

Heavy Quark Symmetry and Chiral Dynamics

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Abstract

We report on recent advances in the understanding of hadrons containing one heavy quark, providing a general introduction to the main ideas behind the heavy quark symmetry. We concentrate on the synthesis of heavy quark and chiral symmetries, describing recent calculations of heavy meson and baryon decays through processes involving strong, electromagnetic and semileptonic weak interactions.

1. Introduction

There is very little direct connection between what will be described here and the main topics of this conference, such as the determination of quark confinement and chiral symmetry breaking in QCD. On the other hand, it will be shown below that by simply accepting these ideas, along with some general properties of QCD, one can make some very precise statements and also extract a lot of physics from a class of hadrons which contain a single heavy quark, even though one does not know the exact solutions of QCD. So perhaps in this connection we can say that once the goal of this conference is achieved, we will be able to remove the assumptions of confinement and chiral symmetry breaking, and everything that will be described here will then be derivable from first principles.

Because of the fact that the fields which enter the QCD Lagrangian are current quarks and gluons, but what we see in the laboratory are hadrons — that is, $q\bar{q}$ and qqq bound states — and that the connection between the field quanta and hadrons is still not completely understood, it is fair to say that all the consequences of QCD which can be demonstrated at present are those which follow from the symmetries. For example, one of the exact symmetries is the gauge symmetry, which because of its nonabelian nature leads to asymptotic freedom, and this in turn

allows us to do calculations at short distance and make a lot of predictions for high-energy phenomena. The other two symmetries are approximate; one is chiral symmetry and is a main topic of this conference. We all know this symmetry comes from the fact that the light quarks have masses which are very small compared with Λ_{QCD} , the typical energy scale of QCD. But recently there has also been a new development in the understanding of a second approximate symmetry of QCD — namely, the heavy quark symmetry — which is manifest in systems containing a heavy quark and light quarks. In this case the symmetry comes from the fact that the heavy quark has mass much greater than Λ_{QCD} . So there are two classes of quarks, one very light $q_i = (u, d, s)$ and one very heavy $Q_i = (c, b, t)$. Because of the light quarks there is an approximate $SU(3)_L \times SU(3)_R$ chiral symmetry and because of the heavy quarks there is an approximate heavy quark symmetry. So we are very fortunate that there is a division between light quarks and heavy quarks and that there exist systems in which the theoretical predictions derived from the two corresponding approximate symmetries can be tested experimentally.

In this paper we give a general introduction to the ideas behind the heavy quark symmetry and present some example calculations which illustrate their use. The work that will be described here is contained in a set of papers by H.Y Cheng, C.Y. Cheung, W. Dimm, G.L. Lin, Y.C. Lin, T.M. Yan and H.L. Yu¹⁻⁴. There are many other groups which work on similar topics; an incomplete set of references is given in Refs. [5-7]. These are simply meant as pointers to the wide literature on the topic of heavy quark symmetry.

The outline of the paper is as follows. In Section 2 the semileptonic weak decay of the heavy quark system will be used as an example to introduce the ideas of the heavy quark symmetry. Then in Section 3 the heavy quark symmetry will be combined with the chiral symmetry of light quarks, and some examples of heavy quark systems involving strong interactions, electromagnetic interactions, and semileptonic decays will be given and the results compared with experimental data when possible. Some final comments will be made in Section 4.

2. Heavy Quark Symmetry

The ideas behind the heavy quark symmetry have long been known to the practitioners in the field, but for some reason for many years they were not clearly formulated. It was due to Isgur and Wise⁸, Voloshin and Shifman⁹, and many others, that the physics of the heavy quark symmetry was given a very precise formulation. The idea is actually very simple; in fact, all that will be described here is so embarrassingly simple in comparison with the other contributions to this conference that most of the physics can be understood by analogy with the simple hydrogen atom. The basic idea, then, is that a $Q\bar{q}$ heavy meson is similar to the hydrogen atom, with the heavy quark playing the role of the nucleus and the light quark playing the role of the electron. Of course, the theory is different: we have QCD instead of QED, but we can proceed by analogy. We know that in QED, as the mass of the proton m_P becomes infinite compared with the electron mass, we have the following exact consequences: (1) The energy spectrum, wave function, and transition matrix elements, etc. are independent of m_P . In

