Hadronic weak decays of the charmed baryon $\Omega_c$

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Abstract: Two-body hadronic weak decays of the charmed baryon $\Omega_c$, including Cabibbo-favored (CF), singly Cabibbo-suppressed (SCS) and doubly Cabibbo-suppressed (DCS) modes, are studied systematically in this work. To estimate nonfactorizable contributions, we work in the pole model for the $P$-wave amplitudes and current algebra for the $S$-wave amplitudes. Among all the channels decaying into a baryon octet and a pseudoscalar meson, $\Omega_c \to \Xi^0 K^0$ is the only allowed CF mode. The predicted branching fraction of order 3.8% and large and positive decay asymmetry of order 0.50 indicate that a measurement of this mode in the near future is promising. Proceeding through purely nonfactorizable contributions, the SCS mode $\Omega_c \to \Lambda^0 K^0$ and DCS mode $\Omega_c \to \Lambda^0 \eta$ are predicted to have branching fractions as large as 0.8% and 0.4%, respectively. The two DCS modes $\Omega_c \to \Sigma^0 \eta$ and $\Omega_c \to \Lambda^0 \pi^0$ are suggested to serve as new physics searching channels for their vanishing SM background.

Keywords: Charmed baryons, weak decays, nonfactorizable contribution, pole model
1 Introduction

The Ω_c comprises the combination of a charm quark and two strange quarks. Among the sextet charmed baryons, Ω_c is the only one that decays weakly. It is also the heaviest one with a mass [1]

\[ m = (2695.2 \pm 1.7) \text{ MeV}, \]  

(1.1)
averaged by Belle[2], CLEO[3] and E687[4] experiments. In 2018, LHCb reported a measurement of Ω_c lifetime [6],

\[ \tau = (2.68 \pm 0.26) \times 10^{-13} \text{ s}, \]  

(1.2)
which is nearly four times larger than previous world-average value [7–9], and is consistent with a recent theoretical prediction[5]. The information on Ω_c decays, however, is less than other weakly decaying charmed baryons (Λ_c^+, Ξ_c^0, Ξ_c^+) due to its low production. So far no absolute branching fraction of Ω_c has been measured. Recently Belle reported a ratio between the modes decaying into Ξ_c^0K^0 and Ω_c^-π^+

\[ \frac{B(\Omega_c^0 \rightarrow \Xi_c^0K^0)}{B(\Omega_c^0 \rightarrow \Omega_c^-\pi^+)} = 1.64 \pm 0.26 \pm 0.12, \]  

(1.3)
using 980 fb$^{-1}$ of $e^+e^-$ annihilation data \cite{10}, while the latter one is taken as a benchmarking channel. As for the channels containing $\Omega^-$ in final states, semileptonic decay $\Omega^- \rightarrow \Omega^- e^+\nu_e$ was firstly observed by CLEO in 2012\cite{11}. More interesting results are anticipated when data is accumulated both in Belle-II and LHCb.

The theoretical studies of hadronic weak decays of singly charmed baryons, including $\Lambda^+_c$, $\Xi^+_c$, and $\Omega^0_c$, had a fast development in the 1990s. It has been widely accepted that nonfactorizable effects play an important role in the hadronic decays. Various methodologies were developed to describe the nonfactorizable effects in charmed baryon decays, including the covariant confined quark model \cite{22, 23}, the pole model \cite{14, 15, 17, 18} and current algebra \cite{16, 18}. Among these works, only few of them \cite{18, 22, 23, 40} were $\Omega^0_c$ related. Recently, due to the measurement of the absolute branching fraction of the decay $\Lambda^+_c \rightarrow pK^-\pi^+$ by both Belle \cite{20} and BESIII \cite{21} and hereafter more experimental results, charmed baryon decays have been studied extensively in theory. However, we realize that the study of $\Omega^0_c$ is with less concerned compared with the works on anti-triplet singly charmed baryons. Hence an investigation on weak decays of $\Omega^0_c$, combining current data, is timely and necessary. In this work, we will continue working in the pole model, following our previous treatment of $\Lambda_c$, $\Xi^+_c$, and $B_c$ baryons \cite{25–27}.

In the pole model, nonfactorizable $S$- and $P$-wave amplitudes for $1/2^+ \rightarrow 1/2^+ + 0^-$ decays are dominated by $1/2^-$ low-lying baryon resonances and $1/2^+$ ground-state baryon poles, respectively. However, the estimation of pole amplitudes is a difficult and nontrivial task since it involves weak baryon matrix elements and strong coupling constants of $1/2^+$ and $1/2^-$ baryon states. As a consequence, the evaluation of pole diagrams is far more uncertain than the factorizable terms. This is the case in particular for $S$-wave terms as they require the information of the troublesome negative-parity baryon resonances which are not well understood in the quark model.

It is well known that the $S$-wave amplitudes in pole model is reduced to current algebra in the soft pseudoscalar-meson limit. In this soft-meson limit, the intermediate excited $1/2^-$ states can be summed up and reduced to a commutator term in the $S$-wave amplitude. Using the relation $[Q^a_\ell, H^\text{PV}_{\text{eff}}] = -[Q^a_\ell, H^\text{PC}_{\text{eff}}]$, the parity-violating amplitude is simplified to a simple commutator term expressed in terms of parity-conserving matrix elements. Therefore, the great advantage of current algebra is that the evaluation of the parity-violating $S$-wave amplitude does not require the information of the negative-parity $1/2^-$ poles. Although the pseudoscalar meson produced in $B_c \rightarrow B + P$ decays is in general not truly soft, current algebra seems to work empirically well for $\Lambda^+_c \rightarrow B + P$ decays \cite{25, 26}. Moreover, the predicted negative decay asymmetries by current algebra for both $\Lambda^+_c \rightarrow \Sigma^+\pi^0$ and $\Sigma^0\pi^+$ agree in sign with the recent BESIII measurements \cite{24} (see \cite{25, 26} for details). In contrast, the pole model or the covariant quark model and its variant always leads to a positive decay asymmetry for aforementioned two modes. Therefore, in this work we shall follow \cite{25–27} to work out the nonfactorizable $S$-wave amplitudes in $\Omega_c$ decays using current algebra and the nonfactorizable contributions, including $W$-exchange contribution, to $P$-wave ones using the pole model.

This paper is organized as follows. In Sec. II we set up the framework for the analysis of hadronic weak decays of the singly charmed baryon $\Omega_c$, including the topological dia-
grams and the formalism for describing factorizable and nonfactorizable terms. We present the explicit expressions of nonfactorizable amplitudes for both S- and P-waves. Baryon matrix elements and axial-vector form factors calculated in the MIT bag model are also summarized. Numerical results and discussions are presented in Sec. III. A conclusion will be given in Sec. 4.

2 Theoretical framework

2.1 Topological diagrams

The general formulation of the topological-diagram scheme for the nonleptonic weak decays of baryons was proposed by Chau, Tseng and Cheng [28] more than two decades ago, which was then applied to all the decays of the antitriplet and sextet charmed baryons. For the weak decays $\Omega_c \rightarrow \mathcal{B} + P$ ($\mathcal{B}$ is baryon octet) of interest in this work, the relevant topological diagrams are the external $W$-emission $T$, the internal $W$-emission $C$, the inner $W$-emission $C'$, and the $W$-exchange diagrams $E_1$ as well as $E_2$ as depicted in Fig. 1. Among them, $T$ and $C$ are factorizable, while $C'$ and $W$-exchange give nonfactorizable contributions. The relevant topological diagrams for all decay modes of $\Omega_c$, including Cabibbo-favored (CF), singly Cabibbo-suppressed (SCS) and doubly Cabibbo-suppressed (DCS) modes, are shown in Table 1.

