HEAVY BARYONS –
STATUS AND OVERVIEW (Theory)

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Abstract

We review recent progress in the understanding of the physics of heavy baryons. We begin our review by presenting some highlights of recent experimental findings on charm and bottom baryons and briefly comment on their theoretical implications. On the theoretical side we review new results on the renormalization of HQET, on the Isgur-Wise function for $\Lambda_b \rightarrow \Lambda_c$ transitions and on the flavour-conserving one-pion transitions between heavy baryons.

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1 Introduction

Three years ago one of the present authors gave a review talk on heavy baryons in this series of conferences in Santa Fe, New Mexico [1]. Since then there has been considerable experimental progress in the study of heavy baryons. Let me just run through a list of news items from the experimental front:

- Altogether 17 charm baryons have been seen to date. This is really quite an impressive figure when you compare this to the 13 charm mesons seen so far. This comparison is indicative of the richness of the heavy baryon spectrum as compared to the heavy meson spectrum.

- The new measurement of the polarization of the \( \Lambda_b \) produced on the \( Z \)-resonance [2]

\[
P_{\Lambda_b} = -0.56^{+0.20}_{-0.13} \pm 0.09 \tag{1}
\]

is in better accord with theoretical expectations \((P_{\Lambda_b} \approx -0.6)\) [3] than the old (96) ALEPH measurement \( P_{\Lambda_b} = -0.23^{+0.24}_{-0.20} \pm 0.08 \) [4].

- An orbitally excited charm-strangeness baryon has been seen [5]

- The one-photon transitions \( \Xi_c' \rightarrow \Xi_c + \gamma \) have finally been detected [6]

- The life-time of the \( \Lambda_b \) is still too small – and even becoming slightly smaller in comparison with the bottom meson life times [7].

- There are no news on the hyperfine splitting of the Heavy Quark Symmetry doublet \( \{\Sigma_b, \Sigma_b^*\} \). In an unpublished '95 paper DELPHI [8] had reported on a disturbingly large hyperfine splitting between the \( \Sigma_b^* \) and the \( \Sigma_b \) which has not been confirmed by other experiments.

- This is unfortunately not a new item. There is presently a large amount of data on heavy baryons on tapes waiting to be analyzed (CLEO, FOCUS, SELEX, SLD, . . . ). The community is hopeful that some or all of this data will be analyzed soon.

A lot of data on heavy baryons can be expected to emerge in the near future when CLEO III, BaBar, Belle, HERA-B and CDF/D0 (with the new injector) begin taking data. Although heavy baryons are not the top priority of these experiments heavy baryons will certainly be seen if only as welcome by-products. Also FOCUS, SELEX and possibly SLD may be back in action soon. Then there is the European project COMPASS at CERN which will certainly see charm baryons. In the more distant future there are the LHC experiments ATLAS, CMS and the dedicated bottom hadron detector LHC-B as well as the detector BTeV at Fermilab. These next generation experiments will feature very high bottom quark production rates with excellent possibilities for the study of bottom baryons and their decays.
One may ask what the interest is in studying heavy baryons (charm and bottom) and the transitions among them? Our favourite answer to this question is quite simple. A heavy baryon is the ideal place to study the dynamics of a diquark system in the environment of a heavy quark. In this regard heavy baryons are far more interesting than heavy mesons where the light system consists of a single light quark only. Apart from some tools designed for the treatment of the heavy degree of freedom (HQET) the theoretical analysis of heavy baryons is done with methods well familiar from the light hadron sector. Among these are lattice simulations, QCD sum rules, the large $N_C$ limit, chiral perturbation theory, light cone sum rules, infinite momentum frame techniques, nonrelativistic potential models, constituent quark models, relativistic quark models, Adler-Weisberger and Cabibbo-Radicati sum rules, etc.. Lack of space prevents us from discussing all of these approaches here. In the main part of this review we will focus our attention on three topics. These are the progress in the renormalization of the HQET Lagrangian, the prospects to determine the Isgur-Wise function in $\Lambda_b \rightarrow \Lambda_c$ transitions and recent results on one-pion transitions between heavy baryon states.

