# Quark and pole models of nonleptonic decays of charmed baryons

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Quark and pole models of nonleptonic decays of charmed baryons are analyzed from the point of view of their symmetry properties. The symmetry structure of the parity-conserving amplitudes that corresponds to the contribution of the ground-state intermediate baryons is shown to differ from the one hitherto employed in the symmetry approach. It is pointed out that the "subtraction" of sea quark effects in hyperon decays leads to an estimate of $W$-exchange contributions in charmed baryon decays that is significantly smaller than naively expected on the basis of SU(4). An $SU(2)_W$ constraint questioning the reliability of the factorization technique is exhibited. Finally, a successful fit to the available data is presented.

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## I. INTRODUCTION

Over the last ten years a description of heavy meson weak decays known as the factorization approach has become available. In the area of heavy baryon decay theoretical progress has been fairly slow, however. To some extent this state of affairs was conditioned by the type and quality of experimental data. Only recently higher statistics data on nonleptonic weak decays of $\Lambda_c^+$ have become available. With several experiments on charmed baryons being now carried out at DESY, Cornell, CERN, and Fermilab the expected growth of the data basis has already started to stimulate a more intensive theoretical effort in this area.

Although views have been expressed that the dynamics of nonleptonic weak decays should become simpler as the decaying quark becomes heavier, a reliable approach to the decays of charmed baryons does not exist thus far. This is hardly surprising in view of the fact that in the long-studied related field of nonleptonic hyperon decays, there is still no consensus as to the relative importance and symmetry structure of various possible contributions to both the parity-violating and parity-conserving amplitudes [1]. In fact not only is it not clear there what the relative size and sign of pole model contributions from various intermediate states is (i.e., mesons, ground-state, and excited baryons—compare Refs. [2–4]), but even the value of the $f/d$ ratio characterizing the SU(3) structure of the soft pion contribution to the parity-violating amplitudes is not agreed upon. The valence quark model predicts $f/d=-1$ while the phenomenological analysis of Pham [5] suggests $f/d=-1.6$, much closer to that needed for a proper description of the parity-conserving amplitudes. Specific models to explain such deviation from the valence quark model have been proposed [6,7].

Given this situation it seems unlikely that in the near future we shall be able to predict through a reliable calculation the corresponding contributions in the decays of charmed baryons. Instead, it is probably the incoming data that will be instrumental in the broadening of our understanding of nonleptonic weak decays for baryons in general and its deepening for hyperons in particular.

To help resolve various emerging questions in a phenomenological way we adopt a framework based on symmetry considerations. The main topic of this paper is the discussion of the implications of various assumptions involving and/or affecting symmetry properties of models of nonleptonic decays of charmed baryons. The symmetry–quark-model approach adopted here is based on papers [8,9] and constitutes their generalization to the charmed baryon sector. Our approach (briefly described in Sec. II), although similar in spirit to the one used previously in this context [10–12], differs from the latter in an essential way. Namely, it turns out that the symmetry structure of the parity-conserving amplitudes of Refs. [10,11] does not correspond to that expected in the pole model with ground-state baryons in the intermediate state. In the present paper the correct symmetry structure of the pole model with ground-state intermediate baryons is used. Thus, our paper essentially replaces the previous symmetry-based papers on charmed baryon decays. Apart from the above difference in the treatment of parity-conserving amplitudes our paper differs from Refs. [10,11] by a more phenomenological treatment of single-quark processes. Furthermore, we point out a couple of uncertainties and corrections hitherto not noticed in the literature (Sec. III). Finally, using a symmetry approach as our framework we fit the existing experimental data (Sec. IV).

## II. THE BASIC QUARK-POLE MODEL

The aim of this paper is to discuss the symmetry structure of the quark and pole models of nonleptonic decays of charmed baryons and the implications of various assumptions involving and/or affecting symmetry properties. These assumptions may be tested by comparing symmetry properties of their predictions for the partial decay widths and asymmetries with the experimental ones.

For the decays with the emission of pseudoscalar mesons these partial decay widths and asymmetries are

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given in terms of the parity-violating ($A_P$) and the parity-conserving ($B_P$) amplitudes by

$$\Gamma=\frac{1}{4\pi}\frac{k(E_f+m_f)}{m_i}(|A_P|^2+|\bar{B}_P|^2)\ ,$$
$$\alpha=\frac{2A_P\bar{B}_P}{|A_P|^2+|\bar{B}_P|^2}\ , \quad (1)$$

where

$$\bar{B}_P=\frac{k}{E_f+m_f}B_P\ . \quad (2)$$

In Eq. (1), $m_i, m_f$ are the masses of the initial and final baryon, $E_f$ is the energy of the final baryon, and $k$ its decay momentum.

For the decays with the emission of vector mesons the corresponding formulas read

$$\Gamma=\frac{1}{4\pi}\frac{k(E_f+m_f)}{m_i}(|A_{V_\perp}|^2+|\bar{B}_{V_\perp}|^2+|A_{V_\parallel}|^2+|\bar{B}_{V_\parallel}|^2)\ ,$$
$$\alpha=\frac{2(A_{V_\perp}\bar{B}_{V_\perp}+A_{V_\parallel}\bar{B}_{V_\parallel})}{(|A_{V_\perp}|^2+|\bar{B}_{V_\perp}|^2+|A_{V_\parallel}|^2+|\bar{B}_{V_\parallel}|^2)}\ , \quad (3)$$

where

$$\bar{B}_{V_{\perp,\parallel}}=\frac{k}{E_f+m_f}B_{V_{\perp,\parallel}} \quad (4)$$

and $A_{V_{\perp,\parallel}}$, etc., are the amplitudes for the emission of transverse ($\perp$) and longitudinal ($\parallel$) vector mesons.

The approach of this paper constitutes an application to the charmed baryon sector of the quark-model technique used in the description of the $\Delta S=1$ hyperon decays in Refs. [8,9,13]. The main idea of these papers was to separate from the dynamics the quark-model-based spin-flavor symmetry relations between the amplitudes. The basic reason for adopting such an approach was the lack of general consensus among theorists on the relative size and sign of various dynamical contributions. The approach of Refs. [8,9] avoids such theoretical uncertainties by lumping various contributions into a few reduced matrix elements to be fitted from experiment. These reduced matrix elements correspond to quark diagrams of Fig. 1. Figures 1(a) and 1(a') correspond to the so-called factorization amplitudes, Figs. 1(b1), 1(b2), and 1(d) are $W$-exchange contributions while Figs. 1(c1) and 1(c2) summarize the effect of quark sea. In fact, Körner and his collaborators (see also Ref. [14]) never consider Fig. 1(c), the inclusion of which is crucial [8] for a proper description of hyperon decays. (An important difference between the hyperon and charmed-baryon sector is the absence of Fig. 1(c) in the latter. The implications of this difference shall be discussed in Sec. III.)

