

Light Scalars in Field Theory

N.N. Achasov *

Sobolev Institute for Mathematics

Academician Koptiug Prospekt, 4, Novosibirsk, 630090, Russia

Abstract

Outline: 1. Introduction, 2. Confinement, chiral dynamics and light scalar mesons, 3. Chiral shielding of $\sigma(600)$, chiral constraints, $\sigma(600)$, $f_0(980)$ and their mixing in $\pi\pi \rightarrow \pi\pi$, $\pi\pi \rightarrow K\bar{K}$, and $\phi \rightarrow \gamma\pi^0\pi^0$, 4. The ϕ meson radiative decays on light scalar resonances. 5. Why $a_0(980)$ and $f_0(980)$ are not the $K\bar{K}$ molecules. 6. Light scalars in $\gamma\gamma$ collisions.

Evidence for four-quark components of light scalars is given. The priority of Quantum Field Theory in revealing the light scalar mystery is emphasized.

1 Introduction

The scalar channels in the region up to 1 GeV became a stumbling block of QCD. The point is that both perturbation theory and sum rules do not work in these channels because there are not solitary resonances in this region.

At the same time the question on the nature of the light scalar mesons is major for understanding the mechanism of the chiral symmetry realization, arising from the confinement, and hence for understanding the confinement itself.

2 Place in QCD

The QCD Lagrangian is given by

$$L = -\frac{1}{2} \text{Tr} (G_{\mu\nu}(x)G^{\mu\nu}(x)) + \bar{q}(x)(i\hat{D} - M)q(x),$$

M is a diagonal matrix of quark masses, $\hat{D} = \gamma^\mu D_\mu$, $D_\mu = \partial_\mu + ig_0 G_\mu(x)$. M mixes left and right spaces. But in chiral limit, $M_{ff} \rightarrow 0$, these spaces separate realize $U_L(3) \times U_R(3)$ flavour symmetry, which, however, is broken by the gluonic anomaly up to $U_{\text{vec}}(1) \times SU_L(3) \times SU_R(3)$. As experiment suggests, confinement forms colourless observable hadronic fields and spontaneous breaking of chiral symmetry with massless pseudoscalar fields. There are two possible scenarios for QCD at low energy. 1. Non-linear σ model. 2. Linear σ model (LSM). The experimental nonet of the light scalar mesons, $f_0(600)$ (or $\sigma(600)$), $\kappa(700-900)$, $a_0(980)$ and $f_0(980)$ mesons, suggests the $U_L(3) \times U_R(3)$ LSM.

3 History

Hunting the light σ and κ mesons had begun in the sixties already and a preliminary information on the light scalar mesons in PDG Reviews had appeared at that time. But long-standing unsuccessful attempts to prove their existence in a conclusive way entailed general disappointment and an information on these states disappeared from PDG Reviews. One of principal reasons against the σ and κ mesons was the fact that both $\pi\pi$ and πK scattering phase shifts do not pass over 90° at putative resonance masses.

*e-mail address: achasov@math.nsc.ru

4 $SU_L(2) \times SU_R(2)$ LSM, $\pi\pi \rightarrow \pi\pi$ [1, 2, 3]

Situation changes when we showed that in LSM there is a negative background phase which hides the σ meson in $\pi\pi \rightarrow \pi\pi$. It has been made clear that shielding wide lightest scalar mesons in chiral dynamics is very natural. This idea was picked up and triggered new wave of theoretical and experimental searches for the σ and κ mesons. Our approximation is as follows (see Fig.1):

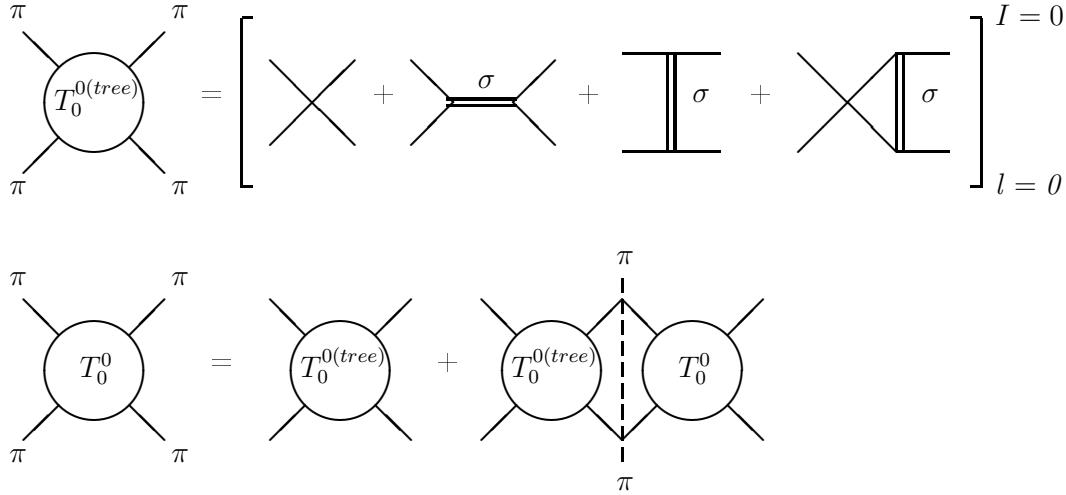


Figure 1. The graphical representation of the S wave $I = 0$ $\pi\pi$ scattering amplitude T_0^0 .

