On heavy–light meson resonances and chiral symmetry

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Abstract

We study heavy–light meson resonances with quantum numbers $J^P = 0^+$ and $J^P = 1^+$ in terms of the non-linear chiral $SU(3)_L$ Lagrangian. At leading order a parameter-free prediction is obtained for the scattering of Goldstone bosons off heavy–light pseudo-scalar and vector mesons. The recently announced narrow open-charm states observed by the BABAR and CLEO Collaborations are reproduced. We suggest the existence of states that form an anti-triplet and a sextet representation of the $SU(3)_L$ group. In particular, so far unobserved narrow isospin-singlet states with negative strangeness are predicted. The open-beauty states with $(I, S) = (0, -1)$ are anticipated at 5761 MeV ($J^P = 0^+$) and 5807 MeV ($J^P = 1^+$). For the anti-triplet states our results differ most significantly from predictions that are based on the chiral quark model in the open-beauty sector. Strongly bound $0^+$- and $1^+$-states with $(I, S) = (0, 1)$ at 5643 and 5690 MeV are predicted.

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1. Introduction

In a recent work [1] it was demonstrated that chiral $SU(3)$ symmetry predicts parameter-free $J^P = 1^+$ light meson resonances. It was observed that the resonance states turn into bound states in the heavy $SU(3)$ limit with $m_{\pi, K, \eta} \simeq 500$ MeV but disappear altogether in the light $SU(3)$ limit with $m_{\pi, K, \eta} \simeq 140$ MeV. In earlier works [2–7] similar results were obtained for light meson resonances with $J^P = 0^+$ quantum numbers. In view of the apparent success of the chiral coupled-channel dynamics to predict the existence of a wealth of meson resonances in the $(u, d, s)$-sector of QCD it is interesting to study whether the same mechanism is able to predict heavy–light meson resonances, i.e., meson resonances with open-charm or beauty. Recently a new narrow state of mass 2.317 GeV that decays into $D^*_s\pi^0$ was announced [8]. This result was confirmed [9] and a second narrow state of mass 2.463 GeV decaying into $D^*_s\pi^0$ was observed. Such states were first predicted in [10,11] based on the spontaneous breaking of chiral symmetry. Since one expects from such studies [10–13] the existence of further so far unobserved states it is important to study the heavy–light resonances in great detail [14].

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The theoretical predictions [10–13] rely on the chiral quark model which predicts the heavy–light $0^+$, $1^+$ resonance states to form an anti-triplet representation of the SU(3) group. If one insists on a non-linear realization of the chiral SU(3) group only excluding any further model assumptions, no a priori prediction can be made for the existence of chiral partners of any given state. For instance in [1] is was shown that the light $1^+$ spectrum predicted by the non-linear chiral representation forms two degenerate octets and an additional singlet consistent with the empirical spectrum. Analogous results were obtained for the light scalar mesons [2–7].

In this Letter we apply the $\chi$-BS(3) approach developed originally for meson–baryon scattering [16–21] but recently also applied to meson–meson scattering [1]. Using the chiral SU(3) Lagrangian involving light–heavy $J^P = 0^-$ and $J^P = 1^-$ fields that transform non-linear under the chiral SU(3) group a coupled-channel description of the meson–meson scattering in the open charm and bottom sector is developed. The possible importance of coupled-channel dynamics for the heavy–light meson states was emphasized recently [15]. The major result of our work is the prediction that there exist states with $J^P = 0^+$, $1^+$ quantum numbers forming anti-triplet and sextet representations of the SU(3) group. This differs from the results implied by the chiral quark model leading to anti-triplet states only. Our result suggests the existence of $J^P = 0^+$, $1^+$ states with unconventional quantum numbers $(I, S) = (1, 1)$ and $(I, S) = (0, -1)$. A particular result concerns the ‘heavy’ SU(3) limit with $m_{\pi, n, K} \sim 500$ MeV and $M_D \sim 1800$ MeV in which the chiral coupled-channel dynamics predicts anti-triplet bound states rather than resonance states only. In the ‘light’ SU(3) limit with $m_{\pi, n, K} \sim 140$ MeV and $M_D \sim 1800$ MeV we do not find anymore resonances or bound states in the $J^P = 0^+$, $1^+$ sectors. Using physical mass parameters we predict narrow $J^P = 0^+$, $1^+$ states in the $(I, S) = (0, 1), (1, 0), (0, -1)$ channels with open-charm and open-beauty. The open-charm (0, 1) states established by the BABAR [8] and CLEO [9] Collaborations are recovered within 20 MeV accuracy. We identify the observed $D(2420)$ resonance with $J^P = 1^+$ and $(I, S) = (\frac{1}{2}, 0)$ to be a member of the sextet. This suggests the existence of isospin zero $\bar{K}D(1867)$- and $KD(2008)$-bound states.

