Axial-vector emitting weak nonleptonic decays of Ω_c^0 baryon

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The axial-vector-emitting weak hadronic decays of the Ω_c^0 baryon are investigated. After employing the factorization and the pole model framework to predict their branching ratios, we derive the symmetry-breaking effects on axial-vector-meson-baryon couplings and effects of flavor dependence on baryon-baryon weak transition amplitudes and, consequently, on their branching ratios. We found that the W-exchange process contributions dominate p-wave-meson emitting decays of the Ω_c^0 baryon.

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I. INTRODUCTION

The production of heavy baryons has always posed experimental challenges, and hence has generated much interest in its studies [1–4]. Many interesting observations by CDF, D0, SELEX, FOCUS, Belle, *BABAR*, CMS, LHCb, etc. [5–14] in the context of mass spectra, lifetimes, and decay rates have been made in resent years. Most recently, the LHCb and CDF collaborations [15–22] have announced more precise measurements of the masses and lifetimes of $(\Xi_c^0, \Xi_c^+, \Lambda_b^{(*)}, \Xi_b^-, \Xi_b^0, \Omega_b^-)$ baryons. Also, LHCb has now identified two new strange-beauty baryonic resonances, denoted by $\Xi_b'^-$ and Ξ_b^{*-} [23], though many doubly and triply heavy states are yet to be confirmed. In the two-body nonleptonic decay sector, the first observation of the $(\Omega_b^- \to \Omega_c^0\pi^-)$ decay process and the measurement of CP asymmetries for $\Lambda_b \to p\pi^-$ and $\Lambda_b \to pK^-$ have been reported by the CDF Collaboration [17,24]. On the other hand, LHCb has reported the first observation of $\Lambda_b \to \Lambda_c^+ D_{(s)}^-$ and $\Lambda_b \to J/\psi p\pi^-$ decays and the measurement of the difference in CP asymmetries between $\Lambda_b \to J/\psi p\pi^-$ and $\Lambda_b \to J/\psi pK^-$ and many other decays involving b-baryons [25–29]. However, little progress has been made in observing decays of the heavy charm meson. All these recent measurements have attracted much needed attention to the heavy baryonic sector.

On the theoretical side, various attempts had been made to investigate weak decays of heavy baryons [30–58]. A number of methods—mainly the current algebra (CA) approach, factorization scheme, pole model technique, nonrelativistic quark model (NRQM), heavy quark effective theory (HQET), framework based on next-to-leading-order QCD improved Hamiltonian, etc.—have been employed. Recent experimental developments have prompted more theoretical efforts in b-baryon decays [58–64]. In all these works, the focus has so far been on s-wave-meson emitting decays of heavy baryons including Ω_c^0 decays. Being heavy, charm and bottom baryons can also emit p-wave mesons. In the past, the p-wave-emitting decays of charm and bottom baryons have been studied using the factorization and pole model approach [65–70]. However, p-wave-emitting decays of the Ω_c^0 baryon remain untouched. The fact that the Ω_c^0 baryon is the heaviest and only doubly strange particle in the charmed baryon sextet that is stable against strong and electromagnetic interactions makes it an interesting candidate for the present analysis. Moreover, study of s-wave-emitting decays of the Ω_c^0 baryon reveals that nonfactorizable W-exchange terms dominate as compared to factorizable contributions [37]. This makes the study of Ω_c^0 decays even more important to understand the mechanism underlying W-exchange processes.

In our previous work [70], we have studied the scalar-meson-emitting decays of bottom baryons employing the pole model. We have shown that such decays can acquire significant pole (W-exchange) contributions to make their branching ratios comparable to s-wave-meson emitting decays. In the present work, we analyze the axial-vector-meson-emitting exclusive nonleptonic decays of the Ω_c^0 baryon. We have already seen that for Ω_c^0 decays the factorization contributions are small in comparison to the pole contributions in the case of s-wave-meson emitting decays. Thus, the factorizable contributions to p-wave-meson emitting decays of the Ω_c^0 baryon are also expected to be suppressed. Therefore, we study weak nonleptonic decays of the Ω_c^0 emitting axial-vector mesons in the factorization and pole model approach. We obtain the factorization contributions using the nonrelativistic quark model (NRQM-based) [33] and heavy quark effective theory (HQET-based) [47] form factors. We employ the effects of symmetry breaking (SB) on strong couplings that may decide crucial pole diagram contributions [50,71]. We use the traditional nonrelativistic approach [72] to evaluate weak matrix elements to obtain flavor-independent pole amplitude contributions at ground-state $\frac{1}{2}^+$ intermediate baryon pole terms. Adding factorizable and pole

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contributions, we predict branching ratios (BRs) of Ω_c^0 decays. Later, we employ the possible flavor dependence via variation of spacial baryon wave function overlap in weak decay amplitude. We find that BRs of all the decay modes are significantly enhanced on the inclusion of flavor-dependent effects.

The article is organized as follows: In Sec. II, we give a general framework including spectroscopy of axial-vector mesons, decay kinematics, and the effective Hamiltonian. Section III deals with weak decay amplitudes, both pole terms and factorization terms, weak transitions, axial-vector-meson-baryon couplings, and baryon to baryon transition matrix elements. Numerical results are given in Sec. IV. We summarize our findings in the last section.

II. GENERAL FRAMEWORK

A. Spectroscopy of axial-vector mesons

Axial-vector-meson spectroscopy has extensively been studied in the literature [73–77]. Here, we list the important facts. Spectroscopically, there are two types of axial-vector mesons: ${}^3P_1(J^{PC}=1^{++})$ and ${}^1P_1(J^{PC}=1^{+-})$. ${}^3P_1$ and ${}^1P_1$ states can mix either within themselves or with one another. Experimentally observed nonstrange and uncharmed axial-vector mesons exhibit the first kind of mixing and can be identified as follows:
The $[{}^3P_1]$ meson 16-plet includes the isovector $a_1(1.230)$ and four isoscalars, namely $f_1(1.285)$, $f_1(1.420)/f_1'(1.512)$, and $\chi_{c1}(3.511)$.[1] The following mixing scheme has been used in isoscalar (1^{++}) mesons:

$$f_1(1.285)=\frac{1}{\sqrt{2}}(u\bar{u}+d\bar{d})\cos\phi_A+(s\bar{s})\sin\phi_A,$$
$$f_1'(1.512)=\frac{1}{\sqrt{2}}(u\bar{u}+d\bar{d})\sin\phi_A-(s\bar{s})\cos\phi_A. \tag{1}$$

The $[{}^1P_1]$ meson multiplet consists of an isovector $b_1(1.229)$ and three isoscalars $h_1(1.170)$, $h_1'(1.380)$, and $h_{c1}(3.526)$, where the spin and parities of the $h_{c1}(3.526)$ and $h_1'(1.380)$ states are yet to be confirmed experimentally. These isoscalar (1^{+-}) mesons can mix in the following manner:

$$h_1(1.170)=\frac{1}{\sqrt{2}}(u\bar{u}+d\bar{d})\cos\phi_{A'}+(s\bar{s})\sin\phi_{A'},$$
$$h_1'(1.380)=\frac{1}{\sqrt{2}}(u\bar{u}+d\bar{d})\sin\phi_{A'}-(s\bar{s})\cos\phi_{A'}. \tag{2}$$

The mixing angles are given by the relation $\phi_{A(A')}=\theta(\text{ideal})-\theta_{A(A')}(\text{physical})$. The experimental observations predominantly favor the ideal mixing for these states, i.e., $\phi_A=\phi_{A'}=0°$.

[1]Here the quantities in parentheses indicate their respective masses (in GeV).

