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# Two-pole structure of the $D_0^*(2400)$

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ABSTRACT

The so far only known charmed non-strange scalar meson is dubbed as $D_0^*(2400)$ in the Review of Particle Physics. We show, within the framework of unitarized chiral perturbation theory, that there are in fact two ($I = 1/2$, $J^P = 0^+$) poles in the region of the $D_0^*(2400)$ in the coupled-channel $D\pi$, $D\eta$ and $D_s\bar{K}$ scattering amplitudes. With all the parameters previously fixed, we predict the energy levels for the coupled-channel system in a finite volume, and find that they agree remarkably well with recent lattice QCD calculations. This successful description of the lattice data is regarded as a strong evidence for the two-pole structure of the $D_0^*(2400)$. With the physical quark masses, the poles are located at $\left(2105^{+6}_{-8} - i\,102^{+10}_{-12}\right)$ MeV and $\left(2451^{+36}_{-26} - i\,134^{+7}_{-8}\right)$ MeV, with the largest couplings to the $D\pi$ and $D_s\bar{K}$ channels, respectively. Since the higher pole is close to the $D_s\bar{K}$ threshold, we expect it to show up as a threshold enhancement in the $D_s\bar{K}$ invariant mass distribution. This could be checked by high-statistic data in future experiments. We also show that the lower pole belongs to the same SU(3) multiplet as the $D_{s0}^*(2317)$ state. Predictions for partners in the bottom sector are also given.



## 1. Introduction

In our quest for understanding the fundamental theory of the strong interactions, Quantum Chromodynamics (QCD), the interpretation of the hadronic spectrum plays a key role. The latter has received a renewed interest with the recent advent of powerful experimental facilities, capable of exploring open and hidden charm (and bottom) energy ranges, as reflected in the Particle Data Group (PDG) [1]. The $D_{s0}^*(2317)$ [2] and $D_0^*(2400)$ [3] are the lightest scalar ($J^P = 0^+$) charm strange and non-strange mesons, respectively, and have attracted much attention [4] as their masses differ from the quark model expectations for $1P$ states. The $D_{s0}^*(2317)$ is around 150 MeV below the predicted mass [5,6] (see, however, Refs. [7,8]), and one would naively expect a significantly larger mass for a $c\bar{s}$ than for a $c\bar{n}$, in contrast with the observation. Clearly, an improved understanding of the nature of these resonances is needed. Different schemes have described them as mostly $c\bar{q}$ states [9–13], as mixture of $c\bar{q}$ with tetraquarks [14] or meson–meson [15] components, as purely tetraquarks [16–21], or as heavy–light meson molecules [22–30] motivated by the closeness of the $D_{s0}^*(2317)$ to the $DK$ threshold.

There have also been lattice QCD (LQCD) simulations to understand the charmed scalar sector. Early studies used only $c\bar{s}$ interpolators for the $D_{s0}^*(2317)$, obtaining masses generally larger than the physical one [31,32]. Masses consistent with the $D_{s0}^*(2317)$ and $D_0^*(2400)$ experimental ones were only obtained after including also meson–meson interpolators [33,34]. The first LQCD study of $D\pi$, $D\eta$, and $D_s\bar{K}$ coupled-channel scattering (for $m_\pi \simeq 391$ MeV) was recently reported [35]. Therein, a bound state with a large coupling to $D\pi$ is found and assigned to the $D_0^*(2400)$.

While the $D_{s0}^*(2317)$ is very narrow and its mass is well measured [1], the situation for the broad $D_0^*(2400)$ is less clear: the reported mass values for the $D_0^*(2400)^0$ at $B$-factories, $(2308 \pm 36)$ MeV (Belle [3]) and $(2297 \pm 22)$ MeV (BaBar [36]), differ from that in $\gamma A$ reactions, $(2407 \pm 41)$ MeV (FOCUS [37]), while the LHCb value for the charged partner lies in between [38]. These analyses use Breit–Wigner parameterizations and assume a single scalar particle.

Moreover, a better understanding of the $D_0^*(2400)$ is also important because its properties influence the shape of the scalar form factor $f_0$ in semileptonic $D \to \pi$ decays [29,39], and indirectly it has some impact in the form factor $f_+$ that determines

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$|V_{cd}|$ [40–44]. The bottom analogue is even more interesting because of the existing tension between the determinations of $|V_{ub}|$ from inclusive and exclusive $\bar{B}$ decays [45,46], and the implications on the unitarity triangle [47,48] and on new physics limits [49].

