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Parity violation and flavor selection rules in charmed-baryon decays

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The Cabibbo-allowed two body nonleptonic decays of charmed baryons are studied. The kinematics of the vector-meson final state is worked out and flavor-SU(3) and nonet selection rules for all the amplitudes are given. A simple dynamical model is used to predict asymmetries and branching ratios for the decay modes $\Lambda_c^+ \to \Lambda\rho^+$ and $\Lambda_c^+ \to p\overline{K}^{*0}$.

I. INTRODUCTION

Up to now the greater part of theoretical efforts to understand charm decay has been devoted to charmed mesons.[1] Here we begin a study of the decay of charmed baryons in a systematic manner. It is our hope that more and more data on baryon decays will become available in the near future, and that this data will provide a new arena in which to test the standard model. Besides the aspect of flavor selection rules, which applies equally to meson and baryon decays, there is the new aspect of parity violation.

It is instructive to recall the history of the original discovery of parity violation (a good account of which is given in the 1968 review article of Pakvasa and Rosen[2]). The first clue came from the so-called τ-θ puzzle in which the same pseudoscalar meson, what we now call the K meson, appeared to decay into both two-pion and three-pion final states, which have to have opposite parities; but the definitive evidence came from baryon decays in which final states of the same total angular momentum but opposite parities could coherently interfere with one another.

The same holds true of charm particles. Two- and three-pseudoscalar meson final states in the decays of charmed mesons have (to very good approximation) definite parities, and it is not until one examines four-pseudoscalar meson final states that one has an opportunity to look for interference effects indicative of parity violation. Such effects are however expected to be present even in the two-body decays of charmed baryons; here we shall limit ourselves to two-body decays consisting of a baryon plus either a pseudoscalar meson (BP) or a vector meson (BV).

In the case of BP decays, the phenomenology is exactly the same as for hyperon decays.[2] Since the spin of the parent is one-half, the final state can be an admixture of S and P orbital angular momentum waves. Interference between the waves leads to an "up-down" asymmetry for the decay meson with respect to the spin direction of the parent baryon, and to a longitudinal polarization for the daughter baryon. If time-reversal invariance also breaks down, the daughter baryon will have a transverse component of polarization in its production plane. Because the charmed-baryon parent will often produce a hyperon as its daughter, we can use the known parameters of hyperon decay to analyze the polarization of the daughter.

For BV decays the analysis is the same in principle, but different in practice because of the unit spin of the daughter meson. The orbital angular momentum of the final state is now an admixture of S, P, and D waves; and there are moreover two independent P-wave amplitudes: one associated with the singlet combination of parent and daughter baryon spins and the other with the triplet. The basic effects of the interference between the parity-violating S and D waves on the one hand with the parity-conserving P-wave amplitude on the other give rise to asymmetries for the daughter particles with respect to the spin of the parent, and to longitudinal polarizations. We give the details of this analysis in a separate section below.

The presence of two or more orbital angular momentum waves in the final states of charmed-baryon decay enriches the predictive power of flavor selection rules. Where these rules lead to sum rules or proportionality relations between the amplitudes for different decay modes, the relations will hold separately for each angular momentum wave. This in turn will then lead to relations not only between the decay rates for the modes involved, but also to relations between the asymmetry and polarization parameters of the modes.

An instructive example of this richness comes from the well-known prediction of the $\Delta T = \frac{1}{2}$ rule for Λ-hyperon decay:[2]

$$A(\Lambda \to p\pi^-) = -\sqrt{2}A(\Lambda \to n\pi^0) \ . \qquad (1)$$

This relation holds for both the S-wave and the P-wave amplitudes and it implies that the rate for the $p\pi^-$ mode is twice that for the $n\pi^0$ mode, but that all the asymmetry parameters for the two modes are equal to one another. In the case of the Σ hyperon, there is a well-known triangular relation which leads to similar types of prediction as long as final-state interactions are neglected.

Whereas the neglect of final-state interactions appears to be a good approximation in hyperon decay, it may not be so for charm decay. In the case of charmed mesons, it is certainly necessary to include large final-state effects in order to satisfy the isospin selection rules of the standard model; moreover, this can be understood as resulting from the existence of strange-meson resonances in the neighborhood of the charmed-meson mass range. To the extent that there may be strange-baryon resonances in the mass range of the stable charmed baryons, there might be large final-state effects in the corresponding decays, which are neglected here.

The three stable charmed baryons of spin $\frac{1}{2}$ form an antitriplet $\mathbf{3}^*$ with respect to flavor SU(3), just as do the charmed mesons. Corresponding to the D^+, D^0, D_s are the $\Xi_c^+, \Xi_c^0, \Lambda_c^+$, respectively, with masses of 2471, 2460, and 2285 MeV.[3] These baryons are formed by replacing a strange quark in the hyperons Ξ^0, Ξ^-, Λ^0 with a charmed quark and antisymmetrizing the two remaining light quarks in flavor space. One obvious mode for the Cabibbo-allowed decay of the charmed baryons is into the corresponding hyperon plus one or more nonstrange mesons; other modes include nonstrange baryons plus strange mesons.

In the SU(3)-flavor space generated by u, d, and s quarks, the effective Hamiltonian for Cabibbo-allowed charm-decay transforms as an admixture of the $\mathbf{6}^*$ and $\mathbf{15}$ representations. As we have done in the case of charmed mesons,[4] so here we construct the effective Hamiltonian by combining the final-state baryon and meson multiplets into definite representations of SU(3), and then combining these representations with the charmed-baryon triplet to form $\mathbf{6}^*$ and $\mathbf{15}$ tensors. The spin-$\frac{1}{2}$ baryon octet combines with the meson nonet to form the usual representations $\mathbf{1}, \mathbf{8}_F, \mathbf{8}_D, \mathbf{10}, \mathbf{10}^*$, and $\mathbf{27}$: the singlet cannot be engendered by the Hamiltonian, the octets can arise from both the $\mathbf{6}^*$ and $\mathbf{15}$, the $\mathbf{10}$ and $\mathbf{27}$ from the $\mathbf{15}$ alone, and the $\mathbf{10}^*$ from the $\mathbf{6}^*$ alone. Nonet symmetry works well for vector mesons, but not for pseudoscalar mesons, and so in the pseudoscalar case we must add additional terms to the Hamiltonian which are constructed solely from the SU(3)-singlet part of the nonet. From this construction we can extract sum rules amongst amplitudes which will hold for all angular momentum waves. We shall explore the implications of these sum rules in the following sections, and we shall also discuss the implications of making various dynamical assumptions as well. But before beginning this part of our work we turn to the phenomenological description of parity violation in the decay amplitudes themselves.