2 Progress in the HQET Lagrangian

The tool to study the physics of heavy baryons and the transitions among them is HQET. The HQET Lagrangian is an expansion of the usual QCD Lagrangian in terms of inverse powers of the heavy quark mass. In the rest frame form the HQET Lagrangian reads (see e.g. [9])

$$\mathcal{L}_{\text{HQET}} = \psi^\dagger_{Q}\left\{ iD_0 \\
+ c_k \frac{\vec{D}^2}{2m_Q} \\
+ c_f \frac{\vec{\sigma} \cdot \vec{B}}{2m_Q} \\
+ c_d \frac{[\vec{D} \cdot \vec{E}]}{8m_Q^2} \\
+ i c_s \frac{\vec{\sigma}(\vec{D} \times \vec{E} - \vec{E} \times \vec{D})}{8m_Q^2} \\
+ \frac{1}{m_Q^3} (\text{eleven terms}) + \ldots \right\} \psi_Q$$

(2)

where $\vec{E}$ and $\vec{B}$ are the chromoelectric and chromomagnetic fields, resp., and where $D_\mu = \partial_\mu - igA_\mu^a T_a = (D_0, -\vec{D})$ is the covariant derivative. $\psi_Q$ is the static heavy quark field.

At tree level the coefficients $c_i$ in the HQET Lagrangian are determined as $c_k = 1$, $c_f = 1$, $c_d = 1$ and $c_s = 1$ [3, 11]. The tree level HQET Lagrangian can be obtained from the QCD Lagrangian through a series of Foldy-Wouthuysen-type transformations
For example, after performing the Foldy-Wouthuysen-type transformations up to $O(1/m_Q^3)$ one obtains \[ L_{\text{HQET}}^v = \bar{\psi}_Q \left( i \not{D}_\parallel - \frac{1}{2m_Q} \not{D}_\perp \right) \not{D}_\parallel \psi_Q \]

where we have now used the covariant representation of the tree level HQET Lagrangian. As before $\psi_Q$ is the heavy quark effective field. $D^\mu_\parallel = D^\mu - v \cdot D v^\mu$ is the transverse component of the covariant derivative and $D^\mu_\parallel = v \cdot D v^\mu$ is its longitudinal component where the transverse and longitudinal components are defined with respect to the arbitrary velocity four-vector $v^\mu = (v_0, \vec{v})$ ($v^2 = 1$).

The HQET Lagrangian possesses a remarkable symmetry, namely reparametrization invariance. Reparametrization invariance can be stated in several equivalent ways. Our favourite way of formulating reparametrization invariance is through Lorentz invariance. Consider the original QCD Lagrangian (which is Lorentz invariant) and expand it to all orders in $1/m_Q$ in terms of two HQET Lagrangians which differ by the velocity parameter $v$ that specifies them. One has

\[ L_{\text{QCD}} = L_{\text{HQET}}^v (\text{all orders}) = L_{\text{HQET}}^{v'} (\text{all orders}) \]

It is quite evident that one must have

\[ L_{\text{HQET}}^v (\text{all orders}) = L_{\text{HQET}}^{v'} (\text{all orders}) \]

On the other hand one can transform $L_{\text{HQET}}^v$ into $L_{\text{HQET}}^{v'}$ by the appropriate Lorentz transformation $v \rightarrow v'$ \[13\]. In this way different coefficients in the HQET Lagrangian become related. For example, one has \[13\,14\]

\[ c_k = 1 \]
\[ c_s = 2c_f - c_k \]

