Open-charm meson resonances with negative strangeness

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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) Lagrangian. Adjusting the free parameters that arise at subleading order to reproduce the mass of the $D(2420)$ resonance as well as the new states established recently by the BABAR, CLEO and BELLE collaborations we obtain refined masses for the anti-triplet and sextet states. Bound states of antikaons at the $D(1867)$ and $D(2008)$ mesons are predicted at 2352 MeV ($J^P = 0^+$) and 2416 MeV ($J^P = 1^+$). In addition we anticipate a narrow scalar state of mass 2389 MeV with $(I, S) = (\frac{1}{2}, 0)$.

1 Introduction

In a recent work [1] it was demonstrated that chiral SU(3) symmetry predicts parameter-free heavy-light $J^P = 0^+$ and $J^P = 1^+$ meson resonances. The states form an anti-triplet and a sextet representation of the SU(3) flavor group. In the open-charm sector the recently discovered narrow state of mass 2.317 GeV [2] and 2.463 GeV [3] were recovered as part of the strongly bound anti-triplet states. The missing $(\frac{1}{2}, 0)$ triplet states were later announced by the BELLE collaboration [4] finding broad resonance structures at 2.308 GeV and 2.427 GeV in the $J^P = 0^+$ and $J^P = 1^+$ channels respectively. Such states were first predicted in [5,6] based on the spontaneous breaking of chiral symmetry and heavy-quark symmetry. The spectacular experimental discoveries triggered a flurry of theoretical publications [5,6,7,8,9].

The purpose of this Letter is to study the properties of the newly predicted sextet states [1] in the open charm sector as chiral correction terms are considered. Since contrary to the open bottom sector in which already the leading
order computation [1] predicts weakly bound $K \bar{B}$ isospin-zero states the binding in the $\bar{K} D$ system is not quite sufficiently strong to form a bound state at leading order. In [1] it was pointed out that chiral correction terms should provide additional attraction to form open charm bound states with negative strangeness also. The argument was based on the identification of the $D(2420)$ resonance as a member of a SU(3)-sextet. Our opinion differs here from the traditional approach [10,11,12,13] which interprets the latter state as part of an anti-triplet. Since the leading order computation underestimates the binding of the latter resonance by about 130 MeV correction terms are expected to provide additional binding for the sextet states in particular in the strangeness minus one sector.

In this work we apply the $\chi$-BS(3) approach developed originally for meson-baryon scattering [14,15,16,17,18,19] but recently also applied successfully to meson-meson scattering [1,20]. For earlier works on meson-meson scattering based on various schemes see [21,22,23,24,25,26]. Using the chiral SU(3) Lagrangian involving heavy-light $J^P = 0^-$ and $J^P = 1^-$ fields that transform non-linear under the chiral SU(3) group a coupled-channel computation of the meson-meson scattering amplitude in the open-charm sector is performed (see also [27]). Leading and subleading terms are considered in the chiral expansion of the interaction kernel. At subleading order three unknown parameters arise that can be adjusted to reproduce accurately the scalar or axial-vector spectrum. The values for the parameters determined independently in the scalar and axial-vector sector are reasonably close consistent with the expectation of heavy-quark symmetry. The central result of this work is the prediction of bound states of antikaons at the $D(1867)$ and $D(2008)$ mesons of mass 2352 MeV ($J^P = 0^+$) and 2416 MeV ($J^P = 1^+$). In addition we anticipate a narrow scalar state of mass 2389 MeV with $(I, S) = (\frac{1}{2}, 0)$.

2 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-orders Lagrangian density [28,29,30,31,32] describing the s-wave interaction of Goldstone bosons with pseudo-scalar and vector-mesons. In order to exploit the heavy-quark symmetry it is convenient to introduce a Dirac valued field [29,30,31,32,13]

$$H(x) = \frac{1}{2} (P^\mu(x) \gamma^\mu - \gamma_5 P(x)),$$

$$\bar{H}(x) = \gamma_0 H^\dagger(x) \gamma_0,$$  

which encodes the massive pseudo-scalar field, $P(x)$, and the vector-meson field, $P^\mu(x)$. We work directly with the relativistic chiral Lagrangian so that the field $H(x)$ carries dimension one and does not involve the 4-velocity $v^\mu$. 