a $Q\bar{q}$ system this is translated into the independence of all these quantities with respect to m_Q — this is known as flavor symmetry. (2) In the hydrogen atom, the spin interaction is inversely proportional to m_P , so that in the limit $m_P \rightarrow \infty$, spin dependence disappears. Thus there is a spin symmetry, which is translated into independence of the spin of Q in the QCD case. (3) QED allows us to include systematic corrections in $1/m_P$. In QCD, in principle we can include corrections in $1/m_Q$, but in practice this is not always possible because this involves matrix elements which we cannot yet determine. (4) A final idea which should be mentioned is that in these heavy-light systems, velocity is a much better variable than momentum. This we can understand in the rest frame of the hydrogen atom, for example, if we replace the proton by another nucleus with the same charge but a very different mass (such as the deuteron), for then the electron will not know any difference. Thus if we boost the whole system to a velocity v , the wave function does not change, even though the momenta of the two systems are very different. Therefore, velocity is a more convenient variable than momentum.

We can summarize all this physics with a very simple picture of a heavy meson as a heavy quark Q surrounded by a light quark cloud \bar{q} . Flavor symmetry then means the structure of the light quark clouds in $Q_i\bar{q}$ and $Q_j\bar{q}$ are the same. The light quark does not know the difference. Spin symmetry means that the light quark clouds are the same for a heavy quark with spin up as that for a heavy quark with spin down. These are very simple statements about the spin and flavor symmetries.

Let us now look a little more quantitatively at the case of QCD. Suppose a heavy quark moves with a large momentum P_Q , and a light \bar{q} goes with it. Then because the mass m_Q is so heavy, the velocity is hard to change, and we can parameterize the momentum as

$$P_Q = m_Q v + k, \quad m_Q \rightarrow \infty.$$

Here, $v^2 = 1$ and the residual momentum k is of the order Λ_{QCD} . In the limit $m_Q \rightarrow \infty$, this parametrization greatly simplifies the Feynman rules in QCD. For example, the propagator of the heavy fermion becomes

$$\frac{\not{P}_Q + m_Q}{P_Q^2 - m_Q^2} = \frac{\not{v} + 1}{2v \cdot k},$$

and the vertex when sandwiched between the projection operators $(\not{v} + 1)/2$ becomes

$$\frac{\not{v} + 1}{2} \gamma_\mu \frac{\not{v} + 1}{2} = \frac{\not{v} + 1}{2} v_\mu \frac{\not{v} + 1}{2}.$$

Moreover, $(\not{v} + 1)/2$ gives 1 when acting on an external leg. So in the effective theory, the quark-gluon vertex is $-igv_\mu T_a$ and the heavy quark propagator is $(v \cdot k)^{-1}$. Thus the spin and mass of the heavy quark disappear, which are expressions of the spin and flavor symmetries of the system. Of course, there are $1/m_Q$ corrections which are well-defined and can in principle be systematically taken into account; perhaps we will be able to compute them after this conference.

Now let us discuss the ground states of the $Q\bar{q}$ system. First of all, there are the pseudoscalar $P = 0^-$, the D mesons ($c\bar{q}$) and the B mesons ($b\bar{q}$). Then there are the spin one $P^* = 1^-$, the D^* and the B^* mesons. These are assumed to be the ground states. At rest, the spin content of the pseudoscalar and spin projection $\lambda = 0$ spin-one heavy meson wavefunctions can be written

$$|P\rangle = \frac{1}{\sqrt{2}}(|Q \uparrow \bar{q} \downarrow\rangle - |Q \uparrow \bar{q} \downarrow\rangle),$$

$$|P^*, \lambda = 0\rangle = \frac{1}{\sqrt{2}}(|Q \uparrow \bar{q} \downarrow\rangle + |Q \uparrow \bar{q} \downarrow\rangle).$$

The two wavefunctions can be easily seen to be related by the spin operator S_Q^z , namely,

$$S_Q^z |P^*, \lambda = 0\rangle = \frac{1}{2} |P\rangle,$$

which is just a statement of the spin symmetry.

