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WEAK SEMILEPTONIC DECAYS OF HEAVY HADRONS*

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WEAK SEMILEPTONIC DECAYS OF HEAVY HADRONS

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ABSTRACT
I describe the phyiscs ideas that go into the Formulation of Heavy Quark Symmetry. I use Heavy Quark Symmetry and the Heavy Quark Effective Theory to discuss the weak semileptonic decays of heavy hadrons where I concentrate on heavy mesons decays.

1. INTRODUCTION

Much of the motivation to study the weak decay properties of heavy hadrons can be traced back to the need to determine one of the fundamental constants of nature, the Kobayashi-Maskawa (KM) matrix element V_{bc}. It was realized in the last few years that exclusive semileptonic decays of bottom to charm hadrons are much better suited for this purpose than, as had been thought originally, the inclusive semileptonic $b \to c$ decays [1]. The reason for this is that the KM matrix element V_{bc} can be regarded as a weak transition charge that can be accurately measured at the zero recoil point, at least in the limit when the bottom and charm quark mass become very heavy. This is so since the associated hadron transition form factor is normalized to one at zero recoil [2] just as in the case of electromagnetic transitions where the charge form factor is normalized to one at $q^2 = 0$. Thus the measurement of the weak transition charge V_{bc} acquires the same status as the measurement of the electric charge, at least in the large mass limit. Much better, when corrections to the large mass limit were studied at a later stage, it was realized that the zero recoil normalization condition remains intact at $0(1/m_Q)$ [3,4], where m_Q is the heavy quark's mass.

Best suited for the determination of the KM element V_{bc} are the mesonic and baryonic ground state to ground state transitions $B \to D, D^*$ and $\Lambda_b \to \Lambda_c$, resp., whose flavour diagrams are drawn in Fig.1.

Other decay candidates in the baryon sector are the $1/2^+ \to 1/2^+$ transitions $\Xi_b(b[su]) \to \Xi_c(c[su])$ and $\Omega_b(b\{ss\}) \to \Omega_c(c\{ss\})$ and the $1/2^+ \to 3/2^+$ ttransition

[1] Supported in part by the BMFT, FRG under constract 06MZ730

$\Omega_b(b\{ss\}) \rightarrow \Omega^*_c(c\{ss\})$, where $[q_1q_2]$ and $\{q_1q_2\}$ refer to flavour-antisymmetric spin 0 and flavour-symmetric spin 1 diquark states, respectively. In this report I will mainly concentrate on heavy meson transitions and leave the subject of heavy baryon transitions to a companion review [5].

Obviously one needs a bridge to connect the physics at the quark level, where theory is formulated and where V_{bc} is defined, to the particle level, where, after all, the experiments are done. Fortunately, there has been significant progress in this program over the last few years (starting with the papers [6-12]) which I want to report on. The progress is related to the fact that now there exists a systematic expansion of QCD in terms of inverse powers of the heavy quark mass termed the "Heavy Quark Effective Theory (HQET)". The leading term in this expansion gives rise to a new symmetry termed the "Heavy Quark Symmetry (HQS)".

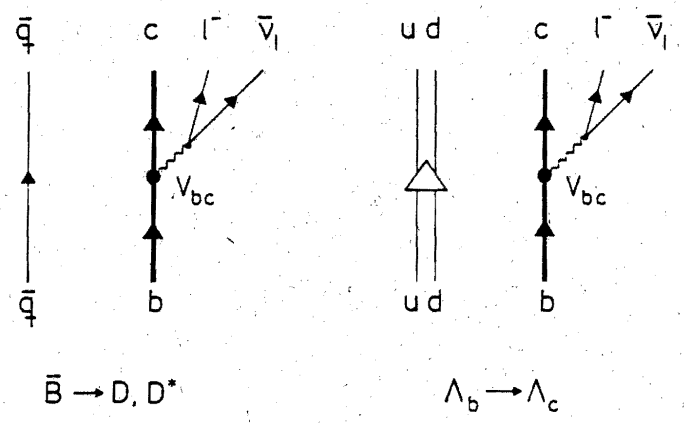

Fig.1: Flavour diagrams of semileptonic decays $\overline{B} \rightarrow (D,D^*) + \ell^-\overline{\nu}_\ell$ and $\Lambda_b \rightarrow \Lambda_c + \ell^-\overline{\nu}_\ell$

Nature has been very kind to us in that it has divided its six flavoured quarks into a heavy and a light quark sector. The "heavy" c-, b-, t-quarks are much heavier than the QCD scale $\Lambda_{QCD} = 300$ MeV whereas the "light" u-, d-, s-quarks are much lighter than Λ_{QCD}, i.e. one has

$$m_c, m_b, m_t >> \Lambda QCD >> m_u, m_d, m_s \qquad (1)$$

In the heavy quark sector it then makes sense to first consider QCD in the limit where the heavy quark masses become very large and then, in the second stage, to consider power corrections to this limit in terms of a systematic $1/m_Q$ expansion. Likewise one can profitably study the light quark sector in the zero mass limit, i.e. in the chiral symmetry limit, and then add corrections to the chiral limit at a later stage.

It is quite intriguing that many of the ideas of HQET date back as far as 1937, then of course in the context of QED [13]. In the Block-Nordsieck approach to soft photon radiation it was the electron that was "infinitely" heavy (on the scale of the soft photons) so that the fermionic degrees of freedom could be treated as a classical source of radiation (no e^+e^- pair creation!). In fact, the Bloch-Nordsieck model was already formulated in terms of an effective theory with the electron degrees of freedom removed from the field theory (see also [14]).

It is quite important to realize that HQS is <u>not</u> a spectrum symmetry but it is a new type of equal veclocity symmetry. That one cannot expect a spectrum symmetry to hold in the heavy quark sector should be quite clear from the fact that there are two orders of magnitude difference between the masses of the c and t quarks! On the other hand, the new type of HQS symmetry at equal velocities takes a little bit of getting used to. But once one has gotten into the habit of thinking in terms of quark and particle velocities the HQS will in fact look quite natural.

We mention that the implications of HQET and HQS have been vigorously studied in the last two-and-a-half years starting with the 1990 papers by Isgur-Wise [6], Bjorken [7], Georgi [10] and our group at Mainz [12]. In the meantime the field is at full blossom with approximately 300 papers published and new papers coming out every week.

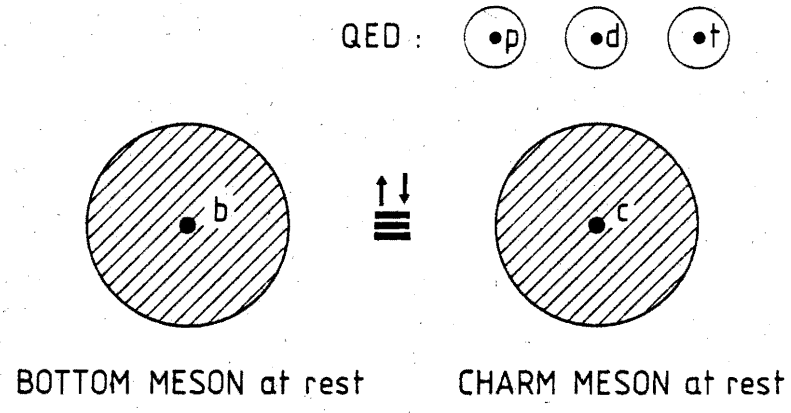

Fig.2: Portrayal of bottom and charm meson wave functions at rest. Upper right corner: wave functions of the hydrogen, deuterium and tritium atoms.

To familiarize one-self with the presence of a spin and flavour symmetry at equal velocity it is quite instructive to consider a bottom and charm meson at rest as

shown in Fig.2. The heavy bottom quark and the charm quark at the center are surrounded by a cloud corresponding to the light antiquark system. The only communication between the cloud and the center is via gluons. But since gluons are flavour blind the light cloud knows nothing about the flavour at the centre. Also, for infinitely heavy quarks, there is no spin communication between the cloud and the center. Thus one concludes that, in the heavy mass limit, a bottom meson at rest is identical to a charm meson at rest regardless of the spin orientation of the heavy quarks, i.e. one has

$$\text{Bottom Meson at rest} \overset{\uparrow\downarrow}{=} \text{Charm Meson at rest} \tag{2}$$

One then just needs to boost the rest configuration by a Lorentz boost from velocity zero to velocity v to conclude

$$\text{Bottom Meson at velocity v} \overset{\uparrow\downarrow}{=} \text{Charm Meson at velocity v} \tag{3}$$

remembering that a Lorentz boost depends only on relative velocities. Eq.(3) exposes the existence of a new spin and flavour symmetry of QCD at equal velocities which holds true in the large mass limit. This is nothing but the advertised Heavy Quark Symmetry HQS.

In fact, everyone should be quite familiar with the existence of such a symmetry in the context of QED. Take a hydrogen, deuterium and tritium atom at rest as also shown in Fig.2. When hyperfine interactions are neglected they possess identical wave functions and thus identical atomic properties. The Coulombic interaction between the electron cloud and the nucleus at the centre is sensitive only to the total charge of the nucleus which is the same for all three isotopes.