II. PARITY VIOLATION IN BP AND BV DECAYS

The description of parity violation in BP decays is well-known,[2] but it is useful to recall it here. Since the initial and final baryons have spin $\frac{1}{2}$, the pseudoscalar meson is in a superposition of S- and P-wave angular momentum states. In the rest-frame of the parent charmed baryon, the decay can be described by the transition matrix

$$M = \Xi_f^\dagger (S + P\sigma \cdot \mathbf{q}) \Xi_i \ , \tag{2}$$

where S and P denote the amplitudes for the corresponding waves and $\mathbf{q}$ is a unit vector in the direction of the daughter baryon; the Ξ_m are the two-spinors for the initial and final baryons. The decay rate for parent and daughter baryons with spins in directions $\mathbf{N}$ and $\mathbf{n}$, respectively, can then be written

$$\begin{aligned} d\Gamma = \frac{\Gamma}{8\pi} \, [1 + \alpha \mathbf{q} \cdot (\mathbf{N} + \mathbf{n}) + (\mathbf{N} \cdot \mathbf{q})(\mathbf{n} \cdot \mathbf{q}) \\ + \beta \mathbf{N} \times \mathbf{q} \cdot \mathbf{n} + \gamma \mathbf{q} \times (\mathbf{N} \times \mathbf{q}) \cdot \mathbf{n}] d\Omega \ , \end{aligned} \tag{3}$$

$$\Gamma = \frac{Q}{8\pi M_c^2} (| S |^2 + | P |^2) \ ,$$

where M_c is the mass of the charmed baryon and Q is the momentum of the decay products in its rest frame.

In the expression for the rate, the various asymmetry parameters are

$$\begin{aligned} \alpha &= 2 \frac{\mathrm{Re} S^* P}{| S |^2 + | P |^2} \ , \\ \beta &= 2 \frac{\mathrm{Im} S^* P}{| S |^2 + | P |^2} \ , \\ \gamma &= \frac{| S |^2 - | P |^2}{| S |^2 + | P |^2} \ . \end{aligned} \tag{4}$$

The parameter α measures both the "up-down" asymmetry and the longitudinal polarization of the daughter baryon, while the parameters β and γ measure transverse components of the daughter polarization. When CP is conserved and final-state interactions are negligible, β vanishes; when final-state interactions are not negligible, the S and P amplitudes acquire the appropriate strong interaction phase shifts by virtue of the Watson theorem and β does not vanish. Therefore, since we expect large final-state effects in charmed-baryon decay, we cannot use the transverse polarization measured by β as an indication of CP violation, without a detailed knowledge of the final-state effects.

If we write the effective Lagrangian density for BP decay in the relativistically covariant form,

$$L_{\mathrm{eff}} = -i [\bar{\psi}_f (A + B\gamma_5) \psi_i] \phi_P \ , \tag{5}$$

then we can relate the dimensionless constants A and B to the parameters S and P of Eq. (2). We find that

$$\begin{aligned} S &= \sqrt{2M_c(E + M_f)} A \ , \\ P &= \sqrt{2M_c(E - M_f)} B \ , \end{aligned} \tag{6}$$

where M_f is the mass of the daughter baryon and E is its total energy in the rest frame of the parent.

Let us turn to BV decay and analyze it in a similar manner. Because the daughter meson now has spin 1, it is in a superposition of S, P, and D waves. The P-wave amplitudes have the additional complication that they may be associated with both spin flip and spin nonflip for the baryons and hence there are two independent amplitudes of this type. The S and D waves have only one independent amplitude apiece. In the rest frame of the decaying charmed baryon, the transition matrix can be written as

$$M = \{\Xi_f^\dagger[S\boldsymbol{\sigma} + P_1\mathbf{p} + iP_2\mathbf{p}\times\boldsymbol{\sigma} + D(\boldsymbol{\sigma}\cdot\mathbf{p})\mathbf{p}]\cdot\boldsymbol{\epsilon}\Xi_i\} \ , \tag{7}$$

where $\boldsymbol{\epsilon}$ is the polarization three-vector of the vector boson and $\mathbf{p}$ is a unit vector in the direction of its momentum. To understand the physical meaning of the amplitudes (S, P_1, P_2, D), it is instructive to construct the transverse and longitudinal forms of $\boldsymbol{\epsilon}$.

We construct a mutually perpendicular triad of unit vectors by taking $\mathbf{p}$, and $\mathbf{s}$, a unit vector perpendicular to $\mathbf{p}$, and the vector product $\mathbf{p}\times\mathbf{s}$, which we denote by $\mathbf{r}$. The corresponding polarization three-vectors are then given by

$$\begin{aligned} \boldsymbol{\epsilon}_p &= \mathbf{p} \ , \\ \boldsymbol{\epsilon}_s &= \mathbf{s} \ , \\ \boldsymbol{\epsilon}_r &= \mathbf{p}\times\mathbf{s} \ . \end{aligned} \tag{8}$$

They can be arranged as eigenvectors of the helicity operator $\mathbf{J}\cdot\mathbf{p}$ with eigenvalues (+1, −1, 0), respectively:

$$\begin{aligned} -\frac{1}{\sqrt{2}}(\boldsymbol{\epsilon}_s + i\boldsymbol{\epsilon}_r) &= (\mathbf{J}\cdot\mathbf{p} = +1) \ , \\ \frac{1}{\sqrt{2}}(\boldsymbol{\epsilon}_s - i\boldsymbol{\epsilon}_r) &= (\mathbf{J}\cdot\mathbf{p} = -1) \ , \\ \boldsymbol{\epsilon}_p &= (\mathbf{J}\cdot\mathbf{p} = 0) \ . \end{aligned} \tag{9}$$

Using these expressions in the general transition matrix of Eq. (7), we can write down the transition matrices for specific helicity states of the vector meson:

$$\begin{aligned} M(+1) &= \frac{P_2 - S}{\sqrt{2}}\{\Xi_f^\dagger[\boldsymbol{\sigma}\cdot(\mathbf{s}+i\mathbf{r})]\Xi_i\} \ , \\ M(-1) &= \frac{P_2 + S}{\sqrt{2}}\{\Xi_f^\dagger[\boldsymbol{\sigma}\cdot(\mathbf{s}-i\mathbf{r})]\Xi_i\} \ , \\ M(0) &= \{\Xi_f^\dagger[(S+D)\sigma\cdot\mathbf{p} + P_1]\Xi_i\} \ . \end{aligned} \tag{10}$$

Thus we find that $(P_2 \pm S)$ are the amplitudes for a transversely polarized vector meson in the final state, and that $(S+D)$ is the amplitude for longitudinal polarization associated with baryonic spin flip, while P_1 is associated with spin nonflip.