There is a very simple application of these ideas to semileptonic weak decays of heavy mesons. Suppose we have a bottom meson P_i or P_i^* with velocity v decaying into a charm meson P_j or P_j^* with velocity v' and leptons. The decay can involve the vector current V_μ^{ij} or the axial current A_μ^{ij} , and we need form factors for the vertices. Using Lorentz covariance we may write

$$\begin{aligned} \langle D(v') | V_\mu | B(v) \rangle &= h_+(v + v')_\mu + h_-(v - v')_\mu; \\ \langle D^*(v', \epsilon') | V_\mu | B(v) \rangle &= i h_V \epsilon_{\mu\nu\alpha\beta} \epsilon'^{\nu\alpha} v'^\beta; \\ \langle D^*(v', \epsilon') | A_\mu | B(v) \rangle &= h_{A1} \epsilon'_\mu - h_{A2} (\epsilon'^* \cdot v) v_\mu - h_{A3} (\epsilon'^* \cdot v) v'_\mu; \end{aligned}$$

and so on. There are a total of 17 form factors h_N needed to describe this process. However, we can use the spin and flavor symmetries to reduce this number.

First, flavor symmetry tells us that the form factors h_N are proportional to $\sqrt{m_{P_j} m_{P_j^*}}$. Then from the spin symmetry, one can relate $\langle D^* | V_\mu | B \rangle$ and $\langle D | V_\mu | B \rangle$, for example, using the relation between the $J = 0$ and $J = 1$ states obtained by applying the spin operator as above:

$$\langle D^* | V_1 | B \rangle = 2 \langle D | S_c^z V_1 | B \rangle = 2 \langle D | [S_c^z, V_1] | B \rangle = i \langle D | V_2 | B \rangle,$$

and so forth. By using these symmetries, then, we can relate all 17 form factors such that only one form factor needs to be computed. This is the Isgur-Wise function $\xi(v \cdot v')$. Now $\xi(v \cdot v')$ involves a lot of complicated dynamics about which we know nothing generally. But if $v = v'$, then one can boost to a system where everything is at rest, and so the static approximation is good, the wave functions overlap completely, and thus we know $\xi(1) = 1$. So the application of heavy quark symmetry tells us two things: first, that the number of form factors necessary to compute is reduced from 17 to 1, and furthermore that this one form factor is normalized to one at a certain kinematical condition.

Now, it is more interesting to look at heavy baryons — that is, one heavy quark plus two light quarks (Qq_1q_2) — because the two light quarks can form a diquark system. This subsystem can be in the flavor $SU(3)$ representation $\mathbf{6}$ or $\bar{\mathbf{3}}$ since the heavy quarks are $SU(3)$ flavor singlets — that is, $\mathbf{3} \times \mathbf{3} = \mathbf{6} + \bar{\mathbf{3}}$. $\mathbf{6}$ is symmetric under interchange of q_1 and q_2 , $\bar{\mathbf{3}}$ is antisymmetric; and so from Fermi statistics, we know that the former is spin 1^+ and the latter is spin 0^+ . The parity is even because the parity of the heavy baryons is always even. For the $\mathbf{6}$ representation, we write $(q_1q_2)_{1^+} = \phi_\mu$, which is an axial vector 1^+ object. From here we can combine this with the heavy quark to form spin $\frac{1}{2}$ and $\frac{3}{2}$ baryons, which we denote B_6 and B_6^* , respectively. For the $\bar{\mathbf{3}}$, we write $(q_1q_2)_{0^+} = \phi$, which is a 0^+ Lorentz scalar, and we can only form spin $\frac{1}{2}$ baryons $B_{\bar{3}}$ when combining ϕ with a heavy quark. Now one may go through a similar although somewhat more complicated argument as above for the heavy mesons. The basic idea is the same, and it turns out that in semileptonic decays of heavy baryons there will be more than 30 form factors needed (including vector and axial vector form factors), but this number can be reduced to three universal Isgur-Wise functions¹⁰. Furthermore, here too we know something about these functions in the limit when $v = v'$.