### A. Parity-violating amplitudes

Our approach to the parity-violating amplitude does not differ in an essential way from that of Refs. [10–12]. The contribution from the Fig. 1(d) is zero. Furthermore, if SU(4) symmetry were exact Figs. 1(a) and 1(a')

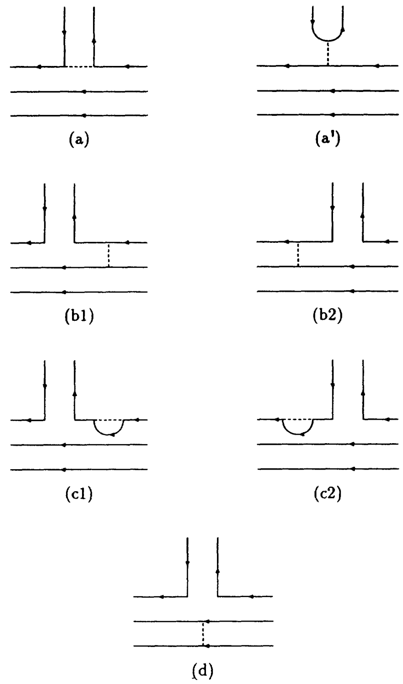


FIG. 1. Quark diagrams for weak decays.

would not contribute to the charmed-baryon decays with pseudoscalar meson emission (see, e.g., Ref. [8]). For the transverse vector mesons they do contribute, however, even in the limit of exact SU(4).

Calculations of the spin-flavor factors corresponding to the parity-violating amplitudes $A_P$, $A_V$ were done using the quark model technique of Ref. [8]. The results are gathered in Tables I and II.

The reduced matrix elements $\tilde{a}, \tilde{a}'$, and $\tilde{b}$ are related to those of hyperon decays by

$$\tilde{b}=b\cot\theta_C\ ,$$
$$\tilde{a}=a\cot\theta_C\ , \quad (5)$$
$$\tilde{a}'=a'\cot\theta_C\ ,$$

where $\theta_C$ is the Cabibbo angle and the parameters $a, a', b$ are the corresponding reduced matrix elements for hyperon decays. Their numerical values have been estimated in Refs. [8,9,15] to be (in units of $10^{-7}$)

$$b=-5.0\ ,$$
$$a=+3.8\ , \quad (6)$$
$$a'=-3.0\ ,$$

and, consequently, we have

$$\tilde{b}=-22.2\ ,$$
$$\tilde{a}=+16.9\ , \quad (7)$$
$$\tilde{a}'=-13.3\ .$$

In addition, the reduced matrix elements corresponding to the emission of longitudinal ($\tilde{b}'$) and transverse ($\tilde{b}$) vector mesons are related in the quark model by

$$\tilde{b}'=\tilde{b}\ . \tag{8}$$

In the SU(4)-broken world, as the factorization approximation indicates, Figs. 1(a) and 1(a′) do contribute to the decays with pseudoscalar meson emission. This has been taken into account in Table I where such contributions have been given strength $g$ and $g'$, respectively.

Estimates of $g$ and $g'$ through factorization [16,17] give for both of them similar values (though with opposite signs) of around (in units of $10^{-7}$)

$$4.0\text{–}6.0\ . \tag{9}$$

Comparing Eq. (7) with Eq. (9) we see that for charmed-baryon decays, as for hyperon decays, the factorization amplitudes still appear to give bigger contributions in the $B\to B'V$ than in the $B\to B'P$ decays although in the latter they are no longer negligible. We shall come back to the discussion of the factorization amplitudes in Secs. III C and III D.

The non-negligible size of the factorization amplitudes $g,g'$, as required by nonvanishing experimental asymmetry in the $\Lambda_c^+\to\Lambda\pi^+$ decay, indicates significant SU(4) breaking effects resulting from a large mass difference between charmed and noncharmed (constituent) quarks:

$$m_c-m_{u,d,s}\approx 1.1\ \text{GeV}\ . \tag{10}$$

Such a large mass difference must lead to significant differences between the standard current-algebra–quark-model approach and the pole model. In the pole model the dominant contribution to the parity-violating amplitudes comes from the lowest-lying negative-parity $\frac{1}{2}^-$ excited baryons propagating between the $W$-exchange and strong decay interactions shown in Figs. 1(b1) and 1(b2). As discussed in Refs. [17,18] the current-algebra and the pole model become equivalent in the SU(4) limit when

$$0\leftarrow m_c-m_{u,d,s}\ll m(\tfrac{1}{2}^-)-m(\tfrac{1}{2}^+)\approx 0.5\ \text{GeV}\ . \tag{11}$$

Then, one can sum the contributions from the intermediate $\frac{1}{2}^-$ baryon resonances and obtain the standard quark-model–current-algebra prescription in which no information on the intermediate $\frac{1}{2}^-$ states is needed.

In reality Eq. (11), is, of course, not satisfied and significant departures from simple current algebra predictions may be anticipated. Such effects were discussed in Ref. [17]. In this paper they are not considered. The reasons behind their neglect are as follows.

First, we want to give a symmetry prediction that, unlike the one given by Körner and collaborators [10–12], does correspond to the standard pole model prescriptions for the parity-conserving amplitudes. Second, we want to point out other ambiguities that as yet have not been dis-

TABLE I. Weak amplitudes for $B_c\to BP$ decays.

