CHIRAL DOUBLING OF HEAVY–LIGHT HADRONS: BABAR 2317 MeV/c^2 AND CLEO 2460 MeV/c^2 DISCOVERIES

Maciej A. Nowak

M. Smoluchowski Institute of Physics, Jagellonian University
Reymonta 4, 30-059 Kraków, Poland
and
Gesellschaft für Schwerionenforschung, Planckstrasse 1
D-64291 Darmstadt, Germany

Mannque Rho

Service de Physique Théorique, C.E. Saclay
91191 Gif-sur-Yvette, France
and
School of Physics, Korea Institute for Advanced Study
Seoul 130-722, Korea

and Ismail Zahed

Department of Physics and Astronomy
State University of New York at Stony Brook, New York 11794, USA

(Received July 14, 2004)

We point out that the very recent discoveries of BaBar (2317) and CLEO II (2460) are consistent with the general pattern of spontaneous breaking of chiral symmetry in hadrons built of heavy *and* light quarks, as originally suggested by us in 1992 (*Phys. Rev.* **D48**, 4370 (1993)), and independently by Bardeen and Hill in 1993 (*Phys. Rev.* **D49**, 409 (1994)). The splitting between the chiral doublers follows from a mixing between the light constituent quark mass and the velocity of the heavy quark, and vanishes for a zero constituent quark mass. The strictures of spontaneous chiral symmetry breaking constrain the axial charges in the chiral multiplet, and yield a mass splitting between the chiral doublers of about 345 MeV when the pion coupling to the doublers is half its coupling to a free quark. The chiral corrections are small. This phenomenon is generic and extends to all heavy–light hadrons. We predict the mass splitting for the chiral doublers of the excited mesons (D_1, D_2). We suggest that the heavy–light doubling can be used to address issues of chiral symmetry restoration in dense and/or hot hadronic matter. In particular, the relative splitting between D and D^* mesons and their chiral partners decreases in matter, with consequences on charmonium evolution at RHIC.

PACS numbers: 12.38.Aw, 12.39.Hg, 12.40.Yx

1. Introduction

On April 12th 2003, the BaBar collaboration announced a narrow peak of mass 2.317GeV/c^2 that decays into $D_s^+\pi^0$ [3]. On May 12th 2003, the CLEO II collaboration confirmed the BaBar result, and also observed a second narrow peak of mass 2.46 GeV/c^2 in the final $D_s^{*+}\pi^0$ state [4]. Both discoveries triggered a flurry of theoretical activity [5, 6], especially in light of the first reports and the press release announcing that the discovery is in disagreement with theoretical predictions.

In this note, we recall that actually the presence of these light states was predicted by theoretical arguments already in 1992 and 1993, and is in fact required from the point of view of symmetries of the QCD interactions. The two particles observed by BaBar and CLEO II are the first chiral partners of hadrons theoretically anticipated built out of light and heavy quarks. As such, they represent rather a *pattern* of spontaneous breakdown of chiral symmetry than isolated events.

Strong interactions involve three light flavors (u, d, s) and three heavy flavors (c, b, t) with respect to the QCD infrared scale. The light sector (l) is characterized by the spontaneous breaking of chiral symmetry, while the heavy sector (h) exhibits heavy-quark (Isgur–Wise) symmetry [7]. In our original work [1] we addressed the question of the form of the heavy–light effective action in the limit where light flavors are massless, while the heavy flavors are infinitely massive. Our chief observation was that a consistent implementation of the spontaneous breaking of chiral symmetry requires in addition to the known $(0^-, 1^-)$ heavy–light D-mesons, new and unknown heavy–light chiral partners $(0^+, 1^+)$ referred to as $\tilde{D}$-mesons. In the heavy-quark limit, the $D\tilde{D}$-splitting is small and of the order of the "constituent quark mass". Surprisingly, the approximate pattern of spontaneous-symmetry breaking observed in light–light systems carries even to heavy–light systems, in contrast to established lore based on Coulomb bound states.

This paper goes beyond recalling the chiral results in [1] in four significant ways: *(i)* it details the physical mechanism behind the splitting between the heavy–light chiral partners; *(ii)* it establishes three novel and generalized chiral Ward identities away from the chiral limit valid for both two and three flavors with massive pseudoscalar mesons. The chiral corrections to the splitting are shown to be of order $m_\pi^2/4m_h$, and therefore small irrespective of an effective Lagrangian analysis. *(iii)* it derives a new constraint between the axial charges of the heavy–light chiral partners; *(iv)* it sets the mass splitting between the chiral partners of the excited states.