2. Chiral coupled-channel dynamics: the $\chi$-BS(3) approach

The starting point to study the scattering of Goldstone bosons off heavy–light mesons is the chiral SU(3) Lagrangian. We identify the leading-order Lagrangian density [22–25] describing the interaction of Goldstone bosons with pseudo-scalar and vector mesons,

$$\mathcal{L}(x) = \frac{1}{8f^2} \text{tr} \left[ P(x) (i\gamma^\nu P^\dagger(x)) - (i\gamma^\nu P(x)) P^\dagger(x) \right] \times \left[ \Phi(x), (\partial_\mu \Phi(x)) \right]_{-}$$

$$- \frac{1}{8f^2} \text{tr} \left[ P^\mu(x) (i\gamma^\nu P^\mu_\nu(x)) - (i\gamma^\nu P^\mu_\nu(x)) P^\mu_\nu(x) \right] \times \left[ \Phi(x), (\partial_\mu \Phi(x)) \right]_{-},$$

where $\Phi$ is the Goldstone bosons field and $P$ and $P_\mu$ are massive pseudo-scalar and vector-meson fields. The parameter $f$ in (1) is known from the weak decay process of the pions. We use $f = 90$ MeV through out this work. Since we will assume perfect isospin symmetry it is convenient to decompose the fields into their isospin multiplets. The fields can be written in terms of isospin multiplet fields like $K = (K^{(+)}, K^{(0)})^T$ and $D = (D^{(+)}, D^{(0)})^T$.

$$\Phi = \tau \cdot \pi(140) + \alpha^T \cdot K(494) + K^\dagger(494) \cdot \alpha + \eta(547) \lambda_8,$n

$$P = \frac{1}{\sqrt{2}} \alpha^T \cdot D(1867) - \frac{1}{\sqrt{2}} D^\dagger(1867) \cdot \alpha$$

$$+ i \tau_2 D^{(s)}(1699),$$

$$P_\mu = \frac{1}{\sqrt{2}} \alpha^T \cdot D_\mu(2008) - \frac{1}{\sqrt{2}} D^{\dagger}_\mu(2008) \cdot \alpha$$

$$+ i \tau_2 D^{(s)}_\mu(2110),$$

$$\alpha^T = \frac{1}{\sqrt{2}} (\lambda_4 + i\lambda_5, \lambda_6 + i\lambda_7),$$

$$\tau = (\lambda_1, \lambda_2, \lambda_3),$$

where the matrices $\lambda_i$ are the standard Gell-Mann generators of the SU(3) algebra. The numbers in the brackets recall the approximate masses of the fields in units of MeV.

The scattering problem decouples into seven orthogonal channels specified by isospin $(I)$ and strange-
heavy–light mesons and small Goldstone bosons. The matrix of coefficients $M$ in coupled-channel space, depend on whether to scatter considered here.