We notice from Table 1 that (i) among all the CF, SCS and DCS decays of $\Omega_c$, there is no purely factorizable mode, (ii) the modes containing $\Xi^0$ or $\Xi^-$ in final states receive both factorizable and nonfactorizable contributions while the modes containing $\Sigma$ and $\Lambda^0$ in final states have purely nonfactorizable contribution (iii) the $W$-exchange contribution is absent from the CF process but manifests in all the SCS and DCS decays, and (iv) the two nucleons $n$ and $p$, as parts of baryon octet, are absent from all the $\Omega_c$ decays.

Figure 1. Topological diagrams contributing to $\Omega_c \rightarrow \mathcal{B} + P$ decays: external $W$-emission $T$, internal $W$-emission $C$, inner $W$-emission $C'$, $W$-exchange diagrams $E_1$ and $E_2$. 
Table 1. Topological diagrams contributing to two-body weak decays $\Omega_c \to BP$, where $B$ is a baryon octet and $P$ is a pseudoscalar meson.

| CF Contributions | SCS Contributions | DSC Contributions |
|------------------|-------------------|-------------------|
| $\Omega_c^0 \to \Xi^K_0$ | $C, C'$ | $T, E_1$ |
| $\Omega_c^0 \to \Sigma^0 K^-$ | $E_2$ | $\Omega_c^0 \to \Sigma^0 \eta$ | $C', E_1, E_2$ |
| $\Omega_c^0 \to \Sigma^0 K^0$ | $C', E_2$ | $\Omega_c^0 \to \Lambda^0 \eta$ | $C', E_1, E_2$ |
| $\Omega_c^0 \to \Lambda^K_0$ | $C', E_2$ | $\Omega_c^0 \to \Sigma^0 \pi^0$ | $E_1, E_2$ |
| $\Omega_c^0 \to \Xi^K_0$ | $C, E_1$ | $\Omega_c^0 \to \Lambda^0 \pi^0$ | $E_1, E_2$ |
| $\Omega_c^0 \to \Xi^0 \pi^0$ | $C, E_1$ | $\Omega_c^0 \to \Xi^- K^+$ | $T, E_1$ |
| $\Omega_c^0 \to \Xi^- K^+$ | $T, E_1$ | $\Omega_c^0 \to \Xi^- \pi^-$ | $E_2$ |
| $\Omega_c^0 \to \Lambda^0 \pi^+$ | $E_1$ | $\Omega_c^0 \to \Sigma^0 \pi^+$ | $E_1$ |

2.2 Kinematics

The amplitude for two-body weak decay $B_i \to B_f P$ generally can be parametrized as

$$M(B_i \to B_f P) = i\bar{u}_f (A - B \gamma_5) u_i,$$  

where $B_i (B_f)$ is the initial (final) baryon and $P$ is a pseudoscalar meson. The decay width and up-down decay asymmetry are given by

$$\Gamma = \frac{p_c}{8\pi} \left[ \frac{(m_i + m_f)^2 - m_i^2}{m_i^2} |A|^2 + \frac{(m_i - m_f)^2 - m_i^2}{m_i^2} |B|^2 \right],$$

$$\alpha = \frac{2\kappa \text{Re}(A^*B)}{|A|^2 + \kappa^2 |B|^2},$$

where $p_c$ is the three-momentum in the rest frame of the mother particle and $\kappa = p_c/(E_f + m_f) = \sqrt{(E_f - m_f)/(E_f + m_f)}$. The $S$- and $P$- wave amplitudes of the two-body decay generally receive both factorizable and non-factorizable contributions

$$A = A^{\text{fac}} + A^{\text{nf}}, \quad B = B^{\text{fac}} + B^{\text{nf}}.$$  

Details of factorizable and non-factorizable amplitudes will be given in the following context.

2.3 Factorizable amplitudes

Here the state-of-the-art effective Hamiltonian approach is adopted to describe factorizable contributions in the charmed baryon decays. Physical observables can be expressed as products of short-distance and long-distance parameters. The short-distance physics is characterized by Wilson coefficients, which is obtained from integrating out high energy part of theory, and long-distance physics resorts to hadronic matrix elements.
2.3.1 General expression of factorizable amplitudes

The effective Hamiltonian for CF process in $\Omega_c$ decay is

$$\mathcal{H}_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{cs} V_{ud}^* (c_1 O_1 + c_2 O_2) + \text{h.c.}, \tag{2.4}$$

where the four-quark operators are given by

$$O_1 = (\bar{s}c)(\bar{u}d), \quad O_2 = (\bar{u}c)(\bar{s}d), \tag{2.5}$$

with $(\bar{q}_1 q_2) \equiv \bar{q}_1 \gamma_\mu (1 - \gamma_5) q_2$. The Wilson coefficients to the leading order are given as $c_1 = 1.346$ and $c_2 = -0.636$ at $\mu = 1.25 \text{ GeV}$ and $\Lambda_{\text{MS}}^{(4)} = 325 \text{ MeV} \ [30]$. Under naive factorization the amplitude can be written down as

$$M = \langle \mathcal{K}^0 \bar{\Xi}^0 | \mathcal{H}_{\text{eff}} | \Omega_c \rangle = \frac{G_F}{\sqrt{2}} V_{cs} V_{ud}^* a_2 \langle \mathcal{K}^0 | (\bar{u}c)(\bar{s}d) | 0 \rangle \langle \bar{\Xi}^0 | (\bar{u}c) | \mathcal{B}_c \rangle, \tag{2.6}$$

where $a_2 = c_2 + \frac{\mu}{M}$. We also give the definition of $a_1$ as $a_1 = c_1 + \frac{\mu}{M}$ for the purpose of describing SCS and DCS decays conveniently. In terms of the decay constants

$$\langle K(q) | \bar{s}c \gamma_\mu (1 - \gamma_5) d | 0 \rangle = i f_K q_\mu, \quad \langle \pi(q) | \bar{u}c \gamma_\mu (1 - \gamma_5) d | 0 \rangle = i f_\pi q_\mu, \tag{2.7}$$

and the form factors defined by

$$\langle \Xi^0(p_2) | \bar{s}c \gamma_\mu (1 - \gamma_5) u | \Omega_c(p_1) \rangle = \bar{u}_2 \left[ f_1(q^2) \gamma_\mu - f_2(q^2) i \sigma_{\mu\nu} q^\nu_M + f_3(q^2) q_\mu_M \right. \tag{2.8}$$

$$- \left. \left( g_1(q^2) \gamma_\mu - g_2(q^2) i \sigma_{\mu\nu} q^\nu_M + g_3(q^2) q_\mu_M \right) \gamma_5 \right] u_1,$$

with the momentum transfer $q = p_1 - p_2$, we obtain the amplitude

$$M(\Omega_c \to \Xi^0(\mathcal{K}^0)) = i \frac{G_F}{\sqrt{2}} a_2 V_{ud}^* V_{cs} f_K \bar{u}_2(p_2) \left[ (m_1 - m_2) f_1(q^2) + (m_1 + m_2) g_1(q^2) \gamma_5 \right] u_1(p_1), \tag{2.9}$$

where contributions from the form factors $f_3$ and $g_3$ can be neglected, with similar reason given in footnote 1 of [27]. The factorizable $S$- and $P$-wave amplitudes finally read

$$A^{\text{fac}}_{\text{CF}} = \frac{G_F}{\sqrt{2}} a_2 V_{ud}^* V_{cs} f_K (m_{\Omega_c} - m_{\Xi^0}) f_1(q^2),$$

$$B^{\text{fac}}_{\text{CF}} = -\frac{G_F}{\sqrt{2}} a_2 V_{ud}^* V_{cs} f_K (m_{\Omega_c} + m_{\Xi^0}) g_1(q^2). \tag{2.10}$$