We study here the heavy–light pseudoscalar meson $J^P = 0^+$ scattering in the strangeness-isospin $(S, I) = (0, 1/2)$ sector, and present a strong case for the existence of two poles in the $D_0^*(2400)$ energy region (and similarly in the bottom sector). The affirmative evidence comes from a remarkably good agreement between our *parameter-free predicted* energy levels and the LQCD results reported in Ref. [35]. This two-pole structure was previously claimed in Refs. [24,26,50]. The leading order chiral Lagrangian was used in Refs. [24,26], while the next-to-leading order one was considered in Ref. [50]. At the next-to-leading order, there appear a few low-energy constants (LECs), which were unknown when the work of Ref. [50] was done, and that were estimated there by means of $1/N_c$ arguments and naturalness assumptions. Here, the values of the LECs are taken from Ref. [51], where they were fitted to lattice data. The predictions for masses and widths of the two poles that will be shown below represent thus a clear improvement over those obtained previously, and more importantly the two-pole structure finds now a strong support. Its dynamical origin is elucidated from the light-flavor SU(3) structure of the interaction, and we find that the lower pole is the SU(3) partner of the $D_{s0}^*(2317)$. Predictions for other $(S, I)$ sectors, including bottom ones, will also be given.

## 2. Formalism

We consider the $S$-wave $D\pi$, $D\eta$, and $D_s\bar{K}$ coupled-channel scattering. A unitary $T$-matrix can be written as [52,53]:

$$T^{-1}(s) = V^{-1}(s) - \mathcal{G}(s)\,, \tag{1}$$

with $s \equiv E^2$, the center-of-mass (CM) energy squared. The diagonal matrix $\mathcal{G}$ is constructed from the two-meson loop function, $\mathcal{G}_{ii}(s) = G(s, m_i, M_i)$ [51], where $m_i$ and $M_i$ are the masses of the light and heavy pseudoscalar mesons in the channel $i$, respectively. It carries the unitarity cut and is regularized with a subtraction constant $a(\mu)$ (at a scale $\mu = 1$ GeV). The matrix $V(s)$ contains the interaction potentials, which are taken from the $\mathcal{O}(p^2)$ chiral Lagrangian of Ref. [56]. They depend on six LECs, $h_{0,\ldots,5}$. For the LECs and $a(\mu)$ we use the values and uncertainties obtained in Ref. [51] from a fit to LQCD results of the $S$-wave charm–light pseudoscalar-meson scattering lengths in several $(S, I)$ sectors. Notice that the channel $(0, 1/2)$ was not included in this fit.

Above threshold, the phase shift $[\delta_i(s)]$ and the inelasticity $[\eta_i(s)]$ of channel $i$ are related to the diagonal elements of $T$, $i\,p_i\,T_{ii} = 4\pi\sqrt{s}\left(\eta_i\, e^{2i\delta_i} - 1\right)$, with $p_i(s)$ the CM momentum. Bound, resonant, and virtual states are associated to poles in different Riemann sheets (RS) of the $T$-matrix. In our three-channel problem, RS are denoted as $(\xi_1\,\xi_2\,\xi_3)$, $\xi_i = 0, 1$, and are defined through analytical continuations:

$$\mathcal{G}_{ii}(s) \to \mathcal{G}_{ii}(s) + i\,\frac{p_i(s)}{4\pi\sqrt{s}}\,\xi_i\,. \tag{2}$$

Thus, (000) is the physical RS. The coupling $g_i$ of a pole to the channel $i$ is obtained from the residue $(g_i^2)$ of $T_{ii}$.