This has been a very brief introduction to heavy quark symmetry, so let us take the opportunity here to summarize why the idea of heavy quark symmetry is so interesting. The main point is that one can develop a precise formulation which allows one to deal with the quantum field theory for a light and heavy quark system. So one can make very precise statements in the limit $m_Q \rightarrow \infty$, such as $\xi(v \cdot v' = 1) = 1$ and relations among the form factors, and then the corrections can be formulated systematically. This provides a new and very intuitive language for describing heavy mesons and baryons — a language which is very useful for understanding physics and for communicating with experimentalists. Furthermore, when one combines the heavy quark symmetry with the chiral symmetry of the light quarks, one can determine the low energy dynamics of ground states of heavy mesons and baryons interacting with Goldstone bosons and also photons. This will be shown next.

3. Synthesis with Chiral Symmetry

We will not dwell on the details of chiral symmetry, since they will be discussed in detail elsewhere in these proceedings. Here we would like only to remind the reader of one relation which will be needed later: namely, that one can use the partial conservation of axial vector current, a consequence of approximate chiral symmetry, and a smoothness assumption, in order to relate the strong coupling constant $g_{\pi NN}$ and the beta decay axial vector coupling constant g_A . This is the Goldberger-Treiman relation

$$g_{\pi NN} = \frac{m_N}{f_\pi} g_A.$$

The axial vector coupling constant is defined as the matrix element of the axial vector current between neutron and proton as

$$g_A = \langle p \uparrow | A_3^{1+i2} | n \uparrow \rangle.$$

Now one can get a prediction for this matrix element from the constituent quark model. This is evidently then not part of the precise statements made above. In the quark model, g_A is computed as a sum of the elementary contributions from the quark transitions $ddu \rightarrow ddu$:

$$g_A = \langle p \uparrow | u^\dagger \sigma_3 d | n \uparrow \rangle.$$

A naive calculation then finds $g_A = \frac{5}{3}$. We know from experiment, however, that $g_A = 1.25$. Georgi, Peris, and Weinberg^{11,12}, among others, attempted to explain the discrepancy by appealing to a renormalization effect of the coupling g_A^{ud} , because the ud vertex can in principle be renormalized by QCD. One does not know how big such an effect is, but in any case Georgi and Manohar¹¹ just dictate that in order to get the right g_A we must assign a quark-quark coupling $g_A^{ud} = 0.75$. So this value will be used later in the applications, but it will be made clear where precise statements are being made and where the quark model is being used to get the numbers.

Now, a heavy meson is $Q\bar{q}$, and a heavy baryon is Qq_1q_2 . We know that for heavy quarks we have the heavy quark symmetry and for light quarks we have chiral symmetry. For these heavy-light systems, then, one should be able to extract a lot of physics from just these two symmetries. The reason this synthesis is useful is because as a consequence of the heavy quark symmetry, the mass difference between different spin multiplets is very small. For example, $M_{D^*} - M_D \sim 145$ MeV, $M_{B^*} - M_B \sim 47$ MeV, and $M_{\Sigma_c} - M_{\Lambda_c} \sim 170$ MeV (where the Σ_c is $c(ud)_{S=1}$ and the Λ_c is $c(ud)_{S=0}$). So the pions and photons are very soft in the decays $D^* \rightarrow D\pi, D\gamma$, $B^* \rightarrow B\gamma$, and $\Sigma_c \rightarrow \Lambda_c\pi, \Lambda_c\gamma$; and low energy theorems can be used to get these matrix elements. The combination of the heavy quark and chiral symmetries thus has very useful applications in these decays.