In Table 1 the channels that contribute in a given sector $(I, S)$ are listed. Heavy–light meson resonances with quantum numbers $I^P = 0^+$ and $I^P = 1^+$ manifest themselves as poles in the $s$-wave scattering amplitudes, $M_{j^P}^{(I, S)}(\sqrt{s})$, which in the $\chi$-BS(3) approach \cite{1,19} take the simple form

$$M_{j^P}^{(I, S)}(\sqrt{s}) = \left[ 1 - V^{(I, S)}(\sqrt{s}) J_{j^P}^{(I, S)}(\sqrt{s}) \right]^{-1} V^{(I, S)}(\sqrt{s}).$$

The effective interaction kernel $V^{(I, S)}(\sqrt{s})$ in (4) is determined by the leading order chiral SU(3) Lagrangian (1),

$$V^{(I, S)}(\sqrt{s}) = \frac{C_{(I, S)}}{8f^2} \left( 3s - M^2 - \bar{M}^2 - m^2 - \bar{m}^2 \right) - \frac{M^2 - m^2}{s} \left( \bar{M}^2 - \bar{m}^2 \right),$$

where $(m, M)$ and $(\bar{m}, \bar{M})$ are the masses of initial and final mesons. We use capital $M$ for the masses of heavy–light mesons and small $m$ for the masses of the Goldstone bosons. The matrix of coefficients $C_{(I, S)}$ that characterize the interaction strength in a given channel is given in Table 2. The $s$-wave interaction kernels are identical for the two scattering problems considered here.

In contrast the loop functions, diagonal in the coupled-channel space, depend on whether to scatter Goldstone bosons off pseudo-scalar or vector heavy–light mesons,

$$J_{0^+}(\sqrt{s}) = I(\sqrt{s}) - I(\mu_{0^+}^{(I, S)}),$$

$$J_{1^+}(\sqrt{s}) = \left( 1 + \frac{p_{cm}^2}{2M^2} \right) I(\sqrt{s}) - I(\mu_{1^+}^{(I, S)}).$$

$$I(\sqrt{s}) = \frac{1}{16\pi^2} \left( \frac{p_{cm}}{\sqrt{s}} \ln \left( 1 - \frac{s - 2p_{cm}\sqrt{s}}{m^2 + M^2} \right) - \ln \left( 1 - \frac{s + 2p_{cm}\sqrt{s}}{m^2 + M^2} \right) \right) + \left( \frac{1}{2m^2 - M^2} - \frac{m^2 - M^2}{2s} \right) \times \ln \left( \frac{m^2}{M^2} \right) + 1 \right) + I(0),$$

where $\sqrt{s} = \sqrt{M^2 + p_{cm}^2 + \sqrt{m^2 + p_{cm}^2}}$. Note however that the two loop functions in (6) differ by a term suppressed with $1/M^2$ only. A crucial ingredient of the $\chi$-BS(3) scheme is its approximate crossing symmetry guaranteed by a proper choice of the subtraction scale $\mu_{j^P}^{(I, S)}$. Referring to the detailed discussions in \cite{1,19–21} we obtain

$$\mu_{0^+}^{(I, 0)} = M_{D(1867)}, \quad \mu_{0^+}^{(I, \pm 1)} = M_{D_s(1969)},$$

$$\mu_{1^+}^{(I, 0)} = M_{D(1008)}, \quad \mu_{1^+}^{(I, \pm 1)} = M_{D_s(2010)},$$

$$\mu_{1^+}^{(I, 2)} = M_{D(2008)},$$

With (4)–(7) the brief exposition of the $\chi$-BS(3) approach as applied to heavy–light meson resonances is completed.


3. Results

To study the formation of meson resonances we generate speed plots as suggested by Hohler [26]. The speed \( \text{Speed}_{ab}^{(I,S)}(\sqrt{s}) \) of a given channel \( ab \) is introduced by [26,27],

\[
\text{Speed}_{ab}^{(I,S)}(\sqrt{s}) = \sum_c \left[ \frac{d}{d\sqrt{s}} \delta_{cbc}^{(I,S)}(\sqrt{s}) \right] \times \left( \delta_{cb} + 2i\delta_{cb}^{(I,S)}(\sqrt{s}) \right)^T , \tag{8}
\]

where we give expressions valid in the \( 0^+ \) sector. A corresponding result for the \( 1^+ \) sector agrees with (8) up to a small correction term vanishing in the heavy-mass limit \( M \to \infty \) (see (6)).