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of the heavy-meson. Neglecting terms that are suppressed in the heavy-quark mass limit we write

\[
\mathcal{L}(x) = \frac{1}{2} \text{tr} \left( \bar{H}(x) \left( \partial^2 + \mathbf{M}^2 + 4 c_0 (\text{tr} \chi_0) - 4 c_1 \chi_0 \right) H(x) \right) \\
+ \frac{1}{8 f^2} \text{tr} \left( \left[ (\partial^\nu H(x)) \bar{H}(x) - H(x) (\partial^\nu \bar{H}(x)) \right] [\Phi(x), (\partial_\nu \Phi(x))] \right)_- \\
+ \frac{g M}{f} \text{tr} \left( \bar{H}(x) \gamma_5 \gamma^\mu (\partial_\mu \Phi(x)) H(x) \right) \\
- \frac{c_0}{f^2} \text{tr} \left( \Phi(x) \chi_0 \Phi(x) \right) \text{tr} \left( \bar{H}(x) H(x) \right) \\
+ \frac{c_1}{4 f^2} \text{tr} \left( \bar{H}(x) \left[ \Phi(x), [\Phi(x), \chi_0]_+ \right]_+ H(x) \right) \\
- \frac{c_2}{f^2} \text{tr} \left( \bar{H}(x) H(x) \right) \text{tr} \left( (\partial^\mu \Phi(x)) (\partial_\mu \Phi(x)) \right) \\
- \frac{c_3}{f^2} \text{tr} \left( \bar{H}(x) (\partial^\mu \Phi(x)) (\partial_\mu \Phi(x)) H(x) \right),
\]

\[
\chi_0 = \frac{1}{3} \left( m_\pi^2 + 2 m_K^2 \right) 1 + \frac{2}{\sqrt{3}} \left( m_\pi^2 - m_K^2 \right) \lambda_8,
\]

where \( \Phi \) is the Goldstone bosons field. The parameter \( f \) in (2) is known from the weak decay process of the pions. We use \( f = 90 \text{ MeV} \) through out this work. The parameter \( g \) can be adjusted the partial decay width \( D^+(2008) \to \pi^+ D(1867) \). Using the latest values [33] of \( 64 \pm 15 \text{ keV} \) one obtains \( g M \approx 1155 \pm 135 \text{ MeV} \) at tree-level. The symmetry breaking parameter \( c_1 \) can be determined by the mass splitting of the open charm pseudo-scalar or vector-meson ground states,

\[
M_{D_0^{(*)}}^2 - M_{D_0^0}^2 = 4 c_1 \left( m_K^2 - m_\pi^2 \right), \quad M_{D_1^{(*)}}^2 - M_{D_1^0}^2 = 4 c_1 \left( m_K^2 - m_\pi^2 \right), \quad \] (3)

which leads to \( c_1 \approx 0.44 \) and \( c_1 \approx 0.47 \) for the pseudo-scalar and vector states respectively. We use the averaged value \( c_1 \approx 0.45 \) in this work. The remaining three dimension less parameters \( c_0, c_2, c_3 \) will be determined in this work. Formally the latter parameters scale with \( c_i \sim M \).

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 \( D = (D^{(+)}, D^{(0)})^t \),

\[
\Phi = \tau \cdot \pi(140) + \alpha^\dagger \cdot K(494) + K^\dagger(494) \cdot \alpha + \eta(547) \lambda_8, \\
P = \frac{1}{\sqrt{2}} \alpha^\dagger \cdot D(1867) - \frac{1}{\sqrt{2}} D^t(1867) \cdot \alpha + i \tau_2 D^{(s)}(1969), \\
P_\mu = \frac{1}{\sqrt{2}} \alpha^\dagger \cdot D_\mu(2008) - \frac{1}{\sqrt{2}} D^t_\mu(2008) \cdot \alpha + i \tau_2 D^{(s)}_\mu(2110), \\
\alpha^\dagger = \frac{1}{\sqrt{2}} (\lambda_4 + i \lambda_5, \lambda_6 + i \lambda_7), \quad \tau = (\lambda_1, \lambda_2, \lambda_3),
\]