Likewise, the $S$- and $P$-wave amplitudes for SCS processes and DCS processes are given, respectively, by

$$A^{\text{fac}}_{\text{SCS}} = \frac{G_F}{\sqrt{2}} a_{1,2} V_{ud}^* V_{cq} f_P (m_{\Omega_c} - m_B) f_1(q^2),$$

$$B^{\text{fac}}_{\text{SCS}} = -\frac{G_F}{\sqrt{2}} a_{1,2} V_{ud}^* V_{cq} f_P (m_{\Omega_c} + m_B) g_1(q^2), \tag{2.11}$$
Table 2. The calculated form factors in MIT bag model at maximum four-momentum transfer squared $q^2 = q^2_{\text{max}} = (m_i - m_f)^2$ and $q^2 = m_P^2$.

| modes          | $(c\bar{q})$ | $f_1(q^2_{\text{max}})$ | $f_1(m_P^2)/f_1(q^2_{\text{max}})$ | $g_1(q^2_{\text{max}})$ | $g_1(m_P^2)/g_1(q^2_{\text{max}})$ | $g_1(m_P^2)$ |
|----------------|--------------|--------------------------|-------------------------------------|--------------------------|-----------------------------------|--------------|
| $\Omega_c^0 \rightarrow \Xi^0 \overline{K}^0$ ($c\bar{s}$) | Y1 | 0.37 | 0.32 | $-\frac{1}{2}Y_2$ | 0.54 | $-0.14$ |
| $\Omega_c^0 \rightarrow \Xi^- \pi^+$ ($c\bar{d}$) | Y1 | 0.29 | 0.25 | $-\frac{1}{2}Y_2$ | 0.46 | $-0.12$ |
| $\Omega_c^0 \rightarrow \Xi^0 \pi^0$ ($c\bar{d}$) | Y1 | 0.28 | 0.25 | $-\frac{1}{2}Y_2$ | 0.46 | $-0.12$ |
| $\Omega_c^0 \rightarrow \Xi^0 K^0$ ($c\bar{d}$) | Y1 | 0.32 | 0.28 | $-\frac{1}{2}Y_2$ | 0.50 | $-0.13$ |
| $\Omega_c^0 \rightarrow \Xi^- K^+$ ($c\bar{d}$) | Y1 | 0.32 | 0.28 | $-\frac{1}{2}Y_2$ | 0.50 | $-0.13$ |

$$A^{\text{Fac}}|_{\text{DCS}} = \frac{G_F}{\sqrt{2}} a_{1,2} V^*_{us} V_{cd} f_P (m_{\Omega_c} - m_B) f_1(q^2),$$

$$B^{\text{Fac}}|_{\text{DCS}} = -\frac{G_F}{\sqrt{2}} a_{1,2} V^*_{us} V_{cd} f_P (m_{\Omega_c} + m_B) g_1(q^2),$$

(2.12)

where the choice of $a_i$ depends on the meson in final states. In the amplitude of SCS processes Eq. (2.11) the flavor of the down-type quark $q, d$ or $s$, depends on the process. Notice that in the case of $\pi^0$ production in the final state, one should replace $a_2$ by $-a_2/\sqrt{2}$ in the factorizable amplitude, where the extra factor of $-1/\sqrt{2}$ comes from the wave function of the $\pi^0$, $\pi^0 = (u\bar{u} - d\bar{d})/\sqrt{2}$.

2.3.2 Baryon transition form factors

In this work we shall work out form factor (FF) for $\Omega_c-B$ transition and baryonic matrix elements in the framework of MIT bag model [31]. Since the decay rates and decay asymmetries are sensitive to the relative sign between factorizable and non-factorizable amplitudes, the estimation of FFs in a globally consistent convention is desirable.

Following [17] we can write down the $q^2$ dependence of FF as

$$f_i(q^2) = \frac{f_i(0)}{1 - q^2/m_V^2}, \quad g_i(q^2) = \frac{g_i(0)}{1 - q^2/m_A^2},$$

(2.13)

where $m_V = 2.01 \text{GeV}$, $m_A = 2.42 \text{GeV}$ for the $(c\bar{d})$ quark content, and $m_V = 2.11 \text{GeV}$, $m_A = 2.54 \text{GeV}$ for $(c\bar{s})$ quark content. In the zero recoil limit where $q^2_{\text{max}} = (m_i - m_f)^2$, FFs can be calculated in the MIT bag model [18], giving

$$f_1^{B_1/B_i}(q^2_{\text{max}}) = \langle B^+_i | b_{q_1} b_{q_2} | B^+_i \rangle \int d^3 r \left( u_{q_1}(r) u_{q_2}(r) + v_{q_1}(r) v_{q_2}(r) \right),$$

$$g_1^{B_1/B_i}(q^2_{\text{max}}) = \langle B^+_i | b_{q_1} b_{q_2} \sigma_2 | B^+_i \rangle \int d^3 r \left( u_{q_1}(r) u_{q_2}(r) - \frac{1}{3} v_{q_1}(r) v_{q_2}(r) \right),$$

(2.14)

where $u(r)$ and $v(r)$ are the large and small components, respectively, of the quark wave function in the bag model. FFs at different $q^2$ are related via

$$f_i(q^2) = \frac{(1 - q^2/m_V^2)^2}{(1 - q^2/m_V^2)^2} f_i(q^2_{\text{max}}), \quad g_i(q^2) = \frac{(1 - q^2/m_A^2)^2}{(1 - q^2/m_A^2)^2} g_i(q^2_{\text{max}}).$$

(2.15)
Table 3. Form factors $f_1(q^2)$ and $g_1(q^2)$ at $q^2 = 0$ for $\Omega_c^0 \rightarrow \Xi$ transitions evaluated in the MIT bag model (this work), non-relativistic quark model (NRQM) [38] and heavy quark effective theory (HQET)[29], which were quoted in [32], as well as light-front quark model (LFQM) [41].

|-transition | $f_1(0)$ | $g_1(0)$ |
|------------|---------|---------|
| $\Omega_c^0 \rightarrow \Xi$ | MIT NRQM HQET LFQM | MIT NRQM HQET LFQM |
| $\Omega_c^0$ | 0.34 | 0.23 | 0.34 | 0.653 | -0.15 | -0.14 | -0.10 | -0.182 |

This allows us to obtain the physical FFs at $q^2 = m_P^2$.

It is obvious that the FF at $q^2_{\text{max}}$ is determined only by the baryons in initial and final states. However, its evolution with $q^2$ is governed by both the final-state meson and relevant quark content. Such dependence is reflected in Table 2, in which the quark contents are shown in the second column. In the zero recoil limit, the FFs at $q^2_{\text{max}}$ calculated from Eq. (2.13) are presented in the third and sixth columns. And then in the fourth and seventh columns, the evolution of FFs from $q^2 = q^2_{\text{max}}$ to $q^2 = m_P^2$ are derived according to Eq. (2.15). The auxiliary quantities $Y_{1,2}$ can be obtained from the calculation in MIT bag model, giving

$$
Y_1 = 4\pi \int r^2 dr (u_u u_c + v_v v_c), \quad Y_1^s = 4\pi \int r^2 dr (u_u u_c + v_v v_c),
$$

$$
Y_2 = 4\pi \int r^2 dr (u_u u_c - \frac{1}{3} v_v v_c), \quad Y_2^s = 4\pi \int r^2 dr (u_u u_c - \frac{1}{3} v_v v_c). \quad (2.16)
$$

The model parameters are adopted from [25] and references therein. Numerically, we have $Y_1 = 0.88, Y_1^s = 0.95, Y_2 = 0.77, Y_2^s = 0.86$, which are consistent with the corresponding numbers in [18].

