Singly Cabibbo suppressed decays of $\Lambda_c^+$ with SU(3) flavor symmetry

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Abstract

We analyze the weak processes of anti-triplet charmed baryons decaying to octet baryons and mesons with the SU(3) flavor symmetry and topological quark diagram scheme. We study the decay branching ratios without neglecting the contributions from $\mathcal{O}(15)$ for the first time in the SU(3) flavor symmetry approach. The fitting results for the Cabibbo allowed and suppressed decays of $\Lambda_c^+$ are all consistent with the experimental data. We predict all singly Cabibbo suppressed decays. In particular, we find that $\mathcal{B}(\Lambda_c^+ \to p\pi^0) = (1.3 \pm 0.7) \times 10^{-4}$, which is slightly below the current experimental upper limit of $2.7 \times 10^{-4}$ and can be tested by the ongoing experiment at BESIII as well as the future one at Belle-II.
Recently, the study of the charmed baryons has been receiving increasing attention both theoretically and experimentally. The main reason for this is the recent measurement of the absolute branching fraction of the golden channel $\Lambda_c^+ \to pK^−\pi^+$ by the Belle Collaboration [1]. This mode and many other $\Lambda_c^+$ ones have also been observed by the BESIII Collaboration [2–10] with using $\Lambda_c^+\bar{\Lambda}_c^-$ pairs produced by $e^+e^−$ collisions at a center-of-mass energy of $\sqrt{s} = 4.6$ GeV, which provides a uniquely clean background to study charmed baryons. Consequently, the Particle Data Group (PDG) [11] has given a new average of $B(\Lambda_c^+ \to pK^−\pi^+) = (6.23 \pm 0.33)\%$. The precision measurement on this mode is very important as it can be used to determine the absolute branching fractions of other $\Lambda_c^+$ decays [11] as well as processes involving $\Lambda_c^+$, such as the extractions of the CKM element from $\Lambda_b \to \Lambda_c^+\mu^−\bar{\nu}_\mu$ [12, 13]. It is clear that a new era of physics for charmed baryons has begun. For a review on the theoretical progress of charmed baryons, please see Ref. [14].

On the other hand, the singly Cabibbo suppressed decays of $\Lambda_c^+ \to p\eta$ and $\Lambda_c^+ \to p\pi^0$ have been recently investigated by BESIII [6]. The branching fraction of the former mode has been measured for the first time with $B(\Lambda_c^+ \to p\eta) = (1.24 \pm 0.28 \pm 0.10) \times 10^{-3}$, whereas that of the later one has also been searched with no significant signal observed, resulting in an upper limit of $B(\Lambda_c^+ \to p\pi^0) < 2.7 \times 10^{-4}$ at the 90% confidence level. These two decays have been extensively studied in the literature based on various dynamical models [15–17] as well as the flavor $SU(3)_F$ symmetry [18–23]. In particular, Cheng, Kang and Xu (CKX) [17] have performed a dynamical calculation based on current algebra to examine the decay of $\Lambda_c^+ \to p\pi^0$ and found that its branching fraction is $0.8 \times 10^{-4}$, which is consistent with the current experimental upper limit. However, those with $SU(3)_F$ have given an inconsistent larger value, e.g., $(5.7 \pm 1.5) \times 10^{-4}$ in Ref. [22].

It is known that it is difficult to make reliable predictions on the charmed baryon decay rates due to the lack of theoretical understanding of underlined dynamics for the charmed baryon structure. Since the Cabbibo allowed decays of $\Lambda_c^+ \to \Sigma^0\pi^+$ and $\Lambda_c^+ \to \Sigma^+\pi^0$ do not receive any factorizable contributions, the nonzero experimental observed values of their branching fractions imply that the factorization approach is not working in charmed baryon decays. Without the use of a dynamical model, it is clear that the most reliable way to analyze charmed baryon processes is to impose $SU(3)_F$ [18–26]. In fact, it has been demonstrated [20, 23] that all the existing data of the Cabbibo favored and suppressed charmed baryon decays except $B(\Lambda_c^+ \to p\pi^0)$ can be fitted well. In these calculations under
The contributions to the decays from the sextet 6 are assumed to be the dominant ones, whereas those from $\mathbf{15}$ are neglected, by taking into account of the enhancements of the QCD running Wilson coefficients associated with the sextet 6 part [27–30] and the vanishing baryonic transition matrix elements from the nonfactorizable contributions with $\mathbf{15}$ [17]. However, it is interesting to ask what the contributions to the decay rates from the factorizable parts of $\mathbf{15}$ are. In this note, we will try to answer this question. Specifically, we examine all possible contributions to the charmed baryon decays under $SU(3)_F$ without neglecting those from $\mathbf{15}$. We examine the singly Cabibbo suppressed $\Lambda^+_c$ decays to check if our results are consistent with the data, in particular, the $\Lambda^+_c \rightarrow p\pi^0$ channel.