These reparametrization invariance relations are expected to hold to all orders in $\alpha_s$ of the renormalized HQET Lagrangian. Whether the conceptually simple derivation of the reparametrization relations through Lorentz invariance can be upheld to any order of $\alpha_s$ remains to be seen. The difficulty is that in the derivation \[13\] it was assumed that the effective fields and operators of the HQET Lagrangian transform as separate entities under Lorentz transformations while they become entangled under renormalization. Needless to say that the reparametrization relations become very useful checks on the correctness of a renormalization calculation. They entail very powerful identities among loop results even at the one-loop level which require quite sophisticated means to understand in detail \[15\].
The renormalization of the HQET Lagrangian is achieved through matching with the corresponding renormalized QCD Lagrangian. The kinetic operator does not get renormalized to any loop order. The one-loop renormalization of the chromomagnetic operator was first done in [16]. The two-loop anomalous dimension of the chromomagnetic operator was obtained in [17] and in [18] where the full renormalization including also finite contributions was carried out. Finite corrections due to the appearance of two different mass scales were obtained in [19]. The one-loop renormalization of the \( O(1/m_Q^2) \) operators (Darwin, spin-orbit) was done in [11, 20] and in [13, 21] where also the mixing with light quark fields was considered. The full two-loop renormalization of these operators is in progress [22]. Finally I mention first attempts at the one-loop renormalization of the set of \( O(1/m_Q^3) \) operators [23].

3 Isgur-Wise function for \( \Lambda_b \to \Lambda_c + l^- + \bar{\nu}_l \)

A great deal of experimental and theoretical effort has been expended on the determination of the Isgur-Wise function in the exclusive semileptonic decays of the \( B \) meson. Exclusive semileptonic \( B \)-decays together with a good understanding of the underlying theory are believed to be one of the key experiments in the determination of the KM matrix element \( V_{bc} \).

It is then quite natural to ask in what way a \( V_{bc} \) determination from exclusive semileptonic heavy bottom baryon decays could complement the \( V_{bc} \) determination from bottom meson decays. The best candidate for such a determination certainly is the \( \Lambda_b \to \Lambda_c \) transition which, in HQET, even has a simpler structure than the corresponding mesonic transitions.

Let us briefly review the \( O(1) \) and \( O(1/m_Q) \) structure of the \( \Lambda_b \to \Lambda_c \) form factors as predicted by HQET (see e.g. [24]). To leading order in the heavy mass expansion the \( b \to c \) current matrix element is given by

\[
\langle \Lambda_c(v_2) \mid J_{\mu}^{V-A} \mid \Lambda_b(v_1) \rangle = F(\omega)\pi_c\gamma_\mu(1 - \gamma_5)u_b
\]

where the \( O(1) \) reduced form factor \( F(\omega) \) satisfies the zero recoil normalization condition \( F(\omega = 1) = 1 \). We define the velocity transfer variable \( \omega \) by \( \omega = v_1 \cdot v_2 \), as usual.

There have been a number of attempts to calculate the \( \omega \)-dependence of the Isgur-Wise function \( F(\omega) \). Unfortunately there is a wide spread in the predictions of the various models for the slope parameter \( \rho^2 \) characterizing its fall-off behaviour at zero recoil (\( \rho^2 \) is defined by the expansion \( F(\omega) = 1 - \rho^2 (\omega - 1) + \ldots \)). These range from \( \rho^2 = 1/3 \) [24, 26] to \( \rho^2 \) around 3 [27, 28]. A recent lattice calculation gives a slope of \( \rho^2 = 1.2^{+0.8}_{-1.1} \) [29]. QCD sum rule determinations suffer from an inherent ambiguity resulting from the fact that there is a two-fold ambiguity in the choice of the interpolating fields of the heavy baryon current. For example, in the corrected version of [30] one has \( \rho^2 = 0.85 \) for both diagonal sum rules and \( \rho^2 = 0.65 \) for the nondiagonal sum rule, both numbers with a theoretical error of \( \approx 0.1 \). It is clear that it would be highly desirable to have some experimental input to clear up the situation. Unfortunately no data has been published thus far on the baryonic Isgur-Wise function. The only available result is from a preprint version of
a DELPHI-analysis \[31\] with the result \(\rho^2 = 1.81^{+0.70}_{-0.67} \pm 0.32\) which, however, has never appeared in print version.

