Combining Vector— and Heavy Quark Symmetry

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Abstract
Assuming that the chiral symmetry of the light degrees of freedom is realized as the vector symmetry proposed by Georgi, we write down a chiral Lagrangian for heavy mesons which incorporates both heavy quark and vector symmetry. Some of the phenomenological implications of this idea are considered.

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1 Introduction

The combination of heavy quark and chiral symmetries has turned out to be a powerful tool to obtain model independent predictions for various processes. In particular, the decays of a heavy meson into light pseudoscalars has been considered in this framework and some interesting predictions for strong processes as well as for electromagnetic transitions have been obtained [1].

This approach to heavy to light decays has been extended to light vector mesons by using the approach of hidden symmetry [2]. In this approach the light vector mesons are introduced as gauge bosons of a hidden symmetry which is broken spontaneously to give masses to the light vector mesons [3, 4, 5]. In this way one may implement chiral and Lorentz invariance in an elegant way; it has been shown [6] that writing down all couplings allowed by these two symmetries yields the hidden symmetry Lagrangian. Including heavy mesons along the lines proposed in [2], the resulting heavy-meson chiral Lagrangian has, however, a large number of unknown coupling parameters, even to lowest order in the chiral expansion.

It has been pointed out by Georgi [7] that the hidden symmetry method has an interesting limit in which additional symmetries occur. In this limit, the so called vector limit, the chiral symmetry is in fact realized in an unbroken way, the vector meson octett becomes massless and the scalar fields corresponding to the longitudinal components of the vector mesons become the chiral partners of the pions. The phenomenology of this model has been worked out in some detail [8].

In fact, there are arguments that this realization of chiral symmetry may be obtained from the dynamics of QCD. Of course, it is impossible to trace what happens to chiral symmetry once the scales become that low that one enters the nonperturbative region of QCD, but at least the enlarged symmetry of the vector limit, the so called vector symmetry, is compatible with the $N_c \to \infty$ limit of QCD, where $N_c$ is the number of colors [7]. Another way such a limit may be realized is the scenario of “mended symmetries” [9], which is envisaged to hold sufficiently close to the chiral symmetry restoration point. Further arguments in favour the vector limit have been given recently [10].

We shall adopt the point of view that the dynamics of QCD is indeed such that reality is sufficiently close to the vector limit and that it may be used as a starting point. In that case the enlarged symmetry in the vector limit yields a reduction of the number of independent coupling constants, which may appear in the chiral Lagrangian at low energies. In fact, in the vector limit the dynamics of the light degrees of freedom (in this case the pions and the longitudinal components of the light vector mesons) is determined solely by $f_\pi$, the pion decay constant.

However, in order to have a good description of data one has to take into account the breaking of vector symmetry. For instance, the relation obtained for the mass of the $\rho$ meson is a factor $\sqrt{2}$ off from the usual KSRF relation,
yielding a $\rho$ mass which is a factor of $\sqrt{2}$ too large. This factor indicates the size of corrections one has to expect when using the vector limit. Other examples are the application of the vector limit to hadronic decays of $D$ mesons \[12\] and to intrinsic parity violating decays \[11\], where an agreement of similar quality has been observed.

The coupling of the transverse components of the $\rho$ meson breaks the enlarged symmetry, and the above prediction for the $\rho$ mass is obtained, if this coupling is the only source of symmetry breaking. The situation may be improved by including systematically other sources of vector symmetry breaking. This has been studied in detail for the hadronic decays of the $\tau$ lepton \[13\] and for the hadronic $D$ meson decays \[14, 15\].

The present note extends the combination of heavy quark and chiral symmetry to the case where the symmetry of the light degrees of freedom is the vector symmetry. Due to the enlarged symmetry relations between the coupling constants are obtained, which reduce the number of independent coupling constant to only one in the vector limit. However, corrections to the vector limit may be sizable and the predictions of the vector limit may obtain corrections factors similar to the one for the $\rho$ mass, namely a factor of $\sqrt{2}$. In the present note, we shall not consider any symmetry breaking terms but shall explore only the consequences of the symmetry limit.