In LQCD, interactions between quarks and gluons are considered in a cubic finite lattice with some boundary conditions. In a finite volume, the momentum running in loops can take only discrete values, $\vec{q} = \frac{2\pi}{L}\vec{n}$, $\vec{n} \in \mathbb{Z}^3$, with $L^3$ the lattice volume. The loop functions are replaced by [57]:

$$\widetilde{\mathcal{G}}_{ii}(s, L) = \mathcal{G}_{ii}(s) + \frac{1}{L^3}\sum_{\vec{n}}^{|\vec{q}|<\Lambda} I_i(\vec{q}\,) - \int_0^\Lambda \frac{q^2 \mathrm{d}q}{2\pi^2}\, I_i(\vec{q}\,)\,, \tag{3}$$

with $\Lambda \to \infty$ and $I_i(\vec{q}\,)$ given in Eq. (13) of Ref. [58]. The potentials $V(s)$ in Eq. (1) do not receive any finite volume correction, and thus the $T$-matrix in a finite lattice, $\widetilde{T}(s, L)$, reads:

$$\widetilde{T}^{-1}(s, L) = V^{-1}(s) - \widetilde{\mathcal{G}}(s, L)\,. \tag{4}$$

The energy levels are obtained as poles of $\widetilde{T}(s, L)$, and they can be directly compared with those obtained in LQCD simulations.

Here we compare with Ref. [35], where energy levels relevant to the $D\pi$, $D\eta$, and $D_s\bar{K}$ channels at different volumes are reported. We employ the meson masses of that reference in the evaluation of $V(s)$ and $\mathcal{G}_{ii}(s)$.

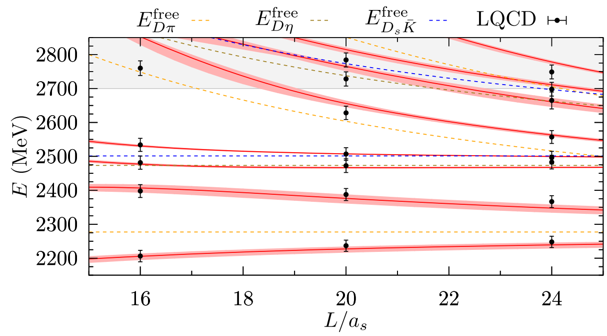


**Fig. 1.** Comparison of Ref. [35] (0, 1/2) energy levels (black dots) with our predictions (red lines and bands). The bands represent the $1\sigma$ uncertainties derived from the LECs fitted in Ref. [51], and $a_s = 0.12$ fm. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

**Table 1**
Position ($\sqrt{s} = M - i\Gamma/2$), couplings (in GeV) and RS of the two poles found in the (0, 1/2) sector using LQCD [35] or physical masses.

| Masses | $M$ (MeV) | $\Gamma/2$ (MeV) | RS | $\lvert g_{D\pi}\rvert$ | $\lvert g_{D\eta}\rvert$ | $\lvert g_{D_s\bar{K}}\rvert$ |
|---|---|---|---|---|---|---|
| lattice | $2264^{+8}_{-14}$ | 0 | (000) | $7.7^{+1.2}_{-1.1}$ | $0.3^{+0.5}_{-0.3}$ | $4.2^{+1.1}_{-1.0}$ |
| | $2468^{+32}_{-25}$ | $113^{+18}_{-16}$ | (110) | $5.2^{+0.6}_{-0.4}$ | $6.7^{+0.6}_{-0.4}$ | $13.2^{+0.6}_{-0.5}$ |
| physical | $2105^{+6}_{-8}$ | $102^{+10}_{-12}$ | (100) | $9.4^{+0.2}_{-0.2}$ | $1.8^{+0.7}_{-0.7}$ | $4.4^{+0.5}_{-0.5}$ |
| | $2451^{+36}_{-26}$ | $134^{+7}_{-8}$ | (110) | $5.0^{+0.7}_{-0.4}$ | $6.3^{+0.8}_{-0.5}$ | $12.8^{+0.8}_{-0.6}$ |

## 3. Results

The (0, 1/2) energy levels as a function of $L$ are shown in Fig. 1. The region above 2.7 GeV (shaded in Fig. 1) is beyond the range of applicability of our $\mathcal{O}(p^2)$ chiral unitary formalism. Below that energy, the agreement of our computed energy levels with those obtained in the LQCD simulation is excellent. This is remarkable, since no fit to the LQCD energy levels is performed.

The level below $D\pi$ threshold is interpreted in Ref. [35] as a bound state associated to the $D_0^*(2400)$. For infinite volume and with the lattice meson masses, our $T$-matrix also presents this pole. The second level, lying between the $D\pi$ and $D\eta$ thresholds, is very shifted with respect to both of them, hinting at the presence of another pole in infinite volume, that we find slightly below the $D\eta$ threshold. Both poles are collected in the upper half of Table 1 and represented with empty red symbols in Fig. 2.