3.1. Strong Interactions. The pion, the Goldstone boson, because of its quantum numbers, is only coupled to the light quarks. It does not couple to the heavy quark spin, and therefore the matrix elements for $D^* \rightarrow D\pi$, and $D^* \rightarrow D^*\pi$ for example, are related. This is just the spin symmetry of the system. Flavor and chiral symmetries then determine the remaining structure of the matrix elements for soft pion emission for the processes $D^*(v, \epsilon) \rightarrow D(v')\pi(q)$ and $D^*(v, \epsilon) \rightarrow D^*(v', \epsilon')\pi(q)$:

$$\frac{f}{f_\pi} \sqrt{M_D M_{D^*}} \epsilon \cdot q \quad \text{and} \quad \frac{f}{f_\pi} M_{D^*} i \epsilon_{\mu\nu\alpha\beta} q^\mu \epsilon'^\nu v^\alpha \epsilon^\beta,$$

respectively. They have the same overall coefficient, the coupling f , to be determined somehow. There is thus only one independent coupling constant for all ground state heavy quarks; and this is a reflection of the fact all these couplings can be thought of as due to the coupling of the elementary light quark to the pion (this is the chiral quark model of Georgi and Manohar¹¹).

Now above we wrote only the coupling f of the heavy meson to a single pion, but of course by chiral symmetry the couplings to any odd number of pions will be related to f . Similarly, there will be another class of couplings involving an even number of pions, but those couplings will not have any new unknown parameters: the only parameter that will come in is f_π , according to the usual chiral symmetry arguments. So there is only one unknown coupling constant needed to describe the coupling of heavy mesons to soft pions.

How about heavy baryons Qq_1q_2 ? It has already been mentioned that the baryon case is more interesting because the two light quarks can form a 1^+ object ϕ_μ or a 0^+ object ϕ , where ϕ_μ leads to spin $\frac{1}{2}$ (B_6) or $\frac{3}{2}$ (B_6^*) baryons and ϕ leads to only spin $\frac{1}{2}$ (B_3) baryons when coupled to a heavy quark Q . Now one can again use the idea that in terms of quantum numbers the pion will couple only to the light quarks. There will be one coupling g_1 for the vertex $\phi_\mu \rightarrow \phi'_\mu + \pi$, and it will show up in the decays

$$B_6 \rightarrow B'_6 + \pi; \quad B_6^* \rightarrow B_6 + \pi; \quad B_6 \rightarrow B_6^* + \pi;$$

and there will be another coupling g_2 for the vertex $\phi_\mu \rightarrow \phi + \pi$, which will show up in the decays

$$B_6 \rightarrow B_3 + \pi; \quad B_6^* \rightarrow B_3 + \pi.$$

A consequence of parity conservation is that the spin zero diquark coupling to the pion $\phi \rightarrow \phi' + \pi$ is not allowed since the ϕ is 0^+ and the π is 0^- . So we have a prediction that the coupling constant describing the transition $B_3 \rightarrow B'_3 + \pi$ will vanish at this order.

3.2. Electromagnetic Interactions. One can use similar ideas to describe the electromagnetic interactions. There will be two types of couplings, the minimal (charge) couplings and M1 transitions. Only the M1 transitions will be discussed here. The photon can come from two sources: if it comes from the light quarks, again the spin symmetry tells us that the coupling is independent of the heavy quark spin; if the photon comes from the heavy quark, then the coupling constant can actually be computed exactly in the soft photon limit, because the light quark cloud is not disturbed. The coupling constant is then related to the Dirac magnetic moment $\mu_Q = \frac{eQ}{2m_Q}$.

So by similar arguments as for the strong interactions, only one coupling constant d is needed to describe heavy meson transitions $P^* \rightarrow P + \gamma$ and $P^* \rightarrow P^* + \gamma$. And then for the baryons, we can use the same arguments to conclude that there will be one coupling constant a_1 for $\phi_\mu \rightarrow \phi'_\mu + \gamma$, which will show up in the decays

$$B_6 \rightarrow B'_6 + \gamma; \quad B_6^* \rightarrow B_6 + \gamma; \quad B_6 \rightarrow B_6^* + \gamma;$$

and there will be another coupling a_2 for the vertex $\phi_\mu \rightarrow \phi + \gamma$, which will show up in the decays

$$B_6 \rightarrow B_3 + \gamma; \quad B_6^* \rightarrow B_3 + \gamma.$$

Once again by parity arguments, there is no transition $\phi \rightarrow \phi' + \gamma$ since the γ is 1^- , and thus there will be no M1 coupling $B_3 \rightarrow B'_3 + \gamma$.