When we calculate the square modulus of the transition matrix and sum over the helicities of the vector meson, we obtain

$$|M|^2 = |M(+1)|^2 + |M(-1)|^2 + \frac{E_v^2}{m_v^2}|M(0)|^2 \ , \tag{11}$$

The factor E_v^2/m_v^2, where E_v and m_v are the energy and mass of the vector meson, respectively, is a relativistic correction coming from the completeness of the polarization *four-vector* for vector mesons. As in the case of BP decays, so here we can sum this expression over the spins of the baryons. For the parent charmed baryon and the daughter baryon with spins pointing in the directions $\mathbf{n}_c$ and $\mathbf{n}_d$, respectively, we find that

$$\begin{aligned} |M|^2 = {} & \{[|S|^2 + |P_2|^2][1 - (\mathbf{n_d}\cdot\mathbf{p})(\mathbf{n_c}\cdot\mathbf{p})] - 2\,\mathrm{Re}(S^*P_2)\mathbf{p}\cdot(\mathbf{n_d}-\mathbf{n_c})\} \\ & + \frac{E_v^2}{2m_v^2}\{|S+D|^2\,[1 - \mathbf{n_d}\cdot\mathbf{n_c} + 2(\mathbf{n_d}\cdot\mathbf{p})(\mathbf{n_c}\cdot\mathbf{p})] \\ & \qquad + |P_1|^2\,(1+\mathbf{n_d}\cdot\mathbf{n_c}) + 2\,\mathrm{Re}(S+D)^*P_1\mathbf{p}\cdot(\mathbf{n_d}+\mathbf{n_c}) - 2\,\mathrm{Im}(S+D)^*P_1\mathbf{p}\times\mathbf{n_d}\cdot\mathbf{n_c}\} \ . \end{aligned} \tag{12}$$

The differential and total decay rates are then given by

$$\begin{aligned} d\Gamma &= \frac{E_d + M_d}{M_c(4\pi)^2}\,|\mathbf{P}_v|\,|\mathcal{M}|^2\,d\Omega \ , \\ \Gamma &= \frac{1}{8\pi}\frac{E_d+M_d}{M_c}\,|\mathbf{P}_v|\,\left((|S+D|^2 + |P_1|^2)\frac{E_v^2}{m_v^2} + 2(|S|^2 + |P_2|^2)\right) \end{aligned} \tag{13}$$

where M_d and E_d are the mass and energy of the daughter baryon in the charmed-baryon rest frame, P_v is the momentum of the vector meson, and M_c and Ω are the parent mass and the decay solid angle.

From the expressions in Eqs. (11) and (12) we can now calculate various asymmetries in the distribution of the decay products. The "up-down" asymmetry of the vector meson with respect to the charmed baryon spin is

$$A(\text{up-down}) = (1 + \alpha \mathbf{p} \cdot \mathbf{n}_c) ,$$
$$\alpha = \frac{2\,\mathrm{Re}(S+D)^* P_1 E_v^2 + 4\,\mathrm{Re}S^* P_2 m_v^2}{(|S+D|^2 + |P_1|^2)E_v^2 + 2(|S|^2 + |P_2|^2)m_v^2} , \tag{14}$$

and the longitudinal-polarization distribution of the decay baryon for unpolarized parent is given by

$$A(\text{long}) = (1 + \alpha' \mathbf{p} \cdot \mathbf{n}_d) ,$$
$$\alpha' = \frac{2\,\mathrm{Re}(S+D)^* P_1 E_v^2 - 4\,\mathrm{Re}S^* P_2 m_v^2}{(|S+D|^2 + |P_1|^2)E_v^2 + 2(|S|^2 + |P_2|^2)m_v^2} . \tag{15}$$

More generally the distribution with respect to the polarization vector of the decay baryon can be written as

$$A(\mathbf{N}_d) = d\Omega(1 + \alpha \mathbf{p} \cdot \mathbf{n}_c)(1 + \mathbf{n}_d \cdot \mathbf{N}) ,$$
$$\mathbf{N} = \frac{(\alpha' + \gamma' \mathbf{p} \cdot \mathbf{n}_c)\mathbf{p} + \beta \mathbf{p} \times \mathbf{n}_c + \gamma \mathbf{p} \times (\mathbf{p} \times \mathbf{n}_c)}{1 + \alpha \mathbf{p} \cdot \mathbf{n}_c} , \tag{16}$$

where the coefficients are

$$\beta = \frac{2\,\mathrm{Im}(S+D)^* P_1 E_v^2}{(|S+D|^2 + |P_1|^2)E_v^2 + 2(|S|^2 + |P_2|^2)m_v^2} ,$$
$$\gamma = \frac{(|S+D|^2 - |P_1|^2)E_v^2}{(|S+D|^2 + |P_1|^2)E_v^2 + 2(|S|^2 + |P_2|^2)m_v^2} , \tag{17}$$
$$\gamma' = \frac{(|S+D|^2 + |P_1|^2)E_v^2 - 2(|S|^2 + |P_2|^2)m_v^2}{(|S+D|^2 + |P_1|^2)E_v^2 + 2(|S|^2 + |P_2|^2)m_v^2} .$$

The transition matrix in Eq. (7) can be thought of either as the nonrelativistic limit of a relativistically covariant expression, or as the reduction of the covariant expression to the rest-frame of the decaying charmed baryon. In either case it is important to establish the covariant transition matrix and the relationship between its independent constants and those in Eq. (7). On grounds of covariance alone, we can use scalar, vector, and tensor couplings between the baryons in the effective decay Lagrangian. Since parity is not conserved, this yields six constants; but only four of them are independent because we can use the Dirac equation to reduce one type of coupling to a linear combination of the other two. For the sake of completeness, however, we shall write down all of the terms and give the connections with the constants in the transition matrix of Eq. (7).

In terms of Dirac spinors, the transition matrix is

$$\begin{aligned} L_{\text{eff}} = {} & [\bar{\psi}_d(x + y\gamma_5)\psi_c](p_c + p_d) \cdot \epsilon \\ & - i[\bar{\psi}_d \gamma_\mu \epsilon_\mu (a + b\gamma_5)\psi_c] \\ & + i[\bar{\psi}_d \sigma_{\mu\nu}(\epsilon_\mu p_\nu - \epsilon_\nu p_\mu)(f + g\gamma_5)\psi_c] \end{aligned} \tag{18}$$

The (p_c, p_d, p) are the four-vectors representing the four-momenta of the parent baryon, the daughter baryon, and the vector meson, respectively, and ϵ is the four-vector polarization of the vector meson. Our notation $p \cdot \epsilon$ denotes the four-scalar product of the two four-vectors. The connections with the parameters in Eq. (7) can now be determined:

$$\begin{aligned} S &= b + 2g(m_c - m_d) , \\ P_1 &= \frac{p}{E_v}\left(\frac{a(m_c + m_d) - 2fm_v^2}{E_d + m_d} - 2xm_c\right) , \\ P_2 &= \frac{p}{E_d + m_d}[-a + 2f(m_c + m_d)] , \\ D &= \frac{p^2}{E_v(E_d + m_d)}[b - 2g(m_c + m_d) + 2ym_c] . \end{aligned} \tag{19}$$

Here p denotes the magnitude of the three-momentum of the vector meson in the parent baryon rest-frame and E_v, E_d denote the energies of the meson and daughter baryon in the same frame.

