Issues in Heavy Flavour Baryons

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ABSTRACT

I present a mini-review on the physics of heavy flavour baryons where I concentrate on the HQET description of their exclusive decay modes. In particular I discuss the structure of current-induced bottom baryon to charm baryon transitions, and the structure of pion and photon transitions between heavy charm or bottom baryons in the Heavy Quark Symmetry limit as $m_Q \to \infty$. The emphasis is on the structural similarity of the Heavy Quark Symmetry predictions for the three types of transitions. The requisite coupling expressions are discussed both in the covariant framework as well as in terms of Clebsch-Gordan coefficients and 6-$j$ symbols. At the end of my review I touch on some unresolved issues in exclusive nonleptonic charm and bottom baryon decays which serve to highlight our present lack of understanding of nonleptonic heavy baryon decays.

*Supported in part by BMBF,FRG under contract 06MZ566
1 Introductory Remarks

Because of the initial abundance of data on heavy charm and bottom mesons the attention of experimentalists and theoreticians had initially been drawn towards applications of the Heavy Quark Effective Theory (HQET) to the meson sector. In the meantime the situation has considerably changed and data on heavy baryons and their decay properties are starting to become available in impressive amounts. In the charm sector the states $\Lambda_c(2285)$ and $\Sigma_c(2453)$ are well established while there is first evidence for the $\Sigma_c^*(2510)$ state. Two excited states $\Lambda_c^{**}(2593)$ and $\Lambda_c^{**}(2627)$ have been seen which very likely correspond to the two lowest lying $p$-wave excitations of the light diquark system making up the $\Lambda_{cK1}$ Heavy Quark Symmetry (HQS) doublet to be discussed in Sec.2. The charm-strangeness states $\Xi_c(2470)$ and $\Omega_c(2720)$ as well as the $\Xi_c^*(2643)$ have been seen. First evidence was presented for the $J^P = \frac{1}{2}^+$ state $\Xi_c^*(2570)$ with the flavour configuration $c\{sq\}$. Thus almost all ground state charm baryons have been seen including two $p$-wave states. In the bottom sector the $\Lambda_b(5640)$ has been identified as well as the $\Sigma_b(5713)$ and the $\Sigma_b^*(5869)$. Some indirect evidence has been presented for the $\Xi_b(5800)$.

Apart from the fact that the existence of the above heavy baryon states has now been established there are also copious experimental data on the production characteristics of heavy baryons and on their exclusive and inclusive decays. Most of the data accumulated so far are on charm baryons. A multitude of different experiments both at collider and fixed target facilities have contributed to our present knowledge of charm baryon physics. New experiments are being planned or have already been set up. For example the SELEX experiment E781 at Fermilab is waiting for beam time and plans to log $5 \times 10^4$ fully reconstructed $\Lambda_c \to pK^+\pi^-$ decays per year with many other decay modes reconstructed. Most of the planned experiments will also see bottom baryons albeit at somewhat reduced rates. The next decade will see the opening of a number of new facilities and experiments among which are HERA-B, LHC-B, COMPASS, CLEO III, CHARM 2000, the US and Japanese B-Factories, and the $\tau$—charm factory project. Heavy baryon physics may not be the prime objective of all of these projects but heavy baryons will certainly be seen at these facilities if only as welcome by-products.

Let us try and gain a historical perspective on how heavy baryon production (and
detection) rates have developed over the past years by taking a look at the strangeness sector. In 1964 V.E. Barnes et al. reported on the first observation of a single \( \Omega^- \) in the BNL 80-in. bubble chamber. Compare this to the \( 4 \times 10^5 \, \Omega^- \) events recently recorded by the E800 Collaboration at Fermilab. Another impressive rate figure is the projected total of \( 10^9 \) reconstructed \( \Xi \) hyperons at the planned HYPERCP experiment E-871 at Fermilab. If these figures can be taken as a foreboding of what lies ahead of us in the charm and bottom baryon sector we are certainly heading for exciting times. As a theoretician I must ruefully admit, though, that our understanding of the dynamics of heavy baryons is far from complete. With all the heavy baryon data expected to come up in the next future there is the acute danger that the experimentalists get ahead of us theoreticians.