The hidden-flavor diagonal states $a_1(1.230)$ and $b_1(1.229)$ cannot mix owing to C- and G-parity considerations. However, there are no such restrictions for the states involving strange partners—namely, K_{1A} and $K_{1A'}$ of the $A(1^{++})$ and $A'(1^{+-})$ mesons, respectively. They mix in the following convention to generate the physical states:

$$K_1(1.270)=K_{1A}\sin\theta_{K_1}+K_{1A'}\cos\theta_{K_1},$$
$$\underline{K}_1(1.400)=K_{1A}\cos\theta_{K_1}-K_{1A'}\sin\theta_{K_1}. \tag{3}$$

Several phenomenological analyses based on the experimental information obtained twofold ambiguous solutions for θ_{K_1}, i.e., $\pm 37°$ and $\pm 58°$ [73–77]. We wish to point out that the experimental measurement of the ratio of $K_1\gamma$ production in B decays and the study of charm meson decays to $K_1(1.270)\pi/K_1(1.400)\pi$ favor negative-angle solutions. Very recently [75], it has been shown that choice of mixing angle θ_{K_1} is intimately related to the choice of angle for f-f' and h-h' mixing schemes. The mixing angle $\theta_{K_1}\sim 35°$ is favored over $\sim 55°$ for near ideal mixing for f-f' and h-h'. Therefore, we use $\theta_{K_1}=-37°$ for our calculation; however, we also give results on $-58°$ for comparison.

B. Kinematics

The matrix element for the baryon decay process, e.g., $B_i(\frac{1}{2}^+,p_i)\to B_f(\frac{1}{2}^+,p_f)+A_k(1^+,q)$, can be expressed as

$$\begin{aligned}\langle B_f(p_f)A_k(q)|H_W|B_i(p_i)\rangle &= i\bar{u}_{B_f}(p_f)\varepsilon^{*\mu}(A_1\gamma_\mu\gamma_5+A_2p_{f\mu}\gamma_5\\ &\quad+B_1\gamma_\mu+B_2p_{f\mu})u_{B_i}(p_i),\end{aligned} \tag{4}$$

where u_B are Dirac spinors for baryonic states B_i and B_f. ε^μ is the polarization vector of the axial-vector-meson state A_k. A_i's and B_i's represent parity-conserving (PC) and parity-violating (PV) amplitudes, respectively. The decay width for the above process is given by

$$\Gamma=\frac{q_\mu}{8\pi}\frac{E_f+m_f}{m_i}\left[2(|S|^2+|P_2|^2)+\frac{E_A^2}{m_A^2}(|S+D|^2+|P_1|^2)\right], \tag{5}$$

where m_i and m_f are the masses of the initial- and final-state baryons, and $q_\mu=(p_i-p_f)_\mu$ is the four-momentum of the axial-vector meson

$$|q_\mu|=\frac{1}{2m_i}\sqrt{[m_i^2-(m_f-m_A)^2][m_i^2-(m_f+m_A)^2]},$$

where m_A is the mass of the emitted p-wave meson [34,41]. The decay amplitude of the final state is now an admixture of S-, P-, and D-wave angular momentum states with

$$S = -A_1, \qquad P_1 = -\frac{q_\mu}{E_A}\left(\frac{m_i + m_f}{E_f + m_f}B_1 + m_i B_2\right),$$

$$P_2 = \frac{q_\mu}{E_f + m_f}B_1, \qquad D = -\frac{q_\mu^2}{E_A(E_f + m_f)}(A_1 - m_i A_2),$$

where E_A and E_f are the energies of the axial-vector meson and the daughter baryon, respectively. Furthermore, there are two independent P-wave amplitudes: one corresponds to the singlet spin combination of the parent and daughter baryon, and the other corresponds to the triplet. The interference between S- and D-wave amplitudes and P-wave amplitudes results in asymmetries for the daughter state with respect to the spin of the parent state. The corresponding asymmetry parameter is

$$\alpha = \frac{4m_A^2 \mathrm{Re}[S * P_2] + 2E_A^2 \mathrm{Re}[(S + D) * P_1]}{2m_A^2(|S|^2 + |P_2|^2) + E_A^2(|S + D|^2 + |P_1|^2)}. \qquad (6)$$

Thus, to determine the decay rate and asymmetry parameters, we are required to estimate amplitudes A and B.

C. Hamiltonian

The QCD modified current $\otimes$ current effective weak Hamiltonian consisting of Cabibbo-favored ($\Delta C = \Delta S = -1$), Cabibbo-suppressed ($\Delta C = -1, \Delta S = 0$), and Cabibo-doubly-suppressed ($\Delta C = -\Delta S = -1$) modes is given by

$$\begin{aligned} H_W^{\mathrm{eff}} = \frac{G_F}{\sqrt{2}}\{&V_{ud}V_{cs}^*[c_1(\bar{u}d)(\bar{c}s) + c_2(\bar{s}d)(\bar{u}c)]_{(\Delta C=\Delta S=-1)} \\ &+ V_{ud}V_{cd}^*[c_1\{(\bar{s}c)(\bar{u}s) - (\bar{d}c)(\bar{u}d)\} + c_2\{(\bar{u}c)(\bar{s}s) - (\bar{u}c)(\bar{d}d)\}]_{(\Delta C=-1,\Delta S=0)} \\ &- V_{us}V_{cd}^*[c_1(\bar{d}c)(\bar{u}s) + c_2(\bar{u}c)(\bar{d}s)]_{(\Delta C=-\Delta S=-1)}\}, \end{aligned} \qquad (7)$$

where V_{ij} and $(\bar{q}_i q_j) \equiv \bar{q}_i\gamma_\mu(1-\gamma_5)q_j$ denote the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements and the weak V-A current, respectively. We use the QCD coefficients $c_1(\mu) = 1.2$, $c_2(\mu) = -0.51$ at $\mu \approx m_c^2$ [35]. Furthermore, nonfactorizable effects may modify c_1 and c_2, thereby indicating that these may be treated as free parameters. The discrepancy between theory and experiment is greatly improved in the large-N_c limit. Interestingly, the charm-conserving decays of Ω_c^0 are also possible, but they are kinematically forbidden in the present analysis.

III. DECAY AMPLITUDES

The hadronic matrix element $\langle B_f A_k|H_W|B_i\rangle$ for the $B_i \to B_f + A_k$ process may be expressed as

$$\langle B_f A_k|H_W|B_i\rangle \equiv \mathcal{A}_{\mathrm{Pole}} + \mathcal{A}_{\mathrm{Fac}}, \qquad (8)$$

where $\mathcal{A}_{\mathrm{Pole}}$ and $\mathcal{A}_{\mathrm{Fac}}$ represent the pole (W-exchange) and factorization amplitudes, respectively. The pole diagram contributions involving the W-exchange process are evaluated using the pole model framework [35]. One may consider factorization as a correction to the pole model, which includes the calculation of possible pole diagrams via the s, u, and t channels, where the t channel virtually implicates tree-level diagrams, i.e., factorizable processes. The contribution of these terms can be summed up in terms of PC and PV amplitudes.

A. Pole amplitudes

The first term, $\mathcal{A}_{\mathrm{Pole}}$, involves the evaluation of the relevant matrix element

$$\langle B_f|H|B_i\rangle = \bar{u}_{B_i}(B + \gamma_5 A)u_{B_j} \qquad (9)$$

between two $\frac{1}{2}^+$ baryon states. A and B are s-wave and p-wave decay amplitudes, respectively. A and B include the contributions of s and u channels for positive-parity intermediate baryon ($J^P = \frac{1}{2}^+$) poles, henceforth given by A^{pole} and B^{pole} as follows:

$$A^{\mathrm{pole}} = -\sum_n\left[\frac{g_{B_f B_n A_k}a_{ni}}{m_i - m_n} + \frac{a_{fn}g_{B_n B_i A_k}}{m_f - m_n}\right], \qquad (10)$$

$$B^{\mathrm{pole}} = \sum_n\left[\frac{g_{B_f B_n A_k}b_{ni}}{m_i + m_n} + \frac{b_{fn}g_{B_n B_i A_k}}{m_f + m_n}\right], \qquad (11)$$

where g_{ijk} are the strong axial-vector-meson-baryon coupling constants; a_{ij} and b_{ij} are weak baryon-baryon matrix elements defined as

$$\langle B_i|H_W|B_j\rangle = \bar{u}_{B_i}(a_{ij} + \gamma_5 b_{ij})u_{B_j}. \qquad (12)$$