III. FLAVOR-SU(3) SELECTION RULES

In the standard model,[1] the effective interaction for Cabibbo-allowed charm-decay transforms as an admixture of the $\mathbf{6}^*$ and $\mathbf{15}$ representations of flavor SU(3). We construct it by first combining the spin-$\frac{1}{2}$ baryon octet with the meson multiplet to form specific representations of SU(3), and then combining these representations with the conjugate of the charmed baryon antitriplet to form the overall $\mathbf{6}^*$ and $\mathbf{15}$ tensors.

For the meson multiplets we make use of nonet symmetry,[5] but with specific limitations. Nonet symmetry is an additional assumption about vector and pseudoscalar mesons, inspired by the almost ideal mixing of the ϕ and ω mesons. It has the effect of relating the coupling constants associated with SU(3)-singlet mesons to those associated with the corresponding octets. Thus it leads to fewer independent amplitudes than are encountered in the most general SU(3) description of charm decay.

In this picture, the product of two nonets yields the same set of representations as the product of two octets, namely one singlet, two octets, one $\mathbf{10}$, one $\mathbf{10}^*$, and one $\mathbf{27}$. Where the nonet symmetry breaks down, we pick up three additional representations, a singlet from the product of the singlet components of the nonets and two octets from the product of the singlet in one nonet times the octet component of the other. As far as is known at this time, nonet symmetry works well for vector mesons but not for pseudoscalar ones. We shall therefore write the effective Hamiltonians in nonet symmetry form, but include an explicit breaking term for BP decays.

The specific forms of the relevant final-state tensor products of the baryon octet times meson nonet can be readily adapted from our earlier papers on charmed-meson decays. The effective Hamiltonian in the nonet symmetry case is

TABLE I. Nonet-symmetry amplitudes for $B_c \to BP$.

Mode	$Sy_6(L)$	$A_6(L)$	$T_6(L)$	$Sy_{15}(L)$	$A_{15}(L)$	$T_{15}(L)$	$T'_{15}(L)$
$\Xi_c^+ \to \Xi^0\pi^+$	0	0	-2	0	0	2	0
$\Xi_c^0 \to \Xi^-\pi^+$	1	1	$-2/3$	1	1	$4/5$	$-2/3$
$\Xi_c^0 \to \Xi^0\pi^0$	$-1/\sqrt{2}$	$-1/\sqrt{2}$	$-(\sqrt{8})/3$	$-1/\sqrt{2}$	$-1/\sqrt{2}$	$(3\sqrt{2})/5$	$(\sqrt{2})/3$
$\Xi_c^+ \to \Sigma^+\bar{K}^0$	0	0	2	0	0	2	0
$\Xi_c^0 \to \Sigma^+\bar{K}^-$	1	-1	$2/3$	1	-1	$4/5$	$2/3$
$\Xi_c^0 \to \Sigma^0\bar{K}^0$	$-1/\sqrt{2}$	$1/\sqrt{2}$	$(\sqrt{8})/3$	$-1/\sqrt{2}$	$1/\sqrt{2}$	$(3\sqrt{2})/5$	$-(\sqrt{2})/3$
$\Lambda_c^+ \to \Xi^0 K^+$	-1	$+1$	$-2/3$	1	-1	$4/5$	$2/3$
$\Lambda_c^+ \to p\bar{K}^0$	-1	-1	$2/3$	1	1	$4/5$	$-2/3$
$\Lambda_c^+ \to \Sigma^+\eta_8$	$-\sqrt{(2/3)}$	0	$-\sqrt{(2/3)}$	$\sqrt{(2/3)}$	0	$-2(\sqrt{6})/5$	$-\sqrt{(2/3)}$
$\Lambda_c^+ \to \Sigma^+\eta_1$	$-2/\sqrt{3}$	0	0	$2/\sqrt{3}$	0	0	0
$\Lambda_c^+ \to \lambda\pi^+$	$-\sqrt{2}/\sqrt{3}$	0	$\sqrt{2}/\sqrt{3}$	$\sqrt{2}/\sqrt{3}$	0	$-(2\sqrt{6})/5$	$\sqrt{2}/\sqrt{3}$
$\Xi_c^0 \to \Xi^0\eta_8$	$-1/\sqrt{6}$	$\sqrt{(3/2)}$	0	$-1/\sqrt{6}$	$\sqrt{(3/2)}$	$(\sqrt{6})/5$	$\sqrt{(2/3)}$
$\Xi_c^0 \to \Xi^0\eta_1$	$2/\sqrt{3}$	0	0	$2/\sqrt{3}$	0	0	0
$\Xi_c^0 \to \lambda\bar{K}^0$	$-1/\sqrt{6}$	$-\sqrt{3}/\sqrt{2}$	0	$-1/\sqrt{6}$	$-\sqrt{3}/\sqrt{2}$	$(\sqrt{6})/5$	$-\sqrt{2}/\sqrt{3}$
$\Lambda_c^+ \to \Sigma^+\pi^0$	0	$\sqrt{2}$	$(\sqrt{2})/3$	0	$-\sqrt{2}$	0	$-(\sqrt{2})/3$
$\Lambda_c^+ \to \Sigma^0\pi^+\rho^0$	0	$-\sqrt{2}$	$-(\sqrt{2})/3$	0	$\sqrt{2}$	0	$(\sqrt{2})/3$

$$H_{\text{nonet}} = Sy_6(L)[6^*;8_{Sy}] + A_6(L)[6^*;8_A] + T_6(L)[6^*;10^*] + Sy_{15}(L)[15;8_{Sy}] + A_{15}(L)[15;8_A] + T'_{15}(L)[15;10] + T_{15}(L)[15;27]\ , \tag{20}$$

where numbers are used to denote representations and the subscripts Sy and A refer to the symmetric and antisymmetric octet products of the baryon and meson multiplets, respectively, and the notation $[X;Y]$ denotes the overall representation X constructed from the charmed-baryon triplet and the representation Y of the final state. The $(S,A,T,T')_{6,15}(L)$ are the invariant amplitudes (coupling constants) for the relevant terms in the effective Hamiltonian with orbital wave (L). The nonet-symmetry-breaking terms, which we assume to apply only to pseudoscalar-meson final states, are written as

$$H_{\text{breaking}} = B_6(L)[6^*;1_P,8_B] + B_{15}(L)[15;1_P,8_B]\ , \tag{21}$$

where the notation indicates that the terms are constructed from the pseudoscalar singlet (1_P). The amplitudes for $B_c \to BP$ derived from this analysis are given in Tables I and II.