$$T_0^0 = \frac{T_0^{0(tree)}}{1 - i\rho_{\pi\pi}T_0^{0(tree)}} = \frac{e^{2i\delta_0^0} - 1}{2i\rho_{\pi\pi}} = \frac{e^{2i(\delta_{bg} + \delta_{res})} - 1}{2i\rho_{\pi\pi}},$$

$$T_0^2 = \frac{T_0^{2(tree)}}{1 - i\rho_{\pi\pi}T_0^{2(tree)}} = \frac{e^{2i\delta_0^2} - 1}{2i\rho_{\pi\pi}}.$$

5 Results in our approximation [3].

$$M_{res} = 0.43 \text{ GeV}, \Gamma_{res}(M_{res}^2) = 0.67 \text{ GeV}, m_\sigma = 0.93 \text{ GeV},$$

$$\Gamma_{res}(s) = \frac{g_{res}^2(s)}{16\pi\sqrt{s}}\rho_{\pi\pi}, g_{res}(M_{res}^2)/g_{\sigma\pi\pi} = 0.33,$$

$$a_0^0 = 0.18 m_\pi^{-1}, a_0^2 = -0.04 m_\pi^{-1}, (s_A)_0^0 = 0.45 m_\pi^2, (s_A)_0^2 = 2.02 m_\pi^2.$$

6 Chiral shielding in $\pi\pi \rightarrow \pi\pi$ [3]

The chiral shielding of the $\sigma(600)$ meson is illustrated in Fig. 2.

7 The σ pole in $\pi\pi \rightarrow \pi\pi$ [3]

$$T_0^0 \rightarrow g_\pi^2/(s - s_R), g_\pi^2 = (0.12 + i0.21) \text{ GeV}^2,$$

$$\sqrt{s_R} = M_R - i\Gamma_R/2 = (0.52 - i0.25) \text{ GeV}.$$

Considering the residue of the σ pole in T_0^0 as the square of its coupling constant to the $\pi\pi$ channel is not a clear guide to understand the σ meson nature for its great obscure imaginary part.

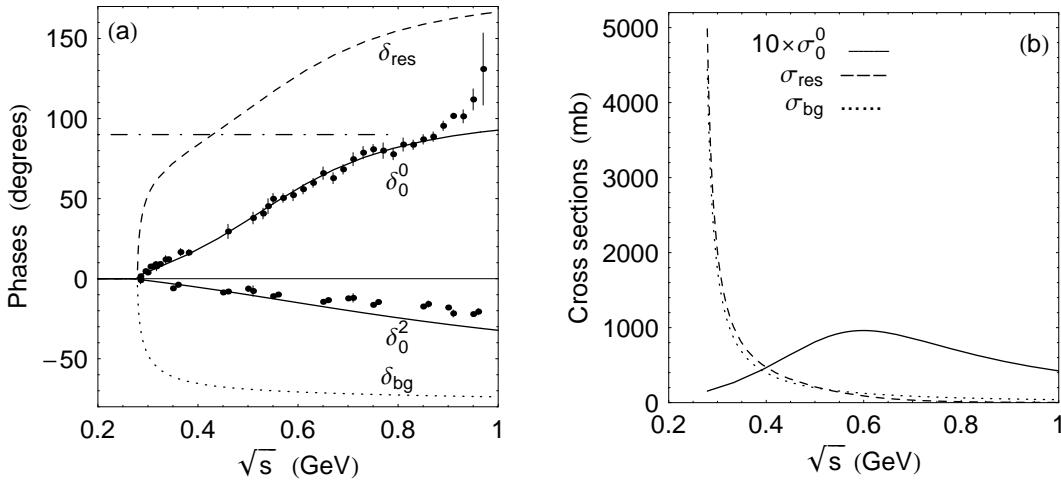


Figure 2. The σ model. Our approximation. $\delta_0^0 = \delta_{res} + \delta_{bg}$. $(\sigma_0^0, \sigma_{res}, \sigma_{bg}) = 32\pi(|T_0^0|^2, |T_{res}|^2, |T_{bg}|^2)/s$.

8 The σ propagator [3]

$$1/D_\sigma(s) = 1/[M_{res}^2 - s + \text{Re}\Pi_{res}(M_{res}^2) - \Pi_{res}(s)].$$

The σ meson self-energy $\Pi_{res}(s)$ is caused by the intermediate $\pi\pi$ states, that is, by the four-quark intermediate states if we keep in mind that the $SU_L(2) \times SU_R(2)$ LSM could be the low energy realization of the two-flavour QCD. This contribution shifts the Breit-Wigner (BW) mass greatly $m_\sigma - M_{res} = 0.50$ GeV. So, half the BW mass is determined by the four-quark contribution at least. The imaginary part dominates the propagator modulus in the region $300 \text{ MeV} < \sqrt{s} < 600 \text{ MeV}$. So, the σ field is described by its four-quark component at least in this energy region.

