THE LIGHT SCALAR $K_0^*(700)$ IN THE VACUUM AND AT NONZERO TEMPERATURE*

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There is mounting evidence toward the existence of a light scalar kaon $\kappa \equiv K_0^*(700)$ with quantum numbers $I(J^P) = 1/2(0^+)$. Here, we recall the results of an effective model with both derivative and non-derivative terms in which only one scalar kaonic field is present in the Lagrangian (the standard quark–antiquark "seed" state $K_0^*(1430)$): a second "companion" pole $K_0^*(700)$ emerges as a dynamically generated state. A related question is the role of $K_0^*(700)$ at nonzero T: since it is the lightest scalar strange state, one would naively expect that it is relevant for π and K multiplicities. However, a repulsion in the πK channel with $I = 3/2$ cancels its effect.

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1. Introduction

The lightest scalar kaonic state listed in the PDG [1] is $K_0^*(700)$ (previously called $K_0^*(800)$, see PDG 2016 [2] and older versions). This state, sometimes called κ, still "*needs confirmation*", but many works do find a pole in that energy region, see Ref. [3] and references therein. The PDG reports at present the following result:

$$\text{pole } \kappa[\text{PDG}] : (630\text{–}730) - i(260\text{–}340) \text{ MeV}\,, \tag{1}$$

(hence, the pole width lies between 520–680 MeV), while the Breit–Wigner (BW) mass and widths are

$$\text{BW [PDG]} : m_{\kappa,\text{BW}} = 824 \pm 30 \text{ MeV}\,, \qquad \Gamma_{\kappa,\text{BW}} = 478 \pm 50 \text{ MeV}\,. \tag{2}$$

The BW and the pole widths are compatible, but the BW mass is somewhat larger. There is, however, no friction, since BW and pole masses are different

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quantities which coincide only when a resonance is narrow. This is definitely not the case for the κ, which is a very broad state with a width-to-mass ratio larger than 0.5.

In a certain sense, the light κ can be regarded as the "brother" of the light $\sigma \equiv f_0(500)$ meson [1]. This state is also very broad and for a long time, it was not clear if there is a pole on the complex plane. Now, its existence is confirmed by many studies and the state is listed in the PDG, see also the review paper [4]. The destiny of the light κ looks somewhat similar: its final confirmation is probably just a matter of time.

Yet, a different issue is the nature of the $\kappa \equiv K_0^*(700)$ and the $\sigma \equiv f_0(500)$. According to mounting evidence, both states are not simple quark–antiquark states, but are rather four-quark objects, either in the form of a tetraquark nonet together with $a_0(980)$ and $f_0(980)$ [5] or as dynamically generated molecular-like states [6]. The κ can be then interpreted as a diquark–antidiquark state ($[u,d][\bar{d},\bar{s}],\dots$) and/or as $K\pi$ state (mixing among these configurations is of course possible and rather probable to occur). If κ is not $\bar{q}q$, where should be the scalar strange quarkonium? According to the quark model [7] and modern chiral approaches [8], the lightest $\bar{q}q$ kaonic state ($u\bar{s},\dots$) is the well-established $K_0^*(1430)$ (similarly, the lightest scalar/isoscalar quarkonium is the state $f_0(1370)$). The question that we review in this work is the link between the standard state $K_0^*(1430)$ and the dynamically generated state $K_0^*(700)$. We find (see Sec. 2) that the πK loops dressing $K_0^*(1430)$ generate $K_0^*(700)$ as a companion pole (a peculiar four-quark object) [9] (similarly, the $a_0(980)$ emerges as a companion pole of $a_0(1450)$ [10]).

There is, however, a related important question: if the light κ is existent, should it be included into thermal hadronic models [11]? At a first sight, the answer is 'yes'. In fact, the light κ is the second-lightest state with nonzero strangeness, thus potentially relevant. Yet, a detailed analysis of the problem [12] shows that one should better *not* include this state into a thermal model (see Sec. 3). Namely, also repulsive channels contribute to the thermodynamics [13–15]. Just as for the $f_0(500)$ whose contribution is cancelled by $\pi\pi$ scattering with $I = 2$, the contribution of the κ is cancelled by the repulsion in πK channel with $I = 3/2$. Thus, the easiest thing to do is to neglect both the $f_0(500)$ and the $K_0^*(700)$ when studying hadronic thermal models for the late stage of heavy-ion collisions.