| Process | $A_P$ | Fig. 1(b1) | Fig. 1(b2) | $B_P$ |
|---|---|---|---|---|
| $\Xi_c^+\to\Xi^0\pi^+$ | $-\frac{1}{2\sqrt{6}}\tilde{b}+\sqrt{3/2}g'$ | 0 | $-\frac{1}{2\sqrt{6}}$ | $\frac{1}{2\sqrt{6}}B-\frac{1}{2\sqrt{6}}M'$ |
| $\Xi_c^+\to\Sigma^+\bar{K}^0$ | $-\frac{1}{2\sqrt{6}}\tilde{b}-\sqrt{3/2}g$ | 0 | $-\frac{1}{2\sqrt{6}}$ | $\frac{1}{2\sqrt{6}}B+\frac{1}{2\sqrt{6}}M$ |
| $\Xi_c^0\to\Xi^0\pi^0$ | $-\frac{1}{2\sqrt{3}}\tilde{b}$ | $\frac{1}{12\sqrt{3}}$ | $-\frac{1}{4\sqrt{3}}$ | $\frac{1}{3\sqrt{3}}B$ |
| $\Xi_c^0\to\Xi^0\eta_8$ | $\frac{1}{6}\tilde{b}$ | $\frac{1}{12}$ | $\frac{1}{12}$ | 0 |
| $\Xi_c^0\to\Xi^0\eta_1$ | $-\frac{1}{6\sqrt{2}}\tilde{b}$ | 0 | $\frac{1}{6\sqrt{2}}$ | $-\frac{1}{6\sqrt{2}}B$ |
| $\Xi_c^0\to\Xi^-\pi^+$ | $\frac{1}{2\sqrt{6}}\tilde{b}+\sqrt{3/2}g'$ | $-\frac{1}{6\sqrt{6}}$ | 0 | $-\frac{1}{6\sqrt{6}}B-\frac{1}{2\sqrt{6}}M'$ |
| $\Xi_c^0\to\Sigma^0\bar{K}^0$ | $-\frac{1}{4\sqrt{3}}\tilde{b}-\frac{\sqrt{3}}{2}g$ | $-\frac{1}{6\sqrt{3}}$ | $-\frac{1}{4\sqrt{3}}$ | $\frac{1}{12\sqrt{3}}B+\frac{1}{4\sqrt{3}}M$ |
| $\Xi_c^0\to\Lambda\bar{K}^0$ | $-\frac{1}{4}\tilde{b}+\frac{1}{2}g$ | 0 | $-\frac{1}{12}$ | $\frac{1}{12}B-\frac{1}{12}M$ |
| $\Xi_c^0\to\Sigma^+K^-$ | 0 | $\frac{1}{3\sqrt{6}}$ | 0 | $\frac{1}{3\sqrt{6}}B$ |
| $\Lambda_c^+\to\Sigma^+\pi^0$ | $-\frac{1}{2\sqrt{3}}\tilde{b}$ | $-\frac{1}{12\sqrt{3}}$ | $-\frac{1}{4\sqrt{3}}$ | $\frac{1}{6\sqrt{3}}B$ |
| $\Lambda_c^+\to\Sigma^+\eta_8$ | $\frac{1}{6}\tilde{b}$ | $-\frac{1}{12}$ | $\frac{1}{12}$ | $-\frac{1}{6}B$ |
| $\Lambda_c^+\to\Sigma^+\eta_1$ | $-\frac{1}{6\sqrt{2}}\tilde{b}$ | 0 | $\frac{1}{6\sqrt{2}}$ | $-\frac{1}{6\sqrt{2}}B$ |
| $\Lambda_c^+\to\Sigma^0\pi^+$ | $\frac{1}{2\sqrt{3}}\tilde{b}$ | $\frac{1}{12\sqrt{3}}$ | $\frac{1}{4\sqrt{3}}$ | $-\frac{1}{6\sqrt{3}}B$ |
| $\Lambda_c^+\to\Lambda\pi^+$ | $-g'$ | $\frac{1}{12}$ | $\frac{1}{12}$ | $\frac{1}{6}M'$ |
| $\Lambda_c^+\to\Xi^0K^+$ | 0 | $\frac{1}{3\sqrt{6}}$ | 0 | $\frac{1}{3\sqrt{6}}B$ |
| $\Lambda_c^+\to p\bar{K}^0$ | $\frac{1}{2\sqrt{6}}\tilde{b}-\sqrt{3/2}g$ | $-\frac{1}{6\sqrt{6}}$ | 0 | $-\frac{1}{6\sqrt{6}}B+\frac{1}{2\sqrt{6}}M$ |

cussed in the literature at all. Third, we think that a reliable inclusion of SU(4) breaking effects might be very difficult. We believe that it will be the experiment that will guide us on our way to a theoretical understanding of how to properly take such effects into account. Accordingly, the simple current-algebra–quark-model approach and its predictions are of great interest themselves since they provide the basis for future discussion of various departures from such simple models.

### B. Parity-conserving amplitudes

Calculation of the parity-conserving amplitudes $B_P, B_V$ is similar to the calculation of the previous subsection. There are two main contributions to the amplitudes. The first is due to the intermediate baryons [Figs. 1(b1), 1(b2), and 1(d)], the second (due to meson poles) is often treated in the factorization approximation [Figs. 1(a) and 1(a′)]. Evaluation of the symmetry structure of the second contribution is straightforward and leads to the pattern exhibited in Tables I and II. In these tables the reduced matrix elements corresponding to Figs. 1(a) and 1(a′) are denoted by $M, M'$ when the emitted meson is an $SU(2)_W$ triplet $(P, V_\perp)$ and by $m, m'$ when it is an $SU(2)_W$ singlet $(V_\parallel)$. In the next section we shall discuss these contributions and their actual size in more detail.

The contribution from the intermediate baryons requires the calculation of the spin-flavor structure of Figs. 1(b1), 1(b2), and 1(d). For the $B_c \to BP$ decays the individual spin-flavor factors corresponding to Figs. 1(b1) and 1(b2) are shown in Table I. In order to obtain the symmetry structure of the baryon pole contributions to the parity-conserving amplitudes, the spin-flavor factors corresponding to Figs. 1(b1) and 1(b2) have to be multiplied by appropriate energy denominators and then added. A closer inspection of these [assume SU(3), i.e., $\Sigma=\Lambda=N=\Xi$; $\Xi_c^+=\Xi_c^0=\Lambda_c^+$] shows that if the intermediate baryons are in the ground state this addition procedure effectively results in the subtraction of the spin-flavor factors of Figs. 1(b1) and 1(b2). The same procedure, when applied to two versions of Fig. 1(d) (with $W$ exchange followed by strong decay and vice versa), leads to the cancellation of these two contributions on account of their identical spin-flavor structure. Thus, no overall contribution from Fig. 1(d) is obtained. The above subtraction procedure may be verified by explicitly calculating all

TABLE II. Weak amplitudes for $B_c \to BV$ decays.