Next, we study the spectroscopic content of our amplitudes when the physical masses are employed. The found poles are collected in Table 1, and shown in Fig. 2. The pole positions and residues for the physical case are similar to those calculated in Ref. [59]. Comparing the couplings, we see that the bound state below the $D\pi$ threshold evolves into a resonance above it when physical masses are used (notice that the threshold decreases from 2277 MeV to 2005 MeV). Such an evolution is typically found for $S$-wave poles (*e.g.*, for the $\sigma$ meson [60,61]). The second pole

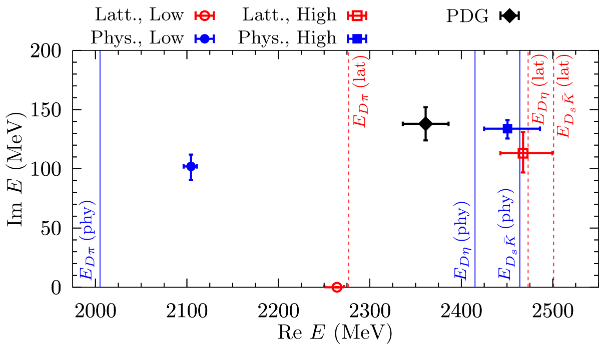


**Fig. 2.** Complex energy plane location of the two-pole-structure collected in Table 1. Empty red (filled blue) symbols stand for the poles obtained when the LQCD [35] (physical) masses are used. The black diamond represents the isospin average of the PDG values for $D_0^*(2400)^0$ and $D_0^*(2400)^+$ [1].

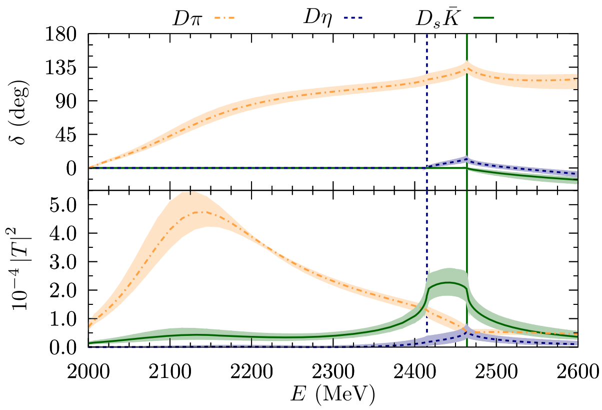


**Fig. 3.** Phase shift and modulus squared of the diagonal $D\pi$, $D\eta$ and $D_s\bar{K}$ amplitudes, $T_{ii}(s)$, in the $(0, 1/2)$ sector.

moves very little and its couplings are rather independent of the meson masses. For physical masses, it is a resonance located between the $D\eta$ and the $D_s\bar{K}$ thresholds in the (110) RS, continuously connected to the physical sheet. The mass of the $D_0^*(2400)$ reported by the PDG [1] lies between those of the two poles found here, whereas the widths are similar (Fig. 2). We conclude the $D_0^*(2400)$ structure is produced by two different states (poles), alongside with complicated interferences with the thresholds. This two-pole structure was previously reported in Refs. [24,26,50], and it receives here a robust support.

Phase shifts and $|T_{ii}(s)|^2$ are shown in Fig. 3. The lower pole causes a very mild effect in the $D\eta$ and $D_s\bar{K}$ amplitudes. It couples mostly to $D\pi$ where a peak around 2.1 GeV is clearly seen, while the phase goes through $\pi/2$ at $\sqrt{s} \simeq 2.2$ GeV. The higher pole manifests in a more subtle way. It produces a small enhancement in the $D\pi$ amplitude, but the strongest effect is a clear peak in the $D_s\bar{K}$ amplitude around 2.45 GeV. However, the shape is quite non-conventional. Despite the relatively large width, the amplitude shows a narrow peak stretched between two cusps at the $D\eta$ and $D_s\bar{K}$ thresholds.[1] Such a behavior provides a possible test of the two-pole structure. The BaBar and Belle data for $B \to D_s^- K\pi$ show an enhancement in the $D_s^- K$ invariant mass distribution [63,64] (see also Ref. [65]). This might confirm the features of the second pole found here, although better statistics data are required.

