

From the κ via the $D_{s0}^*(2317)$ to the χ_{c0} : Connecting light and heavy scalar mesons

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Pole trajectories connecting light and heavy scalar mesons, both broad resonances and quasibound states, are computed employing a simple coupled-channel model. Instead of varying the coupling constant as in previous work, quark and meson masses are continuously changed, so as to have one scalar meson evolve smoothly into another with different flavor(s). In particular, it is shown, among several other cases, how the still controversial $K_0^*(800)$ turns into the established χ_{c0} , via the disputed $D_s(2317)$. Moreover, a $\chi'_{c0}(3946)$ is predicted, which may correspond to the recently observed $Y(3943)$ resonance. These results lend further support to our unified dynamical picture of all scalar mesons, as unitarized $q\bar{q}$ states with important two-meson components.

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After more than four decades, understanding the scalar mesons continues to pose serious difficulties to theorists as well as experimentalists. Still today, no consensus exists about the lightest and oldest structures in the scalar-meson sector, namely, the σ ($f_0(600)$ [1]) [2,3] and the κ ($K_0^*(800)$ [1]) [3–5]. But also the discovery of the surprisingly light charmed scalar $D_{s0}^*(2317)$ [6], though giving a new boost to meson spectroscopy in general, has not contributed to the understanding of scalar mesons, as can be seen from the many different approaches to the $D_{s0}^*(2317)$ in the literature (see Ref. [7] for a representative, albeit not totally exhaustive, list of references). Here, we shall focus on a formalism which successfully describes all mesonic resonances, including the scalar mesons.

In Ref. [8] it was shown that the $D_{s0}^*(2317)$ meson can be straightforwardly explained as a normal $c\bar{s}$ state, but strongly coupled to the nearby DK channel, which is responsible for its low mass. The framework for this calculation was a simple coupled-channel model, which had been employed previously [9] to fit the S -wave $K\pi$ phase shifts, and predict the now listed [1] $K_0^*(800)$, besides reproducing the established $K_0^*(1430)$. Furthermore, another charmed scalar meson was predicted in Ref. [8], i.e., a broad D_0^* resonance above the $D\pi$ threshold, somewhere in the energy region 2.1–2.3 GeV, which may correspond to the $D_0^*(2300–2400)$ [1,10]. Also higher-mass D_{s0}^* and D_0^* resonances were foreseen [8], which have not been observed so far.

The purpose of this paper is to show the interconnection of the scalar mesons $K_0^*(800)$, $D_0^*(2300–2400)$, $D_{s0}^*(2317)$ with one another, and also with the established $\chi_{c0}(3415)$ [1]. Moreover, the same interconnection will be demon-

strated for the higher-mass recurrences of these scalars, thereby finding a candidate for the very recently observed $Y(3943)$ charmonium state [11]. For that purpose, we shall employ the above-mentioned coupled-channel model, but now for fixed, physical coupling, while quark and threshold masses will be varied. Thus, a continuous and smooth transition can be achieved from one scalar meson to another. Crucial here will be a mass scaling [12,13] of the two parameters modeling the off diagonal potential that couples the confined and decay channels. This way, these two parameters, identical to the ones used in Refs. [8,9,12], suffice to reasonably describe a vast range of distinct scalar mesons. On the other hand, the confinement and quark-mass parameters are taken at their usual published values.

The starting point is a simple, intuitive coupled-channel model, describing a confined $q\bar{q}$ system, coupled to one meson-meson channel accounting for the possibility of real or virtual decay via the 3P_0 mechanism. If the transition potential is taken to be a spherical delta function, the 1×1 inverse K matrix can be solved in closed form, reading [9]

$$\cot(\delta_\ell(p)) = \frac{n_\ell(pa)}{j_\ell(pa)} - \left[2\lambda^2 \mu pa j_\ell^2(pa) \times \sum_{n=0}^{\infty} \frac{B_{n\ell_c}}{E - E_{n\ell_c}} \right]^{-1}, \quad (1)$$

where j_ℓ , n_ℓ are spherical Bessel and Neumann functions, respectively, λ is the 3P_0 coupling, a is the delta-shell radius, $E_{n\ell_c}$ are the energies of the bare confinement spectrum, $B_{n\ell_c}$ are the corresponding weight factors, p is the on shell relative momentum in the two-meson channel, given by the kinematically relativistic expression

$$4s p^2 = [s - (M_1 + M_2)^2][s - (M_1 - M_2)^2], \quad (2)$$

and μ is the ensuing relativistic reduced mass

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$$\mu \equiv \frac{1}{2} \frac{dp^2}{d\sqrt{s}} = \frac{\sqrt{s}}{4} \left[1 - \left(\frac{M_1^2 - M_2^2}{s} \right)^2 \right]. \quad (3)$$