It is well known that the PV matrix element b_{ij} vanishes for the hyperons owing to $\langle B_f A_k|H_W^{\mathrm{PV}}|B_i\rangle = 0$ in the SU (3) limit. This also implies for nonleptonic charm meson decays that $b_{ij} \ll a_{ij}$, suppressing s-wave contributions for $\frac{1}{2}^+$ poles. These contributions are further suppressed by presence of sum of the baryon masses in the denominator. Thus, only PC terms survive for nonleptonic decays of charm baryons. It may be noted that the negative-parity intermediate baryon $(J^P = \frac{1}{2}^-)$ may also contribute to these processes and may turn out to be important. However, evaluation of such terms requires knowledge of the axial-vector-meson strong coupling constants for $(\frac{1}{2}^-)$ states. Unfortunately, no such theoretical or experimental information is available in literature. Moreover, *in the leading nonrelativistic approximation*, one can ignore $J^P = \frac{1}{2}^-, \frac{3}{2}^-, \ldots$ and higher resonances, as they would require at least one power of momentum in H_W in order to connect them with the relevant ground state in the overlap integral. That means that one must consider terms of the order v/c. In the same manner, to connect radial excitations with the corresponding ground state, one would need terms of order $(v/c)^2$; otherwise, the overlap integral would be zero due to orthogonality of the wave functions [51]. Therefore, we have restricted our calculation to ground-state $\frac{1}{2}^+$ intermediate baryon pole terms to estimate the pole contributions to the axial-vector-meson-emitting decays of charm baryons.

B. Factorizable amplitudes

Likewise for meson decays, the reduced matrix element (4) can be factorized to obtain decay amplitudes (ignoring the scale factors) in the following form:

$$\langle A_k(q)|A_\mu|0\rangle\langle B_f(p_f)|V^\mu + A^\mu|B_i(p_i)\rangle, \tag{13}$$

where

$$\langle A_k(q)|A_\mu|0\rangle = f_A m_A \varepsilon^*_\mu, \tag{14}$$

and f_A is the decay constant of the emitted axial-vector meson A_k. The baryon-baryon matrix elements of the weak currents are defined as

$$\langle B_f(p_f)|V_\mu|B_i(p_i)\rangle = \bar{u}_f(p_f)\left[f_1\gamma_\mu - \frac{f_2}{m_i}i\sigma_{\mu\nu}q^\nu + \frac{f_3}{m_i}q_\mu\right]u_i(p_i) \tag{15}$$

and

$$\langle B_f(p_f)|A_\mu|B(p_i)\rangle = \bar{u}_f(p_f)\left[g_1\gamma_\mu\gamma_5 - \frac{g_2}{m_i}i\sigma_{\mu\nu}q^\nu\gamma_5 + \frac{g_3}{m_i}q_\mu\gamma_5\right]u_i(p_i), \tag{16}$$

where, f_i and g_i denote the vector and axial-vector form factors as functions of q^2 [35]. The factorizable amplitudes are thus given by

$$\begin{aligned}
A_1^{\rm fac} &= -\frac{G_F}{\sqrt{2}}F_C f_A c_k m_A\left[g_1^{B_i,B_f}(m_A^2) - g_2^{B_i,B_f}(m_A^2)\frac{m_i - m_f}{m_i}\right],\\
A_2^{\rm fac} &= \frac{G_F}{\sqrt{2}}F_C f_A c_k m_A[2g_2^{B_i,B_f}(m_A^2)/m_i],\\
B_1^{\rm fac} &= \frac{G_F}{\sqrt{2}}F_C f_A c_k m_A\left[f_1^{B_i,B_f}(m_A^2) + f_2^{B_i,B_f}(m_A^2)\frac{m_i + m_f}{m_i}\right],\\
B_2^{\rm fac} &= -\frac{G_F}{\sqrt{2}}F_C f_A c_k m_A[2f_2^{B_i,B_f}(m_A^2)/m_i],
\end{aligned} \tag{17}$$

where F_C contains appropriate CKM factors and Clebsch-Gordan (CG) coefficients, and c_k are QCD coefficients.

The baryon-baryon transition form factors f_i and g_i are evaluated in literature using the nonrelativistic quark model (NRQM) [33] and the HQET [47]. The NRQM-based form factors are calculated in the Breit frame and include corrections like the q^2-dependence of the form factors, the hard gluon QCD contributions, and the wave function mismatch. Later, in light of the fact that the form factors for heavy baryon-baryon transitions should also include constraints from the heavy quark symmetry, $1/m_Q$ correction to the form factors was introduced using HQET. It may be noted that the Ω_c decays involving factorizable amplitudes only include $\Omega_c^0 \to \Xi^0$ and $\Omega_c^0 \to \Xi^-$ form factors which come out to be equal, numerically. The evaluated form factors using NRQM [33] are given by

$$\begin{aligned}
\Omega_c \to \Xi:\ & f_1(0) = -0.23, \qquad f_2(0) = 0.21,\\
& g_1(0) = 0.14, \qquad g_2(0) = -0.019.
\end{aligned} \tag{18}$$

Similarly, form factor calculation in HQET [47] yields

$$\begin{aligned}
\Omega_c \to \Xi:\ & f_1(0) = -0.34, \qquad f_2(0) = 0.35,\\
& g_1(0) = 0.10, \qquad g_2(0) = -0.020.
\end{aligned} \tag{19}$$

We wish to remark here that for numerical calculations of the factorizable branching ratios, we use dipole q^2-dependence following HQET constraints.

The decay constants of axial-vector mesons [73–77] used for numerical evaluations are given by

$$f_{a_1} \approx f_{f_1} = 0.221 \text{ GeV}, \qquad f_{K_1}(1270) = 0.175 \text{ GeV},$$

while the decay constant for $\underline{K}_1(1.400)$ may be calculated by using the relation $f_{\underline{K}_1}(1.400)/f_{K_1}(1.270) = \cot\theta_1$, i.e.,

$$f_{\underline{K}_1}(1.400) = -0.099 \text{ GeV} \quad \text{for } \theta_1 = -58°;$$
$$f_{\underline{K}_1}(1.400) = -0.225 \text{ GeV} \quad \text{for } \theta_1 = -37°$$

It may also be noted that decay constants of axial-vector mesons are not so trivial to understand, as these may be effected by factors like C-parity/G-parity conservations, mixing scheme, and SU(3) breaking, etc. For more details, readers are referred to Refs. [73–77].

C. Weak transitions

The flavor-symmetric weak Hamiltonian [40,45] for the quark-level process $q_i + q_j \to q_l + q_m$ can be expressed as

$$H_W \cong V_{il}V^*_{jm}c_-(m_c)[\bar{B}^{[i,j]k}B_{[l,m]k}H^{[l,m]}_{[i,j]}], \tag{20}$$

where $c_- = c_1 + c_2$, and the brackets [ ] represent the antisymmetrization among the indices. The spurion transforms like $H^{[1,3]}_{[2,4]}$. We obtain the weak baryon-baryon matrix elements a_{ij} for CKM-favored and CKM-suppressed modes from the following contraction:

$$H_W \cong a_W[\bar{B}^{[i,j]k}B_{[l,m]k}H^{[l,m]}_{[i,j]}], \tag{21}$$

where a_W is weak amplitude. It is worth remarking here that the enhancement due to hard gluon exchanges, coming through c_-, will affect the weak baryon-baryon matrix element. Also, we ignore the long-distance QCD effects reflected in the bound-state wave functions.

D. Axial-vector-meson-baryon couplings

In SU(4), the Hamiltonian representing the strong transitions is given by

$$H_{\text{strong}} = \sqrt{2}(g_D + g_F)\left(\frac{1}{2}\bar{B}^{[a,b]d}B_{[a,b]c}A^c_d\right) + \sqrt{2}(g_D - g_F)(\bar{B}^{[a,b]d}B_{[a,b]c}A^c_d), \tag{22}$$

where $B_{[a,b]c}$, $\bar{B}^{[a,b]d}$, and A^c_d are the baryon, antibaryon, and axial-vector-meson tensors, respectively, and $g_D(g_F)$ are conventional D-type (F-type) strong coupling constants [30,33,37].