The paper is organized as follows. In the next section we shall recall the concept of hidden and vector symmetry of the light degrees of freedom. In section 3 the coupling of the heavy mesons to the light degrees of freedom is considered and the relations between coupling constants in the vector limit are given. Finally, we shall discuss some phenomenological implications and conclude.

2 Hidden Symmetry and the Vector Limit

The chirally symmetric Lagrangian for the light pseudoscalars is conveniently obtained from a nonlinear representation of the chiral $SU(3)_L \otimes SU(3)_R$ symmetry \[10\]. It is defined in terms of a matrix field $\xi$

$$\xi = \exp \left( \frac{i}{f} \pi \right) \quad \pi = \pi^a T^a \quad \text{Tr} \left( T^a T^b \right) = \frac{1}{2} \delta^{ab} \quad (1)$$

where $T^a$ denote the generators of $SU(3)$ in the fundamental representation. Under chiral $SU(3)_L \otimes SU(3)_R$ the field $\xi$ transforms as

$$\xi \rightarrow L \xi U^\dagger = U \xi R^\dagger \quad L \in SU(3)_L \quad R \in SU(3)_R \quad (2)$$

The Lagrangian to lowest order (i.e. containing the lowest number of derivatives of the fields) for the light pseudoscalar octett is as usual

$$\mathcal{L}_0 = \frac{f^2}{4} \text{Tr} \left\{ (\partial_\mu \Sigma)^\dagger (\partial^\mu \Sigma) \right\} \quad \text{with} \quad \Sigma = \xi^2, \quad (3)$$
where $f = 93$ MeV is the pion decay constant.

To include also the light vector mesons the so called hidden symmetry approach has been proposed in [3]. This method amounts to insert an additional "hidden" $SU(3)_H$ by defining matrix fields $\Sigma_L$ and $\Sigma_R$ which transform as

\[
\Sigma_L \rightarrow L \Sigma_L h^\dagger \quad L \in SU(3)_L \quad h \in SU(3)_H \\
\Sigma_R \rightarrow R \Sigma_R h^\dagger \quad R \in SU(3)_R \quad h \in SU(3)_H
\]

where $SU(3)_H$ is the so called hidden symmetry. The light vector mesons are introduced as gauge fields of the local hidden $SU(3)_H$; since this group is spontaneously broken, the vector mesons will acquire a mass. This becomes evident by expressing the two fields $\Sigma_{L/R}$ in terms of the field $\xi$

\[
\Sigma_L = \xi \exp \left( \frac{i}{f} s \right) \quad s = s^a T^a \\
\Sigma_R = \xi^\dagger \exp \left( \frac{i}{f} s \right)
\]

The additional scalar fields $s$ are the Goldstone bosons of the broken hidden symmetry and will become the longitudinal components of the vector mesons.

The Lagrangian for the light pseudoscalars and the light vector mesons is conveniently formulated in terms of the currents

\[
A_\mu = \frac{i}{2} \left\{ \Sigma_R^\dagger \partial_\mu \Sigma_R - \Sigma_L^\dagger \partial_\mu \Sigma_L \right\} \\
\mathcal{L}_{01} = -f^2 \text{Tr} \left\{ A_\mu A^\mu \right\} - \frac{1}{2} \text{Tr} \left\{ F_\mu\nu F^{\mu\nu} \right\}
\]

and reads

\[
\rho_\mu \rightarrow h \rho_\mu h^\dagger + \frac{i}{g_V} h \partial_\mu h^\dagger.
\]

$F_{\mu\nu}$ is the usual field strength tensor

\[
F_{\mu\nu} = \partial_\mu \rho_\nu - \partial_\nu \rho_\mu + g_V [\rho_\mu, \rho_\nu].
\]

This Lagrangian contains (aside from the pion decay constant $f$) the coupling constant $g_V$ and the parameter $a$. These two parameters are fixed by the value
of the coupling strength of the $\rho$ to the pions $g_{\rho\pi\pi}$ and the mass of the $\rho$ meson. One obtains for the two parameters in terms of the input

$$g_{\rho\pi\pi} = \frac{a}{2}g_V \quad m_\rho^2 = \frac{4}{a^2}f^2g_{\rho\pi\pi}^2$$  \hspace{1cm} (13)

The choice of $a = 2$ yields the KSRF relation $m_\rho^2 = 2f^2g_{\rho\pi\pi}$ which works surprisingly well, while the gauge coupling is fixed for $a = 2$ to be $g_V = g_{\rho\pi\pi} \sim 6.0$.