In order to explore the SU(3) multiplet structure of the resonance states we first study the \( 0^+ \) sector in the ‘heavy’ SU(3) limit [1,21] with \( m_{\pi,K,\eta} = 500 \text{ MeV} \) and \( M_D = 1800 \text{ MeV} \). In this case we obtain an anti-triplet of mass 2204 MeV with poles in the \((0,+1),(1/2,0)\) amplitudes. The sextet channel does not show a bound state signal in this case. However if the attraction is increased slightly by using \( f = 80 \text{ MeV} \) rather than the canonical value 90 MeV poles at mass 2298 MeV arise in the \((1,+1),(1/2,0),(0,-1)\) amplitudes. This finding reflects that the Weinberg–Tomozawa interaction,

\[
\bar{3} \otimes 8 = \bar{3} \otimes 6 \oplus \bar{10} \tag{9}
\]

predicts attraction in the anti-triplet and sextet channel but repulsion for the anti-15-plet. In contrast performing the ‘light’ SU(3) limit [1,21] with \( m_{\pi,K,\eta} \sim 140 \text{ MeV} \) together with \( M_D = 1800 \text{ MeV} \) we do not find any signal of a resonance in any of the channels. Analogous results are found in the \( 1^+ \) sector. If we used identical masses for the \( 1^- \) and \( 0^- \) mesons the differences in the generated spectra are below 1 MeV.

In Figs. 1, 2 the spectrum as it arises with physical masses is shown. We predict a bound state of mass 2303 MeV in the \((0,1)\)-sector (see Fig. 1). According to [12,13] this state should be identified with a narrow resonance of mass 2317 MeV recently observed by the BABAR Collaboration [8]. Since we do not consider isospin violating processes like \( \eta \to \pi_0 \) the latter state is a true bound state in our present scheme. Given the fact that our computation is parameter-free this is a remarkable result. As demonstrated by the real part of the corresponding scattering amplitude of Fig. 1 the state couples dominantly to the \( DK \) channel. In the \((1,+,1)-speeds where we expect a signal from the sextet a strong cusp effect at the \( KD(1867) \)-threshold is seen. The large coupling constant to the \( \pi D_0(1669) \) channel leads to the broad structure seen in the figure. Fig. 2 illustrates that in the \((1/2,0)\)-sector we predict a narrow state of mass 2413 MeV just below the \( \eta D(1867) \)-threshold and a broad state of mass 2138 MeV. Modulo some mixing effects the heavier of the two is part of the sextet the lighter a member of the anti-triplet. The latter \((1/2,0)\)-state was expected to have a large branching ratio into the \( \pi D(1867) \)-channel [13,15]. This is confirmed by our analysis. Finally in the \((0,+,1)\)-speed a pronounced cusp effect at the \( \bar{K} D(1867) \)-threshold is seen.