Experimentally, there are no measurements available for the axial-vector-meson-baryon coupling constants for the charm sector. Since it is difficult to determine g_{Nna} directly, a reasonable estimate could be obtained by using the Goldberger-Treiman (GT) relation:

$$g_{NN\pi} = \frac{g_A m_N}{f_\pi}, \tag{23}$$

which relates the pion-nucleon coupling $g_{NN\pi}$ with the axial-vector coupling g_A [78,79]. The GT relation exhibits the direct relation between spontaneous chiral symmetry breaking and the Partially Conserved Axial-vector Current (PCAC) hypothesis at $SU(2)_L \times SU(2)_R$ or, with a possible extension, $SU(3)_L \times SU(3)_R$. Here, g_A represents the contribution to the dispersion relation of all the axial-vector states higher than the pion. In light of the PCAC, the heavier axial vector states contributing to g_A must reproduce the pion pole at $q^2 = 0$. Thus, the combined contribution of all the heavier states may be replaced by an effective pole a_1—i.e., if we assume axial-vector dominance[2] [81,82] to get

$$g_{NNa_1} \approx \frac{g_A m^2_{a_1}}{f_{a_1}} = 8.60 \tag{24}$$

for $g_A = 1.26$ given by β decay [79]. To proceed forward, we use QCD sum rules analysis [79],

$$g_D = 6.15 \quad \text{and} \quad g_F = 2.45 \quad \text{for } \frac{g_D}{g_F} = 2.5, \tag{25}$$

which in turn yields axial-vector-meson-baryon strong coupling constants based on SU(4) symmetry.

We wish to point out that the SU(4) symmetry is badly broken, hence it would not be wise to use SU(4)-symmetry-based strong coupling constants for charm baryon decays. Therefore, we consider the SU(4)-breaking effects in strong coupling constants by using the Coleman-Glashow null result [83]. The tadpole mechanism can generate breaking effects—namely the medium strong, the electromagnetic, and the weak effects—that transform like an SU(3) octet, via a single symmetry-breaking term. Thus, except for the tadpole term, the hadronic Hamiltonian remains SU(3) invariant. In the SU(3) octet, the strangeness-changing scalar tadpole S_6, transforming as the sixth component of the symmetry-breaking octet, can be rotated away by unitary transformation. These strangeness-changing effects produced by the S_6 tadpole must vanish, leaving behind the electromagnetic and the weak effects. Khanna and Verma [71] exploited the null result to obtain SU(3) broken baryon-baryon-pseudoscalar couplings. After validating the SU(3) case, they extended their results to SU(4), where

[2]As a consequence of spontaneous chiral symmetry breaking, Weinberg sum rules [80] relate m_{a_1} and m_ρ by assuming vector and axial-vector dominance.

symmetry-breaking effects belong to the similar regular representation **15**. In SU(4), the weak interaction Hamiltonian responsible for hadronic weak decays of charm baryons belongs to the representation **20**$''$. The tadpole term of the weak Hamiltonian belongs to the representation **15**. In this case the charm-changing effects generated through the S_9 tadpole must vanish (for details, see Ref. [71]). We wish to remark here that the tadpole-type symmetry-breaking effects do not include any additional parameter. The symmetry-broken (SB) baryon-meson strong couplings are calculated by

$$g_A^{BB'}(\mathrm{SB}) = \frac{M_B + M_B'}{2M_N}\frac{1}{\alpha_P} g_A^{BB'}(\mathrm{Sym}), \tag{26}$$

where $g_A^{BB'}$ (Sym) is the SU(4) symmetric couplings, and the ratio of mass breaking terms α_P is given by [71]

$$\alpha_P = \frac{\delta M_c}{\delta M_s} = \sqrt{\frac{3}{8}}\frac{m_c - m_u}{m_s - m_u}.$$

The obtained values of SB strong axial-vector-meson-baryon coupling constants relevant for our calculation have been given in Table I. For heavy baryon decays, it has been observed [84] that mass independent strong couplings lead to smaller pole contributions. It is quite obvious that symmetry breaking will result in larger values for strong couplings as compared to symmetric ones due to mass dependence. Consequently, higher pole contributions would be expected. We also give the expressions for the $g_A^{BB'}$ in terms of g_D and g_F. However, we have given the absolute numerical values for the strong couplings where the actual sign would depend upon the conventions used and could be determined from their expressions in the present case.

TABLE I. Expressions of strong-coupling constants [$\mathrm{SB} = \mathrm{Sym} \times (\frac{M_B+M_B'}{2M_N\alpha_P})$] and their absolute numerical values at $\theta_{K_1} = -37°(-58°)$.

Strong couplings	$g_A^{BB'} \times (\frac{M_B+M_B'}{2M_N\alpha_P})$	Absolute values $\lvert g_A^{BB'}(\mathrm{SB})\rvert$
$g_{K_1}^{\Lambda p}$	$(\sqrt{3}g_D + \frac{g_F}{\sqrt{3}})\sin\theta_{K_1}$	5.13 (7.23)
$g_{K_1}^{\Sigma^0 p}$	$(-g_D + g_F)\sin\theta_{K_1}$	2.53 (3.56)
$g_{\underline{K}_1}^{\Lambda p}$	$(\sqrt{3}g_D + \frac{g_F}{\sqrt{3}})\cos\theta_{K_1}$	6.81 (4.52)
$g_{\underline{K}_1}^{\Sigma^0 p}$	$(-g_D + g_F)\cos\theta_{K_1}$	3.35 (2.23)
$g_{K_1}^{\Lambda n}$	$-(\sqrt{3}g_D + \frac{g_F}{\sqrt{3}})\sin\theta_{K_1}$	5.14 (7.24)
$g_{K_1}^{\Sigma^0 n}$	$(-g_D + g_F)\sin\theta_{K_1}$	2.53 (3.57)
$g_{\underline{K}_1}^{\Lambda n}$	$-(\sqrt{3}g_D + \frac{g_F}{\sqrt{3}})\cos\theta_{K_1}$	6.82 (4.52)
$g_{\underline{K}_1}^{\Sigma^0 n}$	$(-g_D + g_F)\cos\theta_{K_1}$	3.58 (2.23)
$g_{K_1}^{\Xi^0 \Lambda}$	$(-\sqrt{3}g_D + \frac{g_F}{\sqrt{3}})\sin\theta_{K_1}$	0.54 (0.76)
$g_{K_1}^{\Xi^- \Lambda}$	$(\sqrt{3}g_D - \frac{g_F}{\sqrt{3}})\sin\theta_{K_1}$	0.54 (0.76)
$g_{\underline{K}_1}^{\Xi^0 \Lambda}$	$(-\sqrt{3}g_D + \frac{g_F}{\sqrt{3}})\cos\theta_{K_1}$	0.72 (0.47)
$g_{K_1}^{\Sigma^0 \Xi^{0(-)}}$	$-(g_D + g_F)\sin\theta_{K_1}$	6.92 (9.75)
$g_{\underline{K}_1}^{\Sigma^0 \Xi^{0(-)}}$	$-(g_D + g_F)\cos\theta_{K_1}$	9.18 (6.10)
$g_{K_1}^{\Sigma^+ \Xi^0}$	$-\sqrt{2}(g_D + g_F)\sin\theta_{K_1}$	9.77 (13.76)
$g_{\underline{K}_1}^{\Sigma^+ \Xi^0}$	$-\sqrt{2}(g_D + g_F)\cos\theta_{K_1}$	12.96 (8.60)
$g_{K_1}^{\Omega_c \Xi_c}$	$\frac{-2g_F}{\sqrt{3}}\sin\theta_{K_1}$	11.77 (16.58)
$g_{K_1}^{\Omega_c \Xi_c'}$	$2g_D\sin\theta_{K_1}$	8.27 (11.72)
$g_{\underline{K}_1}^{\Omega_c \Xi_c}$	$\frac{-2g_F}{\sqrt{3}}\cos\theta_{K_1}$	15.62 (10.34)
$g_{\underline{K}_1}^{\Omega_c \Xi_c'}$	$2g_D\cos\theta_{K_1}$	11.00 (7.30)
$g_{f_1}^{\Lambda\Lambda}$	$2(g_D - \frac{2g_F}{3})$	3.92
$g_{f_1'}^{\Lambda\Lambda}$	$-\sqrt{2}(g_D + \frac{g_F}{3})$	7.57
$g_{f_1/f_1'}^{\Sigma^0\Lambda}$	0	0
$g_{a_1}^{\Lambda\Sigma^+}$	$\frac{2g_F}{\sqrt{3}}$	8.72
$g_{a_1}^{\Lambda\Sigma^{0(-)}}$	$-\frac{2g_F}{\sqrt{3}}$	8.74
$g_{a_1}^{\Sigma^0\Sigma^0}$	0	0
$g_{a_1}^{\Sigma^0\Sigma^{+(-)}}$	$2g_D$	6.22
$g_{a_1/f_1}^{\Xi^0\Xi^0}$	$g_D - g_F$	5.18
$g_{a_1^+}^{\Xi^0\Xi^-}$	$\sqrt{2}(g_D - g_F)$	7.35
$g_{a_1/f_1}^{\Omega_c\Omega_c}$	0	0
$g_{f_1'}^{\Omega_c\Omega_c}$	$-2\sqrt{2}g_D$	19.92