2. One-loop results

To one-loop approximation, the order m_h^0 contribution to the heavy–light effective action follows from the diagrams shown in Figs. 1 and 2 in a constituent quark model[1] with light quarks of constituent mass Σ and heavy

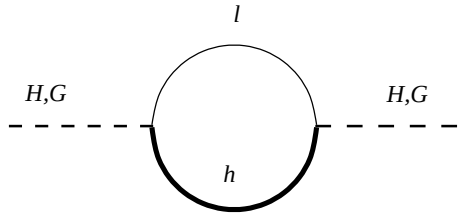


Fig. 1. One-loop contribution to 2-point HH, GG functions. Here l stands for light quark and h for heavy quark.

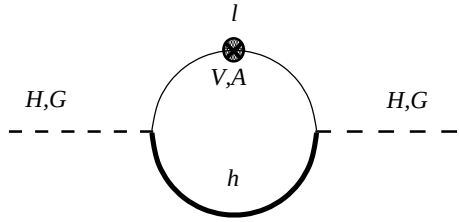


Fig. 2. One-loop contribution to 3-point HHV, HHA, GGV, GGA functions. l and h are as in Fig. 1, V and A stand, respectively, for the external vector and axial-vector sources.

and non-relativistic fields of residual mass set to zero (modulo reparametrization invariance) and a momentum cut-off Λ. The result for the $(0^-, 1^-)$ in the presence of vector and axial vector currents V, A is [1]

$$\begin{aligned}\mathcal{L}^H = \; & -\frac{i}{2}\mathrm{Tr}(\bar{H}v^\mu\partial_\mu H - v^\mu\partial_\mu\bar{H}H) \\ & +\mathrm{Tr}V_\mu\bar{H}Hv^\mu - \boldsymbol{g}_H\mathrm{Tr}A_\mu\gamma^\mu\gamma_5\bar{H}H \\ & -\boldsymbol{m}_H(\Sigma)\,\mathrm{Tr}\bar{H}\,H\,, \end{aligned} \tag{1}$$

where $\boldsymbol{m}_H \approx -\Sigma$ is an induced (cut-off dependent) chiral mass reflecting the dynamical generation of mass ensuing from spontaneously broken chiral symmetry, $\boldsymbol{g}_H$ an induced (cut-off) dependent axial coupling and H the

[1] We have specifically in mind the effective chiral quark model of Manohar and Georgi [8].

dimension 3/2 pseudoscalar-vector multiplet [9],

$$H = \frac{1+\not{v}}{2}(\gamma^\mu D^*_\mu + i\gamma_5 D) \tag{2}$$

with a transverse vector field, *i.e.* $v \cdot D^* = 0$. The Trace in (1) is over flavor and spin. The result is in agreement with known results [10–12] with the exception of the chiral mass contribution missing in these works. The origin and physical implications of the latter is important as we now discuss.

The novel aspect of our original derivation was that consistency with the general principles of spontaneously broken chiral symmetry requires the introduction of *chiral partners* in the form of a $(0^+, 1^+)$ multiplet of pseudoscalars and transverse vectors [1]

$$G = \frac{1+\not{v}}{2}(\gamma^\mu\gamma_5 \tilde{D}^*_\mu + \tilde{D})\,. \tag{3}$$

To leading order in the heavy-quark mass, the one-loop effective action for the $(0^+, 1^+)$ duplicates (1) with a key difference in the sign of the constituent mass contribution. Specifically [1]

$$\begin{aligned}\mathcal{L}^G = \quad & -\frac{i}{2}\mathrm{Tr}(\bar{G}v^\mu\partial_\mu G - v^\mu\partial_\mu\bar{G}G) \\ & +\mathrm{Tr}V_\mu\bar{G}Gv^\mu - \boldsymbol{g}_G\mathrm{Tr}A_\mu\gamma^\mu\gamma_5\bar{G}G \\ & -\boldsymbol{m}_G(\Sigma)\,\mathrm{Tr}\bar{G}\,G\end{aligned} \tag{4}$$

with the induced (cutoff dependent) chiral mass $\boldsymbol{m}_G \approx +\Sigma$ (note the sign flip in comparison to $\boldsymbol{m}_H$). Both chiral mass contributions are invariant under rigid chiral $\mathrm{SU(2)_L} \times \mathrm{SU(2)_R}$ and local $\mathrm{SU(2)}_V$ symmetry [1]. This construction is generic and generalizes for any number of light flavors. In the further part of the paper, we argue, why the chiral shifts are almost identical even in the case of the strange quark mass of order 150 MeV. In Section 3 we also point, that sub-leading chiral effects to Goldberger–Treiman relations, which are crucial for model-independent estimation of the chiral shifts, are strongly suppressed already for the masses of charm quarks.