The framework to treat the dynamics of heavy baryons is Heavy Quark Effective Theory (HQET). In Secs. 2, 3 and 4 we develop in some detail the leading order HQET description of semileptonic, one-pion and photon decays of heavy baryons. The emphasis is on the structural similarity of the HQS description of these decays. In Sec. 5 I discuss possibilities to further constrain the Heavy Quark Symmetry (HQS) structure of the three type of decays by resolving the light diquark transitions in terms of a constituent quark model description of the light diquark transitions with an underlying \( SU(2N_f) \otimes O(3) \) symmetry. Secs. 6 and 7 contain a brief discussion of some multifarious aspects of exclusive nonleptonic heavy baryon decays.

2 Heavy Baryon Spin Wave Functions

Let us begin by constructing the heavy baryon spin wave functions that enter into the description of heavy baryon decays. A heavy baryon is made up of a light diquark system \((qq)\) and a heavy quark \(Q\). The light diquark system has bosonic quantum numbers \(j^P\) with total angular momentum \( j = 0, 1, 2 \ldots \) and parity \( P = \pm 1\). To each diquark system with spin-parity \( j^P\) there is a degenerate heavy baryon doublet with \( J^P = (j \pm \frac{1}{2})^P \) (\( j = 0 \) is an exception). It is important to realize that the HQS structure of the heavy baryon states is entirely determined by the spin-parity \( j^P\) of the light diquark system. The
requisite angular momentum coupling factors can be read off from the coupling scheme

\[ J^P \otimes \frac{1}{2}^+ \Rightarrow J^P. \]  

(1)

Apart from the angular momentum coupling factors the dynamics of the light system is completely decoupled from the heavy quark.

Let us cast these statements into a covariant framework in which the heavy baryon wave function \( \Psi \) describes the amplitude of finding the light diquark system and the heavy quark in the heavy baryon. The covariant equivalent of the coupling scheme Eq. (1) is then given by

\[ \Psi = \phi_{\mu_1 \cdots \mu_j} \psi_{\mu_1 \cdots \mu_j}, \]  

(2)

where \( \phi_{\mu_1 \cdots \mu_j} \) stands for the tensor representation of the spin-parity \( jP \) diquark state and \( \psi_{\mu_1 \cdots \mu_j} \) represents the heavy-side baryon spin wave function (in short: heavy baryon wave function) coupling the heavy quark to the heavy baryon. Let us be more specific. If

\[ | J^P, m_J \rangle = \sum_{m_j+m_Q=m_J} \langle j^P, m_j; \frac{1}{2}^+, m_Q | J^P, m_J \rangle | j^P, m_j \rangle | \frac{1}{2}^+, m_Q \rangle \]  

(3)

defines the light diquark-heavy quark rest-frame wave function, the C.G. coefficients determining the heavy quark - light diquark content of the heavy baryon can be obtained in covariant fashion from the heavy baryon spin wave function by the covariant projection

\[ \langle j^P, m_j; \frac{1}{2}^+, m_Q | J^P, m_J \rangle = \varepsilon^{\ast_{\mu_1 \cdots \mu_j}}(m_j) \bar{u}(m_Q) \psi_{\mu_1 \cdots \mu_j}(m_J). \]  

(4)

The r.h.s. of Eq. (4) can be evaluated for any velocity four-vector \( v_\mu \) of the heavy baryon which, at leading order, equals the velocity of the heavy quark and the diquark system. Details including questions of normalization can be found in [1]. Differing from [1] I have normalized the spinors appearing in Eq. (4) to 1 and not to \( 2M \) and \( 2M_Q \) as in [1]. It is not difficult to construct the appropriate heavy baryon spin wave functions using the heavy quark on-shell constraint \( \psi_{\mu_1 \cdots \mu_j} = \psi_{\mu_1 \cdots \mu_j} \) and the appropriate normalization condition. In Table 1 (fourth column) I have listed a set of correctly normalized heavy baryon spin wave functions that are associated with the diquark states \( j^P = 0^+, 1^+, 0^-, 1^-, 2^- \).

Next I turn my attention to the question of which low-lying heavy baryon states can be expected to exist. From our experience with light baryons and light mesons we know
that one can get a reasonable description of the light particle spectrum in the constituent quark model picture. This is particularly true for the enumeration of states, their spins and their parities. As much as we know up to now, gluon degrees of freedom do not seem to contribute to the particle spectrum. It is thus quite natural to try the same constituent approach to enumerate the light diquark states, their spins and their parities.