| Process | $A_{V_\perp}$ | $B_{V_\perp}$ | $A_{V_\parallel}$ | $B_{V_\parallel}$ |
|---|---|---|---|---|
| $\Xi_c^+ \to \Xi^0 \rho^+$ | $-\frac{1}{2\sqrt{3}}\tilde{b}+\frac{2}{\sqrt{3}}\tilde{a}'$ | $-\frac{1}{2\sqrt{3}}B+\frac{1}{2\sqrt{3}}M'$ | $\frac{1}{2\sqrt{6}}\tilde{b}'$ | $-\frac{1}{2\sqrt{6}}B'+\frac{1}{2\sqrt{6}}m'$ |
| $\Xi_c^+ \to \Sigma^+ \bar{K}^{*0}$ | $-\frac{1}{2\sqrt{3}}\tilde{b}-\frac{2}{\sqrt{3}}\tilde{a}$ | $-\frac{1}{2\sqrt{3}}B-\frac{1}{2\sqrt{3}}M$ | $\frac{1}{2\sqrt{6}}\tilde{b}'$ | $-\frac{1}{2\sqrt{6}}B'-\frac{1}{2\sqrt{6}}m$ |
| $\Xi_c^0 \to \Xi^0 \rho^0$ | $-\frac{1}{3\sqrt{6}}\tilde{b}$ | $-\frac{2}{3\sqrt{6}}B$ | $\frac{1}{3\sqrt{3}}\tilde{b}'$ | $-\frac{1}{2\sqrt{3}}B'$ |
| $\Xi_c^0 \to \Xi^0 \omega$ | $+\frac{2}{3\sqrt{6}}\tilde{b}$ | $+\frac{1}{3\sqrt{6}}B$ | $-\frac{1}{6\sqrt{3}}\tilde{b}'$ | 0 |
| $\Xi_c^0 \to \Xi^0 \phi$ | $-\frac{1}{6\sqrt{3}}\tilde{b}$ | $+\frac{1}{6\sqrt{3}}B$ | $-\frac{1}{6\sqrt{6}}\tilde{b}'$ | $-\frac{1}{2\sqrt{6}}B'$ |
| $\Xi_c^0 \to \Xi^- \rho^+$ | $-\frac{1}{6\sqrt{3}}\tilde{b}+\frac{2}{\sqrt{3}}\tilde{a}'$ | $\frac{1}{6\sqrt{3}}B+\frac{1}{2\sqrt{3}}M'$ | $-\frac{1}{6\sqrt{6}}\tilde{b}'$ | $\frac{1}{2\sqrt{6}}B'+\frac{1}{2\sqrt{6}}m'$ |
| $\Xi_c^0 \to \Sigma^0 \bar{K}^{*0}$ | $-\frac{5}{6\sqrt{6}}\tilde{b}-\sqrt{2/3}\tilde{a}$ | $-\frac{1}{6\sqrt{6}}B-\frac{1}{2\sqrt{6}}M$ | $\frac{1}{12\sqrt{3}}\tilde{b}'$ | $\frac{1}{4\sqrt{3}}B'-\frac{1}{4\sqrt{3}}m$ |
| $\Xi_c^0 \to \Lambda \bar{K}^{*0}$ | $-\frac{1}{6\sqrt{2}}\tilde{b}+\frac{2}{3\sqrt{2}}\tilde{a}$ | $-\frac{1}{6\sqrt{2}}B+\frac{1}{6\sqrt{2}}M$ | $\frac{1}{12}\tilde{b}'$ | $-\frac{1}{4}B'+\frac{1}{12}m$ |
| $\Xi_c^0 \to \Sigma^+ \bar{K}^{*-}$ | $\frac{1}{3\sqrt{3}}\tilde{b}$ | $-\frac{1}{3\sqrt{3}}B$ | $\frac{1}{3\sqrt{6}}\tilde{b}'$ | 0 |
| $\Lambda_c^+ \to \Sigma^+ \rho^0$ | $-\frac{2}{3\sqrt{6}}\tilde{b}$ | $-\frac{1}{3\sqrt{6}}B$ | $\frac{1}{6\sqrt{3}}\tilde{b}'$ | $-\frac{1}{2\sqrt{3}}B'$ |
| $\Lambda_c^+ \to \Sigma^+ \omega$ | $\frac{1}{3\sqrt{6}}\tilde{b}$ | $\frac{2}{3\sqrt{6}}B$ | $-\frac{1}{3\sqrt{3}}\tilde{b}'$ | 0 |
| $\Lambda_c^+ \to \Sigma^+ \phi$ | $\frac{1}{6\sqrt{3}}\tilde{b}$ | $-\frac{1}{6\sqrt{3}}B$ | $\frac{1}{6\sqrt{6}}\tilde{b}'$ | $-\frac{1}{2\sqrt{6}}B'$ |
| $\Lambda_c^+ \to \Sigma^0 \rho^+$ | $\frac{2}{3\sqrt{6}}\tilde{b}$ | $\frac{1}{3\sqrt{6}}B$ | $-\frac{1}{6\sqrt{3}}\tilde{b}'$ | $\frac{1}{2\sqrt{3}}B'$ |
| $\Lambda_c^+ \to \Lambda \rho^+$ | $\frac{1}{3\sqrt{2}}\tilde{b}-\frac{4}{3\sqrt{2}}\tilde{a}'$ | $-\frac{1}{3\sqrt{2}}M'$ | 0 | $-\frac{1}{6}m'$ |
| $\Lambda_c^+ \to \Xi^0 K^{*+}$ | $\frac{1}{3\sqrt{3}}\tilde{b}$ | $-\frac{1}{3\sqrt{3}}B$ | $\frac{1}{3\sqrt{6}}\tilde{b}'$ | 0 |
| $\Lambda_c^+ \to p \bar{K}^{*0}$ | $-\frac{1}{6\sqrt{3}}\tilde{b}-\frac{2}{\sqrt{3}}\tilde{a}$ | $\frac{1}{6\sqrt{3}}B-\frac{1}{2\sqrt{3}}M$ | $\frac{-1}{6\sqrt{6}}\tilde{b}'$ | $\frac{1}{2\sqrt{6}}B'-\frac{1}{2\sqrt{6}}m$ |
| $\Lambda_c^+ \to p \phi$ | $\frac{2}{\sqrt{3}}\tilde{a}\tan\theta_C$ | $\frac{1}{2\sqrt{3}}M\tan\theta_C$ | 0 | $\frac{1}{2\sqrt{6}}m\tan\theta_C$ |

the necessary $B'BM$ strong couplings and weak baryon-to-baryon matrix elements $\langle B|H^{PC}_{weak}|B'\rangle$ and then combining them according to the prescriptions of the pole model. In the process, the contributions from $W$ exchanges between quarks not involved in meson emission get canceled and the symmetry structure of the resulting amplitudes is that obtained from the subtraction of Figs. 1(b1) and 1(b2) (see also Appendix A of Ref. [7]).

If simple symmetry arguments are applied to link (the $W$-exchange contribution to) the parity-conserving hyperon and charmed-baryon decay amplitudes one obtains for the reduced matrix element $B$ of Tables I and II the value

$$B=12\left[1-\frac{F}{D}\right]C\cot\theta_C\ , \tag{12}$$

where $C=-33$ is the value fitted in hyperon nonleptonic decays [9,15] and $F/D=\frac{2}{3}$. In deriving Eq. (12) we took into account the effect discussed in Sec. III B which diminishes the size of $H^{PC}$ matrix elements in the charmed baryon sector.

The estimate of Eq. (12) is not correct, however, since it does not consider the large difference in the size of pole factors $1/(B_{i,f}-B')$ appearing in hyperon and charmed-baryon parity-conserving amplitudes. More properly, Eq. (12) should be replaced by

$$B=12\left[1-\frac{F}{D}\right]C\cot\theta_C\frac{m_{\Sigma,\Lambda}-m_N}{m_{B_c}-m_{\Sigma,\Lambda}}\ . \tag{13}$$

Equation (13) gives as a rough estimate (in units of $10^{-7}$):

$$B\approx-95 \tag{14}$$

which compares well with the value $-73$ of Ref. [16]. The quark model relates the reduced matrix elements $B'$ and $B$ by

$$B'=B\ . \tag{15}$$

We proceed now to the discussion of the implications of various assumptions affecting and/or involving symmetry properties.

## III. DISCUSSION

### A. Symmetry structure of parity-conserving amplitudes

The symmetry structure of the parity-conserving amplitudes in the standard pole model differs from that given by Körner and collaborators [10–12]. In the pole model of Sec. II B flavor symmetry is kept in strong vertices but not in the baryon-to-baryon matrix elements (e.g., masses and weak transition elements). In the case of ground-state intermediate baryons this leads to the effective *subtraction* of the spin-flavor factors corresponding to Fig. 1(b1) and 1(b2). On the other hand, a closer look at Table 10 and Eq. (7) of Ref. [11] reveals that in the approach of Körner these factors are *added*. The net outcome of this difference is probably most easily seen on the example of the $\Lambda_c^+\to\Lambda\pi^+$ decay. Namely, it is well known that the parity-conserving amplitude of this decay receives no contributions from the baryon pole terms in the appropriate symmetry limit [19]. In Refs. [10–12] the total contribution from Figs. 1(b1) and 1(b2) is, however, nonvanishing. In other words, in Refs. [10–12] the intermediate baryons are all assumed to be much heavier than the external ground-state baryons.