As the present paper deals with scalar mesons, we have $\ell = 0$ and $\ell_c = 1$ in Eq. (1). Moreover, since only ground states and first radial excitations are considered here, we shall approximate the infinite sum in Eq. (1) by two confinement-spectrum states plus one rest term, also sticking to the numerical values used in Refs. [8,9,12], namely $B_{01} = 1.0$, $B_{11} = 0.2$, and $B_{21} = E_{21} = \infty$, with $B_{21}/E_{21} = 1$. As for the two confinement levels, we parametrize them by a harmonic oscillator [8,12], i.e.,

$$E_{n1} = (2n + 2.5)\omega + m_{q_1} + m_{q_2}, \quad (4)$$

where $\omega = 0.190$ GeV, $m_n = 0.406$ GeV ($n = u, d$), $m_s = 0.508$ GeV, and $m_c = 1.562$ GeV, as in previous work [8,9,12,14,15]. Finally, we assume a mass scaling of the parameters a and λ given by [12,13]

$$a_{ij}\sqrt{\mu_{ij}} = \text{constant}, \quad \lambda_{ij}\sqrt{\mu_{ij}} = \text{constant}, \quad (5)$$

where the labels ij refer to a particular combination of quark flavors, and $\mu_{ij} \equiv m_{q_i}m_{q_j}/(m_{q_i} + m_{q_j})$ is the corresponding reduced quark mass. This procedure ensures flavor invariance of our equations. Using then the values $\lambda_{ns} = 0.75$ GeV^{-3/2} and $a_{ns} = 3.2$ GeV⁻¹ from the fit to the $K\pi$ S -wave phase shifts in Ref. [9], we have fixed all our parameters,¹ which allows to show the predictive power of our approach. For the required input mesons masses, we take the isospin-averaged values [1] $M_\pi = 0.1373$ GeV, $M_K = 0.4957$ GeV, and $M_D = 1.867$ GeV.

Now we can compute pole trajectories in the complex energy plane for scalar resonances and (virtual) bound states, by searching the values of s for which $\cot\delta_0(p(s)) = i$. However, instead of freely varying λ as in previous work, we shall keep λ_{ns} fixed at its physical value of 0.75 GeV^{-3/2}, while changing instead one of the quark masses, as well as one of the meson masses in the decay channel. This way we can make one scalar meson turn into another. For instance, by letting

$$\left. \begin{aligned} m_{q_1} &= m_n + \alpha(m_c - m_n) \\ M_1 &= M_\pi + \alpha(M_D - M_\pi) \end{aligned} \right\}, \quad 0 \leq \alpha \leq 1, \quad (6)$$

we smoothly change the κ ($n\bar{s}$) meson, coupling to the πK channel, into the $D_{s0}^*(2317)$ ($c\bar{s}$), coupling to DK . The poles themselves are numerically found and checked with two independent methods, i.e., the *MINUIT* package of CERN [16], and *MATHEMATICA* [17].

¹Note that we use here somewhat shifted confinement levels as compared to Ref. [9], namely, the ones following from Eq. (4). This gives rise to a slightly lighter and broader κ meson, and a heavier $K_0^*(1430)$.

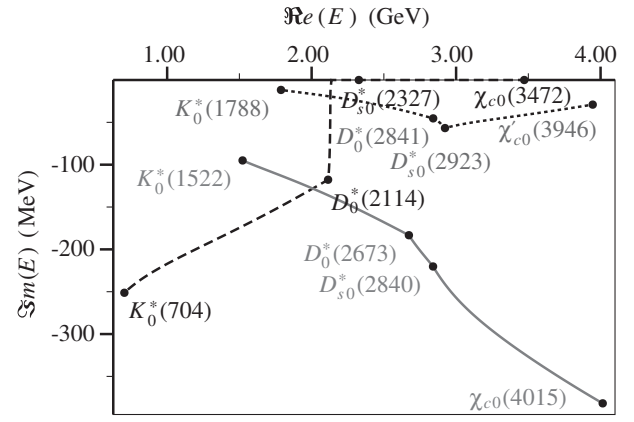


FIG. 1. Scalar-meson pole trajectories in the complex energy plane. Dots represent predicted resonances or bound states. See text and Eqs. (7) and (8) for further details.