There are two approaches to write down the irreducible decay amplitude through $SU(3)_F$. One is to generalize the Wigner Eckart theorem [31] by writing the decay amplitude to be invariant and singlet under $SU(3)_F$. The other is to use topological quark diagrams, where the decay amplitude is represented by all possible diagrams connected by quark lines which satisfy $SU(3)_F$. Both two have their own advantages. For the former, one is able to compare the contributions from different representations of operators. In this case, it is also possible to include the $SU(3)_F$ breaking effect by introducing the strange quark mass [23, 25]. On the other hand, the irreducible amplitude in the later approach is more intuitive and gives an insight on dynamics [32]. In particular, it could shed light for us on distinguishing the nonfactorizable and factorizable contributions in the processes. It is expected that these two approaches should give the same results under $SU(3)_F$. The close connections between the two have been recently examined in Ref. [33].

To study the two-body anti-triplet of the lowest-lying charmed baryon decays of $B_c \rightarrow B_n M$, where $B_c = (\Xi^0_c, -\Xi^+_c, \Lambda^+_c)$ and $B_n$ and $M$ are the baryon and pseudoscalar octet states, given by

\begin{equation}
B_n = \begin{pmatrix}
\frac{1}{\sqrt{6}}\Lambda + \frac{1}{\sqrt{2}}\Sigma^0 & \Sigma^+ & p \\
\Sigma^- & \frac{1}{\sqrt{6}}\Lambda - \frac{1}{\sqrt{2}}\Sigma^0 & n \\
\Xi^- & \Xi^0 & -\sqrt{\frac{2}{3}}\Lambda
\end{pmatrix},
\end{equation}

\begin{equation}
M = \begin{pmatrix}
\frac{1}{\sqrt{6}}\eta + \frac{1}{\sqrt{2}}\eta^0 & \pi^+ & K^+ \\
\pi^- & \frac{1}{\sqrt{6}}\eta - \frac{1}{\sqrt{2}}\eta^0 & K^0 \\
K^- & \bar{K}^0 & -\sqrt{\frac{2}{3}}\eta
\end{pmatrix}.
\end{equation}
Here, we have assumed that the physical state of $\eta$ is solely made of $\eta_8$ due to the small mixing between the weak eigenstates of $\eta_0$ and $\eta_8$ to reduce our fitting parameters.

We start with the effective Hamiltonian responsible for the tree-level $c \to s u \bar{d}$, $c \to u q \bar{q}$ and $c \to d u s \bar{s}$ transitions, given by \cite{34}

$$H_{\text{eff}} = \sum_{i=+, -} \frac{G_F}{\sqrt{2}} C_i \left( V_{cs} V_{ud} O^{ds}_i + V_{cd} V_{us} O^{sq}_i + V_{cd} V_{us} O^{sd}_i \right),$$

with

$$O^{sq}_{\pm i} = \frac{1}{2} \left[ (\bar{u} q_1)(\bar{q}_2 c) \pm (\bar{q}_2 q_1)(\bar{u} c) \right],$$

