

Charmed baryons circa 2015

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This is essentially an update of Ref. [1] [H. Y. Cheng, *Int. J. Mod. Phys. A* 24 (Suppl. 1), 593 (2009)], a review of charmed baryon physics around 2007. Topics covered in this review include the spectroscopy, strong decays, lifetimes, nonleptonic and semileptonic weak decays, and electromagnetic decays of charmed baryons.

Keywords charmed baryons, spectroscopy, decays of baryons, chiral Lagrangian

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1 Introduction

1 Charm baryon spectroscopy provides an excellent ground
2 for studying the dynamics of light quarks in the envi-
3 ronment of a heavy quark. In the past decade, many
4 new excited charmed baryon states have been discov-
5 ered by BaBar, Belle, CLEO and LHCb. B decays and
6 the $e^+e^- \rightarrow c\bar{c}$ continuum are both very rich sources
7 of charmed baryons. Many efforts have been made to
8 identify the quantum numbers of these new states and
9 understand their properties.

10 Consider the strong decays $\Sigma_Q \rightarrow \Lambda_Q\pi$ and $\Xi_Q^* \rightarrow$
11 $\Xi_Q\pi$, where $Q = c, b$. The mass differences between Σ_c
12 and Λ_c and between Ξ_c^* and Ξ_c in the charmed baryon
13 sector were found to be large enough that the strong
14 decays of Σ_c and Ξ_c^* are kinematically allowed. Conse-
15 quently, the charmed baryon system offers a unique and
16 excellent laboratory for testing the ideas and predictions
17 of heavy quark symmetry of the heavy quark and chi-
18 ral symmetry of the light quarks. These tests will have
19 interesting implications for the low-energy dynamics of
20 heavy baryons interacting with Goldstone bosons.

21 Theoretical interest in hadronic weak decays of
22 charmed baryons peaked around the early 1990s and
23 then faded away. To date, we still do not have a good
24 phenomenological model, not mentioning the quantum
25 chromodynamics (QCD)-inspired approach as in meson
26 decays, to describe the complicated physics of baryon
27 decays. We need cooperative efforts from both experi-
28 mentalists and theorists to make progress in this arena.

29 This review is essentially an update of Ref. [1], which
30 described charmed baryon physics around 2007. The out-
31 line of the content is the same as that of Ref. [1] except

*Special Topic: Potential Physics at a Super τ -Charm Factory (Ed. Hai-Bo Li).

that we add discussions of the spectroscopy and lifetimes of doubly charmed baryons.

Several excellent review articles on charmed baryons can be found in Refs. [2–7].

2 Spectroscopy

2.1 Singly charmed baryons

The singly charmed baryon is composed of a charmed quark and two light quarks. Each light quark is a triplet of flavor SU(3). There are two different SU(3) multiplets of charmed baryons: a symmetric sextet **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 , $\Xi_c^{\prime+}$, $\Xi_c^{\prime0}$ and $\Sigma_c^{++,+0}$ form a **6** representation; among them, only Ω_c^0 decays weakly. We follow the Particle Data Group’s (PDG’s) convention [8] of using a prime to distinguish the Ξ_c in the **6** from that in the $\bar{\mathbf{3}}$.

In the quark model, the orbital angular momentum of the light diquark can be decomposed into $\mathbf{L}_\ell = \mathbf{L}_\rho + \mathbf{L}_\lambda$, where \mathbf{L}_ρ is the orbital angular momentum between the two light quarks, and \mathbf{L}_λ is the orbital angular momentum between the diquark and the charmed quark. 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 $\mathbf{J}_\ell = \mathbf{S}_\ell + \mathbf{L}_\ell$, are separately conserved. The total angular momentum is given by $\mathbf{J} = \mathbf{S}_c + \mathbf{J}_\ell$. It is convenient to use S_ℓ , L_ℓ , and J_ℓ to enumerate the spectrum of states. Moreover, one can define two independent relative momenta, $\mathbf{p}_\rho = \frac{1}{\sqrt{2}}(\mathbf{p}_1 - \mathbf{p}_2)$ and $\mathbf{p}_\lambda = \frac{1}{\sqrt{6}}(\mathbf{p}_1 + \mathbf{p}_2 - 2\mathbf{p}_c)$, from the two light quark momenta \mathbf{p}_1 , \mathbf{p}_2 and the heavy quark momentum \mathbf{p}_c . Denoting the quantum numbers L_ρ and L_λ as the eigenvalues of \mathbf{L}_ρ^2 and \mathbf{L}_λ^2 , respectively, the ρ -orbital momentum L_ρ describes relative orbital exci-

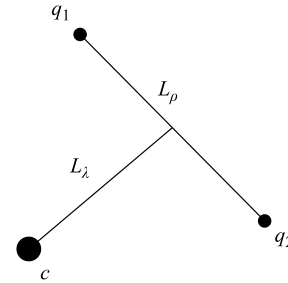


Fig. 1 Singly charmed baryon where L_ρ describes relative orbital excitation of the two light quarks and L_λ the orbital excitation of the center of the mass of the two light quarks relative to the charmed quark.

tations of the two light quarks, and the λ -orbital momentum L_λ describes orbital excitations of the center of mass of the two light quarks relative to the heavy quark (see Fig. 1). The p -wave heavy baryon can be in either the $(L_\rho = 0, L_\lambda = 1)$ λ -state or the $(L_\rho = 1, L_\lambda = 0)$ ρ -state. The orbital λ -state (ρ -state) is clearly symmetric (antisymmetric) under the interchange of \mathbf{p}_1 and \mathbf{p}_2 . In the following, we use the notation $\mathcal{B}_{cJ_\ell}(J^P)$ ($\tilde{\mathcal{B}}_{cJ_\ell}(J^P)$) to denote the states that are 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, the quantum numbers of which are listed in Table 1.

The next orbitally excited states are the positive-parity excitations with $L_\rho + L_\lambda = 2$. There exist multiplets (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 2). We use a hat to distinguish them. Because the orbital $L_\lambda = L_\rho = 1$ states are antisymmetric under the interchange of two light quarks, we use a tilde to denote them. Moreover, the notation $\tilde{\mathcal{B}}_{cJ_\ell}^{L_\ell}(J^P)$ is reserved for tilde states in the $\bar{\mathbf{3}}$, as the quantum number L_ℓ is needed to distinguish different states.

The observed mass spectra and decay widths of charmed baryons are summarized in Table 3 (see also Fig. 2). Note that except for the parity of the lightest Λ_c^+ and the heavier one $\Lambda_c(2880)^+$, none of the other J^P

Table 1 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. The orbital ρ -states with $L_\rho = 1$ and $L_\lambda = 0$ have odd orbital wave functions under the permutation of the two light quarks and are denoted by a tilde.

| 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 ⁻ | $\Sigma_{c0}(\frac{1}{2}^-)$ | 6 | 1 | 1(0,1) | 0 ⁻ |
| $\tilde{\Lambda}_{c0}(\frac{1}{2}^-)$ | $\bar{\mathbf{3}}$ | 1 | 1(1,0) | 0 ⁻ | $\Sigma_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | 6 | 1 | 1(0,1) | 1 ⁻ |
| $\tilde{\Lambda}_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | $\bar{\mathbf{3}}$ | 1 | 1(1,0) | 1 ⁻ | $\Sigma_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ | 6 | 1 | 1(0,1) | 2 ⁻ |
| $\tilde{\Lambda}_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ | $\bar{\mathbf{3}}$ | 1 | 1(1,0) | 2 ⁻ | $\tilde{\Sigma}_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | 6 | 0 | 1(1,0) | 1 ⁻ |
| $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | $\bar{\mathbf{3}}$ | 0 | 1(0,1) | 1 ⁻ | $\Xi'_{c0}(\frac{1}{2}^-)$ | 6 | 1 | 1(0,1) | 0 ⁻ |
| $\tilde{\Xi}_{c0}(\frac{1}{2}^-)$ | $\bar{\mathbf{3}}$ | 1 | 1(1,0) | 0 ⁻ | $\Xi'_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | 6 | 1 | 1(0,1) | 1 ⁻ |
| $\tilde{\Xi}_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | $\bar{\mathbf{3}}$ | 1 | 1(1,0) | 1 ⁻ | $\Xi'_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ | 6 | 1 | 1(0,1) | 2 ⁻ |
| $\tilde{\Xi}_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ | $\bar{\mathbf{3}}$ | 1 | 1(1,0) | 2 ⁻ | $\tilde{\Xi}'_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ | 6 | 0 | 1(1,0) | 1 ⁻ |

Table 2 The first positive-parity excitations of charmed baryons denoted by $\mathcal{B}_{cJ_\ell}(J^P)$, $\hat{\mathcal{B}}_{cJ_\ell}(J^P)$ and $\tilde{\mathcal{B}}_{cJ_\ell}^{L_\ell}(J^P)$. Orbital $L_\rho = L_\lambda = 1$ states with antisymmetric orbital wave functions are denoted by a tilde. States with the symmetric orbital wave functions $L_\rho = 2$ and $L_\lambda = 0$ are denoted by a hat. For convenience, we drop the superscript L_ℓ for tilde states in the sextet.

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

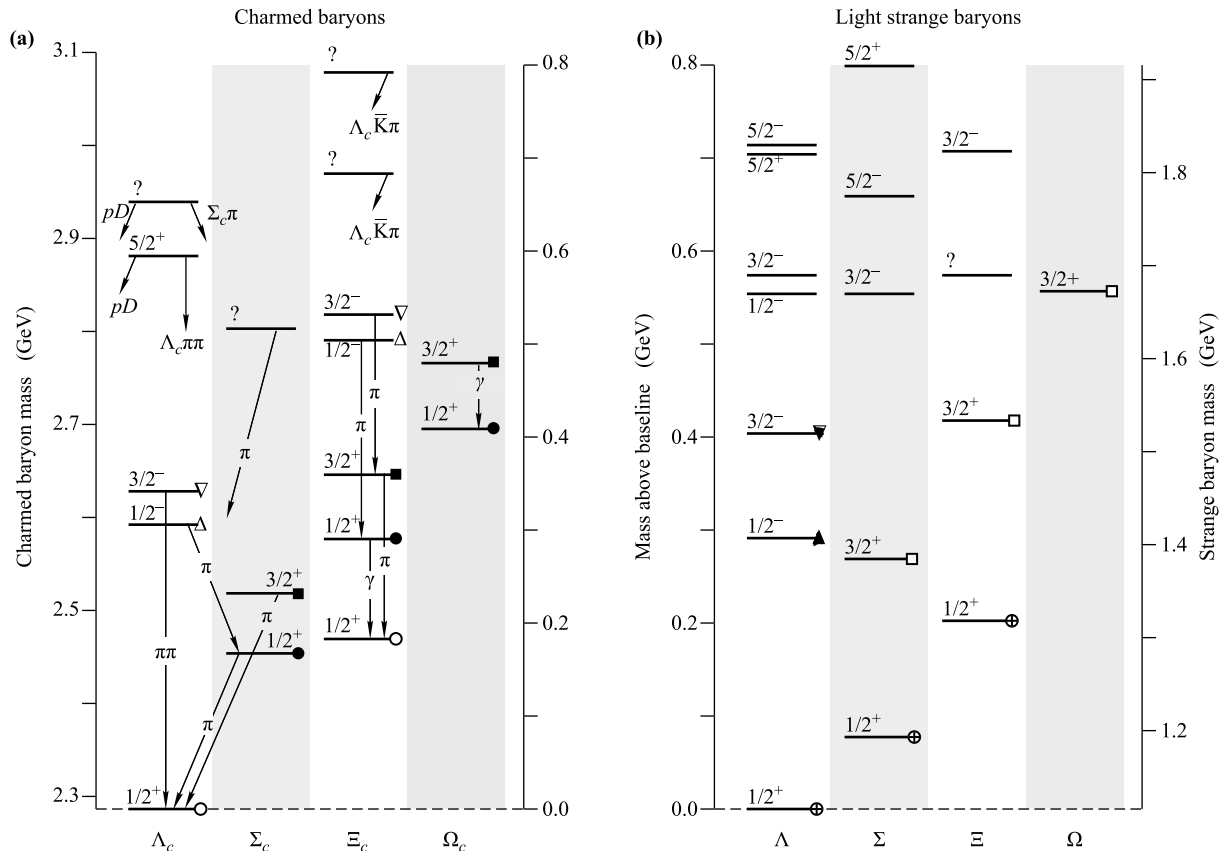


Fig. 2 Charmed baryons and their excitations [8].

Table 3 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 [11] for average. The width of $\Xi_c(2645)^+$ is taken from Ref. [12]. For $\Xi_c(3055)^0$, we quote the preliminary result from Belle [13].

| 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 |
| $\Lambda_c(2595)^+$ | $\frac{3}{2}^-$ | 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^+ | $2467.8^{+0.4}_{-0.6}$ | | weak |
| Ξ_c^0 | $\frac{1}{2}^+$ | 0 | 0 | 0^+ | $2470.88^{+0.34}_{-0.80}$ | | weak |
| $\Xi_c'^+$ | $\frac{3}{2}^+$ | 1 | 0 | 1^+ | 2575.6 ± 3.1 | | $\Xi_c\gamma$ |
| $\Xi_c'^0$ | $\frac{3}{2}^+$ | 1 | 0 | 1^+ | 2577.9 ± 2.9 | | $\Xi_c\gamma$ |
| $\Xi_c(2645)^+$ | $\frac{3}{2}^+$ | 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 | | weak |
| $\Omega_c(2770)^0$ | $\frac{3}{2}^+$ | 1 | 0 | 1^+ | 2765.9 ± 2.0 | | $\Omega_c\gamma$ |

quantum numbers given in Table 3 has been measured. One has to rely on the quark model to determine the spin-parity assignments.

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

2.1.1 Λ_c states

The lowest-lying p -wave Λ_c states are $\tilde{\Lambda}_{c0}(\frac{1}{2}^-)$, $\Lambda_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, $\tilde{\Lambda}_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ and $\tilde{\Lambda}_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$. The doublet $\Lambda_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ is formed by $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ [14]. The allowed strong decays are

$\Lambda_{c1}(1/2^-) \rightarrow [\Sigma_c\pi]_S, [\Sigma_c^*\pi]_D$ and $\Lambda_{c1}(3/2^-) \rightarrow [\Sigma_c\pi]_D, [\Sigma_c^*\pi]_{S,D}, [\Lambda_c\pi\pi]_P$. This explains why the width of $\Lambda_c(2625)^+$ is narrower than that of $\Lambda_c(2595)^+$. Because of isospin conservation in strong decays, Λ_{c1}^+ is not allowed to decay into $\Lambda_c^+\pi^0$.

$\Lambda_c(2765)^+$ is a broad state first seen in $\Lambda_c^+\pi^+\pi^-$ by CLEO [15]. However, it is still not known whether it is a Λ_c^+ or a Σ_c^+ and whether the width might be due to overlapping states. The Skyrme model [16] and the quark model [17, 18] suggest a $J^P = \frac{1}{2}^+$ Λ_c state with a mass of 2742 and 2775 MeV, respectively. Therefore, $\Lambda_c(2765)^+$ could be the first positive-parity excitation of Λ_c . In the

diquark model, it has also been proposed 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 [19].

The state $\Lambda_c(2880)^+$, first observed by CLEO [15] in $\Lambda_c^+\pi^+\pi^-$, was also seen by BaBar in the D^0p spectrum [20]. Belle studied the experimental constraint on the J^P quantum numbers of $\Lambda_c(2880)^+$ [21] and found that $J^P = \frac{5}{2}^+$ is favored by the angular analysis of $\Lambda_c(2880)^+ \rightarrow \Sigma_c^{0,++}\pi^\pm$ together with the ratio $\Sigma_c^*\pi/\Sigma_c\pi$, which was measured to be

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

In the quark model, the candidates for the parity-even spin- $\frac{5}{2}$ state are $\Lambda_{c2}(\frac{5}{2}^+)$, $\hat{\Lambda}_{c2}(\frac{5}{2}^+)$, $\tilde{\Lambda}_{c2}^1(\frac{5}{2}^+)$, $\tilde{\Lambda}_{c2}^2(\frac{5}{2}^+)$, and $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ (see Table 2). The first four candidates, with $J_\ell = 2$, decay to $\Sigma_c\pi$ in an F wave and to $\Sigma_c^*\pi$ in F and P waves. Neglecting the P -wave contribution for the moment,

$$\begin{aligned} \frac{\Gamma(\Lambda_{c2}(5/2^+) \rightarrow [\Sigma_c^*\pi]_F)}{\Gamma(\Lambda_{c2}(5/2^+) \rightarrow [\Sigma_c\pi]_F)} &= \frac{4}{5} \frac{p_\pi^7(\Lambda_c(2880) \rightarrow \Sigma_c^*\pi)}{p_\pi^7(\Lambda_c(2880) \rightarrow \Sigma_c\pi)} \\ &= \frac{4}{5} \times 0.29 = 0.23, \end{aligned} \quad (2.2)$$

where the factor of $4/5$ follows from heavy quark symmetry. At first glance, this appears to be in good agreement with experiment. However, the $\Sigma_c^*\pi$ channel is available via a P -wave and is enhanced by a factor of $1/p_\pi^4$ relative to the F -wave one. However, heavy quark symmetry cannot be applied to calculate the contribution of the $[\Sigma_c^*\pi]_F$ channel to the ratio R , as the reduced matrix elements for the P -wave and F -wave modes differ. In this case, one has to rely on a phenomenological model to compute the ratio R . As for $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$, it decays to $\Sigma_c^*\pi$, $\Sigma_c\pi$ and $\Lambda_c\pi$ all in F waves. Further,

$$\begin{aligned} \frac{\Gamma(\tilde{\Lambda}_{c3}^2(5/2^+) \rightarrow [\Sigma_c^*\pi]_F)}{\Gamma(\tilde{\Lambda}_{c3}^2(5/2^+) \rightarrow [\Sigma_c\pi]_F)} &= \frac{5}{4} \frac{p_\pi^7(\Lambda_c(2880) \rightarrow \Sigma_c^*\pi)}{p_\pi^7(\Lambda_c(2880) \rightarrow \Sigma_c\pi)} \\ &= \frac{5}{4} \times 0.29 = 0.36. \end{aligned} \quad (2.3)$$

Although this deviates from the experimental measurement (2.1) by 1σ , it is a robust prediction. This has motivated us to conjecture that the first positive-parity excited charmed baryon $\Lambda_c(2880)^+$ could be an admixture of $\Lambda_{c2}(\frac{5}{2}^+)$, $\hat{\Lambda}_{c2}(\frac{5}{2}^+)$ and $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ [10].