E. Baryon matrix element

In general, numerical evaluation of W-exchange terms (pole terms) involves a weak matrix element of the form $\langle B_f|H_W^{\mathrm{PC}}|B_i\rangle$. Riazuddin and Fayyazuddin [72] have calculated this matrix element for noncharmed hyperon decays in the nonrelativistic limit. Following their analysis, one can obtain the matrix element for the charm baryons as a *first approximation*. Though Ω_c is heavy and, therefore, outgoing quarks may have large momenta, we use the nonrelativistic approximation to get the first estimates of the baryonic matrix elements. The matrix element for the W-exchange process ($c + d \to s + u$) can be expressed as

$$\mathcal{M} \approx \frac{G_F}{\sqrt{2}} V_{du}V_{cs}[\bar{\psi}_u(p_i')\gamma_\mu(1-\gamma_5)\alpha_i^+\psi_d(p_i)\bar{\psi}_s(p_j') \times \gamma_\mu(1-\gamma_5)\gamma_j^-\psi_c(p_j) + i\leftrightarrow j], \tag{27}$$

where ψ's are Dirac spinors and $q = p_i - p_i' = p_j' - p_j$. The operators α_i^+ convert $d \to u$ and γ_j^- convert $c \to s$. In the leading nonrelativistic approximation, only terms corresponding to γ^0 and $\gamma^i\gamma_5$ have nonzero limits, which are then reduced to the only parity-conserving part of $\mathcal{M}$. Thus, in leading nonrelativistic approximation we have

$$\mathcal{M}^{\mathrm{PC}} = \frac{G_F}{\sqrt{2}} V_{du}V_{cs}\sum_{i>j}(\gamma_i^-\alpha_j^+ + \alpha_i^+\gamma_j^-)(1 - \sigma_i\cdot\sigma_j), \tag{28}$$

where $\mathbf{S}_i = \sigma_i/2$ are Pauli spinors representing the spin of ith quark. Fourier transformation of the above expression gives the parity-conserving weak Hamiltonian

$$H_W^{\rm PC} = \frac{G_F}{\sqrt{2}} V_{du} V_{cs} \sum_{i\neq j} \alpha_i^+ \gamma_j^- (1 - \sigma_i \cdot \sigma_j) \delta^3(r), \quad (29)$$

following which we can get a reasonable estimate of these terms. One can fix the scale by assuming the baryon overlap wave function to be flavor independent such that

$$\langle \psi_f | \delta^3(r) | \psi_i \rangle_c \approx \langle \psi_f | \delta^3(r) | \psi_i \rangle_s, \quad (30)$$

where $\langle \psi_f | \delta^3(r) | \psi_i \rangle$ is the baryon wave function overlap for corresponding flavor. We wish to remark here that ([30]) leads to a well-known SU(4)-based relation that connects nonleptonic charmed baryon decays with hyperon decays. Since SU(4) is badly broken, a large mismatch between charm and strange baryon wave function overlaps would need a correction factor that has been practiced in many models based on different arguments [39,44].

IV. NUMERICAL RESULTS AND DISCUSSIONS

Summing over all the ingredients, the factorization and the pole contributions to different PV and PC amplitudes has been calculated. The numerical values of the possible factorizable contributions to weak decay amplitudes of the Ω_c baryon in CKM-favored, suppressed, and doubly suppressed modes are shown in Table II. Using the symmetry-broken axial-vector-meson-baryon couplings, we obtain the flavor-independent pole amplitudes for all Ω_c baryon decays in CKM-favored, suppressed, and doubly suppressed modes, as shown in column 2 of Table III. We wish to emphasize that we have used only ground-state $\frac{1}{2}^+$ intermediate baryon pole terms to estimate pole contributions. Adding factorizable and pole contributions, the BRs and asymmetry parameters for the flavor-independent case are predicted, as shown in columns 3 and 5 of Table IV. Since factorization contributes to only six of the Ω_c decay modes, the remaining decay modes acquire contributions from pole amplitudes only. These branching ratios are given in column 2 of Table V. We wish to point out that we use $\theta_{K_1} = -37°$ as the reference mixing angle; however, we also give results on mixing angle $-58°$ for comparison. We summarize our results as follows:

(1) The branching ratios of all the decay channels range from 10^{-3} to 10^{-7}. The branching ratios of dominant modes are $\mathcal{O}(10^{-3}) \sim \mathcal{O}(10^{-4})$.
(2) Most of the observed Ω_c decay channels come from W-exchange processes; however, factorization processes contribute to only six of the decay channels.
(3) The factorization contributions obtained from the nonrelativistic quark model (NRQM) and heavy quark effective theory (HQET) compare well without much discrepancy.
(4) Among the Ω_c decays acquiring contributions from both factorizable and pole amplitudes, the only possible CKM-favored ($\Delta C = \Delta S = -1$) decay mode $\Omega_c^0 \to \Xi^0 \bar{K}_1^0$ has a largest branching ratio of 1.15×10^{-3} for $\theta_{K_1} = -37°$. The branching ratio increases further to a value 1.56×10^{-3} for $\theta_{K_1} = -58°$.

TABLE II. Factorizable amplitudes (in units of $\frac{G_F}{\sqrt{2}} V_{uq} V_{cq}^*$) to Ω_c^0 decays for CKM-favored, CKM-suppressed, and CKM-doubly-suppressed modes.

Decays	Model [33,47]	Factorizable amplitudes[a] $A_1^{\rm fac}$	$A_2^{\rm fac}$	$B_1^{\rm fac}$	$B_2^{\rm fac}$
Cabbibo-favored $\Delta C = -1$, $\Delta S = 0$ mode:					
$\Omega_c^0 \to \Xi^0 \bar{K}_1^0$	NRQM	0.033	0.0017	−0.0027	0.052
	HQET	0.025	0.0034	−0.0060	0.087
CKM-suppressed $\Delta C = -1$, $\Delta S = 0$ mode:					
$\Omega_c^0 \to \Xi^0 a_1^0$	NRQM	−0.026	−0.0013	0.021	−0.039
	HQET	−0.018	−0.0026	0.045	−0.065
$\Omega_c^0 \to \Xi^0 f_1$	NRQM	0.029	0.0015	−0.024	0.045
	HQET	0.022	0.0029	−0.052	0.075
$\Omega_c^0 \to \Xi^- a_1^+$	NRQM	−0.090	−0.0046	0.072	−0.14
	HQET	−0.068	−0.0092	0.160	−0.23
CKM-doubly-suppressed $\Delta C = -\Delta S = -1$ mode:					
$\Omega_c^0 \to \Xi^0 K_1^0$	NRQM	0.033	0.017	−0.027	0.052
	HQET	0.025	0.034	−0.062	0.087
$\Omega_c^0 \to \Xi^- K_1^+$	NRQM	−0.082	−0.0042	0.068	−0.128
	HQET	−0.062	−0.0083	0.149	−0.214

[a]The factorizable amplitudes are independent of mixing angle θ_{K_1} for the decays emitting the $K_1(1.270)$ meson, since the decay constant of $K_1(1.270)$ does not depend upon the $K_1(1.270) - \underline{K}_1(1.400)$ mixing angle, which essentially affects the decay constant of $\underline{K}_1(1.400)$.