From the spin degrees of freedom of the two light quarks one obtains a spin 0 and a spin 1 state. The total orbital state of the diquark system is characterized by two angular degrees of freedom which I take to be the two independent relative momenta $k = \frac{1}{2}(p_1 - p_2)$ and $K = \frac{1}{2}(p_1 + p_2 - 2p_3)$ that can be formed from the two light quark momenta $p_1$ and $p_2$ and the heavy quark momentum $p_3$. The $k$-orbital momentum describes relative orbital excitations of the two quarks, and the $K$-orbital momentum describes orbital excitations of the center of mass of the two light quarks relative to the heavy quark. The $(k, K)$-basis is quite convenient in as much as it allows one to classify the diquark states in terms of $SU(2N_f) \otimes O(3)$ representations as will be discussed later on. Table 1 lists all ground state s-wave and excited p-wave heavy baryon wave functions as they occur in the constituent approach to the light diquark excitations. They are grouped together in terms of $SU(2N_f) \otimes O(3)$ representations with $N_f = 2$ for $(u, d)$. The s-wave states are in the $10 \otimes 1$ representation, and the p-wave states are in the $10 \otimes 3$ and $6 \otimes 3$ representation of $SU(4) \otimes O(3)$ for the $K$- and $k$-multiplets, respectively. Apart from the ground state s-wave baryons one thus has altogether seven $\Lambda$-type p-wave states and seven $\Sigma$-type p-wave states. The analysis can easily be extended to the case $SU(6) \otimes O(3)$ bringing in the strangeness quark in addition.

### 3 Generic Picture of Current, Pion and Photon Transitions

In Fig. 1 we have drawn the generic diagrams that describe $b \rightarrow c$ current transitions, and $c \rightarrow c$ pion and photon transitions between heavy baryons in the HQS limit. The heavy-side and light-side transitions occur completely independent of each other (they “factorize”) except for the requirement that the heavy side and the light side have the
same velocity in the initial and final state, respectively, which are also the velocities of the initial and final heavy baryons. The $b \to c$ current transition induced by the flavour-spinor matrix $\Gamma$ is hard and accordingly there is a change of velocities $v_1 \to v_2$, whereas there is no velocity change in the pion and photon transitions. The heavy-side transitions are completely specified whereas the light-side transitions $j_1^{P_1} \to j_2^{P_2}$, $j_1^{P_1} \to j_2^{P_2} + \pi$ and $j_1^{P_1} \to j_2^{P_2} + \gamma$ are described by a number of form factors or coupling factors which parametrize the light-side transitions. The pion and the photon couple only to the light side. In the case of the pion this is due to its flavour content. In the case of the photon the coupling of the photon to the heavy side involves a spin flip which is down by $1/m_Q$ and thus the photon couples only to the light side in the Heavy Quark Symmetry limit.

Referring to Fig. 1 I am now in the position to write down the generic expressions for the current, pion and photon transitions according to the spin-flavour flow depicted in Fig. 1. One has ($\omega = v_1 \cdot v_2$)

**current transitions:**

\[
\overline{\psi}_{2\nu_1\cdots\nu_2} \psi_{1\nu_1\cdots\nu_2} \Gamma \left( \sum_{i=1}^{N} f_i(\omega) t_{\nu_1\cdots\nu_2;\nu_1\cdots\nu_2} \right)
\]

\[n_1 \cdot n_2 = 1 \quad N = j_{\text{min}} + 1\]
\[n_1 \cdot n_2 = -1 \quad N = j_{\text{min}}\]

**pion transitions:**

\[
\overline{\psi}_{2\nu_1\cdots\nu_2} \psi_{1\nu_1\cdots\nu_2} \left( \sum_{i=1}^{N} f_i(\omega) t_{\nu_1\cdots\nu_2;\nu_1\cdots\nu_2} \right)
\]

\[n_1 \cdot n_2 = 1 \quad N = j_{\text{min}}\]
\[n_1 \cdot n_2 = -1 \quad N = j_{\text{min}} + 1\]

