Open-charm meson systems in the hadrogenesis conjecture

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Abstract

We study the spectrum and isospin-violating strong decays of charmed mesons with strangeness. The scalar $\mathrm{D}_{s0}^*(2317)^\pm$ and the axial vector $\mathrm{D}_{s1}^*(2460)^\pm$ states are generated by coupled-channel dynamics based on the leading order chiral Lagrangian. The effect of chiral corrections is investigated. We show that taking into account large-N_c relations implies a measurable signal for an exotic axial vector state in the ηD^* invariant mass distribution. The hadronic decay widths of the scalar $\mathrm{D}_{s0}^*(2317)^\pm$ and the axial vector $\mathrm{D}_{s1}^*(2460)^\pm$ are predicted to be 140 keV.

1 Introduction

The observation [1, 2] of two narrow, positive parity, strange charmed mesons at masses lower than expected in quark models [3, 4] may provide new insight into the way hadrons are generated. The properties of these two mesons, the scalar $\mathrm{D}_{s0}^*(2317)^\pm$ and the axial vector $\mathrm{D}_{s1}^*(2460)^\pm$, appear indeed sensitive to strong interaction symmetries as well as to the degrees of freedom building up hadronic excitations [5, 6, 7].

In this talk we discuss some aspects of a recent work [8] in which we computed the hadronic decay widths of the $D_{s0}^*(2317)$ and $D_{s1}^*(2460)$ states in the hadrogenesis conjecture, but also explored further the possibility of the formation of exotic multiplets. Earlier works [9, 10, 11] showed that such states can be generated dynamically by coupled-channel interactions based on the leading order chiral interaction of the Goldstone bosons with the open-charm pseudoscalar and vector ground states. This approach for heavy-light mesons exploits both heavy-quark and spontaneously-broken chiral symmetries.

A striking prediction of the studies [10, 11] is the existence of an exotic multiplet implied by the same coupled-channel dynamics that reproduces the scalar $\mathrm{D}_{s0}^*(2317)^\pm$ and the axial vector $\mathrm{D}_{s1}^*(2460)^\pm$ mesons. The existence of such states is amenable to experimental investigation. Chiral dynamics together with heavy-quark symmetry and large-N_c considerations do indeed predict a measurable exotic signal in the ηD^* invariant mass distribution [8]. Such effects are not expected in approaches where the constraints of the chiral symmetry of QCD and of the heavy-quark mass limit are not implemented [12, 13].

2 Open-charm states from chiral dynamics at leading order

We consider first the scalar $D_{s0}^*(2317)^{\pm}$ state. It is dynamically generated as a direct consequence of the leading order chiral interaction [10, 11]. We recall the relevant terms of the chiral Lagrangian density at leading order:

$$\begin{aligned}\mathcal{L} &= \frac{1}{4}\,\mathrm{tr}\,(\partial_\mu\Phi)\,(\partial^\mu\Phi) - \frac{1}{4}\,\mathrm{tr}\,\chi_0\,\Phi^2 + \;(\partial_\mu \mathrm{D})\,(\partial^\mu\bar{\mathrm{D}}) - \mathrm{D}\,\mathrm{M}_{0^-}^2\,\bar{\mathrm{D}} \\ &+ \frac{1}{8\,f^2}\Big\{(\partial^\mu D)\;[\Phi,(\partial_\mu\Phi)]_-\,\bar{D} - D\;[\Phi,(\partial_\mu\Phi)]_-\,(\partial^\mu\bar{D})\Big\}\,, \end{aligned} \tag{1}$$

where Φ and D are pseudoscalar octet and triplet fields. We use the notation $\bar{D} = D^\dagger$. In the particle representation the Goldstone and ground state open-charm meson fields are

$$\Phi = \begin{pmatrix} \pi^0 + \frac{1}{\sqrt{3}}\,\eta & \sqrt{2}\,\pi^+ & \sqrt{2}\,K^+ \\ \sqrt{2}\,\pi^- & -\pi^0 + \frac{1}{\sqrt{3}}\,\eta & \sqrt{2}\,K^0 \\ \sqrt{2}\,K^- & \sqrt{2}\,\bar{K}^0 & -\frac{2}{\sqrt{3}}\,\eta \end{pmatrix}\,, \qquad D = (D^0, -D^+, D_s^+)\,. \tag{2}$$