9 Chiral shielding in $\gamma\gamma \rightarrow \pi\pi$ [3]

$$T_S(\gamma\gamma \rightarrow \pi^+\pi^-) = T_S^{Born}(\gamma\gamma \rightarrow \pi^+\pi^-) + 8\alpha I_{\pi^+\pi^-} T_S(\pi^+\pi^- \rightarrow \pi^+\pi^-),$$

$$T_S(\gamma\gamma \rightarrow \pi^0\pi^0) = 8\alpha I_{\pi^+\pi^-} T_S(\pi^+\pi^- \rightarrow \pi^0\pi^0),$$

$$T_S^{Born}(\gamma\gamma \rightarrow \pi^+\pi^-) = (8\alpha/\rho_{\pi^+\pi^-}) \text{Im} I_{\pi^+\pi^-} I_{\pi^+\pi^-} = \frac{m_\pi^2}{s} (\pi + i \ln \frac{1+\rho_{\pi\pi}}{1-\rho_{\pi\pi}})^2 - 1, \quad s \geq 4m_\pi^2.$$

Our results are shown in Fig. 3.

$$\Gamma(\sigma \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma, s) = \frac{1}{16\pi\sqrt{s}} |g(\sigma \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma, s)|^2,$$

where $g(\sigma \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma, s) = (\alpha/2\pi) \times I_{\pi^+\pi^-} g_{res \pi^+\pi^-}(s)$; see Fig. 4. So, the the $\sigma \rightarrow \gamma\gamma$ decay is described by the triangle $\pi^+\pi^-$ loop diagram $res \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma$. Consequently, it is due to the four-quark transition because we imply a low energy realization of the two-flavour QCD by means of the the $SU_L(2) \times SU_R(2)$ LSM.

As Fig. 4 suggests, the real intermediate $\pi^+\pi^-$ state dominates in $g(res \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma)$ in the σ region $\sqrt{s} < 0.6$ GeV. Thus the picture in the physical region is clear and informative. But, what about the pole in the complex s plane? Does the pole residue reveal the σ indeed?

10 The σ pole in $\gamma\gamma \rightarrow \pi\pi$ [3]

$$\frac{1}{16\pi} \sqrt{\frac{3}{2}} T_S(\gamma\gamma \rightarrow \pi^0\pi^0) \rightarrow g_\gamma g_\pi / (s - s_R),$$

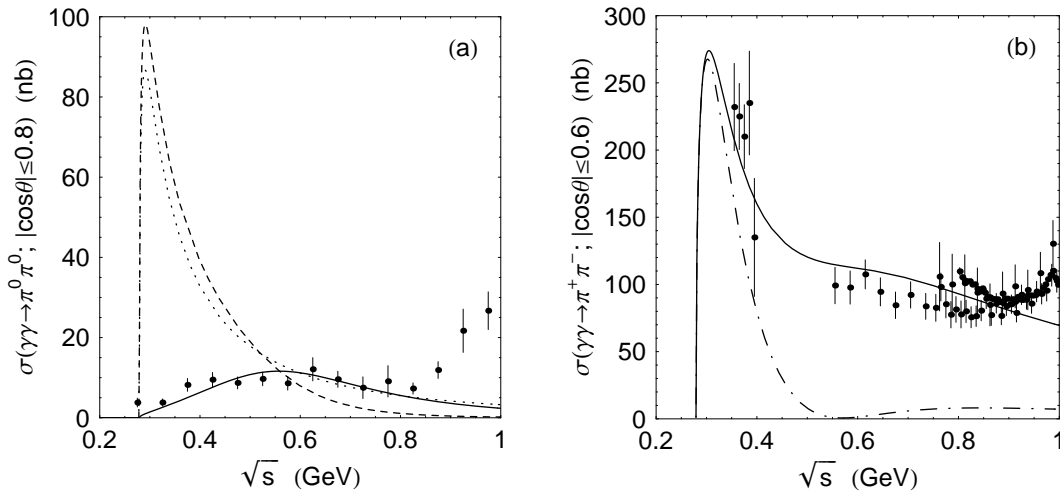


Figure 3. (a) The solid, dashed, and dotted lines are $\sigma_S(\gamma\gamma \rightarrow \pi^0\pi^0)$, $\sigma_{res}(\gamma\gamma \rightarrow \pi^0\pi^0)$, and $\sigma_{bg}(\gamma\gamma \rightarrow \pi^0\pi^0)$. (b) The dashed-dotted line is $\sigma_S(\gamma\gamma \rightarrow \pi^+\pi^-)$, the solid line includes the higher waves from $T^{Born}(\gamma\gamma \rightarrow \pi^+\pi^-)$.