(4)
Table 1
Coupled-channel states with \((I, S)\).

\[
\begin{array}{c|ccc}
\left(\frac{1}{2}, +2\right) & (0, +1) & (1, +1) & \left(\frac{1}{2}, 0\right) \\
(D_s K) & \left(\frac{1}{\sqrt{2}} D^t i \sigma_2 K\right) & (D_s \pi) & \left(\frac{1}{\sqrt{3}} \pi \cdot \sigma D\right) \\
(D_s \eta) & (D_s D^t i \sigma_2 \sigma K) & (\eta D) & (D_s i \sigma_2 \bar{K}^t) \\
\left(\frac{3}{2}, 0\right) & (0, -1) & (1, -1) & (\frac{1}{\sqrt{2}} \bar{K} D) & (\frac{1}{\sqrt{2}} \bar{K} \sigma D) \\
(\pi \cdot T D) & (\frac{1}{\sqrt{2}} \bar{K} D) & (\frac{1}{\sqrt{2}} \bar{K} \sigma D) & \\
\end{array}
\]

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. Even though we did not write down the relevant term in (2) describing the mass splitting of the \(0^-\) and \(1^-\) states we will use physical masses as given in (4) throughout this work.

The scattering problem decouples into seven orthogonal sectors specified by isospin \((I)\) and strangeness \((S)\) quantum numbers,

\[
(I, S) = \left(\left(\frac{1}{2}, +2\right), (0, +1), (1, +1), \left(\frac{1}{2}, 0\right), (\frac{3}{2}, 0), (0, -1), (1, -1)\right) .
\]

(5)

In Tab. 1 the channels that contribute in a given sector \((I, S)\) are listed. Heavy-light meson resonances with quantum numbers \(J^P = 0^+\) and \(J^P = 1^+\) manifest themselves as poles in the \(s\)-wave scattering amplitudes, \(M^{(I,S)}_{jp}(\sqrt{s})\), which in the \(\chi-\text{BS}(3)\) approach \([17,20]\) take the simple form

\[
M^{(I,S)}_{jp}(\sqrt{s}) = \left[1 - V^{(I,S)}_{jp}(\sqrt{s}) J^{(I,S)}_{jp}(\sqrt{s})\right]^{-1} V^{(I,S)}(\sqrt{s}) .
\]

(6)

We first specify the contributions to the effective interaction kernel \(V^{(I,S)}(\sqrt{s})\) in (6) from the Weinberg-Tomozawa term and the subleading interaction terms proportional to the parameters \(c_i\) as introduced in (2),

\[
V^{(I,S)}(\sqrt{s}) = \frac{C^{(I,S)}}{8 f^2} \left(3 s - M^2 - \bar{M}^2 - m^2 - \bar{m}^2 \right)
- \frac{M^2 - m^2}{s} (M^2 - \bar{m}^2) \\
+ 2 \frac{m_\pi^2}{f^2} \left( c_0 C^{(I,S)}_{\pi,0} + c_1 C^{(I,S)}_{\pi,1} \right) + 2 \frac{m_K^2}{f^2} \left( c_0 C^{(I,S)}_{K,0} + c_1 C^{(I,S)}_{K,1} \right) \\
+ \left( c_2 \frac{C^{(I,S)}_2}{s f^2} + c_3 \frac{C^{(I,S)}_3}{s f^2} \right) \left( s - \bar{M}^2 + \bar{m}^2 \right) \left( s - M^2 + m^2 \right) ,
\]

(7)
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 \(s\)-wave interaction kernels are identical for the two scattering problems considered here. Following [20] we neglect here the mixing of the \(s\)- with a \(d\)-wave channel in the \(1^+\) sector. Such effects are largely suppressed. We continue with a presentation of the contributions of the \(s\)- and \(u\)-channel exchange terms that are proportional to \(g^2\). In this case one has to discriminate the \(0^+\) and \(1^+\) channels,