In Fig. 1, one sees in one glimpse the nine trajectories

- a : $K_0^*(704) \rightarrow D_0^*(2114) \rightarrow D_{s0}^*(2327) \rightarrow \chi_{c0}(3472)$,
 - b : $K_0^*(1522) \rightarrow D_0^*(2673) \rightarrow D_{s0}^*(2840) \rightarrow \chi_{c0}(4015)$,
 - c : $K_0^*(1788) \rightarrow D_0^*(2841) \rightarrow D_{s0}^*(2923) \rightarrow \chi_{c0}'(3946)$,
- (7)

where the numbers between parentheses are the real parts (in MeVs) of the respective resonance/bound-state poles, the corresponding imaginary parts being

- a : $K_0^*(-251) \rightarrow D_0^*(-118) \rightarrow D_{s0}^*(0) \rightarrow \chi_{c0}(0)$,
 - b : $K_0^*(-95) \rightarrow D_0^*(-183) \rightarrow D_{s0}^*(-220) \rightarrow \chi_{c0}(-382)$,
 - c : $K_0^*(-12) \rightarrow D_0^*(-45) \rightarrow D_{s0}^*(-57) \rightarrow \chi_{c0}'(-29)$.
- (8)

Before discussing the actual trajectories, a few remarks are due concerning the precise values found for the pole positions. Clearly, for such a simple model without any fitting freedom, moreover covering a vast energy range, a very accurate reproduction of the masses and widths of all experimentally observed mesons cannot, and should not even be expected. In particular, the inclusion of only the lowest, dominant decay channel for each state will certainly reflect itself in one way or another. For instance, the much too small width of our $K_0^*(1788)$, which should correspond to the observed [18] $K_0^*(1820)$, is probably due to the neglect of the important $K\eta'$ channel. Furthermore, the somewhat too large mass of our $\chi_{c0}(3472)$, as compared to the established [1] $\chi_{c0}(3415)$, may very well be due to the omission of vector-vector decay channels, which are relevant for charmonium ground states [19]. Note, however, that the latter discrepancy of 57 MeV is quite insignificant when compared to the huge coupled-channel shifts in charmonium recently found in Refs. [20,21]. Notwithstanding, a clear identification can be made of our broad $K_0^*(704)$, $D_0^*(2114)$, and $K_0^*(1522)$ states with the listed [1] $K_0^*(800)$, $D_0^*(2300-2400)$, and

$K_0^*(1430)$ resonances, respectively. Here, one should also notice that we give the real parts of the pole positions of our resonances, which usually do not coincide with the experimental masses resulting from Breit-Wigner fits when the widths are large. As for the remaining observed mesons, our $D_{s0}^*(2327)$ is very close to the $D_{s0}^*(2317)$, while our $\chi_{c0}'(3946)$, with a width of about 60 MeV, seems a good candidate for the brand new [11] charmonium state $Y(3943)$. Finally, we predict the two medium-broad charmed mesons $D_0^*(2841)$ and $D_{s0}^*(2923)$, so far undetected, as well as the very broad states $D_0^*(2673)$, $D_{s0}^*(2840)$, and $\chi_{c0}(4015)$, which will be extremely hard to observe at all. In any case, the predictions for the latter higher-mass states may change significantly when additional decay channels are taken into account.

Turning now to the trajectories themselves, it is remarkable to observe that physical states with radically disparate widths can be continuously connected to one another in flavor. This is one of the reasons why scalar-meson spectroscopy is so intricate. Moreover, as we shall see below, states on the same mass trajectory can have different origins when viewed as $q\bar{q}$ states distorted by meson loops, which point will become clearer when we study Fig. 2. Anyway, the first radial excitations of the $n\bar{s}$, $c\bar{n}$, $c\bar{s}$, and $c\bar{c}$ systems are all on the same trajectory in Fig. 1, i.e., the one connecting the $K_0^*(1788)$ and $\chi_{c0}'(3946)$.

In Fig. 2, the lowest states for the various flavor combinations are displayed again, but now also show how the corresponding poles move when the coupling λ is reduced from its fixed value. We see that the $K_0^*(704)$ and the $D_0^*(2114)$ appear to find their origin in the continuum, corresponding to infinitely negative imaginary parts of their pole positions, while the $D_{s0}^*(2327)$ and $\chi_{c0}(3472)$ are connected to the confinement spectrum, with poles on the real axis. This is quite surprising for the nearby pair $D_0^*(2114)$ - $D_{s0}^*(2327)$. However, even the physical $D_{s0}^*(2317)$ itself can be either interpreted as a “confinement” state [12,22], or a “continuum” state [8], depending