2. The light κ in the vacuum

As a first step, we write down a Lagrangian that contains *only one* scalar state K_0^*, to be identified with $K_0^*(1430)$, coupled to $K\pi$ pairs

$$\mathcal{L}_{K_0^*} = aK_0^{*+}K^-\pi^0 + bK_0^{*+}\partial_\mu K^-\partial^\mu\pi^0 + \dots\,, \tag{3}$$

where dots refer to other isospin channels. Note, there is *no* $\kappa \equiv K_0^*(700)$ in the model (yet). There are both derivative and non-derivative terms: the former naturally dominates in the context of chiral perturbation theory and also emerge from the extended Linear Sigma Model [8]. The decay width reads

$$\Gamma_{K_0^* \to K\pi}(m) = 3\frac{\left|\vec{k}_1\right|}{8\pi m^2}\left[a - b\frac{m^2 - M_K^2 - M_\pi^2}{2}\right]^2 F_\Lambda(m)\,, \tag{4}$$

with the vertex function $F_\Lambda(m) = \exp(-2\vec{k}_1^2/\Lambda^2)$. Here, Λ is an energy scale describing the nonlocal nature of mesons [16], $\vec{k}_1$ the three-momentum of one outgoing particle, M_K the kaon mass, and M_π the pion mass. (For details and phenomenology of the spectral function, see Ref. [17].)

The propagator of K_0^* is given by $\Delta_{K_0^*}(m^2) = [m^2 - M_0^2 + \Pi(m^2) + i\varepsilon]^{-1}$, M_0 being the bare mass of $K_0^*(1430)$ and $\Pi(m^2)$ the one-loop contribution. The spectral function $d_{K_0^*}(m) = \frac{2m}{\pi}|\mathrm{Im}\Delta_{K_0^*}(p^2 = m^2)|$ is the mass probability density (its integral is normalized to unity). Typically, for the Breit–Wigner value M_{BW} determined as $M_{\mathrm{BW}}^2 - M_0^2 + \mathrm{Re}\,\Pi(M_{\mathrm{BW}}^2) = 0$, the spectral function has a peak's width $\Gamma_{\mathrm{BW}} = \mathrm{Im}\,\Pi(M_{\mathrm{BW}})/M_{\mathrm{BW}}$. A useful approximation, valid if the width is sufficiently small, is the relativistic Breit–Wigner expression

$$d_{K_0^*}(m) \approx d_{K_0^*}^{\mathrm{BW}}(m) = N\left[\left(m^2 - M_{\mathrm{BW}}^2\right)^2 + M_{\mathrm{BW}}^2\Gamma_{\mathrm{BW}}^2\right]^{-1}. \tag{5}$$

Under this approximation, there is only one pole in the complex plane at $m^2 \simeq M_{\mathrm{BW}}^2 - iM_{\mathrm{BW}}\Gamma_{\mathrm{BW}}$ (hence, $m \simeq M_{\mathrm{BW}} - i\Gamma_{\mathrm{BW}}/2$). However, when a resonance is broad, these approximations are not valid anymore.

We now turn to πK scattering. Within our framework, the pion–kaon phase shift is given by [9]

$$\delta_{\pi K,\text{S-wave}}(m) = \delta_{(I=1/2,J=0)}(m) = \tfrac{1}{2}\arccos\left[1 - \pi\Gamma_{K_0^*}(m)d_{K_0^*}(m)\right], \tag{6}$$

where $\delta_{(I,J)}(m)$ is the general phase shift for a given isospin I and total spin J. The amplitude of the process and the phase-shift are linked by $a_{(I,J)} = (e^{i\delta_{(I,J)}(m)} - 1)/(2i)$. The parameters (a, b, M_0, Λ) entering in Eq. (3) were determined via a fit to πK phase-shift data [18], see Ref. [9] for details. A very good description of data is achieved. A study of the complex plane shows an interesting fact: besides the pole corresponding to the well-known $K_0^*(1430)$ state $(1.413 \pm 0.057) - i(0.127 \pm 0.011)$ GeV, there is a second pole which correspond to $K_0^*(700)$

$$(0.745 \pm 0.029) - i(0.263 \pm 0.027)\ \mathrm{GeV}\,. \tag{7}$$

The numerical value is compatible with the PDG value of Eq. (1). A large-$N_{\rm c}$ study confirms that, while the first pole tends to the real axis (and hence is a $\bar{q}q$ state), the second one moves away from it, as it is expected for a dynamically generated state.