**photon transitions:**

\[
\overline{\psi}_{2\nu_1\cdots\nu_2} \psi_{1\nu_1\cdots\nu_2} \left( \sum_{i=1}^{N} f_i(\omega) t_{\nu_1\cdots\nu_2;\nu_1\cdots\nu_2} \right)
\]

\[j_1 = j_2 \quad N = 2j_1\]
\[j_1 \neq j_2 \quad N = 2j_{\text{min}} + 1\]

where the $\psi^{\nu_1\cdots\nu_j}$ are the heavy baryon spin wave functions introduced in Sec. 2. The pattern of the above decomposition parallels the corresponding decomposition in lepton-hadron interactions where the transition amplitude is written as $j^\mu_{\text{lepton}} \cdot j^h_{\text{hadron}}$. The
structure of the leptonic current \( j_{\mu \text{lepton}} \) is known and the unknown hadronic current \( j_{\mu \text{hadron}} \) is expanded along a set of covariants with the familiar invariant amplitudes as coefficient functions.

In each of the above cases we have also given the result of counting the number \( N \) of independent form factors or coupling factors. These are easy to count by using either helicity amplitude counting or \( LS \) partial wave amplitude counting. In the case of current and pion transitions the counting involves the normalities of the light-side diquarks which is defined by \( n = (-1)^j P \).

The tensors \( t_{\nu_1 \cdots \nu_j \mu_1 \cdots \mu_j} \) appearing in Eq. (5) have to be build from the vectors \( v_1^{\nu_1} \) and \( v_2^{\nu_j} \), the metric tensor \( g_{\mu_\nu} \), the pion or the photon momentum and, depending on parity, from Levi-Civita objects such as \( \varepsilon(\mu_i \nu_k v_1 v_2) := \varepsilon_{\mu_\nu \alpha \beta} v_1^\alpha v_2^\beta \). The number of independent tensors that can be written down in each of the three cases is necessarily identical to the numbers listed in Eqs. (5), (6) and (7). Lack of space prevents us from giving the explicit forms of these tensors. They can be found in [1].

The generic expressions Eq. (5), Eq. (6) and Eq. (7) completely determine the HQS structure of the current, pion and photon transition amplitudes. It is not difficult to work out relations between rates, angular decay distributions etc. from these expressions.

4 6 \(- j\) Symbols in Heavy Baryon Transitions

It is well worth mentioning that all three covariant coupling expressions in Sec.3 (current, pion, photon) can also be written down in terms of Wigner’s 6-\( j \) symbol calculus [1,2] as can be appreciated from the discussion in Sec. 2 (see Eqs. (2) and (3)). For example, looking at the pion transition in Fig. 1 one sees that one has to perform altogether three angular couplings. They are

\[
\begin{align*}
(i) & \quad j_1^{P_1} \otimes \frac{1}{2}^+ \Rightarrow J_1^{P_1} \\
(ii) & \quad j_2^{P_2} \otimes \frac{1}{2}^+ \Rightarrow J_2^{P_2} \\
(iii) & \quad J_2^{P_2} \otimes L_\pi \Rightarrow J_1^{P_1}
\end{align*}
\]
where $L_\pi = l_\pi$ is the orbital momentum of the pion and $J_1^{P_1}$ and $J_2^{P_2}$ denote the $J^P$ quantum numbers of the initial and final baryons. This is a coupling problem well-known from atomic and nuclear physics and the problem is solved by Wigner’s 6-$j$ symbol calculus. One finds [1,2]

$$M_\pi^{J_1 J_2^z \rightarrow J_2^z + L_\pi m} \quad (9)$$

\[
M_{L_\pi}(-1)^{L_\pi+j_2+\frac{1}{2}+J}(2j_1+1)^{1/2}(2J_2+1)^{1/2} \left\{ \begin{array}{ccc} j_2 & j_1 & L_\pi \\ J & J_2 & \frac{1}{2} \end{array} \right\} \langle LmJ_2J_2^2|J_1J_1^2 \rangle,
\]

where the expression in curly brackets is Wigner’s 6-$j$ symbol and $\langle L_\pi M_2J_2^2|J_1J_1^2 \rangle$ is the Clebsch-Gordan coefficient coupling $L_\pi$ and $J_2$ to $J_1$. $M_{L_\pi}$ is the reduced amplitude of the one-pion transition. It is proportional to the invariant coupling $f_{\pi}$ occurring in the covariant expansion in Eq. (8).