\[
V_{0^+}^{(I,S)}(\sqrt{s}) = g^2 \frac{C_{s-ch}^{(I,S)}}{4 s f^2} \left( s - \bar{M}^2 + \bar{m}^2 \right) \left( s - M^2 + m^2 \right) \\
+ g^2 \frac{C_{u-ch}^{(I,S)}}{f^2} \int_{-1}^1 \frac{dx}{2} \left( \bar{q} \cdot q \right) \mu^2 - \left( \bar{m}^2 - \bar{q} \cdot p \right) \left( m^2 - p \cdot q \right),
\]

\[
V_{1^+}^{(I,S)}(\sqrt{s}) = -g^2 \frac{C_{u-ch}^{(I,S)}}{2 f^2} \int_{-1}^1 \frac{dx}{2} \frac{\mu^2 \bar{p} cm p cm x (1 - x^2)}{M^2 + \bar{M}^2 - \mu^2 - s + 2 \bar{q} \cdot q},
\]

\[
\bar{q} \cdot q = \sqrt{\bar{m}^2 + \bar{p}^2 cm} \sqrt{m^2 + p^2 cm - \bar{p} cm p cm x},
\]

\[
\bar{q} \cdot p = -\bar{q} \cdot q + \frac{s - \bar{M}^2 + \bar{m}^2}{2}, \quad \bar{p} \cdot q = -\bar{q} \cdot q + \frac{s - M^2 + m^2}{2},
\]

\[
\sqrt{s} = \sqrt{M^2 + \bar{p}^2 cm + \sqrt{m^2 + p^2 cm}} = \sqrt{M^2 + \bar{p}^2 cm + \sqrt{\bar{m}^2 + \bar{p}^2 cm}},
\]

where \(\mu\) is the appropriate mass of the heavy-meson exchanged in the \(u\)-channel. Note that we identified \(\bar{M} \simeq \mu\) in (8). As is evident from the representation (8) the contribution of the \(s\)- and \(u\)-channel terms scale with \(\sim (\bar{M})^0\) as compared to the linear scaling \(\sim (\bar{M})^1\) of the terms in (7), which follows from \(c_i \sim \bar{M}\). Therefore we expect the contributions (8) to be of subsubleading importance. The matrix of coefficients \(C^{(I,S)}\) that characterize the interaction strength in any given channel are presented in Tab. 2.

We turn to the loop functions introduced in (6). The latter are diagonal in the coupled-channel space and depend on whether to scatter Goldstone bosons off pseudo-scalar or vector heavy-light mesons,

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

\[
J_{1^+}(\sqrt{s}) = \left( 1 + \frac{p_{cm}^2}{3 M^2} \right) \left( I(\sqrt{s}) - I(\mu_{1^+}^{(I,S)}) \right),
\]

\[
I(\sqrt{s}) = \frac{1}{16 \pi^2} \left( \frac{p_{cm}}{\sqrt{s}} \ln \left( 1 - \frac{s - 2 p_{cm} \sqrt{s}}{m^2 + M^2} \right) - \ln \left( 1 - \frac{s + 2 p_{cm} \sqrt{s}}{m^2 + M^2} \right) \right) \\
+ \left( \frac{1}{2} \frac{m^2 + M^2}{m^2 - M^2} - \frac{m^2 - M^2}{2 s} \right) \ln \left( \frac{m^2}{M^2} \right) + 1 + I(0). \tag{9}
\]
The coefficients $C^{(I,S)}$ that characterize the interaction of Goldstone bosons with heavy-meson fields $H$ as introduced in (7,8) for given isospin ($I$) and strangeness ($S$). The channel ordering is specified in Tab. 1.

As expected from heavy-quark symmetry the two loop functions $J_{0^+}(\sqrt{s})$ and $J_{1^+}(\sqrt{s})$ in (9) 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 [17,18,19,20] we obtain

$$
\begin{align*}
\mu_{0^+}^{(I,0)} &= M_{D(1867)}, & \mu_{0^+}^{(I,\pm 1)} &= M_{D_s(1969)}, & \mu_{0^+}^{(I,2)} &= M_{D(1867)}, \\
\mu_{1^+}^{(I,0)} &= M_{D(2008)}, & \mu_{1^+}^{(I,\pm 1)} &= M_{D_s(2110)}, & \mu_{1^+}^{(I,2)} &= M_{D(2008)}.
\end{align*}
\tag{10}
$$

With (6,7,8,9,10) the brief exposition of the $\chi$–BS(3) approach as applied to heavy-light meson resonances is completed.
Table 3
Coupling constants $c_i$ (see (2)) for scalar and axial-vector open-charm mesons. The resonance coupling constants $g_i$ are defined in (11) with respect to channel labelling of Tab. 1.