The Weinberg-Tomozawa term in (1), which is proportional to f^{-2}, is obtained by chirally gauging the kinetic term of the D-mesons. The parameter $f \simeq f_\pi = 92.4$ MeV is approximatively known from the weak decay of the charged pions. It predicts the leading s-wave interaction of the Goldstone bosons with the open-charm meson fields. A precise determination of f requires a chiral SU(3) extrapolation of some data set. In [14] the value $f \simeq 90$ MeV was obtained from a detailed study of pion- and kaon-nucleon scattering data. It is a universal and flavour-independent constant, defining the scale of chiral symmetry breaking. We use the value $f = 90$ MeV throughout this work.

The ground-state D-meson mass matrix is denoted by M_{0^-}. The mass term of the Goldstone bosons is proportional to the quark-mass matrix

$$\chi_0 = 2\,B_0 \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_d & 0 \\ 0 & 0 & m_s \end{pmatrix} = \frac{1}{3}\left(m_\pi^2 + 2\,m_K^2\right)\,1 + \frac{2}{\sqrt{3}}\,\left(m_\pi^2 - m_K^2\right)\lambda_8\,. \tag{3}$$

At leading order the latter can be expressed in terms of the pion and kaon masses as indicated in (3).

If we admit isospin breaking effects, i.e. $m_u \neq m_d$, there is a term in (1) proportional to $(m_u - m_d)\,\pi^0\,\eta$, inducing $\pi^0\,\eta$ mixing. A unitary transformation is required such that the transformed fields $\tilde{\pi}^0$ and $\tilde{\eta}$ with

$$\pi^0 = \tilde{\pi}^0\,\cos\epsilon - \tilde{\eta}\,\sin\epsilon\,, \qquad \eta = \tilde{\pi}^0\,\sin\epsilon + \tilde{\eta}\,\cos\,\epsilon\,, \tag{4}$$

decouple. According to [15] the ratio of quark masses relevant for the determination of the mixing angle takes the value

$$\frac{2}{\sqrt{3}}\,\frac{\sin(2\,\epsilon)}{\cos(2\,\epsilon)} = \frac{m_d - m_u}{m_s - (m_u + m_d)/2} = \frac{1}{43.7 \pm 2.7}\,, \tag{5}$$

which implies $\epsilon = 0.010 \pm 0.001$.

Heavy-light meson resonances with quantum numbers $J^P = 0^+$ manifest themselves as poles in the s-wave scattering amplitude. We consider the four isospin states $\langle K\,D, I|$, $\langle \pi\,D_s, 1|$ and $\langle \eta\,D_s, 0|$. In the presence of isospin mixing all channels couple. The mixing of the two isospin sectors is of order ϵ. We introduce the four states as follows

$$\begin{aligned} \langle 1| &= \langle\,\tilde{\pi}^0\,D_s^+| = \cos\epsilon\,\langle \pi^0\,D_s, 1| + \sin\epsilon\,\langle \eta\,D_s, 0|\,, \\ \langle 2| &= \langle\,\tilde{\eta}\;\;D_s^+| = \cos\epsilon\,\langle \eta\,D_s, 0| - \sin\epsilon\,\langle \pi^0\,D_s, 1|\,, \\ \langle 3| &= \langle K^0 D^+| = +\tfrac{1}{\sqrt{2}}\,(\langle K\,D, 0| - \langle K\,D, 1|)\,, \\ \langle 4| &= \langle K^+ D^0| = -\tfrac{1}{\sqrt{2}}\,(\langle K\,D, 0| + \langle K\,D, 1|)\,, \end{aligned} \tag{6}$$

where we use the phase convention introduced for the isospin states in [10, 11]. The Weinberg-Tomozawa interaction (1) implies an s-wave scattering amplitude of the simple form [10, 11]

$$M^{(0^+)}(s) = \Big[1 - V^{(0^+)}(s)\, J^{(0^+)}(s)\Big]^{-1} V^{(0^+)}(s)\,. \tag{7}$$