Let us, for example, calculate the doublet to doublet transition rates for e.g. $\{\Lambda_{Qk}^\ast\} \rightarrow \{\Sigma_Q\} + \pi$. The rates are in the ratios $4:14:9:9$ as represented in Fig. 2 [1,3]. This result can readily be calculated using the 6-$j$ formula Eq. (9) and some standard orthogonality relations for the 6-$j$ symbols. The corresponding calculation in the covariant approach involves considerably more labour. Also, the result “4+14 = 9+9” for doublet to doublet one-pion transitions is a general result which again can easily be derived using the 6-$j$ approach [1].
5 Constituent Quark Model Approach to Light-Side Transitions

Interest in the constituent quark model has recently been rekindled by the discovery (or rediscovery) that two-body spin-spin interactions between quarks are non-leading in $1/N_C$, at least in the baryon sector [4]. Thus, to leading order in $1/N_C$, light quarks behave as if they were heavy as concerns their spin interactions. In the constituent quark model approach one further assumes that spin and orbital degrees of freedom decouple. One can therefore classify the light diquark system in terms of $SU(2N_f) \otimes O(3)$ symmetry multiplets. Transitions between light quark systems are parametrized in terms of a set of one-body operators whose matrix elements are then evaluated between the $SU(2N_f) \otimes O(3)$ multiplets.

The constituent quark model light-side spin wave functions are constructed according to the coupling scheme

$$\frac{1}{2}^+ \otimes \frac{1}{2}^+ \otimes l_K \otimes l_k \Rightarrow j^P$$

where $l_K$ and $l_k$ denote the two possible orbital angular momenta of the light diquark with parities $(-1)^{l_K}$ and $(-1)^{l_k}$, respectively. The construction of the light-side spin wave functions proceeds in complete analogy to the atomic helium problem only that one has to take into account the extra colour and flavour degrees of freedom present in the quark case. Table 1 lists the appropriate light-side spin wave functions in covariant form. Again, the corresponding C.G. coupling expressions can easily be obtained from the covariant expressions by the appropriate $m$-quantum number projections.

Let us illustrate how the constituent quark model for the light-side transitions leads to predictions that go beyond the HQS predictions by calculating the ground state to ground state current transitions. The relevant light-side one-body transition operator is given by

$$O = A(\omega) \cdot 1 \otimes 1$$

which has to be evaluated between the ground state diquark spin wave functions. There are altogether three ground state to ground state heavy baryon form factors or Isgur-Wise functions, one for the $\Lambda_b \rightarrow \Lambda_c$ transition and two for the $\{\Sigma_b\} \rightarrow \{\Sigma_c\}$ transitions.
Equation (11) tells us that they can all be expressed in terms of the single form factor $A(\omega)$, where $A(1) = 1$ at zero recoil. One then finds that the current transition amplitudes are given by \[\Lambda_b \rightarrow \Lambda_c: M_\lambda = \bar{u}_2 \Gamma^\lambda u_1 \frac{\omega + 1}{2} A(\omega)\] \[\{\Sigma_b\} \rightarrow \{\Sigma_c\}: M_\lambda = \bar{\psi}_2 \Gamma^\lambda \psi_1 \left(-\frac{\omega + 1}{2} g_{\mu\nu} + \frac{1}{2} g_{\nu} \Gamma^\lambda \right) A(\omega)\]

The same result has been obtained by C.K. Chow by analyzing the large $N_C$ limit of QCD \[\text{[4]}.\]

For the current transitions from the bottom baryon ground states to the $p$-wave charm baryon states one similarly reduces the number of reduced form factors when invoking $SU(2N_f) \otimes O(3)$ symmetry in addition to HQS. For the transition into the $K$-multiplet one has a reduction from five HQS reduced form factors to two constituent quark model form factors whereas for transitions into the $k$-multiplet one can relate two HQS reduced form factors to one single spin-orbit form factor \[\text{[5]}.\] These are testable predictions in as much as the population of helicity states in the daughter baryon is fixed resulting in a characteristic decay pattern of its subsequent decay.