A quick first estimate of the slope of the baryonic Isgur-Wise function can be obtained by relating it to the mesonic Isgur-Wise function assuming that the two light quarks in the heavy baryon move independently of each other. In this way one obtains \[24, 32\]

\[
F(\omega) = \frac{\omega + 1}{2} - \xi^2(\omega)
\]

(9)

where \(\xi^2(\omega)\) is the mesonic Isgur-Wise function. The factor \((\omega + 1)/2\) is a purely relativistic effect and guarantees the correct threshold behaviour in the crossed \(e^+e^-\)-channel \[24, 32\]. For the slope parameter one then obtains

\[
\rho_{\text{baryon}}^2 = 2\rho_{\text{meson}}^2 - \frac{1}{2}
\]

(10)

where the term 1/2 results from the above relativistic effect. Given that the slope of the mesonic Isgur-Wise function is \(\approx 1\) one would then obtain \(\rho_{\text{baryon}}^2 \approx 1.5\). Incidentally, this value is quite close to the results of the dipole model in \[33\] (1.77), and the values of the IMF model \[33\] (1.44) and the relativistic three-quark model \[34\] (1.35) using their favoured sets of model parameters.

At \(\mathcal{O}(1/m_Q)\) all three vector and axial vector form factors \(f_i^V\) and \(f_i^A\) of the process become populated. They are defined according to

\[
\langle \Lambda_c(v_2) | J^V_\mu | \Lambda_b(v_1) \rangle = \pi_c(v_2)(f_1^V \gamma_\mu + f_2^V v_{1\mu} + f_3^V v_{2\mu})u_b(v_1)
\]

(11)

\[
\langle \Lambda_c(v_2) | J^A_\mu | \Lambda_b(v_1) \rangle = \pi_c(v_2)(f_1^A \gamma_\mu + f_2^A v_{1\mu} + f_3^A v_{2\mu})\gamma_5u_b(v_1)
\]

(12)

The \(\mathcal{O}(1/m_Q)\) prediction for these form factors read (see e.g. \[24, 33\])

\[
f_1^V(\omega) = F(\omega) + \frac{1}{2} \left[ \frac{1}{m_c} + \frac{1}{m_b} \right] \left( \eta(\omega) + \bar{\Lambda}F(\omega) \right)
\]

\[
f_1^A(\omega) = F(\omega) + \frac{1}{2} \left[ \frac{1}{m_c} + \frac{1}{m_b} \right] \left( \eta(\omega) + \bar{\Lambda}F(\omega) \frac{\omega - 1}{\omega + 1} \right)
\]

\[
f_2^V(\omega) = f_2^A(\omega) = -\frac{1}{m_c} \frac{\bar{\Lambda}F(\omega)}{1 + \omega}
\]

\[
f_3^V(\omega) = -f_3^A(\omega) = -\frac{1}{m_b} \frac{\bar{\Lambda}F(\omega)}{1 + \omega}
\]

(13)

where \(\eta(\omega)\) satisfies the zero recoil normalization condition \(\eta(\omega = 1) = 0\). Eq. (13) shows that, up to \(\mathcal{O}(1/m_Q)\), the six form factors are given in terms of the \(\mathcal{O}(1)\) form factor function \(F(\omega)\), a new \(\mathcal{O}(1/m_Q)\) form factor function \(\eta(\omega)\) and the constant \(\bar{\Lambda} \approx m_Q - m_Q \approx 600\ MeV\). The \(\mathcal{O}(1/m_Q)\) results can be seen to satisfy Luke’s theorem which reads

\[
f_1^V(1) + f_2^V(1) + f_3^V(1) = 1
\]

\[
f_1^A(1) = 1
\]