In this sense the heavy mesons can be regarded as the hydrogen atom of QCD. This analogy should of course not be taken too literal since the interaction in the true hydrogen atom is Coulombic whereas the confinement potential of QCD remains to be explored. Nevertheless, one may draw a qualitative picture of the level scheme of heavy mesons in analogy to that of the hydrogen atom (see Fig.3). At the fine splitting level one disregards the interaction of the light antiquark system with the heavy quark and consequently one obtains the "atomic" levels of the heavy mesons by coupling the spin of the light antiquark ($J^P = 1/2^-$) with its orbital angular momentum ℓ according to $j(\text{light}) = \ell \otimes 1/2 = (\ell + 1/2) \oplus (\ell - 1/2)$. One obtains e.g. an s-wave state with $j^P(\text{light}) = 1/2^-$ and two p-wave states with $J^P(\text{light}) = 1/2^+$ and $3/2^+$. These states have a two - fold degeneracy each as there is no interaction with the heavy quark's spin at this stage.

Bringing in the hyperfine interaction at $0(1/m_Q)$ one now has spin interactions and the levels will split as shown in Fig.3. The total spin of the hyperfine levels result from the

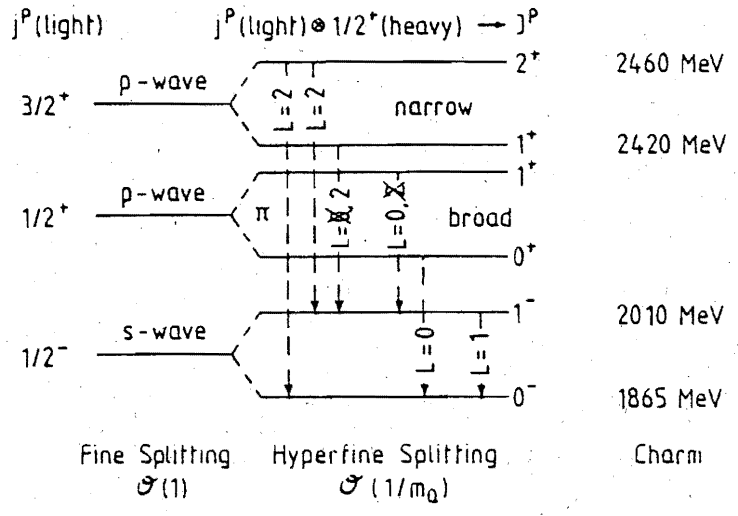

Fig.3: Heavy meson s- and p-wave levels at $0(1)$ (fine splitting) and at $0(1/m_Q)$ (hyperfine splitting). Dashed lines represent one-pion transitions. Orbital angular momenta transitions are crossed out that are $0(1)$ forbidden. To the right: experimentally detected charm meson states.

decomposition of j(light) $\otimes$ 1/2. In the charm sector the s-wave states are the familiar $J^P = 0^-$ and 1^- D and D* states at 1865 MeV and 2010 MeV, respectively. As regards the four p-wave states two rather narrow states with $J^P = 1^+$ and 2^+ have been identified at 2420 MeV and 2460 MeV with widths of $\Gamma \cong 20$ MeV [15]. They would naturally be associated with the $j^P(\text{light}) = 3/2^+$ fine splitting levels as these levels are expected to be narrow on account of the fact that the associated light-side one-pion transitions to the ground-state level ($3/2^+ \rightarrow 1/2^- + \pi$!) are d-wave transitions and thus angular momentum suppressed [16,17,18].

In contradistinction the two $J^P = 0^+$ and 1^+ p-wave states associated with the level $j^P(\text{light}) = 1/2^+$ are expected to have rather large widths ($\Gamma \cong 400$ MeV) as the associated light side one-pion transitions ($1/2^+ \rightarrow 1/2^- + \pi$!) are s-wave transitions [16,17,18]. This is probably the reason why they have escaped detection so far.

This qualitative picture of the spectroscopy of heavy meson states has found another remarkable confirmation through the experimental determination of the B*-B mass splitting [15]. As emphasized above, the level splitting between the two ground state mesons is expected to be a $1/m_Q$ hyperfine splitting effect. Taking the experimental hyperfine splitting values in the bottom and charm sector, i.e. [15]

$$M_{B^*} - M_B = 46 \text{ MeV} \tag{4}$$

$$M_{D^*} - M_D = 142 \text{ MeV} \tag{5}$$

one finds

$$\frac{M_{B^*} - M_B}{M_{D^*} - M_D} = \frac{46}{142} \cong \frac{1/m_b}{1/m_c}, \tag{6}$$

clearly establishing that a $1/m_Q$ type hyperfine splitting effect is in operation for the ground state splitting in the charm and bottom sector.

2. SPIN COMPLEXITY OF TRANSITION FORM FACTORS AND ANGULAR DECAY DISTRIBUTIONS

To start with let us first enumerate the number of form factors that describe semileptonic $B \to D$ and $B \to D^*$ transitions. This is easily done in the usual covariant expansion. One has ($q = p_1 - p_2$)

<u>$B \to D$:</u>

$$\langle D(p_2)|V_\mu|B(p_1)\rangle = F_+^V (p_1 + p_2)_\mu + F_-^V q_\mu \tag{7a}$$

<u>$B \to D^*$:</u>

$$\langle D^*(p_2)|V_\mu + A_\mu|B(p_1)\rangle =$$
$$\varepsilon_2^{*\alpha} (F_1^A g_{\mu\alpha} + F_2^A p_{1\mu} p_{1\alpha} + F_3^A q_\mu p_{1\alpha} + F^V i\, \varepsilon_{\mu\alpha\rho\sigma} p_1^\rho p_2^\sigma) \tag{7b}$$

There are thus (1+1) and (3+1) form factors for $B \to D$ and $B \to D^*$, respectively. The first number in the brackets counts the number of form factors that can be measured in the zero lepton mass case (typically e and μ) whereas a measurement of the form factors multiplying q_μ (F_-^V and F_3^A) require non-zero lepton masses (typically the τ). When one wants to define physical observables it is more advantageous to linearly transform the invariant amplitude F_i defined in (7) to helicity amplitudes H_i (see e.g. [19,20]). These again split into the two (1+1) and (3+1) sets mentioned above.

It is quite remarkable that, in the infinite mass limit, HQS tells us that all the six form factors in (7) are related to one reduced form factor $F(\omega)$ which is a function of the ("scaling") velocity transfer variable $\omega = v_1 \cdot v_2$ and which is normalized at zero recoil $F(\omega=1)=1$. I have intentionally chosen the phrase "reduced form factor" in analogy to the corresponding phrase "reduced matrix element" used in the Wigner-Eckart theorem. We shall later on describe how one actually determines the "Clebsch-Gordan" coefficients that project the six form factors onto the reduced form factor $F(\omega)$. HQS by itself can say nothing about the actual functional form of $F(\omega)$ except for its normalization at zero recoil $\omega=1$. To obtain its functional form one needs additional dynamical input.

Before I derive the HQS relation for the mesonic ground state transitions I want to delineate how the form factors can actually be measured in the semileptonic decay processes and how the predictions of HQS can thus be tested. In particular I want to discuss some recent experimental angular distribution data which already bear witness to the relevance of the HQS predictions. Thereby I hope to convince the reader that the predictions of HQS are no longer just a theoretical fancy but already an experimental factuality.

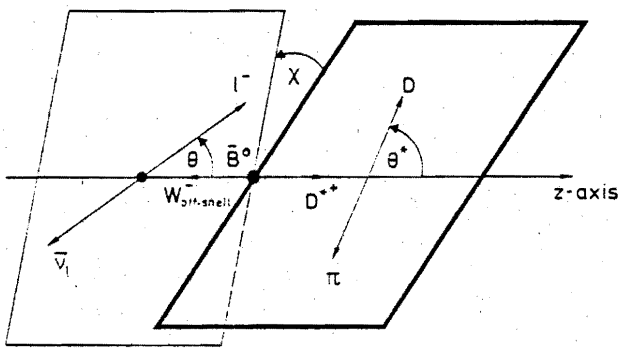

Fig.4: Definition of hadron side polar angle θ^*, lepton-side polar angle θ, and azimuthal angle χ in the decay $\bar{B}^o \to D^{*+} (\to D\pi) + \ell^- \bar{\nu}_\ell$.

Fig.4 shows the decay configuration $\bar{B} \to D^*(\to D\pi) + W^-_{off\text{-}shell}\ (\to \ell^- \bar{\nu}_\ell)$ in the $\bar{B}$ rest system. I view the decay process as a two step process. In the first step the $\bar{B}$ decays into the $W^-_{off\text{-}shell}$ on the one side and the D^* on the other side (back-to-back to the W^-).[2] In the second step these further cascade via $W^- \to \ell^- \bar{\nu}_\ell$ (lepton side) and, via $D^* \to D\pi$ (hadron side)[3]. The second step decays are again analyzed in their respective rest systems in terms of a lepton-side polar angle θ and a hadron-side polar angle θ^*. Finally, the relative orientation of the two decay planes defines an azimuthal angle χ as Fig.4 shows. The first step is governed by the weak decay amplitudes H_i or F_i describing the "decay" $\bar{B} \to D^* + W^-$. The decay products D^* and W^- emerge in highly polarized states. Their polarization density matrices are given in terms of bilinear forms of the decay amplitudes. The second step decays $D^* \to D\pi$ and $W^- \to \ell^- \bar{\nu}_\ell$ can then in turn be used to analyze the polarization states of the D^* and the W^-. In this regard the decay $W^- \to \ell^- \bar{\nu}_\ell$ is an optimal analyzer since it posseses 100% analyzing power.