The one-pion and photon transitions can be treated in a similar manner. Again one finds a significant simplification of the HQS structure, i.e. the number of coupling factors is reduced from those listed in Eqs. \[\text{[6] and [7]}\] when $SU(2N_f) \otimes O(3)$ is invoked in addition to HQS. Results for the one-pion transitions can be found in \[\text{[8]}.\] Corresponding results for the photon transitions are presently being worked out. We mention that the constituent quark model approach leads to a one-pion width of $\Gamma \approx 1$ MeV for the recently observed charm baryon state $\Xi_c^{0*}(2643)$. This width is consistent with the experimental upper width limit of 5.5 MeV but unfortunately is too small to be measured with present techniques \[\text{[9]}.\]

6 Asymmetry Parameters in $\Lambda_c \rightarrow \Lambda_s$ Transitions

Recently the ARGUS and CLEO collaborations have determined the asymmetry parameters in the semileptonic transition $\Lambda_c \rightarrow \Lambda_s e^+ \nu_e$ \[\text{[10]}\] and in the nonleptonic one-pion
transition $\Lambda_c \rightarrow \Lambda_s + \pi^+$ [11]. In both cases the measured asymmetry parameter $\alpha$ (or, equivalently, the polarization of the daughter baryon $\Lambda_s$) turns out to be rather close to $-1$. The question is whether these two results have a common theoretical explanation. Since the literature contains some wrong statements on this issue I want to take the opportunity to clarify the situation.

To begin with let me remind you that there is a remarkable prediction of HQS for the heavy to light semileptonic transition $\Lambda_c \rightarrow \Lambda_s$ at zero momentum transfer $q^2 = 0$. The $\Lambda_s$ is predicted to emerge with 100% negative polarization at this point [12]. Within error bars this is borne out by experiment [10]. All what is needed in this prediction of HQS is a heavy $\Lambda_c$ while the $\Lambda_s$ can be taken to be light.

Let me assume for the moment that the nonleptonic decay $\Lambda_c \rightarrow \Lambda_s + \pi^+$ is dominated by the so-called factorizable contribution (diagrams I in Fig.3). If this were the case the asymmetries in the two decays would in fact become related. Let me, however, hasten to add beforehand that the nonleptonic charm baryon decays are not dominated by the factorizable diagram as we shall presently see. Returning to the factorizable contribution in Fig.3 one might wonder why there would be a relation at all between two different components of the weak $c \rightarrow s$ current: the nonleptonic one-pion decay tests the scalar current component whereas in the semileptonic transition one is testing the longitudinal current component. A priori these two components are not related except at the point $q^2 = 0$. This can be seen by projecting the relevant current components using the appropriate polarization four-vectors. For these one has

$$\begin{align*}
\text{longitudinal:} & \quad \epsilon_\mu(0) = \frac{1}{\sqrt{q^2}}(|\vec{q}|, 0, 0, q_0) \\
\text{scalar:} & \quad \epsilon_\mu(s) = \frac{1}{\sqrt{q^2}}(q_0, 0, 0, |\vec{q}|)
\end{align*}$$

(13)

where one should keep in mind that the transverse pieces of the vector current transitions decouple at $q^2 = 0$. From Eq. (13) it is evident that $\epsilon_\mu(0) = \epsilon_\mu(s)$ at $q^2 = 0$ where $|\vec{q}| = q_0$ and thus the scalar and longitudinal components become related at this point. The pion point $q^2 = m_\pi^2$ is so close to $q^2 = 0$ that the extrapolation to $q^2 = 0$ is perfectly save. With what has been said up to now it is then very tempting to (erroneously!) invoke a common theoretical Heavy Quark Symmetry origin for the near equality of the above two
asymmetries.

As concerns the mesonic sector one knows that the one-pion transitions and semileptonic transitions close to \( q^2 = 0 \) are in fact related. However, nonleptonic baryon decays are quite different from nonleptonic mesonic decays in that there are more contributing diagrams in the baryon case. In addition to the factorizing diagrams I in Fig. 3 there are the nonfactorizing diagrams II\(_{a,b}\) and III in Fig. 3. That the nonfactorizing diagrams cannot be neglected in charm baryon decays can be surmized from the fact that some of the observed nonleptonic charm baryon decays can only proceed via the nonfactorizing diagrams. As a sample decay take the decay \( \Lambda_c \to \Xi^0 + K^+ \) which proceeds through diagrams II\(_a\) and III and yet has a sizeable experimental branching fraction. From all what has been said one must conclude that the observation of a near maximal negative polarization in the decay \( \Lambda_c \to \Lambda_s + \pi \) does not have a simple explanation but must be considered to be a dynamic accident resulting from the interplay of a number of contributing diagrams.