It is worth mentioning that the Peking group [22] has studied the strong decays of charmed baryons on the basis of the so-called 3P_0 recombination model. For the $\Lambda_c(2880)$, the Peking group found that (i) $\Lambda_c(2880)$ cannot be a radial excitation, as its decay into D^0p is prohibited in the 3P_0 model if $\Lambda_c(2880)$ is the first radial excitation of Λ_c , and (ii) the states $\Lambda_{c2}(\frac{5}{2}^+)$, $\tilde{\Lambda}_{c2}^1(\frac{5}{2}^+)$ and $\hat{\Lambda}_{c2}(\frac{5}{2}^+)$ are excluded, as they do not decay to D^0p according to the 3P_0 model. Moreover, the predicted ratios $\Sigma_c^*\pi/\Sigma_c\pi$ are either too large or too small compared to experiment, for example,

$$\begin{aligned} \frac{\Gamma(\Lambda_{c2}(5/2^+) \rightarrow \Sigma_c^*\pi)}{\Gamma(\Lambda_{c2}(5/2^+) \rightarrow \Sigma_c\pi)} &= 89, \\ \frac{\Gamma(\hat{\Lambda}_{c2}(5/2^+) \rightarrow \Sigma_c^*\pi)}{\Gamma(\hat{\Lambda}_{c2}(5/2^+) \rightarrow \Sigma_c\pi)} &= 0.75. \end{aligned} \quad (2.4)$$

Both symmetric states Λ_{c2} and $\hat{\Lambda}_{c2}$, are thus ruled out. Hence, it appears that $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ dictates the inner structure of $\Lambda_c(2880)$.¹⁾ However, there are several problems with this assignment: (i) the quark model indicates a $\Lambda_{c2}(\frac{5}{2}^+)$ state around 2910 MeV, which is close to the mass of $\Lambda_c(2880)$, whereas the mass of $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ is even higher [17, 18], (ii) $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ can decay to an F -wave $\Lambda_c\pi$, and this has not been seen by BaBar and Belle, and (iii) the calculated width, 28.8 MeV, is too large compared to the measured one, 5.8 ± 1.1 MeV. One may argue that the 3P_0 model's prediction can easily differ from the experimental measurement by a factor of 2–3 owing to its inherent uncertainties [22].

Interestingly, the quantum numbers $J^P = \frac{5}{2}^+$ for the $\Lambda_c(2880)$ were correctly predicted on the basis of the diquark concept in Refs. [24, 25] before the Belle experiment.

The highest state $\Lambda_c(2940)^+$ was first discovered by BaBar in the D^0p decay mode [20] and confirmed by Belle in the decays $\Sigma_c^0\pi^+$, $\Sigma_c^{++}\pi^-$, which subsequently decay into $\Lambda_c^+\pi^+\pi^-$ [21]. Its spin-parity assignment is quite diverse. For example, it has been argued that $\Lambda_c(2940)^+$ is the radial excitation of $\Lambda_c(2595)$ with $J^P = \frac{1}{2}^-$, but the predicted mass is too large by the order of 40 MeV. Alternatively, it could be the first radial excitation of Σ_c (not Λ_c !) with $J^P = 3/2^+$ [26]. The latter assignment has the advantage that the predicted mass is in better agreement with experiment. Because the mass of $\Lambda_c(2940)^+$ is barely below the threshold of $D^{*0}p$, this observation motivated the authors of Ref. [27] to suggest

¹⁾ It has been argued in Ref. [23] that in the chiral quark model $\Lambda_c(2880)$ favors to be the state $|\Lambda_c^{2S+1}L_\sigma J^P\rangle = |\Lambda_c^2 D_{\lambda\lambda} \frac{3}{2}^+\rangle$ with $L_\rho = 0$ and $L_\lambda = 2$ rather than $|\Lambda_c^2 D_A \frac{5}{2}^+\rangle$ with $L_\rho = L_\lambda = 1$ as the latter cannot decay into D^0p . However, this is not our case as $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ does decay to D^0p and can reproduce the measured value of R .

an exotic molecular state of D^{*0} and p with a binding energy of the order of 6 MeV and $J^P = \frac{1}{2}^-$ for $\Lambda_c(2940)^+$. 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 [17, 18]. A similar result of 2906 MeV for $\frac{3}{2}^+$ Λ_c was also obtained in the relativistic quark model [28].

2.1.2 Σ_c states

The highest isotriplet charmed baryons, $\Sigma_c(2800)^{++,+,0}$, which decay to $\Lambda_c^+\pi$, were first measured by Belle [29] with widths of the order of 70 MeV. 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 $\Sigma_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$. The states Σ_{c1} and $\tilde{\Sigma}_{c1}$ are ruled out because their decays to $\Lambda_c^+\pi$ are prohibited in the heavy quark limit. Now the $\Sigma_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ baryon decays primarily into the $\Lambda_c\pi$ system in a D -wave, whereas $\Sigma_{c0}(\frac{1}{2}^-)$ decays into $\Lambda_c\pi$ in an S -wave. Because heavy hadron chiral perturbation theory (HHChPT) implies a very broad Σ_{c0} with a width of the order of 885 MeV (see Section 3.2 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 a mixture of the two. In the quark-diquark model [26], both of them have very close masses compatible with experiment. Given that for light strange baryons, the first orbital excitation of the Σ also has the quantum numbers $J^P = 3/2^-$ (see Fig. 2), we will advocate a $\Sigma_{c2}(3/2^-)$ state for $\Sigma_c(2800)$.

2.1.3 Ξ_c states

The states $\Xi_c(2790)$ and $\Xi_c(2815)$ form the doublet $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$. Because the diquark transition $1^- \rightarrow 0^+ + \pi$ is prohibited, $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$ cannot decay to $\Xi_c\pi$. The dominant decay modes are $[\Xi'_c\pi]_S$ for $\Xi_{c1}(\frac{1}{2}^-)$ and $[\Xi_c^*\pi]_S$ for $\Xi_{c1}(\frac{3}{2}^-)$.

Many excited charmed baryon states $\Xi_c(2980)$, $\Xi_c(3055)$, $\Xi_c(3080)$, and $\Xi_c(3123)$ have been seen at B factories [12, 30, 31]. Another state $\Xi_c(2930)^0$, which is omitted from the PDG summary table, has been seen only by BaBar in the $\Lambda_c^+K^-$ mass projection of $B^- \rightarrow \Lambda_c^+\bar{\Lambda}_c^-K^-$ [33]. However, as we shall see below, it may form a sextet with $\Sigma_c(2800)$ and $\Omega_c(3050)$. The states $\Xi_c(2980)$, $\Xi_c(3055)$, $\Xi_c(3080)$, and $\Xi_c(3123)$ could be the first positive-parity excitations of the Ξ_c . The study of Regge phenomenology is very useful for the J^P assignment of charmed baryons [26, 34]. The Regge analysis suggests that $J^P = 3/2^+$ for $\Xi_c(3055)$ and $5/2^+$ for $\Xi_c(3080)$ [26]. From Table 5 below, we shall see that $\Xi_c(3080)$ and $\Lambda_c(2880)$ form a nice $J^P = 5/2^+$

antitriplet.

In the relativistic quark-diquark model [26], $\Xi_c(2980)$ is a sextet $J^P = \frac{1}{2}^+$ state. According to Table 2, possible candidates are $\Xi'_{c1}(\frac{1}{2}^+)$, $\hat{\Xi}'_{c1}(\frac{1}{2}^+)$, $\tilde{\Xi}'_{c0}(\frac{1}{2}^+)$, and $\tilde{\Xi}'_{c1}(\frac{1}{2}^+)$. As pointed out in Ref. [32], strong decays of these four states, which were studied in Ref. [22] using the 3P_0 model show that $\tilde{\Xi}'_{c1}(\frac{1}{2}^+)$ 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 $J_\ell = L_\ell = 1$.

The possible quark states for the $J^P = \frac{5}{2}^+$ $\Xi_c(3080)$ baryon in an antitriplet are $\Xi_{c2}(\frac{5}{2}^+)$, $\hat{\Xi}_{c2}(\frac{5}{2}^+)$, $\tilde{\Xi}_{c2}(\frac{5}{2}^+)$, $\tilde{\Xi}_{c2}^2(\frac{5}{2}^+)$, and $\tilde{\Xi}_{c3}^2(\frac{5}{2}^+)$ (see Table 2). Because $\Xi_c(3080)$ is above the $D\Lambda$ threshold, the two-body mode $D\Lambda$ should exist although 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 [13]. According to the 3P_0 model, the first four states are excluded as they do not decay into $D\Lambda$ [22]. The only remaining possibility is $\tilde{\Xi}_{c3}^2(\frac{5}{2}^+)$. This is the analog of $\tilde{\Lambda}_{c3}^2(\frac{5}{2}^+)$ for $\Lambda_c(2880)$. Nevertheless, the identification of $\tilde{\Xi}_{c3}^2(\frac{5}{2}^+)$ with $\Xi_c(3080)$ encounters two potential problems: (i) its width is dominated by the $\Xi_c\pi$ and $\Lambda_c^+\bar{K}$ modes, which have not been seen experimentally, and (ii) the predicted width of the order of 47 MeV [22] is too large compared to the measured one, which is of the order of 5.7 MeV.

2.1.4 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 $\Omega_c(2770) \rightarrow \Omega_c\gamma$ [39].

Molecular picture

Because $\Lambda_c(2940)^+$ and $\Sigma_c(2800)$ are barely below the $D^{*0}p$ and DN thresholds, respectively, it is tempting to conjecture an exotic molecular structure of D^{*0} and p for the former and a molecular state of DN for the latter [27, 40–44]. Likewise, $\Xi_c(2980)$ could be a molecular state of $D\Lambda$.

A coupled-channel calculation of the baryon-meson ND system has been performed to look for an isospin-spin channel that is attractive enough to form a molecular state [42, 45]. $(I)J^P = (0)\frac{1}{2}^-$ was found to be the most attractive one, followed by $(I)J^P = (1)\frac{3}{2}^-$. This suggests that $\Sigma_c(2800)$ might be an s -wave DN molecular state with $(I)J^P = (0)\frac{1}{2}^-$ and $\Lambda_c(2940)$ an s -wave D^*N molecular state with $(I)J^P = (1)\frac{3}{2}^-$ (see Fig. 3

of [45]). Another possibility is a DN molecular state with $(I)J^P = (1)\frac{3}{2}^-$ for $\Sigma_c(2800)$ and a D^*N state with $(I)J^P = (0)\frac{1}{2}^-$ for $\Lambda_c(2940)$. Because $\Sigma_c(2800)$ has isospin 1, and moreover, as we have noted in passing, $\Sigma_c(2800)$ will be too broad if it is assigned to $J^P = 1/2^-$, we conclude that the second possibility is preferable (see also [44]).

The possible spin-parity quantum numbers of the higher excited charmed baryon resonances that have been suggested in the literature are partially summarized in Table 4. Some of the predictions are already ruled out by experiment. For example, $\Lambda_c(2880)$ has $J^P = \frac{5}{2}^+$, as seen by Belle. More experimental studies are certainly to pin down the quantum numbers.

Charmed baryon spectroscopy has been studied extensively in various models. It appears that the spectroscopy is well described by the heavy quark-light diquark picture elaborated by Ebert, Faustov, and Galkin (EFG) [26] (see also Ref. [35]). As noted in passing, the quantum numbers $J^P = \frac{5}{2}^+$ of $\Lambda_c(2880)$ were correctly predicted in a model based on the diquark idea before the Belle experiment [24, 25]. Moreover, EFG have shown that all the available experimental data on heavy baryons fit the linear Regge trajectories nicely, namely, the trajectories in the (J, M^2) and (n_r, M^2) planes for orbitally and radially excited heavy baryons, respectively,

$$J = \alpha M^2 + \alpha_0, \quad n_r = \beta M^2 + \beta_0, \quad (2.5)$$

where n_r is the radial excitation quantum number, α and β are the slopes, and α_0 and β_0 are intercepts. 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. [26] can be regarded as a theoretical benchmark (see Fig. 3).

Antitriplet and sextet states

The antitriplet and sextet states of charmed baryons are listed in Table 5. To date, the following states

are established: the $J^P = \frac{1}{2}^+, \frac{1}{2}^-$, and $\frac{3}{2}^- \bar{\mathbf{3}}$ states, $(\Lambda_c^+, \Xi_c^+, \Xi_c^0)$, $(\Lambda_c(2595)^+, \Xi_c(2790)^+, \Xi_c(2790)^0)$, and $(\Lambda_c(2625)^+, \Xi_c(2815)^+, \Xi_c(2815)^0)$, respectively, and the $J^P = \frac{1}{2}^+$ and $\frac{3}{2}^+ \mathbf{6}$ states, $(\Omega_c, \Sigma_c, \Xi_c')$ and $(\Omega_c^*, \Sigma_c^*, \Xi_c'^*)$, respectively. The mass difference $m_{\Xi_c} - m_{\Lambda_c}$ in the antitriplet states clearly lies between 180 and 200 MeV. We note in passing that $\Xi_c(3080)$ should carry the quantum numbers $J^P = 5/2^+$. From Table 5, we see that $\Xi_c(3080)$ and $\Lambda_c(2880)$ form a nice $J^P = 5/2^+$ antitriplet as the mass difference between $\Xi_c(3080)$ and $\Lambda_c(2880)$ is consistent with that observed in other antitriplets. Likewise, the mass differences in the $J^P = 3/2^-$ sextet $(\Omega_c(3050), \Xi_c'(2930), \Sigma_c(2800))$ predicted by the quark-diquark model are consistent with that measured in $J^P = 1/2^+$ and $3/2^+$ sextets. Note that there is no $J^P = \frac{1}{2}^-$ sextet as the $\Sigma_c(2800)$ with these spin-parity quantum numbers will be too broad to be observed.

On the basis of the QCD sum rules, many charmed baryon multiplets classified according to $[\mathbf{6}_F(\text{or } \bar{\mathbf{3}}_F), J_\ell, S_\ell, \rho/\lambda]$ were recently studied in Ref. [36]. Three sex-

Charmed baryon spectra

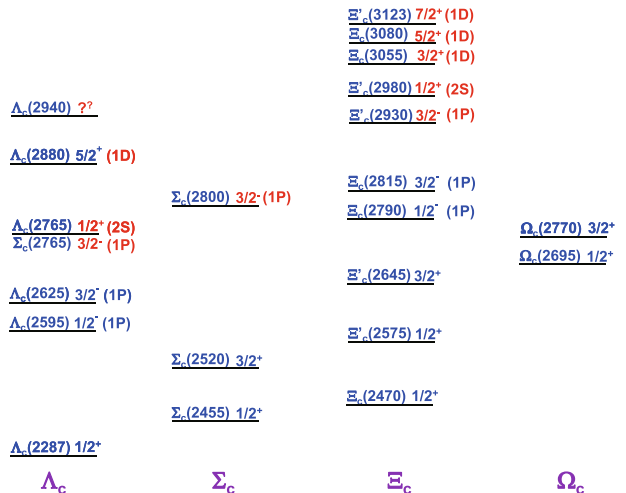


Fig. 3 Singly charmed baryon states where the spin-parity quantum numbers in red are taken from Ref. [26].