The spectrum predicted for the \( 1^+ \) states is very similar to the spectrum of the \( 0^+ \) states. Figs. 1, 2 demonstrate that it is shifted up by approximatively
Fig. 1. Open-charm resonances with $J^P = 0^+, 1^+$ and $(I, S) = (0, 1), (1, 1)$ as seen in the scattering of Goldstone bosons of $D(1867), D_s(1969)$ and $D(2008), D_s(2110)$ mesons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude $f_{ab}$, with $t_{ab} = f_{ab} |P_{cm} P_{cm}^{\dagger}|^{1/2}$ (see (8)).
Fig. 2. Open-charm resonances with $J^P = 0^+, 1^+$ and $(J, S) = (\frac{1}{2}, 0), (0, -1)$ as seen in the scattering of Goldstone bosons of $D(1867), D_s(1969)$ and $D(2008), D_s(2110)$ mesons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, $f_{ab}$, with $s_{ab} = f_{ab} P_{cm} P_{cm}^{-1/2}$ (see (8)).
Fig. 3. Open-beauty resonances with $J^P = 0^+, 1^+$ and $(I,S) = (0,1), (1,1)$ as seen in the scattering of Goldstone bosons of $B(5279), B_s(5370)$ and $B(5325), B_s(5417)$ mesons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, $f_{ab}$, with $t_{ab} = f_{ab}(\sqrt{s_{cm}} p_{cm})^{1/2}$ (see (8)).
Fig. 4. Open-beauty resonances with $J^P = 0^+, 1^+$ and $(I,S) = (1/2, 0), (0, -1)$ as seen in the scattering of Goldstone bosons of $B(5279), B_s(5370)$ and $B(5325), B_s(5417)$ mesons. Shown are speed plots together with real and imaginary parts of reduced scattering amplitude, $f_{ab}$, with $l_{ab} = f_{ab} P_{cm}^{(a)} P_{cm}^{(b)}$ (see (8)).
140 MeV with respect to the $0^+$ spectrum. The bound state in the $(0,1)$-sector comes at 2440 MeV. Thus the mass splitting of the $1^+$ and $0^+$ states in this channel agrees very well with the empirical value of about 140 MeV measured by the BABAR and CLEO Collaborations [8,9]. A narrow structure at 2552 MeV is predicted in the $\bar{D}(2420)$-channel which should be identified with the $\bar{D}(2420)$-resonance [28]. Even though the resonance mass is overestimated by about 130 MeV our result is consistent with its small width of about 20 MeV. The triplet state in this sector of mass 2325 MeV has again a quite large width reflecting the strong coupling to the $\pi D(2008)$-channel. Finally we obtain strong cusp effects at the $\bar{K} D(2008)$- and $K D(2008)$-thresholds in the $(0,-1)$- and $(1,+1)$-sectors. It is interesting to speculate whether chiral corrections in the $(0,1)$-sector since due to heavy-quark symmetry chiral correction terms conspire to slightly increase the net attraction in these sectors. This would lead to a $(0,-1)$-bound state. The fact that we overestimate the mass of the sextet state $D(2420)$ by about 130 MeV we take as a strong indication that this should indeed be the case. An analogous statement holds for the $0^+$ sector since due to heavy-quark symmetry chiral correction effects in the $0^+$ and $1^+$ are identical at leading order.

Our predictions for the heavy–light resonance spectrum differ significantly from the analyses [12,13] that were based on the linear realization of the chiral SU(3) symmetry. Besides the additional sextet states predicted by the more general non-linear realization of the chiral SU(3) group, most notable are the differences in the open-beauty sector. Using 5279 MeV and 5370 MeV for the $0^-$ ground states [28] we obtain the following spectrum (see Figs. 3, 4). The $(0,1)$-bound state comes at 5643 MeV, a value about 70 MeV lower than predicted in [12,13]. In the $(1,1)$-channel we predict a broad resonance with mass of about 5750 MeV. In the $(1,0)$-channel a broad state of mass 5526 MeV and a narrow resonance of mass 5760 MeV and width of about 30 MeV is predicted. Most spectacular is a bound-state just below the $\bar{K}B$-threshold at 5761 MeV predicted in the $(0,-1)$ channel.

The $1^+$ spectrum we compute in terms of the ground state masses 5325 MeV and 5417 MeV [12, 28]. The results resemble the spectrum found in the $0^+$ sector. We predict the masses 5690 MeV (0, 1), 5790 MeV (1, 1), 5590 MeV $\left(\frac{1}{2}^-, 0\right)$, 5810 MeV $\left(\frac{1}{2}^+, 0\right)$, and 5807 MeV (0, −1).

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**References**

E. van Beveren, G. Rupp, hep-ph/0305035;
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S. Godfrey, hep-ph/0305122;
P. Colangelo, F. De Fazio, hep-ph/0305140;
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