7 Some Selected Remarks on Exclusive Nonleptonic Bottom Baryon Decays

At the Brüssel ’95 EPS meeting the ALEPH [13] and DELPHI [14] collaborations presented preliminary evidence for the nonleptonic decay \( \Lambda_b \to \Lambda_c + \pi^- \). Projecting into the future one can imagine that, given enough statistics, the full decay chain

\[
\Lambda_b \to \Lambda_c + \pi^- \\
\leftrightarrow \Lambda_s + \pi^+ \\
\leftrightarrow p + \pi^-
\]

(14)
can eventually be reconstructed. The angular decay distribution in the decay chain can be seen to be given by [13]

\[
W(\theta_2, \theta_3) = 1 + \alpha_1 \alpha_2 \cos \theta_2 + \alpha_3 (\alpha_2 + \alpha_1 \cos \theta_2) \cos \theta_3
\]

(15)

where \( \alpha_1, \alpha_2 \) and \( \alpha_3 \) are the asymmetry parameters in the decays \( \Lambda_b \to \Lambda_c + \pi^- \), \( \Lambda_c \to \Lambda_s + \pi^+ \) and \( \Lambda_s \to p + \pi^- \), respectively. The polar angles \( \theta_2 \) and \( \theta_3 \) are defined through the
momenta of the $\Lambda_s$ in the $\Lambda_c$ rest frame, and the proton in the $\Lambda_s$ rest frame, respectively. If one integrates over $\cos \theta_3$ one arrives at the the angular decay distribution

$$W(\theta_2) = 1 + \alpha_1 \alpha_2 \cos \theta_2$$  \hspace{1cm} (16)$$

Since $\alpha_2$ has been measured ($\alpha_2 = -0.89^{+0.19}_{-0.12}$) the decay distribution Eq. (16) can be used to determine the unknown asymmetry parameter $\alpha_1$ in the decay $\Lambda_b \to \Lambda_c + \pi^-$. The quality of this measurement depends on the accuracy with which the $c \to s$ asymmetry parameter $\alpha_2$ is known. This simple observation invites a general comment: the quality of the analysis of future $b \to c$ data depends on the quality of the present $c \to s$ physics analysis. This holds both for polarization type variables and also for branching ratios. The message for experimentalists is obvious: try to improve on the error bars in charm decays if only for the sake of improving future bottom decay analysis. One step further down in the decay chain the situation is quite satisfactory in this regard. The asymmetry parameter $\alpha_3$ in the decay $\Lambda_s \to p + \pi^-$ (which is used as an analyzer to determine the asymmetry parameter $\alpha_2$ in the decay $\Lambda_c \to \Lambda_s + \pi^+$) is known with sufficient accuracy ($\alpha_3 = 0.642 \pm 0.013$) not to limit the accuracy of the $\alpha_2$ determination.

Where do we stand at the moment in our understanding of exclusive nonleptonic bottom baryon decays such as the decay $\Lambda_b \to \Lambda_c + \pi^-$? As a theoretician I must ruefully admit that there exist no satisfactory theory for exclusive nonleptonic bottom baryon decays at present. This is in marked difference to exclusive nonleptonic bottom meson decays where one has achieved a basic understanding over the last few years. The reason for this was stated before: bottom meson decays can be described by calculable factorizing contributions whereas heavy baryon decays involve also nonfactorizing contributions which are difficult to calculate.

One may turn to the strange and charm baryon sector for advice. Nonleptonic hyperon decays have been traditionally calculated using the three ingredients ”soft pion theorem + current algebra + nearest pole dominance”. This approach is not easily carried over to the $c \to s$ and $b \to c$ sectors in as much as the pion is no longer really soft in the latter two cases. For example, in the decay $\Lambda_c \to \Lambda_s + \pi^+$ one has $|\vec{p}_\pi| = 0.86$ GeV and in the decay $\Lambda_b \to \Lambda_c + \pi^-$ one has $|\vec{p}_\pi| = 2.21$ GeV. The soft-pion approach may barely be justified in $c \to s$ decays but certainly does not make sense for the $b \to c$ decays.
A related problem is that the energy released in the decays is so large that one is far away from the region where the ground state baryons can be used for the pole dominance approximation. Nevertheless this approach has been applied with reasonable success to the $c \to s$ decays but certainly should not be used for bottom baryon decays.