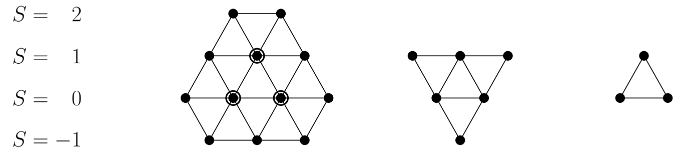


**Fig. 4.** Weight diagrams of the $\overline{\mathbf{15}}$, $\mathbf{6}$ and $\overline{\mathbf{3}}$ irreps.

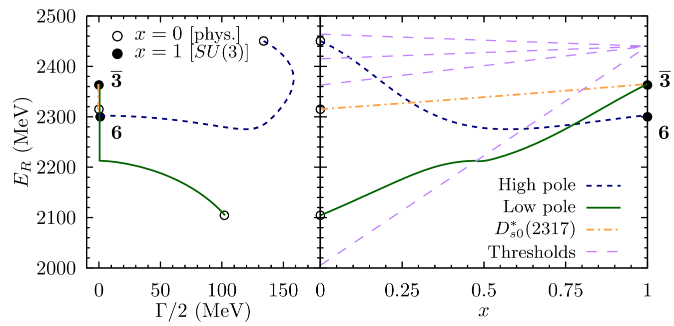


**Fig. 5.** Evolution from the physical to the flavor SU(3) symmetric cases of $(0, 1/2)$ and $(1, 0)$ poles $\sqrt{s} = E_R - i\Gamma/2$. The blue short dashed, green solid and the orange dashed dotted lines represent the higher and lower $D_0^*(2400)$ and the $D_{s0}^*(2317)$, respectively. Left: path followed by the poles in the complex plane. Right: evolution of $E_R$ with $x$. The purple long dashed lines stand for the $D\pi$, $D\bar{K}$, $D\eta$, and $D_s\bar{K}$ thresholds (from bottom to top).

## 4. SU(3) study

To gain further insights, we study the evolution of the two poles in the light-flavor SU(3) limit, *i.e.*, when all light and heavy meson masses take common values, $m_i = m$ and $M_i = M$, respectively. Similar analyses were done in Refs. [24,27]. In this limit, the heavy–light meson scattering decomposes into irreducible representations (irreps) as $\bar{\mathbf{3}} \otimes \mathbf{8} = \overline{\mathbf{15}} \oplus \mathbf{6} \oplus \bar{\mathbf{3}}$ (Fig. 4), and the potential matrix can be diagonalized, $V_d(s) = D^\dagger V(s) D = \mathrm{diag}\left(V_{\overline{\mathbf{15}}}, V_{\mathbf{6}}, V_{\bar{\mathbf{3}}}\right)$. Since the three channels have a common subtraction constant [51], the $T$-matrix is diagonalizable, $T_d(s) = D^\dagger T(s) D = \mathrm{diag}(T_{\overline{\mathbf{15}}}, T_{\mathbf{6}}, T_{\bar{\mathbf{3}}})$, where $T_A^{-1}(s) = V_A^{-1}(s) - G(s, m, M)$, $A \in \{\overline{\mathbf{15}}, \mathbf{6}, \bar{\mathbf{3}}\}$. At chiral $\mathcal{O}(p)$, $V_d(s) = f(s)\,\mathrm{diag}(1, -1, -3)$, with $f(s)$ positive in the scattering region, showing that the interaction in the $\mathbf{6}$ and $\bar{\mathbf{3}}$ ($\overline{\mathbf{15}}$) irreps is attractive (repulsive). The most attractive irrep, $\bar{\mathbf{3}}$, admits a $c\bar{q}$ ($q = u, d, s$) configuration. At $\mathcal{O}(p^2)$, the potentials receive corrections, but these qualitative features remain.