Table 4 Possible spin-parity quantum numbers for excited charmed baryon resonances.

| | $\Lambda_c(2765)$ | $\Lambda_c(2880)$ | $\Lambda_c(2940)$ | $\Sigma_c(2800)$ | $\Xi_c(2930)$ | $\Xi_c(2980)$ | $\Xi_c(3055)$ | $\Xi_c(3080)$ | $\Xi_c(3123)$ |
|---------------------------------|---------------------|--------------------------------|--|------------------------------------|---|------------------------------------|---------------------|---------------------|------------------------------------|
| Capstick <i>et al.</i> [17, 18] | $\frac{1}{2}^+$ | | $\frac{3}{2}^+, \frac{5}{2}^-$ | $\frac{3}{2}^-, \frac{5}{2}^-$ | | | | | |
| B. Chen <i>et al.</i> [35] | $\frac{1}{2}^+(2S)$ | $\frac{5}{2}^+(1D)$ | $\frac{1}{2}^-(2P)$ | | | $\frac{1}{2}^+(2S)$ | $\frac{3}{2}^+(1D)$ | $\frac{5}{2}^+(1D)$ | $\frac{1}{2}^-(2P)$ |
| H. Chen <i>et al.</i> [36] | | | $\frac{1}{2}^+, \frac{1}{2}^-$ | $\frac{1}{2}^-, \frac{3}{2}^-$ | | $\frac{1}{2}^-, \frac{3}{2}^-$ | | $\frac{5}{2}^-$ | |
| Cheng <i>et al.</i> [10] | | | | $\frac{3}{2}^-$ | | $\frac{1}{2}^+$ | | $\frac{5}{2}^+$ | |
| Ebert <i>et al.</i> [26] | $\frac{1}{2}^+(2S)$ | $\frac{5}{2}^+(1D)$ | $\frac{1}{2}^-(2P), \frac{3}{2}^+(2S)$ | $\frac{1}{2}^-, \frac{3}{2}^-(1P)$ | $\frac{1}{2}^-, \frac{3}{2}^-, \frac{5}{2}^-$ | $\frac{1}{2}^+(2S)$ | $\frac{3}{2}^+(1D)$ | $\frac{5}{2}^+(1D)$ | $\frac{7}{2}^+(1D)$ |
| Garcilazo <i>et al.</i> [28] | $\frac{1}{2}^+$ | $\frac{1}{2}^-, \frac{3}{2}^-$ | $\frac{3}{2}^+$ | $\frac{1}{2}^-, \frac{3}{2}^-$ | | | | | |
| Gerasyuata <i>et al.</i> [37] | $\frac{5}{2}^-$ | $\frac{1}{2}^-$ | | $\frac{5}{2}^-$ | | | | | |
| Liu <i>et al.</i> [38] | | | | | $\frac{1}{2}^-(1P)$ | $\frac{1}{2}^-, \frac{3}{2}^-(1P)$ | $\frac{3}{2}^+(1D)$ | $\frac{1}{2}^+(2S)$ | $\frac{3}{2}^+, \frac{5}{2}^+(1D)$ |
| Wilczek <i>et al.</i> [24, 25] | | $\frac{5}{2}^+$ | | | | | | | |
| Zhong <i>et al.</i> [23] | $\frac{1}{2}^-(1P)$ | $\frac{3}{2}^+(1D)$ | $\frac{5}{2}^+(1D)$ | $\frac{1}{2}^-, \frac{5}{2}^-(1P)$ | | | | | |

Table 5 Antitriplet and sextet states of charmed baryons. The J^P quantum numbers of $\Xi_c(3080)$, $\Xi'_c(2930)$, $\Sigma_c(2800)$ are not yet established and the $\Omega_c(3/2^-)$ state has not been observed. Mass differences $\Delta m_{\Xi_c \Lambda_c} \equiv m_{\Xi_c} - m_{\Lambda_c}$, $\Delta m_{\Xi'_c \Sigma_c} \equiv m_{\Xi'_c} - m_{\Sigma_c}$, $\Delta m_{\Omega_c \Xi'_c} \equiv m_{\Omega_c} - m_{\Xi'_c}$ are in units of MeV.

| | J^P | States | Mass difference | Status |
|--------------------|---------|--|---|--------|
| $\bar{\mathbf{3}}$ | $1/2^+$ | $\Lambda_c(2287)^+$, $\Xi_c(2470)^+$, $\Xi_c(2470)^0$ | $\Delta m_{\Xi_c \Lambda_c} = 183$ | estab |
| | $1/2^-$ | $\Lambda_c(2595)^+$, $\Xi_c(2790)^+$, $\Xi_c(2790)^0$ | $\Delta m_{\Xi_c \Lambda_c} = 198$ | estab |
| | $3/2^-$ | $\Lambda_c(2625)^+$, $\Xi_c(2815)^+$, $\Xi_c(2815)^0$ | $\Delta m_{\Xi_c \Lambda_c} = 190$ | estab |
| | $5/2^+$ | $\Lambda_c(2880)^+$, $\Xi_c(3080)^+$, $\Xi_c(3080)^0$ | $\Delta m_{\Xi_c \Lambda_c} = 196$ | [26] |
| $\mathbf{6}$ | $1/2^+$ | $\Omega_c(2695)^0$, $\Xi'_c(2575)^{+,0}$, $\Sigma_c(2455)^{+,+,0}$ | $\Delta m_{\Xi'_c \Sigma_c} = 124$, $\Delta m_{\Omega_c \Xi'_c} = 119$ | estab |
| | $3/2^+$ | $\Omega_c(2770)^0$, $\Xi'_c(2645)^{+,0}$, $\Sigma_c(2520)^{+,+,0}$ | $\Delta m_{\Xi'_c \Sigma_c} = 128$, $\Delta m_{\Omega_c \Xi'_c} = 120$ | estab |
| | $3/2^-$ | $\Omega_c(3050)^0$, $\Xi'_c(2930)^{+,0}$, $\Sigma_c(2800)^{+,+,0}$ | $\Delta m_{\Xi'_c \Sigma_c} = 131$, $\Delta m_{\Omega_c \Xi'_c} = 119$ | [26] |

tets 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 $J^P = 5/2^-$. Note that $\Xi'_c(2980)$ and $\Xi'_c(3080)$ were treated as p -wave baryons rather than the first positive-parity excitations, as we have discussed before. The results for the multiplet $[6_F, 1, 0, \rho]$ led the authors of Ref. [36] 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 $\Omega_c(1/2^-, 3/2^-)$ is around 3250 ± 200 MeV. Using the central value of the predicted masses to label the states in the multiplet $[6_F, 1, 0, \rho]$ (see Table I of Ref. [36]), one obtains ($\Omega_c(3250)$, $\Xi'_c(2960)$, $\Sigma_c(2730)$) for $J^P = 1/2^-$ and ($\Omega_c(3260)$, $\Xi'_c(2980)$, $\Sigma_c(2750)$) for $J^P = 3/2^-$. One can check that $\Delta m_{\Xi'_c \Sigma_c} = 230 \pm 234$ MeV, and $\Delta m_{\Omega_c \Xi'_c}$ is of the order of 285 ± 250 MeV. Owing to the large theoretical uncertainties in the masses, it is not clear whether the QCD sum rule calculations are compatible with the mass differences measured in the $J^P = 1/2^+$ and $3/2^+$ sextets. In any event, it will be interesting to test these two different model predictions for $J^P = 3/2^-$ and $1/2^-$ sextets in the future.

2.2 Doubly charmed baryons

Evidence of doubly charmed baryon states has been reported by SELEX in $\Xi_{cc}(3520)^+ \rightarrow \Lambda_c^+ K^- \pi^+$ [46]. Further observation of $\Xi_{cc}^+ \rightarrow p D^+ K^-$ was also announced by SELEX [47]. However, none of the doubly charmed states discovered by SELEX has been confirmed by FOCUS [48], BaBar [49], Belle [12] and LHCb [50], although $10^6 \Lambda_c$ events are produced in B factories, for example, versus 1630 Λ_c events observed at SELEX.

The doubly charmed baryons $\Xi_{cc}^{(*)++}$, $\Xi_{cc}^{(*)+}$, $\Omega_{cc}^{(*)+}$ with the quark contents ccu, ccd, ccs form an SU(3) triplet. They have been studied extensively using many different approaches: the quark model, light quark-heavy diquark model, QCD sum rules, and lattice simu-

lation. The predicted doubly charmed baryon masses calculated in various models are tabulated in Refs. [51, 52]. For recent QCD sum rule calculations, see e.g. [53–55]. Chiral corrections to the masses of doubly heavy baryons up to N³LO were presented in Ref. [56].

Figure 4 shows the results of recent lattice studies of doubly and triply charmed baryon spectra by different groups: RQCD [57], HSC [58], Brown *et al.* [59], ETMC [60], ILGTI [61], PACS-CS [62], Durr *et al.* [63], Briceno *et al.* [64], Liu *et al.* [65], and Na *et al.* [66]. A new lattice calculation of $\Omega_{cc}^{(*)}$ and Ω_{ccc} was available in Ref. [67]. The various lattice results are consistent with each other and they fall into the ranges

$$\begin{aligned}
 M(\Xi_{cc}) &= 3.54\text{--}3.68 \text{ GeV}, \\
 M(\Xi_{cc}^*) &= 3.61\text{--}3.72 \text{ GeV}, \\
 M(\Omega_{cc}) &= 3.57\text{--}3.76 \text{ GeV}, \\
 M(\Omega_{cc}^*) &= 3.68\text{--}3.85 \text{ GeV},
 \end{aligned}
 \tag{2.6}$$

and

$$M(\Omega_{ccc}) = 4.70\text{--}4.84 \text{ MeV}.
 \tag{2.7}$$

Although lattice studies suggest that the mass of the low-lying Ξ_{cc} exceeds 3519 MeV, it is interesting to note that the authors of [51] calculated the masses of doubly and triply charmed baryons on the basis of the Regge phenomenology and found $M(\Xi_{cc}^+) = 3520.2^{+40.6}_{-39.8}$ MeV, in good agreement with SELEX.

3 Strong decays

Owing to the rich mass spectrum and 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 $\Sigma_c \rightarrow \Lambda_c \pi$ are soft. Therefore, heavy quark symmetry of the heavy

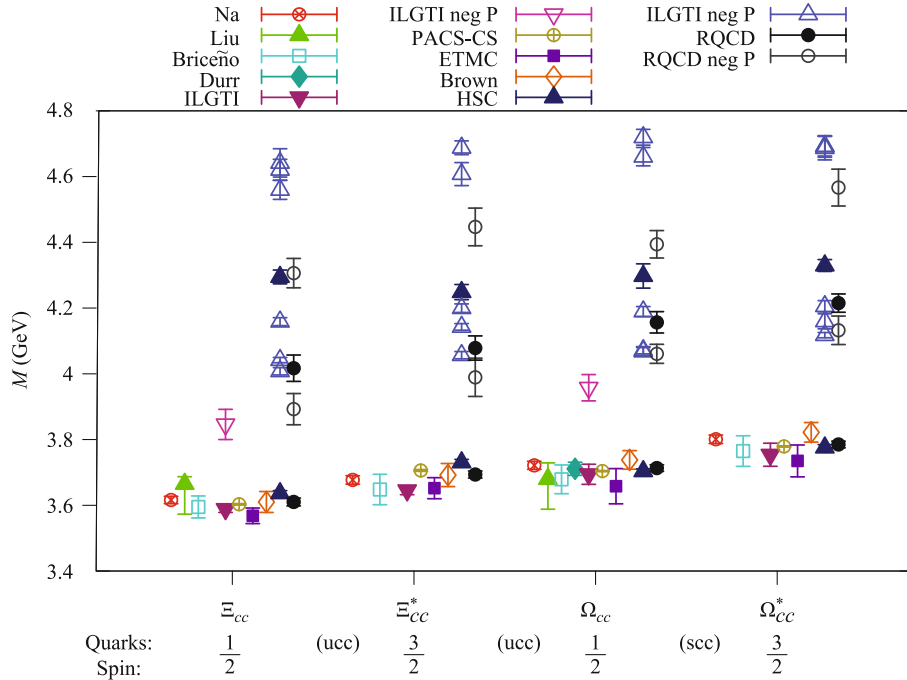


Fig. 4 Doubly charmed low-lying baryon spectra taken from Ref. [57].

quark and chiral symmetry of the light quarks will have interesting implications for the low-energy dynamics of heavy baryons interacting with Goldstone bosons.

The strong decays of charmed baryons are most conveniently described by the HHChPT, into which heavy quark symmetry and chiral symmetry are incorporated [68–70]. Heavy baryon chiral Lagrangians were first constructed in Ref. [68] for strong decays of s -wave charmed baryons and in Refs. [9, 14] for p -wave ones. Previous phenomenological studies of the strong decays of p -wave charmed baryons based on HHChPT can be found in Refs. [9, 10, 14, 71, 72]. The chiral Lagrangian involves two coupling constants, g_1 and g_2 , for P -wave transitions between s -wave and s -wave baryons [68]; 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 [9]. The general chiral Lagrangian for heavy baryons coupled to pseudoscalar mesons can be expressed compactly in terms of superfields. We will not write the relevant Lagrangians here; instead the reader is referred to Eqs. (3.1) and (3.3) of Ref. [9]. The partial widths relevant for our purposes are [9]

$$\Gamma(\Sigma_c^* \rightarrow \Sigma_c \pi) = \frac{g_1^2}{2\pi f_\pi^2} \frac{m_{\Sigma_c}}{m_{\Sigma_c^*}} p_\pi^3,$$

$$\Gamma(\Sigma_c \rightarrow \Lambda_c \pi) = \frac{g_2^2}{2\pi f_\pi^2} \frac{m_{\Lambda_c}}{m_{\Sigma_c}} p_\pi^3,$$

$$\Gamma(\Lambda_{c1}(1/2^-) \rightarrow \Sigma_c \pi) = \frac{h_2^2}{2\pi f_\pi^2} \frac{m_{\Sigma_c}}{m_{\Lambda_{c1}}} E_\pi^2 p_\pi,$$

$$\Gamma(\Sigma_{c0}(1/2^-) \rightarrow \Lambda_c \pi) = \frac{h_3^2}{2\pi f_\pi^2} \frac{m_{\Lambda_c}}{m_{\Sigma_{c0}}} E_\pi^2 p_\pi,$$

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

$$\Gamma(\Sigma_{c1}(3/2^-) \rightarrow \Sigma_c^* \pi) = \frac{h_9^2}{9\pi f_\pi^2} \frac{m_{\Sigma_c^*}}{m_{\Sigma_{c1}(3/2^-)}} p_\pi^5,$$

$$\Gamma(\Sigma_{c2}(3/2^-) \rightarrow \Lambda_c \pi) = \frac{4h_{10}^2}{15\pi f_\pi^2} \frac{m_{\Lambda_c}}{m_{\Sigma_{c2}}} p_\pi^5,$$

$$\Gamma(\Sigma_{c2}(3/2^-) \rightarrow \Sigma_c^* \pi) = \frac{h_{11}^2}{10\pi f_\pi^2} \frac{m_{\Sigma_c^*}}{m_{\Sigma_{c2}}} p_\pi^5,$$

$$\Gamma(\Sigma_{c2}(5/2^-) \rightarrow \Sigma_c \pi) = \frac{2h_{11}^2}{45\pi f_\pi^2} \frac{m_{\Sigma_c}}{m_{\Sigma_{c2}}} p_\pi^5,$$

$$\Gamma(\Sigma_{c2}(5/2^-) \rightarrow \Sigma_c^* \pi) = \frac{7h_{11}^2}{45\pi f_\pi^2} \frac{m_{\Sigma_c^*}}{m_{\Sigma_{c2}}} p_\pi^5, \quad (3.1)$$

where $f_\pi = 132$ 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 $g_1, g_2, h_2, \dots, h_7$ are dimensionless, whereas h_8, \dots, h_{15} have canonical dimension E^{-1} .

3.1 Strong decays of s -wave charmed baryons

Because the strong decay $\Sigma_c^* \rightarrow \Sigma_c \pi$ is kinematically prohibited, the coupling g_1 cannot be extracted directly

Table 6 Decay widths (in units of MeV) of *s*-wave charmed baryons where the measured rates are taken from 2006 PDG [73].

| Decay | Expt. | HHChPT |
|---|-----------------|----------------|
| $\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+$ | 2.23 ± 0.30 | input |
| $\Sigma_c^+ \rightarrow \Lambda_c^+ \pi^0$ | < 4.6 | 2.6 ± 0.4 |
| $\Sigma_c^0 \rightarrow \Lambda_c^+ \pi^-$ | 2.2 ± 0.4 | 2.2 ± 0.3 |
| $\Sigma_c(2520)^{++} \rightarrow \Lambda_c^+ \pi^+$ | 14.9 ± 1.9 | 16.7 ± 2.3 |
| $\Sigma_c(2520)^+ \rightarrow \Lambda_c^+ \pi^0$ | < 17 | 17.4 ± 2.3 |
| $\Sigma_c(2520)^0 \rightarrow \Lambda_c^+ \pi^-$ | 16.1 ± 2.1 | 16.6 ± 2.2 |
| $\Xi_c(2645)^+ \rightarrow \Xi_c^{0,+} \pi^{+,0}$ | < 3.1 | 2.8 ± 0.4 |
| $\Xi_c(2645)^0 \rightarrow \Xi_c^{+,0} \pi^{-,0}$ | < 5.5 | 2.9 ± 0.4 |

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 $\Sigma_c^{(*)} \rightarrow \Lambda_c \pi$, $\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+$ is the most well-measured. Hence, we shall use this mode to extract the coupling g_2 . Using the 2006 data [73] $\Gamma(\Sigma_c^{++}) = \Gamma(\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+) = 2.23 \pm 0.30$ MeV, the coupling g_2 is extracted as

$$|g_2|_{2006} = 0.605_{-0.043}^{+0.039} \quad (3.2)$$

The predicted rates of other modes are shown in Table 6. The agreement between theory and experiment is clearly excellent, except that the predicted width for $\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+$ is slightly too large.

Using the new data from the 2014 Particle Data Group [8] in conjunction with the new measurements of the Σ_c and Σ_c^* widths by Belle [11], we obtain the new average $\Gamma(\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+) = 1.94_{-0.16}^{+0.08}$ MeV (see Table 3). Therefore, the coupling g_2 is reduced to

$$|g_2|_{2015} = 0.565_{-0.024}^{+0.011} \quad (3.3)$$

From Table 7 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 [12], and the new calculated width for $\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+$ is now in agreement with 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.

3.2 Strong decays of *p*-wave charmed baryons

Because $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ form the doublet $\Lambda_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, it appears from Eq. (3.1) that the couplings h_2 and h_8 can in principle be extracted from $\Lambda_c(2595) \rightarrow \Sigma_c \pi$ and from $\Lambda_c(2625) \rightarrow \Sigma_c \pi$, respectively. 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 [9].