It has been pointed out by Georgi that for the choice $a = 1$ and in the limit $g_V \to 0$, the symmetry of the system becomes enlarged to the so called vector symmetry. The Lagrangian becomes in this limit

$$\mathcal{L}_{VL} = -f^2\text{Tr} \left\{ A_\mu A^\mu + V_\mu V^\mu \right\}$$

$$= \frac{f^2}{2}\text{Tr} \left\{ (\partial_\mu \Sigma_L)(\partial^\mu \Sigma_L)^\dagger + (\partial_\mu \Sigma_R)(\partial^\mu \Sigma_R)^\dagger \right\}$$  \hspace{1cm} (15)

which has a global

$$SU(3)_L \otimes SU(3)_H \otimes SU(3)_R \otimes SU(3)_H$$

$$SU(3)_{L+H_L} \otimes SU(3)_{R+H_R}$$

symmetry which acts in the following way on the fields

$$\Sigma_L \to L\Sigma_L h_L^\dagger \quad L \in SU(3)_L \quad h_L \in SU(3)_H$$

$$\Sigma_R \to R\Sigma_R h_R^\dagger \quad R \in SU(3)_R \quad h_R \in SU(3)_H.$$  \hspace{1cm} (16)

This means in particular, that for the light degrees of freedom there is an unbroken chiral symmetry with parity doublets consisting of the pions and the longitudinal components of the light vector mesons. This enlarged symmetry is broken by the coupling of the transverse components of the vector mesons which act as gauge bosons of the diagonal group $H_L + H_R$. The resulting Lagrangian is the same as (10) with the choice $a = 1$. However, if this coupling is the only source of vector symmetry breaking one obtains a mass for the vector mesons, which is a factor 1.4 too large. This indicates that there have to be additional sources of vector symmetry breaking, which will contribute substantially; still the limiting case may be of some use to obtain an idea of the size of the coupling constants considered below.

### 3 Coupling of Heavy Mesons

We are now ready to consider the coupling of the heavy mesons to the light pseudoscalar and vector mesons. We shall start by writing down the Lagrangian for light pseudoscalar and vector mesons for the general case, including the coupling of these light states to the heavy ground state spin symmetry doublet, and to an
excited spin symmetry doublet, consisting of a $0^+$ and a $1^+$ state. We combine the two heavy spin symmetry doublets into two fields:

$$H(v) = \frac{\sqrt{m_H}}{2}(1 + \not{\phi}) \left( -\gamma_5 \not{P} + \not{P} \gamma_\mu \right) \not{P}^i \cdot v = 0$$ \hspace{1cm} (18)

for the $(0^-, 1^-)$ spin symmetry doublet

$$K(v) = \frac{\sqrt{m_H}}{2}(1 + \not{\phi}) \left( -\not{S} + \not{S} \gamma_\mu \gamma_5 \right) \not{S}^i \cdot v = 0$$ \hspace{1cm} (19)

which transform as

$$H(v) \rightarrow H(v)h^\dagger \hspace{1cm} K(v) \rightarrow K(v)h^\dagger$$ \hspace{1cm} (20)

under $SU(3)_L \otimes SU(3)_R \otimes SU(3)_H$.

In leading order the most general Lagrangian consistent with hidden symmetry and spin symmetry is then given by

$$\mathcal{L}_2 = \text{Tr} \left[ \bar{H}^\alpha (iv \cdot D^\beta \alpha) H^\beta \right] - \text{Tr} \left[ \bar{K}^\alpha (iv \cdot D^\beta \alpha) K^\beta \right] + g_H \text{Tr} \left[ \bar{H}^\alpha H^\beta \gamma_5 A^\beta \right] + g_K \text{Tr} \left[ \bar{K}^\alpha K^\beta \gamma_5 A^\beta \right] + g'_H \text{Tr} \left[ \bar{H}^\alpha H^\beta (\not{\psi}^\beta - g_\psi^\beta \not{\psi'}) \right] + g'_K \text{Tr} \left[ \bar{K}^\alpha K^\beta (\not{\psi}^\beta - g_\psi^\beta \not{\psi'}) \right] + \lambda_V \text{Tr} \left[ \bar{H}^\alpha K^\beta (\not{\psi}^\beta - g_\psi^\beta \not{\psi'}) \right] + \lambda_A \text{Tr} \left[ \bar{H}^\alpha K^\beta A^\beta \right] + h.c.$$ \hspace{1cm} (21)