In conclusion, the simple model of Eq. (3) is able to describe πK scattering data and naturally gives rise to the pole of $K_0^*(700)$ as a companion pole of the predominantly quark–antiquark resonance $K_0^*(1430)$.

3. The light κ at nonzero temperature

The partition function of an hadronic gas can be expressed as the sum of the contributions of stable particles and their mutual interactions

$$\ln Z = \ln Z_{\rm pions} + \ln Z_{\rm kaons} + \cdots + \ln Z^{\rm int}\,, \qquad \ln Z^{\rm int} = \sum_{I,J} \ln Z_{(I,J)}\,. \quad (8)$$

The first term $\ln Z_{\rm pions} = 3F_1(m_\pi)$ refers to pions and $\ln Z_{\rm kaons} = 4F_1(m)$ to kaons, where $F_1(m) = \int \frac{{\rm d}^3p}{(2\pi)^3} \ln\left[1 - e^{-\sqrt{\vec{p}^{\,2}+m_\pi^2}/T}\right]$ is the contribution of a free particle with mass m. The term $\ln Z_{IJ}$ refers to the contribution of the interactions in the (I,J) channel [13]

$$\ln Z_{(I,J)} = (2I+1)(2J+1)\int_0^\infty \frac{{\rm d}m}{\pi}\frac{{\rm d}\delta_{(I,J)}(m)}{{\rm d}m}F_1(m)\,. \quad (9)$$

When in a certain channel a narrow resonance is present, one finds its standard contribution. For instance, for $I=J=1$, the ρ meson is produced. In the nonrelativistic BW-limit, $\frac{1}{\pi}\frac{{\rm d}\delta_{(1,1)}(m)}{{\rm d}m} \simeq \frac{\Gamma_\rho}{2\pi}[(m-M_\rho)^2+\Gamma_\rho^2/4]^{-1}$. (Moreover, for $\Gamma_\rho \to 0$, $\delta(m-M_\rho)$ emerges: the contribution of a stable ρ is obtained.)

However, Eq. (9) is very general and can describe also broad resonances as well as non-resonant channels, such as repulsive ones. This is important for the κ. In the resonant $I=1/2, J=0$ channel in which the κ is formed, one has (upon integrating up to 1 GeV) $\ln Z_{(1/2,0)} = \int_0^{1\ {\rm GeV}} \frac{2{\rm d}m}{\pi}\frac{{\rm d}\delta_{(1/2,0)}(m)}{{\rm d}m} F_1(m)$. This is sizable. However, one should also consider the repulsion in the $I=3/2$, $J=0$ channel. Remarkably, the sum

$$\ln Z_{(1/2,0)} + \ln Z_{(3/2,0)} = \int_0^{1\ {\rm GeV}} {\rm d}m \left(\frac{2}{\pi}\frac{{\rm d}\delta_{(1/2,0)}(m)}{{\rm d}m} + \frac{4}{\pi}\frac{{\rm d}\delta_{(3/2,0)}(m)}{{\rm d}m}\right)F_1(m) \quad (10)$$

is *small*. Namely, while $\frac{\mathrm{d}\delta_{(1/2,0)}(m)}{\mathrm{d}m} > 0$ (attraction), $\frac{4}{\pi}\frac{\mathrm{d}\delta_{(3/2,0)}(m)}{\mathrm{d}m} < 0$ (repulsion). Note: $\frac{1}{\pi}\frac{\mathrm{d}\delta_{(1/2,0)}(m)}{\mathrm{d}m} \neq d_{K_0^*}(m)$. (This would be true only in the BW-limit). In conclusion, the light κ can be safely neglected in the construction of thermal hadronic models.

4. Conclusions

We have described the emergence of the state $\kappa \equiv K_0^*(700)$ as a companion pole of $K_0^*(1430)$ by using an effective hadronic model [9]. The numerical value of the pole (7) is in agreement with the present PDG estimate of Eq. (1). On the other hand, contrary to the naive expectations, the light κ is not relevant in a thermal hadronic gas. Namely, its influence on thermodynamical properties is cancelled by a repulsion in the $I = 3/2$ channel. Either one includes both the light κ and the repulsion, or — even easier — neglects both of them.

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