Although the coupling h_2 can be inferred from the two-body decay $\Lambda_c(2595) \rightarrow \Sigma_c \pi$, this method is less accurate because this decay is kinematically barely allowed or even prohibited, depending on the mass of $\Lambda_c(2595)^+$. For the old mass measurement $m(\Lambda_c(2595)) = 2595.4 \pm 0.6$ MeV

Table 8 Decay widths (in units of MeV) of *p*-wave charmed baryons where the measured rates are taken from 2006 PDG [73].

| Decay | Expt. | HHChPT |
|---|------------------------|------------------------|
| | [73] | [10] |
| $\Lambda_c(2595)^+ \rightarrow (\Lambda_c^+ \pi \pi)_R$ | $2.63_{-1.09}^{+1.56}$ | input |
| $\Lambda_c(2595)^+ \rightarrow \Sigma_c^{++} \pi^-$ | $0.65_{-0.31}^{+0.41}$ | $0.72_{-0.30}^{+0.43}$ |
| $\Lambda_c(2593)^+ \rightarrow \Sigma_c^0 \pi^+$ | $0.67_{-0.31}^{+0.41}$ | $0.77_{-0.32}^{+0.46}$ |
| $\Lambda_c(2593)^+ \rightarrow \Sigma_c^+ \pi^0$ | | $1.57_{-0.65}^{+0.93}$ |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^{++} \pi^-$ | < 0.10 | 0.029 |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^0 \pi^+$ | < 0.09 | 0.029 |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^+ \pi^0$ | | 0.041 |
| $\Lambda_c(2625)^+ \rightarrow \Lambda_c^+ \pi \pi$ | < 1.9 | 0.21 |
| $\Sigma_c(2800)^{++} \rightarrow \Lambda_c \pi, \Sigma_c^{(*)} \pi$ | 75_{-17}^{+22} | input |
| $\Sigma_c(2800)^+ \rightarrow \Lambda_c \pi, \Sigma_c^{(*)} \pi$ | 62_{-40}^{+60} | input |
| $\Sigma_c(2800)^0 \rightarrow \Lambda_c \pi, \Sigma_c^{(*)} \pi$ | 61_{-18}^{+28} | input |
| $\Xi_c(2790)^+ \rightarrow \Xi_c^{0,+} \pi^{+,0}$ | < 15 | $8.0_{-3.3}^{+4.7}$ |
| $\Xi_c(2790)^0 \rightarrow \Xi_c^{'+,0} \pi^{-,0}$ | < 12 | $8.5_{-3.5}^{+5.0}$ |
| $\Xi_c(2815)^+ \rightarrow \Xi_c^{*+,0} \pi^{0,+}$ | < 3.5 | $3.4_{-1.4}^{+2.0}$ |
| $\Xi_c(2815)^0 \rightarrow \Xi_c^{*+,0} \pi^{-,0}$ | < 6.5 | $3.6_{-1.5}^{+2.1}$ |

Table 7 Decay widths (in units of MeV) of *s*-wave charmed baryons. Data are taken from 2014 PDG [8] together with the new measurements of Σ_c , Σ_c^* [11] and $\Xi_c(2645)^+$ widths [12]. Theoretical predictions of [74] are taken from Table IV of [75].

| Decay | Expt. | HHChPT | Tawfiq | Ivanov | Huang | Albertus |
|---|------------------------|----------------------|--------------------|--------------------|--------------------|--------------------|
| | [8] | | <i>et al.</i> [74] | <i>et al.</i> [75] | <i>et al.</i> [76] | <i>et al.</i> [78] |
| $\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+$ | $1.94_{-0.16}^{+0.08}$ | input | 1.51 ± 0.17 | 2.85 ± 0.19 | 2.5 | 2.41 ± 0.07 |
| $\Sigma_c^+ \rightarrow \Lambda_c^+ \pi^0$ | < 4.6 | $2.3_{-0.2}^{+0.1}$ | 1.56 ± 0.17 | 3.63 ± 0.27 | 3.2 | 2.79 ± 0.08 |
| $\Sigma_c^0 \rightarrow \Lambda_c^+ \pi^-$ | $1.9_{-0.2}^{+0.1}$ | $1.9_{-0.2}^{+0.1}$ | 1.44 ± 0.16 | 2.65 ± 0.19 | 2.4 | 2.37 ± 0.07 |
| $\Sigma_c(2520)^{++} \rightarrow \Lambda_c^+ \pi^+$ | $14.8_{-0.4}^{+0.3}$ | $14.5_{-0.8}^{+0.5}$ | 11.77 ± 1.27 | 21.99 ± 0.87 | 8.2 | 17.52 ± 0.75 |
| $\Sigma_c(2520)^+ \rightarrow \Lambda_c^+ \pi^0$ | < 17 | $15.2_{-1.3}^{+0.6}$ | | | 8.6 | 17.31 ± 0.74 |
| $\Sigma_c(2520)^0 \rightarrow \Lambda_c^+ \pi^-$ | $15.3_{-0.5}^{+0.4}$ | $14.7_{-1.2}^{+0.6}$ | 11.37 ± 1.22 | 21.21 ± 0.81 | 8.2 | 16.90 ± 0.72 |
| $\Xi_c(2645)^+ \rightarrow \Xi_c^{0,+} \pi^{+,0}$ | 2.6 ± 0.5 | $2.4_{-0.2}^{+0.1}$ | 1.76 ± 0.14 | 3.04 ± 0.37 | | 3.18 ± 0.10 |
| $\Xi_c(2645)^0 \rightarrow \Xi_c^{+,0} \pi^{-,0}$ | < 5.5 | $2.5_{-0.2}^{+0.1}$ | 1.83 ± 0.06 | 3.12 ± 0.33 | | 3.03 ± 0.10 |

Table 9 Decay widths (in units of MeV) of p -wave charmed baryons. Data are taken from 2014 PDG [8] together with the new measurements of Σ_c, Σ_c^* [11] and $\Xi_c(2645)^+$ widths [12]. Theoretical predictions of [74] are taken from Table IV of [75].

| Decay | Expt. [8] | HHChPT | Tawfiq <i>et al.</i> [74] | Ivanov <i>et al.</i> [75] | Huang <i>et al.</i> [76] | Zhu <i>et al.</i> [22] |
|---|------------------|------------------------|------------------------------|------------------------------|-----------------------------|---------------------------|
| $\Lambda_c(2595)^+ \rightarrow (\Lambda_c^+ \pi\pi)_R$ | 2.59 ± 0.56 | input | | | 2.5 | |
| $\Lambda_c(2595)^+ \rightarrow \Sigma_c^{++} \pi^-$ | | | 1.47 ± 0.57 | 0.79 ± 0.09 | $0.55^{+1.3}_{-0.55}$ | 0.64 |
| $\Lambda_c(2595)^+ \rightarrow \Sigma_c^0 \pi^+$ | | | 1.78 ± 0.70 | 0.83 ± 0.09 | 0.89 ± 0.86 | 0.86 |
| $\Lambda_c(2595)^+ \rightarrow \Sigma_c^+ \pi^0$ | | $2.74^{+0.57}_{-0.60}$ | 1.18 ± 0.46 | 0.98 ± 0.12 | 1.7 ± 0.49 | 1.2 |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^{++} \pi^-$ | < 0.10 | $\lesssim 0.028$ | 0.44 ± 0.23 | 0.076 ± 0.009 | 0.013 | 0.011 |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^0 \pi^+$ | < 0.09 | $\lesssim 0.040$ | 0.47 ± 0.25 | 0.080 ± 0.009 | 0.013 | 0.011 |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^+ \pi^0$ | | $\lesssim 0.029$ | 0.42 ± 0.22 | 0.095 ± 0.012 | 0.013 | 0.011 |
| $\Lambda_c(2625)^+ \rightarrow \Lambda_c^+ \pi\pi$ | < 0.97 | $\lesssim 0.35$ | | | 0.11 | |
| $\Sigma_c(2800)^{++} \rightarrow \Lambda_c \pi, \Sigma_c^{(*)} \pi$ | 75^{+22}_{-17} | input | | | | |
| $\Sigma_c(2800)^+ \rightarrow \Lambda_c \pi, \Sigma_c^{(*)} \pi$ | 62^{+60}_{-40} | input | | | | |
| $\Sigma_c(2800)^0 \rightarrow \Lambda_c \pi, \Sigma_c^{(*)} \pi$ | 72^{+22}_{-15} | input | | | | |
| $\Xi_c(2790)^+ \rightarrow \Xi_c'^{0,+} \pi^{+,0}$ | < 15 | $16.7^{+3.6}_{-3.6}$ | | | | |
| $\Xi_c(2790)^0 \rightarrow \Xi_c'^{+,0} \pi^{-,0}$ | < 12 | $17.7^{+2.9}_{-3.8}$ | | | | |
| $\Xi_c(2815)^+ \rightarrow \Xi_c^{*+,0} \pi^{0,+}$ | < 3.5 | $7.1^{+1.5}_{-1.5}$ | 2.35 ± 0.93 | 0.70 ± 0.04 | | |
| $\Xi_c(2815)^0 \rightarrow \Xi_c^{*+,0} \pi^{-,0}$ | < 6.5 | $7.7^{+1.7}_{-1.7}$ | | | | |

[73], $\Lambda_c(2595)^+ \rightarrow \Sigma_c^{++} \pi^-, \Sigma_c^0 \pi^+$ and $\Lambda_c(2595)^+ \rightarrow \Sigma^+ \pi^0$ are kinematically barely allowed. However, for the new measurement by the CDF, $m(\Lambda_c(2595)) = 2592.25 \pm 0.28$ MeV [79], only the last mode is allowed. Moreover, the finite width effect of the intermediate resonant states could become important [71].

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 . As shown in Ref. [10], the 2006 data $\Gamma(\Lambda_c(2595)) = 3.6^{+2.0}_{-1.3}$ MeV [73] and for the $\Lambda_c(2595)$ mass lead to the resonant rate [10]

$$\Gamma(\Lambda_c(2593)^+ \rightarrow \Lambda_c^+ \pi\pi)_R = (2.63^{+1.56}_{-1.09}) \text{ MeV}, \quad (3.4)$$

as shown in Table 9. Assuming the pole contributions to $\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \pi\pi$ due to the intermediate states Σ_c and Σ_c^* , the resonant rate for the process $\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \pi^+ \pi^-$ can be calculated in the framework of HHChPT [9]. Numerically, we found

$$\begin{aligned} \Gamma(\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \pi\pi)_R &= 13.82h_2^2 + 26.28h_8^2 - 2.97h_2h_8, \\ \Gamma(\Lambda_c(2625)^+ \rightarrow \Lambda_c^+ \pi\pi)_R &= 0.617h_2^2 + 0.136 \times 10^6 h_8^2 - 27h_2h_8, \end{aligned} \quad (3.5)$$

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 the order of 10^{-3} (in units of MeV^{-1}), whereas the decay width of $\Lambda_c(2595)$ is entirely governed by the coupling h_2 . Specifically, we have [10]

$$\begin{aligned} |h_2|_{2006} &= 0.437^{+0.114}_{-0.102}, \\ |h_8|_{2006} &< 3.65 \times 10^{-3} \text{ MeV}^{-1}. \end{aligned} \quad (3.6)$$

It has been noted [72] 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 that will lower the $\Lambda_c(2595)^+$ mass by 2 – 3 MeV compared to the observed mass. A more sophisticated treatment of the mass lineshape of $\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \pi^+ \pi^-$ by the CDF yields $m(\Lambda_c(2595)) = 2592.25 \pm 0.28$ MeV [79], which is 3.1 MeV smaller than the 2006 world average. Therefore, the strong decay $\Lambda_c(2595) \rightarrow \Lambda_c \pi\pi$ is very close to the threshold. With the new measurement of $m(\Lambda_c(2595))$, we have (in units of MeV) [32]

$$\begin{aligned} \Gamma(\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \pi\pi)_R &= g_2^2(20.45h_2^2 + 43.92h_8^2 - 8.95h_2h_8), \\ \Gamma(\Lambda_c(2625)^+ \rightarrow \Lambda_c^+ \pi\pi)_R &= g_2^2(1.78h_2^2 + 4.557 \times 10^6 h_8^2 - 79.75h_2h_8). \end{aligned} \quad (3.7)$$

By fitting to the measured $M(pK^- \pi^+ \pi^+) - M(pK^- \pi^+)$ mass difference distributions and using $g_2^2 = 0.365$, CDF obtained $h_2^2 = 0.36 \pm 0.08$ or $|h_2| = 0.60 \pm 0.07$ [79]. This corresponds to a decay width $\Gamma(\Lambda_c(2595)^+) = 2.59 \pm 0.30 \pm 0.47$ MeV [79]. For the width of $\Lambda_c(2625)^+$, CDF observed a value consistent with zero and therefore calculated an upper limit of 0.97 MeV using a Bayesian approach. From the CDF measurements $\Gamma(\Lambda_c(2595)^+) = 2.59 \pm 0.56$ MeV and $\Gamma(\Lambda_c(2625)^+) < 0.97$ MeV, we obtain

$$\begin{aligned} |h_2|_{2015} &= 0.63 \pm 0.07, \\ |h_8|_{2015} &< 2.32 \times 10^{-3} \text{ MeV}^{-1}. \end{aligned} \quad (3.8)$$

Hence, the magnitude of the coupling h_2 is greatly en-

hanced from 0.437 to 0.63. Our h_2 differs slightly from the value of 0.60 obtained by CDF. This is because the CDF used $|g_2| = 0.604$ to calculate the mass dependence of $\Gamma(\Lambda_c^+ \pi \pi)$, whereas we used $|g_2| = 0.565$.

The large difference between the values of the coupling h_2 obtained in 2006 and 2015 is ascribed to the fact that the mass of $\Lambda_c(2595)^+$ is 3.1 MeV lower than the previous world average because of the threshold effect. To illustrate this, 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. 5. $\Gamma(\Lambda_c^+ \pi \pi)/h_2^2$ at $\Delta M(\Lambda_c(2595)) = 305.79$ MeV is clearly 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 the doublet $\Xi_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$. The $\Xi_c(2790)$ and $\Xi_c(2815)$ widths predicted using the coupling h_2 obtained from Eq. (3.8) and assuming SU(3) flavor symmetry are shown in Table 9. The predicted two-body decay rates of $\Xi_c(2790)^0$ and $\Xi_c(2815)^+$ clearly exceed the current experimental limits because of the enhancement of h_2 (see Table 9). Hence, there is a tension for the coupling h_2 , as its value extracted from $\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \pi \pi$ will imply $\Xi_c(2790)^0 \rightarrow \Xi_c' \pi$ and $\Xi_c(2815)^+ \rightarrow \Xi_c^* \pi$ rates slightly above current limits. It is conceivable that SU(3) flavor symmetry breaking can help account for the discrepancy.

Some information on the coupling h_{10} can be inferred from the strong decays of $\Sigma_c(2800)$. From Eq. (3.1) and the quark model relation $|h_3| = \sqrt{3}|h_2|$ from Ref. [9], we obtain, for example, $\Gamma(\Sigma_c^{++} \rightarrow \Lambda_c^+ \pi^+) \approx 885$ MeV. Hence, $\Sigma_c(2800)$ cannot be identified with $\Sigma_{c0}(1/2^-)$. Using the quark model relation $h_{11}^2 = 2h_{10}^2$ and the measured widths of $\Sigma_c(2800)^{+,+,0}$ (Table 3), we obtain

$$|h_{10}| = (0.85_{-0.08}^{+0.11}) \times 10^{-3} \text{ MeV}^{-1}. \quad (3.9)$$

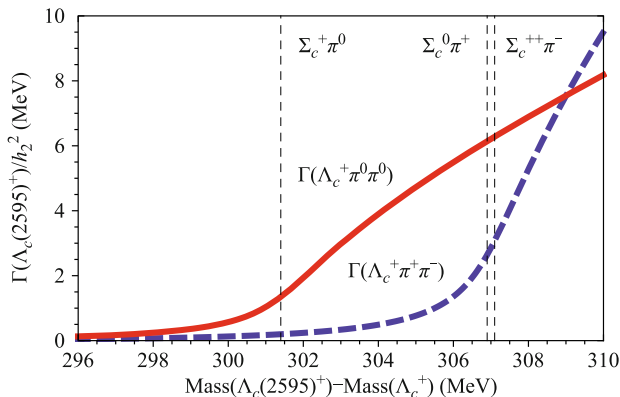


Fig. 5 Calculated dependence of $\Gamma(\Lambda_c^+ \pi^0 \pi^0)/h_2^2$ (full curve) and $\Gamma(\Lambda_c^+ \pi^+ \pi^-)/h_2^2$ (dashed curve) on $m(\Lambda_c(2595)^+) - m(\Lambda_c^+)$, where we have used the parameters $g_2 = 0.565$, $h_2 = 0.63$, and $h_8 = 0.85 \times 10^{-3} \text{ MeV}^{-1}$.

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

$$|h_8| \approx (0.85_{-0.08}^{+0.11}) \times 10^{-3} \text{ MeV}^{-1}, \quad (3.10)$$

which improves the previous limit (3.8) by a factor of 3. The calculated partial widths of $\Lambda_c(2625)^+$ shown in Table 9 are consistent with experimental limits.