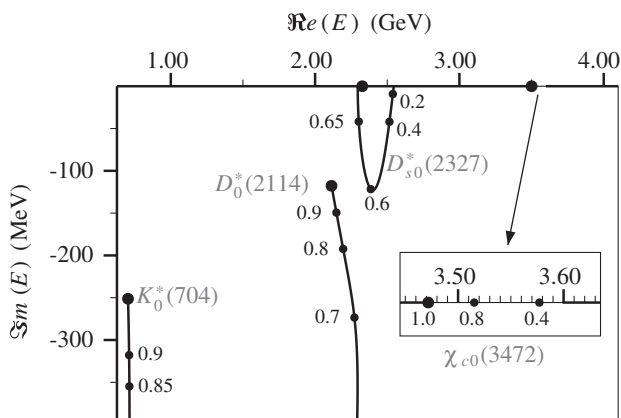


FIG. 2. Pole trajectories of lowest states as a function of λ . Numbers indicate reductions relative to the maximum value.

on tiny changes in e.g. the parameter a . What this figure also shows is an extremely delicate balance of coupling effects. With a small decrease of λ , the $D_0^*(2300\text{--}2400)$ and especially the κ meson would become even broader and thus almost impossible to observe experimentally, while the $D_{s0}^*(2317)$ would be a resonance or a virtual state instead of a quasibound state.

Finally, in Fig. 3 a direct transition of the $K_0^*(704)$ into the $D_0^*(2327)$ is displayed, by letting $m_n \rightarrow m_c$, $M_\pi \rightarrow M_D$ as in Eq. (6), and moreover in a different fashion. Namely, instead of giving the pole positions in the complex energy plane, we now plot the corresponding real and imaginary parts as a function of the varying quark mass, as well as the proportionally changing threshold value. It is striking to see how the $K_0^*(704)$ resonance quickly turns into a virtual bound state, while its real part remains almost constant. Here, we probably see the kinematical Adler zero [22] at work, which rapidly moves away as one of the decay masses increases from M_π , thus allowing the pole to approach the real axis. Then, the pole moves along the real axis as a virtual state, until it touches the threshold at about 1.76 GeV, after which it becomes a bound state. Notice again the tiny margin, at least on this scale, by which the $D_{s0}^*(2327)$ is bound.

To conclude, in the present paper we have shown how several light and heavy scalar mesons can be linked to one another, by continuously varying some of the involved flavor and decay masses. This way, the common dynamical nature of the studied—and probably all—scalar mesons, as ordinary $q\bar{q}$ states but strongly distorted due to coupled channels, is further substantiated. Thus, tetraquarks and other exotic configurations are not needed in this context. Moreover, we deduce that labeling scalar mesons as $q\bar{q}$ states as *opposed* to dynamical meson-meson resonances

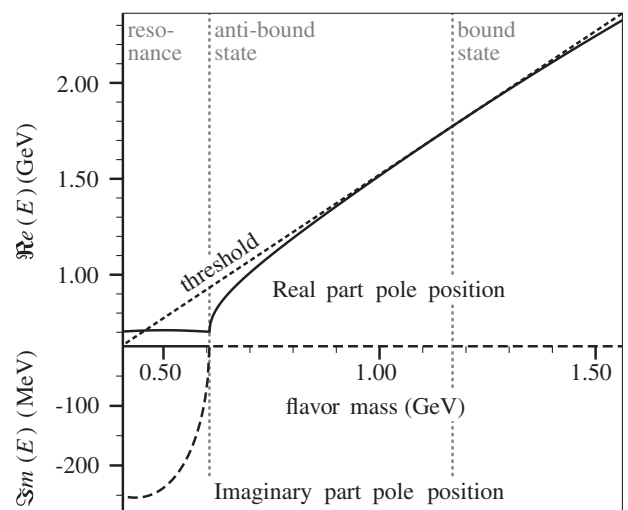


FIG. 3. Real and imaginary parts of the $K_0^*(704)$ pole turning into the $D_{s0}^*(2327)$, as a function of varying quark mass. The straight dashed line stands for decay threshold (real).

makes no sense, in view of the tiny parameter variations needed to turn one kind of pole into another. Rather, scalar mesons should be considered nonperturbatively dressed $q\bar{q}$ systems, with large meson-meson components, no matter if one uses a coupled-channel quark model [23] or e.g. the quark-level linear sigma model [24]. As a consequence, the spectroscopy of scalar mesons is much more complex than for ordinary mesons, with the total number of potentially observable states being different from the number of confined, bare $q\bar{q}$ states.

In the course of this analysis, we have also found a candidate for the new charmonium state $Y(3943)$ [11]. It is true that such a resonance, if indeed a scalar, should

dominantly decay to $D\bar{D}$, a mode which has not been observed yet. However, the reported decay $Y(3943) \rightarrow \omega J/\Psi$ is Okubo-Zweig-Iizuka (OZI)-forbidden, so that it cannot account for the measured sizable width of $\Gamma = 87(\pm 22 \pm 26)$ MeV.

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