The sign flip follows from the γ_5 difference in the definition of the fields H and G, in other words the parity assignment. Indeed, the mass contribution arising from Fig. 1 has the generic structure (constant H)

$$\mathrm{Tr}\left(\boldsymbol{P}_2\,\frac{\not{Q}+\Sigma}{Q^2-\Sigma^2}\,H\boldsymbol{P}_3\frac{\not{v}}{(v\cdot Q)}\bar{H}\right) \tag{5}$$

and similarly for $H \to G$. The trace is over 4-momentum Q, spin and flavor with $\boldsymbol{P}_2 = \mathrm{diag}(1,1,0)$ and $\boldsymbol{P}_3 = \mathrm{diag}(0,0,1)$. The range in Q is $0 < Q < \Lambda$

where Λ is an ultraviolet cut-off. We note that in (5) only the contribution

$$\mathrm{Tr}\left(\boldsymbol{P}_2\,\frac{\Sigma}{Q^2-\Sigma^2}\,H\boldsymbol{P}_3\frac{\not v}{v\cdot Q}\bar H\right) \tag{6}$$

is sensitive to the parity content of the heavy–light field since $H\not v = -H$ and $G\not v = +G$. The result is a split between the heavy–light mesons of opposite chirality. This unusual contribution of the chiral quark mass stems from the fact that it tags to the *velocity* $H\not v\bar H$ of the heavy field and is therefore sensitive to *parity*. It is not affected by a shift Δ in the heavy quark mass, which amounts to the substitution

$$\frac{\not v}{v\cdot Q}\to\frac{\not v\, v\cdot Q+\Delta}{(v\cdot Q)^2-\Delta^2} \tag{7}$$

which is seen to shift H and G in the same direction. The reparametrization invariance (invariance under velocity shifts of the heavy quark to order one) introduces mass shifts that are parity insensitive to leading order in $1/m_h$ [13].

The HG-mass difference is dictated by the spontaneous breaking of chiral symmetry, modulo the U(1) anomaly through instantons which will be discussed elsewhere. If we recall that the H, G fields carry mass dimension 3/2 through a rescaling of the complex dimension 1 fields by $\sqrt{m_h}$, it follows from the normalizations of the kinetic and mass term in (1) and in (4) that

$$\begin{aligned} m_H &= \quad m_h+\boldsymbol{m}_H\,,\\ m_G &= \quad m_h+\boldsymbol{m}_G \end{aligned} \tag{8}$$

in the chiral and heavy-quark limit. In retrospect, this result can be arrived at simply as follows: *(i)* the light quark contributes a mass shift of order of an induced cut-off dependent constituent mass Σ; *(ii)* it is repulsive in the scalars (no $i\gamma_5$) and attractive in the pseudoscalars (with $i\gamma_5$). In this limit, the spontaneous breakdown of chiral symmetry enforces the mass relation [1]

$$m(\tilde D^*)-m(D*)=m(\tilde D)-m(D)=\boldsymbol{m}_G-\boldsymbol{m}_H \tag{9}$$

since the dispersion relation is linear after the heavy mass reduction. The interaction term is given by

$$\mathcal{L}_{HG}=+\sqrt{\frac{\boldsymbol{g}_G}{\boldsymbol{g}_H}}\,\mathrm{Tr}(\gamma_5\bar G H\gamma^\mu A_\mu)-\sqrt{\frac{\boldsymbol{g}_H}{\boldsymbol{g}_G}}\,\mathrm{Tr}(\gamma_5\bar H G\,\gamma^\mu A_\mu) \tag{10}$$

with no vector mixing because of parity. We note that (10) follows from the expansion of a Dirac operator with external vector (V) and axial-vector (A)

sources and is in general complex since the Dirac operator is not self-adjoint (this point is at the origin of flavor anomalies).

These results are expected to hold qualitatively in the presence of non-zero current quark masses, modulo the re-summation of standard chiral logs (chiral perturbation theory) and the U(1) anomaly (instantons). The interaction term accounts for the strong decay of heavy mesons via emission of Goldstone bosons $\tilde{D} \to D\,\pi$.

In [1] we used a constituent quark model to one-loop to estimate the pertinent parameters in (1), (4) and (10) which were found to be sensitive to the cut-off procedure used in regulating the one-loop of Figs. 1 and 2. For a general covariant cutoff, the mass splitting for a large cutoff limit is

$$\boldsymbol{m}_G - \boldsymbol{m}_H = 2\Sigma\left(1 - \frac{1}{4\pi^2}\ln\frac{\Lambda}{\Sigma}\right) \tag{11}$$

with equal and finite axial charges $\boldsymbol{g}_H \approx \boldsymbol{g}_G \approx 1/3$, such that (10) reduces to

$$\mathcal{L}_{HG} = \mathrm{Tr}(\gamma_5(\bar{G}H - \bar{H}G)\gamma^\mu A_\mu)\,. \tag{12}$$