We can connect the physical and SU(3) symmetric cases by continuously varying the meson masses as:

$$m_i = m_i^{\mathrm{phy}} + x(m - m_i^{\mathrm{phy}})\,, \tag{5}$$

and analogously for $M_i$. Thus, $x = 0$ ($x = 1$) corresponds to the physical (SU(3) symmetric) case. Numerically, we take $m = 0.49$ GeV and $M = 1.95$ GeV. The evolution of the poles with $x$ is shown in Fig. 5. The lower $D_0^*(2400)$ pole found in the physical case (in the (100) RS) connects with a bound state of $T_{\bar{\mathbf{3}}}$ in the SU(3) limit, whereas the higher pole (in the (110) RS) connects with a virtual ($V_{\mathbf{6}}$ is not attractive enough to bind) state generated in $T_{\mathbf{6}}$.[2]

In the $(S, I) = (1, 0)$ sector involving the $DK$ and $D_s\eta$ channels, and using the same inputs and physical masses, we find a

[1] A similar interplay of thresholds and poles was also observed in the case of the $N(1535)$ and $N(1650)$ in Ref. [62], where the relevant thresholds are $K\Lambda$, $K\Sigma$ and $\eta N$.

[2] In the physical case, the RS are specified by $(\xi_1\,\xi_2\,\xi_3)$, with $\xi_i = 0$ or 1 [Eq. (2)]. In the SU(3) symmetric case, since all channels have the same threshold, there are only two RS, the physical [(000)] and the unphysical [(111)] sheets. To connect the lower pole in the physical case, located in the (100) RS, with the $T_{\bar{\mathbf{3}}}$ pole, in the (000) RS, we vary the parameter $\xi_1 = 1 - x$. For the higher pole, one has to evolve $\xi_3 = x$ to connect (110) (physical case) with (111) (SU(3) limit).

bound state at $2315^{+18}_{-28}$ MeV [51] which is naturally identified with the $D^*_{s0}(2317)$. Its evolution is also shown in Fig. 5, and we see it emerges from the $T_{\bar{3}}$ pole. Hence, the $D^*_{s0}(2317)$ and the lower $D^*_0(2400)$ poles are flavor SU(3) partners. This in fact solves the puzzle mentioned in the Introduction about why the charm-strange $D^*_{s0}(2317)$ is lighter than the nonstrange $D^*_0(2400)$ were the PDG mass taken for it. The $D^*_{s0}(2317)$ has two nonstrange SU(3) partners, and their masses are indeed lighter.[3]

We now discuss other sectors in the physical case. The $(-1, 0)$ involves only the $D\bar{K}$ channel, and it is part of the **6** irrep, which is weakly attractive. Indeed, we find a virtual pole, at $2342^{+13}_{-41}$ MeV, roughly 20 MeV below threshold, which has a sizable influence on the $D\bar{K}$ scattering length [51]. The $(1, 1)$ sector, involving the $D_s\pi$ and $DK$ channels, has contributions from the **6** and the repulsive $\overline{\mathbf{15}}$. Because of this, we do not find any pole that can be associated to a physical state.

In the bottom sector, due to the heavy-flavor symmetry,[4] we foresee a similar pattern. In the $(0, 1/2)$ sector there is also a two-pole structure, located at $(5537^{+9}_{-11}, 116^{+14}_{-15})$ MeV and $(5840^{+12}_{-13}, 25^{+6}_{-5})$ MeV. For $(S, I) = (1, 0)$, we find a $B^*_{s0}$ state with a mass of $5724^{+17}_{-24}$ MeV, bound by about 50 MeV, as the $D^*_{s0}(2317)$ in the charm sector. All these pole positions are very similar to those found already at $\mathcal{O}(p)$ [26]. In the $(-1, 0)$ sector, we find a virtual state located almost at threshold, which can also appear as a bound state considering the $\mathcal{O}(p^2)$ parameter uncertainties. As in the charm case, we do not find physical poles in the $(1, 1)$ sector that could be identified with the $X(5568)$, recently reported by the D0 Collaboration [67], but not seen in other experiments [68, 69]. We conclude the $X(5568)$ is not generated by the $\bar{B}_s\pi - \bar{B}K$ rescattering [70] (see also Refs. [71–73]).

Finally, we remind that heavy quark spin symmetry relates the $0^+$ and $1^+$ sectors, and thus in the latter we find a similar pattern of bound, resonant and virtual states [24–26,74,66]. We highlight the predictions for the $\bar{\mathbf{3}}$ multiplet, where we find $2436^{+16}_{-22}$ MeV and $\left(2240^{+5}_{-6} - i\,93^{+9}_{-9}\right)$ MeV, for the $D_{s1}(2460)$ and a new $D_1$ resonance, respectively. In the bottom sector, we predict $5768^{+17}_{-23}$ MeV and $\left(5581^{+9}_{-11} - i\,115^{+13}_{-15}\right)$ MeV for the $B_{s1}$[5] and the lowest $B_1$. The higher $D_1$ and $B_1$ poles stemming from the **6** will presumably be affected by channels involving $\rho$ mesons [76,77]. For $(-1, 0)$, we find, as in the $0^+$ case, an axial state located almost at threshold in the bottom sector, while for charmed mesons the pole (virtual) moves deep in the complex plane.