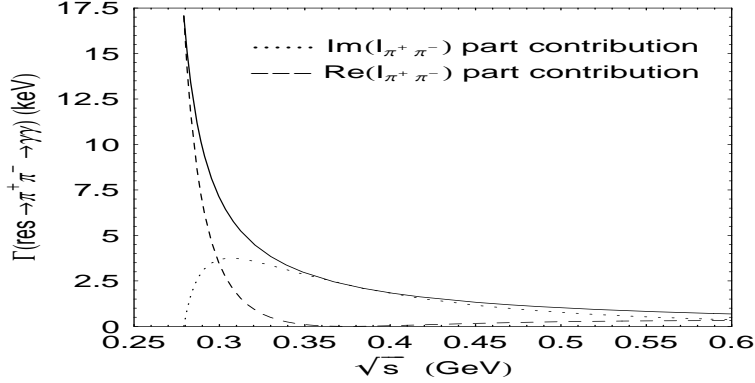


Figure 4. The energy dependent width of the $\sigma \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma$ decay.

$$g_\gamma g_\pi = (-0.45 - i0.19) \cdot 10^{-3} \text{ GeV}^2, \quad g_\gamma/g_\pi = (-1.61 + i1.21) \cdot 10^{-3}, \quad \Gamma(\sigma \rightarrow \gamma\gamma) = \frac{|g_\gamma|^2}{M_R} \approx 2 \text{ keV}.$$

It is hard to believe that anybody could learn the complex but physically clear dynamics of the $\sigma \rightarrow \gamma\gamma$ decay from the residues of the σ pole.

11 Discussion [3, 4, 5, 6]

Leutwyler and collaborators obtained $\sqrt{s_R} = M_R - i\Gamma_R/2 = (441_{-8}^{+16} - i272_{-9}^{+12.5}) \text{ MeV}$. Our result agrees with the above only qualitatively, $\sqrt{s_R} = M_R - i\Gamma_R/2 = (518 - i250) \text{ MeV}$. It is natural, for our approximation.

Could the above scenario incorporates the primary lightest scalar Jaffe four-quark state? Certainly the direct coupling of this state to $\gamma\gamma$ via neutral vector pairs ($\rho^0\rho^0$ and $\omega\omega$), contained in its wave function, is negligible, $\Gamma(q^2\bar{q}^2 \rightarrow \rho^0\rho^0 + \omega\omega \rightarrow \gamma\gamma) \approx 10^{-3} \text{ keV}$, as we showed in 1982 [6]. But its coupling to $\pi\pi$ is strong and leads to $\Gamma(q^2\bar{q}^2 \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma)$ similar to $\Gamma(res \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma)$ in the above Fig. 4.

Let us add to $T_S(\gamma\gamma \rightarrow \pi^0\pi^0)$ the amplitude for the the direct coupling of σ to $\gamma\gamma$ conserving unitarity $T_{direct}(\gamma\gamma \rightarrow \pi^0\pi^0) = sg_{\sigma\gamma\gamma}^{(0)} g_{res}(s) e^{i\delta_{bg}} / D_{res}(s)$. Fitting the $\gamma\gamma \rightarrow \pi^0\pi^0$ data gives a negligible value of $g_{\sigma\gamma\gamma}^{(0)}$, $\Gamma_{\sigma\gamma\gamma}^{(0)} = |M_{res}^2 g_{\sigma\gamma\gamma}^{(0)}|^2 / (16\pi M_{res}) \approx 0.0034 \text{ keV}$, in astonishing agreement with our prediction [6].

12 Phenomenological chiral shielding [7]

$g_{\sigma\pi^+\pi^-}^2/4\pi=0.99 \text{ GeV}^2$, $g_{\sigma K^+K^-}^2/4\pi=2\cdot 10^{-4} \text{ GeV}^2$, $g_{f_0\pi^+\pi^-}^2/4\pi=0.12 \text{ GeV}^2$, $g_{f_0K^+K^-}^2/4\pi=1.04 \text{ GeV}^2$. The BW masses and width: $m_{f_0}=989 \text{ MeV}$, $m_\sigma=679 \text{ MeV}$, $\Gamma_\sigma=498 \text{ MeV}$. The $l=I=0$ $\pi\pi$ scattering length $a_0^0=0.223 m_{\pi^+}^{-1}$. Figure 5 illustrates the excellent agreement our phenomenological treatment with the experimental and theoretical data.