(14)
The new unknown $O(1/m_Q)$ form factor function $\eta(\omega)$ has been found to be negligibly small in two recent theoretical evaluations using QCD sum rules \cite{33} and Infinite Momentum Frame (IMF) methods \cite{33}. Some arguments have been presented in \cite{36} that $\eta(\omega)$ is zero at the leading order of the $1/N_C$-expansion. If one assumes that $\eta(\omega)$ can be entirely neglected then the $O(1/m_Q)$ behaviour of exclusive semileptonic $\Lambda_b$ decays is solely determined by the leading order $O(1)$ Isgur-Wise function $F(\omega)$ ($\bar{\Lambda}$ can be determined from elsewhere). This observation opens the way to a meaningful comparison of experimental data analyzed at $O(1/m_Q)$ with theoretical evaluations done at $O(1)$.

4 One-pion transitions between heavy baryons

In this section we will be concerned with flavour-conserving one-pion transitions between heavy baryons (in contradistinction to flavour-changing transitions as e.g. in $\Lambda_b \to \Lambda_c + \pi$). We will discuss ground-state $S$-wave heavy baryons as well as excited $P$-wave heavy baryons and the one-pion transitions between them. There are two types of $P$-wave states depending on whether the two light quarks are in relative $P$-wave or whether the two light quarks as a whole are in a $P$-wave state relative to the heavy quark. The latter we call $K$-excitations while we call the former $k$-excitations. The $K$-excitations lie $\approx 150$ MeV below the $k$-excitations according to a potential model calculation using harmonic oscillator forces \cite{37}. The two experimentally observed excited $\Lambda_c$ states $\Lambda_c(2593)$ and $\Lambda_c(2625)$ are very likely the $J = 1/2$ and $J = 3/2$ members of the lowest lying Heavy Quark Symmetry doublet $\{\Lambda_cK\}$. They have been discovered through their pion transitions. In due course other $P$-wave heavy baryon states and their pion transitions (such as the evidence for the orbitally excited charm-strangeness baryon \cite{5}) will be discovered. It is the purpose of this section to describe some of the progress which has been made in the description of one-pion transitions between heavy baryons. The language to be used in this description will be the very compact language of the $3nj$-formalism.

At the particle level the transition $J_1 \to J_2 + \pi(l)$ is described in terms of the reduced matrix elements

$$\langle J_2 \parallel O^I \parallel J_1 \rangle$$

where $O^I$ is the total transition operator between the initial state with spin $J_1$ and the final state $J_2 + \pi(l)$ with a pion in the orbital state $l$. To leading order in the heavy quark mass expansion these transitions can be viewed as a pion transition between the light diquarks $j_1 \to j_2 + \pi(l)$ in the presence of a noninteracting static heavy quark. In this limit the transition is thus governed by the reduced matrix elements

$$\langle j_2 \parallel O^I \parallel j_1 \rangle$$

(16)

Since the number of reduced matrix elements (or coupling constants) at the diquark level is less than that at the particle level one has achieved a reduction in the number of independent coupling constants that describe the one-pion transitions.

The number of independent coupling constants can be further reduced if one invokes in addition the constituent quark model for the light-side transitions together with the
assumption that the one-pion transition is a one-body operator (justified in the $1/N_C$ approach [38]). In the constituent quark model the light-side transitions are given in terms of the product of reduced matrix elements

$$
\langle s_{q_2} \parallel O^\sigma \parallel s_{q_1} \rangle \langle L_2 \parallel O^L \parallel L_1 \rangle
$$

where $s_{q_1}$ and $s_{q_2}$ are the active light quarks in the one-pion transition and $O^\sigma$ is the spin-1 one-body operator that induces the transition between the two. $L_1$ and $L_2$ are the orbital angular momenta of the light quark system and $O^L$ is the orbital angular momentum operator that induces the orbital transition.

Technically the reduction from the particle level to the diquark level and then further to the constituent quark level involves a recoupling analysis of the various angular momenta involved in the transition. One is therefore naturally led to the use of 6$j$- and 9$j$-symbols.