The matrix of loop functions, $J^{(0+)}(s)$, is diagonal and given by [10, 11]

$$\begin{aligned} J^{(0^+)}(s) &= I(s) - I(\mu_M^2)\,, \\ I(s) &= \frac{1}{16\,\pi^2}\left(\frac{p_{cm}}{\sqrt{s}}\left(\ln\left(1-\frac{s-2\,p_{cm}\sqrt{s}}{m^2+M^2}\right) - \ln\left(1-\frac{s+2\,p_{cm}\sqrt{s}}{m^2+M^2}\right)\right)\right. \\ &\quad \left. + \left(\frac{1}{2}\frac{m^2+M^2}{m^2-M^2} - \frac{m^2-M^2}{2\,s}\right)\ln\left(\frac{m^2}{M^2}\right) + 1\right) + I(0)\,, \end{aligned} \tag{8}$$

where $\sqrt{s} = \sqrt{M^2+p_{cm}^2} + \sqrt{m^2+p_{cm}^2}$. Each diagonal element depends on the masses of the light (Goldstone) and heavy (open charm) mesons m and M which constitute the associated coupled-channel state. The matching scale μ_M in (8) should be identified with the ground-state mass of the D_s meson, i.e. $\mu_M \simeq 1968$ MeV. In this case s- and u-channel unitarized scattering amplitudes may be smoothly matched around the matching scale as to define a full scattering amplitude that is crossing symmetric by construction [9, 10].

One may vary the matching scale slightly around its natural value. As shown in [9] the resulting effects are small for reasonable variations. A large variation is excluded since it would make the matching of u- and s-channel unitarized amplitudes possible only at the price of introducing a strong discontinuity, which is at odds with causality. The crucial issue is to identify at which order in a given scheme such a residual variation is justified. If one expands the coupled-channel amplitude in strict chiral perturbation theory the residual dependence on μ_M can be absorbed unambiguously into local counter terms of order Q_χ^3. This is to be contrasted with the correction terms to the Weinberg-Tomozawa Lagrangian (1) which start at order Q_χ^2. We argue therefore that it is reasonable to vary the matching scale μ_M only if the effective interaction kernel is evaluated to order Q_χ^3. In this case however, the amplitude should show an even smaller sensitivity on μ_M since its effect can be balanced by readjusting the Q_χ^3 counter terms.

At leading order the coupled-channel interaction kernel $V_{ij}^{(0^+)}(s)$ in (7) is determined by the chiral SU(3) Lagrangian (1),

$$V^{(0^+)}_{WT,ij}(s) = \frac{C_{ij}}{8\,f^2}\left(3\,s - M^2 - \bar{M}^2 - m^2 - \bar{m}^2 - \frac{M^2-m^2}{s}\,(\bar{M}^2 - \bar{m}^2)\right), \tag{9}$$

where (m, M) and $(\bar{m}, \bar{M})$ are the masses of the initial and final mesons. The matrix C_{ij} is specified in [8].

Given the coupled-channel scattering amplitude (7) with the effective interaction (9) it is straightforward to determine the mass and width of possible resonances. The Weinberg-Tomozawa interaction is strongly attractive in the isospin strangeness $(I, S) = (0, 1)$ sector, leading to the formation of a scalar resonance of mass M_{0^+} identified with the $D^*_{s0}(2317)$ state. The latter manifests itself as a pole in the scattering amplitude. The resonance properties can be determined from the factorized form of this amplitude close to the pole.