where $O^{sq}_{\pm i}$ and $O^{sq}_{\pm} \equiv O^{dd}_{\pm} - O^{ss}_{\pm}$ are the four-quark operators, $(q_1 q_2) \equiv \bar{q}_1 \gamma_{\mu}(1 - \gamma_5) q_2$, $G_F$ is the Fermi constant, $V_{ij}$ are the CKM matrix elements, and $(c_+, c_-) = (0.76, 1.78)$, corresponding to the scale-dependent Wilson coefficients with the QCD corrections. By using $(V_{cs} V_{ud}, V_{cd} V_{us}, V_{cd} V_{us}) \simeq (1, -s_c, -s_c^2)$ in Eq. (3) with $s_c \equiv \sin \theta_c = 0.2248$ \cite{11} representing the well-known Cabibbo angle $\theta_c$, the decays associated with $O^{ds}_{\pm}, O^{sq}_{\pm}$ and $O^{sd}_{\pm}$ are the so-called Cabibbo-allowed, singly Cabibbo-suppressed and doubly Cabibbo-suppressed processes, respectively.

Under SU(3)$_F$, the operators in Eq. (4) correspond to $(\bar{q}^i q_k)(\bar{q}^j c)$ with $q_i = (u, d, s)$ as the triplet of 3, which can be decomposed as the irreducible forms of $(3 \times 3 \times 3)c = (3 + \bar{3} + 6 + \bar{1}5)c$ with $c$ as a flavor singlet. As a result, $(O_-, O_+)$ fall into the irreducible presentations of $(O_6, O_{15})$ \cite{24}. In analogy to octet baryons and mesons, we can write down the operators related to $O_6$ and $O_{15}$ in tensor forms, given by

$$H(6)_{ij} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 2 & -2s_c \\ 0 & -2s_c & 2s_c^2 \end{pmatrix},$$

$$H(15)_{ij} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & s_c \\ 0 & s_c & 0 \end{pmatrix}, \begin{pmatrix} 0 & s_c & 1 \\ 0 & -s_c^2 & -s_c \\ -s_c^2 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 0 & 0 \\ 1 & 0 & 0 \\ -s_c & 0 & 0 \end{pmatrix},$$

with $(i, j, k) = 1, 2$ and 3, where $H(15)$ is traceless and symmetric in upper indices, while $H(6)_{ij}$ is symmetric in lower indices. One can also write the matrix elements of $H(6)_{ij}$ and $H(15)_{ij}$ in Eq. (5) as a single one, given by

$$H_{k}^{ij} = \frac{1}{2} \left( H(15)_{ij} + \frac{1}{2} \epsilon^{ijkl} H(6)_{kl} \right).$$

(6)
Now, we can write down the SU(3) irreducible amplitude for $B_c \to B_n M$ as

$$A(B_c \to B_n M) = \langle B_n M|\mathcal{H}_{\text{eff}}|B_c \rangle \equiv \frac{G_F}{\sqrt{2}} (T_{C_6} + T_{C_{15}}) ,$$

where

$$T_{C_6} = a_1 H_{ij}(6)(B'_c)^{jk}(B_n)^{il}_k(M)^l_j + a_2 H_{ij}(6)(B'_c)^{jk}(M)^l_k(B_n)^{il}_j + a_3 H_{ij}(6)(B_n)^{il}_k(M)^l_j(B'_c)^{jk}$$

$$T_{C_{15}} = a_4 H^u_k(15)(B_c)^{ij}_k(M)^l_i + a_5 (B_n)^{ij}_k(M)^l_i H(15)^k_j(B_c)^l_j + a_6 (B_n)^{ij}_k(M)^l_i H(15)^k_j(B_c)^l_k ,$$

with $(B'_c)^{jk} \equiv (B_c)_i e^{ijk}$. Here, the Wilson coefficients have been absorbed in the parameters $a_i$.