A comparison is made in Table 3 among results of $\Omega_c \rightarrow \Xi$ FF in various approaches at $q^2 = 0$. From Eq. (2.13) and Table 2, FFs at $q^2 = q^2_{\text{max}}$ are identical for $\Omega_c \rightarrow \Xi^0, P$. However, the values would be changed at different energy scale and also depend on different meson states. In this comparison, we ignore such slight differences. Note that all the nonperturbative quantities involving in this work are calculated in MIT bag model, thus our convention of signs should be consistent globally. For this reason we correct the signs in NRQM and HQET cases. It is apparent that NRQM gives small values for both $f_1$ and $g_1$, while the predictions from LFQM are the largest.

2.4 Non-factorizable amplitudes

Nonfactorizable contribution play an essential role in charmed baryon decays. Various methods have been developed to study nonfactorizable contribution, here we continue working in the pole model. The general formula for $S$- and $P$-wave amplitudes in the
pole model is given as [25–27]

\[
A_{\text{Pole}} = - \sum_{B^*_n(1/2^-)} \left[ \frac{g_{B_i B^*_n} f_{n^* i} b_{n^* i}}{m_i - m_{n^*}} + \frac{b_{f n^*} g_{B_i B^*_n} M_{i n^*}}{m_f - m_{n^*}} \right],
\]

\[
B_{\text{Pole}} = \sum_{B_n} \left[ \frac{g_{B_f B_n} M_{a n} a_{n f}}{m_i - m_n} + \frac{a_{f n} g_{B_n B_i} M_{n^* i}}{m_f - m_n} \right],
\]

with the baryonic matrix elements

\[
\langle B_n | H | B_i \rangle = \bar{u}_n (a_{n i} + b_{n i} \gamma_5) u_i, \quad \langle B_i^\ast (1/2^-) | H | B_j \rangle = \bar{u}_i b^*_{i j} u_j.
\]

(2.17)

To estimate the \( S \)-wave amplitudes in the pole model is a difficult and nontrivial task as it involves the matrix elements and strong coupling constants of 1/2\(^-\) baryon resonances which is less known [17]. Nevertheless, provided a soft emitted pseudoscalar meson, the intermediate excited baryons can be summed up, leading to a commutator term

\[
A_{\text{com}} = - \frac{\sqrt{2}}{f_{P_a}} \langle B_f | [Q_a^\ast, H_{\text{PV}}^{\text{eff}}] | B_i \rangle = \frac{\sqrt{2}}{f_{P_a}} \langle B_f | [Q_a, H_{\text{PC}}^{\text{eff}}] | B_i \rangle,
\]

(2.19)

with

\[
Q^a = \int d^3 x \bar{q} \gamma^0 \frac{\lambda^a}{2} q, \quad Q_5^a = \int d^3 x \bar{q} \gamma^5 \gamma^0 \frac{\lambda^a}{2} q.
\]

(2.20)

Likewise, the \( P \)-wave amplitude is reduced in the soft-meson limit to

\[
B_{\text{ca}} = \frac{\sqrt{2}}{f_{P_a}} \sum_{B_n} \left[ g_{B_f B_n}^A \frac{m_f + m_n a_{n f}}{m_i - m_n} + \frac{a_{f n} g_{B_n B_i}^A M_{n^* i}}{m_f - m_n} \right],
\]

(2.21)

where we have applied the generalized Goldberger-Treiman relation

\[
g_{B'B_a} = \frac{\sqrt{2}}{f_{P_a}} (m_B + m_{B'}) g_{BB}^A.
\]

(2.22)

Eqs. (2.19) and (2.21) are the master equations for nonfactorizable amplitudes in the pole model under the soft meson approximation.

2.4.1 S-wave amplitudes

Now we know the nonfactorizable \( S \)-wave amplitude is determined by the commutator terms of conserving charge \( Q^a \) and the parity-conserving part of the Hamiltonian, shown in Eq. (2.19). We further present the \( A_{\text{com}} \) terms for various meson production more
explicitly:

\[ A^{\text{com}}(B_i \to B_f \pi^\pm) = \frac{1}{f_\pi} \langle B_f | [I_\mp, H_{\text{eff}}^{PC}] | B_i \rangle, \quad (2.23) \]

\[ A^{\text{com}}(B_i \to B_f \pi^0) = \frac{\sqrt{2}}{f_\pi} \langle B_f | [I_3, H_{\text{eff}}^{PC}] | B_i \rangle, \quad (2.24) \]

\[ A^{\text{com}}(B_i \to B_f \eta_8) = \sqrt{\frac{3}{2}} \frac{1}{f_{\eta_8}} \langle B_f | [Y, H_{\text{eff}}^{PC}] | B_i \rangle, \quad (2.25) \]

\[ A^{\text{com}}(B_i \to B_f K^\pm) = \frac{1}{f_K} \langle B_f | [V_\mp, H_{\text{eff}}^{PC}] | B_i \rangle, \quad (2.26) \]

\[ A^{\text{com}}(B_i \to B_f \overline{K}^0) = \frac{1}{f_K} \langle B_f | [U_+, H_{\text{eff}}^{PC}] | B_i \rangle, \quad (2.27) \]

\[ A^{\text{com}}(B_i \to B_f K^0) = \frac{1}{f_K} \langle B_f | [U_-, H_{\text{eff}}^{PC}] | B_i \rangle. \quad (2.28) \]

In Eq. (2.23), \( \eta_8 \) is the octet component of the \( \eta \) and \( \eta' \)

\[ \eta = \cos \theta \eta_8 - \sin \theta \eta_0, \quad \eta' = \sin \theta \eta_8 + \cos \theta \eta_0, \quad (2.29) \]

with \( \theta = -15.4^\circ \) [33]. For the decay constant \( f_{\eta_8} \), we shall follow [33] to use \( f_{\eta_8} = f_8 \cos \theta \) with \( f_8 = 1.26 f_\pi \). The convention for hypercharge \( Y \) is taken to be \( Y = B + S - C \) [25].