We consider now the axial vector $\mathrm{D}^*_{s1}(2460)^\pm$ state. It is important to study scalar and axial vector mesons in the open-charm sector on equal footing. The properties of spin 0 and spin 1 mesons are indeed closely related by the heavy-quark symmetry of QCD [16, 17]. We represent massive vector fields by antisymmetric tensors $D_{\mu\nu} = -D_{\nu\mu}$ and $\bar{D}_{\mu\nu} = D^\dagger_{\mu\nu}$ with $D_{\mu\nu} = (D^0_{\mu\nu}, -D^+_{\mu\nu}, D^+_{s,\mu\nu})$ describing the (vector) heavy-quark multiplet partners of the field D introduced in (2). Consider the Lagrangian density,

$$\mathcal{L} = -(\partial_\mu D^{\mu\alpha})\,(\partial^\nu \bar D_{\nu\alpha}) + \frac{1}{2}\, D^{\mu\alpha}\, M^2_{1^-}\, \bar D_{\mu\alpha} - \frac{1}{8\,f^2}\Big\{(\partial^\nu D_{\nu\alpha})\,[\Phi,(\partial_\mu\Phi)]_-\,\bar D^{\mu\alpha} - D_{\nu\alpha}\,[\Phi,(\partial_\nu\Phi)]_-\,(\partial_\mu \bar D^{\mu\alpha})\Big\}\,, \tag{10}$$

involving the kinetic term and its associated Weinberg-Tomozawa interaction. The sum of the Lagrangian densities (1) and (10) respects the heavy-quark mass limit of QCD if the mass parameters M_{0^-} and M_{1^-} are assumed to be degenerate in that limit.

The Lagrangian (10) predicts the s-wave interaction of the Goldstone bosons with the open-charm vector states unambiguously in terms of the known parameter $f = 90$ MeV. A corresponding s-wave partial-wave amplitude $M^{(1^+)}(s)$ is computed in terms of the effective interaction $V^{(1^+)}(s)$ and loop function $J^{(1^+)}(s)$, with

$$M^{(1^+)}(s) = \Big[1 - V^{(1^+)}(s)\, J^{(1^+)}(s)\Big]^{-1} V^{(1^+)}(s)\,,$$
$$J^{(1^+)}(s) = \Big(\frac{3}{2}\,M^2 + \frac{p^2_{cm}}{2}\Big)\Big\{I(s) - I(\mu_M^2)\Big\}\,, \tag{11}$$

where the universal integral $I(s)$ was introduced in (8). The specification of the matching scale μ_M was discussed in [10]. We identify $\mu_M = 2012$ MeV with the mass of the vector-meson ground state.

The effective interaction $V^{(1+)}_{ij}(s)$ implied by (10) is derived using a convention for the coupled-channel states analogous to (6). A straightforward application of [9] leads to the result

$$V^{(1+)}_{WT,ij}(s) = \frac{\bar M^2 + M^2}{3\,\bar M^2\, M^2}\, V^{(0^+)}_{WT,ij}(s) - \frac{(\bar M^2 - M^2)}{12\,f^2\,\bar M^2\, M^2}\,(\bar m^2 - m^2)\, C_{ij}\,, \tag{12}$$

in terms of the matrix $V^{(0^+)}_{WT,ij}(s)$ and the C_{ij} coefficients already encountered in (9). In (12) the parameters M and $\bar M$ denote the masses of open-charm vector meson of the initial and final state respectively. Like in (9) the masses m and $\bar m$ stand for the masses of the initial and final Goldstone bosons.

We point out that in the particular limit $\bar M = M$ we recover the expressions obtained before in [10] using the vector field representation of the spin 1 states, i.e. the invariant amplitude $M^{(1+)}(s)$ is identical to that of [10] within a normalization factor $2\,M^2/3$. This observation implies in particular that the predictions for the axial-vector spectrum are consistent with the expectation of the heavy-quark symmetry, as emphasized in [10]. The $D^*_{s1}(2460)$ state is generated dynamically. The scattering amplitude develops a pole at $s = M^2_{1^+}$.