4 Lifetimes

4.1 Singly charmed baryons

Among singly charmed baryons, the antitriplet states Λ_c^+ , Ξ_c^+ , Ξ_c^0 , and the Ω_c^0 baryon in the sextet decay weakly. In 2006, the world averages of their lifetimes were [73]

$$\begin{aligned} \tau(\Lambda_c^+) &= (200 \pm 6) \times 10^{-15} \text{ s}, \\ \tau(\Xi_c^+) &= (442 \pm 26) \times 10^{-15} \text{ s}, \\ \tau(\Xi_c^0) &= (112_{-10}^{+13}) \times 10^{-15} \text{ s}, \\ \tau(\Omega_c^0) &= (69 \pm 12) \times 10^{-15} \text{ s}. \end{aligned} \quad (4.1)$$

These results remain the same even in 2014 [8]. As we shall see below, the lifetime hierarchy $\tau(\Xi_c^+) > \tau(\Lambda_c^+) > \tau(\Xi_c^0) > \tau(\Omega_c^0)$ is understood qualitatively but not quantitatively in the operator product expansion (OPE) approach.

On the basis of the OPE approach to analysis of inclusive weak decays, the inclusive rate of the charmed baryon is schematically represented by

$$\Gamma(\mathcal{B}_c \rightarrow f) = \frac{G_F^2 m_c^5}{192\pi^3} V_{\text{CKM}} \cdot \left(A_0 + \frac{A_2}{m_c^2} + \frac{A_3}{m_c^3} + \mathcal{O}\left(\frac{1}{m_c^4}\right) \right), \quad (4.2)$$

where V_{CKM} is the relevant Cabibbo-Kobayashi-Maskawa matrix element. The A_0 term comes from the c quark decay and is common to all charmed hadrons. There are no linear $1/m_Q$ corrections to the inclusive decay rate owing to the lack of gauge-invariant dimension-four operators [81–84], a consequence known as Luke’s theorem [85]. Nonperturbative corrections start at order $1/m_Q^2$ and are model-independent. Spectator effects in inclusive decays due to the Pauli interference and W -exchange contributions account for the $1/m_c^3$ corrections and they have two noteworthy features: First, the estimate of the spectator effects is model dependent; the hadronic four-quark matrix elements are usually evaluated by assuming the factorization approximation for mesons and the quark model for baryons. Second, there is a two-body phase-space enhancement factor of $16\pi^2$ for spectator effects relative to the three-body phase space

Table 10 Various contributions to the decay rates (in units of 10^{-12} GeV) of singly charmed baryons [86]. Experimental values are taken from Ref. [8].

| | Γ^{dec} | Γ^{ann} | Γ_{-}^{int} | Γ_{+}^{int} | Γ_{SL} | Γ^{tot} | $\tau(10^{-13} \text{ s})$ | $\tau_{\text{expt}}(10^{-13} \text{ s})$ |
|---------------|-----------------------|-----------------------|---------------------------|---------------------------|----------------------|-----------------------|----------------------------|--|
| Λ_c^+ | 1.006 | 1.342 | -0.196 | | 0.323 | 2.492 | 2.64 | 2.00 ± 0.06 |
| Ξ_c^+ | 1.006 | 0.071 | -0.203 | 0.364 | 0.547 | 1.785 | 3.68 | 4.42 ± 0.26 |
| Ξ_c^0 | 1.006 | 1.466 | | 0.385 | 0.547 | 3.404 | 1.93 | $1.12_{-0.10}^{+0.13}$ |
| Ω_c^0 | 1.132 | 0.439 | | 1.241 | 1.039 | 3.851 | 1.71 | 0.69 ± 0.12 |

for heavy quark decay. This implies that spectator effects, which are of the order of $1/m_c^3$, are comparable to and even exceed the $1/m_c^2$ terms.

The total width of the charmed baryon \mathcal{B}_c generally receives contributions from inclusive nonleptonic and semileptonic decays: $\Gamma(\mathcal{B}_c) = \Gamma_{\text{NL}}(\mathcal{B}_c) + \Gamma_{\text{SL}}(\mathcal{B}_c)$. The nonleptonic contribution can be decomposed into

$$\Gamma_{\text{NL}}(\mathcal{B}_c) = \Gamma^{\text{dec}}(\mathcal{B}_c) + \Gamma^{\text{ann}}(\mathcal{B}_c) + \Gamma_{-}^{\text{int}}(\mathcal{B}_c) + \Gamma_{+}^{\text{int}}(\mathcal{B}_c), \quad (4.3)$$

corresponding to the c -quark decay, W -exchange contribution, and destructive and constructive Pauli interference. The inclusive decay rate is known to be governed by the imaginary part of an effective nonlocal forward transition operator T . Therefore, Γ^{dec} corresponds to the imaginary part of Fig. 6(a) sandwiched between the same \mathcal{B}_c states. At the Cabibbo-allowed level, Γ^{dec} represents the decay rate of $c \rightarrow su\bar{d}$, and Γ^{ann} denotes the contribution from the W -exchange diagram $cd \rightarrow us$. The interference Γ_{-}^{int} (Γ_{+}^{int}) arises from destructive (constructive) interference between the u (s) quark produced in c -quark decay and the spectator u (s) quark in the charmed baryon \mathcal{B}_c . Note that the constructive Pauli interference is unique to the charmed baryon sector, as it does not occur in the bottom sector. From the quark content of the charmed baryons, it is clear that at the Cabibbo-allowed level, destructive interference occurs in Λ_c^+ and Ξ_c^+ decays (Fig. 6(c)), whereas Ξ_c^+ , Ξ_c^0 and Ω_c^0 can have constructive interference Γ_{+}^{int} (Fig. 6(d)). Because Ω_c^0

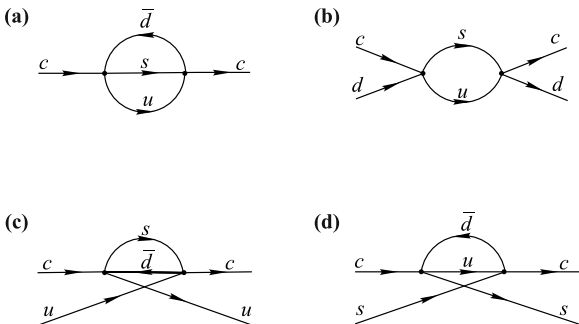


Fig. 6 Contributions to nonleptonic decay rates of charmed baryons from four-quark operators: (a) c -quark decay, (b) W -exchange, (c) destructive Pauli interference, and (d) constructive interference.

contains two s quarks, it is natural to expect that $\Gamma_{+}^{\text{int}}(\Omega_c^0) \gg \Gamma_{+}^{\text{int}}(\Xi_c)$. The W -exchange contribution (Fig. 6(b)) occurs only for Ξ_c^0 and Λ_c^+ at the same Cabibbo-allowed level. In the heavy quark expansion approach, the above-mentioned spectator effects can be described in terms of the matrix elements of local four-quark operators.

The inclusive nonleptonic rates of charmed baryons in the valence quark approximation and in the limit $m_s/m_c = 0$ are expressed as [86]:

$$\begin{aligned} \Gamma_{\text{NL}}(\Lambda_c^+) &= \Gamma^{\text{dec}}(\Lambda_c^+) + \cos^2\theta_C \Gamma^{\text{ann}} + \Gamma_{-}^{\text{int}} + \sin^2\theta_C \Gamma_{+}^{\text{int}}, \\ \Gamma_{\text{NL}}(\Xi_c^+) &= \Gamma^{\text{dec}}(\Xi_c^+) + \sin^2\theta_C \Gamma^{\text{ann}} + \Gamma_{-}^{\text{int}} + \cos^2\theta_C \Gamma_{+}^{\text{int}}, \\ \Gamma_{\text{NL}}(\Xi_c^0) &= \Gamma^{\text{dec}}(\Xi_c^0) + \Gamma^{\text{ann}} + \Gamma_{+}^{\text{int}}, \\ \Gamma_{\text{NL}}(\Omega_c^0) &= \Gamma^{\text{dec}}(\Omega_c^0) + 6 \sin^2\theta_C \Gamma^{\text{ann}} + \frac{10}{3} \cos^2\theta_C \Gamma_{+}^{\text{int}}, \end{aligned} \quad (4.4)$$

where θ_C is the Cabibbo angle.

The results of a model calculation in Ref. [86] are shown in Table 10. The lifetime pattern

$$\tau(\Xi_c^+) > \tau(\Lambda_c^+) > \tau(\Xi_c^0) > \tau(\Omega_c^0) \quad (4.5)$$

clearly agrees with experiment. This lifetime hierarchy is qualitatively understandable. The Ξ_c^+ baryon is the longest-lived among charmed baryons because of the smallness of W -exchange and partial cancellation between constructive and destructive Pauli interference, whereas Ω_c is the shortest-lived owing to the presence of two s quarks in the Ω_c , which greatly enhances the contribution of Γ_{+}^{int} . Because Γ_{+}^{int} is always positive, Γ_{-}^{int} is negative, and the constructive interference has a larger magnitude than the destructive interference, this explains why $\tau(\Xi_c^+) > \tau(\Lambda_c^+)$. It is also clear from Table 10 that, although the qualitative feature of the lifetime pattern is comprehensive, the quantitative estimates of charmed baryon lifetimes and their ratios are still rather poor.

4.2 Doubly charmed baryons

The inclusive nonleptonic rates of doubly charmed baryons in the valence quark approximation and in the

Table 11 Predicted lifetimes of doubly charmed baryons in units of 10^{-13} s.

| | Kiselev <i>et al.</i> [87, 88] | Guberina <i>et al.</i> [89] | Chang <i>et al.</i> [90] | Karliner <i>et al.</i> [52] |
|-----------------|--------------------------------|-----------------------------|--------------------------|-----------------------------|
| Ξ_{cc}^{++} | 4.6 ± 0.5 | 10.5 | 6.7 | 1.85 |
| Ξ_{cc}^+ | 1.6 ± 0.5 | 2.0 | 2.5 | 0.53 |
| Ω_{cc}^+ | 2.7 ± 0.6 | 3.0 | 2.1 | |

limit $m_s/m_c = 0$ are expressed as

$$\begin{aligned} \Gamma_{\text{NL}}(\Xi_{cc}^+) &= \Gamma^{\text{dec}}(\Xi_{cc}^+) + \cos \theta_C^2 \Gamma^{\text{ann}} + \sin \theta_C^2 \Gamma_+^{\text{int}}, \\ \Gamma_{\text{NL}}(\Xi_{cc}^{++}) &= \Gamma^{\text{dec}}(\Xi_{cc}^{++}) + \Gamma_-^{\text{int}}, \\ \Gamma_{\text{NL}}(\Omega_{cc}^+) &= \Gamma^{\text{dec}}(\Omega_{cc}^+) + \sin \theta_C^2 \Gamma^{\text{ann}} + \cos \theta_C^2 \Gamma_+^{\text{int}}. \end{aligned} \quad (4.6)$$

Because Γ_+^{int} is positive and Γ_-^{int} is negative, it is obvious that Ξ_{cc}^{++} is longest-lived, whereas Ξ_{cc}^+ (Ω_{cc}^+) is the shortest-lived if $\Gamma_+^{\text{int}} > \Gamma^{\text{ann}}$ ($\Gamma_+^{\text{int}} < \Gamma^{\text{ann}}$). In general, we have

$$\tau(\Xi_{cc}^{++}) \gg \tau(\Omega_{cc}^+) \sim \tau(\Xi_{cc}^+). \quad (4.7)$$

The predictions available in the literature are summarized in Table 11. Note that the lifetime of Ξ_{cc}^+ was measured by SELEX to be $\tau(\Xi_{cc}^+) < 0.33 \times 10^{-13}$ s [46].

Because the mass splitting between Ξ_{cc}^* and Ξ_{cc} and between Ω_{cc}^* and Ω_{cc} is less than 100 MeV (see also Eq. (2.6) for the lattice calculations),

$$\begin{aligned} m_{\Xi_{cc}^*} - m_{\Xi_{cc}} &= m_{\Sigma_c^*} - m_{\Sigma_c} \approx 65 \text{ MeV}, \\ m_{\Omega_{cc}^*} - m_{\Omega_{cc}} &= m_{\Omega_c^*} - m_{\Omega_c} \approx 71 \text{ MeV}, \end{aligned} \quad (4.8)$$

it is clear that only electromagnetic decays are allowed for Ω_{cc}^* and Ξ_{cc}^* .

5 Hadronic weak decays

5.1 Nonleptonic decays

In contrast to the significant advances made over the last 10 years or so in the study of hadronic weak decays in the bottom baryon sector, progress in the arena of charmed baryons, both theoretical and experimental, has been very slow.

In the naive factorization approach, the coefficients a_1 for the external W -emission amplitude and a_2 for internal W -emission are given by $(c_1 + \frac{c_2}{N_c})$ and $(c_2 + \frac{c_1}{N_c})$, respectively. However, we have learned from charmed meson decays that the naive factorization approach never works for the decay rate of color-suppressed decay modes, although it usually works for color-allowed decays. Empirically, it was learned in the 1980s that if the Fierz-transformed terms characterized by $1/N_c$ are

dropped, the discrepancy between theory and experiment is greatly reduced [91–93]. This leads to the so-called large- N_c approach to describing hadronic D decays [94]. Theoretically, explicit calculations based on QCD sum-rule analysis [95–97] indicate that the Fierz terms are indeed largely compensated by the nonfactorizable corrections.

As the $1/N_c$ expansion method greatly reduces the discrepancy between theory and experiment for charmed meson decays, it is natural to ask if this scenario also works in the baryon sector. This issue can be settled by experimental measurement of the Cabibbo-suppressed mode $\Lambda_c^+ \rightarrow p\phi$, which receives contributions only from the factorizable diagrams. As pointed out in Ref. [98], the rate predicted by the large- N_c approach is in good agreement with the measured value. In contrast, its decay rate predicted by the naive factorization approximation will be too small by a factor of 15. Therefore, the $1/N_c$ approach also works for the factorizable amplitude of charmed baryon decays. This also implies that the inclusion of nonfactorizable contributions is inevitable and necessary. If nonfactorizable effects amount to a redefinition of the effective parameters a_1 , a_2 , and are universal (i.e., channel-independent) in charm decays, then we still have a new factorization scheme with the universal parameters a_1 , a_2 to be determined from experiment.

It is known that for heavy mesons, the nonfactorizable contributions will render color suppression of internal W -emission ineffective. However, W -exchange in baryon decays is not subject to color suppression even in the absence of nonfactorizable terms. A simple way to see this is to consider the large- N_c limit. Although the W -exchange diagram is down by a factor of $1/N_c$ relative to the external W -emission one, this difference is compensated by the fact that the baryon contains N_c quarks in the limit of large N_c , thus allowing N_c different possibilities for W exchange between heavy and light quarks [100]. That is, the pole contribution can be as important as the factorizable one. Experimental measurement of the decay modes $\Lambda_c^+ \rightarrow \Xi^0 K^+, \Delta^{++} K^-$, which proceed only through the W -exchange contributions, indicates that W -exchange indeed plays an essential role in charmed baryon decays.

Various theoretical approaches to weak decays of heavy baryons have been investigated, including the current algebra approach [101–117], factorization scheme, pole model technique [98, 118–122], relativistic quark model [100, 123], and quark diagram scheme [124, 125]. Model predictions of the branching fractions and decay asymmetries can be found in Tables VI–VII of [1] for $\mathcal{B}_c \rightarrow \mathcal{B} + P$ decays, Table VIII for $\mathcal{B}_c \rightarrow \mathcal{B} + V$ decays and Table IX for $\mathcal{B}_c \rightarrow \mathcal{B}(\frac{3}{2}^+) + P(V)$ decays.

Table 12 Branching fractions of the Cabibbo-allowed two-body decays of Λ_c^+ in units of %. Data are taken from PDG [8] except that the absolute branching fraction $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+) = (5.0 \pm 1.3)\%$ is replaced by the new measurement of $(6.84 \pm 0.24_{-0.27}^{+0.21})\%$ by Belle [126].

| Decay | \mathcal{B} | Decay | \mathcal{B} | Decay | \mathcal{B} |
|---|-----------------|--|-----------------|---|-----------------|
| $\Lambda_c^+ \rightarrow \Lambda\pi^+$ | 1.46 ± 0.13 | $\Lambda_c^+ \rightarrow \Lambda\rho^+$ | < 6.5 | $\Lambda_c^+ \rightarrow \Delta^{++}K^-$ | 1.16 ± 0.07 |
| $\Lambda_c^+ \rightarrow \Sigma^0\pi^+$ | 1.44 ± 0.14 | $\Lambda_c^+ \rightarrow \Sigma^0\rho^+$ | | $\Lambda_c^+ \rightarrow \Sigma^{*0}\pi^+$ | |
| $\Lambda_c^+ \rightarrow \Sigma^+\pi^0$ | 1.37 ± 0.30 | $\Lambda_c^+ \rightarrow \Sigma^+\rho^0$ | < 1.9 | $\Lambda_c^+ \rightarrow \Sigma^{*+}\pi^0$ | |
| $\Lambda_c^+ \rightarrow \Sigma^+\eta$ | 0.75 ± 0.11 | $\Lambda_c^+ \rightarrow \Sigma^+\omega$ | 3.7 ± 1.0 | $\Lambda_c^+ \rightarrow \Sigma^{*+}\eta$ | 1.16 ± 0.35 |
| $\Lambda_c^+ \rightarrow \Sigma^+\eta'$ | | $\Lambda_c^+ \rightarrow \Sigma^+\phi$ | 0.42 ± 0.07 | $\Lambda_c^+ \rightarrow \Sigma^{*+}\eta'$ | |
| $\Lambda_c^+ \rightarrow \Xi^0 K^+$ | 0.53 ± 0.13 | $\Lambda_c^+ \rightarrow \Xi^0 K^{*+}$ | 0.53 ± 0.19 | $\Lambda_c^+ \rightarrow \Xi^{*0} K^+$ | 0.36 ± 0.10 |
| $\Lambda_c^+ \rightarrow p\bar{K}^0$ | 3.2 ± 0.3 | $\Lambda_c^+ \rightarrow p\bar{K}^{*0}$ | 2.1 ± 0.3 | $\Lambda_c^+ \rightarrow \Delta^+\bar{K}^0$ | 1.36 ± 0.44 |

5.2 Discussion

5.2.1 Λ_c^+ decays

Experimentally, nearly all the branching fractions of the Λ_c^+ are measured relative to the $pK^-\pi^+$ mode. On the basis of ARGUS and CLEO data, the PDG has made a model-dependent determination of the absolute branching fraction, $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+) = (5.0 \pm 1.3)\%$ [8]. Recently, Belle reported a value of $(6.84 \pm 0.24_{-0.27}^{+0.21})\%$ [126] from the reconstruction of $D^*\pi\pi$ recoiling against the Λ_c^+ production in e^+e^- annihilation. Hence, the uncertainties are much smaller, and, most importantly, this measurement is model independent! More recently, BESIII has also measured this mode directly with the preliminary result $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+) = (5.84 \pm 0.27 \pm 0.23)\%$ [127]. Its precision is comparable to that of Belle's result. Another approach is to exploit a particular decay, $B^+ \rightarrow p\pi^+\pi^+\bar{\Sigma}_c^{*-}$, and its charge conjugate to measure $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+)$, also in a model independent manner [128].