The remaining task is then to evaluate different groups of commutators. We shall present our results after a straightforward calculation, for the S-wave amplitudes, as follows:

\[ A^{\text{com}}(\Omega_c^0 \to \Xi^0 \overline{K}^0) = -\frac{\sqrt{2}}{f_K} a_{\Xi^0 \overline{K}^0}, \quad (2.30) \]

for Cabibbo-favored (CF) process,

\[ A^{\text{com}}(\Omega_c^0 \to \Xi^- \pi^+) = \frac{1}{f_\pi} a_{\Xi^- \pi^+}, \quad A^{\text{com}}(\Omega_c^0 \to \Xi^+ \pi^-) = \frac{1}{f_K} (a_{\Xi^0 \pi^0} - \sqrt{2} a_{\Sigma^0 \Xi^\prime}^0), \]

\[ A^{\text{com}}(\Omega_c^0 \to \Sigma^0 \overline{K}^0) = \frac{1}{f_K} (-\sqrt{\frac{2}{3}} a_{\Xi^0 \overline{K}^0} - \sqrt{2} a_{\Sigma^0 \Xi^\prime}^0), \quad A^{\text{com}}(\Omega_c^0 \to \Xi^0 \pi^0) = \frac{1}{f_\pi} \sqrt{\frac{2}{3}} a_{\Xi^0 \pi^0}, \]

\[ A^{\text{com}}(\Omega_c^0 \to \Lambda^0 \Xi^0) = \frac{1}{f_K} (\sqrt{\frac{6}{5}} a_{\Xi^0 \overline{K}^0} - \sqrt{2} a_{\Lambda^0 \Xi^0}), \quad (2.31) \]

for singly Cabibbo-suppressed (SCS) processes, and

\[ A^{\text{com}}(\Omega_c^0 \to \Xi^0 K^0) = \frac{1}{f_K} (\sqrt{\frac{\sqrt{2}}{2}} a_{\Xi^0 \Sigma^0} + \sqrt{\frac{\sqrt{2}}{2}} a_{\Lambda^0 \Xi^0}), \quad A^{\text{com}}(\Omega_c^0 \to \Sigma^0 \eta) = \frac{\sqrt{6}}{f_\eta} a_{\Sigma^0 \eta}, \]

\[ A^{\text{com}}(\Omega_c^0 \to \Lambda^0 \eta) = \frac{\sqrt{6}}{f_\eta} a_{\Lambda^0 \eta}, \quad A^{\text{com}}(\Omega_c^0 \to \Sigma^- \pi^+) = \frac{\sqrt{\frac{2}{3}}}{f_\eta} a_{\Sigma^0 \eta}, \]

\[ A^{\text{com}}(\Omega_c^0 \to \Xi^- K^+) = -\frac{1}{f_K} (\sqrt{\frac{2}{3}} a_{\Sigma^0 \Xi^0} + \sqrt{\frac{2}{3}} a_{\Lambda^0 \Xi^0}), \quad A^{\text{com}}(\Omega_c^0 \to \Sigma^+ \pi^-) = -\sqrt{\frac{2}{3}} a_{\Sigma^0 \Xi^0}, \]

\[ A^{\text{com}}(\Omega_c^0 \to \Sigma^0 \pi^0) = 0, \quad A^{\text{com}}(\Omega_c^0 \to \Lambda^0 \pi^0) = 0, \quad (2.32) \]

for the doubly Cabibbo-suppressed (DCS) processes, where the baryonic matrix element \( (B'|H_{\text{eff}}^{PC}|B) \) is denoted by \( a_{B'B} \). We can find that a straightforward calculation of the last two terms directly leads to vanished results. This can be easily understood as \( I_3(\Lambda^0) = I_3(\Sigma^0) = 0 \).
2.4.2 P-wave amplitudes

Now we turn to the nonfactorizable P-wave amplitudes given by Eq. (2.21). By substituting explicit hadron states, we have

\[
B^{ca}(\Omega_c^0 \to \Xi^0 K^0) = \frac{1}{f_K} \left( a_{\Xi^0 z e} \frac{m_{\Xi^0} + m_{\Xi^0}}{m_{\Xi^0} - m_{\Xi^0}} g_{\Xi^0 \Omega_c^0} + a_{\Xi^0 z e} \frac{m_{\Xi^0} + m_{\Xi^0}}{m_{\Xi^0} - m_{\Xi^0}} g_{\Xi^0 \Omega_c^0} \right),
\]

for CF decays,

\[
B^{ca}(\Omega_c^0 \to \Xi^- \pi^+) = \frac{1}{f_{\pi}} \left( g_{A(\pi^-)} \frac{m_{\Xi^-} + m_{\Xi^-}}{m_{\Xi^-} - m_{\Xi^-}} a_{\Xi^- \Omega_c^0} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Sigma^+ K^-) = \frac{1}{f_K} \left( g_{A(K^-)} \frac{m_{\Sigma^+} + m_{\Sigma^+}}{m_{\Sigma^+} - m_{\Sigma^+}} a_{\Sigma^+ \Omega_c^0} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Sigma^0(1385)^0) = \frac{1}{f_K} \left( a_{\Sigma^0(1385)^0} \frac{m_{\Sigma^0} + m_{\Sigma^0}}{m_{\Sigma^0} - m_{\Sigma^0}} g_{\Sigma^0(1385)^0} + a_{\Sigma^0(1385)^0} \frac{m_{\Sigma^0} + m_{\Sigma^0}}{m_{\Sigma^0} - m_{\Sigma^0}} g_{\Sigma^0(1385)^0} + g_{\Sigma^0(1385)^0} m_{\Sigma^0} a_{\Sigma^0(1385)^0} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Lambda^0 K^0) = \frac{1}{f_K} \left( a_{\Lambda^0} \frac{m_{\Lambda^0} + m_{\Lambda^0}}{m_{\Lambda^0} - m_{\Lambda^0}} g_{\Lambda^0(\Xi^0)^0} + a_{\Lambda^0(\Xi^0)^0} \frac{m_{\Lambda^0} + m_{\Lambda^0}}{m_{\Lambda^0} - m_{\Lambda^0}} g_{\Lambda^0(\Xi^0)^0} + g_{\Lambda^0(\Xi^0)^0} m_{\Lambda^0} a_{\Lambda^0(\Xi^0)^0} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Xi^- K^0) = \frac{1}{f_K} \left( g_{A(K^-)} \frac{m_{\Xi^-} + m_{\Xi^-}}{m_{\Xi^-} - m_{\Xi^-}} a_{\Xi^- \Omega_c^0} + g_{A(\Xi^-)} \frac{m_{\Xi^-} + m_{\Xi^-}}{m_{\Xi^-} - m_{\Xi^-}} a_{\Xi^- \Omega_c^0} \right),
\]

for SCS processes, and

\[
B^{ca}(\Omega_c^0 \to \Xi^- \pi^+) = \frac{1}{f_{\pi}} \left( g_{A(\pi^-)} \frac{m_{\Xi^-} + m_{\Xi^-}}{m_{\Xi^-} - m_{\Xi^-}} a_{\Xi^- \Omega_c^0} + g_{A(\Xi^-)} \frac{m_{\Xi^-} + m_{\Xi^-}}{m_{\Xi^-} - m_{\Xi^-}} a_{\Xi^- \Omega_c^0} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Sigma^+ K^-) = \frac{1}{f_K} \left( a_{\Sigma^+ K^-} \frac{m_{\Sigma^+} + m_{\Sigma^+}}{m_{\Sigma^+} - m_{\Sigma^+}} g_{\Sigma^+ K^-} + a_{\Sigma^+ K^-} \frac{m_{\Sigma^+} + m_{\Sigma^+}}{m_{\Sigma^+} - m_{\Sigma^+}} g_{\Sigma^+ K^-} + g_{\Sigma^+ K^-} m_{\Sigma^+} a_{\Sigma^+ K^-} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Sigma^0(1385)^0) = \frac{1}{f_K} \left( a_{\Sigma^0(1385)^0} \frac{m_{\Sigma^0} + m_{\Sigma^0}}{m_{\Sigma^0} - m_{\Sigma^0}} g_{\Sigma^0(1385)^0} + a_{\Sigma^0(1385)^0} \frac{m_{\Sigma^0} + m_{\Sigma^0}}{m_{\Sigma^0} - m_{\Sigma^0}} g_{\Sigma^0(1385)^0} + g_{\Sigma^0(1385)^0} m_{\Sigma^0} a_{\Sigma^0(1385)^0} \right),
\]

\[
B^{ca}(\Omega_c^0 \to \Lambda^0 K^0) = \frac{1}{f_K} \left( a_{\Lambda^0} \frac{m_{\Lambda^0} + m_{\Lambda^0}}{m_{\Lambda^0} - m_{\Lambda^0}} g_{\Lambda^0 K^0} + a_{\Lambda^0 K^0} \frac{m_{\Lambda^0} + m_{\Lambda^0}}{m_{\Lambda^0} - m_{\Lambda^0}} g_{\Lambda^0 K^0} + g_{\Lambda^0 K^0} m_{\Lambda^0} a_{\Lambda^0 K^0} \right),
\]

for DCS decay processes. Next we need to derive explicit expressions for \(a_{BG}^c\) and \(g_{BG}^c\).
2.5 Hadronic matrix elements and axial-vector form factors

There are two types of non-perturbative quantities involved in the nonfactorizable amplitudes: hadronic matrix elements and axial-vector form factors. The estimation of the two types of parameters are carried out within the framework of the MIT bag model [31].