3 Chiral correction terms

Following [11] we construct chiral correction terms to the leading order interactions (1) and (10). We take into account the chiral correction terms to the effective interactions $V^{(0^+)}(s)$ and $V^{(1^+)}(s)$ relevant at chiral order Q^2_χ. There will be two types of contributions for s-wave scattering. On the one hand we include s- and u-channel exchanges of the D-meson ground states based on the leading order vertices involving a Goldstone boson and two D-mesons. On the other hand local 2-body counter terms (breaking chiral symmetry and chiral symmetric respectively) will be constructed.

We identify the leading order 3-point vertices involving the Goldstone bosons

$$\mathcal{L} = i\,\frac{g_P}{f}\Big\{D_{\mu\nu}\,(\partial^\mu\Phi)\,(\partial^\nu \bar D) - (\partial^\nu D)\,(\partial^\mu\Phi)\,\bar D_{\mu\nu}\Big\} + \frac{\tilde g_P}{4\,f}\,\epsilon^{\mu\nu\alpha\beta}\Big\{D_{\mu\nu}\,(\partial_\alpha\Phi)\,(\partial^\tau \bar D_{\tau\beta}) + (\partial^\tau D_{\tau\beta})\,(\partial_\alpha\Phi)\,\bar D_{\mu\nu})\Big\}\,, \tag{13}$$

that are responsible for the s- and u-channel exchange contributions. As detailed in [8] the decay of the charged D^*-mesons implies

$$|g_P| = 0.57 \pm 0.07\,, \tag{14}$$

using $f = 90$ MeV. The parameter $\tilde{g}_P$ in (13) can not be extracted from empirical data directly. An accurate evaluation within unquenched QCD lattice simulations would be helpful but at present not available. Relying on the heavy-quark symmetry of QCD, the size of that parameter is estimated to be $\tilde{g}_P = g_P$ at leading order (see e.g. [8]).

The coupling constant g_P contributes to the effective interactions $V^{(0^+)}(s)$ and $V^{(1^+)}(s)$ introduced in (7) and (11) via s- and u-channel exchange of the D-meson ground states. In the scalar sector only the u-channel exchange of the $J^P = 1^-$ contributes. In contrast to the findings of Ref. [11], the s-channel exchange of the 1^- state does not contribute here. This is a consequence of the tensor field representation of the spin-one field. As already noted in [11], the influence of the u-channel process is of very minor importance for the formation of the $D^*_{s0}(2317)$. At $f = 90$ MeV and $g_P = 0$ we obtain a mass of 2304 MeV which is pulled down by 1 MeV only if we switch on $g_P = 0.57$. The situation is different for the axial-vector state $D^*_{s1}(2460)$. In this case there are three processes contributing: the s-channel exchange of the 1^- state and the u-channel exchanges of the 0^- and 1^- charmed mesons. The s-channel exchange of the 1^- state is a consequence of the tensor-field approach and was not present when using the vector-field representation of the spin-one D-mesons [11]. The influence of the u-channel exchange interaction on the formation of the $D^*_{s1}(2460)$ state is somewhat more important than it is in the scalar sector. For $f = 90$ MeV and $g_P = \tilde{g}_P = 0$ we obtain 2441 MeV, a mass which is pushed up by 5 MeV for $g_P = \tilde{g}_P = 0.57$. This effect is dominated largely by the exchange of the 1^- state. Including the s-channel exchange in addition yields a resonance mass of 2433 MeV for $f = 90$ MeV and $g_P = \tilde{g}_P = 0.57$. Clearly, further correction terms are needed in order to reproduce quantitatively the masses of the $D^*_{s0}(2317)$ and $D^*_{s1}(2460)$ states.