Numerous sum rules relating the amplitudes for different decay modes can be extracted from these tables. Some of these rules follow from the isospin selection rule $\Delta T{=}1$, others from the U-spin rule $\Delta U{=}1$, and others require the full power of SU(3). One simple example is the relation for the BP decay modes,

$$A(L)(\Lambda_c^+ \to \Sigma^+\pi^0) = -A(L)(\Lambda_c^+ \to \Sigma^0\pi^+)\ , \tag{22}$$

and a similar relation for the corresponding BV modes (using Table III). This happens to be a statement that the final state of the Λ_c^+ decay must be a pure isovector and it will hold for all orbital waves L. It implies that the partial rates for the two decay modes and all the asymmetry parameters must be equal to one another.

More complicated relations hold for the BP decays of the Ξ_c doublet and the BV counterparts: for example,

$$A(L)(\Xi_c^0 \to \Xi^-\pi^+) + \sqrt{2}A(L)(\Xi_c^0 \to \Xi^0\pi^0) = A(L)(\Xi_c^+ \to \Xi^0\pi^+)\ . \tag{23}$$

In this case the final state is an admixture of isospins $\frac{1}{2}$ and $\frac{3}{2}$, and hence we can extract only one relation amongst the three amplitudes for each L wave. Because of final-state interactions we must treat each $A(L)$ as a complex number, and hence Eq. (23) represents a triangle in a multidimensional space. The implications of this will depend on the orientation of the triangle. Many of these sum rules were written down some time ago.[6]

TABLE II. Nonet-symmetry breaking amplitudes for $B_c \to BP$. The amplitudes for $B_c \to BV$ satisfy nonet symmetry and are given in Table III.

Mode	$B_6(L)$	$B_{15}(L)$
$\Lambda_c^+ \to \sigma^+\eta_1$	-1	1
$\Xi_c^0 \to \Xi^0\eta_1$	1	1

IV. MODEL FOR $B_C \to BP$ AND $B_C \to BV$

We attempt to predict the decay rates and asymmetry parameters of charmed-baryon decays with a concrete model for Λ_c^+ decays to exclusive two-body channels where some data are now available. We concentrate particularly on the baryon-pseudoscalar channels $\Lambda_c^+ \to \Lambda\pi^+, p\overline{K}^0$ and the baryon-vector meson channels $\Lambda_c^+ \to \Lambda\rho^+, p\overline{K}^{*0}$.

In general the nonleptonic hyperon decay amplitudes

TABLE III. Nonet-symmetry amplitudes for $B_c \to BV$.

Mode	$Sy_6(L)$	$A_6(L)$	$T_6(L)$	$Sy_{15}(L)$	$A_{15}(L)$	$T_{15}(L)$	$T'_{15}(L)$
$\Xi_c^+ \to \Sigma^+ \bar{K}^{*0}$	0	0	-2	0	0	2	0
$\Xi_c^0 \to \Sigma^+ \bar{K}^{*-}$	1	1	$-2/3$	1	1	$4/5$	$-2/3$
$\Xi_c^0 \to \Sigma^0 \bar{K}^{*0}$	$-1/\sqrt{2}$	$-1/\sqrt{2}$	$(\sqrt{8})/3$	$-1/\sqrt{2}$	$-1/\sqrt{2}$	$(3\sqrt{2})/5$	$(\sqrt{2})/3$
$\Xi_c^+ \to \Xi^0 \rho^+$	0	0	2	0	0	2	0
$\Xi_c^0 \to \Xi^- \rho^+$	1	-1	$2/3$	1	-1	$4/5$	$2/3$
$\Xi_c^0 \to \Xi^0 \rho^0$	$-1/\sqrt{2}$	$1/\sqrt{2}$	$(\sqrt{8})/3$	$-1/\sqrt{2}$	$1/\sqrt{2}$	$(3\sqrt{2})/5$	$-(\sqrt{2})/3$
$\Lambda_c^+ \to p \bar{K}^{*0}$	-1	$+1$	$-2/3$	1	-1	$4/5$	$2/3$
$\Lambda_c^+ \to \Xi^0 K^{*+}$	-1	-1	$2/3$	1	1	$4/5$	$-2/3$
$\Lambda_c^+ \to \Lambda \rho^+$	$-\sqrt{(2/3)}$	0	$-\sqrt{(2/3)}$	$\sqrt{(2/3)}$	0	$-2(\sqrt{6})/5$	$-\sqrt{(2/3)}$
$\Lambda_c^+ \to \Sigma^+ \phi$	0	0	$-2/3$	0	0	$4/5$	$-2/3$
$\Lambda_c^+ \to \Sigma^+ \omega$	$-\sqrt{2}$	0	$(\sqrt{2})/3$	$\sqrt{2}$	0	$-(2\sqrt{2})/5$	$(\sqrt{2})/3$
$\Xi_c^0 \to \Lambda \bar{K}^{*0}$	$-1/\sqrt{6}$	$\sqrt{(3/2)}$	0	$-1/\sqrt{6}$	$\sqrt{(3/2)}$	$(\sqrt{6})/5$	$\sqrt{(2/3)}$
$\Xi_c^0 \to \Xi^0 \omega$	$1/\sqrt{2}$	$-1/\sqrt{2}$	0	$1/\sqrt{2}$	$-1/\sqrt{2}$	$(\sqrt{2})/5$	$-(\sqrt{2})/3$
$\Xi_c^0 \to \Xi^0 \phi$	1	1	0	1	1	$-2/5$	$2/3$
$\Lambda_c^+ \to \Sigma^0 \rho^+$	0	$\sqrt{2}$	$(\sqrt{2})/3$	0	$-\sqrt{2}$	0	$-(\sqrt{2})/3$
$\Lambda_c^+ \to \Sigma^+ \rho^0$	0	$-\sqrt{2}$	$-(\sqrt{2})/3$	0	$\sqrt{2}$	0	$(\sqrt{2})/3$

will involve contributions from the factorization approximation, current-algebra (soft-meson) term,[7] and the pole contribution. However in charmed-meson nonleptonic decays, factorization approximation alone gives a good account of exclusive 2-body decay modes.[8] We therefore exploit the same technique for charmed baryon decays, except that the equal-time-commutator (ETC) term is included in addition for the parity violating (PV) piece.