So the physics is pretty simple in these pictures, and by using the spin symmetry one can substantially reduce the number of independent parameters. For the strong interactions involving the heavy mesons there is a reduction from 2 to 1 and for baryons from 6 to 2 couplings. For the electromagnetic M1 transitions there is the same reduction in the number of couplings. Up to now all these have been general statements, independent of any model, just consequences

of symmetries of QCD. Now we shall describe how these coupling constants can be computed in the constituent quark model, and this will be in the same spirit as in the calculation of g_A described above. Here, of course, we arrive at the subject of this conference, namely, how to determine whether these quark model calculations are valid or not.

Thus, if one uses the same technique as for computing the axial coupling g_A for beta decay in the nucleon, one finds that f — the coupling of pions to heavy mesons — is related to g_A^{ud} of the light quarks, namely, $f = -2g_A^{ud}$. Then for the two heavy baryon coupling constants, one can use the analogue of the Goldberger-Treiman relation to express g_1 and g_2 in terms of this single coupling constant: $g_1 = \frac{1}{3}g_A^{ud}$ and $g_2 = -\sqrt{\frac{2}{3}}g_A^{ud}$. Now as mentioned before, naively we have $g_A^{ud} = 1$, but because of the discrepancy with experiment Georgi and Manohar suggest that we should assign $g_A^{ud} = 0.75$. We shall just adopt this assumption here. The electromagnetic coupling constants d , a_1 , and a_2 can also be computed in the constituent quark model, and they will be related to the magnetic moments of the light quarks. So using the heavy quark symmetry we have reduced the number of coupling constants to a small number, and we can then compute these in the constituent quark model.

With this combination of first principle calculations and the quark model results for the coupling constants, then, one can get some numbers. Taking the quark constituent masses from the particle data book¹³: $m_u = 338$ MeV, $m_d = 322$ MeV, $m_s = 510$ MeV, $m_c = 1.6$ GeV, and $m_b = 5$ GeV, and using $g_A^{ud} = 0.75$, one can make predictions for these decays. Unfortunately, the experiments cannot measure the absolute widths because they are too small. But one can compute and compare the branching ratios.

These results for strong and electromagnetic interactions with heavy mesons are given in Table ?? . CLEO¹⁴ re-measured the branching ratios just recently for the charm mesons, and they came up with the numbers given in the fourth column. Now if you look at the 1992 Particle Data book¹⁵ (fifth column), you will find that there is a discrepancy with the recent CLEO measurements. The theory agrees very well with the CLEO measurements. Most interesting is the decay $D^{*0} \rightarrow D^0\gamma$, whose branching ratio is not small; it agrees very well with the CLEO measurement. So there is an indication that something is working here. It is important to test experimentally the validity of our predictions for the absolute widths.

One can also compute widths and branching ratios for D_s^{*+} , B_u^{*+} and B_d^{*0} decays, but there are no experimental data yet with which to compare these numbers. Note that in Table ?? the charm meson mass is kept finite. In fact, if $m_c \rightarrow \infty$, one finds

$$D^{*+} \rightarrow D^+\gamma \quad 6\text{keV} \quad (1)$$

$$D^{*0} \rightarrow D^0\gamma \quad 23\text{keV} \quad (2)$$

$$D_s^{*+} \rightarrow D_s^+\gamma \quad 2.4\text{keV} \quad (3)$$

as compared with the decay widths for finite m_c above. The difference is substantial, because $m_c = 1.6\text{GeV}$ is not very heavy, and also since the charm quark has charge $\frac{2}{3}$. So one can see that it is quite important to keep the charm quark mass finite. Of course, this is not quite