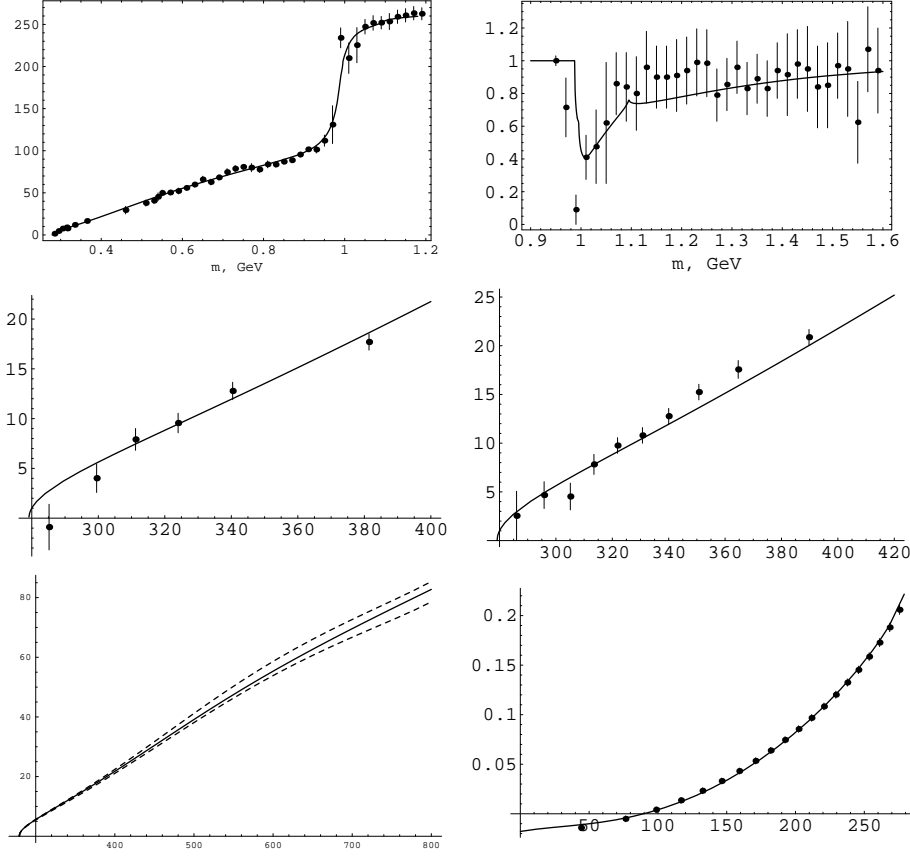


Figure 5. The phenomenological chiral shielding. $\delta_0^0=\delta_B^{\pi\pi}+\delta_{res}$. The comparison with the CERN-Munich data for δ_0^0 and inelasticity η_0^0 , with the BNL and NA48 data for δ_0^0 , with the CGL band for δ_0^0 [8], and with Leutwyler's calculation of T_0^0 for $s < 4m_\pi^2$, respectively.

13 Four-quark model [5, 9, 10, 11]

There are numerous evidences in favour of the $q^2\bar{q}^2$ structure of $f_0(980)$ and $a_0(980)$. As for the nonet as a whole, even a dope's look at PDG Review gives an idea of the four-quark structure of the light scalar meson nonet, $\sigma(600)$, $\kappa(700-900)$, $a_0(980)$, and $f_0(980)$, inverted in comparison with the classical P wave $q\bar{q}$ tensor meson nonet $f_2(1270)$, $a_2(1320)$, $K_2^*(1420)$, and $f_2'(1525)$. Really, it can be easy understood for the $q^2\bar{q}^2$ nonet, where $\sigma(600)$ has no strange quarks, $\kappa(700-900)$ has the s quark, $a_0(980)$ and $f_0(980)$ have the $s\bar{s}$ pair.

14 Radiative decays of ϕ meson [7, 9, 10, 11, 12, 13, 14]

Twenty years ago we showed [12] that the study of the radiative decays $\phi \rightarrow \gamma a_0 \rightarrow \gamma\pi\eta$ and $\phi \rightarrow \gamma f_0 \rightarrow \gamma\pi\pi$ can shed light on the problem of $a_0(980)$ and $f_0(980)$ mesons. Now these decays have been studied not only theoretically but also experimentally. Note that $a_0(980)$

is produced in the radiative ϕ meson decay as intensively as $\eta'(958)$ containing $\approx 66\%$ of $s\bar{s}$, responsible for $\phi \approx s\bar{s} \rightarrow \gamma s\bar{s} \rightarrow \gamma\eta'(958)$. It is a clear qualitative argument for the presence of the $s\bar{s}$ pair in the isovector $a_0(980)$ state, i.e., for its four-quark nature.

15 K^+K^- loop mechanism [7, 9, 10, 11, 12, 13, 14]

When basing the experimental investigations, we suggested [12] one-loop model $\phi \rightarrow K^+K^- \rightarrow \gamma[a_0(980)/f_0(980)]$; see Fig. 6.