In order to reduce the fitting parameters for the processes based on the amplitudes in Eqs. (7) and (8), as mentioned earlier, the contributions related to $O(15)$ in Eq. (7) have been neglected due to the fact that $c_-/c_+ \approx 2.5$ and the vanishing contributions of $O(15)$ from the nonfactorizable part to the amplitude. To see the later reason, we write the amplitude of $B_c \to B_n M$ due to $O(15)$ in terms of the matrix element

$$A(O(15)) = \langle B_n M|O(15)|B_c \rangle = \frac{1}{2} \langle B_n M|(\bar{u}q_1)(\bar{q}2c) + (\bar{q}2q_1)(\bar{uc})|B_c \rangle .$$

Since the operator $O(15) \sim (\bar{u}q_1)(\bar{q}2c) + (\bar{q}2q_1)(\bar{uc})$ is symmetric in color indices, whereas the baryon states $B_i$ are antisymmetric, one easily arrives that $\langle B_i|O(15)|B_j \rangle = 0$. From the calculations of the nonfactorizable (NF) contributions in terms of the baryon poles $(B^*)_n$, one has that $A(O_{NF(15)})$ is related to the combination of $g_{B_n B^* M}(B^*|O(15)|B_c)$ and $g_{B_n M}(B_n|O(15)|B_c)$ as illustrated in Fig. 1 and 2, respectively. Note that the quark indices

\[ \text{It is clear that we have ignored the soft gluon interactions whenever the factorization problem is discussed.} \]
FIG. 1: Pole diagrams of the nonfactorizable amplitude for $B_c \rightarrow B_n M$.

represent the light quark lines of hadrons or operators with $q_i = (u, d, s)$. From Fig. 2 we obtain that

$$A_F(B_c \rightarrow B_n M) = T(B'_n)_{ijk} (B'_c)^{jk} H_{ml} M_i^m + C(B'_n)_{ijk} (B'_c)^{jk} H_{ml} M_i^m$$

(11)

where $(B'_n)_{ijk} \equiv (B_n)^i_i \epsilon_{njk}$, $(B'_c)^{jk} \equiv (B_c)^m_{m} \epsilon^{mjk}$, and $T(C)$ represents the color allowed (suppressed) amplitude. By using Eq. (6) and the tensor identity $\epsilon_{njk} \epsilon^{mjk} = 2 \delta_{mn}$, we find that

$$A_F(B_c \rightarrow B_n M) = T(B'_n)^{n}(B'_c)^{n} H(\overline{15})_{ml} M_i^m + C(B'_n)^{n}(B'_c)^{n} H(\overline{15})_{ml} M_i^m + A_F(O(6))$$

$$= (T + C)(B'_n)^{n}(B'_c)^{n} H(\overline{15})_{ml} M_i^m + A_F(O(6))$$

(12)

where $A_F(O(6))$ corresponds to the factorizable amplitude from $O(6)$. Here, we have used that $H(\overline{15})$ is symmetry in upper indices in the second line of Eq. (12). By comparing Eq. (12) with Eq. (8), we immediately identify that only the $a_6$ term in Eq. (8) contains the factorizable amplitude of $O(\overline{15})$. Consequently, we can safely neglect the $a_4$, $a_5$ and $a_7$ terms in $T_{O(6)}$ of Eq. (8) as they do not have the factorizable contributions to the processes. We remark that in Eq. (12), if $T$ and $C$ both exist in $A_F(O(\overline{15}))$, one of them should be

FIG. 2: Topological diagram for color allowed and suppressed processes.

2 In general, the term associated with $a_6$ also contribute the non-factorizable part.
canceled out by the corresponding term in $A_F(O(6))$, resulting in the process to be either color allowed or color suppressed. This can be explicitly demonstrated by the recent work in Ref. \[33\] on the connection between the topological and $SU(3)_F$ approaches.

To illustrate the effect of the only $a_6$ term from $O(15)$, we show the decay amplitudes of $\Lambda_c^+ \to p\pi^0$ and $\Lambda_c^+ \to n\pi^+$, given by \[23\]

$$A(\Lambda_c^+ \to p\pi^0) \propto \sqrt{2} \left( a_2 + a_3 - \frac{a_6 - a_7}{2} \right) = \sqrt{2} \left( a_2 + a_3 - \frac{a_6}{2} \right),$$

$$A(\Lambda_c^+ \to n\pi^+) \propto 2 \left( a_2 + a_3 + \frac{a_6 + a_7}{2} \right) = 2 \left( a_2 + a_3 + \frac{a_6}{2} \right). \quad (13)$$