## 5. Conclusions

We have studied the $D\pi$, $D\eta$ and $D_s\bar{K}$ scattering in the $J^P = 0^+$ and $(S, I) = (0, 1/2)$ sector. Although so far only one meson, the $D^*_0(2400)$, has been reported in experiments [1], we present a strong support for the existence of two poles in the $D^*_0(2400)$ mass region: the physical amplitudes, that contain two poles, when put in a finite volume produce energy levels that successfully describe the recent LQCD results in Ref. [35] without adjusting any parameter. The two poles are located at $\left(2105^{+6}_{-8} - i\,102^{+10}_{-12}\right)$ MeV and $\left(2451^{+36}_{-26} - i\,134^{+7}_{-8}\right)$ MeV, with the largest couplings to the $D\pi$ and $D_s\bar{K}$ channels, respectively. A group theoretical analysis shows that the lower pole and the $D^*_{s0}(2317)$ complete the $\bar{\mathbf{3}}$ multiplet, being thus flavor SU(3) partners. We expect the two-pole structure to produce distinctive features in $D\pi$, $D\eta$ and $D_s\bar{K}$ invariant mass spectra in high-energy reactions, such as $\bar{B}$ decays. In particular, despite of its large width, the higher pole shows up as a narrow peak in the $D_s\bar{K} \to D_s\bar{K}$ amplitude and should produce a sizable near-threshold enhancement. Note that clear $D_s\bar{K}$ threshold enhancements have been already observed in $B$ decays [63, 64]. Future data from better statistics experiments, such as LHCb and Belle-II, will shed light into their origin.

A similar resonance pattern is also found in the bottom sector. Besides the two-pole structure, we stress the possible existence of a near-threshold bound or virtual state in the $\bar{B}\bar{K}$ (or $BK$) channel, both for $0^+$ and $1^+$ sectors. These exotic states, with quark content $bs\bar{d}\bar{u}$, will have a large impact in the scattering length, and if bound they could only decay through weak and/or electromagnetic interactions.

The predicted phase shifts, both in the charm and bottom sectors, can be used as inputs to the Omnès representation of the scalar form factors describing heavy meson semileptonic decays [29,39]. Thus, the special two-pole structure discussed here is also of interest to achieve an accurate determination of the Cabibbo–Kobayashi–Maskawa matrix elements.

It is also worthwhile to notice the resemblance between the results obtained here for the $D^*_0(2400)$ and the widely discussed two-pole structure of the $\Lambda(1405)$ linked to $\Sigma\pi$ and $N\bar{K}$ [1,53, 78–80]. The existence of such a two-pole structure is rooted in both cases in chiral dynamics, which on one hand determines the interaction strength, and on the other hand ensures the lightness of pions and kaons. The latter is important to separate the two poles from higher hadronic channels. A two-pole structure driven by chiral dynamics is also found in Refs. [81–83] for the $K_1(1270)$.


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[3] In Ref. [54], in order to solve this puzzle, the $D^*_0(2400)$ was proposed to be an excited state, while the ground state in that work had a mass around 1980 MeV, smaller than both the $D\pi$ threshold and the $D^*$ mass. A state with a mass 2114 MeV similar to our lower pole was predicted in the quark model of Ref. [55].

[4] The LECs $h_i$ appearing in the Lagrangian of Refs. [51,56] depend on the heavy quark mass $m_Q$. By comparing with the Lagrangian of Ref. [66], one deduces that the LECs scale as $h_i \sim m_Q$ for $i = 0, \ldots, 3$ and $h_i \sim 1/m_Q$ for $i = 4, 5$. Furthermore, we follow Ref. [26] to translate the value of the subtraction constant from the charm to the bottom sector.

[5] We comment here that our predictions for the $B^*_{s0}$ and $B_{s1}$ nicely agree with those in the LQCD simulation of Ref. [75].

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