2.5.1 Hadronic matrix elements

The baryonic matrix elements \( a_{B'B} \) get involved both in \( S \)-wave and \( P \)-wave amplitudes. Their general expression in terms of the effective Hamiltonian Eq. (2.4) is given by

\[
a_{B'B} = \langle B' | H_{\text{eff}}^{\text{PC}} | B \rangle = \begin{cases} \frac{G_F}{2\sqrt{2}} V_{cs} V_{ud}^* c_- \langle B' | O_- | B \rangle, & \text{CF} \\ \frac{G_F}{2\sqrt{2}} \sum_q V_{cq} V_{uq}^* c_- \langle B' | O^0 \langle B | B \rangle, & \text{SCS} \\ \frac{G_F}{2\sqrt{2}} V_{cd} V_{us}^* c_- \langle B' | O^D \langle B | B \rangle, & \text{DCS} \end{cases}
\]

where \( O_\pm = (\bar{s}c)(\bar{u}d) \pm (\bar{sd})(\bar{uc}), O^0_\pm = (\bar{q}c)(\bar{u}q) \pm (\bar{q}q)(\bar{uc}) \) (with \( q = d, s \)) and \( O^D_\pm = (\bar{d}c)(\bar{u}s) \pm (\bar{d}s)(\bar{uc}) \) and \( c_\pm = c_1 \pm c_2 \). The matrix element of \( O^{(q,D)}_\pm \) vanishes as this operator is symmetric in color indices. We shall calculate relevant baryon matrix elements in MIT bag model. The results for CF processes are

\[
\langle \Xi^0 | O_- | \Xi^0_c \rangle = - \frac{2\sqrt{2}}{3} (X_1 + 9X_2)(4\pi), \quad \langle \Xi^0 | O_- | \Xi^0_c \rangle = - \frac{2\sqrt{6}}{3} (X_1 - 3X_2)(4\pi),
\]

Likewise, the matrix elements for SCS decays are calculated to be

\[
\langle \Sigma^0 | O^d | \Sigma^0_c \rangle = 0, \quad \langle \Sigma^0 | O^s | \Sigma^0_c \rangle = - \frac{4}{3} (X_1^s + 9X_2^s)(4\pi),
\]

\[
\langle \Sigma^+ | O^d | \Sigma^+ c \rangle = 0, \quad \langle \Sigma^+ | O^s | \Sigma^+ c \rangle = \frac{2\sqrt{2}}{3} (X_1^s - 9X_2^s)(4\pi),
\]

\[
\langle \Sigma^0 | O^d | \Sigma^0 c \rangle = \frac{4}{3} X_1^d(4\pi), \quad \langle \Sigma^0 | O^s | \Sigma^0 c \rangle = - \frac{2}{3} (X_1^s - 9X_2^s)(4\pi),
\]

\[
\langle \Lambda^0 | O^d | \Lambda^0 c \rangle = - 4\sqrt{2} X_2^d(4\pi), \quad \langle \Lambda^0 | O^s | \Lambda^0 c \rangle = - \frac{2\sqrt{3}}{3} (X_1^s + 3X_2^s)(4\pi),
\]

\[
\langle \Sigma^0 | O^d | \Xi^0 c \rangle = - 4\sqrt{3} X_2^d(4\pi), \quad \langle \Sigma^0 | O^s | \Xi^0 c \rangle = - \frac{2\sqrt{3}}{3} (X_1^s + 3X_2^s)(4\pi),
\]

\[
\langle \Lambda^0 | O^d | \Xi^0 c \rangle = - 4X_2^d(4\pi), \quad \langle \Lambda^0 | O^s | \Xi^0 c \rangle = - 2(X_1^s - X_2^s)(4\pi),
\]

and for DCS processes are

\[
\langle \Sigma^0 | O^D | \Omega^0 c \rangle = \frac{4}{3} \sqrt{2} X_1^D(4\pi), \quad \langle \Lambda | O^D | \Omega^0 c \rangle = - 4\sqrt{6} X_2^D(4\pi).
\]
where we have introduced the bag integrals $X_1$ and $X_2$

$$X_1 = \int_0^R r^2 dr (u_s v_u - v_s u_u)(u_c v_d - v_c u_d), \quad X_2 = \int_0^R r^2 dr (u_s u_u + v_s v_u)(u_c u_d + v_c v_d),$$

$$X_1^d = \int_0^R r^2 dr (u_s v_u - v_s u_u)(u_c v_d - v_c u_d), \quad X_2^d = \int_0^R r^2 dr (u_s u_u + v_s v_u)(u_c u_d + v_c v_d),$$

$$X_1^s = \int_0^R r^2 dr (u_s v_u - v_s u_u)(u_c v_s - v_c u_s), \quad X_2^s = \int_0^R r^2 dr (u_s u_u + v_s v_u)(u_c u_s + v_c v_s),$$

$$X_1^D = \int_0^R r^2 dr (u_s v_u - v_s u_u)(u_c v_s - v_c u_s), \quad X_2^D = \int_0^R r^2 dr (u_s u_u + v_s v_u)(u_c u_s + v_c v_s),$$

(2.40)

with the numbers $X_1 = 3.56 \times 10^{-6}$, $X_2 = 1.74 \times 10^{-4}$, $X_1^d = 0$, $X_2^d = 1.60 \times 10^{-4}$, $X_1^s = 2.60 \times 10^{-6}$, $X_2^s = 1.96 \times 10^{-4}$, $X_1^D = 0$, $X_2^D = 1.78 \times 10^{-4}$. To obtain numerical results, we have employed the following bag parameters

$$m_u = m_d = 0, \quad m_s = 0.279 \text{ GeV}, \quad m_c = 1.551 \text{ GeV}, \quad R = 5 \text{ GeV}^{-1},$$

(2.41)

where $R$ is the radius of the bag.