It is clear that the relation of $A(\Lambda_c^+ \to n\pi^+) = \sqrt{2} A(\Lambda_c^+ \to p\pi^0)$ \[19\] is violated with the contributions from $a_6$. This violation has been explicitly pointed out in Ref. \[17\] based on a dynamical model. On the other hand, some direct relations still exist in some modes. For example, one has that

$$A(\Lambda_c^+ \to \Sigma^0 K^+) \propto \sqrt{2} \left( a_1 - a_3 - \frac{a_4 + a_5}{2} \right) = \sqrt{2} (a_1 - a_3),$$

$$A(\Lambda_c^+ \to \Sigma^+ K_S^0) \propto \sqrt{2} \left( a_1 - a_3 - \frac{-a_4 + a_5}{2} \right) = \sqrt{2} (a_1 - a_3). \quad (14)$$

Future experimental searches for these decays will confirm if the discussions based on $SU(3)_F$ are right or not.

We are now ready to perform our numerical calculation. Since the $SU(3)_F$ flavor symmetry does not involve the dynamical details, we have to determine the parameters in the irreducible amplitude by the experimental data, which can be found in the PDG \[11\] along with the recent measurements by BESIII \[2, 6\]. Currently, there are 9 data points from the absolute branching fractions, along with the original data of $B(\Lambda_c^+ \to p\pi^0) = (0.8 \pm 1.3) \times 10^{-4}$ by BESIII \[36\], which are summarized in Table \[1\]. In addition, we include the relative branching ratio of $R_{\Xi_c} \equiv B(\Xi_c^0 \to \Lambda^0 K^0)/B(\Xi_c^0 \to \Xi^- \pi^+) = 0.420 \pm 0.056$ in our fitting. Altogether, there are seven $SU(3)_F$ parameters $(a_1, |a_2|e^{i\delta_{a2}}, |a_3|e^{i\delta_{a3}}, |a_6|e^{i\delta_{a6}})$ to fit with eleven data points in Table \[1\]. Here, we have set $a_1$ to be a real parameter due to the removal of an overall phase. We use the minimum $\chi^2$ fit as shown in Ref. \[22\]. Explicitly, we obtain

$$(a_1, |a_2|, |a_3|, |a_6|) = (0.271 \pm 0.006, 0.126 \pm 0.010, 0.051 \pm 0.012, 0.055 \pm 0.030) GeV^3,$$

$$\delta_{a2}, \delta_{a3}, \delta_{a6} = (82 \pm 6, -20 \pm 24, 40 \pm 36) \degree,$$

$$\chi^2/d.o.f = 0.5, \quad (15)$$
TABLE 1: Decay amplitudes related to the $SU(3)_F$ parameters and the experimental data for the absolute branching fractions and $R_{\Xi^0}$ \cite{2, 6, 11, 36}.

| Channel | Amplitude | Data |
|---------|-----------|------|
| $\Lambda_c^+ \to \Sigma^+ \pi^0$ | $\sqrt{2}(a_1 - a_2 - a_3)$ | $(12.4 \pm 1.0) \times 10^{-3}$ |
| $\Lambda_c^+ \to \Sigma^+ \eta$ | $\frac{\sqrt{6}}{3}(-a_1 - a_2 + a_3)$ | $(7.0 \pm 2.3) \times 10^{-3}$ |
| $\Lambda_c^+ \to \Sigma^0 \pi^0$ | $\sqrt{2}(-a_1 + a_2 + a_3)$ | $(12.9 \pm 0.7) \times 10^{-3}$ |
| $\Lambda_c^+ \to \Xi^0 K^+$ | $-2a_2$ | $(5.9 \pm 1.0) \times 10^{-3}$ |
| $\Lambda_c^+ \to p\bar{K}^0$ | $-2a_1 + a_6$ | $(31.6 \pm 1.6) \times 10^{-3}$ |
| $\Lambda_c^+ \to \Lambda^0 \pi^+$ | $\frac{\sqrt{6}}{3}(-a_1 - a_2 - a_3 - a_6)$ | $(13.0 \pm 0.7) \times 10^{-3}$ |
| $\Lambda_c^+ \to \Xi^0 K^+$ | $\sqrt{2}(a_1 - a_3)$ | $(5.2 \pm 0.8) \times 10^{-4}$ |
| $\Lambda_c^+ \to \Sigma^+ K^0$ | $2(a_1 - a_3)$ | - |
| $\Lambda_c^+ \to p\pi^0$ | $\sqrt{2}(a_2 + a_3 - \frac{a_6}{2})$ | $(0.8 \pm 1.3) \times 10^{-4}$ |
| $\Lambda_c^+ \to n\pi^+$ | $2(a_2 + a_3 + \frac{a_6}{2})$ | - |
| $\Lambda_c^+ \to p\eta$ | $\frac{\sqrt{6}}{3}(-2a_1 + a_2 - a_3 + \frac{3}{2}a_6)$ | $(12.4 \pm 3.0) \times 10^{-4}$ |
| $\Lambda_c^+ \to \Lambda^0 K^+$ | $\frac{\sqrt{6}}{3}(a_1 - 2a_2 + a_3 + a_6)$ | $(6.1 \pm 1.2) \times 10^{-4}$ |
| $\Xi^0 \to \Xi^- \pi^+$ | $2a_1 + a_6$ | - |
| $\Xi^0 \to \Lambda^0 \bar{K}^0$ | $\sqrt{6}(2a_1 + a_2 + a_3 + \frac{a_6}{2})$ | - |
| $R_{\Xi^0}$ | $0.420 \pm 0.056$ | |