One encounters the latter situation, e.g., in the parity-violating hyperon decay amplitudes. There, the intermediate $\frac{1}{2}^-$ excited baryons are indeed heavier than the external ground-state $\frac{1}{2}^+$ baryons. However, for parity-conserving hyperon decay amplitudes the assumptions of Refs. [10–12] correspond to neglecting the dominant contribution arising from the intermediate ground-state baryons (which is singular in the flavor symmetry limit). Consequently, it is the prescription of the previous subsection (i.e., Sec. II B) and not that of Refs. [10–12] that corresponds to the symmetry structure of the standard pole model of the parity-conserving amplitudes.

The agreement of the symmetry structure of the parity-violating hyperon decay amplitudes as calculated in the quark and pole models is thus, to some extent, accidental. Namely, had the "excited" $B^*(\frac{1}{2}^-)$ baryons been degenerate with the ground-state $B(\frac{1}{2}^+)$ baryons [but assuming broken SU(3), i.e., $\Lambda^*=\Lambda=\Sigma^*=\Sigma>N^*=N$] we would have ended up with an analogous situation in the parity-violating sector (see also Ref. [20]).

As it is obvious from the above discussion, in general both the parity-violating and the parity-conserving amplitudes may contain pieces with symmetry structure of both the sum and the difference of spin-flavor factors of Figs. 1(b1) and 1(b2). Which of the two is dominant (if any) depends on the dynamics. Similar considerations apply, of course, also to Fig. 1(d). The smallness of its contribution to the parity-conserving amplitudes, as obtained in the fit of Ref. [11], should perhaps be understood as an indication of the dominance of the "difference" structure in Fig. 1(d), in agreement with the prescriptions of the standard pole model with intermediate ground-state baryons. Clearly, the smallness of Fig. 1(d) obtained in Ref. [11] cannot be understood as a complete phenomenological "proof" of the dominance of this "difference" structure since in their fit Körner and Kramer used the "sum" structure for Fig. 1(b).

It is very unfortunate that the highlighted above essential difference between the (naively applied) arguments of symmetry and the structure of the standard pole model, although recognized already in the classical treatise of Marshak, Riazuddin, and Ryan [21], has been forgotten in various later papers and books on the subject (see, e.g., Ref. [22]).

### B. The SU(4) link between the hyperon and charmed-baryon decays

To calculate the absolute size of the nonleptonic decays of charmed baryons some authors (e.g., Ref. [19], for other references see Ref. [16]) use SU(4) symmetry to get the relevant information from hyperon decays. The way in which SU(4) symmetry is applied in such approaches is equivalent to the consideration of symmetry relationships between the baryon-to-baryon matrix elements of the parity-conserving part of the $W$-exchange contribution. The relevant diagrams are shown in Fig. 2(a).

The diagrams of Fig. 2(a) lead to the well-known SU(4) relation which connects charmed-baryon and hyperon nonleptonic decays:

$$\langle \Sigma^+ | H^{\rm PC}_{\rm weak} | \Lambda_c^+ \rangle = \frac{1}{\sqrt{6}} \cot\theta_C \langle p | H^{\rm PC}_{\rm weak} | \Sigma^+ \rangle \ . \quad (16)$$

It has been argued [16] that SU(4) symmetry breaking due to the large mass difference between $c$ and $s$ quarks should lead to a large mismatch in the baryon wave functions used in the overlap integrals in Eq. (16). As a result the baryon matrix elements of the $\Delta C = +1$ weak Hamiltonian should be smaller than that given by Eq. (16). Estimates in the bag model [23,24,16] yield a correction factor of around 0.5.

Here we point out another reason why these matrix elements should be smaller than expected on the basis of Eq. (16). Namely, in the quark-model–symmetry approach of Refs. [8,9] there is a large contribution to the $\langle p | H^{\rm PC}_{\rm weak} | \Sigma^+ \rangle$ matrix element that comes from the "sea-quark" diagrams [Fig. 2(b) or 1(c)]. On the other hand, in charmed-baryon decays the ($c$)-type diagrams are absent. Consequently, one has to "subtract" from the experimental value of the $\langle p | H^{\rm PC}_{\rm weak} | \Sigma^+ \rangle$ matrix element this part of it that is due to Fig. 2(b). This leads to the replacement of formula (16) by

$$\langle \Sigma^+ | H^{\rm PC}_{\rm weak} | \Lambda_c^+ \rangle = \frac{2}{1-(f/d)_{\rm soft\ meson}} \frac{1}{\sqrt{6}} \cot\theta_C \langle p | H^{\rm PC}_{\rm weak} | \Sigma^+ \rangle \ , \quad (17)$$

where $(f/d)_{\rm soft\ meson}$ is the ratio of the invariant SU(3) couplings $f$ and $d$ in the soft meson approximation to the parity-violating amplitudes of nonleptonic hyperon decays (or in the baryon-to-baryon matrix elements of the parity-conserving part of the $\Delta S = 1$ weak Hamiltonian). Estimates of $(f/d)_{\rm soft\ meson}$ vary. If one uses the estimate of Refs. [5,7] ($f/d = -1.6$) one gets a suppression factor of

$$\frac{2}{1-f/d} \rightarrow 0.77 \ . \quad (18)$$

If, on the other hand, one assigns the whole experimentally observed deviation from $f/d = -1$ in the parity-violating amplitudes [$(f/d)_{\rm PV} = -2.5$] to the soft-meson term (and nothing to other possible terms) one obtains

$$\frac{2}{1-f/d} \rightarrow 0.56 \ . \quad (19)$$

Apparently, "subtraction" of this part of the $f$ coupling that does not come from the $W$-exchange diagram leads to a substantial correction to the naive SU(4) formula (16).

The origin of the deviation of $f/d$ from its naive quark model value of $-1$ has not been agreed upon yet. We think that the main correction is due to the sea-quark effects discussed in Refs. [6,7]. Such effects not only renormalize the $f/d$ ratio but, on account of the large mass of the charmed quark, might renormalize differently the $W$-exchange diagrams in the $\Sigma^+ \rightarrow p$ and $\Lambda_c^+ \rightarrow \Sigma^+$ transitions (a part of the sea contribution in the latter, the $c\bar{c}$ sea, should be negligible). We have checked by explicit calculation, however, that in the framework of the hadron-loop model for the quark sea (Refs. [7,25]), the resulting renormalizations of both transitions are identical precisely when the $c\bar{c}$ sea is neglected.

The effect discussed above has been taken into account in Sec. II B where we related the size of the reduced matrix elements in the hyperon and charmed-baryon parity-conserving amplitudes [Eqs. (12) and (13)].

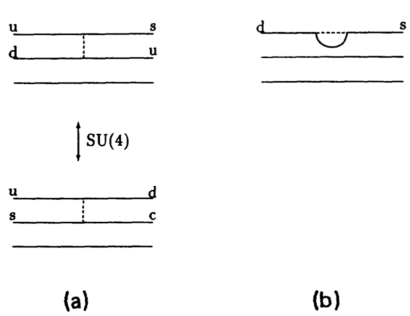


FIG. 2. Quark diagrams for the baryon-to-baryon matrix elements of the parity-conserving part of the weak Hamiltonian.