where we have written the $SU(3)_H$ indices explicitly,

$$D^\mu_{\alpha\beta} = \delta^\alpha_\beta \partial_\mu + ig_\psi^\beta \rho^{\alpha\beta}$$

is a covariant derivative of $SU(3)_H$ and the trace is to be taken with respect to the dirac matrix structure. This is the Lagrangian as it was written down in [2]. It is expressed in terms of six independent couplings $g_H, g_K, g'_H, g'_K, \lambda_V, \lambda_A$, which need to be fixed from some experimental input.

In the vector limit the chiral $SU(3)_L \otimes SU(3)_R$ symmetry is unbroken and hence chiral partners of the heavy particles have to show up. Obviously one may not use a similar construction as for the light states, since it makes no sense to consider the massless limit of the heavy states. The way we shall proceed here is to introduce the chiral partners directly, i.e. we shall combine the $0^-$ and the $0^+$ state into a parity doublet, while the $1^-$ and the $1^+$ state form another parity doublet. From the heavy quark limit it follows that the splitting between spin symmetry partners is of order $\Lambda_{QCD}$, while from heavy quark symmetry alone one expects that the splitting between the $0^-$ and the $0^+$ state should be a constant, which is independent of the heavy quark mass and related to vector symmetry breaking. Thus in the combined heavy quark and vector limit, all four states should be degenerate.
It is convenient to define left and right handed heavy fields as

\[ H_L = (H + K)^\frac{1 - \gamma_5}{2} \]
\[ H_R = (H + K)^\frac{1 + \gamma_5}{2} \]  

(22) 

(23)

which transform under \( SU(3)_R \otimes SU(3)_H \otimes SU(3)_L \otimes SU(3)_H \) as

\[ H_L \rightarrow H_L h_L^\dagger \]
\[ H_R \rightarrow H_R h_R^\dagger \]  

(24)

The most general lagrangian which embodies the enlarged symmetry of the vector limit is then simply given by

\[ L_{VL}^2 = -\text{Tr} \left[ \bar{H}_L^\alpha (iv \cdot \partial) H_L^\alpha \right] - \text{Tr} \left[ \bar{H}_R^\alpha (iv \cdot \partial) H_R^\alpha \right] + g \text{Tr} \left[ H_L^\alpha H_R^\beta (\not{v} - \not{A}) \right] + g \text{Tr} \left[ H_R^\alpha H_R^\beta (\not{v} + \not{A}) \right] \]  

(25)

where \( g \) is given by the \( H^* H \pi \) coupling. Thus the enlarged symmetry forces all the six coupling constants appearing in (21) to become equal

\[ g = g_H = g_K = g_H' = g_K' = \lambda_V = \lambda_A \]  

(26)

The relations (26) remain true, if the vector limit is only broken by the gauge coupling of the light vector mesons. However, it is known that there are other sources of vector symmetry breaking, which are large and will affect the relation (26). Thus (26) is likely to receive large corrections, but it still may be useful for a first guess.

4 Phenomenology

In this section we shall discuss the phenomenological implications of the combined heavy quark and vector symmetry. First we consider strong decays. Off the vector limit the coupling constant \( g \) is determined from the strong decays \( D^+ \rightarrow D \pi \). The total rates for these decays are given by

\[ \Gamma(D^{*+} \rightarrow D^0 \pi^+) = \frac{g^2}{96m_D^3 \pi f^2} \left[ (m_{D^*}^2 - (m_D + m_\pi)^2)(m_{D^*}^2 - (m_D - m_\pi)^2) \right]^{3/2} \]  