3 Results

We begin with a presentation of results for the open-charm axial-vector mesons. At leading order [1] with $c_i = 0$ and $g = 0$ using physical masses for the intermediate states a narrow $(0, 1)$-state at 2440 MeV is generated by coupled-channel dynamics. Furthermore in the $(\frac{1}{2}, 0)$-sector a narrow and a broad state at 2552 MeV and 2300 MeV are predicted. The leading order results are already rather close to the empirical values of $2463 \pm 2$ MeV [3] for the $(0, 1)$-state and $2421.1 \pm 2.7$ MeV and $2427 \pm 61$ MeV with widths $23.7 \pm 6.9$ MeV and $384 \pm 94$ MeV [4] for the $(\frac{1}{2}, 0)$-states. It is straightforward to find the optimal values for the three unknown parameters $c_0$, $c_2$ and $c_3$ that arise at subleading order as to reproduce the empirical pattern quite accurately. In Tab. 3 the resulting parameters are given. To be precise we use the averaged mass $\mu = 1918$ MeV in (8). The typical size of these parameters appears consistent with the naturalness assumption $c_i \sim \mathcal{M}/m_\rho$ suggesting that the chiral expansion is well converging. The results do not sensitively depend on the value of $g$ within the range $0.4 < g < 0.7$. The empirical masses of all three states can be reproduced within experimental errors. The chiral corrections terms pull the narrow $(\frac{1}{2}, 0)$-state down and the broad $(\frac{1}{2}, 0)$-state up close to their empirical masses (see Tab. 3). In Fig. 1 we confront the spectral distribution of the $\pi D(2008)$-channel measured recently by the BELLE collaboration [4] with the theoretical distribution obtained from the imaginary part of the $\pi D(2008)$-scattering amplitude in the $(\frac{1}{2}, 0)$-sector. Within empirical errors we are able to reproduce the shape of the distribution. The small contribution of a state with $J = 2$ shown in the figure by the histograms [4] is not considered in this work. In Fig. 2 the scattering amplitude of the $(\frac{1}{2}, 0)$-sector are shown. The figure demonstrates that the $D(2420)$ resonance couples dominantly to the closed $\eta D(2008)$- and $K D_s(2110)$-channels which are the driving chan-
channels for the dynamic generation of the latter resonance. Having adjusted all parameters at subleading order we can now turn to the $(0, -1)$- and $(1, 1)$-sectors unconstrained yet by data. We find a $(1, 1)$-resonance at 2445 MeV of width 70 MeV (see Fig. 2) that couples dominantly to the $K D(2008)$-channel with $|g_2| = 3.0$ and a narrow $(0, -1)$-state at 2416 MeV. As anticipated in [1] the chiral correction terms increase the amount of attraction in the sextet channel predicting the existence of a $\bar{K} D(2008)$-bound state.

We analyze the properties of the anti-triplet and sextet states in more detail by extracting coupling constants. If in a given coupled-channel scattering amplitude $M_{ab}(\sqrt{s})$ a bound or resonance state of mass $M_R$ is present we determine the coupling constants $g_a$ of that state to the channel $a$ by the condition

$$M_{ab}(\sqrt{s}) \simeq -\frac{g_a^* g_b M_R}{\sqrt{s} - M_R + i \Gamma/2}, \quad (11)$$
Fig. 2. Open charm resonances with $J^P = 1^+$ and $(I, S) = (\frac{1}{2}, 0), (0, 1)$ as seen in the scattering of Goldstone bosons of $D(2008)$- and $D_s(2110)$-mesons. Shown are real and imaginary parts of reduced scattering amplitude, $f_{aa}(\sqrt{s}) = M_{aa}(\sqrt{s})/(8\pi\sqrt{s})$.

close to the pole. The resulting parameters are collected in Tab. 3. Whereas the chiral correction terms affect the mass of the $(0, 1)$-state by 24 MeV only, the change in the coupling constant is more important. At leading order with $c_i = 0$ and $g = 0$ one finds lower values of $|g_1| = 3.3$ and $|g_2| = 2.0$ as compared to the values 4.2 and 5.3 given in Tab. 3.