where $d.o.f$ represents the degree of freedom. The value of $\chi^2/d.o.f$ indicates that our fit is good. In the previous studies of the $\Lambda_c^+$ decays based on $SU(3)_F$ \cite{19, 23}, the contributions of $\mathcal{O}(15)$ have been neglected. Our results in Eq. \cite{15} show that the absolute value of $a_6$ is about $1/6$ compared to that of the leading one $a_1$ in $T(\mathcal{O}_6)$, so that the ignorance of $\mathcal{O}(15)$ is indeed valid.

In Table 2, we list our fitting results for the branching ratios of the Cabibbo allowed and singly Cabibbo suppressed $\Lambda_c^+$ decays. In the table, we have also included the previous results based on $SU(3)_F$ \cite{22} without $\mathcal{O}(15)$ along with the data as well as those from the dynamical model calculations by CKX \cite{17}. As seen in Table 2, our results for the Cabibbo
TABLE 2: Branching ratios for the Cabibbo allowed and singly Cabibbo suppressed decays of $\Lambda_c^+$.

| Decay branching ratio  | This work | Data    | $SU(3)_F$ [22] | CKX [17] |
|------------------------|-----------|---------|----------------|----------|
| $10^3\mathcal{B}(\Lambda_c^+ \to \Sigma^+\pi^0)$ | 12.6 ± 2.1 | 12.4 ± 1.0 | 12.8 ± 2.3 | -        |
| $10^3\mathcal{B}(\Lambda_c^+ \to \Sigma^+\eta)$   | 5.4 ± 1.0  | 7.0 ± 2.3 | 7.1 ± 3.8  | -        |
| $10^3\mathcal{B}(\Lambda_c^+ \to \Sigma^0\pi^0)$  | 12.6 ± 2.1 | 12.9 ± 0.7 | 12.8 ± 2.3 | -        |
| $10^3\mathcal{B}(\Lambda_c^+ \to \Xi^0 K^+)$      | 5.9 ± 1.0  | 5.9 ± 0.9 | 5.5 ± 1.4  | -        |
| $10^3\mathcal{B}(\Lambda_c^+ \to p\bar{K}^0)$     | 31.3 ± 1.6 | 31.6 ± 1.6 | 32.7 ± 1.5 | -        |
| $10^3\mathcal{B}(\Lambda_c^+ \to \Lambda^0\pi^0)$  | 13.1 ± 1.6 | 13.0 ± 0.7 | 12.8 ± 1.7 | -        |
| $10^4\mathcal{B}(\Lambda_c^+ \to \Sigma^+ K^0)$   | 11.4 ± 2.0 | -       | 8.0 ± 1.6  | 14.4     |
| $10^4\mathcal{B}(\Lambda_c^+ \to \Sigma^0 K^+)$   | 5.7 ± 1.0  | 5.2 ± 0.8 | 4.0 ± 0.8  | 7.18     |
| $10^4\mathcal{B}(\Lambda_c^+ \to p\pi^0)$         | 1.3 ± 0.7  | < 2.7   | 5.7 ± 1.5  | 0.8      |
| $10^4\mathcal{B}(\Lambda_c^+ \to p\eta)$          | 13.0 ± 1.0 | 12.4 ± 3.0 | 12.5$^{+3.8}_{-3.6}$ | 12.8 |
| $10^4\mathcal{B}(\Lambda_c^+ \to n\pi^0)$         | 6.1 ± 2.0  | -       | 11.3 ± 2.9 | 2.7      |
| $10^4\mathcal{B}(\Lambda_c^+ \to \Lambda^0 K^+)$  | 6.4 ± 0.9  | 6.1 ± 1.2 | 4.6 ± 0.9  | 10.6     |