One can then ask oneself whether there is any reason to believe that in bottom baryon decays the nonfactorizing contributions are suppressed. In such a case one could then hope to have a theoretical handle on nonleptonic bottom baryon decays. Turning to $1/N_C$ arguments does not help. Although the nonfactorizing diagrams $\Pi_{a,b}$ and $\Pi_3$ appear to be colour suppressed relative to the factorizing diagram $\Pi_a$ this is true only for $N_C = 3$. Considering the fact that baryons contain $N_C$ quarks as $N_C \to \infty$ with $O(N_C)$ light flavoured quarks there is a combinatorial factor proportional to $N_C$ which cancels the explicit diagrammatic $1/N_C$ factor. This result is in agreement with the analysis in the charm and strangeness sector where the nonfactorizing contributions are certainly needed. However, there do in fact exist qualitative arguments for a suppression of the nonfactorizing diagrams in nonleptonic bottom baryon decays. First of all diagram $\Pi_b$ can be seen to be helicity suppressed [16]. Second, in diagrams $\Pi_a$ and $\Pi_3$ one needs to create an energetic light quark-antiquark pair from the vacuum. Since there is a considerable amount of energy released e.g. in the decay $\Lambda_b \to \Lambda_c + \pi^-$ both suppression mechanisms should be quite effective. For example, from the remaining factorizing contribution one would then predict that the asymmetry parameter $\alpha_1$ in the decay $\Lambda_b \to \Lambda_c + \pi^-$ is maximally negative following the HQS arguments presented in Sec.6. Needless to say that it would be highly desirable to put these qualitative arguments on a more quantitative basis.

8 Concluding Remarks

In this review we have limited our attention to the exclusive decay modes of heavy baryons. This choice was dictated by space limitations. There have certainly been some important theoretical advances in the understanding of the inclusive decays of heavy baryons which I did not have time to cover. These advances are important since they have a bearing on
two unresolved puzzles in bottom baryon physics that have been widely publicized in the last year. The first puzzle concerns the unexpectedly small polarization of the Λ_b from Z-decays as measured by the ALEPH collaboration [7]. The second puzzle concerns the sizeable deviation of the lifetime of the Λ_b from the lifetime average of the other bottom hadrons. The latter topic was covered by Bijan Nemati at this workshop.

Acknowledgement:
Much of the material presented in this review is drawn from work done in collaboration with F. Hussain, M. Krämer, J. Landgraf, D. Pirjol and S. Tawfiq. I would like to express my gratitude to S. Groote for help and advice on the LATEX version of this report.

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Figure Captions

Fig. 1: Generic picture of bottom to charm current transitions, and pion and photon transitions in the charm sector in the HQS limit $m_Q \to \infty$

Fig. 2: One-pion transition strengths for the transitions $\{\Lambda_{QK}^{**}\} \to \{\Sigma_Q\} + \pi$. Degeneracy levels are split for illustrative purposes.

Fig. 3: Quark flow diagrams for exclusive nonleptonic baryon decays. The explicit quark flavour labelling is for the decay $\Lambda_c \to \Lambda_s + \pi^+$.
|        | light side s.w.f.  | $j^P$ | heavy side s.w.f.  | $j^P$ |
|--------|---------------------|------|---------------------|------|
| $\Lambda_Q$ | $\hat{\chi}$      | $0^+$ | $u$                 | $1^+$ |
| $\{\Sigma_Q\}$ | $\chi^{1\mu_1}$  | $1^-$ | $\frac{1}{\sqrt{3}}\gamma_{\mu_1}^5 u$ | $\frac{1}{2}^-$ |
|                           |                    |      | $u_{\mu_1}$         |      |

$s$-wave states ($l_K = 0$, $l_K = 0$)

$p$-wave states ($l_K = 0$, $l_K = 1$)

$p$-wave states ($l_K = 1$, $l_K = 0$)

Table 1: Spin wave functions (s.w.f.) of heavy $\Lambda$-type and $\Sigma$-type $s$- and $p$-wave heavy baryons ($\hat{\chi}^0 = \frac{1}{\sqrt{2}}[(\psi + 1)\gamma_5 C]; \hat{\chi}^{1\mu} = \frac{1}{\sqrt{2}}[(\psi + 1)\gamma_{\mu}^5 C]$).
Figure 1
Figure 2
Figure 3