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[2] For reasons of conciseness the $W_{off\text{-}shell}$ will be referred to as W^- in the following.

[3] The semileptonic decay $\bar{B}^o \to D + W^-(\to \ell^- \bar{\nu}_\ell)$ can be discussed along similar lines. However, the subsequent decays of the D as e.g. $D \to K\pi$ provide no further information on the $B \to D$ dynamics since the D is a spin zero particle.

What has been described in words can be surmised in the form of a joint angular decay distribution for the decay $\bar{B} \to D^*(\to D\pi) + W^-(\to \ell^- \bar{\nu}_\ell)$. The joint angular decay distribution will involve the lepton side polar angle θ, the hadron side polar angle θ^* and the relative azimuth χ of the two decay planes. Collecting all kinematical factors one has [19-21]

$$
\begin{aligned}
&\frac{d\Gamma(\bar{B} \to D^*(\to D\pi) + W^-(\to \ell^- \bar{\nu}_\ell)}{dq^2 d\cos\theta\, d\cos\theta^*\, d\chi} = \frac{G^2}{(2\pi)^4} |V_{bc}|^2 \frac{q^2 p}{12 M_1^2} B(D^* \to D\pi) \\
&\Big\{ \frac{3}{8}(1+\cos^2\theta) \frac{3}{4}\sin^2\theta^* (|H_+|^2 + |H_-|^2) \\
&+ \frac{3}{4}\sin^2\theta \frac{3}{2}\cos^2\theta^* |H_o|^2 \\
&- \frac{3}{4}\sin^2\theta \cos 2\chi \frac{3}{4}\sin^2\theta^* \, \mathrm{Re}(H_+ H_-^*) \\
&- \frac{9}{16}\sin 2\theta \cos\chi \sin 2\theta^* \frac{1}{2}\mathrm{Re}(H_+ H_o^* + H_- H_o^*) \\
&+ \frac{3}{4}\cos\theta \frac{3}{4}\sin^2\theta^* (|H_+| - |H_-|^2) \\
&- \frac{9}{8}\sin\theta \cos\chi \sin 2\theta^* \frac{1}{2}\mathrm{Re}(H_+ H_o^* + H_- H_o^*) \Big\}
\end{aligned}
\tag{8}
$$

where $q^2 = (p_1 - p_2)^2$ is the invariant momentum transfer squared and p is the CM momentum of the D*. The H_{λ_W} are the aforementioned helicity amplitudes of the decay $\bar{B} \to D^* + W^-$ where λ_W is the helicity of the W^-. For $\bar{B} \to D + W^-$ one has a corresponding formula where now only the second term in the decay distribution survives. For the decays $B \to \bar{D}^* + W^+$ involving $\bar{b} \to \bar{c} + W^+$ one has to switch the signs of the last two contributions in (8). The decay distribution (8) holds for zero lepton masses. If lepton mass effects are included there are seven more terms in (8) [20,22]. Furthermore, if one includes also the so-called T-odd contributions that could arise from CP and/or final state interaction effects there are even five more additional terms in (8) when $m_\ell \neq 0$ [23,24]. Thus, when lepton mass and T-odd effects are included, there are altogether 18 observables in the decay distribution $\bar{B} \to D^*(\to D\pi) + W^-(\to \ell^- \bar{\nu}_\ell)$. Since there are only four independent amplitudes in the decay process a complete or even a partial measurement of the observables would considerably overdetermine the form factor amplitudes.

Let me remind the reader that the analysis of joint angular decay distributions such as the one given in Eq.(8) has by now become a standard fare in the analysis of weak decays of heavy hadrons. For example, the well-known amplitude analysis of the decay $D \to K^* + \ell\nu_\ell$ by E691 was based on a full three-fold angular fit to

an event sample of $\cong 200$ events [25]. A similar analysis was done by E653 for the same decay $D \to K^*(\to K\pi)+\mu\nu_\mu$ ($\cong 300$ events) where lepton mass effects ($m_\mu \neq 0$) were included in the analysis. For $b \to c$ decays, ARGUS ($\cong 400$ events) [27] and CLEO ($\cong 200$ events) [28] have done a full amplitude analysis based on the threefold angular decay distribution in the decay $B \to D^*(\to D\pi)+ \ell\nu_\ell$. Some results of their analysis will be presented in a moment's time.

On the theoretical side various aspects of joint angular decay distributions in semi-leptonic decays have been discussed in the literature. I cite refs. [19-24,29,30] for semileptonic meson decays and refs. [31-33] for semileptonic baryon decays. Recently there has been a very comprehensive, almost encyclopedic analysis of joint angular decay distributions in the weak semileptonic and nonleptonic decays of heavy mesons and baryons including non-zero lepton mass effects as well as polarization effects [34].

Instead of analyzing the full three-fold angular decay distribution (8) one can also consider single angle distributions. For example, the longitudinal/transverse composition of the polarization of the D* (or "alignment" polarization) is determined by the hadron side polar angle distrubution

$$\frac{d\Gamma}{dq^2 d\cos\theta^*} \propto 1 + \alpha_{\theta^*}\cos^2\theta^* \tag{9}$$

The asymmetry parameter α_{θ^*} in (9) is determined by [19,20]

$$\alpha_{\theta^*} = \frac{2|H_o|^2 - |H_+|^2 - |H_-|^2}{|H_+|^2 + |H_-|^2} \tag{10}$$

as an inspection of Eq.(8) shows.

In Fig.5 we show a plot of the asymmetry parameter plotted against the lower momentum (or energy) cut-off of the lepton's momentum. The lower momentum cut-off dependence is included because of experimental requirements: a lepton in the decay $B \to D^* + \ell + \nu_\ell$ is detectable only for momentum values above a detector dependent lower momentum threshold. Typically these momentum cut-offs were 1.0 GeV and 1.4 GeV for the first stage ARGUS and CLEO detectors. The 1989 ARGUS [35] and CLEO [36] measurments are included in the plot as well as the predictions of the so-called KS [19,20] model. The experimental values can be seen to fall nicely on top of the theoretical curve. In as much as the KS model incorporates most of the relevant features of HQS, the agreement shows that the HQS predictions work well.

Recently there has been a new measurement of the asymmetry parameter by the ARGUS Collab. who quote [27]

$$\alpha_{\theta^*} = 1.46 \pm 0.62 \pm 0.24 \tag{11}$$

The quoted value now applies to the whole lepton momentum range. This is the result of a new type of analysis of the ARGUS group where the blind regions of the detector have been filled out by smooth extrapolation from the active region of the detector. The new ARGUS value is included in Fig.5 at the appropriate value of zero lower p_ℓ cut-off and can again be seen to agree with the predictions of the KS model.

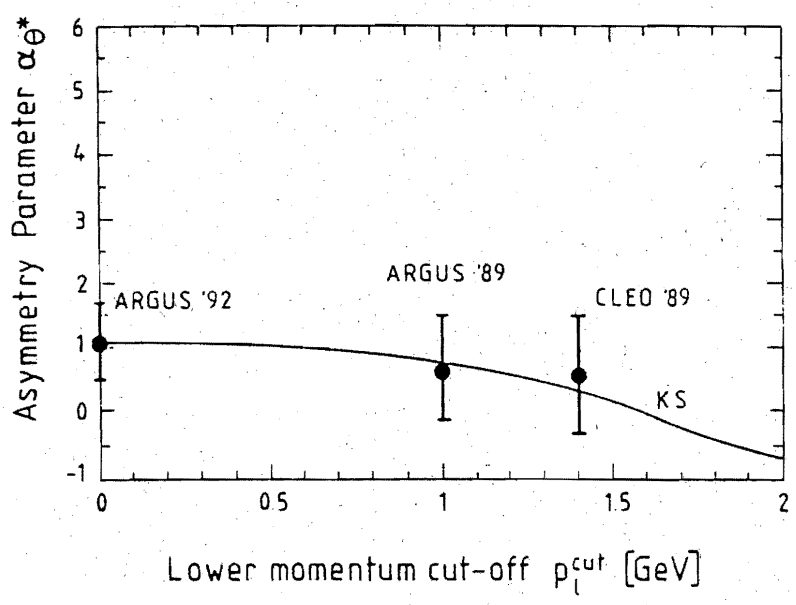

Fig.5: Hadron-side asymmetry parameter $\alpha_{\theta*}$ vs. lower lepton momentum cut-off in the KS model [19,20]. Included are the 1989 measurements of ARGUS [35] and CLEO [36] and the recent ARGUS measurment [27].