TABLE III. Pole amplitudes (in units of $\frac{G_F}{\sqrt{2}}V_{uq}V_{cq}^*$) of all Ω_c^0 decays for CKM-favored, CKM-suppressed and CKM- doubly-suppressed modes at $\theta_{K_1} = -37°(-58°)$. Flavor dependent pole contributions include effects of $|\psi(0)|^2$ variation.

	Pole amplitudes	
Decays	Flavor independent	Flavor dependent
CKM-favored ($\Delta C = \Delta S = -1$) mode:		
$\Omega_c^0 \to \Xi^0 \bar{K}_1^0$	−0.026(−0.036)	−0.054(−0.076)
CKM-suppressed ($\Delta C = -1, \Delta S = 0$) mode:		
$\Omega_c^0 \to \Xi^0 a_1^0$	−0.20	−0.42
$\Omega_c^0 \to \Xi^0 f_1$	−0.20	−0.42
$\Omega_c^0 \to \Xi^- a_1^+$	−0.28	−0.59
$\Omega_c^0 \to \Lambda \bar{K}_1^0$	0.12 (0.17)	0.27 (0.36)
$\Omega_c^0 \to \Lambda \underline{\bar{K}}_1^0$	−0.16(−0.11)	−0.34(−0.23)
$\Omega_c^0 \to \Sigma^+ K_1^-$	0.049 (0.069)	0.10 (0.14)
$\Omega_c^0 \to \Sigma^+ \underline{K}_1^-$	−0.065(−0.043)	−0.14(−0.090)
$\Omega_c^0 \to \Sigma^0 \bar{K}_1^0$	0.034 (0.048)	0.072 (0.10)
$\Omega_c^0 \to \Sigma^0 \underline{\bar{K}}_1^0$	−0.046(−0.030)	−0.096(−0.064)
CKM-doubly-suppressed ($\Delta C = -\Delta S = -1$) mode:		
$\Omega_c^0 \to \Xi^0 K_1^0$	0.015 (0.021)	0.031 (0.044)
$\Omega_c^0 \to \Xi^- K_1^+$	−0.015(−0.021)	−0.031(−0.044)
$\Omega_c^0 \to p K_1^-$	0.22 (0.30)	0.45 (0.64)
$\Omega_c^0 \to p \underline{K}_1^-$	−0.29(−0.19)	−0.60(−0.40)
$\Omega_c^0 \to n K_1^0$	−0.22(−0.30)	−0.45(−0.64)
$\Omega_c^0 \to n \underline{K}_1^0$	0.29 (0.19)	0.60 (0.40)
$\Omega_c^0 \to \Lambda f_1$	−0.11	−0.22
$\Omega_c^0 \to \Lambda f_1'$	0.34	0.71
$\Omega_c^0 \to \Sigma^+ a_1^-$	0.24	0.50
$\Omega_c^0 \to \Sigma^0 a_1^0$	−0.24	−0.50
$\Omega_c^0 \to \Sigma^- a_1^+$	−0.24	−0.50

(5) For $\Omega_c^0 \to \Xi^0 \bar{K}_1^0$ decay, we find that the dominant contribution comes from factorizable amplitudes with pole contributions as low as ~20%-25%. It may be noted that color-suppressed factorizable amplitude interferes constructively with pole amplitude, resulting in a large branching ratio.

(6) In CKM-suppressed ($\Delta C = -1, \Delta S = 0$) mode, the most dominant decay has $\mathrm{Br}(\Omega_c^0 \to \Xi^- a_1^+) \sim 1.00 \times 10^{-3}$ in HQET (though all the decays in this mode occur at comparable footing). We wish to point out that despite the CKM suppression and destructive interference betwen pole and factorization contributions, the $\Omega_c^0 \to \Xi^- a_1^-$ decay is overly compensated by QCD enhancement (c_1).

(7) The next-order dominant decays in CKM-suppressed mode are $\Omega_c^0 \to \Xi^0 a_1^0/\Xi^0 f_1$ with roughly comparable branching ratios. Here also, factorization and pole terms interfere constructively and destructively for the decays involving f_1 and a_1, respectively, though both are suppressed due to color suppression and small CG coefficients. It may be noted that $\Omega_c^0 \to \Xi^- a_1^+/\Xi^0 a_1^0/\Xi^0 f_1$ decays have dominant pole term contributions as compared to factorizable term contributions ~(20%-40%).

(8) As expected, the decay channels in Cabibbo doubly suppressed ($\Delta C = \Delta S = -1$) modes have relatively smaller BRs of $\mathcal{O}(10^{-6}\text{-}10^{-7})$. We observe an increase in the branching ratio of the color-suppressed $\Omega_c \to \Xi^0 K_1^0$ decay despite the expected destructive interference between pole and factorization terms. We wish to remark here that the change of angle θ_{K_1} to −58° leads to smaller BR for $\Omega_c \to \Xi^0 K_1^0$, though it increases for the HQET case.

TABLE IV. Branching ratios and asymmetry parameters of Ω_c^0 decays acquiring contributions from both factorization and pole amplitudes for CKM-favored, CKM-suppressed and CKM-doubly-suppressed modes at $\theta_{K_1} = -37°(-58°)$. Flavor-dependent results include effects of $|\psi(0)|^2$ variation.

		Branching ratios		Asymmetry 'α'	
Decays	Model [33,47]	Flavor independent	Flavor dependent	Flavor independent	Flavor dependent
CKM-favored $\Delta C = -1$, $\Delta S = 0$ mode:					
$\Omega_c^0 \to \Xi^0 \bar{K}_1^0$	NRQM	$1.15 \times 10^{-3}(1.56 \times 10^{-3})$	$2.42 \times 10^{-3}(3.78 \times 10^{-3})$	0.39 (0.34)	0.28 (0.22)
	HQET	$0.98 \times 10^{-3}(1.34 \times 10^{-3})$	$2.11 \times 10^{-3}(3.37 \times 10^{-3})$	0.63 (0.55)	0.46 (0.36)
CKM-suppressed $\Delta C = -1$, $\Delta S = 0$ mode:					
$\Omega_c^0 \to \Xi^0 a_1^0$	NRQM	5.90×10^{-4}	3.00×10^{-3}	−0.13	−0.058
	HQET	6.40×10^{-4}	3.11×10^{-3}	−0.20	−0.091
$\Omega_c^0 \to \Xi^0 f_1$	NRQM	7.96×10^{-4}	3.05×10^{-3}	0.091	0.045
	HQET	7.51×10^{-4}	2.96×10^{-3}	0.15	0.076
$\Omega_c^0 \to \Xi^- a_1^+$	NRQM	7.58×10^{-4}	4.88×10^{-3}	−0.38	−0.15
	HQET	9.90×10^{-4}	5.37×10^{-3}	−0.51	−0.23
CKM-doubly-suppressed $\Delta C = -\Delta S = -1$ mode:					
$\Omega_c^0 \to \Xi^0 K_1^0$	NRQM	$5.15 \times 10^{-7}(3.42 \times 10^{-7})$	$2.41 \times 10^{-7}(3.40 \times 10^{-7})$	0.74 (0.69)	0.16(−0.70)
	HQET	$6.10 \times 10^{-7}(5.16 \times 10^{-7})$	$5.54 \times 10^{-7}(8.15 \times 10^{-7})$	0.55 (0.23)	−0.38(−0.78)
$\Omega_c^0 \to \Xi^- K_1^+$	NRQM	$3.40 \times 10^{-6}(2.35 \times 10^{-6})$	$4.93 \times 10^{-6}(4.13 \times 10^{-6})$	0.66 (0.69)	0.72 (0.74)
	HQET	$4.70 \times 10^{-6}(4.10 \times 10^{-6})$	$3.67 \times 10^{-6}(3.04 \times 10^{-6})$	0.77 (0.73)	0.64 (0.41)

TABLE V. Branching ratios of Ω_c^0 for CKM-suppressed and CKM-doubly-suppressed modes at $\theta_{K_1} = -37°(-58°)$ acquiring contributions from pole amplitudes only. Flavor-dependent branching ratios include effects of $|\psi(0)|^2$ variation.