The branching fractions of the Cabibbo-allowed two-body decays of Λ_c^+ are listed in Table 12. The central values of data taken from the PDG [8] are scaled up by a factor of 1.37 because of the new measurement of $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+)$ by Belle [126]. BESIII recently measured 2-body, 3-body, and 4-body decay modes of Λ_c^+ with significantly improved precision [127]. For example, the result $\mathcal{B}(\Lambda_c^+ \rightarrow \Lambda\pi^+) = (1.24 \pm 0.07 \pm 0.03)\%$ obtained by BESIII has much better precision than the value of $(1.07 \pm 0.28)\%$ quoted by the PDG [8].

Many of the Λ_c^+ decay modes such as $\Sigma^+K^+K^-$, $\Sigma^+\phi$, $\Xi^{(*)}K^{(*)+}$, and $\Delta^{++}K^-$ can proceed only through W -exchange. Experimental measurement of them implies the importance of W -exchange, which is not subject to color suppression in charmed baryon decays.

Some Cabibbo-suppressed modes such as $\Lambda_c^+ \rightarrow \Lambda K^+$ and $\Lambda_c^+ \rightarrow \Sigma^0 K^+$ have been measured by Belle [129] and

BaBar [130], respectively. Their branching fractions are of the order of $10^{-3} - 10^{-4}$. The first measured Cabibbo-suppressed mode $\Lambda_c^+ \rightarrow p\phi$, is of particular interest because it receives contributions only from the factorizable diagram and is expected to be color suppressed in the naive factorization approach. A calculation in Refs. [131, 132] yields

$$\begin{aligned} \mathcal{B}(\Lambda_c^+ \rightarrow p\phi) &= 2.26 \times 10^{-3} a_2^2, \\ \alpha(\Lambda_c^+ \rightarrow p\phi) &= -0.10. \end{aligned} \quad (5.1)$$

From the experimental measurement $\mathcal{B}(\Lambda_c^+ \rightarrow p\phi) = (11.2 \pm 2.3) \times 10^{-4}$ [8]²⁾, it follows that

$$|a_2|_{\text{expt}} = 0.70 \pm 0.07. \quad (5.2)$$

This is consistent with the result obtained by the $1/N_c$ approach, $a_2 = c_2(m_c) \approx -0.59$.

All the models except that of Ref. [121] predict a positive decay asymmetry α for the decay $\Lambda_c^+ \rightarrow \Sigma^+\pi^0$ (see Table VII of Ref. [1]). Therefore, the measurement of $\alpha = -0.45 \pm 0.31 \pm 0.06$ by CLEO [133] is very surprising. If the negative sign of α is confirmed in the future, this will imply that the s -wave and p -wave amplitudes for this decay have opposite signs, contrary to the model expectation. The implication of this has been discussed in detail in Refs. [98, 99]. Because the error of the previous CLEO measurement is very large, it is crucial to more accurately measure the decay asymmetry for $\Lambda_c^+ \rightarrow \Sigma^+\pi^0$.

5.2.2 Ξ_c^+ decays

No absolute branching fractions have been measured. The branching ratios listed in Tables VI and VIII of Ref. [1] are those relative to $\Xi_c^+ \rightarrow \Xi^-\pi^+\pi^+$. Several Cabibbo-suppressed decay modes such as $p\bar{K}^{*0}$, $\Sigma^+\phi$, $\Sigma^+\pi^+\pi^-$, $\Sigma^-\pi^+\pi^+$, and $\Xi(1690)K^+$ have been observed [8].

The Cabibbo-allowed decays $\Xi_c^+ \rightarrow \mathcal{B}(3/2^+) + P$ have been studied and are believed to be forbidden, as they

²⁾ We have scaled up the PDG number $(8.7 \pm 2.7) \times 10^{-4}$ [8] by a factor of 1.37 for its central value.

do not receive factorizable and $1/2^\pm$ pole contributions [100, 120]. However, the $\Sigma^{*+} \bar{K}^0$ mode was seen earlier by FOCUS [134], and this may indicate the importance of pole contributions beyond low-lying $1/2^\pm$ intermediate states.

5.2.3 Ξ_c^0 decays

No absolute branching fractions have been measured to date. However, there are several measurements of the ratios of branching fractions, for example [8],

$$R_1 = \frac{\Gamma(\Xi_c^0 \rightarrow \Lambda K_S^0)}{\Gamma(\Xi_c^0 \rightarrow \Xi^- \pi^+)} = 0.21 \pm 0.02 \pm 0.02,$$

$$R_2 = \frac{\Gamma(\Xi_c^0 \rightarrow \Omega^- K^+)}{\Gamma(\Xi_c^0 \rightarrow \Xi^- \pi^+)} = 0.297 \pm 0.024. \quad (5.3)$$

The decay modes $\Xi_c^0 \rightarrow \Omega^- K^+$ and $\Sigma^+ K^-$ and $\Sigma^+ \pi^-$ proceed only through W -exchange. The measured branching ratio of $\Omega^- K^+$ relative to $\Xi^- \pi^+$ implies that the W -exchange mechanism plays a significant role. The model of Körner and Krämer [100] predicts $R_2 = 0.33$ (see Table IX of Ref. [1]), in agreement with experiment, but its prediction $R_1 = 0.06$ is too small compared to the data.

5.2.4 Ω_c^0 decays

A unique feature of Ω_c^0 decays is that the decay $\Omega_c^0 \rightarrow \Omega^- \pi^+$ proceeds only via external W -emission, whereas $\Omega_c^0 \rightarrow \Xi^{*0} \bar{K}^0$ proceeds via the factorizable internal W -emission diagram. Various model predictions of Cabibbo-allowed $\Omega_c^0 \rightarrow \mathcal{B}(3/2^+) + P(V)$ are listed in Table IX of [1] with the unknown parameters a_1 and a_2 . From the decay $\Lambda_c^+ \rightarrow p \phi$ we learn that $|a_2| = 0.70 \pm 0.07$. The hadronic weak decays of the Ω_c^0 were recently studied in great detail in Ref. [135], where most of the decay channels in Ω_c^0 decays were found to proceed only through the W -exchange diagram; moreover, the W -exchange contributions dominate in the rest of the processes, with some exceptions. Observation of such decays would shed light on the mechanism of W -exchange effects in these decay modes.

5.3 Charm-flavor-conserving nonleptonic decays

There is a special class of weak decays of charmed baryons that can be studied reliably, namely, heavy-flavor-conserving nonleptonic decays. Some examples are the singly Cabibbo-suppressed decays $\Xi_c \rightarrow \Lambda_c \pi$ and $\Omega_c \rightarrow \Xi'_c \pi$. The idea is simple: In these decays, only the light quarks inside the heavy baryon will partici-

pate in weak interactions; that is, while the two light quarks undergo weak transitions, the heavy quark behaves as a “spectator”. As the emitted light mesons are soft, the $\Delta S = 1$ weak interactions among light quarks can be handled by the well known short-distance effective Hamiltonian. This special class of weak decays can usually be calculated more reliably than the conventional charmed baryon weak decays. The synthesis of the heavy quark and chiral symmetries provides a natural setting for investigating these reactions [136]. The weak decays $\Xi_Q \rightarrow \Lambda_Q \pi$ with $Q = c, b$ were also studied in Refs. [137–139].

The combined symmetries of heavy and light quarks severely restrict the weak interactions allowed. In the symmetry limit, it is found that $\mathcal{B}_3 - \mathcal{B}_6$ and $\mathcal{B}_6^* - \mathcal{B}_6$ nonleptonic weak transitions [136] cannot occur. Symmetries alone permit three types of transitions: $\mathcal{B}_3 - \mathcal{B}_3$, $\mathcal{B}_6 - \mathcal{B}_6$ and $\mathcal{B}_6^* - \mathcal{B}_6$ transitions. However, in both the MIT bag and diquark models, only $\mathcal{B}_3 - \mathcal{B}_3$ transitions have nonzero amplitudes. The general amplitude for $\mathcal{B}_i \rightarrow \mathcal{B}_f + P$ is given by

$$M(\mathcal{B}_i \rightarrow \mathcal{B}_f + P) = i\bar{u}_f(A - B\gamma_5)u_i, \quad (5.4)$$

where A and B are the S - and P -wave amplitudes, respectively. The S -wave amplitude can be evaluated using current algebra in terms of the parity-violating commutator term. For example, the S -wave amplitude of $\Xi_c^+ \rightarrow \Lambda_c^+ \pi^0$ is given by

$$A(\Xi_c^+ \rightarrow \Lambda_c^+ \pi^0) = -\frac{1}{\sqrt{2}f_\pi} \langle \Lambda_c^+ \uparrow | \mathcal{H}_{\text{eff}}^{\text{PC}} | \Xi_c^+ \uparrow \rangle, \quad (5.5)$$

whereas the P -wave amplitude arises from the ground-state baryon poles [136]:

$$B(\Xi_c^+ \rightarrow \Lambda_c^+ \pi^0) = \frac{g_2}{2f_\pi} \frac{m_{\Xi_c} + m_{\Xi'_c}}{m_{\Lambda_c} - m_{\Xi_{c2}}} \langle \Lambda_c^+ \uparrow | \mathcal{H}_{\text{eff}}^{\text{PC}} | \Xi_c^+ \uparrow \rangle \sin \phi, \quad (5.6)$$

where ϕ is the mixing angle of Ξ_c , and Ξ'_c and Ξ_{c1}, Ξ_{c2} are their mass eigenstates. The matrix element $\langle \Lambda_c^+ \uparrow | \mathcal{H}_{\text{eff}}^{\text{PC}} | \Xi_c^+ \uparrow \rangle$ was evaluated in Ref. [136] using two different models: the MIT bag model [140, 141] and the diquark model.

The predicted rates are [136]

$$\begin{aligned} \Gamma(\Xi_c^0 \rightarrow \Lambda_c^+ \pi^-) &= 1.7 \times 10^{-15} \text{ GeV}, \\ \Gamma(\Xi_c^+ \rightarrow \Lambda_c^+ \pi^0) &= 1.0 \times 10^{-15} \text{ GeV}, \\ \Gamma(\Omega_c^0 \rightarrow \Xi'^+ \pi^-) &= 4.3 \times 10^{-17} \text{ GeV}, \end{aligned} \quad (5.7)$$

and the corresponding branching fractions are

$$\begin{aligned} \mathcal{B}(\Xi_c^0 \rightarrow \Lambda_c^+ \pi^-) &= 2.9 \times 10^{-4}, \\ \mathcal{B}(\Xi_c^+ \rightarrow \Lambda_c^+ \pi^0) &= 6.7 \times 10^{-4}, \end{aligned}$$

$$\mathcal{B}(\Omega_c^0 \rightarrow \Xi_c^{\prime+} \pi^-) = 4.5 \times 10^{-6}. \quad (5.8)$$

As stated above, the $\mathcal{B}_6 - \mathcal{B}_6$ transition $\Omega_c^0 \rightarrow \Xi_c^{\prime+} \pi^-$ vanishes in the chiral limit. It receives a finite factorizable contribution as a result of a symmetry-breaking effect. At any rate, the predicted branching fractions for the charm-flavor-conserving decays $\Xi_c^0 \rightarrow \Lambda_c^+ \pi^-$ and $\Xi_c^+ \rightarrow \Lambda_c^+ \pi^0$ are of the order of $10^{-3} - 10^{-4}$ and should be readily accessible in the near future.

5.4 Semileptonic decays

Exclusive semileptonic decays of charmed baryons: $\Lambda_c^+ \rightarrow \Lambda e^+ (\mu^+) \nu_e$, $\Xi_c^+ \rightarrow \Xi^0 e^+ \nu_e$ and $\Xi_c^0 \rightarrow \Xi^- e^+ \nu_e$ have been observed experimentally. Their rates depend on the $\mathcal{B}_c \rightarrow \mathcal{B}$ form factors $f_i(q^2)$ and $g_i(q^2)$ ($i = 1, 2, 3$), which are defined as

$$\begin{aligned} & \langle \mathcal{B}_f(p_f) | V_\mu - A_\mu | \mathcal{B}_c(p_i) \rangle \\ &= \bar{u}_f(p_f) [f_1(q^2) \gamma_\mu + i f_2(q^2) \sigma_{\mu\nu} q^\nu + f_3(q^2) q_\mu \\ & - (g_1(q^2) \gamma_\mu + i g_2(q^2) \sigma_{\mu\nu} q^\nu + g_3(q^2) q_\mu) \gamma_5] u_i(p_i). \end{aligned} \quad (5.9)$$

These form factors have been evaluated using the non-relativistic quark model [131, 132, 142–144, 151], MIT bag model [142, 143], relativistic quark model [145, 149], light-front quark model [146], and QCD sum rules [147, 148, 150]. Experimentally, the only information available to date is the form-factor ratio measured in the semileptonic decay $\Lambda_c \rightarrow \Lambda e \bar{\nu}$. In the heavy quark limit, the six $\Lambda_c \rightarrow \Lambda$ form factors are reduced to two:

$$\begin{aligned} \langle \Lambda(p) | \bar{s} \gamma_\mu (1 - \gamma_5) c | \Lambda_c(v) \rangle &= \bar{u}_\Lambda (F_1^{\Lambda_c \Lambda}(v \cdot p) \\ &+ \not{v} F_2^{\Lambda_c \Lambda}(v \cdot p)) \gamma_\mu (1 - \gamma_5) u_{\Lambda_c}. \end{aligned} \quad (5.10)$$

Assuming that the form factors exhibit dipole q^2 behavior, the ratio $R = \tilde{F}_2^{\Lambda_c \Lambda} / \tilde{F}_1^{\Lambda_c \Lambda}$ is measured by CLEO to be [152]

$$R = -0.31 \pm 0.05 \pm 0.04. \quad (5.11)$$

Various model predictions of the charmed baryon semileptonic decay rates and decay asymmetries are shown in Table 13. Dipole q^2 dependence of the form factors is assumed whenever the form factor momentum dependence is not available in the model. Four different sets of predictions for $\Lambda_c^+ \rightarrow n e^+ \nu_e$, which are not listed in Table 13, were presented in the sum rule calculations of Ref. [153]. The semileptonic decays of Ω_c were treated in Ref. [154] within the framework of a constituent quark model. From Table 13, we see that the computed branching fractions of $\Lambda_c^+ \rightarrow \Lambda e^+ \nu$, which fall in the range $1.4\% \sim 2.6\%$ are slightly smaller than the experimental values, $(2.9 \pm 0.5)\%$ [$(2.1 \pm 0.6)\%$], of the PDG [8]. The branching fractions of $\Xi_c^0 \rightarrow \Xi^- e^+ \nu$

and $\Xi_c^+ \rightarrow \Xi^0 e^+ \nu$ are predicted to lie in the ranges $(0.8 - 2.0)\%$ and $(3.3 - 8.1)\%$, respectively, except that the QCD sum rule calculation in Ref. [150] predicts a much larger rate for $\Xi_c \rightarrow \Xi e^+ \nu_e$. Experimentally, only the ratios of the branching fractions are available to date [8]:

$$\begin{aligned} \frac{\Gamma(\Xi_c^+ \rightarrow \Xi^0 e^+ \nu)}{\Gamma(\Xi_c^+ \rightarrow \Xi^- \pi^+ \pi^+)} &= 2.3 \pm 0.6_{-0.6}^{+0.3}, \\ \frac{\Gamma(\Xi_c^0 \rightarrow \Xi^- e^+ \nu)}{\Gamma(\Xi_c^0 \rightarrow \Xi^- \pi^+)} &= 3.1 \pm 1.0_{-0.5}^{+0.3}. \end{aligned} \quad (5.12)$$

There have been active studies of semileptonic decays of doubly charmed baryons. The interested reader can consult [155–159] for further references.