The pertinent loop integrals can be found in [1]. All integrals were evaluated with Minkowski metric and a covariant 4-dimensional cutoff to preserve reparametrization invariance. For a finite cutoff the results were quoted in [1], with small effects on the splitting and somehow larger effects on one of the axial coupling. The logarithmic sensitivity of the mass splitting is weak but expected in field theory. The approach advocated in [1] is the Wilsonian approach with a finite and physical cutoff Λ to separate between the hard modes of order m_h and the soft modes. Clearly, the present results are also sensitive to a residual mass shift Δ in the heavy quark mass. These effects are harder to track down in the Wilsonian approach we have followed due to the reparametrization invariance of the formulation. These are easier to track *e.g.* in dimensional regularization scheme, however both schemes differ due to the presence of strong power divergences in the loop integrations. In general, these ambiguities call for a first principle calculation using lattice QCD simulations or the instanton vacuum model [14,16,19].

The effects of a light current quark mass m_l can be estimated *e.g.* in the aforementioned instanton vacuum model [16]. A simple parametrization with good comparison to lattice data [17] was quoted in [18]

$$\Sigma(m_l) \approx m_l + \Sigma\left(\sqrt{1 + \left(\frac{m_l}{d}\right)^2} - \frac{m_l}{d}\right) \tag{13}$$

with $\Sigma \approx 345\,\mathrm{MeV}/c^2$, $d \approx \sqrt{0.08/2N_c}8\pi\rho/R^2 \approx 198\,\mathrm{MeV}/c^2$ for a standard instanton size $\rho \approx 1/3$ fm and interinstanton distance $R \approx 1$ fm. For a

strange quark mass $m_s \approx 150\,\mathrm{MeV}/c^2$, the second term is reduced to $\Sigma/2$, making the combination (13) weakly dependent on m_l and of order Σ all the way up to the strange quark mass. Thus, both mass splittings are about the same for (u,d,s) heavy–light mesons. The width of the non-strange heavy light partners is however not restricted by kinematics as in the case of D_s, hence these particles may be much broader and harder to detect.

3. Goldberger–Treiman relations

However, in our case chiral symmetry offers further important constraints on the spontaneous generation of mass and the ensuing pion-H-G interactions. This allows for model independent relations between the cut-off dependent parameters discussed above. Indeed, in the pure pion model discussed originally in [1] (no vector dominance) the axial-vector current A_μ in (1), (4), (12) is purely pionic and reads

$$A_\mu = \frac{i}{2}\left(\xi\partial_\mu\xi^\dagger - \xi^\dagger\partial_\mu\xi\right) \tag{14}$$

with $\xi = e^{i\pi/2f_\pi}$. Inserting (14) into (1), (4), (12) yields the pseudovector π–HG interactions

$$\begin{aligned}
\mathcal{L}_{\pi H} &= \frac{\boldsymbol{g}_H}{2f_\pi}\,\mathrm{Tr}\left(\gamma^\mu\gamma^5\bar{H}H\,\partial_\mu\pi\right)\,,\\
\mathcal{L}_{\pi G} &= \frac{\boldsymbol{g}_G}{2f_\pi}\,\mathrm{Tr}\left(\gamma^\mu\gamma^5\bar{G}G\,\partial_\mu\pi\right)\,,\\
\mathcal{L}_{\pi HG} &= \frac{\boldsymbol{g}_{HG}}{2f_\pi}\,\mathrm{Tr}\left(\gamma^\mu\gamma^5(\bar{G}H-\bar{H}G)\,\partial_\mu\pi\right)\,,
\end{aligned} \tag{15}$$

where for generality we introduced the axial transition coupling $\boldsymbol{g}_{HG}$ which is 1 in (12). Integrating by parts in (15) and using the transversality of the heavy-vector fields result in a single Goldberger–Treiman relation from the last of the tree couplings in (15) [2]

$$\frac{1}{2}(m_G - m_H) \approx \frac{1}{2}(\boldsymbol{m}_G - \boldsymbol{m}_H) = \frac{f_\pi\,\boldsymbol{g}_{\pi HG}}{\boldsymbol{g}_{HG}}\,, \tag{16}$$

with $\boldsymbol{g}_{\pi HG}$ the pion pseudoscalar coupling to the chiral doublet in the heavy and chiral limit. This relation was originally observed in [2] up to a missing factor of 1/2. It involves the splitting between the even–odd partners which is less sensitive to Δ. This relations is slightly modified if vector dominance

[2] The transversality of the vectors in H and G yields zero pseudoscalar couplings from the first two relations in (15). This point will be clarified below.

is enforced in the heavy–light sector [1], *i.e.* the rhs of (16) is divided by $(m_\rho/m_{a_1})^2$ since part of the pion field is eaten up by the a_1 through a Higgs-like mechanism. The corrections due to a finite pion mass m_π and a large but finite heavy quark mass m_h will be discussed below on general grounds.