Decays	Flavor-independent BRs	Flavor-dependent BRs
CKM-suppressed ($\Delta C = -1, \Delta S = 0$) mode:		
$\Omega_c^0 \to \Lambda \bar{K}_1^0$	$3.96 \times 10^{-4}(7.89 \times 10^{-4})$	$1.74 \times 10^{-3}\,(3.47 \times 10^{-3})$
$\Omega_c^0 \to \Lambda \underline{\bar{K}}_1^0$	$4.83 \times 10^{-4}(2.12 \times 10^{-4})$	$2.13 \times 10^{-3}(9.37 \times 10^{-4})$
$\Omega_c^0 \to \Sigma^+ K_1^-$	$5.50 \times 10^{-5}(1.09 \times 10^{-4})$	$2.42 \times 10^{-4}(4.82 \times 10^{-4})$
$\Omega_c^0 \to \Sigma^+ \underline{K}_1^-$	$5.94 \times 10^{-5}(2.62 \times 10^{-4})$	$2.62 \times 10^{-4}(1.15 \times 10^{-4})$
$\Omega_c^0 \to \Sigma^0 \bar{K}_1^0$	$2.72 \times 10^{-5}(5.41 \times 10^{-5})$	$1.20 \times 10^{-4}(2.39 \times 10^{-4})$
$\Omega_c^0 \to \Sigma^0 \underline{\bar{K}}_1^0$	$2.92 \times 10^{-5}(1.28 \times 10^{-5})$	$1.29 \times 10^{-4}(5.67 \times 10^{-5})$
CKM-doubly-suppressed ($\Delta C = -\Delta S = -1$) mode:		
$\Omega_c^0 \to p K_1^-$	$7.75 \times 10^{-5}(1.54 \times 10^{-4})$	$3.42 \times 10^{-4}(6.79 \times 10^{-4})$
$\Omega_c^0 \to p \underline{K}_1^-$	$1.04 \times 10^{-4}(4.60 \times 10^{-5})$	$4.60 \times 10^{-4}(2.03 \times 10^{-4})$
$\Omega_c^0 \to n K_1^0$	$7.72 \times 10^{-5}(1.53 \times 10^{-4})$	$3.41 \times 10^{-4}(6.77 \times 10^{-4})$
$\Omega_c^0 \to n \underline{K}_1^0$	$1.04 \times 10^{-5}(4.58 \times 10^{-5})$	$4.58 \times 10^{-4}(2.02 \times 10^{-4})$
$\Omega_c^0 \to \Lambda f_1$	1.55×10^{-5}	6.83×10^{-5}
$\Omega_c^0 \to \Lambda f_1'$	1.00×10^{-4}	4.38×10^{-4}
$\Omega_c^0 \to \Sigma^+ a_1^-$	7.75×10^{-5}	3.42×10^{-4}
$\Omega_c^0 \to \Sigma^0 a_1^0$	7.74×10^{-5}	3.41×10^{-4}
$\Omega_c^0 \to \Sigma^- a_1^+$	7.72×10^{-5}	3.40×10^{-4}

The relative (signs) strengths of the S-, P-, and D-wave amplitudes may be attributed for the observed behavior. Similarly, we observe an increase in the branching ratio of the color favored mode $\Omega_c \to \Xi^- K_1^+$ where the factorization term appears to be dominant. It may also be noted that P-wave amplitudes acquire larger magnitude in both these decays.

(9) For Ω_c decays acquiring contributions from pole terms (W-exchange diagrams) only, the CKM-suppressed mode has BRs of $\mathcal{O}(10^{-4}\text{-}10^{-5})$. The dominant modes are $\Omega_c \to \Lambda \bar{K}_1^0/\Lambda \underline{\bar{K}}_1^0$ with BRs of $\mathcal{O}(10^{-4})$. It may be noted that among decays arising from pole diagrams only, $\Omega_c \to \Lambda \bar{K}_1^0/\Lambda \underline{\bar{K}}_1^0$ decays acquire most dominant pole amplitude contributions.

(10) In the case of CKM-suppressed modes coming via pole diagrams only, the BRs are comparable to CKM-suppressed modes of the same category. The dominant decays are $\Omega_c \to p\underline{K}_1^-/\Lambda f_1'$ with BR $\sim 1.00 \times 10^{-4}$. The branching ratios of all the remaining decays are of $\mathcal{O}(10^{-5})$. Despite CKM suppression, BRs of these decays compete well with CKM-suppressed modes due to higher pole contributions.

(11) The absence of weak PV transition amplitudes (b_{ij}'s) lead to zero decay asymmetries for the decay modes coming from W-exchange processes only.

(12) Also, we observe that mass dependence of strong couplings coming through SB effects result in larger strong couplings, and hence, higher BRs.

(13) The overall trend shows that the BRs of the decay modes involving ${}^3P_1$ (${}^1P_1$) axial-vector states tend to increase (decrease) for $\theta_{K_1} = -58°$.

(14) All the decays involving nonstrange mesons in the final state have zero u-pole contributions except for $\Omega_c^0 \to \Lambda f_1'$ decay, which acquires contributions from both the u and s channels. The highly suppressed decay modes $\Omega_c^0 \to \Xi^0 K_1^0/\Xi^- K_1^+$ have only s-pole contributions.

(15) The decay modes consisting of $b_1/h_1/h_1'$ mesons in the final states are forbidden in the isospin limit.

In the literature, several attempts have been made to establish the fact that the lifetimes of semileptonic and nonleptonic decays of heavy flavor baryons show a strong dependence on square of the baryon wave function overlap at the origin, $|\psi(0)|^2$ [43,85–87]. In order to lower the discrepancy in theory and experiment, one needs to take into account the variation of $|\psi(0)|^2$ (being a dimensional quantity). For example, in the case of nonleptonic decays, the inclusion of the flavor dependence of the hadron wave function at the origin has resulted in good agreement between theory and experiment [43,88]. Following the analysis given in Ref. [70], we consider the variation of $|\psi(0)|^2$ with flavor. It has been long advocated that a reliable estimate of wave function at the origin of the ground-state baryon can be obtained by experimental hyperfine splitting [89]. A straightforward hyperfine splitting calculation, using the constituent quark model, between Σ_c and Λ_c reveals

$$m_{\Sigma_c} - m_{\Lambda_c} = \frac{16\pi}{9}\alpha_s(m_c)\frac{(m_c - m_u)}{m_c m_u^2}|\psi(0)|_c^2, \quad (31)$$

where we assume $|\psi(0)|^2_{\Sigma_c} = |\psi(0)|^2_{\Lambda_c}$. We obtain the flavor enhancement scale in the strange and charm sectors from the following expression:

$$\frac{m_{\Sigma_c} - m_{\Lambda_c}}{m_\Sigma - m_\Lambda} = \frac{\alpha_s(m_c)}{\alpha_s(m_s)}\frac{m_s(m_c - m_u)|\psi(0)|_c^2}{m_c(m_s - m_u)|\psi(0)|_s^2}, \quad (32)$$

which yields

$$r \equiv \frac{|\psi(0)|_c^2}{|\psi(0)|_s^2} \approx 2.1 \quad (33)$$

for the choice $\alpha_s(m_c)/\alpha_s(m_s) \approx 0.53$ [42,63]. Finally, we discuss the effects of this scale enhancement due to variation of the spatial baryon wave function overlap on branching ratios. The flavor-dependent BRs for CKM-favored, suppressed, and doubly suppressed modes are evaluated using $|\psi(0)|^2$ variation. The numerical values of pole amplitudes only are given in column 3 of Table III. Consequently, the obtained numerical results for branching ratios and asymmetry parameters involving both factorizable and pole contributions are given in columns 4 and 6 of Table IV, whereas the branching predictions for the processes involving pole contributions only are shown in column 3 of Table V. We wish to point out that the implications of variation of the spatial baryon wave function overlap lead to a flavor enhancement scale ratio $(r) \sim 2$. This may also be seen simply as a variation in r from 1 to 2 for flavor-dependent $|\psi(0)|^2$ owing to dimensionality arguments. It may also be noted that the factorization hypothesis does not involve flavor-dependent effects. In the absence of any experimental and theoretical information, we compare our results with flavor-independent BRs. We observe the following:

(1) The variation of $|\psi(0)|^2$ has enhanced BRs of most of the decays by roughly a factor of 4 as compared to flavor-independent BRs. Consequently, the number of decay modes with BRs of $\mathcal{O}(10^{-3}) \sim \mathcal{O}(10^{-4})$ has become large.

(2) Since the factorization amplitudes remain unaffected by flavor-dependent effects, the change in BRs in all the cases may be attributed due to flavor-dependence effects on pole contributions.

(3) In the case of the Ω_c decays involving both factorizable and pole amplitudes, the dominant decay channels $\Omega_c^0 \to \Xi^0 \bar{K}_1^0$ (for CKM-favored) and $\Omega_c^0 \to \Xi^- a_1^+/\Xi^0 a_1^0/\Xi^0 f_1$ (for CKM-suppressed mode) have BRs of $\mathcal{O}(10^{-3})$, which make them viable candidates for the experimental search. The highest $\mathrm{Br}(\Omega_c^0 \to \Xi^- a_1^+) = 5.37 \times 10^{-3}$, where color enhancement has overcome CKM suppression. However, the branching ratio of color-suppressed $\Omega_c^0 \to \Xi^0 \bar{K}_1^0$ decay comes out to be smaller, i.e., 2.42×10^{-3}. It is worth remarking that in spite of constructive interference between factorizable and pole amplitudes, the BR of $\Omega_c^0 \to \Xi^0 \bar{K}_1^0$ decay tends to be small in comparison to that of CKM-suppressed mode. The reason is that the magnitude of the pole amplitude for CKM-favored mode is smaller by an order when compared to CKM-suppressed modes. However, the pole contributions to $\Omega_c^0 \to \Xi^0 \bar{K}_1^0$ decay arise up to 40%-50% because of flavor dependence which may further increase to 3.78×10^{-3} for $\theta_{K_1} = -58°$.

(4) Unlike flavor independent case, CKM-doubly-suppressed decay modes $\Omega_c \to \Xi^0 K_1^0/\Xi^- K_1^+$ show little change in BRs when flavor-dependent effects are included. Comparable factorizable and pole terms add to the ambiguity of these decay modes. The relative magnitude and signs of the S-, P-, and D-wave amplitudes become more important, as may be seen from variation in the asymmetry parameter (both in sign and magnitude). Only experimental observation of these modes can provide a clear picture.

(5) Among the Ω_c decay modes arising through pole contributions only, the dominant decay modes with BRs of $\mathcal{O}(10^{-3})$ are $\Omega_c^0 \to \Lambda\bar{K}_1^0/\Lambda\underline{\bar{K}}_1^0$ with higher $\mathrm{Br}(\Omega_c^0 \to \Lambda\underline{\bar{K}}_1^0) = 2.13 \times 10^{-3}$. The BRs of all the remaining decay channels in CKM-suppressed decay mode are enhanced to $\mathcal{O}(10^{-4})$. However, the BRs may further increase or decrease with $\theta_{K_1} = -58°$ for corresponding K_1 and $\underline{K}_1$ modes, respectively. The comparable BRs of $\Omega_c^0 \to \Lambda\bar{K}_1^0/\Lambda\underline{\bar{K}}_1^0$ to that of CKM-favored mode can be explained by dominant pole contributions to the former.

(6) The flavor-enhanced pole amplitudes have placed CKM-doubly-suppressed modes well in competition with CKM-suppressed modes. The BRs of all these decays, namely $\Omega_c^0 \to p\underline{K}_1^-/n\underline{\bar{K}}_1^0/\Lambda f_1'/pK_1^-/n\bar{K}_1^0/\Sigma^+ a_1^-/\Sigma^0 a_1^0/\Sigma^- a_1^+$, have increased by an order of magnitude, i.e., to $\mathcal{O}(10^{-4})$. All these decay channels possess experimentally observable decay widths.

V. SUMMARY

We have analyzed axial-vector-meson-emitting exclusive two-body nonleptonic weak decays of Ω_c^0 baryon for CKM-favored and suppressed modes in the factorization and pole model approach. We have obtained the factorizable contributions by using the nonrelativistic quark model (NRQM) [33] and heavy quark effective theory (HQET) [47] to evaluate the form factors f_i and g_i. We expected that W-exchange diagrams could dominate Ω_c^0 weak decays, and these are evaluated using pole model. The relevant baryon matrix elements of the weak Hamiltonian have been calculated which determine the pole term with short-distance QCD corrections. Also, we have observed that the mass dependence (SB effects) of strong couplings turns out to be crucial in deciding pole contributions to heavy baryon decays. These effects can be important specifically in the decays coming from the W-exchange process (pole diagram) only. Nonrelativistic evaluation of weak matrix elements involving the PC weak Hamiltonian has been carried out for flavor-independent and flavor-dependent cases. We have predicted BRs of Ω_c^0 decays for the cases a) involving both factorization and pole amplitudes and b) arising via pole amplitudes (W-exchange diagram) only. We list our results as follows:

(1) For the flavor-independent case, the only dominant decay mode $\Omega_c^0 \to \Xi^- a_1^+$ has a branching ratio of $\mathcal{O}(10^{-3})$. The next-order dominant modes are $\Omega_c^0 \to \Xi^0 \bar{K}_1^0/\Xi^0 a_1^0/\Xi^0 f_1$. All these decay modes consist of the interference of pole and factorizable contributions. In $\Omega_c^0 \to \Xi^- a_1^+$ decay, the dominant contribution comes from the factorization term while in the rest of the decay channels pole contributions dominate. For the decay arising from pole amplitudes only, the $\Omega_c^0 \to \Lambda \bar{K}_1^0/\Lambda f_1'$ has branching ratios of $\mathcal{O}(10^{-4})$.

(3) For the flavor-dependent case, we consider the variation of spatial baryon wave function overlap at the origin, i.e., $|\psi(0)|^2$ with flavor. We observe that the introduction of flavor dependence has raised the BRs of all the decays by roughly a factor of 4. A number of dominant modes $\Omega_c^0 \to \Xi^- a_1^+/\Xi^0 \bar{K}_1^0/\Xi^0 a_1^0/\Xi^0 f_1/\Lambda \bar{K}_1^0/\Lambda \underline{\bar{K}}_1^0$ now have BRs of $\mathcal{O}(10^{-3})$. All these decay channels fall within the limit of experimental reach.

(3) We wish to remark here that most of the decay channels in Ω_c^0 decay only through the W-exchange diagram; moreover, the W-exchange contributions dominate in the rest of the process, with some exceptions. Observation of such decays would shed some light on the mechanism of W-exchange effects in these decay modes.

A conventional concept expects the p-wave-emitting decays to be kinematically suppressed; however, we find that BRs of axial-vector-emitting decays of Ω_c^0 are comparable to the experimentally observed two-body s-wave-meson emitting decays of charm baryons. We hope this would generate ample interest in experimental search of these decay modes.

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