### C. Factorization and $SU(2)_W$ in parity-conserving amplitudes

Evaluation of the Figs. 1(a) and 1(a′) is most widely performed through the use of the factorization technique. In this approach one starts with the QCD-corrected effective weak Hamiltonian which, for the $\Delta C = \Delta S = +1$ processes in question takes the form

$$H_{\rm weak} = \frac{G \cos^2\theta_C}{\sqrt{2}} (c_1 O_1 + c_2 O_2) \ , \quad (20)$$

where

$$O_1 = [\bar{s}\gamma_\mu(1-\gamma_5)c][\bar{u}\gamma_\mu(1-\gamma_5)d] \ ,$$
$$O_2 = [\bar{u}\gamma_\mu(1-\gamma_5)c][\bar{s}\gamma_\mu(1-\gamma_5)d] \ . \quad (21)$$

The Wilson coefficients $c_1, c_2$ include the short-range QCD effects and for charm decays they have the values

$$c_1 \approx 1.3 \ ,$$
$$c_2 \approx -0.6 \ . \quad (22)$$

In accordance with the factorization idea the $(\bar{u}d)$ current in $O_1$ [$(\bar{s}d)$ in $O_2$] generates $\pi^+$ or $\rho^+$ [$\bar{K}^0$ or $\bar{K}^{*0}$] out of hadronic vacuum. [This is the so-called "new factorization" in which the Fierz-transformed contributions from Eq. (21) are simply dropped. Such an assumption has now considerable experimental support [11,16,17].] Operator $O_1$ corresponds to Fig. 1(a′), while $O_2$ to Fig. 1(a) after its "customization" by Fierz transformation to the needs of factorization technique.

Let us consider the factorization contribution to the parity-conserving amplitudes of the $\Lambda_c^+ \rightarrow \Lambda\pi^+$ and $\Lambda_c^+ \rightarrow \Lambda\rho^+$ decays. From Eq. (20) one then obtains

$$\langle \pi^+ \Lambda | H^{\rm PC}_{\rm weak} | \Lambda_c^+ \rangle_{\rm fact} = \frac{G\cos^2\theta_C}{\sqrt{2}} c_1 \langle \pi^+ | A^\mu | 0 \rangle \langle \Lambda | A_\mu | \Lambda_c^+ \rangle ,$$
$$\langle \rho^+ \Lambda | H^{\rm PC}_{\rm weak} | \Lambda_c^+ \rangle_{\rm fact} = \frac{G\cos^2\theta_C}{\sqrt{2}} c_1 \langle \rho^+ | V^\mu | 0 \rangle \langle \Lambda | V_\mu | \Lambda_c^+ \rangle . \quad (23)$$

The matrix elements of the currents in Eq. (23) are given by

$$\langle \pi^+ | A^\mu | 0 \rangle = f_\pi q^\mu ,$$
$$\langle \rho^+ | V^\mu | 0 \rangle = \epsilon^\mu f_\rho ,$$
$$\langle \Lambda | A_\mu | \Lambda_c^+ \rangle = g^A_{\Lambda\Lambda_c^+}(m_\pi^2) \bar{u}_\Lambda \gamma_\mu \gamma_5 u_{\Lambda_c^+} ,$$
$$\langle \Lambda | V_\mu | \Lambda_c^+ \rangle = f^V_{\Lambda\Lambda_c^+}(m_\rho^2) \bar{u}_\Lambda \gamma_\mu u_{\Lambda_c^+} , \quad (24)$$

where $f_\pi = 0.13$ GeV, $f_\rho = 0.17$ GeV$^2$, and $g^A, f^V$ are axial-vector and vector form factors.

Let us now see if the factorization prescription is consistent with the $SU(2)_W$ symmetry between the $\pi^+$ and $\rho^+$ couplings as employed in the previous section. Application of the requirement of $SU(2)_W$ symmetry to the couplings of Eq. (23) leads to the condition

$$f_\rho = f_\pi (m_\Lambda + m_{\Lambda_c^+}) \frac{g^A_{\Lambda\Lambda_c^+}(m_\pi^2)}{f^V_{\Lambda\Lambda_c^+}(m_\rho^2)} . \quad (25)$$

The ratio of $g^A/f^V$ is 1 in the simplest approach. If the bag model calculations of these form factors are employed (Ref. [26]) one obtains instead (with $g^A_{\Lambda\Lambda_c^+} = 0.50$, $f^V_{\Lambda\Lambda_c^+} = 0.46$)

$$g^A_{\Lambda\Lambda_c^+}(m_\pi^2) = g^A_{\Lambda\Lambda_c^+} \left[ 1 - \frac{m_\pi^2}{m_A^2} \right]^{-2} \approx 0.50 \quad (m_A = 2.54 \text{ GeV}) ,$$
$$f^V_{\Lambda\Lambda_c^+}(m_\rho^2) = f^V_{\Lambda\Lambda_c^+} \left[ 1 - \frac{m_\rho^2}{m_*^2} \right]^{-2} \approx 0.61 \quad (m_* = 2.11 \text{ GeV}) , \quad (26)$$

and the relevant ratio of axial-vector and vector form factors becomes smaller:

$$\frac{g^A_{\Lambda\Lambda_c^+}(m_\pi^2)}{f^V_{\Lambda\Lambda_c^+}(m_\rho^2)} = 0.82 . \quad (27)$$

Using the above value of $g^A/f^V$ Eq. (25) then reads

$$0.17 \text{ GeV}^2 = 0.36 \text{ GeV}^2 . \quad (28)$$

There is therefore a factor of 2 discrepancy (2.5 if $g^A/f^V = 1$ is used) between the $SU(2)_W$ symmetry predictions and the standard factorization technique. Similar discrepancy exists between the $SU(2)_W$ and factorization predictions for the $\bar{K}^0$ and $\bar{K}^{*0}$ production amplitudes of Fig. 1(a). One has to keep in mind, however, that, in principle, the factorization amplitude constitutes but a single contribution to the meson-pole terms [1]. Unfortunately, direct theoretical estimates of these contributions do not seem to be reliable [1]. If one believes in the accuracy of the $SU(2)_W$ symmetry predictions, the discrepancy of Eq. (28) shows that the factorization technique may be trusted here to within a factor of 2 only. Such accuracy is insufficient for making reliable predictions. On the other hand, if the contributions from the $f_2 \bar{u}_f \sigma_{\mu\nu} q^\nu u_i$ and $g_2 \bar{u}_f \sigma_{\mu\nu} \gamma_5 q^\nu u_i$ terms to the current matrix elements are considered (as in Ref. [17]) the disagreement in question is much milder ($\approx 30\%$).