(27)

\[ \Gamma(D^{*+} \rightarrow D^+ \pi^0) = \frac{g^2}{192m_D^3 \pi f^2} \left[ (m_{D^*}^2 - (m_D + m_\pi)^2)(m_{D^*}^2 - (m_D - m_\pi)^2) \right]^{3/2} \]  

(28)

Experimentally only an upper bound for these decays is known [19]

\[ \Gamma(D^{*+} \rightarrow D^0 \pi^+) < 72 \text{ keV(90\% CL)} \]  

(29)
from which one obtains $|g| < 0.63$ [1]. In the vector limit one predicts that this coupling is the same for the $D^*D\rho$ vertex, at least for the longitudinal components. However, this vertex cannot mediate a real decay due to phase space.

From the above Lagrangian one also predicts that the strong decay $D_1(2420) \to D^*(\pi)$ is governed by the same coupling constant $g$. In terms of this coupling constant we have

$$
\Gamma(D_1(2420)^0 \to D^+\pi^-) = \Gamma(D_1(2420)^0 \to D^{*+}\pi^-)
$$

$$
= \frac{g^2m_{D^*}^2}{32m_{D_1}^4\pi f^2} \left[ (m_{D_1}^2 - (m_D^* + m_{\pi})^2)(m_{D_1}^2 - (m_D^* - m_{\pi})^2) \right]^{3/2}
$$

This decays has been seen, but no branching fraction has been measured. Based on present data we have [18]

$$
\Gamma(D_1(2420) \to D^*(\pi)) < \Gamma_{tot}(D_1(2420)) = 20^{+7}_{-6} \text{ MeV}
$$

If we use the measured values for the masses, we extract a limit on $|g|$ which is stronger than the one from $D^{**} \to D^0\pi^+$, namely $|g| < 0.35$. This, however, includes already symmetry breaking effects due to the mass difference between the $D^*$ and the $D_1$. The result depends on the third power of the pion three momentum $|\vec{p}_{\text{cms}}|$ in the cms frame, which is a factor of nine larger in the $D_1$ decay than in the $D^*$ decay. The standard procedure to correct for phase space effects is to divide out a factor $2|\vec{p}_{\text{cms}}|/m_D$, which gives in the present comparison $|g| < 1.0$. However, this is somewhat arbitrary; a comparison between the two results for the decay constants needs to be more sophisticated and has to take into account symmetry breaking.

The enlarged symmetry has also consequences for weak decays. For the decay constants defined as

$$
\langle 0|\bar{q}\gamma_\mu(1-\gamma_5)Q_v|H_v(0^-)\rangle = iF_H m_H v_\mu
$$

$$
\langle 0|\bar{q}\gamma_\mu(1-\gamma_5)Q_v|K_v(0^+)\rangle = -iF_K m_K v_\mu
$$

one obtains the prediction $F_H = F_K$ in the vector limit. This prediction is somewhat counterintuitive, if one has a wave function model for the mesons in mind. The $0^+$ and $1^+$ states are both $P$-wave states in a wave function picture of the heavy-light meson, while the $0^-$ and $1^-$ states are $S$-waves. Furthermore, in such a picture the decay constants are proportional to the wave function at the origin, in other words, it should vanish for the $P$ meson states. However, in real life this means that $F_H$ is larger than $F_K$ and a difference between the two constants from the point of vector symmetry has to be attributed to symmetry breaking, which is known to be sizable.

Let us now consider semileptonic decays of heavy mesons into light ones. In the vector limit, the transverse components of the light vector particles decouple, and hence the decays into transversely polarized light vector mesons should be
Table 1: Measurements of $\Gamma_L/\Gamma_T$ in $D \to K^\ast \ell \nu$.

|        | E653 [21]    | E691 [22]    | MARK III [23] |
|--------|--------------|--------------|---------------|
| $\Gamma_L/\Gamma_T$ | $1.18 \pm 0.18 \pm 0.08$ | $1.8^{+0.6}_{-0.4} \pm 0.3$ | $0.5^{+1.0+0.1}_{-0.1-0.2}$ |

suppressed compared to the decay rate into longitudinally polarized light vectors. Furthermore, the rate for longitudinally polarized vector mesons should become equal to the rate into pseudoscalar mesons [20].

