

# Charmed meson resonances from chiral coupled-channel dynamics

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**Abstract.** Charmed meson resonances with quantum numbers  $J^P = 0^+$  and  $J^P = 1^+$  are generated in terms of chiral coupled-channel dynamics. At leading order in the chiral expansion a parameter-free prediction is obtained for the scattering of Goldstone bosons off charmed 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) group. In particular, so far unobserved narrow isospin-singlet states with negative strangeness are predicted.

Recently a new narrow state of mass 2.317 GeV that decays into  $D_s^+ \pi^0$  was announced [1]. This result was confirmed [2] and a second narrow state of mass 2.463 GeV decaying into  $D_s^* \pi^0$  was observed. Such states were first predicted in [3, 4] based on the spontaneous breaking of chiral symmetry. The theoretical predictions [3, 4, 5, 6] 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 and excludes any further model assumptions, no a priori prediction can be made for the existence of chiral partners of any given state.

Recently it was shown [7] that solving the coupled-channel Bethe-Salpeter equation with the interaction kernel following from a non-linear chiral SU(3) Lagrangian one is able to predict two octets and a singlet multiplets of the light  $1^+$  mesons consistent with the empirical spectrum. Similar results were obtained for light meson resonances with  $J^P = 0^+$  quantum numbers [8, 9, 10, 11, 12, 13]. In view of the evident success of the chiral coupled-channel dynamics to predict the existence of a wealth of meson and baryon resonances in the  $(u, d, s)$ -sector of QCD [7, 14, 15] it is expectable that the chiral SU(3) symmetry should also predict spectra of hadrons with open charm or beauty (see also [16]). In this talk we review the  $\chi$ -BS(3) approach [17, 18, 19, 20, 7, 14, 15] as applied to open-charm meson resonances [21, 22]. We will not have space to discuss further exciting results concerning open-beauty meson resonances [21], open-charm baryon resonances [23] or results in the  $(u, d, s)$ -sector of QCD [7, 14, 15].

Heavy-light meson states with quantum numbers  $J^P = 0^+$  and  $J^P = 1^+$  may be studied by considering the s-wave scattering of Goldstone bosons off the heavy-light ground state mesons with  $J^P = 0^-$  and  $J^P = 1^-$ . If scalar or axial vector resonances exist they should manifest themselves as poles in the corresponding scattering amplitudes. The starting point to describe low-energy scattering processes is the chiral SU(3) Lagrangian

[24, 25, 26, 27] including heavy-light  $0^-$  and  $1^-$  fields. A systematic approximation scheme arises due to a successful scale separation justifying chiral power counting rules [28]. Our effective field theory for the scattering of Goldstone bosons off any heavy field is based on the assumption that the scattering amplitudes are perturbative at subthreshold energies with the expansion parameter  $Q/\Lambda_\chi$ . The small scale  $Q$  is to be identified with any small momentum of the system. The chiral symmetry-breaking scale is

$$\Lambda_\chi \simeq 4\pi f \simeq 1.13 \text{ GeV},$$

with the parameter  $f \simeq 90 \text{ MeV}$  determined by the pion decay process. Once the available energy is sufficiently high to permit elastic two-body scattering a further typical dimensionless parameter  $m_K^2/(8\pi f^2) \sim 1$  arises [17, 18, 19] if strangeness is considered explicitly. This extra parameter invalidates any perturbative calculation within chiral SU(3) effective theory. Since this ratio is uniquely linked to two-particle reducible diagrams it is sufficient to sum those diagrams keeping the perturbative expansion of all irreducible diagrams, i.e. the coupled-channel Bethe-Salpeter equation has to be solved. This is the basis of the  $\chi$ -BS(3) approach developed in [17, 18, 19, 7].