We turn to the local 2-body interaction terms, considering the chiral symmetry breaking and the chiral symmetric terms successively [11]. At chiral order Q_χ^2 the following terms break chiral symmetry explicitly,

$$\begin{aligned}
\mathcal{L}_{\chi-SB} = & -2\,c_1\,D\,\chi_0\,\bar{D} - (4\,c_0 - 2\,c_1)\,(D\,\bar{D})\,\mathrm{tr}\,\chi_0 \\
& + \frac{2\,c_0 - c_1}{f^2}\,D\,\bar{D}\,\mathrm{tr}\left(\Phi\,\chi_0\,\Phi\right) + \frac{c_1}{4\,f^2}\,D\left\{\Phi, \left\{\Phi\,, \chi_0\right\}\right\}\bar{D} \\
& + \tilde{c}_1\,D_{\alpha\beta}\,\chi_0\,\bar{D}^{\alpha\beta} + (2\,\tilde{c}_0 - \tilde{c}_1)\,(D_{\alpha\beta}\,\bar{D}^{\alpha\beta})\,\mathrm{tr}\,\chi_0 \\
& - \frac{2\,\tilde{c}_0 - \tilde{c}_1}{2\,f^2}\,D_{\alpha\beta}\,\bar{D}^{\alpha\beta}\,\mathrm{tr}\left(\Phi\,\chi_0\,\Phi\right) - \frac{\tilde{c}_1}{8\,f^2}\,D_{\alpha\beta}\left\{\Phi, \left\{\Phi\,, \chi_0\right\}\right\}\bar{D}^{\alpha\beta}\,,
\end{aligned} \tag{15}$$

where the matrix χ_0 was introduced in (3). The parameters c_1 and $\tilde{c}_1$ are determined by the empirical mass differences of the $J^P = 0^-$ and $J^P = 1^-$ charmed mesons. According to [11] we have

$$c_1 \simeq 0.44\,, \qquad \tilde{c}_1 \simeq 0.47\,. \tag{16}$$

The parameters c_0 and $\tilde{c}_0$ could in principle be determined by unquenched lattice QCD simulation upon studying the pion- and kaon-mass dependence of the D-meson ground states. So far they are unknown.

Clearly, the number of unknown parameters ($c_{0,2,3}$ and $\tilde{c}_{0,2,3}$) appears large at first. A free fit to the masses of the $D^*_{s0}(2317)$ and $D^*_{s1}(2460)$ states only would not be significant. Additional constraints from QCD should be used. According to [11] the parameters c_i and $\tilde{c}_i$ are degenerate in the heavy-quark mass limit, i.e. we expect

$$c_i \simeq \tilde{c}_i\,. \tag{17}$$

It is reassuring that the values for c_1 and $\tilde{c}_1$ given in (16) are quite compatible with the expectation of the heavy-quark symmetry relations (17). We consider further constraints from QCD as they arise in the limit of large number of colors N_c [18]. Since at leading order in a $1/N_c$ expansion single-flavour trace interactions are dominant, we anticipate the relations

$$c_0 \simeq \frac{c_1}{2}\,, \qquad c_2 \simeq -\frac{c_3}{2}\,, \qquad \tilde{c}_0 \simeq \frac{\tilde{c}_1}{2}\,, \qquad \tilde{c}_2 \simeq -\frac{\tilde{c}_3}{2}\,. \tag{18}$$

In the combined heavy-quark and large-N_c limit, we are left with one free parameter only. We may vary $c_3 = \tilde{c}_3$. The optimal value $c_3 = \tilde{c}_3 = 1.2$ together with $c_1, \tilde{c}_1$ as dictated by (16), $g_P = \tilde{g}_P = 0.57$ and $f = 90$ MeV predicts 2330 MeV and 2449 MeV for the $D^*_{s0}(2317)$ and $D^*_{s1}(2460)$ states respectively.