However, this cannot be true for over the whole phase space available, since the total rates for semileptonic $D$ decays do not support this picture. In table [4] we show the data for the ratio $\Gamma_L/\Gamma_T$ for the decay $D \to K^\ast e\nu$ from different experiments. In the scenario considered above we would expect $\Gamma_L/\Gamma_T \gg 1$, which is not consistent with the measurements. In addition, the ratio of decay rates

$$R = \frac{\Gamma(D \to K^\ast e\nu)}{\Gamma(D \to Ke\nu)} = \frac{\Gamma_T + \Gamma_L}{\Gamma(D \to Ke\nu)}$$

is experimentally $R = 0.51 \pm 0.18$ in the neutral $D$ decays and $R = 0.74 \pm 0.19$ in the charged $D$ decays. From the vector limit one would conclude that $R \sim 1$, with the real value being less than unity due to the small ratio $\Gamma_L/\Gamma(D \to Ke\nu)$, but even for the measured values of $\Gamma_L/\Gamma_T \sim 1.2$ the ratio $R$ is inconsistent with the above picture.

However, chiral symmetry is expected to be valid only for sufficiently soft light particles. This assumption is not valid over the whole phase space in semileptonic heavy to light decays. Neglecting the mass of the light particle, the maximal energy of the light meson is $E_{max} = m_H/2$, which becomes large in the heavy mass limit. The description based on the chiral limit is certainly valid only close to the kinematic point, where the energy of the light particle is small, $E \sim \Lambda_{QCD}$, and the problems discussed above for the total rate may be related to the inadequacy of the chiral limit in most of the phase space. In order to test this one has to compare the lepton spectra of the decays into pseudoscalar and vector mesons.

Defining the form factors for the semileptonic decays according to

$$\langle K(p) | \bar{s} \gamma_{\mu} (1 - \gamma_5) h_{\nu} | H(v) \rangle = F_+(m_H v_{\mu} + p_{\mu})$$

$$\langle K(p, \varepsilon) | \bar{s} \gamma_{\mu} (1 - \gamma_5) h_{\nu} | H(v) \rangle = m_H \varepsilon_{\mu} F_1^A + m_H (v \cdot \varepsilon) v_{\mu} F_2^A + i \varepsilon_{\mu\nu\rho\sigma} \varepsilon^\nu v^\rho p^\sigma R$$

where form factors proportional to $q_{\mu} = m_H v_{\mu} - p_{\mu}$ have been omitted. Neglecting the mass of the light meson one obtains for the differential rates

$$\frac{d\Gamma_0(D^0 \to K^+ e\nu)}{dz} = \frac{G_F^2 m_D^5}{192\pi^3} |V_{cs}|^2 z^3 |F_+|^2$$

8
where \( z = 2(\nu p)/m_H \). Note that in the longitudinal rate the dependence on \( r_* = m_H/m_{K^*} \) has to be kept, since the longitudinal rate behaves as \( 1/m_{K^*}^2 \).

Approaching the vector limit, the longitudinal rate has to have a finite limit and one obtains the following relation between the form factors \[ m_D F_1^A + (\nu \cdot p) F_2^A = 2g_V f F_+ . \] (40)

Together with the form factor \( F_+ \) which may be derived from the chiral Lagrangian \[ F_+ = \frac{F_D}{2F_K} \left( 1 + g \frac{m_D - \nu \cdot p}{\nu \cdot p + m_D} \right) \] (41)

one obtains a prediction for the shape of the spectra, which should hold for not too large energies of the light meson.

Finally, it is worth to mention that the implications of vector symmetry for nonleptonic weak decays of heavy mesons have also been considered \[ \text{[12, 14, 15]}, \] with the result that the relations between the various decay rates are within a factor 2 in agreement with data.

## 5 Conclusions

Combining heavy quark and chiral symmetry in its conventional form has lead to many interesting predictions. However, in the conventional formulation of the chiral Lagrangian for the light degrees of freedom only the pions appear, but for phenomenological applications it is desirable to have also the light vector mesons in the Lagrangian for the light degrees of freedom.

The light vector mesons may be introduced by writing down all coupling terms allowed by chiral and Lorentz invariance; this, however, leads to a proliferation of unknown coupling constants in the heavy-meson chiral Lagrangian, which need to be fixed from experimental input.