We identify the leading-order Lagrangian density [24, 25, 26, 27] describing the interaction of Goldstone bosons with pseudo-scalar and vector mesons,

$$\begin{aligned} \mathcal{L}(x) = & \frac{1}{8f^2} \text{tr} \left[ \mathbf{P}(x) (\partial^\nu \mathbf{P}^\dagger(x)) - (\partial^\nu \mathbf{P}(x)) \mathbf{P}^\dagger(x) \right] [\Phi(x), (\partial_\nu \Phi(x))]_- \\ & - \frac{1}{8f^2} \text{tr} \left[ \mathbf{P}^\mu(x) (\partial^\nu \mathbf{P}_\mu^\dagger(x)) - (\partial^\nu \mathbf{P}^\mu(x)) \mathbf{P}_\mu^\dagger(x) \right] [\Phi(x), (\partial_\nu \Phi(x))]_-, \end{aligned} \quad (1)$$

where  $\Phi$  is the octet of Goldstone boson fields and  $P$  and  $P_\mu$  are the triplets of massive pseudo-scalar and vector-meson fields in the matrix representation.

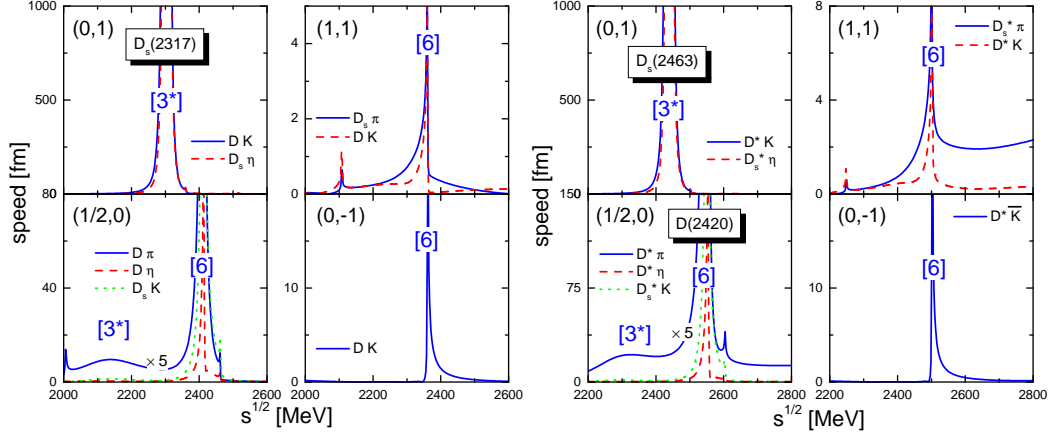
Within the  $\chi$ -BS(3) approach [19, 7] the s-wave scattering amplitudes,  $M_{JP}^{(I,S)}(\sqrt{s})$  take the simple form

$$M_{JP}^{(I,S)}(\sqrt{s}) = \left[ 1 - V^{(I,S)}(\sqrt{s}) J_{JP}^{(I,S)}(\sqrt{s}) \right]^{-1} V^{(I,S)}(\sqrt{s}). \quad (2)$$

The effective interaction kernel  $V^{(I,S)}(\sqrt{s})$  in (2) is determined by the leading order chiral SU(3) Lagrangian (1),

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

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, and the loop functions  $J_{JP}^{(I,S)}(\sqrt{s})$  are given [7]. As expected from heavy-quark symmetry the interaction kernels as well as the loop functions are identical for the  $0^-$  and  $1^-$  sectors in the limit  $M \rightarrow \infty$ .



**FIGURE 1.** Speed plots for heavy-light scalar (left panel) and axial-vector (right panel) mesons with isospin ( $I$ ) and strangeness ( $S$ ).

In order to guarantee the perturbative nature of the scattering amplitude at subthreshold energies the  $\chi$ -BS(3) approach insists on a renormalization condition of the form

$$M^{(I,S)}(\sqrt{s} = \mu^{(I,S)}) = V^{(I,S)}(\sqrt{s} = \mu^{(I,S)}) \quad (4)$$

with the natural subtraction scales

$$\begin{aligned} \mu_{0+}^{(I,0)} &= M_{D(1867)}, & \mu_{0+}^{(I,\pm 1)} &= M_{D_s(1969)}, & \mu_{0+}^{(I,2)} &= M_{D(1867)}, \\ \mu_{1+}^{(I,0)} &= M_{D(2008)}, & \mu_{1+}^{(I,\pm 1)} &= M_{D_s(2110)}, & \mu_{1+}^{(I,2)} &= M_{D(2008)}. \end{aligned} \quad (5)$$