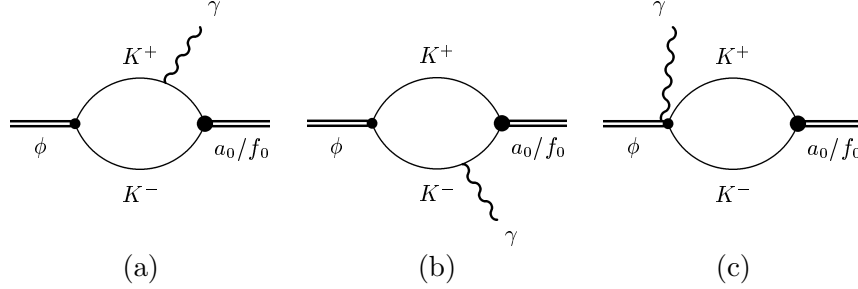


Figure 6. The K^+K^- loop model.

This model is used in the data treatment and is ratified by experiment, see Figs. 7 (a) and 7 (b).

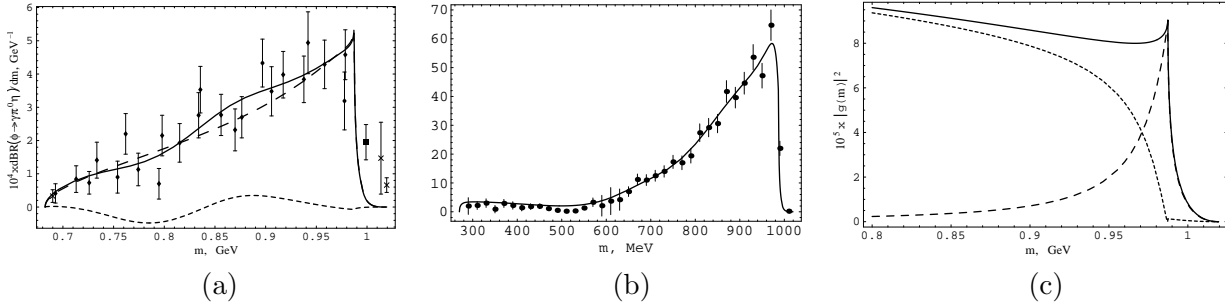


Figure 7. (a) The fit to the KLOE data for the $\pi^0\eta$ mss spectrum in the $\phi \rightarrow \gamma\pi^0\eta$ decay. (b) The fit to the KLOE data for the $\pi^0\pi^0$ mss spectrum in the $\phi \rightarrow \gamma\pi^0\pi^0$ decay. (c) The universal in the K^+K^- loop model function $|g(m)|^2 = |g_R(m)/g_{RK^+K^-}|^2$ is shown by the solid curve. The contribution of the imaginary (real) part is shown by dashed (dotted) curve.

For $dBR[\phi \rightarrow \gamma(a_0/f_0) \rightarrow \gamma(\pi^0\eta/\pi^0\pi^0), m]/dm \sim |g(m)|^2\omega(m)$, the function $|g(m)|^2$ should be smooth at $m \leq 0.99$ GeV. But gauge invariance requires that $g(m)$ is proportional to the photon energy $\omega(m)$. Stopping the function $(\omega(m))^3$ at $\omega(990\text{ MeV})=29$ MeV is the crucial point. The K^+K^- loop model solves this problem in the elegant way, see Fig. 7 (c).

16 Four-quark transition and OZI rule [10]

So, we are dealing here with the four-quark transition. A radiative four-quark transition between two $q\bar{q}$ states requires creation and annihilation of an additional $q\bar{q}$ pair, i.e., is forbidden according to OZI rule, while a radiative four-quark transition between $q\bar{q}$ and $q^2\bar{q}^2$ states requires only creation of an additional $q\bar{q}$ pair, i.e., is allowed according to the OZI rule. The large N_C expansion support this conclusion.

17 Why $a_0(980)$ and $f_0(980)$ are not the $K\bar{K}$ molecules [15]

Every diagram contribution in

$$T\{\phi(p) \rightarrow \gamma[a_0(q)/f_0(q)]\} = (a) + (b) + (c) \text{ (see Fig. 6)}$$

is divergent and hence should be regularized in a gauge invariant manner, for example, in the Pauli-Villars one.

$$\overline{T}\{\phi(p) \rightarrow \gamma[a_0(q)/f_0(q)], M\} = \overline{(a)} + \overline{(b)} + \overline{(c)},$$

$$\overline{T}\{\phi(p) \rightarrow \gamma[a_0(q)/f_0(q)], M\} = \epsilon^\nu(\phi)\epsilon^\mu(\gamma)\overline{T}_{\nu\mu}(p, q) = \epsilon^\nu(\phi)\epsilon^\mu(\gamma) [\overline{a}_{\nu\mu}(p, q) + \overline{b}_{\nu\mu}(p, q) + \overline{c}_{\nu\mu}(p, q)],$$

$$\overline{a}_{\nu\mu}(p, q) = -\frac{i}{\pi^2} \int \left\{ \frac{(p-2r)_\nu(p+q-2r)_\mu}{(m_K^2-r^2)[m_K^2-(p-r)^2][m_K^2-(q-r)^2]} - \frac{(p-2r)_\nu(p+q-2r)_\mu}{(M^2-r^2)[M^2-(p-r)^2][M^2-(q-r)^2]} \right\} dr,$$