For comparison, we recall that the constituent quark mass obeys the Goldberger–Treiman relation [8]

$$\Sigma = \frac{f_\pi \, \boldsymbol{g}_{\pi qq}}{\boldsymbol{g}_A} \tag{17}$$

with $\boldsymbol{g}_A \approx 0.75$ and $\boldsymbol{g}_{\pi qq} \approx \boldsymbol{g}_{\pi NN}/3 \approx 3.3$. If we were to use $\boldsymbol{g}_{\pi HG} \approx \boldsymbol{g}_{\pi qq}/2$ and $\boldsymbol{g}_{HG} \approx \boldsymbol{g}_A$ it follows from the last relation in (16) that the splitting in the chiral multiplet would be one constituent quark mass

$$m_G - m_H \approx \boldsymbol{m}_G - \boldsymbol{m}_H \approx \Sigma \tag{18}$$

which speaks for a large cutoff in (11). That the pion coupling to the chiral multiplet is $1/2$ its coupling to the free light quark is forced upon us by the BaBar and CLEO II results. This may be understood as a sign of nontrivial screening mechanism in action in the presence of the heavy quark, that is the pion is "busy" half the time with the massive quark.

The deviations from the heavy and chiral limits to (16) can be assessed using the general framework for spontaneous breaking of chiral symmetry developed in [20]. Within this approach, the one-pion reduced axial transition $\tilde{D} \to D\pi$ reads

$$\langle D(p_2)|\boldsymbol{j}^a_{A\mu}(0)|\tilde{D}(p_1)\rangle = \left(\frac{(p_1-p_2)_\mu}{m_D+m_{\tilde{D}}}\,\boldsymbol{G}_1(t) + \frac{(p_1+p_2)_\mu}{m_D+m_{\tilde{D}}}\,\boldsymbol{G}_2(t)\right) D^\dagger \frac{\tau^a}{2}\tilde{D}\,, \tag{19}$$

where $\boldsymbol{j}^a_{A\mu}$ is the one-pion reduced axial vector current satisfying [20]

$$\partial^\mu \boldsymbol{j}^a_{A\mu}(x) = f_\pi \left(\Box + m_\pi^2\right) \pi^a(x)\,. \tag{20}$$

The first form factor in (19) is one-pion reduced, and the $D, \tilde{D}$ on the RHS are unit isospinors. For $p_1 = p_2$ (at rest) the axial charge follows from $\mu = 0$ as

$$\boldsymbol{G}_2(0)\, D^\dagger \frac{\tau^a}{2}\, \tilde{D}$$

which identifies $\boldsymbol{G}_2(0)$ with the properly normalized axial charge in the transition matrix element. Inserting (20) into (19) gives

$$\begin{aligned}\langle D(p_2)|\pi^a(0)|\tilde{D}(p_1)\rangle \;&=\; \frac{1}{f_\pi}\frac{1/(m_D+m_{\tilde{D}})}{m_\pi^2 - t} \\ &\times \left(t\,\boldsymbol{G}_1(t) + (m^2_{\tilde{D}} - m^2_D)\,\boldsymbol{G}_2(t)\right) D^\dagger \frac{\tau^a}{2}\,\tilde{D}\,. \end{aligned} \tag{21}$$

By definition, the π–$D\tilde{D}$ coupling is

$$\langle D(p_2)|\pi^a(0)|\tilde{D}(p_1)\rangle = \boldsymbol{g}_{\pi D\tilde{D}}(t)\,\frac{1}{m_\pi^2 - t}\,D^\dagger\tau^a\,\tilde{D} \tag{22}$$

which corresponds to

$$\boldsymbol{g}_{\pi D\tilde{D}}\,\pi^a\,\left(\tilde{D}^\dagger\tau^a\,D + \text{h.c.}\right).$$

A comparison of (22) with (21) gives at the pion pole $t \approx m_\pi^2$

$$f_\pi\,\boldsymbol{g}_{\pi D\tilde{D}}(m_\pi^2) = +\frac{1}{2}(m_{\tilde{D}} - m_D)\,\boldsymbol{G}_2(m_\pi^2) + \frac{1}{2}\frac{m_\pi^2}{(m_D + m_{\tilde{D}})}\boldsymbol{G}_1(m_\pi^2) \tag{23}$$

which is the general form of the Goldberger–Treiman relation for the transition amplitude $\tilde{D} \to D\pi$. In the (double) heavy and chiral limit it reduces to (16) with the identifications $\boldsymbol{g}_{\pi D\tilde{D}} = \boldsymbol{g}_{\pi HG}$ and $\boldsymbol{G}_2(0) = \boldsymbol{g}_{HG}$. The second chiral correction in (23) is the analog of the πN sigma term. In our case this amounts to a chiral correction of order $m_\pi^2/4m_h \ll m_\pi$ to (16) which is negligible.