We turn to our results for the scalar open-charm mesons. Relying on the heavy-quark symmetry we could use the values for $c_i$ obtained by reproducing the properties of the axial-vector mesons. However, since we will make a prediction for the mass of the scalar $(0, 1)$-bound state a more accurate result is expected if we adjust the parameters $c_i$ in the scalar sector. Therewith sub-leading interaction terms not displayed in (2) which lead to the independence of the parameters in the scalar and axial-vector sectors are probed. Of course, consistency requires that the parameters should turn out not too different. We adjust the parameters to reproduce the narrow $(0, 1)$-state at $2317 \pm 3$ MeV.
Fig. 3. Open-charm resonances with $J^P = 0^+$ and $(I, S) = (\frac{1}{2}, 0), (0, 1)$ as seen in the scattering of Goldstone bosons of $D(1867)$- and $D_s(1969)$-mesons. Shown are real and imaginary parts of reduced scattering amplitude, $f_{aa}(\sqrt{\sigma}) = M_{aa}(\sqrt{\sigma})/(8\pi\sqrt{\sigma})$.

[2] and the $(\frac{1}{2}, 0)$-state at $2308 \pm 60 \text{ MeV}$ of width $276 \pm 99 \text{ MeV}$. The fit is unique since besides the $(\frac{1}{2}, 0)$-state we include also the empirical $\pi D(1867)$-spectrum obtained recently by the BELLE collaboration [4]. The fact that the latter spectrum does not show an additional peak implies that the sextet state couples only very weakly to the $\pi D(1867)$-channel. As shown in Fig. 1 it is possible to adjust the parameters such that the $\pi D(1867)$-spectrum decouples from the sextet state. The theoretical spectrum is again compatible with the empirical one. Note that the figure includes the contribution of a state with $J = 2$ and mass $2461.6 \pm 5.9 \text{ MeV}$ shown by the dashed histograms [4] but not considered in this work. Complementary is Fig. 3 which shows the various scattering amplitude in the $(\frac{1}{2}, 0)$-sector. The sextet state couples dominantly to the closed $\eta D(1867)$- and $K D_s(1969)$-channels. This result is analogous to the corresponding result for the axial-vector state analyzed in Fig. 2 only that in the scalar sector the coupling constant of the sextet state to the $\pi D(1867)$ is even further suppressed. The optimal set of parameters obtained are col-
lected in Tab. 3, which also includes the resonance parameters. The chiral correction terms push up the \((\frac{1}{2}, 0)\)-state by 15 MeV to its empirical value. Its coupling constants increase somewhat as compared to the leading order values \(|g_1| = 3.3\) and \(|g_2| = 2.0\). The values obtained for the coupling constants \(c_{0,2,3}\) are reasonably close to the values obtained in the axial-vector sector. To be precise we should mention that we used the averaged mass \(\mu = 2059\) MeV for the s- and u-channel contributions as specified in (8).

Having fixed all parameters we discuss the predictions for the scalar mesons so far not observed. The sextet state in the \((\frac{1}{2}, 0)\)-sector at 2389 MeV is quite narrow with a width of below 1 MeV since its coupling to the \(\pi D(1867)\) is much suppressed. We emphasize that it is difficult to make a precise prediction for the width of that state. If we only slightly change the set of parameters the sextet state shows up as a narrow peak in the \(\pi D(1867)\)-spectrum of Fig. 3. Depending on the precise values of the parameters the sextet state may be detected most efficiently via its coupling to the \(\eta D(1867)\) channel utilizing the \(\eta - \pi_0\)-mixing effect. We do not find a clear signal for the sextet state in the \((1, 1)\)-sector reflecting the fact that the amount of attraction for the sextet state is weaker in the scalar sector as compared to the attraction in axial-vector sector. However, a \((0, -1)\)-bound state of mass 2353 MeV is predicted.

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