Contrary to the hadron-side polar angle distribution the lepton side polar angle distribution is sensitive also to parity-violating effects, i.e. sensitive to the difference of the transverse helicity amplitudes $|H_+|^2 - |H_-|^2$. The polar angle distribution reads [18,20]

$$\frac{d\Gamma}{dq^2 d\cos\theta} \propto 1 + \alpha'_\theta \cos\theta + \alpha_\theta \cos^2\theta \tag{12}$$

with the asymmetry parameters

$$\alpha'_\theta = 2\frac{|H_+|^2 - |H_-|^2}{|H_+|^2 + |H_-|^2 + 2|H_o|^2} \quad \alpha_\theta = \frac{|H_+|^2 + |H_-|^2 - 2|H_o|^2}{|H_+|^2 + |H_-|^2 + 2|H_o|^2} \tag{13}$$

Alternatively one can define a forward-backward (FB) asymmetry ($\theta \geq \pi/2$)

$$A_{FB} = \frac{d\Gamma(\theta) - d\Gamma(\pi - \theta)}{d\Gamma(\theta) + d\Gamma(\pi - \theta)} \tag{14}$$

which, when taken over the whole interval $\pi/2 < \theta < \pi$, is given by

$$A_{FB} = -\frac{3}{4}\frac{|H_+|^2 - |H_-|^2}{|H_+|^2 + |H_-|^2 + |H_0|^2} \tag{15}$$

In Fig.6 we show a plot of the FB asymmetry again plotted against the lower lepton momentum cut-off. The recent ARGUS [27] and CLEO [28] FB asymmetry values can be seen to agree nicely with the predictions of the KS model [19,20] which again demonstrates that the HQS predictions work well.

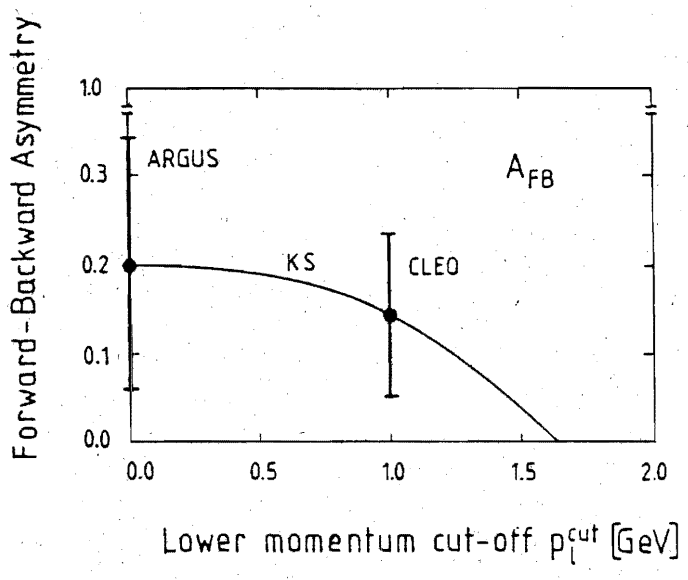

Fig.6: Lepton-side forward-backward asymmetry A_{FB} vs. lower lepton momentum cut-off in the KS model [19,20] compared to the recent ARGUS [27] and CLEO [28] measurements.

The significance of these new asymmetry measurements lies in the fact that they constitute the first measurments of the handedness (or the chirality) of the b→c transition. The positive value of A_{FB} confirms the Standard Model prediction that the transition is left-chiral. That a positive value for A_{FB} provides evidence for the left-chirality of the b→c transition can be easily appreciated by simple helicity arguments. In a left-chiral b→c transition the c and thereby the transverse state of the D* emerges with dominant negative helicity. Since the decaying particle $\overline{B}$ has spin zero the W^- has to follow suit and consequently also has a dominant negative transverse helicity. As the lepton ℓ^- wants to align its helicity with that of the W^- it preferentially emerges in the forward direction giving a positive forward-backward asymmetry. Vice versa, if the b→c transition were right chiral then A_{FB} would be negative. One concludes that the measured sign of the forward-backward asymmetry provides evidence for the b→c transition to be left-chiral as predicted by the Standard Model.

Note that the heavy quark symmetry predictions for the asymmetry parameters in Eqs.(10), (13) and (15) no longer depend on the unknown reduced form factor $F(\omega)$ as this dependence drops out for any given value of q^2 (or ω) when the ratio is formed. Even when one is considering integrated asymmetries (integrating numerator and denominator separately!) the ratios' dependence on the reduced form factor is weak as long as the q^2 (or ω) dependence of the reduced form factor is reasonably smooth. In fact, the authors of Ref.[37] calculated the forward-backward asymmetry using various model form factors and demonstrated that the model dependence of the predicted value of the foward-backward asymmetry is quite weak.

However, the authors of Ref.[37] raise an important and interesting issue concerning the utility of the forward-backward asymmetry measure to determine the chirality of the $b \to c$ transitions. In fact, if one replaces the left-handed SM gauge boson W_L by a non standard model right handed gauge boson W_R an identical asymmetry ratio obtains. Again, one can use simple helicity arguments to convince oneself of this fact. In as much as the SM left-chiral nature of the lepton-side coupling is used as input to analyze the chiral nature of the $b \to c$ transition the FB asymmetry measure by itself cannot provide complete evidence of the handedness of the $b \to c$ transition. They go on to present a viable model of $b \to c$ transitons involving a right chiral boson W_R that is in accord with all present data including the present A_{FB} measurements [37].

The FB asymmetry measure defined in (14) constitutes a momentum-momentum correlation measure $\langle \vec{\ell} \cdot \vec{p} \rangle$ which clearly is not a truly parity-violating measure. For example, it is well-known that in e^+e^--annihilation the two-photon exchange contribution also gives rise to nonvanishing FB asymmetries despite of the fact that QED is parity conserving. In order to determine the handedness of the $b \to c$ transitions in an unambiguous manner one needs to define a truly parity violating spin-momentum correlation measure as e.g. $\langle \vec{\sigma} \cdot \vec{p} \rangle$. Ideal in this regard would be the semileptonic cascade decays $\Lambda_b \to \Lambda_c(\to \Lambda_s \pi) + \ell^- + \bar{\nu}_\ell$ using the known asymmetry structure [38] of the nonleptonic $\Lambda_c \to \Lambda_s \pi$ decays to anaylze the polarization of the Λ_c. In as much as the Λ_c is made from a c-quark and a s-wave spin-zero diquark the Λ_c carries the full spin information of the c-quark as it emerges from the $b \to c$ transition. In this sense such an analysis is sensitive to the parity violating measure $\langle \vec{p}_{\Lambda_s} \cdot \vec{\sigma}_{\Lambda_c} \rangle$ and would allow one to unambiguously conclude for the handedness of the $b \to c$ transitions. There is no doubt that such an analysis will be done in the near future.

3. HEAVY QUARK SYMMETRY AND HEAVY MESON TRANSITION FORM FACTORS

Consider the semileptonic decay of a bottom meson to a charm meson as drawn in Fig.7. The bottom quark at four-velocity v_1 makes a transition to a charm quark at four-velocity v_2 by emitting a virtual W^-. The light "spectator" quark system, which propagates independently, is dragged along to expedite it from the velocity v_1 to v_2 without, however, touching its spin.[4] The spin neutral velocity kick (or alignment)

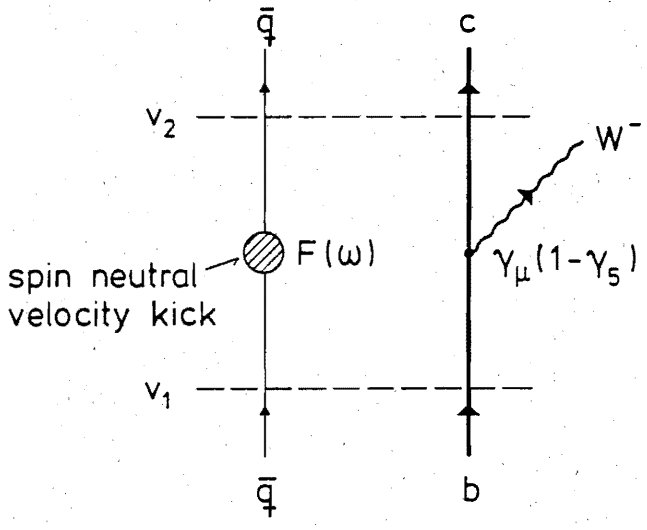

Fig.7: Current induced transition between heavy mesons $\overline{B} \to D, D^*$. The $\overline{B}$ at velocity v_1 makes a transition to the D, D* at velocity v_2. Spin neutral velocity kick form factor $F(\omega)$ of the light spectator quark system is function of Lorentz scalar $\omega = v_1 v_2$.

can be conceived of to result from the exchange of many soft gluons between the c-quark at velocity v_2 and the spectator system which starts off at velocity v_1 and ends up with velocity v_2 in order to align its velocity with the c-quark. Compared to the time scale of the b→c transition the alignment process is slow. The exchanged gluons are all of the longitudinal non flip type, i.e. there is no spin information transferred from the heavy side to the light side. This can be made manifest on the heavy side by splitting the gluon's γ_μ coupling into its spin flip and spin non flip components, viz.

$$\gamma_\mu = \underbrace{(\gamma_\mu - v_\mu)}_{\text{flip}} + \underbrace{v_\mu}_{\text{non flip}} \tag{16}$$

[4] I always like to paraphrase the circumstance that the final quarks move with the same one velocity v_2 by saying "Quarks that fly together stay together", freely borrowing from the American saying "Families that pray together stay together".

The spin flip coupling (γ_μ-v_μ) vanishes in the heavy mass limit and one remains with the Bloch-Nordsieck type non flip coupling v_μ.