Similar arguments can be employed to analyze the Goldberger–Treiman relations corresponding to the πHH and the πGG couplings in (15). For instance, the one-pion reduced axial transition $D^* \to D\pi$ yields

$$\begin{aligned}\langle D(p_2)|\boldsymbol{j}^a_{A\mu}(0)|D^*(p_1,\epsilon)\rangle \;=\; & (\epsilon_\mu(m_D + m_{D^*})\,\boldsymbol{H}_1(t) \\ & +(p_1 - p_2)_\mu\,\epsilon\cdot(p_1 - p_2)\boldsymbol{H}_2(t) \\ & +(p_1 + p_2)_\mu\epsilon\cdot(p_1 - p_2)\boldsymbol{H}_3(t)) \\ & \times(m_D + m_{D^*})^{-1}\,D^\dagger\frac{\tau^a}{2}D^*\,,\end{aligned} \tag{24}$$

where ϵ is the covariantly transverse vector polarization of the D^*. Again, $\boldsymbol{H}_2$ is one-pion reduced, and the D and D^* on the RHS are unit isospinors. Using the π–DD^* coupling given by

$$\frac{\boldsymbol{g}_{\pi DD^*}}{(m_D + m_{D^*})}\,\pi^a\,\left((\partial_\mu\,D^\dagger)\,\tau^a\,D^{*\mu} + \text{h.c.}\right) \tag{25}$$

which is

$$\langle D(p_2)|\pi^a(0)|D^*(p_1,\epsilon)\rangle = \frac{\boldsymbol{g}_{\pi DD^*}}{(m_D + m_{D^*})}\,\frac{\epsilon\cdot(p_1 - p_2)}{m_\pi^2 - t}D^\dagger\tau^a\,D^* \tag{26}$$

and a rerun of the preceding arguments yield

$$\begin{aligned}2f_\pi\,\boldsymbol{g}_{\pi DD^*}(m_\pi^2) \;=\; & (m_D + m_{D^*})\,\boldsymbol{H}_1(m_\pi^2) \\ & +m_\pi^2\,\boldsymbol{H}_2(m_\pi^2) + (m_{D^*}^2 - m_D^2)\boldsymbol{H}_3(m_\pi^2)\,.\end{aligned} \tag{27}$$

In the heavy and chiral limit, we have

$$\begin{aligned} f_\pi \, \boldsymbol{g}_{\pi HH} &= m_H \, \boldsymbol{g}_H \, , \\ f_\pi \, \boldsymbol{g}_{\pi GG} &= m_G \, \boldsymbol{g}_G \end{aligned} \tag{28}$$

with $\boldsymbol{H}_1(0) = \boldsymbol{g}_H$. The last relation follows from an identical reasoning. The first relation in (28) was noted by Nussinov and Wetzel [21] and yields semileptonic decay widths that are consistent with data. Equation (27) gives its general chiral corrections. Note that the mass of the heavy quark m_h appears explicitly in (28) which is the chief reason for why these relations were not *a priori* accessible from (15) through an integration by part as (15) involves solely the soft scales[3]. Combining (28) with (16) leads to a relation between the various axial couplings

$$\frac{\boldsymbol{g}_{\pi GG}}{\boldsymbol{g}_G} - \frac{\boldsymbol{g}_{\pi HH}}{\boldsymbol{g}_H} = 2 \, \frac{\boldsymbol{g}_{\pi HG}}{\boldsymbol{g}_{HG}} \, . \tag{29}$$

The π–GG and π–HH couplings are fixed by semi-leptonic decays, thereby constrain the axial charges in (28), (29). Clearly, the results (19)–(29) are properties of QCD and should be reproduced by *any* attempt to explain the strong decay $\tilde{D} \to D\pi$, *i.e.* the BaBar and CLEO results. The original approaches [1,2] fulfill these constraints by construction.

For completeness, we note that the systematics of subleading $(1/m_h)$ corrections for the couplings between the particles from H multiplet, for the couplings between the particles from G multiplet and the axial current mediated G–H transitions could be addressed following and adapting the approach outlined in [22].