$$\overline{b}_{\nu\mu}(p, q) = -\frac{i}{\pi^2} \int \left\{ \frac{(p-2r)_\nu(p-q-2r)_\mu}{(m_K^2-r^2)[m_K^2-(p-r)^2][m_K^2-(p-q-r)^2]} - \frac{(p-2r)_\nu(p-q-2r)_\mu}{(M^2-r^2)[M^2-(p-r)^2][M^2-(p-q-r)^2]} \right\} dr,$$

$$\overline{c}_{\nu\mu}(p, q) = -\frac{i}{\pi^2} 2g_{\nu\mu} \int dr \left\{ \frac{1}{(m_K^2-r^2)[m_K^2-(q-r)^2]} - \frac{1}{(M^2-r^2)[M^2-(q-r)^2]} \right\},$$

where M is the regulator field mass. $M \rightarrow \infty$ in the end

$$\overline{T}[\phi \rightarrow \gamma(a_0/f_0), M \rightarrow \infty] \rightarrow T^{Phys}[\phi \rightarrow \gamma(a_0/f_0)].$$

We can shift the integration variables in the regularized amplitudes and easily check the gauge invariance condition

$$\epsilon^\nu(\phi)k^\mu\overline{T}_{\nu\mu}(p, q) = \epsilon^\nu(\phi)(p-q)^\mu\overline{T}_{\nu\mu}(p, q) = 0.$$

It is instructive to consider how the gauge invariance condition

$$\epsilon^\nu(\phi)\epsilon^\mu(\gamma)\overline{T}_{\nu\mu}(p, p) = 0$$

holds true,

$$\epsilon^\nu(\phi)\epsilon^\mu(\gamma)\overline{T}_{\nu\mu}(p, p) = \epsilon^\nu(\phi)\epsilon^\mu(\gamma)T_{\nu\mu}^{m_K}(p, p) - \epsilon^\nu(\phi)\epsilon^\mu(\gamma)T_{\nu\mu}^M(p, p) = (\epsilon(\phi)\epsilon(\gamma))(1-1) = 0.$$

The superscript m_K refers to the nonregularized amplitude and the superscript M refers to the regulator field amplitude. So, the contribution of the (a), (b), and (c) diagrams does not depend on a particle mass in the loops (m_K or M) at $p = q$ ¹. But, the physical meaning of these contributions is radically different. The $\epsilon^\nu(\phi)\epsilon^\mu(\gamma)T_{\nu\mu}^{m_K}(p, p)$ contribution is caused by intermediate momenta (a few GeV) in the loops, whereas the regulator field contribution is caused fully by high momenta ($M \rightarrow \infty$) and teaches us how to allow for high K virtualities in a gauge invariant way.

Needless to say the integrand of $\epsilon^\nu(\phi)\epsilon^\mu(\gamma)\overline{T}_{\nu\mu}(p, p)$ is not equal to 0.

It is clear that

$$\epsilon^\nu(\phi)\epsilon^\mu(\gamma)T_{\nu\mu}^{M \rightarrow \infty}(p, q) \rightarrow \epsilon^\nu(\phi)\epsilon^\mu(\gamma)T_{\nu\mu}^{M \rightarrow \infty}(p, p) \equiv \epsilon^\nu(\phi)\epsilon^\mu(\gamma)T_{\nu\mu}^M(p, p) \equiv (\epsilon(\phi)\epsilon(\gamma)).$$

So, the regulator field contribution tends to the subtraction constant when $M \rightarrow \infty$.

The finiteness of the subtraction constant hides its high momentum origin and sometimes gives rise to an illusion of a nonrelativistic physics in the K^+K^- model with the pointlike interaction.

¹A typical example of such integrals is $2 \int_0^\infty \frac{m^2 x}{(x+m^2)^3} dx = 1$.

18 $a_0(980)/f_0(980) \rightarrow \gamma\gamma$ & $q^2\bar{q}^2$ model

Twenty six years ago we predicted [6] the suppression of $a_0(980)/f_0(980) \rightarrow \gamma\gamma$ decays basing on $q^2\bar{q}^2$ model. Experiment supported this prediction. The $a_0 \rightarrow K^+K^- \rightarrow \gamma\gamma$ model [16] describes adequately data and corresponds to the four-quark transition $a_0 \rightarrow q^2\bar{q}^2 \rightarrow \gamma\gamma$.

$$\langle \Gamma(a_0 \rightarrow K^+K^- \rightarrow \gamma\gamma) \rangle \approx 0.3 \text{ keV}, \Gamma_{a_0 \rightarrow \gamma\gamma}^{\text{direct}} \ll 0.1.$$

19 $\gamma\gamma \rightarrow \pi\pi$ from Belle [17, 18]

Recently, we analyzed the new high statistics Belle data on the reactions $\gamma\gamma \rightarrow \pi\pi$ and clarified the current situation around the $\sigma(600)$, $f_0(980)$, and $f_2(1270)$ resonances in $\gamma\gamma$ collisions, see Fig. 9.

