to 1^- + 0^-$	2	$\frac{3}{2}^- \to \{\frac{1}{2}^-, \frac{3}{2}^+\} + 0^-$	$\Xi_c(2790, 2815)\pi$, $\Lambda_c(2593, 2625)K$
	$1^+ \to 0^+ + 0^- + 0^-$	2	$\frac{3}{2}^+ \to \frac{1}{2}^+ + 0^- + 0^-$	$\Xi_c\pi\pi$, $\Lambda_cK\pi$
	$1^+ \to 1^+ + 0^- + 0^-$	0	$\frac{3}{2}^+ \to \frac{3}{2}^+ + 0^- + 0^-$	$\Xi_c^*\pi\pi$
	$1^+ \to 1^+ + 0^- + 0^-$	2	$\frac{3}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^- + 0^-$	$\Xi'_c\pi\pi$, $\Xi_c^*\pi\pi$
	$2^+ \to 1^+ + 0^-$	1	$\frac{3}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^-$	$\Xi'_c\pi$, $\Xi_c^*\pi$, $\Sigma_c^{(*)}K$
	$2^+ \to 1^- + 0^-$	2	$\frac{3}{2}^- \to \{\frac{1}{2}^-, \frac{3}{2}^+\} + 0^-$	$\Xi_c(2790, 2815)\pi$, $\Lambda_c(2593, 2625)K$
	$2^+ \to 0^+ + 0^- + 0^-$	2	$\frac{3}{2}^+ \to \frac{1}{2}^+ + 0^- + 0^-$	$\Xi_c\pi\pi$, $\Lambda_cK\pi$
	$2^+ \to 1^+ + 0^- + 0^-$	2	$\frac{3}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^- + 0^-$	$\Xi'_c\pi\pi$, $\Xi_c^*\pi\pi$
$\frac{5}{2}^+$	$2^+ \to 1^+ + 0^-$	1	$\frac{5}{2}^+ \to \frac{3}{2}^+ + 0^-$	$\Xi_c^*\pi$, Σ_c^*K
	$2^+ \to 1^+ + 0^-$	3	$\frac{5}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^-$	$\Xi'_c\pi$, $\Xi_c^*\pi$, $\Sigma_c^{(*)}K$
	$2^+ \to 1^- + 0^-$	2	$\frac{5}{2}^- \to \{\frac{1}{2}^-, \frac{3}{2}^+\} + 0^-$	$\Xi_c(2790, 2815)\pi$, $\Lambda_c(2593, 2625)K$
	$2^+ \to 0^+ + 0^- + 0^-$	2	$\frac{5}{2}^+ \to \frac{1}{2}^+ + 0^- + 0^-$	$\Xi_c\pi\pi$, $\Lambda_cK\pi$
	$2^+ \to 1^+ + 0^- + 0^-$	2	$\frac{5}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^- + 0^-$	$\Xi'_c\pi\pi$, $\Xi_c^*\pi\pi$
	$3^+ \to 0^+ + 0^-$	3	$\frac{5}{2}^+ \to \frac{1}{2}^+ + 0^-$	$\Xi_c\pi$, Λ_cK, $D\Lambda$
	$3^+ \to 1^+ + 0^-$	3	$\frac{5}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^-$	$\Xi'_c\pi$, $\Xi_c^*\pi$, $\Sigma_c^{(*)}K$
	$3^+ \to 1^- + 0^-$	2	$\frac{5}{2}^- \to \{\frac{1}{2}^-, \frac{3}{2}^+\} + 0^-$	$\Xi_c(2790, 2815)\pi$, $\Lambda_c(2593, 2625)K$
	$3^+ \to 1^+ + 0^- + 0^-$	2	$\frac{5}{2}^+ \to \{\frac{1}{2}^+, \frac{3}{2}^+\} + 0^- + 0^-$	$\Xi'_c\pi\pi$, $\Xi_c^*\pi\pi$

$\tilde{\Lambda}''_{c1}(\frac{3}{2}^+)$ and $\tilde{\Lambda}''_{c2}(\frac{3}{2}^+)$. Ratios of $\Lambda_c(2940)$ partial widths are expected in HHChPT to be

$$\frac{\Gamma(\tilde{\Lambda}_{c2}(5/2^-)\to[\Sigma_c^*\pi]_D)}{\Gamma(\tilde{\Lambda}_{c2}(5/2^-)\to[\Sigma_c\pi]_D)}=\frac{7}{2}\frac{p_\pi^5(\Lambda_c(2940)\to\Sigma_c^*\pi)}{p_\pi^5(\Lambda_c(2940)\to\Sigma_c\pi)}=\frac{7}{2}\times 0.48=1.68,$$

$$\frac{\Gamma(\Lambda_{c2}(3/2^+)\to[\Sigma_c^*\pi]_P)}{\Gamma(\Lambda_{c2}(3/2^+)\to[\Sigma_c\pi]_P)}=\frac{1}{5}\frac{p_\pi^3(\Lambda_c(2940)\to\Sigma_c^*\pi)}{p_\pi^3(\Lambda_c(2940)\to\Sigma_c\pi)}=\frac{1}{5}\times 0.65=0.13,$$

$$\frac{\Gamma(\tilde{\Lambda}'_{c1}(3/2^+)\to[\Sigma_c^*\pi]_P)}{\Gamma(\tilde{\Lambda}'_{c1}(3/2^+)\to[\Sigma_c\pi]_P)}=5\frac{p_\pi^3(\Lambda_c(2940)\to\Sigma_c^*\pi)}{p_\pi^3(\Lambda_c(2940)\to\Sigma_c\pi)}=5\times 0.65=3.25. \quad (3.46)$$

Since the predicted ratios differ significantly for different J^P quantum numbers, the measurements of the ratio of $\Sigma_c^*\pi/\Sigma_c\pi$ will enable us to discriminate the J^P assignments for $\Lambda_c(2940)$.