The first stage of the reduction from the particle level to the diquark level involves a recoupling of the six angular momenta $j_1$ (initial light diquark spin), $j_2$ (final light diquark spin), $J_1$ (initial heavy baryon spin), $J_2$ (final heavy baryon spin), $l$ (orbital angular momentum of pion) and the heavy quark spin $s_Q = 1/2$. The number of angular momenta already suggests that the desired reduction can be achieved with the help of the 6$j$-symbol. In fact one has the two coupling schemes

I. $l + j_2 = j_1$, $j_1 + 1/2 = J_1$  
II. $j_2 + 1/2 = J_2$, $J_2 + l = J_1$

The two coupling schemes are related via recoupling coefficients which in this case are given by the 6$j$-symbol

$$
\left\{ \begin{array}{ccc}
  l & j_2 & j_1 \\
  1/2 & J_1 & J_2 \\
\end{array} \right\}
$$

Figure 1: Recoupling diagram representing a 6$j$-symbol in the recoupling of six angular momenta in the HQS limit.
A pictorial representation of the 6j-symbol is shown in Fig. 1. Each of the six links in Fig. 1 represent angular momenta while the four nodes represent the coupling of angular momenta.

Using standard orthogonality relations involving C.G. coefficients and the 6j-symbol one can relate the reduced matrix elements at the particle and diquark level. One has

$$\langle J_2 || O^l || J_1 \rangle = (-1)^{j_1 + 1/2 + l - j_1} \sqrt{(2 J_1 + 1)(2 J_2 + 1)} \times \left\{ \begin{array}{ccc} l & j_2 & j_1 \\ 1/2 & J_1 & J_2 \end{array} \right\} \langle j_2 || O^l || j_1 \rangle \quad (21)$$

It is evident that one has thereby achieved a reduction in the coupling constant complexity.

In the second stage one resolves the diquark transitions into constituent quark transitions. At this stage one has to recouple altogether twelve angular momenta. The first six are \( s_{q_1} = 1/2 \) and \( s_{q_2} = 1/2 \) (initial and final active light quarks), \( s_{q_s} = 1/2 \) (passive spectator quark), \( S_1 \) and \( S_2 \) (initial and final sum of spins in the light diquark) and \( \sigma = 1 \) (angular momentum of one-pion transition operator). In addition one has the orbital angular momenta \( L_1, L_2 \) and \( L \) from the orbital transition operator \( O^L \), and the angular momenta \( j_1, j_2 \) and \( l \), as before. At first sight one would presume that the one-pion transitions are now described in terms of a 12j-symbol. However, in the constituent quark model one neglects spin-orbit coupling. The spin and orbital spaces decouple and factorise in the transition. This implies that the one-pion transitions are determined by a product of a 6j- and 9j-symbol acting separately in spin space and orbital space, respectively.

In spin space one has the two coupling schemes

I. \( s_{q_s} + s_{q_2} = S_2, \quad S_2 + \sigma = S_1 \) \quad (22)

II. \( s_{q_2} + \sigma = s_{q_1}, \quad s_{q_s} + s_{q_1} = S_1 \) \quad (23)

with the recoupling coefficient (6j-symbol)

\[
\left\{ \begin{array}{ccc} s_{q_s} & s_{q_2} & S_2 \\ \sigma & S_1 & s_{q_1} \end{array} \right\} \quad (24)
\]

In orbital angular momentum space one has the two coupling schemes

I. \( L + \sigma = l, \quad L_2 + S_2 = j_2, \quad l + j_2 = j_1 \) \quad (25)

II. \( L + L_2 = L_1, \quad \sigma + S_2 = S_1, \quad L_1 + S_1 = j_1 \), \quad (26)

and thus the recoupling coefficient (9j-symbol)

\[
\left\{ \begin{array}{ccc} L & \sigma & l \\ L_2 & S_2 & j_2 \\ L_1 & S_1 & j_1 \end{array} \right\} \quad (27)
\]