$$\langle \Gamma(\sigma \rightarrow \pi^+\pi^- \rightarrow \gamma\gamma) \rangle \approx 0.45 \text{ keV}, \langle \Gamma(f_0 \rightarrow K^+K^- \rightarrow \gamma\gamma) \rangle \approx 0.2 \text{ keV},$$

$$\Gamma_{\sigma \rightarrow \gamma\gamma}^{\text{direct}} \ll 0.1 \text{ keV}, \Gamma_{f_0 \rightarrow \gamma\gamma}^{\text{direct}} \ll 0.1 \text{ keV}.$$

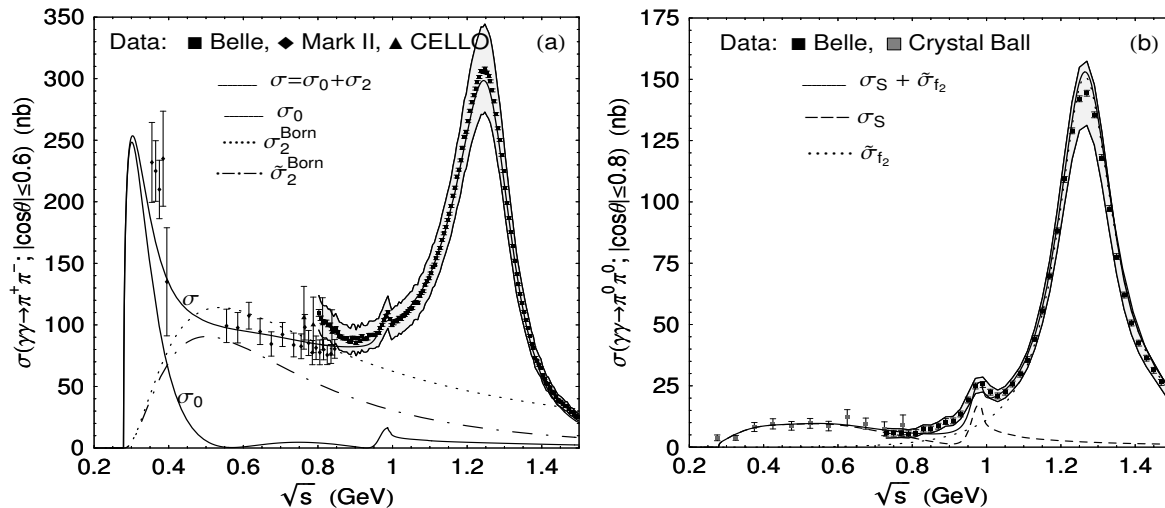


Figure 8. (a) Cross section for $\gamma\gamma \rightarrow \pi^+\pi^-$. (b) Cross section for $\gamma\gamma \rightarrow \pi^0\pi^0$.

20 The lessons

The majority of current investigations of the mass spectra in scalar channels do not study particle production mechanisms. That is why such investigations are nothing more than pre-processing experiments, and the derivable information is very relative. And only the progress in understanding the particle production mechanisms can essentially advance us in revealing the light scalar meson nature, as it is evident from the foregoing.

Theoretical investigations of light scalar mesons, using effective lagrangians in tree level approximation, are very preliminary ones if not exercises. Real investigations require describing real processes $\pi\pi \rightarrow \pi\pi$, $\gamma\gamma \rightarrow \pi\pi$, $\phi \rightarrow \gamma\pi\pi$, $\phi \rightarrow \gamma\pi^0\eta$, and so on, as it is evident from the foregoing.

The K^+K^- loop model, $\phi \rightarrow K^+K^- \rightarrow \gamma(f_0/a_0) \rightarrow \gamma(\pi\pi/\pi^0\eta)$, describes the relativistic physics and strongly supports a compact $q^2\bar{q}^2$ nature of $f_0(980)$ and $a_0(980)$. Consideration of

Section 17 is transferred to the $(f_0/a_0) \rightarrow K^+K^- \rightarrow \gamma\gamma$ decays easily. So, there are no grounds to speak about a nonrelativistic molecular nature of $f_0(980)$ and $a_0(980)$.

The classic P wave $q\bar{q}$ tensor mesons $f_2(1270)$, $a_2(1320)$, and $f_2'(1525)$ are produced by the direct transitions $\gamma\gamma \rightarrow q\bar{q}$ in the main, whereas the light scalar mesons $\sigma(600)$, $f_0(980)$, and $a_0(980)$ are produced by the rescattering $\gamma\gamma \rightarrow \pi^+\pi^- \rightarrow \sigma$, $\gamma\gamma \rightarrow K^+K^- \rightarrow f_0$, and $\gamma\gamma \rightarrow K^+K^- \rightarrow a_0$. The direct transitions $\gamma\gamma \rightarrow \sigma$, $\gamma\gamma \rightarrow f_0$, and $\gamma\gamma \rightarrow a_0$ are negligible, as it is expected in four-quark model.

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