If chiral symmetry is indeed realized in an unbroken way and the vector limit may be used as a starting point, the symmetry becomes larger than in the conventional picture. Although it is still obscure, whether and how this enlarged symmetry is generated from QCD, it is still a useful tool to reduce the number of independent coupling constants once the light vector mesons are introduced.

Vector symmetry leads to a relation between the matrix elements involving light pseudoscalars and longitudinal components of light vector mesons. However, for the heavy particles, vector symmetry has to be realized in a different way, since
the vector limit corresponds to the massless limit for the light mesons. We have chosen to have an explicit representation of the unbroken symmetry for the heavy mesons, such that we have degenerate parity doublets for the heavy states. We have included the heavy ground state spin symmetry doublet for the mesons and a spin symmetry doublet of excited, positive parity mesons, which we identified with the chiral partners of the ground state spin symmetry doublet. Due to the enlarged symmetry, the Lagrangian with this set of fields still has only one unknown coupling constant, which is the same as the one appearing in the usual chiral Lagrangian and which is related to the $H^*H\pi$ vertex.

The comparison of this idea with phenomenology shows that large symmetry breaking will be present. The symmetry breaking effects have been parameterized for the light degrees of freedom $\mathbf{8}$, $\mathbf{13}$ and a similar approach may be taken for heavy light systems as well. However, a more detailed comparison in order to determine the size of symmetry breaking has to wait until better data become available.

References

[1] For a review see: M. Wise, Caltech Report CALT-68-1860 (1993), Lectures given at CCAST Symposium on Particle Physics at the Fermi scale, May 27 - Jun 4, 1993.

[2] R. Casalbuoni, A. Deandrea, N. Di Bartolomeo, R. Gatto, F. Feruglio and G. Nardulli, Phys. Lett. B292 (1992) 371.

[3] M. Bando, T. Kugo and K. Yamawaki, Phys. Rep. 164 (1988) 217.

[4] M. Harada and K. Yamawaki, Phys. Lett. B297 (1992) 151.

[5] M. Harada, T Kugo and K. Yamawaki, Phys. Rev. Lett. 71 (1993) 1299.

[6] S. Weinberg, Phys. Rev. 166 (1968) 1568.

[7] H. Georgi, Phys. Rev. Lett. 63 (1989) 1917, Nucl. Phys. B331 (1990) 311.

[8] P. Cho, Nucl. Phys. B358 (1991) 383.

[9] S. Weinberg, Phys. Rev. Lett. 65 (1990) 1177.

[10] G. Brown and M. Rho, Stony Brook preprint SUNY NTG-94-44.

[11] A. Falk and M. Luke, Nucl. Phys. B337 (1990) 49.

[12] H. Georgi and F. Uchiyama, Phys. Lett. B238 (1990) 395.

[13] T. Mannel, Phys. Lett. B244 (1990) 502.
[14] H. Georgi and F. Uchiyama, Phys. Lett. B247 (1990) 394.

[15] H. Georgi, F. Uchiyama and A. Yamada, Nucl. Phys. B382 (1992) 3.

[16] C. Callan, S. Coleman, J. Wess and B. Zumino, Phys. Rev. 177 (1969) 2247.

[17] K. Kawarabayashi, M. Suzuki, Phys. Rev. Lett. 16 (1966) 255; Riazuddin, Fayyazuddin, Phys. Rev. 147 (1966) 1071.

[18] P. Avery et al. (CLEO Collaboration), Phys. Lett. B331 (1994) 236.

[19] S. Barlag et al., (ACCMOR Collaboration), Phys. Lett. B278 (1992) 480.

[20] P. Manakos and T. Mannel, Z. Phys. C52 (1991) 659.

[21] K. Kodama et al., (E653 Collaboration), Phys. Lett. B286 (1992) 187.

[22] J. Anjos et al., (E691 Collaboration), Phys. Rev. Lett. 62 (1989) 1587; Phys. Rev. Lett. 65 (1990) 2630.

[23] Z. Bai et al., (MARK III Collaboration), Phys. Rev. Lett. 66 (1991) 1011.