Just as with the hadronic decays discussed in the last subsection, there are also heavy-flavor-conserving semileptonic processes, for example, $\Xi_c^0 \rightarrow \Lambda_c^+ (\Sigma_c^+) e^- \bar{\nu}_e$ and $\Omega_c^0 \rightarrow \Xi_c^+ e^- \bar{\nu}_e$. In these decays only the light quarks inside the heavy baryon will participate in weak interactions, while the heavy quark behaves as a spectator. This topic was recently investigated in Ref. [139]. Owing to the severe phase-space suppression, the branching fractions are of order 10^{-6} in the best cases, and typically 10^{-7} to 10^{-8} .

6 Electromagnetic and weak radiative decays

Although radiative decays are well measured in the charmed meson sector, e.g., $D^* \rightarrow D \gamma$ and $D_s^{*+} \rightarrow D_s^+ \gamma$, only three of the radiative modes in the charmed baryon sector have been seen, namely, $\Xi_c^{0\prime} \rightarrow \Xi_c^0 \gamma$, $\Xi_c^{+\prime} \rightarrow \Xi_c^+ \gamma$ and $\Omega_c^{*0} \rightarrow \Omega_c^0 \gamma$. This is understandable because $m_{\Xi_c} - m_{\Xi_c} \approx 108$ MeV, and $m_{\Omega_c^*} - m_{\Omega_c} \approx 71$ MeV. Hence, Ξ_c^{\prime} and Ω_c^* are governed by the electromagnetic decays. However, it will be difficult to measure the rates of these decays because these states are too narrow to be experimentally resolvable. Nevertheless, we shall systematically study the two-body electromagnetic decays of charmed baryons and also weak radiative decays.

6.1 Electromagnetic decays

In the charmed baryon sector, the following two-body electromagnetic decays are of interest:

$$\begin{aligned} \mathcal{B}_6 &\rightarrow \mathcal{B}_3 + \gamma : \Sigma_c \rightarrow \Lambda_c + \gamma, \quad \Xi_c^{\prime} \rightarrow \Xi_c + \gamma, \\ \mathcal{B}_6^* &\rightarrow \mathcal{B}_3 + \gamma : \Sigma_c^* \rightarrow \Lambda_c + \gamma, \quad \Xi_c^{\prime*} \rightarrow \Xi_c + \gamma, \\ \mathcal{B}_6^* &\rightarrow \mathcal{B}_6 + \gamma : \Sigma_c^* \rightarrow \Sigma_c + \gamma, \quad \Xi_c^{\prime*} \rightarrow \Xi_c^{\prime} + \gamma, \\ \Omega_c^* &\rightarrow \Omega_c + \gamma, \end{aligned} \quad (6.1)$$

where we denote the spin $\frac{1}{2}$ baryons as \mathcal{B}_6 and \mathcal{B}_3 for

Table 13 Predicted semileptonic decay rates (in units of $10^{10}s^{-1}$) and decay asymmetries (second entry) in various models. The absolute branching fraction $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+) = (5.0 \pm 1.3)\%$ is replaced by the new measurement of $(6.84 \pm 0.24_{-0.27}^{+0.21})\%$ by Belle [126] for the data of $\Gamma(\Lambda_c^+ \rightarrow \Lambda^0 \ell^+ \nu_\ell)$ taken from the PDG [8]. Predictions of [142, 143] are obtained in the non-relativistic quark model and the MIT bag model (in parentheses).

| Process | [131, 132] | [142, 143] | [144] | [145] | [146] | [147] | [148] | [149] | [150] | [151] | Expt. [8] |
|---|------------|-------------|-------|-------|-------|----------------|----------------|-------|-----------------|------------|------------------|
| $\Lambda_c^+ \rightarrow \Lambda^0 e^+ \nu_e$ | 7.1 | 11.2 (7.7) | 9.8 | 7.22 | 7.0 | 13.2 ± 1.8 | 10.9 ± 3.0 | | | | 14.4 ± 2.6 |
| $\Lambda_c^+ \rightarrow \Lambda^0 \mu^+ \nu_e$ | 7.1 | 11.2 (7.7) | 9.8 | 7.22 | 7.0 | 13.2 ± 1.8 | 10.9 ± 3.0 | | | | -0.86 ± 0.04 |
| $\Lambda_c^+ \rightarrow n e^+ \nu_e$ | | | | 1.32 | | | | 1.01 | | 0.96, 1.37 | 13.3 ± 2.8 |
| $\Xi_c \rightarrow \Xi e^+ \nu_e$ | 7.4 | 18.1 (12.5) | 8.5 | 8.16 | 9.7 | | | | 64.8 ± 22.6 | | seen |
| $\Xi_c \rightarrow \Sigma e^+ \nu_e$ | | | | | | | | | 3.3 ± 1.7 | | |

the symmetric sextet **6** and antisymmetric antitriplet $\bar{\mathbf{3}}$, respectively, and the spin $\frac{3}{2}$ baryon by \mathcal{B}_6^* .

An ideal theoretical framework for studying the above-mentioned electromagnetic decays is provided by the formalism in which the heavy quark symmetry and the chiral symmetry of light quarks are combined [68–70]. When supplemented by the nonrelativistic quark model, the formalism determines completely the low energy dynamics of heavy hadrons. The electromagnetic interactions of heavy hadrons consist of two distinct contributions: one from gauging electromagnetically the chirally invariant strong interaction Lagrangians for heavy mesons and baryons given in Refs. [68–70], and the other from the anomalous magnetic moment couplings of the heavy particles. The heavy quark symmetry reduces the number of free parameters needed to describe the magnetic couplings to the photon. There are two undetermined parameters for the ground-state heavy baryons. All of these parameters are related simply to the magnetic moments of the light quarks in the nonrelativistic quark model. However, the charmed quark is not particularly heavy ($m_c \simeq 1.6$ GeV), and it carries a charge of $\frac{2}{3}e$. Consequently, the contribution from its magnetic moment cannot be neglected. The chiral and electromagnetic gauge-invariant Lagrangian for heavy baryons can be found in Eqs. (3.8) and (3.9) of Ref. [160], and are denoted by $\mathcal{L}_B^{(1)}$ and $\mathcal{L}_B^{(2)}$, respectively.

The most general gauge-invariant Lagrangian Eq. (3.9) of Ref. [160] for magnetic transitions of heavy baryons can be recast in terms of superfields [161, 162]:

$$\begin{aligned} \mathcal{L}_B^{(2)} = & -i3a_1 \text{tr}(\bar{S}^\mu Q F_{\mu\nu} S^\nu) \\ & + \sqrt{3}a_2 \epsilon_{\mu\nu\alpha\beta} \text{tr}(\bar{S}^\mu Q v^\nu F^{\alpha\beta} T) + h.c. \\ & + 3a'_1 \text{tr}(\bar{S}^\mu Q' \sigma \cdot F S_\mu) - \frac{3}{2}a'_1 \text{tr}(\bar{T} Q' \sigma \cdot FT), \end{aligned} \quad (6.2)$$

where $\sigma \cdot F \equiv \sigma_{\mu\nu} F^{\mu\nu}$, $Q = \text{diag}(2/3, -1/3, -1/3)$ is the charge matrix for the light u, d , and s quarks, and $Q' = e_Q$ is the charge of the heavy quark. In the above equation,

$$T = \mathcal{B}_{\bar{\mathbf{3}}}, \quad S^\mu = \mathcal{B}_6^{*\mu} - \frac{1}{\sqrt{3}}(\gamma^\mu + v^\mu)\gamma_5 \mathcal{B}_6. \quad (6.3)$$

It follows that [161, 162]

$$\begin{aligned} & A[S_{ij}^\mu(v) \rightarrow S_{ij}^\nu + \gamma(\varepsilon, k)] \\ & = i\frac{3}{2}a_1 \bar{U}^\nu (Q_{ii} + Q_{jj})(k_\nu \varepsilon_\mu - k_\mu \varepsilon_\nu) \mathcal{U}^\mu - i6a'_1 Q' \bar{U}^\mu \not{k} \not{\varepsilon} \mathcal{U}_\mu, \\ & A[S_{ij}^\mu(v) \rightarrow T_{ij} + \gamma(\varepsilon, k)] \\ & = -2\sqrt{3/2}a_2 \epsilon_{\mu\nu\alpha\beta} \bar{u}_{\bar{\mathbf{3}}} v^\nu k^\alpha \varepsilon^\beta (Q_{ii} - Q_{jj}) \mathcal{U}^\mu \quad (i < j), \end{aligned} \quad (6.4)$$

where k_μ is the photon 4-momentum, and ε_μ is the polarization 4-vector. As stressed in Ref. [160], SU(3) breaking effects due to light-quark mass differences can be incorporated by replacing the charge matrix Q with

$$Q \rightarrow \tilde{Q} = \text{diag}\left(\frac{2}{3}, -\frac{\alpha}{3}, -\frac{\beta}{3}\right) \quad (6.5)$$

where $\alpha = M_u/M_d$ and $\beta = M_u/M_s$. To avoid any confusion with the current quark mass m_q , we have used capital letters to denote the constituent quark masses. In the quark model, the coefficients a_1 and a_2 are simply related to the Dirac magnetic moments of light quarks,

$$a_1 = -\frac{e}{3} \frac{1}{M_u}, \quad a_2 = \frac{e}{2\sqrt{6}} \frac{1}{M_u}, \quad (6.6)$$

whereas a'_1 is connected to those of heavy quarks. Explicitly, a'_1 is fixed by heavy quark symmetry as

$$a'_1 = \frac{e}{12} \frac{1}{M_Q}. \quad (6.7)$$

Within the framework of HHChPT, the authors of [163] proceeded to construct chiral Lagrangians at the level $\mathcal{O}(p^2)$ and $\mathcal{O}(p^3)$ and then calculated the electromagnetic decay amplitudes of charmed baryons up to $\mathcal{O}(p^3)$. It is not clear whether their $\mathcal{O}(p^2)$ Lagrangian (see Eq. (12) of Ref. [163]), characterized by the four couplings f_2, f_3, \tilde{f}_3 and f_4 , is equivalent to the first two terms of the $\mathcal{O}(p)$ Lagrangian given by Eq. (6.2). The

unknown couplings there were also estimated using the quark model.

The general amplitudes of electromagnetic decays are given by [160]

$$\begin{aligned} A(\mathcal{B}_6 \rightarrow \mathcal{B}_3 + \gamma) &= i\eta_1 \bar{u}_3 \sigma_{\mu\nu} k^\mu \varepsilon^\nu u_6, \\ A(\mathcal{B}_6^* \rightarrow \mathcal{B}_3 + \gamma) &= i\eta_2 \varepsilon_{\mu\nu\alpha\beta} \bar{u}_3 \gamma^\nu k^\alpha \varepsilon^\beta u^\mu, \\ A(\mathcal{B}_6^* \rightarrow \mathcal{B}_6 + \gamma) &= i\eta_3 \varepsilon_{\mu\nu\alpha\beta} \bar{u}_6 \gamma^\nu k^\alpha \varepsilon^\beta u^\mu. \end{aligned} \quad (6.8)$$

The corresponding decay rates are [160]

$$\begin{aligned} \Gamma(\mathcal{B}_6 \rightarrow \mathcal{B}_3 + \gamma) &= \eta_1^2 \frac{k^3}{\pi}, \\ \Gamma(\mathcal{B}_6^* \rightarrow \mathcal{B}_3 + \gamma) &= \eta_2^2 \frac{k^3}{3\pi} \frac{3m_i^2 + m_f^2}{4m_i^2}, \\ \Gamma(\mathcal{B}_6^* \rightarrow \mathcal{B}_6 + \gamma) &= \eta_3^2 \frac{k^3}{3\pi} \frac{3m_i^2 + m_f^2}{4m_i^2}, \end{aligned} \quad (6.9)$$

where m_i (m_f) is the mass of the parent (daughter) baryon. The coupling constants η_i can be calculated using the quark model for a_1 , a_2 , and a'_1 [160, 164]:

$$\begin{aligned} \eta_1(\Sigma_c^+ \rightarrow \Lambda_c^+) &= \frac{e}{6\sqrt{3}} \left(\frac{2}{M_u} + \frac{1}{M_d} \right), \\ \eta_1(\Xi_c'^+ \rightarrow \Xi_c^+) &= \frac{e}{6\sqrt{3}} \left(\frac{2}{M_u} + \frac{1}{M_s} \right), \\ \eta_1(\Xi_c'^0 \rightarrow \Xi_c^0) &= \frac{e}{6\sqrt{3}} \left(\frac{1}{M_s} - \frac{1}{M_d} \right), \\ \eta_2(\Sigma_c^{*+} \rightarrow \Lambda_c^+) &= \frac{e}{3\sqrt{6}} \left(\frac{2}{M_u} + \frac{1}{M_d} \right), \\ \eta_2(\Xi_c'^{*+} \rightarrow \Xi_c^+) &= \frac{e}{3\sqrt{6}} \left(\frac{2}{M_u} + \frac{1}{M_s} \right), \\ \eta_2(\Xi_c'^{*0} \rightarrow \Xi_c^0) &= \frac{e}{3\sqrt{6}} \left(-\frac{1}{M_d} + \frac{1}{M_s} \right), \\ \eta_3(\Sigma_c^{*++} \rightarrow \Sigma_c^{++}) &= \frac{2\sqrt{2}e}{9} \left(\frac{1}{M_u} - \frac{1}{M_c} \right), \\ \eta_3(\Sigma_c^{*0} \rightarrow \Sigma_c^0) &= \frac{2\sqrt{2}e}{9} \left(-\frac{1}{2M_d} - \frac{1}{M_c} \right), \\ \eta_3(\Sigma_c^{*+} \rightarrow \Sigma_c^+) &= \frac{\sqrt{2}e}{9} \left(\frac{1}{M_u} - \frac{1}{2M_d} - \frac{2}{M_c} \right), \\ \eta_3(\Omega_c^{*0} \rightarrow \Omega_c^0) &= \frac{2\sqrt{2}e}{9} \left(-\frac{1}{2M_s} - \frac{1}{M_c} \right), \\ \eta_3(\Xi_c'^{*+} \rightarrow \Xi_c'^+) &= \frac{\sqrt{2}e}{9} \left(\frac{1}{M_u} - \frac{1}{2M_s} - \frac{2}{M_c} \right), \\ \eta_3(\Xi_c'^{*0} \rightarrow \Xi_c'^0) &= \frac{\sqrt{2}e}{9} \left(-\frac{1}{2M_d} - \frac{1}{2M_s} - \frac{2}{M_c} \right). \end{aligned} \quad (6.10)$$

The results calculated using the constituent quark

masses, $M_u = 338$ MeV, $M_d = 322$ MeV, $M_s = 510$ MeV [8], and $M_c = 1.6$ GeV, are summarized in the second column of Table 14. Some other model predictions are also listed there for comparison.

Radiative decays of s -wave charmed baryons are considered in Ref. [165] in the quark model, and the predictions are similar to ours. A similar procedure is followed in Ref. [166] where the heavy quark symmetry is supplemented with light-diquark symmetries to calculate the widths of $\Sigma_c^+ \rightarrow \Lambda_c^+ \gamma$ and $\Sigma_c^* \rightarrow \Sigma_c \gamma$. The authors of Ref. [75] apply the relativistic quark model to predict various electromagnetic decays of charmed baryons. In addition to the magnetic dipole (M1) transition, the author of Ref. [167] also considered and estimated the electric quadrupole (E2) amplitude for $\Sigma_c^{*+} \rightarrow \Lambda_c^+ \gamma$ arising from the chiral loop correction. The E2 contributions were analyzed in detail in Ref. [168]. The E2 amplitudes appear at different higher orders for the three types of decays: $\mathcal{O}(1/\Lambda_\chi^2)$ for $\mathcal{B}_6^* \rightarrow \mathcal{B}_6 + \gamma$, $\mathcal{O}(1/m_Q \Lambda_\chi^2)$ for $\mathcal{B}_6^* \rightarrow \mathcal{B}_3 + \gamma$ and $\mathcal{O}(1/m_Q^3 \Lambda_\chi^2)$ for $\mathcal{B}_6 \rightarrow \mathcal{B}_3 + \gamma$. Therefore, the E2 contribution to $\mathcal{B}_6 \rightarrow \mathcal{B}_3 + \gamma$ is completely negligible. The electromagnetic decays were calculated in Refs. [169–171] using the QCD sum rule method, and they were studied within the framework of the modified bag model in Ref. [172].

It is evident from Table 14 that the predictions in Refs. [160, 164] and [163] all based on HHChPT are quite different for the following three modes: $\Sigma_c^{*++} \rightarrow \Sigma_c^{++} \gamma$, $\Sigma_c^{*+} \rightarrow \Lambda_c^+ \gamma$ and $\Xi_c'^{*+} \rightarrow \Xi_c^+ \gamma$. Indeed, the results for the last two modes in Ref. [163] are larger than all the other existing predictions by one order of magnitude! It is naively expected that all HHChPT approaches should agree with each other to the lowest order of chiral expansion provided that the coefficients are inferred from the nonrelativistic quark model. The lowest order predictions $\Gamma(\Sigma_c^{*+} \rightarrow \Lambda_c^+ \gamma) = 756$ keV and $\Gamma(\Xi_c'^{*+} \rightarrow \Xi_c^+ \gamma) = 403$ keV obtained in Ref. [163] are still very large. Note that a recent lattice calculation in Ref. [174] yields $\Gamma(\Omega_c^* \rightarrow \Omega_c \gamma) = 0.074 \pm 0.008$ keV which is much smaller than the value of $\Gamma(\Omega_c^*) = 4.82$ keV predicted in Ref. [163].