The effective $\Delta C = 1, \Delta S = 1$ Hamiltonian including short-distance QCD corrections is[9]

$$H_w = \frac{G_F}{\sqrt{2}} \cos^2\theta_C \; [f_\Lambda \bar{s}\gamma_\mu(1+\gamma_5)c \; \bar{u}\gamma_\mu(1+\gamma_5)d + f_p \; \bar{u}\gamma_\mu(1+\gamma_5)c \; \bar{s}\gamma_\mu(1+\gamma_5)d] \; , \tag{24}$$

where f_Λ and f_p are dimensionless constants in the range 1.1 to 1.4 and 0.4 to 0.7, respectively.

The factorization approximation for $B_c \to BM$ (i.e., charmed-baryon decay to baryon B and meson M) is just the statement that

$$\langle BM \mid H_w \mid B_c \rangle = f_B \left(\frac{G_F \cos^2\theta_C}{\sqrt{2}} \right) \langle B \mid J_\mu \mid B_c \rangle \langle M \mid J_\mu \mid 0 \rangle \; . \tag{25}$$

Together with the use of the ETC term for the PV piece, we have the following amplitudes for $\Lambda_c^+ \to \Lambda\pi^+$ and $\Lambda_c^+ \to p\overline{K}^0$:

$$\mathcal{M}(\Lambda_c^+ \to \Lambda\pi^+) = f_\Lambda f_\pi \Big(\frac{G_F \cos^2\theta_C}{\sqrt{2}} \Big) \langle \Lambda \mid \bar{s}\gamma_\mu(1+\gamma_5)c \mid \Lambda_c^+ \rangle q_\mu - \frac{1}{f_\pi} \langle \Lambda \mid [Q^{\pi^+}, H_w^{\rm PV}] \mid \Lambda_c^+ \rangle \tag{26}$$

and

$$\mathcal{M}(\Lambda_c^+ \to p\overline{K}^0) = f_p f_K \left(\frac{G_F \cos^2\theta_C}{\sqrt{2}} \right) \langle p \mid \bar{u}\gamma_\mu(1+\gamma_5)c \mid \Lambda_c^+ \rangle q_\mu - \frac{1}{f_K} \langle p \mid [Q^{\overline{K}^0}, H_w^{\rm PV}] \mid \Lambda_c^+ \rangle \tag{27}$$

where q_μ is the meson four-momentum. However $[Q^{\pi^+}, H_w^{\rm PV}] = 0$ and by SU(3) rotation

$$\langle p \mid [Q^{\overline{K}^0}, H_w^{\rm PV}] \mid \Lambda_c^+ \rangle = -\frac{1}{\sqrt{2}} \langle \Sigma^+ \mid H_w^{\rm PC} \mid \Lambda_c^+ \rangle \; . \tag{28}$$

In the limit of SU(4) symmetry, right-hand side (RHS) of (28) becomes

$$-\left(\frac{1}{\sqrt{2}} \frac{1}{f_K} \right) \cot\theta_C \frac{1}{\sqrt{6}} R \langle p \mid H_w^{\rm PC} \mid \Sigma^+ \rangle \tag{29}$$

where R~1, and f_K=1.28 f_π, f_π ≅0.130 GeV. The matrix element $\langle p \mid H_w^{\rm PC} \mid \Sigma^+ \rangle$ has been estimated in a number of ways, e.g., by model calculations,[10] by fitting S-wave hyperon decay and by fitting P-wave hyperon decay. For numerical illustration, we shall take the value

$$\langle p \mid H_w^{\rm PC} \mid \Sigma^+ \rangle = -1.2 \times 10^{-4} \text{ MeV} \tag{30}$$

obtained from P-wave hyperon decay which is in fair agreement with the model estimate.[10]

For the baryonic matrix element of the weak current $\langle B \mid J_\mu \mid B_c \rangle$, we use the fit to semileptonic decays $\Lambda_c^+ \to \Lambda e^+ \nu$ given recently by Avila-Aoki *et al.* and

Perez-Marcial *et al.*[11] With the parametrization

$$\langle \Lambda \mid \bar{s}\gamma_\mu(1+\gamma_5)c \mid \Lambda_c^+\rangle = \bar{u}_\Lambda \left(-i\gamma_\mu(f_1+g_1\gamma_5) + i\sigma_{\mu\nu}q_\nu \frac{f_2+g_2\gamma_5}{m_1} + q_\mu \frac{f_3+g_3\gamma_5}{m_1}\right) u_{\Lambda_c} , \quad (31)$$

where $m_1 = M_{\Lambda_c^+}$. The best fit to the experimental semileptonic branching ratio is obtained when the matrix element is evaluated with the bag model and dipole form factors. To wit

$$f_i(q^2) = \frac{f_i(0)}{(1-q^2/m^{*2})^2} , \quad g_i(q^2) = \frac{g_i(0)}{(1-q^2/m_A^2)^2} \quad (32)$$

with m^*=2.112 GeV, m_A=2.5 GeV for $\Lambda_c^+ \xrightarrow{\pi} \Lambda$ (and m^*=2.01 GeV, m_A=2.4 GeV for $\Lambda_c^+ \xrightarrow{\overline{K}^0} p$). For $f_i(0)$ and $g_i(0)$ in the bag model with n=2 (dipole) form factor, we have

$f_1(0)$	$f_2(0)$	$f_3(0)$	$g_1(0)$	$g_2(0)$	$g_3(0)$
0.46	0.19	0	0.5	−0.05	−0.44

(33)

The baryonic matrix element $\langle p \mid J_\mu \mid \Lambda_c^+\rangle$ has form factors f_i' and g_i' (i=1,2,3) defined analogously to Eq. (31). They are related to f_i, and g_i by an SU(3) factor, viz.,

$$f_i'(0) = \sqrt{\tfrac{3}{2}} f_i(0) , \quad g_i'(0) = \sqrt{\tfrac{3}{2}} g_i(0) \quad (34)$$

and $f_i'(q^2)$, $g_i'(q^2)$ are related to $f_i'(0)$, $g_i'(0)$ via relations analogous to (32) with m^*=2.01 GeV and m_A=2.4 GeV.