allowed $\Lambda_c^+$ decays with the consideration of $O(\overline{15})$ are slightly better than those without $O(\overline{15})$, but they all fit the data well. On the other hand, the decay branching ratios for singly Cabibbo suppressed modes of $\Lambda_c^+$ with and without $O(\overline{15})$ are quite different. In particular, we predict that $\mathcal{B}(\Lambda_c^+ \to p\pi^0) = (1.3 \pm 0.7) \times 10^{-4}$, which is consistent with the experiments upper limit of $2.7 \times 10^{-4}$ as well as the result of $0.8 \times 10^{-4}$ calculated by the pole model with current algebra in Ref. [17]. It is clear that the inconsistent branching ratio of $(5.7 \pm 1.5) \times 10^{-4}$ in the previous study with $SU(3)_F$ [22] results from the ignorance of $O(\overline{15})$, in which a large destructive interference occurs between $O(\overline{15})$ and $O(6)$. It is also interesting to note that $\mathcal{B}(\Lambda_c^+ \to n\pi^0)$ is found to be $(6.1 \pm 2.0) \times 10^{-4}$, which is reduced by almost a factor 2 in comparing with that in Ref. [22]. Although the signs for the contributions from $a_6$ to $\Lambda_c^+ \to p\pi^0$ and $\Lambda_c^+ \to n\pi^0$ in Eq. (13) are opposite, the resulting values are both reduced due to the complex numbers of $a_{2,3}$ and $a_6$ in Eq. (13).
In addition, from Table 2 we have that

\[ B(\Lambda_c^+ \to \Sigma^+ K_S^0) = (5.7 \pm 1.0) \times 10^{-4}, \]  

which agrees with the experimental value of \( B(\Lambda_c^+ \to \Sigma^0 K^+) = (5.2 \pm 0.8) \times 10^{-4} \) \[11\]. The future search for \( \Lambda_c^+ \to \Sigma^+ K_S^0 \) is a good test for \( SU(3)_F \).

Finally, we remark that we are unable to discuss the \( SU(3)_F \) breaking effects after including the contributions of \( \mathcal{O}(15) \) in the fit due to the insufficient experimental data points. Once more experimental data are available in the future, the studies of these effects along with the \( \eta' \) channels would be possible.

In sum, we have studied the two-body decays of \( \Lambda_c^+ \to B_nM \) based on the approach with the \( SU(3)_F \) flavor symmetry, which is a powerful tool to examine charmed baryon physics and allows us to connect the physical quantities without knowing the underlined dynamics. We have successfully fitted all the existing experimental data from the Cabibbo allowed and suppressed decays of \( \Lambda_c^+ \). By considering the approach with the topological quark diagrams, for the first time, the contributions from \( \mathcal{O}(15) \) have been included in the calculations with the \( SU(3)_F \) method. As a result, we have predicted all singly Cabibbo suppressed decays. In particular, we have found that \( B(\Lambda_c^+ \to p\pi^0) = (1.3 \pm 0.7) \times 10^{-4} \), which is slightly below the current experimental upper limit of \( 2.7 \times 10^{-4} \). This result can be tested by the experiments at BESIII and Belle-II.

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