### 2.5.2 Axial-vector form factors

The axial-vector form factor in the static limit can be expressed in the bag model as

$$g_{B' B}^{A(\nu)} = \langle B' \| b_1^\dagger b_2 \sigma_z | B \rangle \int d^3 r \left( u_{q_1} u_{q_2} - \frac{1}{3} v_{q_1} v_{q_2} \right).$$

(2.42)

The relevant results are

$$g_{\Sigma(\Lambda)^0\Xi^0}^{A(\nu)} = -\frac{\sqrt{6}}{3}(4\pi Z_2), \quad g_{\Xi^0(\Lambda)^0}^{A(\nu)} = \frac{2\sqrt{2}}{3}(4\pi Z_2), \quad g_{\Xi^0(\Lambda)^0}^{A(\pi^+)} = -\frac{1}{3}(4\pi Z_1),$$

$$g_{\Xi^0(\Lambda)^0}^{A(\pi^-)} = \frac{5}{3}(4\pi Z_2), \quad g_{\Xi^0(\Lambda)^0}^{A(\pi^\nu)} = 0, \quad g_{\Xi^0(\Lambda)^0}^{A(\pi^\sigma)} = \frac{1}{6}(4\pi Z_1),$$

$$g_{\Lambda^0\Lambda^0}^{A(\nu)} = \frac{\sqrt{6}}{3}(4\pi Z_1), \quad g_{\Lambda^0\Lambda^0}^{A(\nu)} = 0, \quad g_{\Lambda^0\Lambda^0}^{A(\pi^\nu)} = 0,$$

(2.43)
where the auxiliary bag integrals are given by

\[
Z_1 = \int r^2 dr \left( u_u^2 - \frac{1}{3} v_u^2 \right), \quad Z_2 = \int r^2 dr \left( u_u v_u - \frac{1}{3} v_u v_s \right).
\]

(2.44)

Numerically, \((4\pi)Z_1 = 0.65\) and \((4\pi)Z_2 = 0.71\).

3 Results and discussion

3.1 Numerical results and discussions

In this section, we shall numerically calculate branching fractions and up-down decay asymmetries. The decay asymmetries rely on \(S\)- and \(P\)-wave amplitudes, which have been calculated analytically yet. One more parameter, lifetime, enters the calculation of branching fractions based on decay width in Eq. (2.2). The value of lifetime quoted in this work is reported by LHCb in 2018 (see Eq. (1.2)), which is nearly four times larger than world-average value before 2018 and supported by a recent theoretical study on charmed baryon lifetime [5].

Factorizable and nonfactorizable amplitudes, branching fractions and decay asymmetries of all the two-body weak decays of \(\Omega_c\), including CF, SCS and DCS processes, are summarized in Table 4. The channel \(\Omega_c \rightarrow \Xi^0 K^0\) is the unique CF mode among all the \(\Omega_c \rightarrow B P\) decays, where \(B\) is baryon octet. In both \(S\)- and \(P\)-wave amplitudes, the nonfactorizable contributions are large and give destructive interference between factorizable ones. The branching fraction with full factorizable and nonfactorizable contributions is predicted to be 3.78%. The benchmarking channel \(\Omega_c \rightarrow \Omega^- \pi^+\), which is also classified into

| Channel          | \(A^{\text{fac}}\) | \(A^{\text{com}}\) | \(A^{\text{tot}}\) | \(B^{\text{fac}}\) | \(B^{\text{ca}}\) | \(B^{\text{tot}}\) | \(B^{\text{theo}}\) | \(B^{\text{expt}}\) | \(\alpha^{\text{theo}}\) |
|------------------|-------------------|-------------------|-------------------|-------------------|-------------------|-------------------|-------------------|-------------------|-------------------|
| \(\Omega_c^0 \rightarrow \Xi^0 K^0\) | -2.15             | 10.92             | 8.78              | -2.64             | 10.12             | 7.48              | 3.78              | -                 | 0.51              |
| \(\Omega_c^0 \rightarrow \Sigma^+ K^-\) | 0                 | -0.01             | -0.01             | 0                 | -6.10             | -6.10             | 2.32              | -                 | 0.01              |
| \(\Omega_c^0 \rightarrow \Sigma^0 K^0\) | 0                 | 0.01              | 0.01              | 0                 | -1.21             | -1.21             | 0.09              | -                 | -0.03             |
| \(\Omega_c^0 \rightarrow \Lambda^0 K^0\) | 0                 | -4.21             | -4.21             | 0                 | 0.04              | 0.04              | 8.05              | -                 | -0.01             |
| \(\Omega_c^+ \rightarrow \Xi^0 \pi^0\) | -0.88             | -2.43             | -3.31             | -1.21             | 1.00              | -0.21             | 5.46              | -                 | 0.04              |
| \(\Omega_c^0 \rightarrow \Xi^- \pi^+\) | -0.89             | -3.44             | -4.33             | -1.22             | 1.42              | 0.20              | 9.34              | -                 | -0.03             |
| \(\Omega_c^0 \rightarrow \Xi^- K^+\) | 0.10              | 1.34              | 1.43              | 0.13              | 0.49              | 0.62              | 9.58              | -                 | 0.27              |
| \(\Omega_c^0 \rightarrow \Xi^0 K^0\) | 0.10              | -1.34             | -1.24             | 0.131             | -0.49             | -0.36             | 7.04              | -                 | 0.18              |
| \(\Omega_c^0 \rightarrow \Lambda \eta\) | 0                 | -2.66             | -2.66             | 0                 | -2.56             | -2.56             | 36.28             | -                 | 0.66              |
| \(\Omega_c^0 \rightarrow \Sigma^0 \pi^0\) | 0                 | 0                 | 0                 | 0                 | -1.03             | -1.03             | 0.77              | -                 | 0                 |
| \(\Omega_c^0 \rightarrow \Sigma^+ \pi^-\) | 0                 | 0                 | 0                 | -1.03             | -1.03             | 0.77              | -                 | 0                 |
| \(\Omega_c^0 \rightarrow \Sigma^- \pi^+\) | 0                 | 0                 | 0                 | -1.03             | -1.03             | 0.77              | -                 | 0                 |

Table 4. Decays \(\Omega_c \rightarrow B P\): the amplitudes are in units of \(10^{-2}\)\(G_F\)GeV\(^2\), branching fractions for CF(SCS, DCS) process(es) is (are) in unit(s) of \(10^{-2}(10^{-3}, 10^{-4})\), and the asymmetry parameters \(\alpha\) are shown in the last column.
Although no explicit measurement of the mode $\Omega_c \rightarrow \Xi^0 \overline{K}^0$ are turned off, the predicted value of its branching fraction would be 2.44%. This partially helps to understand the Belle measurement of large relative ratio between $\Omega_c \rightarrow \Xi^0 \overline{K}^0$ and $\Omega_c \rightarrow \Omega^- \pi^+$, see Eq. (1.3). A detailed consideration for these channels decaying into baryon decuplet will be presented in a separate work. $^1$ Although no explicit measurement of the mode $\Omega_c \rightarrow \Xi^0 \overline{K}^0$ has been given, a large branching fraction prediction indicates a direct measurement is promising in the near future. The decay asymmetry $\alpha$ is predicted to be positive and with a measurable value 0.51, which is also testable when more data is available.

The three decay modes $\Omega_c \rightarrow \Sigma^+ K^-, \Sigma^0 \overline{K}^0, \Lambda^0 \overline{K}^0$ in SCS channels, which do not receive factorizable contributions, are typical examples for the essential role of nonfactorizable contribution in charmed baryon decays. Due to the breaking SU(3) flavor symmetry, see parameter $X_1^s$, the $S$-wave amplitudes for the modes with $\Sigma$ baryon final states are tiny but not vanishing. On the other hand, $S$-wave amplitude for $\Omega_c \rightarrow \Lambda^0 \overline{K}^0$ is significantly enhanced from $\Sigma^0 \overline{K}^0$ for its typical size is described by $X_2^s$, which is two orders of magnitude larger than $X_1^s$. Among the $P$-wave terms of the three modes, $\Sigma^+ K^-$ is the largest one as cancellation occurs in the other two modes. Since $S$-wave amplitude dominates branching fraction, according to Eq. (2.2), $\Omega_c \rightarrow \Lambda^0 \overline{K}^0$ is predicted with the largest branching fraction among the three modes. However, the decay asymmetries for all the three modes is tiny, which is a natural consequence of tiny value for either $A$ or $B$. The remaining two modes in SCS processes, $\Omega_c \rightarrow \Xi^0 \pi^0$ and $\Xi^- \pi^+$, share almost same factorizable contributions while in both channels the nonfactorizable terms contribute constructively in $S$-wave and destructively in $P$-wave terms, leading to tiny decay asymmetries again.