### D. Factorization and sextet dominance

The relative size of the factorization contribution to the nonstrange ($\pi^+, \rho^+$) and strange ($\bar{K}^0, \bar{K}^{*0}$) meson emission is fixed by (22) and SU(3) symmetry-breaking factors such as $f_K/f_\pi$ as well as by the $q^2$ dependence of the form factors $g^A$ and $f^V$. Calculations along these lines are straightforward (e.g., see Ref. [17]). To relate such calculations to the parametrization of this paper we express below the results of Refs. [16,17] in terms of our reduced matrix elements.

For the parity-conserving amplitudes the estimates of the factorization amplitudes of Cheng and Tseng [17] correspond to the following values of the $M, M'$ parameters of Sec. II (in units of $10^{-7}$).

(a) For the pseudoscalar mesons,

$$M \approx 75 , \quad M' \approx -120 . \quad (29)$$

(b) For the transverse vector mesons,

$$M \approx 61 , \quad M' \approx -88 . \quad (30)$$

One observes that $(M'/M)_{\rm CT} \approx -1.5$, not very far from the sextet-dominance relation $M'/M = -1$.

For the longitudinal vector mesons sextet dominance requires similarly $m'/m = -1$, while the quark model relates the reduced matrix elements $m$, $M$ by

$$m = -M . \quad (31)$$

For the parity-conserving amplitudes the calculations of Ref. [17] correspond to (in units of $10^{-7}$)

$$g_{\rm CT} \approx 5 , \quad g'_{\rm CT} \approx -6 , \quad (32)$$

while those of Ref. [16] yield

$$g_{XK} \approx 3.4 \pm 1 , \quad g'_{XK} \approx -6.5 . \quad (33)$$

In the vector-meson sector results of Ref. [17] are translated into our scheme as

$$\tilde{a}_{\rm CT} = 9.6 , \quad \tilde{a}'_{\rm CT} = -14.3 . \quad (34)$$

Again, the ratios $\tilde{a}'/\tilde{a}$ or $g'/g$ are around $-1.5$, not very far from the sextet-dominance value of $-1$. The estimates of Eq. (34) correspond to

$$a_{\rm CT}=+2.2\ ,\qquad a'_{\rm CT}=-3.2\ . \tag{35}$$

These numbers should be compared with an estimate of Desplanques, Donoghue, and Holstein [8],

$$a'_{\rm DDH}\approx -3.0\ , \tag{36}$$

and with the result of the fit to the weak radiative hyperon decays [15]

$$a_Z=+3.8\ . \tag{37}$$

In view of the inherent uncertainties of the factorization technique all these estimates suggest that the sextet-dominance assumption may be a good approximation for the "factorization" amplitudes of the nonleptonic decays of charmed baryons. A similar view has been expressed by Savage and Springer [27].

## IV. FITS AND CONCLUSIONS

In the preceding section it was pointed out that in the symmetry approach of Körner and collaborators the symmetry structure of the parity-conserving amplitudes does not correspond to the symmetry structure of the standard pole model. This fact plus the appearance of various uncertainties in the reduced matrix elements under consideration (as also discussed in the last section) means that the fits in the symmetry-based approach should be done anew. In the following we will present such a fit. We stress very strongly, however, that, on account of many simplifications involved, the fit should not be considered overly seriously. Rather it should be regarded as purporting the thesis that the present data on charmed baryon decays can be well accommodated in the symmetry-based approach. The fairly limited set of data now available does not warrant a detailed consideration of various symmetry breaking effects. It is only when more data are gathered that the phenomenological determination and discussion of such effects will become possible within the generic framework of this paper. Since at present there are only a few experimental numbers to be fitted we must reduce the number of free parameters of the fit if it is to be meaningful. To this end we make the following simplifying assumptions.

(1) We assume that the connection between the longitudinal and transverse vector meson emission is that given by the quark model, i.e.,

$$\tilde{b}'=\tilde{b}\ ,\qquad B'=B\ ,\qquad m=-M\ . \tag{38}$$

Our fit is based on four fairly accurate data points characterizing the decays with the emission of pseudoscalar mesons and on the not so well-determined branching ratio for the $\Lambda_c^+\to p\phi$ process. Consequently, the above assumption does not affect the predictions of the fit for the decays with the emission of pseudoscalar mesons.

(2) We assume that the sextet-dominance rule holds for the factorization contributions to the parity-conserving amplitudes with both pseudoscalar and vector meson emission, i.e.,

$$M'=-M\ ,\qquad m'=-m\,(=M)\ , \tag{39}$$

as well as for the factorization contributions to the parity-violating amplitudes with pseudoscalar meson emission,

$$g'=-g\ . \tag{40}$$

For the factorization pieces in the parity-violating amplitudes with (transverse) vector-meson emission we use the values extracted from hyperon decays [Eqs. (36) and (37)]. No significant change in the quality of the fit to the

TABLE III. Fit to branching ratios (BR's) and asymmetries.

| $\Lambda_c^+\to$ | BR(%) | Asymmetry |
|---|---|---|
| $\Sigma^+\pi^0$ | 0.43 | −0.76 |
| $\Sigma^+\eta$ | 0.25 | −0.91 |
| $\Sigma^+\eta'$ | 0.05 | +0.72 |
| $\Sigma^0\pi^+$ | 0.43 | −0.76 |
| $\Lambda\pi^+$ | 0.59 | −0.86 |
| $\Xi^0K^+$ | 0.07 | 0.00 |
| $p\bar{K}^0$ | 1.90 | −0.90 |
| $\Sigma^+\rho^0$ | 0.53 | +0.10 |
| $\Sigma^+\omega$ | 0.36 | +0.57 |
| $\Sigma^+\phi$ | 0.04 | −0.87 |
| $\Sigma^0\rho^+$ | 0.53 | +0.10 |
| $\Lambda\rho^+$ | 0.51 | +0.79 |
| $\Xi^0K^{*+}$ | 0.09 | −0.54 |
| $p\bar{K}^{*0}$ | 2.27 | +0.83 |
| $p\phi$ | 0.10 | +0.54 |
| $\Xi_c^0\to$ | BR(%) | Asymmetry |
| $\Xi^0\pi^0$ | 0.29 | −0.99 |
| $\Xi^0\eta$ | 0.04 | −0.32 |
| $\Xi^0\eta'$ | 0.03 | +0.90 |
| $\Xi^-\pi^+$ | 0.88 | −0.78 |
| $\Sigma^0\bar{K}^0$ | 0.05 | −0.89 |
| $\Lambda\bar{K}^0$ | 0.40 | −0.84 |
| $\Sigma^+K^-$ | 0.07 | 0.00 |
| $\Xi^0\rho^0$ | 0.31 | −0.17 |
| $\Xi^0\omega$ | 0.14 | +0.73 |
| $\Xi^0\phi$ | 0.03 | +0.17 |
| $\Xi^-\rho^+$ | 0.64 | +0.80 |
| $\Sigma^0\bar{K}^{*0}$ | 0.12 | +0.62 |
| $\Lambda\bar{K}^{*0}$ | 0.49 | +0.58 |
| $\Sigma^+K^{*-}$ | 0.14 | −0.81 |
| $\Xi_c^+\to$ | BR(%) | Asymmetry |
| $\Xi^0\pi^+$ | 0.31 | +0.65 |
| $\Sigma^+\bar{K}^0$ | 0.28 | +0.68 |
| $\Xi^0\rho^+$ | 1.72 | −0.61 |
| $\Sigma^+\bar{K}^{*0}$ | 2.63 | −0.48 |

$\Lambda_c^+ \to p\phi$ is observed if one accepts sextet dominance for these amplitudes with $\tilde{a} = -\tilde{a}' \approx 13$. To further diminish the number of free parameters we use a single value of $g$ in the range suggested in Eqs. (32) and (33):

$$g = 4.5 \ . \tag{41}$$

The above assumption of sextet dominance seems acceptable in view of the inherent uncertainty of the factorization estimates (29) and (30). Furthermore, it reduces the number of free parameters significantly.