From what was said it is clear that the weak amplitude $\overline{B} \rightarrow (D, D^*) + W^-$ factorizes into a heavy-side and a light-side transition amplitude. The only information that is exchanged between the heavy and the light-side is velocity information necessitated by the requirement to reassemble the final charm quark and the light antiquark system in the same final charm hadron. The dynamics of the heavy-side transition $b \rightarrow c + W^-$ is known. It is specified by the usual SM left-chiral weak coupling with a coupling strength proportional to V_{bc}. The light-side transition involves the one unknown transition probability

$$|\text{antiquark}; v_1, \lambda_1\rangle \rightarrow |\text{antiquark}; v_2, \lambda_2 = \lambda_1\rangle$$

which we parametrize by the form factor function $F(\omega)$.[5] It can only be a function of $\omega = v_1 \cdot v_2$ since the velocity transfer variable ω is the only Lorentz invariant variable that can be constructed in the light-side transition. At zero recoil, when $v_1 = v_2$ and $\omega=1$, the antiquark goes through unhindered with amplitude1 and thus we have the normalization condition $F(\omega=1)=1$. It is clear that one has to identify $F(\omega)$ with the reduced form factor function $F(\omega)$ of Sect.2. One expects $F(\omega)$ to fall when ω moves away from the zero recoil limit as it costs to provide the velocity kick. Pole-type form factors and explicit model calculations confirm this expectation.

Since the zero recoil normalization condition of HQS is of such central importance let us have another look at it from a different point of view. Replace the final state c-quark in Fig.7 by a b-quark with the same velocity. This is a symmetry operation as shown in Sect.1. At zero recoil, and for the vector current part of the transition, the normalization $F(\omega=1)=1$ now is nothing but the well-familiar charge form factor normalization at $q^2=0$ applied to the elastic $\overline{B} \rightarrow \overline{B}$ transition. Still another way of looking at the zero recoil normalization condition is afforded by considering the b-quark rest configuration in Fig.2. Replacing the b-quark at rest by a c-quark at rest, as happens in the decay at zero recoil, will not affect the wave function of the light antiquark system. Thus, the overlap between the wave functions before and after the $b \rightarrow c$ transition is complete giving again the zero recoil normalization condition.

Let us now turn to the spin properties of bottom to charm transitions as implied by the spectator quark picture Fig.7. As has been emphasized before there is complete spin factorization of the heavy-side and light-side transitions. It is therefore clear that, for a given heavy side spin configuration, the product amplitudes involving light-side spin up and spin down transitions are identical to one another (up to phases).

[5] Parity relates the two helicity transitions $\lambda_1=\lambda_2=1/2$ and $\lambda_1=\lambda_2=-1/2$.

In this way all product spin configurations become related to one another and, in the final step, one projects the spin configurations onto the relevant heavy meson states. This is essentially how the HQS predictions were originally derived in the helicity matching approach of [19,20,39] (see also [40]) and by Isgur and Wise [6].

Needless to say that it is a cumbersome and error-prone venture to obtain the correct HQS form factor relations using the above approach. In particular, one has to be careful to use a consistent set of phases for the angular momentum coupling factors. Much more convenient in this regard is the so-called covariant trace formula [7,19,20,41-44]. In this method the requisite angular momentum projections onto the heavy meson states are all done covariantly. The transition matrix element for $\bar{B} \to D, D^*$ can then be written very compactly as

$$M_\mu^{V-A} = \frac{1}{4}\sqrt{M_1 M_2}\,\mathrm{Tr}\left\{(\underbrace{\gamma_5}_{D} + \underbrace{\not{\varepsilon}_2^{\,*}}_{D^*})(\not{v}_2 + 1)\gamma_\mu(1-\gamma_5)(\not{v}_1 + 1)\underbrace{\gamma_5}_{\bar{B}}\right\}F(\omega) \qquad (17)$$

where the γ_5 and $\not{\varepsilon}$ covariantly project onto the $J^P=0^-$ and 1^- states, respectively, as indicated in Eq.(17). The positive energy projectors $(\not{v}_1 + 1)$ and $(\not{v}_2 + 1)$, on the other hand project onto the on-mass-shell heavy quark states. $F(\omega)$ is the reduced HQS form factor, as before, with $F(\omega=1)=1$. The mass factors in front of the trace, finally, provide the appropriate heavy mass power scaling factors. Let me mention that analogous covariant "trace" formulas have been written down for baryon ground state transitions [44,45-48] and for transitions to mesonic p-wave states [17]. Ref.[49] lists covariant wave functions for any spin and HQS "trace" formulas for transitions between them.

It is quite instructive to also consider the case when only the initial meson in the transition $1 \to 2$ is heavy, as for example in the heavy to light transition $B \to \pi$. The heavy quark symmetry now only applies to the initial meson (at velocity v_1). The corresponding trace formula for the $0^- \to 0^-$ case is obtained from (17) by the replacement

$$F(\omega)\gamma_5(\not{v}_2 + 1) \to A(p_2 v_1)\gamma_5 \not{v}_2 + B(p_2 v_1)\gamma_5 \qquad (18)$$

The two covariants $\gamma_5 \not{v}_2$ and γ_5 no longer combine in the form of the mass-shell positive energy projector $\gamma_5(\not{v}_2 + 1)$ since the light active quark in the final state can no longer be considered to be on its mass-shell. The transition is now described by two unnormalized form factors $A(p_2 v_1)$ and $B(p_2 v_1)$ instead of the one normalized HQS form factor $F(\omega)$. As the $0^- \to 0^-$ transition altogether involves only two transition form factors (see Eq.(7a)) one might ask what has one gained? In fact, in terms of the spin symmetry nothing has been gained, but then one has not yet exploited the flavour symmetry implicit in (18). This is at the basis of the suggestion made in Ref.[50] to relate the $D \to \pi$ and $D \to K$ form factors to the $B \to \pi$ form factors. In the latter case one

first uses heavy flavour symmetry in the form of Eq.(18) to relate $D \to K$ to $B \to K$, and then to relate $B \to K$ to $B \to \pi$ via an SU(3) rotation. The practical implimentation of this suggestion appears to be rather difficult as dicussed in [51].

The corresponding heavy to light transition formula for $0^- \to 1^-$ transitions can be obtained from Eq.(17) by the replacement

$$F(\omega)\not{\epsilon}_2^*(\not{v}_2+1) \to A'(p_2v_1)\not{\epsilon}_2^*\not{v}_2 + B'(p_2v_1)\not{\epsilon}_2^* + C'(p_2v_1)\not{\epsilon}_2^*\not{v}_2\not{v}_1 + D'(p_2v_1)\not{\epsilon}_2^*\not{v}_1 \tag{19}$$

i.e. there are now four unnormalized form factors instead of the one reduced form factor $F(\omega)$. Note also that the four form factors A',B',C' and D' are not related to the form factors A and B in Eq.(18). Considering the general structure Eq.(7b) there is again no gain as far as spin symmetry is concerned. However, exploiting the flavour symmetry implicit in (19), one can again attempt to relate $D \to \rho$ and $D \to K^*$ to $B \to \rho$ by the above chain of arguments.

Let us summarize our results in the form of Table 1 which shows how the general form factor structure with six form factors eventually reduces to one HQS form factor when both the initial and final meson are heavy. We have also included the predictions of the HQET when "$1/m_c$" corrections are included for the final charm quark. This will be the subject of the next section.

As regards the reduction of form factor complexity from six to one one can draw an analogy to deep inleastic scattering where the Callan-Gross relation reduces the two structure functions W_1 and W_2 to one (transverse) structure function. Just as is the case in deep inelastic scattering, where the reduction to the transverse structure function tells us about the spin of the active quark, the HQS form factor structure in the last column of Table 1 reflects on the spin 1/2 nature of the "active" heavy quarks and their specific form of interaction that participate in the current transitions.

4. HEAVY QUARK EFFECTIVE THEORY, THE $1/m_Q$ EXPANSION AND LUKE'S THEOREM

Our aim in this section is two-fold. First we want to describe how one can rewrite the QCD Lagrangian of a heavy quark field in terms of a $1/m_Q$ power expansion by using a series of Foldy-Wouthuysen-type field redefinitions. Second by using diagrammatic techniques, we want to demonstrate that the $O(1/m_Q)$ contributions to current-induced transitions leave the zero recoil normalization condition $F(\omega=1)=1$ intact.

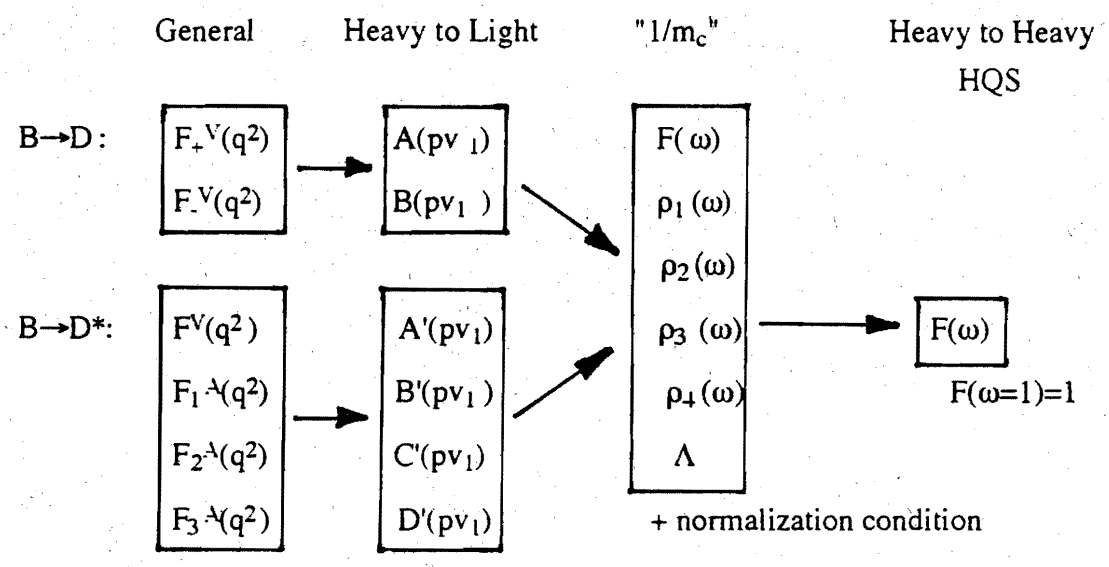

Table 1: Form factor complexity for meson transitions. Listed are sets of independent form factors in the general case (first column), for heavy to light transitions (second column), for heavy to light transitions with $1/m_Q$ corrections for the secondary meson (third column) and for heavy to heavy transitions in the HQS limit (fourth column).