Chiral loop corrections to the M1 electromagnetic decays and to the strong decays of heavy baryons were computed at the one loop order in Refs. [161, 162]. The leading chiral-loop effects we found are nonanalytic in the forms of m/Λ_χ and $(m^2/\Lambda_\chi^2) \ln(\Lambda^2/m^2)$ (or $m_q^{1/2}$ and $m_q \ln m_q$, where m_q is the light quark mass). Some results are [161, 162]

$$\begin{aligned} \Gamma(\Sigma_c^+ \rightarrow \Lambda_c^+ \gamma) &= 112 \text{ keV}, \\ \Gamma(\Xi_c'^+ \rightarrow \Xi_c^+ \gamma) &= 29 \text{ keV}, \\ \Gamma(\Xi_c'^0 \rightarrow \Xi_c^0 \gamma) &= 0.15 \text{ keV}, \end{aligned} \quad (6.11)$$

Table 14 Electromagnetic decay rates (in units of keV) of s -wave charmed baryons. Among the four different results listed in Refs. [165] and [173], we quote those denoted by $\Gamma_\gamma^{(0)}$ and “Present (ecqm)”, respectively.

| Decay | HHChPT [160, 164] | HHChPT [163] | Dey <i>et al.</i> [165] | Ivanov <i>et al.</i> [75] | Tawfiq <i>et al.</i> [166] | Bañuls <i>et al.</i> [168] | Aliev <i>et al.</i> [169] | Wang [170, 171] | Bernotas <i>et al.</i> [172] | Majethiya <i>et al.</i> [173] |
|---|----------------------|-----------------|----------------------------|------------------------------|-------------------------------|-------------------------------|------------------------------|------------------------|---------------------------------|----------------------------------|
| $\Sigma_c^+ \rightarrow \Lambda_c^+ \gamma$ | 91.5 | 164.16 | 120 | 60.7 ± 1.5 | 87 | | | | 46.1 | 60.55 |
| $\Sigma_c^{*+} \rightarrow \Lambda_c^+ \gamma$ | 150.3 | 892.97 | 310 | 151 ± 4 | | | 130 ± 45 | | 126 | 154.48 |
| $\Sigma_c^{*++} \rightarrow \Sigma_c^{*+} \gamma$ | 1.3 | 11.60 | 1.6 | | 3.04 | | 2.65 ± 1.20 | $6.36^{+6.79}_{-3.31}$ | 0.826 | 1.15 |
| $\Sigma_c^{*+} \rightarrow \Sigma_c^+ \gamma$ | 0.002 | 0.85 | 0.001 | 0.14 ± 0.004 | 0.19 | | 0.40 ± 0.16 | $0.40^{+0.43}_{-0.21}$ | 0.004 | $< 10^{-4}$ |
| $\Sigma_c^{*0} \rightarrow \Sigma_c^0 \gamma$ | 1.2 | 2.92 | 1.2 | | 0.76 | | 0.08 ± 0.03 | $1.58^{+1.68}_{-0.82}$ | 1.08 | 1.12 |
| $\Xi_c'^+ \rightarrow \Xi_c^+ \gamma$ | 19.7 | 54.31 | 14 | 12.7 ± 1.5 | | | | | 10.2 | |
| $\Xi_c'^{*+} \rightarrow \Xi_c^+ \gamma$ | 63.5 | 502.11 | 71 | 54 ± 3 | | | 52 ± 25 | | 44.3 | 63.32 |
| $\Xi_c'^{*+} \rightarrow \Xi_c'^+ \gamma$ | 0.06 | 1.10 | 0.10 | | | | | $0.96^{+1.47}_{-0.67}$ | 0.011 | |
| $\Xi_c'^0 \rightarrow \Xi_c^0 \gamma$ | 0.4 | 0.02 | 0.33 | 0.17 ± 0.02 | | 1.2 ± 0.7 | | | 0.0015 | |
| $\Xi_c'^{*0} \rightarrow \Xi_c^0 \gamma$ | 1.1 | 0.36 | 1.7 | 0.68 ± 0.04 | | 5.1 ± 2.7 | 0.66 ± 0.32 | | 0.908 | 0.30 |
| $\Xi_c'^0 \rightarrow \Xi_c'^0 \gamma$ | 1.0 | 3.83 | 1.6 | | | | | $1.26^{+0.80}_{-0.46}$ | 1.03 | |
| $\Omega_c'^0 \rightarrow \Omega_c^0 \gamma$ | 0.9 | 4.82 | 0.71 | | | | | $1.16^{+1.12}_{-0.54}$ | 1.07 | 2.02 |

Table 15 Electromagnetic decay rates (in units of keV) of p -wave charmed baryons.

| Decay | Ivanov <i>et al.</i> [75] | Tawfiq <i>et al.</i> [166] | Aziza Baccouche <i>et al.</i> [176] | Zhu [177] | Chow [178] | Gamer mann <i>et al.</i> [179] |
|--|------------------------------|-------------------------------|--|--------------|---------------|-----------------------------------|
| $1/2^- \rightarrow 1/2^+(3/2^+) \gamma$ | | | | | | |
| $\Lambda_c(2595)^+ \rightarrow \Lambda_c^+ \gamma$ | 115 ± 1 | | 25 | 36 | 16 | 274 ± 52 |
| $\Lambda_c(2595)^+ \rightarrow \Sigma_c^+ \gamma$ | 77 ± 1 | 71 | | 11 | | 2.1 ± 0.4 |
| $\Lambda_c(2595)^+ \rightarrow \Sigma_c^{*+} \gamma$ | 6 ± 0.1 | 11 | | 1 | | |
| $3/2^- \rightarrow 1/2^+(3/2^+) \gamma$ | | | | | | |
| $\Lambda_c(2625)^+ \rightarrow \Lambda_c^+ \gamma$ | 151 ± 2 | | | 48 | 21 | |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^+ \gamma$ | 35 ± 0.5 | 130 | | 5 | | |
| $\Lambda_c(2625)^+ \rightarrow \Sigma_c^{*+} \gamma$ | 46 ± 0.6 | 32 | | 6 | | |
| $\Xi_c(2815)^+ \rightarrow \Xi_c^+ \gamma$ | 190 ± 5 | | | | | |
| $\Xi_c(2815)^0 \rightarrow \Xi_c^0 \gamma$ | 497 ± 14 | | | | | |

which should be compared with the corresponding quark-model results: 92, 20, and 0.4 keV, respectively (Table 14).

The electromagnetic decays $\Xi_c'^{*0} \rightarrow \Xi_c^0 \gamma$ and $\Xi_c'^0 \rightarrow \Xi_c^0 \gamma$ are of special interest. It was advocated in Ref. [175] that a measurement of their branching fractions will allow us to determine one of the coupling constants in HHChPT, namely, g_1 . They are forbidden at tree level in the SU(3) limit [see Eq. (6.10)]. In heavy baryon chiral perturbation theory, this radiative decay is induced via chiral loops where SU(3) symmetry is broken by the light current quark masses. By identifying the chiral loop contribution to $\Xi_c'^{*0} \rightarrow \Xi_c^0 \gamma$ with the quark model prediction given in Eq. (6.10), it was found in Ref. [164] that one of the two possible solutions agrees with the quark model expectation for g_1 .

For the electromagnetic decays of p -wave charmed baryons, the search for $\Lambda_c(2593)^+ \rightarrow \Lambda_c^+ \gamma$ and $\Lambda_c(2625)^+ \rightarrow \Lambda_c^+ \gamma$ has not yet succeeded. The interested reader is referred to Refs. [14, 75, 166, 175–180] for more

details on the theoretical side. Some predictions are collected in Table 15, and they are more diversified than the s -wave case. For electromagnetic decays of doubly charmed baryons, see, e.g., Refs. [172, 181].

The electromagnetic decays considered to date do not critically test the heavy quark symmetry or chiral symmetry. The results follow simply from the quark model. There are examples in which both the heavy quark symmetry and chiral symmetry enter in a crucial way. These are the radiative decays of heavy baryons involving an emitted pion. Some examples that are kinematically allowed are

$$\Sigma_c \rightarrow \Lambda_c \pi \gamma, \quad \Sigma_c^* \rightarrow \Lambda_c \pi \gamma, \quad \Sigma_c^* \rightarrow \Sigma_c \pi \gamma, \quad \Xi_c^* \rightarrow \Xi_c \pi \gamma. \quad (6.12)$$

The contact interaction dictated by the Lagrangian $\mathcal{L}_B^{(1)}$ can be nicely tested by the decay $\Sigma_c^0 \rightarrow \Lambda_c^+ \pi^- \gamma$, whereas a test of the chiral structure of $\mathcal{L}_B^{(2)}$ is provided by the process $\Sigma_c^+ \rightarrow \Lambda_c^+ \pi^0 \gamma$; see Ref. [160] for the analysis.

6.2 Weak radiative decays

At the quark level, three different types of processes can contribute to the weak radiative decays of heavy hadrons, namely, single-, two-, and three-quark transitions [183]. The single-quark transition mechanism arises from the so-called electromagnetic penguin diagram. Unfortunately, the penguin process $c \rightarrow u\gamma$ is highly suppressed, so it plays no role in radiative decays of charmed hadrons. There are two contributions from the two-quark transitions: one from the W -exchange diagram accompanied by photon emission from the external quark, and the other from the same W -exchange diagram but with a photon radiated from the W boson. The latter is typically suppressed by a factor of $m_q k/M_W^2$ (where k is the photon energy) compared to the former bremsstrahlung process [182]. For charmed baryons, the Cabibbo-allowed decay modes via $c\bar{u} \rightarrow s\bar{d}\gamma$ (Fig. 7) or $cd \rightarrow us\gamma$ are

$$\Lambda_c^+ \rightarrow \Sigma^+\gamma, \quad \Xi_c^0 \rightarrow \Xi^0\gamma. \quad (6.13)$$

Finally, the three-quark transition involving W -exchange between two quarks and photon emission by the third quark is quite strongly suppressed because of the very small probability of finding three quarks that are adequately kinematically matched with the baryons [183, 184].

The general amplitude of weak radiative decays of baryons is given by

$$A(\mathcal{B}_i \rightarrow \mathcal{B}_f\gamma) = i\bar{u}_f(a + b\gamma_5)\sigma_{\mu\nu}\epsilon^\mu k^\nu u_i, \quad (6.14)$$

where a and b are the parity-conserving and -violating amplitudes, respectively. The corresponding decay rate is

$$\Gamma(\mathcal{B}_i \rightarrow \mathcal{B}_f\gamma) = \frac{1}{8\pi} \left(\frac{m_i^2 - m_f^2}{m_i}\right)^3 (|a|^2 + |b|^2). \quad (6.15)$$

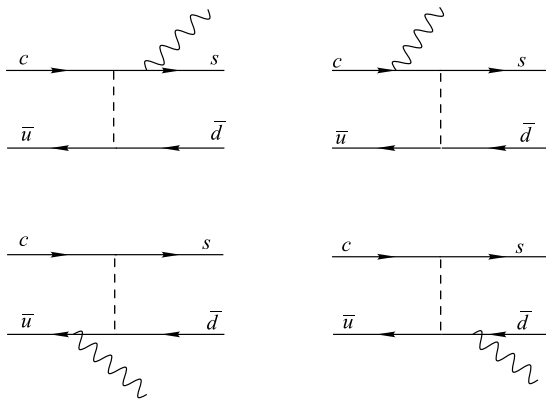


Fig. 7 W -exchange diagrams contributing to the quark-quark bremsstrahlung process $c + \bar{u} \rightarrow s + \bar{d} + \gamma$. The W -annihilation type diagrams are not shown here.

Nonpenguin weak radiative decays of charmed baryons, such as those in Eq. (6.13), are characterized by emission of a hard photon and the presence of a highly virtual intermediate quark between the electromagnetic and weak vertices. It was shown in Ref. [185] that these features should make it possible to analyze these processes by perturbative QCD; that is, these processes can be described by an effective local and gauge invariant Lagrangian:

$$\mathcal{H}_{\text{eff}}(c\bar{u} \rightarrow s\bar{d}\gamma) = \frac{G_F}{2\sqrt{2}} V_{cs} V_{ud}^* (c_+ O_+^F + c_- O_-^F), \quad (6.16)$$

with

$$O_{\pm}^F(c\bar{u} \rightarrow s\bar{d}\gamma) = \frac{e}{m_i^2 - m_f^2} \left\{ \left(e_s \frac{m_f}{m_s} + e_u \frac{m_i}{m_u} \right) \times \left(\tilde{F}_{\mu\nu} + iF_{\mu\nu} \right) O_{\pm}^{\mu\nu} - \left(e_d \frac{m_f}{m_d} + e_c \frac{m_i}{m_c} \right) \times \left(\tilde{F}_{\mu\nu} - iF_{\mu\nu} \right) O_{\mp}^{\mu\nu} \right\}, \quad (6.17)$$

where $m_i = m_c + m_u$, $m_f = m_s + m_d$, $\tilde{F}_{\mu\nu} \equiv \frac{1}{2}\epsilon_{\mu\nu\alpha\beta}F^{\alpha\beta}$ and

$$O_{\pm}^{\mu\nu} = \bar{s}\gamma^\mu(1 - \gamma_5)c\bar{u}\gamma^\nu(1 - \gamma_5)d \pm \bar{s}\gamma^\mu(1 - \gamma_5)d\bar{u}\gamma^\nu(1 - \gamma_5)c. \quad (6.18)$$

For radiative decays of charmed baryons, one needs to evaluate the matrix element $\langle \mathcal{B}_f | O_{\pm}^{\mu\nu} | \mathcal{B}_i \rangle$. Because the quark-model wave functions most resemble the hadronic states in the frame where both baryons are static, the static MIT bag model was adopted in Ref. [185] for the calculation. The predictions are³⁾

$$\begin{aligned} \mathcal{B}(\Lambda_c^+ \rightarrow \Sigma^+\gamma) &= 4.9 \times 10^{-5}, \\ \alpha(\Lambda_c^+ \rightarrow \Sigma^+\gamma) &= -0.86, \\ \mathcal{B}(\Xi_c^0 \rightarrow \Xi^0\gamma) &= 3.5 \times 10^{-5}, \\ \alpha(\Xi_c^0 \rightarrow \Xi^0\gamma) &= -0.86. \end{aligned} \quad (6.19)$$

A different analysis of the same decays in Ref. [186] yielded the following results:

$$\begin{aligned} \mathcal{B}(\Lambda_c^+ \rightarrow \Sigma^+\gamma) &= 2.8 \times 10^{-4}, \\ \alpha(\Lambda_c^+ \rightarrow \Sigma^+\gamma) &= 0.02, \\ \mathcal{B}(\Xi_c^0 \rightarrow \Xi^0\gamma) &= 1.5 \times 10^{-4}, \\ \alpha(\Xi_c^0 \rightarrow \Xi^0\gamma) &= -0.01. \end{aligned} \quad (6.20)$$

Evidently, these predictions (especially the decay asymmetry) differ greatly from those obtained in Ref. [185].

Finally, it is worth remarking that, analogous to the heavy-flavor-conserving nonleptonic weak decays dis-

³⁾ The branching fraction of $\Xi_c^0 \rightarrow \Xi^0\gamma$ has been updated using the current lifetime of Ξ_c^0 .

cussed in Section 6.3, there is a special class of weak radiative decays in which heavy flavor is conserved, for example, $\Xi_c \rightarrow \Lambda_c \gamma$ and $\Omega_c \rightarrow \Xi_c \gamma$. In these decays, weak radiative transitions arise from the light quark sector of the heavy baryon, whereas the heavy quark behaves as a spectator. However, the dynamics of these radiative decays is more complicated than that of their counterpart in nonleptonic weak decays, e.g., $\Xi_c \rightarrow \Lambda_c \pi$. In any event, it merits an investigation.

7 Conclusions

In this report, we began with a brief overview of the spectroscopy of charmed baryons and discussed their possible structure and spin-parity assignments in the quark model. For p -wave baryons, 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 $\Xi_c(3080)$, although the first and last assignments may encounter some potential problems.

It should be stressed that mass analysis alone is usually not sufficient to pin down the spin-parity quantum numbers of higher excited charmed baryon states, a study of their strong decays is necessary. For example, $\Sigma_{c0}(\frac{1}{2}^-)$, $\Sigma_{c1}(\frac{1}{2}^-, \frac{3}{2}^-)$, and $\Sigma_{c2}(\frac{3}{2}^-, \frac{5}{2}^-)$ for $\Sigma_c(2800)$ all have similar masses. Analysis of strong decays allows us to exclude the first two possibilities. It should be stressed that Regge phenomenology and the mass relations for antitriplet and sextet multiplets also provide very useful guidance for the spin-parity quantum numbers.

On the basis of various theoretical tools such as lattice QCD and the QCD sum rule method, there are many theoretical studies of charmed baryon spectroscopy, especially for doubly and triply charmed baryons. However, progress on the hadronic decays, radiative decays, and lifetimes has been very slow. Experimentally, nearly all the branching fractions of the Λ_c^+ are measured relative to the $pK^-\pi^+$ mode. The recent measurements $\mathcal{B}(\Lambda_c^+ \rightarrow pK^-\pi^+) = (6.84 \pm 0.24^{+0.21}_{-0.27})\%$ by Belle and $(5.84 \pm 0.27 \pm 0.23)\%$ by BESIII are very encouraging. Moreover, BESIII recently measured 2-body, 3-body and 4-body nonleptonic decay modes of Λ_c^+ with significantly improved precision. It is conceivable that many new data emerging from LHCb, BESIII, and B factories in the near future and from the experiments at J-PARC and PANDA in the future can be used to test the underlying mechanism of hadronic weak decays.

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