Table 1: Heavy meson decay widths and branching ratios

Decay Mode	Γ (keV)	BR(Th) (%)	BR(CLEO)	BR(PDG)
$D^{*+} \rightarrow D^0\pi^+$	102	68	$68.1 \pm 1.0 \pm 1.3$	55 ± 4
$D^{*+} \rightarrow D^+\pi^0$	46	31	$30.8 \pm 0.4 \pm 0.8$	27.2 ± 2.5
$D^{*+} \rightarrow D^+\gamma$	2	1.3	$1.1 \pm 1.4 \pm 1.6$	18 ± 4
$D^{*+} \rightarrow \text{all}$	150	100		
$D^{*0} \rightarrow D^0\pi^0$	70	66.7	$63.6 \pm 2.3 \pm 3.3$	55 ± 6
$D^{*0} \rightarrow D^0\gamma$	34	33.3	$36.4 \pm 2.3 \pm 3.3$	45 ± 6
$D^{*0} \rightarrow \text{all}$	104	100		
$D_s^{*+} \rightarrow D_s^+\gamma$	0.3	100		
$B_u^{*+} \rightarrow B_u^+\gamma$	0.84	100		
$B_d^{*0} \rightarrow B_d^0\gamma$	0.28	100		

Table 2: Heavy baryon decay widths and branching ratios

Decay Mode	Γ (keV)	BR(Th) (%)
$\Sigma_c^+ \rightarrow \Lambda_c^+\pi^0$	2430	96.2
$\Sigma_c^+ \rightarrow \Lambda_c^+\gamma$	93	3.8
$\Sigma_c^+ \rightarrow \text{all}$	2540	
$\Xi_c'^+ \rightarrow \Xi_c^+\gamma$	16	100
$\Xi_c'^0 \rightarrow \Xi_c^0\gamma$	0.3	100

consistent, since we have included only one source of $1/m_Q$ corrections. Recall that the quark masses are used for determining the magnetic moments for the electromagnetic M1 transitions.

Calculations have also been performed for the heavy baryons, as shown in Table ???. Unfortunately, there are no experimental data yet with which we can compare these numbers. In the table, the baryons not yet introduced in the text are $\Xi_c'^+ = c(us)_{S=1}$, $\Xi_c^+ = c(us)_{S=0}$, $\Xi_c'^0 = c(ds)_{S=1}$, and $\Xi_c^0 = c(ds)_{S=0}$.

3.3. Semileptonic Decays. The previous predictions actually have little to do with heavy quark symmetry. They follow just from combining the quark model with chiral symmetry. Now we shall show one example where, at least theoretically, one has to use both chiral and heavy quark symmetries. This is the semileptonic decay of a heavy meson including a soft pion. Some motivation for doing this analysis comes from the existing experimental data, which give the branching ratios for $B^0 \rightarrow D^-l^+\nu$ to be $(1.8 \pm 0.5)\%$ and for $B^0 \rightarrow D^{*-}l^+\nu$ to be $(4.9 \pm 0.8)\%$, and similar results for B^+ decay, whereas the branching ratio for the inclusive decay $B \rightarrow \text{hadrons } e^\pm\nu_e$ is $(10.7 \pm 0.5)\%$. Since the first two do not add up to 10.7%, this indicates that there will be other states besides D and D^* important for this decay. We want to find out if soft pions could be important.

Theoretically, the problem is quite interesting. As an example, consider the decay $B(v) \rightarrow$

Table 3: Semileptonic decay integrated rates and branching ratios

Frame	cutoff (MeV)	Rate (MeV)	BR	L_0 (%)	L (%)	R (%)
B	100	3.20×10^{-15}	0.63×10^{-5}	26	57	17
B	200	1.84×10^{-14}	0.36×10^{-4}	30	52	18
B and D^*	100	9.62×10^{-16}	0.19×10^{-5}	27	54	19
B and D^*	200	1.04×10^{-14}	0.20×10^{-4}	33	48	19

$D^*(v')\pi(q)l\nu$. This decay may proceed in three ways: (i) $B \rightarrow B^*(\sim v)\pi \rightarrow D^*\pi l\nu$, (ii) $B \rightarrow D^*(\sim v')l\nu \rightarrow D^*\pi l\nu$, or (iii) $B \rightarrow D(\sim v')l\nu \rightarrow D^*\pi l\nu$. As discussed before, we can use the heavy quark symmetry to relate the different weak decays, so that there is just one unknown form factor $\xi(v \cdot v')$. And for soft pions the different strong interaction vertices all depend on just the one coupling f because of chiral symmetry, as described before. Therefore, even though these amplitudes are in principle independent, we can use spin, flavor, and chiral symmetries to determine their relative phases and magnitudes for soft pions, and thus there is only an overall normalization which we do not know. One then just needs a model for the form factor $\xi(v \cdot v')$, and so Burdman's fit¹⁶ was used for the numerical calculation.