In order to get a handle on the physics of the $1/m_Q$ expansion let us first discuss heuristically how the QCD Lagrangian for a heavy quark field can be made to look quite simple in the heavy mass limit. Consider the QCD Lagrangian for a heavy quark field

$$L = \bar{Q}(i\gamma_\mu D^\mu - m_Q)Q \tag{20}$$

In the heavy mass limit one can safely drop the three-momenta relative to the heavy mass, viz.

$$E = \sqrt{m_Q^2 + \vec{p}^2} = m_Q(1 + \frac{1}{2}\frac{\vec{p}^2}{m_Q^2} + \ldots) \tag{21}$$

That is, one can drop the three-derivative in the Lagrangian (20) and obtains

$$L = \bar{Q}'(i\gamma_o D^o - m_Q)Q' \tag{22}$$

The new fields entering the modified Lagrangian (22) are denoted by Q'. Further one can shift the energy scale $E \rightarrow E' = E - m_Q$ so that energies are now measured from m_Q rather than from zero. The Lagrangian (22) then simplifies to

$$L = \bar{Q}'' i\gamma_o D^o Q'' \tag{23a}$$

If one chooses to work only in terms of the positive energy components $\tilde{Q} = (1+\gamma_o)Q$ of the heavy quark field the Lagrangian (23a) further simplifies to

$$L = \tilde{Q}''^{+} iD^o \tilde{Q}'' \tag{23b}$$

Using the positive energy fields $\tilde{Q}$ may be opportune since the quark and antiquark sectors of the theory decouple for infinitely heavy quark masses. The Lagrangian (23b) is the celebrated static Lagrangian of Eichten and Hill [11].

Sometimes one sees the Lagrangians (23a) and (23b) in somewhat different guises. One introduces an auxiliary velocity four-vector v_μ with $v^2=1$ which, in the static limit, reads $v_\mu=(1,0,0,0)$. Using the four vector v_μ the Lagrangians (23a) and (23b) now read

$$L = \bar{Q}'' i\not{v} D\cdot v\, Q'' \tag{24a}$$

$$L = \tilde{Q}''^{+}\, iD\cdot v\, \tilde{Q}'' \tag{24b}$$

The step-wise reduction of the QCD Lagrangian (20) to the final form (24) can be achieved by a series of Foldy-Wouthuysen-type field redefinitions which eventually yields the leading term results (24) plus all higher dimension operators in the $1/m_Q$ expansion [4]. The first transformation (change of variable) is (j=1,2,3)

$$\begin{aligned} Q &= \exp(i\gamma_j \overleftarrow{D}_j / 2m_Q) Q' \\ \bar{Q} &= \bar{Q}' \exp(-i\gamma_j \overleftarrow{D}_j / 2m_Q) \end{aligned} \tag{25}$$

where the arrow on the derivative indicates in which direction the derivative acts ($\int f\overleftarrow{D}g = -\int f\overrightarrow{D}g$). The heavy quark action

$$S = \int d^4x\, \bar{Q}(i\not{D} - m_Q)\, Q \tag{26}$$

now becomes

$$S_{Q'} = \int d^4x \left\{ \bar{Q}'(i\gamma_o D_o - m_Q) Q' + \sum_{i=1}^{\infty} \left(\frac{1}{2m_Q}\right)^i \bar{Q}'\, \theta_i Q' \right\} \tag{27}$$

giving a form of the action which makes explicit the mass perturbations.

The second transformation which removes the remaining mass dependence in the numerator factors is

$$\begin{aligned} Q' &= \exp(-i\, m_Q \gamma_o t)\, Q'' \\ \bar{Q}' &= \bar{Q}'' \exp(i\, m_Q \gamma_o t) \end{aligned} \tag{28}$$

It is easy to see that the transformation (28) does in fact remove the mass dependence of the leading term in (27), as desired. However, a little bit of more work has to be done on the higher dimension operators 0_i in (27) after the first transformation to bring them into a form suitable for the final transformation (28). This intermediate step is rather technical and will not be discussed any further here. More details can be found in Ref.[4].

Using again the covariant form with the velocity v_μ the final form of the HQET action reads[6]

$$S_{Q''} = \int d^4x \left\{ \bar{Q}''i\not{v}D\cdot v\, Q'' + \sum_{i=1}^{\infty} \left(\frac{1}{2m_Q}\right)^i \bar{Q}''0_i Q'' \right\} \tag{29}$$

The higher dimension operators are easily determined from the transformations, for example[7]

$$\begin{aligned} 0_1 = & -\frac{1}{2}D^2 + \frac{1}{2}(D\cdot v)^2 + \frac{i}{4}g\sigma_{\mu\nu}F^{\mu\nu} \\ & -\frac{i}{2}g\gamma_\mu \not{v} F^{\mu\nu}v_\nu \end{aligned} \tag{30}$$

The higher order terms in the $1/m_Q$ expansion are obtained by using the Baker-Hausdorff formula [4]. Since one is working within the path integral framework to derive (29) one thing that must be checked is whether the path integral measure is invariant under the change of variable. Using dimensional regularization this is shown to be the case in Ref.[4], however, other regularization schemes may require the addition of local counterterms.

In other formulations of the effective theory, where the operators 0_i are determined on dimensional grounds, the relative coefficients of the operators that make up the 0_i are determined by "matching conditions". These matching conditions arise on a comparison of the effective theory with the true theory. In our approach no matching is required, it is built in, as we derive the operators directly from the QCD Lagrangian. Various issues, such as the gauge dependence and renormalizability of the effective theory compared to the complete theory are addressed in Ref.[4].

All the heavy mass dependence of the action (29) has been put into the explicit $1/m_Q$ power dependencies. From the heavy quark action one can derive explicit Feynman rules. For example, the 0(1) term generates a mass independent propagator and the Bloch-Nordsieck-type v_μ $Q\bar{Q}G$ velocity coupling (the $\not{v}$ in the 0(1) term is always placed next to a $(\not{v}+1)$ projector and is thus irrelevant). The $0(1/m_Q)$ term involves, among others, $Q\bar{Q}G$ and $Q\bar{Q}GG$ vertices with vertex strength proportional to $1/m_Q$ and couplings specified by (30). The heavy fields know nothing about the mass

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[6] The expansion (29) with its subsequent higher and higher dimension operators may be likened to the operator product expansion in deep inelastic scattering. Although the expansion (29) is quite useful in that it provides all the higher dimension operators and their relative weights the toil does not stop here. One still has to calculate or parametrize their matrix elements for a given transition. This involves unknown soft physics, just as in deep inelastic scattering. Furthermore, when loop effects are taken into account the higher dimension operators aquire anomalous dimension factors which have to be calculated.

[7] Note that the last term in (30) vanishes when taken between positive energy fields $\tilde{Q} = (1+\not{v})Q$.

that is associated with their particular flavour. This makes the equal velocity heavy flavour symmetry associated with the 0(1) term in (29) quite explicit.

After having written down the HQET action (29) I am now in a position to discuss current transitions in a diagrammatic language using the first few operator terms in (29). I am particularly interested in the zero recoil configuration where $v_1=v_2$. The general case $v_1 \neq v_2$ is discussed in [3] and leads to the $1/m_c$ symmetry breaking pattern given in Table 1.

In Fig.8 I have drawn some low order diagrams for inelastic $b \to c$ and elastic $c \to c$ current transitions in the $1/m_c$ expansion, meaning that I choose to treat the bottom sector at 0(1). I have already specified to the zero recoil configuration $v_1 = v_2 = v$. I have omitted drawing any gluon lines in Fig.8 as they are not relevant for the following arguments. They could e.g. connect the insertions with the light spectator system.

The 0(1) contributions corresponding to the Feynman diagrams (8a) and (8b) are identical to one another as emphasized before. The b and c propagators as well as the $b \to c$ and $c \to c$ current vertices are the same (in fact at any $v_1 \neq v_2$). From charge normalization one knows that the vector current transition in the elastic $c \to c$ case is normalized to 1, so is the inelastic $b \to c$ transition. This is the Voloshin-Shifman zero recoil normalization previously referred to [2].