With the procedure outlined above [(28)–(34)], the amplitude for $\mathcal{M}(\Lambda_c^+ \to \Lambda\pi^+)$ given by (26) with vanishing ETC and $\sigma_{\mu\nu}q_\nu$ terms as well as negligible contributions from the f_3, g_3 term in (31), yields

$$\mathcal{M}(\Lambda_c^+ \to \Lambda\pi^+) = i\, f_\Lambda f_\pi \left(\frac{G_F \cos^2\theta_C}{\sqrt{2}}\right) \bar{u}_\Lambda[-0.538 \text{ GeV} + (1.7 \text{ GeV})\gamma_5]u_{\Lambda_c} . \quad (35)$$

The amplitude $\mathcal{M}(\Lambda_c^+ \to p\overline{K}^0)$ needs to take into account both the factorization term and the ETC term[7] given by Eqs. (27)–(30). Using Eqs. (32)–(34), we find

$$\mathcal{M}(\Lambda_c^+ \to p\overline{K}^0) = i\, f_K \left(\frac{G_F \cos^2\theta_C}{\sqrt{2}}\right) \bar{u}_\Lambda[-(f_p+1.0)0.76 \text{ GeV} + f_p(1.97 \text{ GeV}\gamma_5)]u_{\Lambda_c} . \quad (36)$$

Recasting the principal features of Eqs. (2)–(6) in covariant form, we have (in the rest frame of B_c)

$$\mathcal{M} = i\bar{u}_B(A + B\gamma_5)u_{B_c} ,$$
$$S = A, \quad P = B \mid \mathbf{p}_B \mid /(E_B + m_B) ,$$
$$\Gamma = \frac{\mid \mathbf{p}_B \mid}{4\pi M_{\Lambda_c^+}}(E_B+m_B)[\mid S \mid^2 + \mid P \mid^2] , \quad (37)$$
$$\alpha = \frac{2\,\mathrm{Re}(A^*B) \mid \mathbf{p}_B \mid /(E_B+m_B)}{\mid A \mid^2 + \mid B \mid^2 [\mathbf{p}_B/(E_B+m_B)]^2} .$$

Equations (35) and (36) have of course ignored final-state interaction and CP violation; hence, the appropriate A and B are real. The ratio of $\Lambda_c^+ \to \Lambda\pi$ and $\Lambda_c^+ \to p\overline{K}^0$ decay rate from (37) is

$$\frac{\Gamma(\Lambda_c^+ \to \Lambda\pi^+)}{\Gamma(\Lambda_c^+ \to p\overline{K}^0)} = \frac{0.746 f_\Lambda^2}{(f_p+1.0)^2 + f_p^2} . \quad (38)$$

For $f_p \sim 0.5, f_\Lambda \sim 1.1$, the ratio is 0.36 and thus consistent with experimental data which hover around 0.34 ± 0.13[12] and 0.33 ± 0.19.[13] The predicted asymmetry parameters are, from Eq. (37),

$$\alpha_{\Lambda\pi^+} \approx -1, \quad \alpha_{p\overline{K}^0} \approx -0.61 . \quad (39)$$

The calculations of absolute rate from (37) is however less satisfactory. We have $\Gamma(\Lambda_c^+ \to \Lambda\pi^+) \sim 8.3 \times 10^{10} f_\Lambda^2 \text{ sec}^{-1} = 10 \times 10^{10} \text{ sec}^{-1}$ compared with $\Gamma_{\text{expt}}(\Lambda_c^+ \to \Lambda\pi^+) \approx 2.8 \times 10^{10} \text{ sec}^{-1}$. Likewise $\Gamma(\Lambda_c^+ \to p\overline{K}^0) = 2.78 \times 10^9 \text{ sec}^{-1}$ compared with $\Gamma_{\text{expt}}(\Lambda_c^+ \to p\overline{K}^0) \approx 0.83 \times 10^9 \text{ sec}^{-1}$. Given the simplicity of our model, we feel that discrepancies of a factor of 3 or so in absolute rates are to be tolerated. The asymmetry predictions for α given by Eq. (39) are expected to be more reliable and the deviation of $\alpha(p\overline{K}^0)$ from -1 is a test of the ETC contributions. In asymmetry predictions, (uncertain) absolute normalization is not involved. Since final-state phases are neglected α's can be smaller than our predictions in Eq. (39) by $\cos(\delta_s - \delta_p)$ factor.

The calculation of $B_c \to BV$ for $\Lambda_c^+ \to \Lambda\rho^+, p\overline{K}^{*0}$ then proceeds in a straightforward manner taking into account the subsidiary condition for vector mesons $q_\mu\epsilon_\mu(q)$=0 and evaluating $f_i(q^2), g_i(q^2), f_i'(q^2), g_i'(q^2)$ at $q^2 = m_\rho^2$ (for $\Lambda_c^+ \to \Lambda\rho^+$) and $q^2 = m^2_{\overline{K}^{*0}}$ (for $\Lambda_c^+ \to p\overline{K}^{*0}$) from Eqs. (32)–(34). Equal-time-commutator contributions play no role here. The relevant matrix elements are

$$\mathcal{M}(\Lambda_c^+ \to \Lambda\rho^+) = f_\Lambda f_\rho \left(\frac{G_F}{\sqrt{2}}\cos^2\theta_C\right)\overline{u}_\Lambda[i(0.948 - 0.58\gamma_5)\gamma_\mu + (0.22\ \text{GeV}^{-1} - 0.053\ \text{GeV}^{-1}\gamma_5)P_{1\mu}]u_{\Lambda_c}\epsilon_\mu\ , \tag{40}$$

$$\mathcal{M}(\Lambda_c^+ \to p\overline{K}^{*0}) = f_p f_{K^*}\left(\frac{G_F}{\sqrt{2}}\cos^2\theta_C\right)\overline{u}_p[(1.386 - 0.867\gamma_5)\gamma_\mu + (0.32\ \text{GeV}^{-1} - 0.070\ \text{GeV}^{-1}\gamma_5)P_{1\mu}]u_{\Lambda_c}\epsilon_\mu\ , \tag{41}$$

where $f_\rho \approx f_{K^*} \cong 0.11\ \text{GeV}^2$,[14] ϵ_μ is the polarization of the appropriate vector meson, and $P_{1\mu}$ is the four-momentum of Λ_c^+.

The calculations of decay rate and α asymmetry parameter then follow from Eqs. (13) and (14) of Sec. II. The results are

$$\Gamma(\Lambda_c^+ \to \Lambda^0\rho^+) = f_\Lambda^2(1.37)\times 10^{11}\ \text{sec}^{-1}\ ,$$
$$\Gamma(\Lambda_c^+ \to p\overline{K}^{*0}) = f_p^2(2.68)\times 10^{11}\ \text{sec}^{-1}\ . \tag{42}$$

Taking $f_p \sim 0.5, f_\Lambda \sim 1.1$, we have

$$\Gamma(\Lambda_c^+ \to \Lambda^0\rho^+) = 1.66\times 10^{11}\ \text{sec}^{-1}\ ,$$
$$\Gamma(\Lambda_c^+ \to p\overline{K}^{*0}) = 0.67\times 10^{11}\ \text{sec}^{-1}\ . \tag{43}$$