The relevant recoupling diagrams representing the recoupling of the respective two coupling schemes are depicted in Fig. 2. Fig. 2a is a pictorial representation of the 6j-symbol acting in spin space with four nodes (couplings) and six links (angular momenta) while Fig. 2b represents the 9j-symbol acting in orbital space with six nodes and nine links.
The diquark reduced matrix element can thus be expressed in terms of the product of spin and orbital reduced matrix elements of the constituent quark model. After a little bit of algebra using identities involving C.G. coefficients, $6j$- and $9j$-symbols one obtains

$$\langle j_2 || \mathcal{O}_l || j_1 \rangle = (-1)^{S_{q_2}+S_{q_1}+l+S_1+j_1-j_2} \times \sqrt{(2l+1)(2j_1+1)(2j_2+1)(2S_1+1)(2S_2+1)}$$

$$\times \left\{ \begin{array}{c} s_{q_2} \times s_{q_1} \\ S_1 \\ s_{q_2} \end{array} \right\} \left\{ \begin{array}{c} \sigma \\ S_1 \\ \sigma \end{array} \right\} \left\{ \begin{array}{c} L \\ L_1 \\ L \end{array} \right\} \left\{ \begin{array}{c} l \\ j_1 \\ j_2 \end{array} \right\} \langle S_{q_2} || \mathcal{O}_\sigma || S_{q_1} \rangle \langle L_2 || \mathcal{O}_l || L_1 \rangle .$$

This is our master formula giving the predictions of the constituent quark model for the one-pion transitions between two orbitally excited heavy baryon states for any general transition $L_1 \rightarrow L_2$. For the two cases discussed here, namely $L_1 = L_2 = 0$ and $L_1 = 1$, $L_2 = 0$ the structure of the master formula considerably simplifies since the $9j$-symbol reduces to a Kronecker-$\delta$ in the first case and to a $6j$-symbol in the second case.

In Table 1 we enumerate the number of independent couplings using various model assumptions starting from the particle level. We then count the number of independent couplings using Heavy Quark Symmetry (HQS) and further using the constituent quark model (CQM). It is gratifying to see how each additional symmetry reduces the number of independent couplings. In the case of elastic transitions one even obtains an absolute prediction in the CQM approach since the orbital overlap becomes normalized in the elastic case and the coupling of the pion to the constituent quarks is known from PCAC. On the other hand, in order to obtain absolute predictions for the $P$-wave to $S$-wave transitions one needs to bring in further dynamics. First dynamical model calculations in this direction.
| Number of couplings | particle level | HQS | CQM | PCAC |
|---------------------|----------------|-----|-----|------|
| S-wave to S-wave    | 8              | 2   | 1   | predicted |
| P-wave to S-wave    | 19             | 7   | 2   | –    |

Table 1: Enumeration of the number of independent couplings (or reduced matrix elements) in one-pion transitions between heavy baryons using various model assumptions. The $P$-wave states refer either to $K$- or to $k$-excitations.

have been done using a light-front quark model \[1\] and a relativistic three-quark model \[12\] with some promising results. There is certainly need for more data to compare the model predictions with.

All what was done in this section could have been done using chirally invariant couplings and explicit quark model spin wave functions \[13\] or covariant quark model wave functions \[39\]. We chose to present our results in terms of the very elegant $3nj$-symbol approach partly for the reason that parts of the audience at the Bonn meeting might find it amusing that their upbringing in nuclear and/or atomic physics, where $3nj$-symbols are heavily used, now would allow them to quickly grasp the physics of one-pion transitions, and, for that matter, one-photon transitions \[24, 44, 45\] between heavy baryons.

5 Concluding remarks

We provided a brief review of a few selected topics in heavy baryon physics. Lack of space prevented us from covering more details. We are looking forward to more data on heavy baryons which will hopefully be forthcoming soon. These data will certainly stimulate further theoretical progress.

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