4. BaBar and CLEO results

As a whole, the experimental results of BaBar and CLEO are overall consistent with the chiral doubling proposal:

(i) The even–odd parity mass shifts are the same in the spin 0 and 1 channels and of the order of the constituent quark mass of ~ 345 MeV,

$$\begin{aligned} m(\tilde{D}_s^+(2316.8)) - m(D_s^+) &= 348.3 \text{ MeV}/c^2 \, , \\ m(\tilde{D}_s^+(2316.8)) - m(D_s^+) &= 350.4 \pm 1.2 \pm 1.0 \text{ MeV}/c^2 \, , \\ m(\tilde{D}_s^{*+}(2463)) - m(D_s^{*+}) &= 351.6 \pm 1.7 \pm 1.0 \text{ MeV}/c^2 \, , \end{aligned}$$

where we used our original "tilde" notation, for the two new particles. The first quote is from BaBar, while the last two quotes are from CLEO II.

[3] If we were to assume an arbitrary momentum for the mass shell condition, (28) could be arrived at from (15) through a simple integration by part.

(ii) The decay widths of the strange even parity states are very small owing to the lightness of Σ, shutting off the natural kaon decay mode $\tilde{D} \to DK$, and operating chiefly through the isospin violating mode $\tilde{D} \to D(\eta \to \pi^0)$. This is overall consistent with our interaction term (10).

(iii) No photonic (vector) channels were found for transitions between the pairs $D_s(2317)$, $D_s(1969)$ or $D_s(2460)$, $D_s^*(2112)$.

In [13], we further pointed out that chiral partners are also expected for the excited $(1^+, 2^+) = (D_1, D_2)$ mesons in the form of a $(1^-, 2^-)$ chiral pair. Our prediction for the masses are:

$$\begin{aligned} m(\tilde{D}_{s1}) &= 2721 \pm 10 \text{ MeV}\,, \\ m(\tilde{D}_{s2}) &= 2758 \pm 10 \text{ MeV}\,, \end{aligned} \tag{30}$$

where we used as an input the observed BaBar and CLEO splitting for the chiral multiplet $(0^+, 1^+)$ and the mass formulae obtained in [13]. Generalized Goldberger–Treiman relations for the excited states can also be derived using the general arguments presented above.

We expect a similar splitting for the non-strange heavy–light mesons, in particular a splitting of about 368 MeV between the $D_{u,d}$ and their chiral partners $\tilde{D}_{u,d}$. The chiral doubling should be even more pronounced for bottom mesons, since the $1/m_h$ corrections are three times smaller. For $m_s = 150$ MeV, we expect the chiral partners of B_s and B_s^* to be 323 MeV heavier, while the chiral partners of B and B^* to be 345 MeV heavier. We note that any observation of chiral doubling for B mesons would be a strong validation for our proposal. Indeed, in the recently proposed alternative scenarios [5] (multiquark states, hadronic molecules, modifications of quark potential, unitarization) a repeating pattern from charm to bottom calls for additional assumptions.

Bardeen and Hill [2] suggested a "solvable toy field-theoretical model" and arrived at totally analogous results for chiral partners of D and D^* mesons, by using a similar one-loop calculation. Their Nambu–Jona-Lasinio model after Fierz transformation and to one-loop approximation reduces to our effective action construction, hence the consistency between our results and theirs. As far as we know [1] and [2] were the only early predictions of the phenomenon of chiral doubling for charmed and bottomed hadrons involving light quarks. This idea was later developed further in other papers [23].

Soon after the BaBar announcement, several theoretical papers appeared [5,6] suggesting a variety of explanations for the newly observed state. In particular, Bardeen, Eichten and Hill [6] adapted the effective chiral action [1,2] to three light flavors, exploiting a constituent-quark version of Goldberger–Treiman relation and fixing the unknown parameter of the effective Lagrangian to the experimentally observed splitting. The results

of their calculations which are in remarkable agreement with experiments provide a good confirmation to our and their early suggestions for a chiral doubling in the heavy–light sector of QCD [1,2].

Many issues regarding heavy–light systems in the QCD instanton vacuum were discussed in [14,15] including constituent heavy-baryons such as qqQ and qQQ and exotics such as $\bar{Q}\bar{q}qq$, $\bar{Q}qqqq$. In particular, it was suggested that the heavy–light H-dibaryon ($Qqq\,Qqq$) with $Q = c, b$ is *bound* owing to the smallness of the three-body force in the presence of the heavy quark (about 10% the value of the two-body force). It may even have a bound chiral partner. Finally, it is worth pointing out that the successful treatment of heavy–light baryons as solitons [24,25] could be readily extended to the chiral doubling now revealed in the heavy–light systems.