If the invariant mass of the $D\pi$ system is less than three widths from the D^* invariant mass — that is, $|M(D\pi) - M(D^*)| < 3\Gamma_{D^*}$ — then we call it a part of the D^* resonance. Everything else is non-resonance, and it turns out that the non-resonant part is very small compared with the resonant part. In Table ??, we list the integrated decay rates and branching ratios for $\bar{B}^0 \rightarrow D^{*+}\pi^0 e^- \bar{\nu}_e$. We also give contributions from different polarizations of the D^* (longitudinal, left-handed, and right-handed polarization are labelled L_0 , L , and R , respectively). The relative magnitudes of these can be understood from the fact that the underlying weak interaction is $V - A$. It turns out that the branching ratio is not very big. For the D it is $\sim(0.5-1)\%$, and for the D^* only $\sim 10^{-4} - 10^{-5}$. So the difference between the inclusive and exclusive B decay cannot be explained by soft pion contributions. It must be due to something else, perhaps an excited state.

4. Concluding Remarks

This has been a very short summary of work contained in Ref. [1-4] that has been done on heavy quark systems involving strong and electromagnetic interactions and semileptonic weak decays with soft pion emission. But besides the particular calculations which have been described here, this work has also included some other computations which we would like to at least mention. This includes work on $1/m_Q$ corrections¹⁷, which however could not produce reliable numerical estimates because the coefficients involve matrix elements of light quark operators which will require the success of this conference to be worked out. There is also work on chiral symmetry breaking¹⁷ due to the fact that the light quark masses are non-zero. This has been more successful: it is possible to compute non-analytic contributions (depending on the

square root and logarithm of the mass, and so on). Some heavy flavor conserving weak decays — for example, $\Xi_c \rightarrow \Lambda_c \pi$ — have also been studied. And recently work on weak radiative decays¹⁷ — flavor changing electromagnetic decays — has been completed. This calculation was motivated by the observation at CLEO of the so-called penguin decay $B \rightarrow K^* \gamma$. The question is then whether the decay $B \rightarrow D^* \gamma$ is as small or bigger than this. This calculation can be done in a similar framework to those described here, and it turns out that the branching ratio for $B \rightarrow B^* \gamma$ is only 10^{-6} , and so the dominant decay is $B \rightarrow K^* \gamma$.

To summarize, then, it has been shown that by combining the heavy quark and chiral symmetries, one can reduce the number of unknown parameters for low energy dynamics of heavy mesons and baryons with Goldstone bosons and photons. Furthermore, if one applies the constituent quark model, then everything is determined. Now there is a curious fact which should be mentioned: namely, suppose one does not assume any heavy quark symmetry and just goes ahead and computes all coupling constants and so forth using the quark model, then one finds that the coupling constants all satisfy the heavy quark and spin symmetry relations. So the constituent quark model is consistent with the heavy quark symmetry. We leave it to the reader to judge the relative profundity or triviality of this result.

Finally, we remark that just as the hydrogen atom is a simple system whose study allows us to begin to understand QED, we have long searched for a similar relatively simple starting point for understanding QCD. When charmonium was discovered, it was said by some that the $Q\bar{Q}$ system, the heavy quark quarkonium, could be the hydrogen atom of QCD. But it seems that the $Q\bar{q}$ system provides a better analogy with the hydrogen atom because as $m_Q \rightarrow \infty$ the system is essentially reduced by one degree of freedom, and one only has to deal with the light quark. So in the sense of providing a simple starting point for understanding the difficult aspects of the theory, then, we may say that the heavy meson is the hydrogen atom of QCD.

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