A crucial ingredient of the  $\chi$ -BS(3) approach is a matching of s- and u-channel unitarized scattering amplitudes at subthreshold energies [19, 7]. This construction reflects our basic assumption that diagrams showing an s-channel or u-channel unitarity cut need to be summed to all orders at least at energies close to where the diagrams develop their imaginary part. By construction, a matched scattering amplitude satisfies crossing symmetry exactly at energies where the scattering process takes place. At subthreshold energies crossing symmetry is implemented approximatively only, however, to higher and higher accuracy when more chiral correction terms are considered. Insisting on the renormalization condition (4,5) guarantees that subthreshold amplitudes match smoothly and therefore the final 'matched' amplitudes comply with the crossing-symmetry constraint to high accuracy. A conceivable small variation of the subtraction scales around their natural values (5) has very little effect on the results. In fact chiral correction terms modify the effective interaction  $V(\sqrt{s})$  rather than giving rise to a modification of the subtraction scale [19, 7, 22]. Changing the optimal subtraction scale (5) would deteriorate the quality of the matching of u- and s-channel unitarized amplitudes [19, 7].

We turn to the results of the  $\chi$ -BS(3) approach for charmed mesons. It is instructive to explore first the SU(3) multiplet structure of the resonance states formed by the chiral coupled-channel dynamics. First the  $0^+$  sector is discussed in a 'heavy' SU(3) limit [15, 7] with  $m_{\pi,K,\eta} = 500$  MeV and  $M_D = 1800$  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  $f = 80$  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,

$$\bar{3} \otimes 8 = \bar{3} \oplus 6 \oplus \overline{15} \quad (6)$$

predicts attraction in the anti-triplet and sextet channel but repulsion for the anti-15-plet. In contrast performing a 'light' SU(3) limit [15, 7] with  $m_{\pi,K,\eta} \sim 140$  MeV together with  $M_D = 1800$  MeV we do not find any signal of a resonance in both anti-triplet and sextet channels. Analogous results are found in the  $1^+$  sector.

To study the formation of meson resonances we generate speed plots [29] as suggested by Höhler [30] (for definitions cf. [21]). Fig. 1 shows the spectra of  $0^+$  (left panel) and  $1^+$  (right panel) as they arise in calculations with physical masses. We predict a bound state of mass 2303 MeV in the  $(0, 1)$ -sector of  $0^+$  mesons (see Fig. 1, left panel). According to [5, 6] this state can be identified with a narrow resonance of mass 2317 MeV recently observed by the BABAR collaboration [1]. Since we do not consider isospin violating processes like  $\eta \rightarrow \pi_0$  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. 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  $\pi D_s(1969)$  channel leads to the broad structure seen in the figure. Fig. 1 (left panel) illustrates that in the  $(\frac{1}{2}, 0)$ -sector we predict a narrow  $0^+$  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  $(\frac{1}{2}, 0)$ -state was expected to have a large branching ratio into the  $\pi D(1867)$ -channel [16, 6]. 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. Fig. 1 (right panel) demonstrates that it is shifted up by approximately 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 [1, 2]. A narrow structure at 2552 MeV is predicted in the  $(\frac{1}{2}, 0)$ -channel which may be identified with the  $D(2420)$ -resonance [31]. 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  $KD(2008)$ -thresholds in the  $(0, -1)$ - and  $(1, +1)$ -sectors. It is interesting to speculate whether 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 prediction that this should indeed be the case [22]. 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.

To summarize: We presented a coupled-channel description of the meson-meson scattering in the open charm sector using the chiral SU(3) Lagrangian involving light-heavy  $J^P = 0^-$  and  $J^P = 1^-$  fields that transform non-linearly under the chiral SU(3) group. The major result of our study is the prediction of the charmed mesonic 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  $(I, S) = (0, -1)$ .

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