4 Isospin-violating width parameters and exotics

To determine the strong width parameters corresponding to the isospin-violating decays $D^*_{s0}(2317) \to D_s(1968)\,\pi^0$ and $D^*_{s1}(2460) \to D^*_s(2112)\,\pi^0$, it is important to reproduce the masses accurately. Consequently we allow for small variations of the parameters around the heavy-quark scenario, leaving the large-N_c relations (18) untouched. A precise reproduction of the scalar and axial-vector states is achieved for $c_3 = 1.0$ and $\tilde{c}_3 = 1.4$. For the mixing angle $\epsilon = 0.01$ determined by (5), the strong decay widths of the $D^*_{s0}(2317)$ and $D^*_{s1}(2460)$ are the same and equal to 140 keV. Our predictions for the widths are roughly an order of magnitude larger than the values obtained in previous coupled-channel calculations based on the chiral Lagrangian [19, 20]. As detailed in [8] the differences are due to the neglect of isospin breaking effects in the kaon and D-meson masses and to a lesser extent to the absence of chiral correction terms in [19, 20]. The importance of isospin breaking effects in the D-meson masses was also found in [21] using a different approach.

A striking prediction of the large-N_c scenario is a clear measurable signal of the exotic axial state in the ηD^* invariant mass distribution. The exotic axial state at mass 2568 MeV lies above the ηD^* threshold, giving it a width of about 18 MeV. The heavy-quark partner of this state lies at 2410 MeV with a width of 2 MeV below the ηD threshold. Since such states have not been seen so far in experiment, it is important to verify the compatibility of these predictions with available empirical constraints. The BELLE collaboration [22] measured indeed the invariant πD and πD^* mass distributions.

In Fig. 1 we confront the imaginary part of the $\pi D \to \pi D$ amplitude with the empirical πD mass distribution [22]. The latter is dominated by the broad $(\frac{1}{2}, 0)$ state, a member of the triplet to which the $D^*_{s0}(2317)$ belongs, as well as the tensor state $D^*_2(2460)$, the contribution of which is illustrated by the histogram. The possible presence of a narrow $(\frac{1}{2}, 0)$ state is not excluded by the present data. It is interesting to observe that the exotic state leads to a dip in the mass distribution rather than a peak. This is a consequence of the nearby ηD channel that couples strongly to that state. We note that, with the exception of a strong cusp effect at the $\bar{K}D$ threshold in the $(0, -1)$ sector, there is no further strong signal of any sextet state.

In Fig. 1 we compare also the imaginary part of the $\pi D^* \to \pi D^*$ amplitude with the empirical πD^* mass distribution [22]. The empirical distribution shows two axial and one tensor states. The contribution of the tensor state is illustrated by the histogram. In conventional approaches the tensor state $D^*_2(2460)$ is grouped together with the $D^*_1(2420)$ state to form a heavy-quark multiplet. Within the hadrogenesis conjecture we would expect to generate that multiplet dynamically via coupled-channel effects once the light vector mesons are considered as additional and explicit degrees of freedom. The theoretical amplitude of the present work describes only the broad state, which has a width of about 300 MeV. In contrast to the πD distribution shown in the left panel of Fig. 1 we do not predict any significant signal in the πD^* distribution that one may use to discover the exotic axial state. This reflects a coupling constant of that state to the πD^* channel that is almost compatible with zero.