For the charmed states $\Xi_c(2980)$ and $\Xi_c(3077)$, they could be the first positive-parity excitations of Ξ_c in viewing of their large masses. 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 $\Xi_c(3077)$ 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 an even-parity excitation as the quark model [9] and the Skyrme model [13] suggest a $J^P=\frac{1}{2}^+$ state with a mass 2742 and 2775 MeV, respectively. It is thus tempting to assign $J^P=\frac{1}{2}^+$ for $\Xi_c(2980)$ and $\frac{5}{2}^+$ for $\Xi_c(3077)$. Of course, the assignment of $J^P=\frac{3}{2}^+$ is also possible. The possible strong decays of the first positive-parity excitations of the Ξ_c states are summarized in Table VII. Since the two-body modes $\Xi_c\pi$, Λ_cK, $\Xi'_c\pi$ and Σ_cK are in P (F) waves and the three-body modes $\Xi_c\pi\pi$ and $\Lambda_cK\pi$ are in S (D) waves in the decays of $\frac{1}{2}^+$ ($\frac{5}{2}^+$), this explains why $\Xi_c(2980)$ is broader than $\Xi_c(3077)$. Since both $\Xi_c(2980)$ and $\Xi_c(3077)$ are above the $D\Lambda$ threshold, it is important to search for them in the $D\Lambda$ spectrum as well.

IV. CONCLUSIONS

Strong decays of charmed baryons are analyzed in the framework of heavy hadron chiral perturbation theory in which heavy quark symmetry and chiral symmetry are synthesized. Our main conclusions are the following:

(i) For s-wave charmed baryons, we use the channel $\Sigma_c^{++}\to\Lambda_c^+\pi^+$ to fix the coupling constant g_2. The value of $|g_2|=0.591\pm 0.023$ are in good agreement with the quark model expectation. The predictions for the strong decays $\Sigma_c^*\to\Lambda\pi$ and $\Xi_c^*\to\Xi_c\pi$ are in excellent agreement with experiment.

(ii) For $L=1$ orbitally excited baryons, two of the unknown couplings, namely, h_2 and h_{10}, are determined from the resonant $\Sigma_c^+\pi\pi$ mode produced in the $\Lambda_c(2593)$ decay and the width of $\Sigma_c(2800)$, respectively. The results are $|h_2|=0.437^{+0.114}_{-0.102}$ and $|h_{10}|\lesssim(0.86^{+0.08}_{-0.10})\times 10^{-3}$ MeV^{-1}. Since the two-pion system $\Lambda_c^+\pi\pi$ in $\Lambda_c(2593)^+$ decays receives nonresonant contributions, our value for h_2 is smaller than the previous estimates. Applying the quark model relation $h_8^2=h_{10}^2$, predictions for the strong decays of other p-wave charmed baryons such as $\Lambda_c(2625)\to\Sigma_c\pi$, $\Lambda_c\pi\pi$, $\Xi_c(2790)\to\Xi'_c\pi$ and $\Xi_c(2815)\to\Xi_c^*\pi$ are presented in Table VI. Since the decays $\Lambda_c(2593)\to\Lambda_c\pi\pi$ and $\Lambda_c(2593)\to\Sigma_c\pi$ occur very close to the threshold, they are very sensitive to the pion's mass and hence isospin symmetry is violated, for example, $\Gamma(\Sigma_c^+\pi^0)\approx 2\Gamma(\Sigma_c^{++}\pi^-)$ and $\Gamma(\Lambda_c^+\pi^0\pi^0)\approx\Gamma(\Lambda_c^+\pi^+\pi^-)$ in $\Lambda_c(2593)$ decays.

(iii) We have examined the first positive-parity excited charmed baryons. We conjecture that the state $\Lambda_c(2880)$ with $J^P=\frac{5}{2}^+$ is an admixture of $\Lambda_{c2}(\frac{5}{2}^+)$ with an $\tilde{\Lambda}''_{c3}(\frac{5}{2}^+)$; both are $L=2$ orbitally excited states. Potential models suggest the possible allowed J^P numbers of the $\Lambda_c(2940)^+$ to be $\frac{5}{2}^-$ and $\frac{3}{2}^+$. We have demonstrated that the measurements of the ratio of $\Sigma_c^*\pi/\Sigma_c\pi$ will enable us to discriminate the J^P assignments for $\Lambda_c(2940)$. We advocate that the J^P quantum numbers of $\Xi_c(2980)$ and $\Xi_c(3077)$ are $\frac{1}{2}^+$ and $\frac{5}{2}^+$, respectively. Under this J^P assignment, it is easy to understand why $\Xi_c(2980)$ is broader than $\Xi_c(3077)$.

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