5. Summary

In this note, we have pointed out that the newly discovered charmed mesons by BaBar and CLEO are chiral partners of charmed and bottomed hadrons that include at least one light quark, a pattern suggested a decade ago [1,2]. The result is a chiral splitting between the even and odd parity partners of about a constituent quark mass as reported recently by BaBar and CLEO. More chiral partners are expected. It may be a bit of surprise that the pion coupling to the heavy–light chiral multiplet comes out to be $1/2$ its coupling to a free quark. The experimental results are telling us that it should be so. Although we do not have a rigorous argument to justify it, we conjecture that the "screening" results since the pion is "busy" half of the time with the heavy quark in the chiral multiplet. A consistent treatment of the parity doubling in the heavy–light systems — which can answer this as well as other questions — can be achieved in the QCD instanton vacuum which is parameter-free, since the vacuum dynamics is totally fixed in the light–light systems. We have shown that the chiral corrections are small.

QCD implies chiral Ward identities in the heavy–light systems in the form of generalized Goldberger–Treiman relations. The even–odd splitting is constrained by one of them. Any explanation of the strong decay $\tilde{D} \to D\pi$ should abide by these constraints, in particular (23). The chiral doubling approach used in [1,2] fulfills these identities by construction in the heavy and chiral limit. For a plausible axial charge of unity for the $D\tilde{D}$-transition amplitude, the observed small splitting of about 345 MeV by BaBar and CLEO is uniquely explained by a small π–$D\tilde{D}$ coupling of about half its value to a constituent light quark. This conclusion is generic to QCD and should therefore be reached by all the recently proposed alternative scenarios [5] if they were to be viable. Chiral doubling is then an immediate consequence of rigid chiral symmetry from quantum numbers.

Particularly relevant to the on-going effort to gain a deeper understanding of strong interactions is the question: To what extent can the newly discovered chiral partners shed light on the changes of the QCD vacuum caused by external parameters such as temperature and/or baryon density? This is an important issue in light of the current and future experiments at RHIC and LHC as well as at SIS 300 [26]. It is also an interesting possibility for lattice simulations. Since the chiral partners are split by the dynamically generated chiral quark mass, it is likely that through a chiral phase transition D and D^* should move towards their chiral partners $\tilde{D}$ and $\tilde{D}^*$ to reduce to a degenerate chiral multiplet. This should prove particularly important for charmonium absorption/regeneration in thermal models with medium effects. Also, this can serve as a "litmus gauge" for the size of the chiral condensate in varying temperature and/or density as manifested in the properties of hadrons in hot/dense medium [27]. In the case of the D_s partners the restoration will not be complete due to the substantial current mass of the strange quark. The restoration of chiral symmetry in light–light systems has spurred many activities in the past (for a recent phenomenological discussion see [28]) and we expect this to extend now to the heavy–light systems.

We are pleased that the BaBar and CLEO II results are generating so much excitement in both the experimental and theoretical high energy/nuclear physics community, and it is gratifying that our old ideas have come full circle, with so many new theoretical venues and experimental possibilities.

6. Note added

After submitting the paper to the database, we became aware of the new experimental results:

1. Belle collaboration [29], announced two new $c\bar{u}$ states D_0^* $(2308 \pm 17 \pm 15 \pm 20)$ MeV and D_1^{*0} $(2427 \pm 26 \pm 20 \pm 15)$ MeV, with the spin-parity assignment $(0^+, 1^+)$. They are likely to be the chiral partners of the nonstrange D and D^* $(0^-, 1^-)$ multiplet. As expected they are much broader compared to the D_s states. The intriguing pattern observed, *i.e.*, that these two new states are almost as heavy as the corresponding strange multiplet discovered by BaBar and CLEO, is again in qualitative agreement with our chiral doubling arguments above.

2. Belle collaboration [30] has confirmed observations of BaBar and CLEO states, and provided substantial evidence for the spin-parity assignment $0^+, 1^+$.

3. SELEX has announced a new very narrow state $D_s(2632)$ [31]. It is tempting to speculate, that this state may be a chiral partner of the $D_{s1}(2536)$. If indeed this is the case, the data tell that the chiral shift for excited mesonic states is only of order 100 MeV, *i.e.* even smaller than the $O(m_c^0)$ estimate of order $\Sigma/2 \sim 175$ MeV given in this paper. An immediate consequence of such shift for excited mesonic chiral doublers is the presence of the chiral doubler for D_{s2}. Since $1/m_c$ split between the known D_{s1} and D_{s2} states is equal to 37 MeV, we expect a new state at $2632+37 = 2669$ MeV, with spin-parity assignment 2^-, with similar decay patterns to SELEX state, *i.e.* preferably decaying into D_s^* and η.

We would like to thank Henryk Pałka and James Russ for drawing our attention to the recent Belle [29] and SELEX [31] results. This work was partially supported by the Polish State Committee for Scientific Research (KBN) grant 2P03B 09622 (2002-2004), and by the US DOE grant DE-FG-88ER40388.

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