The normalization survives even in the presence of $0(1/m_c)$ corrections [3,4,52,53]. This statement is generally referred to as Luke's theorem. There are two sources of $0(1/m_c)$ corrections. The first is represented in Fig.8e and comes from the transformation of the current

$$\overline{Q}_b J_\mu Q_c \to \overline{Q}_b^{\cdot} J_\mu Q_c^{\cdot} + \overline{Q}_b^{\cdot} J_\mu (D_i \gamma_i / 2m_c) Q_c^{\cdot} + 0(1/m_c^2) \qquad (31)$$

However, it is easy to see that the "current insertions" are zero at zero recoil. The reason is that the current vertex in Fig.8e is sandwiched between the projectors $(\gamma + 1) \hat{=} (\gamma_o + 1)$. Now, at zero recoil, one has contributions only from J_o^V and J_j^A (see also the discussion in Sec.5). But then, both $(\gamma_o+1)\gamma_o\gamma_i(\gamma_o+1)$ and $(\gamma_o+1)\gamma_j\gamma_i\gamma_5(\gamma_o+1)$ vanish.

The second source of $1/m_c$-corrections comes from the fact that there are the $0(1/m_c)$ "propagator insertions" stemming from 0_1 in (30). For $(c\overline{q}) \to (c\overline{q})$ this is shown in Fig.8d and can be seen to be twice the $(b\overline{q}) \to (c\overline{q})$ contribution Fig.8c. But as this $0(1/m_c)$ insertion must vanish for $(c\overline{q}) \to (c\overline{q})$ from charge normalization it must also vanish for $(b\overline{q}) \to (c\overline{q})$. At $0((1/m_c)^2)$ the diagonal and nondiagonal charm propagator insertions are no longer proportional to one another as the sample diagram Fig.8f shows..

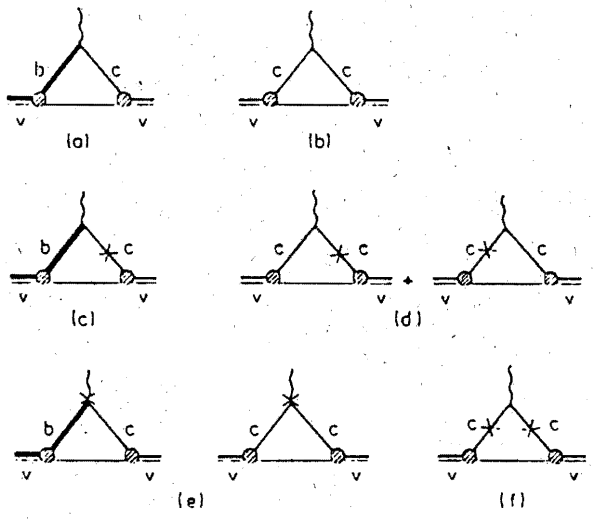

Fig.8: 0(1) ((a) and (b)), $0(1/m_c)$ ((c), (d) and (e)) and $0(1/m_c^2)$ (f) contributions to mesonic b→c and c→c transitions at zero recoil $v_1 = v_2 \equiv v$.

The above proof of the vanishing of the $1/m_c$ propagator insertions in $(b\bar{q}) \to (c\bar{q})$ holds for vector current transitons between alike meson states, for example for $B \xrightarrow{V_\mu} D$. It can be extended to the general case involving also axial current transitions between different degenerate HQS states, as e.g. $B \xrightarrow{A_\mu} D^*$, by using also the spin symmetry of HQS (see e.g. [54]). It is clear that the above diagrammatic proof of Luke's theorem does not depend on the composition of the light constituents at the bottom of the diagrams Fig.7. They could be light diquarks, as in the case of heavy baryons [54,55] or even light supersymmetric particles [52]. The above proof of the absence of $1/m_c$ corrections to the zero recoil normalization condition is also valid in these cases.

5. NORMAL AND ANOMALOUS DIMENSIONS AND HEAVY MASS SCALING

In physics it is often possible to obtain information on the qualitative behaviour of physical quantities by applying dimensional scaling arguments. A good example at hand is given by the cross section behaviour of $e^+e^- \to \mu^+\mu^-$ or $e^+e^- \to q\bar{q}$ at high energies. On purely dimensional grounds one finds

$$\sigma(e^+e^- \to \mu^+\mu^-) \sim 1/s\,. \tag{32}$$

The argument goes as follows. The dimension of cross section (area, as the name says) is length squared, or, in terms of mass dimension, $[m^{-2}]$. As the CM energy squared s is the only relevant mass scale in the problem at high energies one concludes for the 1/s high energy behaviour of the e^+e^--cross section, viz. (32). The high energy behaviour can of course be checked by doing an explicit calculation, but the scaling argument is

much quicker. Mind you, the rapid fall-off of the e^+e^--cross section given by (32) is of course a serious concern for machine builders and their funding agencies.

Similar scaling arguments apply to heavy hadron physics in terms of heavy mass scaling. Take for example the mass dependence of the pseudoscalar decay constant f_P which is defined by

$$\langle 0|A_\mu|P(p)\rangle = f_P p_\mu \tag{33}$$

as usual. We choose the conventional state normalization

$$\langle P(p')|P(p)\rangle = 2E\,\delta^3(\vec{p}-\vec{p}\,') \,. \tag{34}$$

This choice is really not convenient when one wants to discuss the infinite mass limit since the normalization factor E tends to infinity at the same time. A better choice is to normalize to E/M, i.e. to work in terms of states $|P(p')\rangle'$ with

$$|P(p)\rangle' = |P(p)\rangle / \sqrt{2M} \tag{35}$$

The primed states $|P(p)\rangle'$ are then independent of the heavy mass scale.

Let us accordingly divide both sides of (33) by $\sqrt{2M_P}$. One obtains

$$\langle 0|A_\mu|\frac{P(p)}{\sqrt{2M_P}}\rangle = \frac{1}{\sqrt{2}} f_P \sqrt{M_P}\, v_\mu \tag{36}$$

In the heavy mass limit the left-hand side of Eq.(36) is heavy mass scale indenpendent and thus one concludes that $f_P\sqrt{M_P} \sim$ constant.[8] The heavy mass scale independence of the l.h.s. of (36) comes about since the leading contribution of the full axial vector current A_μ is given by the HQET 0(1) current $A_\mu^{(0)}$, at least at the tree level. But, as shown in Sec.4, the leading term in the $1/m_Q$ expansion of the axial vector current Eq.(31) is heavy mass scale independent.

The scaling behaviour of f_P relates the decay constants of pseudoscalar mesons with different flavours as e.g. f_B and f_D. One obtains [2,56]

$$\frac{f_P}{f_{P'}} = \sqrt{\frac{M_{P'}}{M_P}}\left(1 + 0\left(\frac{1}{M_P}, \frac{1}{M_{P'}}\right)\right) \left(\frac{\alpha_s(M_P)}{\alpha_s(M_{P'})}\right)^{-\frac{6}{33-2N_f}} \tag{37}$$

The first factor in (37) represents the "normal" scaling behaviour of f_P according to the "normal" power behaved heavy mass dimensions of f_P.[9] The last factor in (37)

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[8] One recognizes that $f_P\sqrt{M_P}$ is numerically proportional to the wave function at the origin of the heavy bound state system. But, as shown in Sect.1, the wave function of the heavy bound system does not depend on the heavy mass located at the centre. This is another way of ascertaining the heavy mass scale independence of $f_P\sqrt{M_P}$.

[9] At present the question of whether f_D and f_B fall into the heavy mass scaling regime is being widely discussed. In fact, lattice and QCD sum rule calculations indicate that $0(1/m_Q)$ corrections are still quite important in the charm and bottom region. Unfortunately the discussion remains rather academic as long as f_D and f_B have not been measured, and the prospects of doing so in the next few years are rather dim.

represents the "anomalous" scaling behaviour of f_P arising from logarithmic scaling violations (or anomalous dimensions of f_P) whose source are loop corrections to the 0(1) HQET axial vector current $A_\mu^{(0)}$. The loop corrections are always factorizable and lead to a modification of the 0(1) axial vector current by an anomalous dimension factor $A_\mu^{(0)} \rightarrow C(M_P/\mu)A_\mu^{(0)}$ [2,56]. The anomalous dimension factor $C(M_P/\mu)$ has been calculated, at least approximately, by summing the leading log loop contributions to the axial vector current giving the anomalous dimension factor in Eq.(37). Clearly one can again draw an analogy to the physics of deep inelastic scattering where one also has normal and logarithmic anomalous contributions in the scaling regime.

Similar dimensional arguments may be used to extract the "normal" heavy mass dimensions of the transition form factors discussed in the previous sections that give rise to the square root mass factors in (17). As regards the logarithmic anomalous dimensions there are now also velocity dependent contributions in addition to the anomalous contributions encountered in the calculation of f_P [8]. Although of considerable theoretical interest one must say, though, that the anomalous contributions are generally quite small numerically.

6. DETERMINATION OF THE KM MATRIX ELEMENT V_{bc}

Let us return to where we have set out at the beginning of this review, namely the measurement of the KM matrix element V_{bc} using the exclusive semileptonic $B \rightarrow D, D^*$ transitions.

As has been particularly emphasized by Neubert and his Heidelberg collaborators the $B \rightarrow D^*$ transitions are much better suited for a determination of the KM matrix element V_{bc} than $B \rightarrow D$ transitons [57,58]. Why this is so can be understood in quite simple terms by analyzing the partial wave content of the final states in the two decays $B \rightarrow D+W$ and $B \rightarrow D^*+W$. I shall present the argument in quite some detail as this way of looking at the decay process sheds a lot of light on the physics of the decays.