Predictions for DCS channels are also summarized in Table 4. The mode $\Omega_c \rightarrow \Lambda^0 \eta$ is of particular interest in all the DCS channels. Its $S$- and $P$-wave amplitudes, which both are depicted by $X_2^D$, are substantial and hence lead to a large branching fraction 0.36%. The large and positive decay asymmetry 0.66 is also predicted. Among all the modes which contain net nonfactorizable contribution, one of the two of the reasons $I_3(\Sigma^0) = I_3(\Lambda) = 0$ and $X_1^D = 0$ determines the vanishing $S$-wave amplitudes for the channels $\Omega_c \rightarrow \Sigma \pi / \eta$ and $\Omega_c \rightarrow \Lambda \pi^0$. And the properties $g_{A^0 A^0}^{A(\pi \eta)} = g_{\Sigma^0 A^0}^{A(\pi \eta)} = 0$ forbid the two decays $\Omega_c \rightarrow \Sigma^0 \eta$ and $\Lambda^0 \pi^0$ furthermore.

3.2 Comparison with other works

In the early 1990s, there were many efforts to study charmed baryon decays, among which few were $\Omega_c$ involved [18, 22, 23, 40]. Later semileptonic decays of heavy $\Omega$ baryons, including $\Omega_c$, was studied in [41, 42]. In recent years there have been some interests on its hadronic weak decays [32, 41, 43], in which [32] focused on modes with axial-vector final state and only modes with decuplet baryon final state were partially involved in

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$^1$ A prediction for branching fraction of $\Omega_c \rightarrow \Omega^- \pi^+$, based on an early work in [40] with updated $a_1 = 1.26$ and latest $\Omega_c$ lifetime, is of order 9%. The incompatibility among our prediction of $\Omega_c \rightarrow \Xi^0 \overline{K}^0$, the prediction of $\Omega_c \rightarrow \Omega^- \pi^+$ in [40] and Belle measurement Eq. (1.3) will also be discussed therein.
Table 5. Predicted branching fractions in the unit of $10^{-2}, 10^{-3}$ and $10^{-4}$ (upper entry in each mode) and decay asymmetry $\alpha$ (lower entry) of $\Omega_c$ decays by different groups.

| Mode          | Our | Cheng et al. CA [18] | Cheng et al. pole model [18] | Körner et al. [22] | Ivanov et al. [23] | Zhao [41] |
|---------------|-----|----------------------|------------------------------|-------------------|----------------|---------|
| $\Omega_c \rightarrow \Xi^0 K^0$ | 3.78 | 2.63                 | 0.35                         | 4.69              | 0.09          |         |
|               | 0.51 | 0.44                 | -0.93                        | 0.51              | -0.81         |         |
| $\Omega_c \rightarrow \Xi^- \pi^+$ | 9.34 |                      |                              |                   |                | 0.7     |
|               | -0.03 |                     |                              |                   |                |         |
| $\Omega_c \rightarrow \Xi^- K^+$ | 9.58 |                      |                              |                   |                | 0.6     |
|               | 0.27 |                      |                              |                   |                |         |

In Table 5 a comparison with other groups, whose predictions have been updated by incorporating current $\Omega_c$ lifetime, is summarized in available channels.

The CF channel $\Omega_c \rightarrow \Xi^0 K^0$ attracts more attention in the past [18, 22, 23] and nonfactorizable contributions have been incorporated by all groups. Based on the pole mode combing current algebra, our results both for branching fraction and decay asymmetry can be confirmed by the early calculation within the same approach [18]. However, the prediction for branching fraction is around 10 times larger than an early estimation relied on pure pole model [18], and the sign of decay asymmetry is opposite. Small branching fraction and negative asymmetry were also predicted within a relativistic three-quark model with a Gaussian shape for the momentum dependence of baryon-three-quark vertex [23]. Such situation occurred in the studies on anti-triplet charmed baryons [25, 26]. Taking the mode $\Lambda_c \rightarrow \Sigma^+ \pi^0$ as an example, the pure pole model in [18] and quark model calculation in [23] both predicted positive decay asymmetry while current algebra predicted $\alpha = -0.76$, which is consistent with experimental value $\alpha = -0.55 \pm 0.11$ [1]. Interestingly, working in an independent approach, Körner-Krämer gave a consistent prediction for both branching fraction and decay asymmetry in covariant quark model [22].

The branching fractions of SCS process $\Omega_c \rightarrow \Xi^- \pi^+$ and DCS process $\Omega_c \rightarrow \Xi^- K^+$ were estimated in [41], where baryon-baryon transition form factors were calculated in light-front quark model and only factorizable contribution was taken into account. It has been widely accepted that nonfactorizable contribution should play an essential role in the hadronic decays. The numerical results for individual term in Table 4 shows that nonfactorizable terms even give dominated contributions, which helps to explain why our prediction is more than 10 times larger.

4 Conclusions

In this work we have systematically studied the branching fractions and up-down decay asymmetries of CF, SCS and DCS decays of $\Omega_c$, the heaviest charmed baryon which receiv-
ing weak decays dominately. Both factorizable and nonfactorizable terms have been taken into account in the calculation of $S$- and $P$-wave amplitudes. To estimate nonfactorizable contribution, we work in the pole model for $P$-wave amplitudes and current algebra for $S$-wave ones. All the non-perturbative parameters, including baryon-baryon transition form factors, baryon matrix elements and axial-vector form factors, are evaluated within MIT bag model throughout the whole calculations.

Some conclusions can be drawn from our analysis as follows.

- The channel $\Omega_c \to \Xi^0 K^0$ is the unique mode for CF decay. Although no absolute branching fraction has been measured up to now, the predicted large value for branching fraction indicates this mode is quite promising to be measured in the near future. Meanwhile its decay asymmetry is predicted to be large in magnitude and positive in sign.

- Among all the SCS modes, the channel $\Omega_c \to \Lambda^0 K^0$ is special as the decay amplitudes are contributed from purely nonfactorizable contribution. Though $P$-wave amplitude is small, its large $S$-wave amplitude leads to a large branching ratio. Hence a measurement of $\Omega_c \to \Lambda^0 K^0$ in the future will demonstrate the essential role of nonfactorizable contribution in charmed baryon weak decay.

- The decay asymmetries of all SCS modes are small, which is a natural consequence of either small $S$- or $P$-wave amplitude. In other words, it is difficult to measure decay asymmetries of SCS $\Omega_c$ weak decay in experiment.

- The measurement of $\Omega_c \to \Lambda \eta$ would also be interesting. Although this mode is classified as the DCS mode, not only the branching fraction is predicted to be large, but also its decay asymmetry is predicted to be large in magnitude and positive in sign, which makes the measurement workable in experiment. Such features can be explained by simultaneous large $S$- and $P$-wave amplitudes.

- The two DCS modes $\Omega_c \to \Sigma^0 \eta$ and $\Omega_c \to \Lambda^0 \pi^0$ are forbidden, for both $S$- and $P$-wave amplitudes are found to be zero. On the other hand, these two modes can serve as golden channels for new physics searching.

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