(3) The values of parameters corresponding to the $W$-exchange diagrams [Fig. 1(b)] should be taken from hyperon decays [i.e., $b = -5.0$, $B = -97.5$, Eqs. (6) and (14)]. However, as discussed by Xu and Kamal [16], one expects a mismatch in the baryon wave functions of charmed and noncharmed baryons due to the large mass of the charmed quark. Thus, one expects the $W$-exchange contributions to be smaller than the simple estimates of Eqs. (6) and (14). We take this into account by introducing an overlap parameter $r$ such that the reduced matrix elements $b, B$ are replaced in our formulas by

$$\begin{aligned} b &\to rb \ , \\ B &\to rB \ . \end{aligned} \tag{42}$$

In the following we fit the absolute branching ratios given by Particle Data Group [28]. One has to remember, though, that these are measured relative to the $\Lambda_c^+ \to pK^-\pi^+$ [12]. Thus, the fitted value of $r$ does not correspond to the overlap suppression factor alone—it takes care of the uncertainty in the absolute size of the experimental branching ratios as well.

In summary, we have two parameters $M$ and $r$ and five experimental data points to be fitted. These are the branching ratios of $\Lambda_c^+ \to \Sigma^0\pi^+, \Lambda\pi^+, p\bar{K}^0, p\phi$ and the asymmetry of the $\Lambda_c^+ \to \Lambda\pi^+$ process. The fit achieves $\chi^2 \approx 1.2$ with three degrees of freedom. Apparently, the data are not restrictive enough as yet. Results of the fit are presented in Table III and compared with other papers in Table IV. The fitted values of parameters are

$$\begin{aligned} r_{\text{fit}} &= 0.63 \ , \\ M_{\text{fit}} &= +45 \ . \end{aligned} \tag{43}$$

As expected $r$ is smaller than 1. The fitted value of the reduced matrix element $M$ is about half of that predicted in the factorization approach [cf. Eqs. (29) and (30)]. One has to remember, however, that (1) as it was argued in Sec. III C factorization may be trusted to within a factor of 1.5 or 2 and (2) the uncertainty in the absolute size of the branching ratios has not been taken into account here (as it was the case for the reduced matrix elements $\tilde{b}$ and $B$) since $M$ is a free parameter, anyway.

Although the presented fit is obtained under several simplifying assumptions, it suggests that factorization amplitudes are not as big as one might expect. That factorization prescription seems to give too large contributions has been already noticed by Ebert and Kallies [23]. Furthermore, the fit indicates that nonvanishing contributions from the $W$-exchange diagram [Fig. 1(b)] are needed. Their presence thwarts all attempts to describe nonleptonic decays of charmed baryons with the help of the factorization contribution alone. This conclusion was stressed in Refs. [11,12] as well.

In summary, we have shown that the parity-conserving

TABLE IV. Comparison of model predictions for selected decays.

| | This work | | Ref. [11] | | Ref. [17] | | Experiment | |
|---|---|---|---|---|---|---|---|---|
| $\Lambda_c^+ \to$ | BR | Asym. | BR | Asym. | BR | Asym. | BR | Asym. |
| $\Sigma^+\pi^0$ | 0.43 | −0.76 | 0.31 | +0.71 | 0.72 | +0.83 | | |
| $\Sigma^+\eta$ | 0.25 | −0.91 | 0.15 | +0.33 | | | | |
| $\Sigma^+\eta'$ | 0.05 | +0.72 | 1.22 | −0.45 | | | | |
| $\Sigma^0\pi^+$ | 0.43 | −0.76 | 0.31 | +0.70 | 0.72 | +0.83 | 0.55±0.26 | |
| $\Lambda\pi^+$ | 0.59 | −0.86 | 0.71 | −0.70 | 0.87 | −0.96 | 0.58±0.16 | −1.03±0.29 |
| $\Xi^0K^+$ | 0.07 | 0.00 | 0.25 | 0.00 | | | | |
| $p\bar{K}^0$ | 1.90 | −0.90 | 2.01 | −1.00 | 1.20 | −0.49 | 1.60±0.40 | |
| $\Sigma^+\rho^0$ | 0.53 | +0.10 | 3.0 | | ∼0.1 | +0.10 | | |
| $\Sigma^+\omega$ | 0.35 | +0.57 | 3.8 | | | | | |
| $\Lambda\rho^+$ | 0.51 | +0.79 | 18.2 | | 2.3–2.6 | −0.2 | | |
| $\Xi^0\bar{K}^{*+}$ | 0.09 | −0.54 | 0.11 | | | | | |
| $p\bar{K}^{*0}$ | 2.27 | +0.83 | 2.9 | | 1.8–3.3 | −0.1 | | |
| $p\phi$ | 0.10 | +0.54 | 0.20 | | 0.19 | | 0.13±0.9 | |

| | This work | | Ref. [11] | |
|---|---|---|---|---|
| $\Xi_c^+ \to$ | BR | Asym. | BR | Asym. |
| $\Xi^0\pi^+$ | 0.30 | +0.65 | 2.4 | −0.78 |
| $\Sigma^+\bar{K}^0$ | 0.28 | +0.69 | 4.4 | −1.0 |
| $\Xi^0\rho^+$ | 1.72 | −0.61 | 65.0 | |
| $\Sigma^+\bar{K}^{*0}$ | 2.63 | −0.48 | 1.6 | |

amplitudes in the symmetry approach of Refs. [10–12] do not possess the symmetries of the standard pole model with ground-state intermediate baryons. The proper symmetry structure of these amplitudes that does correspond to this standard assumption of the pole model has been given. In addition, a couple of uncertainties inherent in present approaches to nonleptonic decays of charmed baryons have been identified and discussed. Finally, a fit to the existing data has been carried out.

We would like to stress once again that the fit itself should not be taken overly seriously. There are many unanswered questions concerning the validity of the adopted assumptions. For example, one may worry about the contributions to the parity-conserving amplitudes from intermediate baryons other than the ground-state ones, such as members of the radially excited $(56,\frac{1}{2}^{+*})$ multiplet [29]. Another questionable assumption is that of the SU(4) current algebra used in the description of parity-violating amplitudes: it is only the limit of exact SU(4) that current-algebra and the standard $\frac{1}{2}^{-}$ pole model become equivalent. Although further theoretical studies of various such symmetry-breaking effects in the general framework adopted in this paper are clearly needed, we believe that it will be the experiment that will guide us in our attempts to understand theoretically the nonleptonic decays of charmed baryons.

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