In Fig.9 I have drawn a decay diagram for the decay process in terms of the J^P quantum numbers of the particles involved. The two processes are $0^- \rightarrow 1^- + W(1_A^+)$ and $0^- \rightarrow 1^- + W(1_V^-, 1_A^+)$ where the J^P values in the brackets refer to the vector and axial vector pices of the off-shell W. The spin 0 components of the off-shell W (0^+ and 0^- for V_μ and A_μ, respectively) are not active in the decays since the lepton masses m_e and m_μ are negligibly small on the scale of the Q-value of the process.

Let us now determine the partial wave content of the decays. One has

$B \to D$:	$0^- \to 0^- + 1_V^-$	p-wave	"first forbidden Fermi transition"	(38a)
$B \to D^*$:	$0^- \to 1^- + 1_V^-$	p-wave	"first forbidden Fermi transition"	(38b)
	$0^- \to 1^- + 1_A^+$	s-wave;	"allowed Gamow-Teller transition"	
		d-wave	"second forbidden Gamow-Teller transition"	

There are one and three partial wave amplitudes for the $B \to D$ and $B \to D^*$ transitions, i.e. the same number as counted in Sec.2. As $\omega \to 1$ and the phase space closes[10] only the s-wave contribution in $0^- \to 0^- + 1_A^+$ survives. Let me mention that in nuclear physics parlance the vector and axial vector transitions are referred to as Fermi and Gamow-Teller transitions, respectively. They are further classified according to their partial wave threshold behaviour by "allowed", "first forbidden" etc. The nuclear physics classification has been included in Eq.(38). One might then ask why one does not see s-wave "allowed Fermi transitions" in particular in the $B \to D$ case. The answer is that these are mediated by the 0^+ vector current component which is not active in the $B \to D, D^*$ transitions. This is different in nuclear transitions where the Q-values of the transitions are comparable to the electron mass and thus the 0_V^+ component does in fact become active. The reason that the $B \to D$ transition is not well suited for the measurement of V_{bc} relates to the absence of an "allowed" or s-wave transition in this transition as will be explained in the following.

After this brief excursion into nuclear physics I return to the partial wave amplitudes in (38). It is very important to realize that the zero recoil normalization condition at $\omega=1$, true at $O(1)$ and at $O(1/m_c)$ (and also at $O(1/m_b)$ for that matter), holds true only for the s-wave amplitudes. In this sense there is no zero recoil normalization condition for $B \to D$ and therefore no way to measure the "weak transition charge" V_{bc} in this process.[11] This is different for $B \to D^*$ where one has an s-wave contribution normalized at $\omega=1$ according to the "allowed Gamow-Teller transition" in (38b). In addition, there is sufficient rate close to $\omega=1$ in $B \to D^*$ transitions due to the presence of the s-wave contribution. The presence of such a pseudothreshold behaviour at large q^2 is strinkingly evident in the q^2 spectra of the $\overline{B} \to D$ and $\overline{B} \to D^*$ transitions as shown in Fig.10. An extrapolation of the experimental data into the zero recoil point at the boundary of phase space will thus be tractable if there is enough data available in the pseudo-threshold region.

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[10] Those readers that are old enough will remeber that the zero recoil point $\omega=1$ used to be also called the pseudo-threshold. "Pseudo" because phase space closes when the zero recoil point is approached in constrast to the normal threshold where phase space opens.

[11] Here we are always talking about semileptonic $B \to D$ decays involving electrons and muons. The decays $B \to D + \tau + \nu_\tau$ do involve "allowed" s-wave contributions but are hard to identify experimentally.

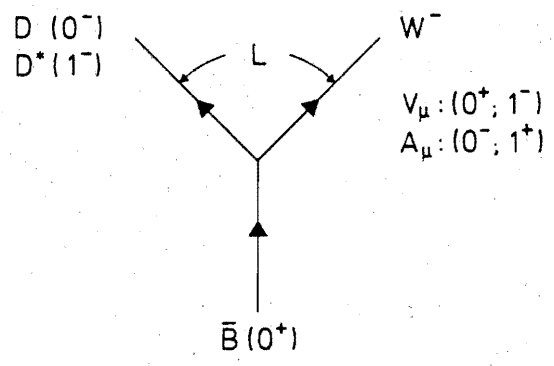

Fig.9: Quasi two-body transitions $\overline{B} \to D^*$ and $\overline{B} \to D^* + W^-$. Round brackets contain J^P quantum numbers of states. Vector and axial vector current coupling pieces of the off-shell W^- have $(0_V^+, 1_V^-)$ and $(0_A^-, 1_A^+)$ quantum numbers, respectively. In the zero lepton mass limit only the 1_V^- and 1_A^+ components are active.

Such an analysis has been done in Ref.[58] using the data on $B \to D^* + \ell\nu_\ell$. The analysis incorporates new results on the renormalization of the heavy quark currents (the logarithmic "anomalous" dimension effects described in Sec.5) at the next-to-leading order level [59]. The outcome of the calculation of [59] is that the heavy quark current gets renormalized downward by 1% at ω=1 as compared to an upward renormalization of 13% when only leading log-effects are included in the renormalization [8]. Using an updated value for the total branching ratio of $Br(B \to D^* + \ell\nu_\ell) = 4.4\% \pm 0.5\%$ [60] the analysis of the $B \to D^* + \ell\nu_\ell$ data leads to

$$|V_{bc}| = 0.04 \pm 0.005 \tag{39}$$

for τ_B = 1.3 ps. At present one is incurring large errors when one is extrapolating the data into the zero recoil point because present data close to the zero recoil point is not very good. Undoubtedly the situation will improve much in the next few years as more data is collected. One can thus hope to considerably reduce the above error in Eq.(39).

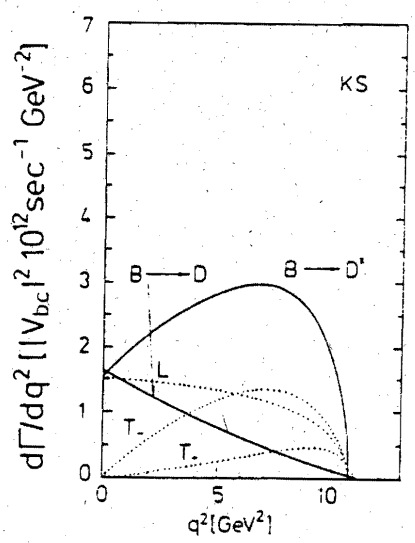

Fig.10: q^2-spectra of $\overline{B} \rightarrow D$ and $\overline{B} \rightarrow D^*$ transitions in the KS model [19,20].

7. CONCLUDING REMARKS

There are many more applications of Heavy Quark Symmetry and the Heavy Quark Effective Theory that have not been touched on in this review talk. For example, there was only little mention of heavy baryons and current-induced transitions between them. Among the heavy baryons, the (Qqq) baryons with one heavy quark provide an ideal stage to study the interaction of the two accompanying light quarks with one another and with the vacuum. Also quite interesting from the theoretical point of view are (QQq) baryons with two heavy quarks. When the two heavy quarks are in the lower-lying antisymmetric 3* colour state, the heavy-light system as a whole is identical to a ($\overline{Q}$q) heavy-light meson system, and one would have an exact supersymmetry relating the two! Heavy Λ-type baryons are well suited to study polarization effects in the context of HQS where some intriguing quantum coherence effects are at work when polarized heavy quarks hadronize into heavy hadrons. This is of great topicality since b-quarks from Z-decays at LEP are predicted to be $\approx$ 97% polarized [61].

In this review talk I have only been concerned with heavy hadron transitions at low recoil. The large recoil regime can be conveniently studied in the Brodsky-Lepage hard scattering picture. At large recoil one has a new type of heavy quark symmetry which is reminiscent but not identical to the HQS at low recoil [62].

There was not much discussion on excited heavy meson nor on excited heavy baryon states. Their contribution to heavy quark sum rules and closure relations provide some interesting constraints on the reduced HQS form factors [7,16,17,63,64].

Neither did I discuss the interesting implications of HQS for nonleptonic decays of heavy mesons and baryons.

Recently there have been some interesting developments in studying one- and multi-pion transitions between heavy hadrons by unifying heavy quark symmetry and spontaneously broken chiral symmetry [18,65]. Also photon transitions between heavy hadrons can be conveniently dealt with by unifying heavy quark symmetry with the gauge symmetry of electromagnetism [66].

Present model calculations based on QCD sum rules, on lattice simulations and on explicit wave function models have to be pushed much further. Thereby one could accurately determine the much needed functional form of the ω-dependence of the HQS reduced form factors. The model results would allow one to determine the size of the $1/m_Q$ and $1/m_Q^2$ corrections to the heavy quark symmetry limit and would lend credibility to applying heavy quark physics ideas in the bottom and charm regime. There is also the possibility of a surprise that some trace of heavy quark symmetry survives in the strangeness sector.

On the formal side one wants to make the statements and definitions in the formulation of heavy quark symmetry more precise [54].

The field is still quite alive and I am sure we shall see a lot of activity in heavy quark physics in the next few years. As experience has shown real progress is achieved when theoretical and experimental advances go hand in hand. In this sense I am looking forward to a lot of new experimental data on heavy quark physics in the coming years.

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