Using the experimental lifetime of Λ_c^+, $\tau(\Lambda_c^+) = 1.8\times 10^{-13}$ sec, we have the predicted branching ratios

$$B(\Lambda_c^+ \to \Lambda^0\rho^+) \cong 3\times 10^{-2} = 3\%\ ,$$
$$B(\Lambda_c^+ \to p\overline{K}^{*0}) \cong 1.2\times 10^{-2} = 1.2\%\ . \tag{44}$$

The branching ratio $B(\Lambda_c^+ \to p\overline{K}^{*0})$ is somewhat higher than the preliminary experimental number of $\sim 0.56\%$ but is in the right ballpark. The asymmetry parameters are predicted to be

$$\alpha_{\Lambda\rho^+} \cong 0.55\ ,\quad \alpha_{p\overline{K}^{*0}} \cong 0.51\ . \tag{45}$$

Again we stress that the asymmetry predictions are more reliable since they do not depend on the uncertainties in our estimate f_p and f_Λ, and the prediction from (42) that

$$\frac{B(\Lambda_c^+ \to \Lambda\rho^+)}{B(\Lambda_c^+ \to p\overline{K}^{*0})} \cong \left(\frac{f_\Lambda}{f_p}\right)^2 \frac{1}{2} \cong 2 \text{ to } 3 \tag{46}$$

is a clean test of the model. Neglected final-state phases can only decrease the values in Eq. (45).

The methods developed here are clearly applicable also to Ξ_c^+ decays to two-body baryon-meson channels, and we shall take up this matter in a subsequent paper.

V. CONCLUSIONS

Our aim in this paper has been to emphasize parity violation in charm decay as a new arena in which to test the standard electroweak model. The consequences of parity violation are most likely to be detected in the decays of charmed baryons because the final state will, in general, be admixture of waves of opposite parity, and because the sequential decay of daughter hyperons will provide a measure of the polarization of the daughter.

We have concentrated on the two-body decays of the lowest-lying spin-$\frac{1}{2}$ charmed baryons into the baryon octet of spin $\frac{1}{2}$ plus either a pseudoscalar meson or a vector meson. The kinematical analysis of final-state polarizations and angular distributions is well known in the pseudoscalar-meson case, but not in the vector-meson case, and we have worked it out here for the first time, to the best of our knowledge.

Flavor-SU(3) selection rules have been used to express the amplitudes for individual decay modes in terms of a set of invariant amplitudes. These expressions lead to relations between the amplitudes for different decay modes and these relations manifest themselves in relationships between the branching ratios and polarization parameters for these modes. Care is required in the handling of final-state interactions, which are likely to be much more important than in the case of hyperon decays.

Finally we have considered a simple model for the decay modes of the Λ_c^+ based upon current algebra and the factorization approximation. We are able to make predictions for branching ratios and asymmetry parameters; in view of our approximations the predictions for asymmetry parameters are likely to be more reliable than those for branching ratios.

We believe that there is much elegant and interesting physics to be found in charmed-baryon decays and we look forward to the time when many new data become available.

Note added. In the text we have neglected contributions from $\frac{1}{2}^+$ and $\frac{1}{2}^-$ baryons in s and u channels. We have now taken these into account. We find that the largest contributions come from the ground-state $\frac{1}{2}^+$ baryon octet to the parity conserving amplitudes.

For the decay mode $\Lambda_c^+ \to \Lambda\pi^+$, the contributions from the s-channel Σ^+ pole and u-channel Σ_c^0 pole tend to cancel, yielding [in the SU(4) limit, when $g_{\Sigma^+\Lambda\pi^+} = g_{\Lambda_c^+\Sigma^0\pi^+}$ and $\langle\Sigma^+ \mid H_w \mid \Lambda_c^+\rangle = \langle\Lambda \mid H_w \mid \Sigma_c^0\rangle = \langle p \mid H_w \mid \Sigma^+\rangle$] for the additional PC pole term

$$g_{\Sigma^+\Lambda\pi^+}\left(\frac{1}{\Lambda_c^+ - \Sigma^+} - \frac{1}{\Sigma_c^0 - \Lambda}\right)\langle p \mid H_w \mid \Sigma^+\rangle \tag{47}$$

With the value of $g_{\Sigma^+\Lambda\pi^+} \sim 9$ (for $D/F \sim 1.8$) and $\langle p \mid H_w \mid \Sigma^+\rangle \sim 1.2\times 10^{-7}$ GeV as in the text, we find for this a value of 2×10^{-7} which is about 10% of the nonpole part of the PC amplitude. Hence the neglect of the pole term is justified. For the decay $\Lambda_c^+ \to p\overline{K}^0$, the s-channel

Σ^+ pole gives

$$\frac{1}{\Lambda_c^+ - \Sigma^+}\langle p \mid H_w \mid \Sigma^+\rangle g_{\Sigma^+ p\overline{K}^0} , \quad (48)$$

which for $g_{\Sigma^+ p\overline{K}^0} \sim 3$ gives 3×10^{-7}, which is about 22% of the nonpole PC amplitude. Although slightly larger than $\Lambda\pi^+$ case it is still small enough to justify its neglect.

For the vector meson decay modes things are rather different. For the $\Lambda_c^+ \to \Lambda\rho^+$ case the pole terms are zero because the $\Sigma^+\Lambda\rho^+$ and $\Lambda_c^+\Sigma_c^0\rho^+$ couplings vanish in the approximation that ρ generates isospin. Since ρ couplings are pure F type and so $g_{\Sigma^+ p\overline{K}^{*0}} = -g_{pn\rho^+} \cong -5.6$, hence the pole contribution to PC amplitude for $\Lambda_c^+ \to p\overline{K}^{*0}$, which is given by

$$\frac{1}{\Lambda_c^+ - \Sigma^+}\langle p \mid H_w \mid \Sigma^+\rangle g_{\Sigma^+ p\overline{K}^{*0}} \quad (49)$$

becomes -8.64×10^{-7} to be compared to 6.31×10^{-7} from the nonpole piece in Eq. (41) corresponding to $f_p f_{K^*}(G_F/\sqrt{2})\cos^2\theta_C(1.386)$ for the γ_μ coupling. The net overall amplitude now becomes -2.33×10^{-7} and the rate and asymmetry are altered drastically. We find a new rate for $\Lambda_c^+ \to p\overline{K}^{*0}$:

$$\Gamma(\Lambda_c^+ \to p\overline{K}^{*0}) = 0.265\times 10^{11}\ \text{sec}^{-1} \quad (50)$$

corresponding to a branching ratio

$$B(p\overline{K}^{*0}) = 0.48\% \quad (51)$$

in much better agreement with the experimental value of 0.56%. The asymmetry parameter is now found to be

$$\alpha(p\overline{K}^{*0}) = 0.14 . \quad (52)$$

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