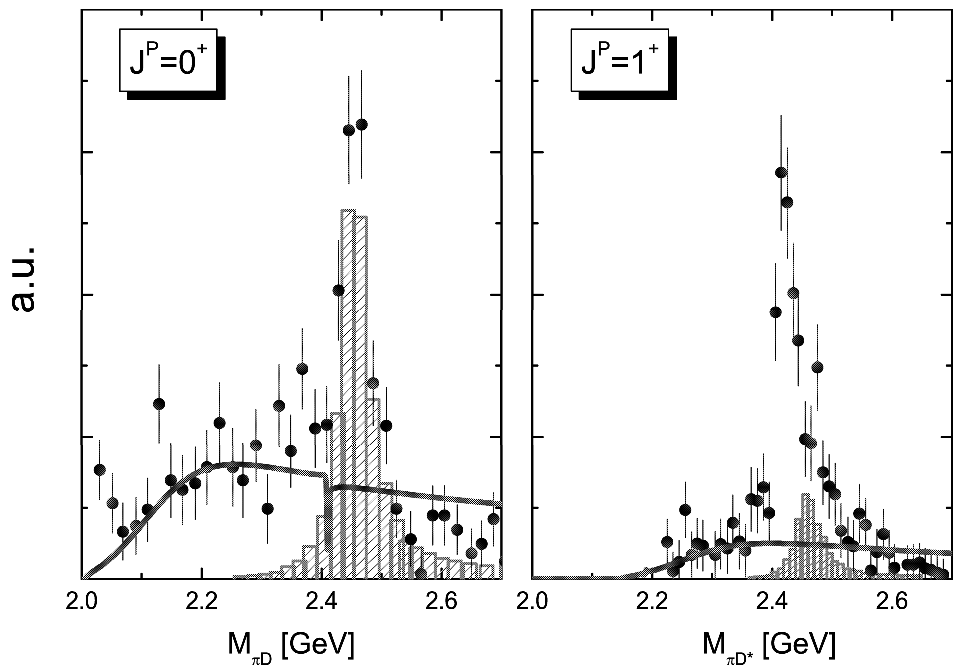

Figure 1: Mass spectra of the (I,S) $= (\frac{1}{2}, 0)$-resonances as seen in the $\pi\, D(1867)$-channel (l.h. figure $J^P = 0^+$) and $\pi\, D^*(2008)$-channel (r.h. figure $J^P = 1^+$). The solid lines show the theoretical mass distributions. The data are taken from [22] as obtained from the $B \to \pi\, D(1867)$ and $B \to \pi\, D^*(2008)$ decays. The histograms indicate the contribution from the $J = 2$ resonances $D_2^*(2460)$ as given in [22].

Nonetheless, we deem the exotic axial state to be easily discovered by ongoing experiments once the invariant ηD^* mass distribution is analyzed. The discovery of the scalar state, in contrast, would require a measurement of the πD invariant mass with an energy resolution of a few MeV as may be possible with the PANDA experiment at FAIR.

We briefly comment on the previous results of [11]. In that work a different scenario was investigated. Using the conventional vector-field representation of the 1^- charmed mesons, it was assumed that the axial-vector resonance $D_1^*(2420)$ was a member of the exotic sextet, predicted at leading order by chiral coupled-channel dynamics [10, 11]. In [11] chiral correction terms were tuned in such a way as to pull down the exotic axial state with $(I, S) = (\frac{1}{2}, 0)$ to match the properties of the $D_1^*(2420)$. The invariant πD and πD^* mass distributions as measured by the BELLE collaboration [22] were used as an additional constraint. It was argued that the scalar heavy-quark partner of the exotic axial state decouples from the πD channel, and therefore is not seen in the data. Based on the large-N_c relations (18), that were not considered in [11], we would deem this scenario unlikely.

5 Summary

Based on the chiral Lagrangian, properties of scalar and axial-vector meson molecules with open-charm content were studied. Chiral correction terms were incorporated systematically in the coupled-channel dynamics and calculated relying on constraints from large-N_c QCD and heavy-quark symmetry. We focused on the $D_{s0}^*(2317)$ and $D_{s1}^*(2460)$ states, computing both their masses and strong isospin-violating widths. The latter are predicted to be 140 keV. Chiral correction terms are important contributions to the decay widths. The invariant $\eta\, D^*$ invariant mass distribution shows a signal of a member of an exotic axial-vector sextet with a mass of 2568 MeV and a width of 18 MeV. While that state decouples from the πD^* spectrum, its heavy-quark partner defines a narrow dip at a mass of 2410 MeV with a width of 2 MeV in the $\pi\, D$ mass distribution.

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