

## Exotic mesons, locking states, and their role in the formation of the confinement barrier

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We discuss a mechanism by which a set of  $q\bar{q}$  states mixes with an exotic one, so that the exotic state accumulates the widths of the overlapping  $\bar{q}q$  resonances. The broad state created this way acts as a “locking” state for the others. Using results of a previous analysis of the  $(IJ^{PC}=00^{++})$  wave, we estimate the mean radius squared of the broad state  $f_0(1530_{-250}^{+90})$ ; it appears to be distinctly larger than the mean squared radii of the locked narrow states. This supports the idea about the constructive role of broad states in forming the confinement barrier. [S0556-2821(98)50315-3]

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The objective of this Rapid Communication is to introduce an idea having two components. The idea is that an exotic meson resonance that overlaps the standard  $q\bar{q}$  resonances may become broad via mixing and in turn reduce the widths of the other states. This mixing is not suppressed because of the absence of the orthogonality of exotic and  $q\bar{q}$  states; the mixing is also not suppressed according to the rules of  $1/N$  expansion [1]. The broad resonance, descendant of the exotic state, thus acts as a confinement barrier for narrow ones. We call the broad state a “locking” state. We draw attention to experimental evidence for this idea and suggest that the formation of broad states may be a general phenomenon involving exotic states: the existence of a broad state and narrow ones in the same mass region should be considered as a signature of exotics.

### EFFECT OF ACCUMULATION OF WIDTHS IN THE $K$ -MATRIX APPROACH

The mechanism of accumulation of the widths of overlapping resonances was discussed for decades (see, for example, Ref. [2]). To illustrate the mixing of nonstable states in a pure form, let us consider an example of  $K$ -matrix amplitude with three resonances decaying into the same channel. An application of the  $K$ -matrix approach to meson resonances in the mass range up to 1900 MeV was performed in Ref. [3]; in Refs. [4,5] it was shown that the use of the  $N/D$ -amplitude, which takes into account the real parts of the transition loop diagrams, does not change qualitatively the picture of mixing.

In the  $K$ -matrix approach, the three-resonance amplitude we consider reads

$$A = K(1 - i\rho K)^{-1}, \quad K = g^2 \sum_{a=1,2,3} (M_a^2 - s)^{-1}. \quad (1)$$

Here, for the purpose of illustration, we take  $g^2$  to be the same for all three resonances, and make the approximations that (i) the phase space factor  $\rho$  is constant, and (ii)  $M_1^2 = m^2 - \delta$ ,  $M_2^2 = m^2$ ,  $M_3^2 = m^2 + \delta$ . Figure 1 shows the location of poles in the complex- $M$  plane ( $M = \sqrt{s}$ ) as the coupling  $g$  increases. At large  $g$ , which corresponds to a strong overlapping of the resonances, one resonance accumulates the widths of the others, while two counterparts of the broad state become nearly stable.

### RESONANCE STRUCTURE IN THE $(IJ^{PC}=00^{++})$ WAVE

This simple model may be compared with what is known about the actual situation for the  $\pi\pi S$  wave. The analysis of

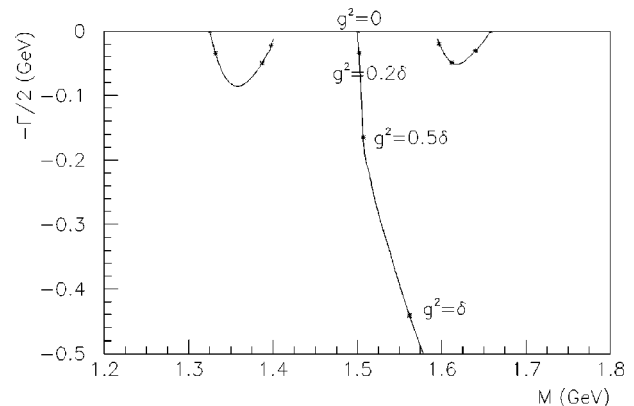


FIG. 1. Position of the poles of the amplitude of Eq. (1) in the complex  $\sqrt{s}$  plane ( $\sqrt{s} = M - i\Gamma/2$ ) with an increase of  $g^2$ . In this example  $m = 1.5$  GeV,  $\delta = 0.5$  GeV<sup>2</sup>, and the phase space factor is fixed,  $\rho = 1$ .

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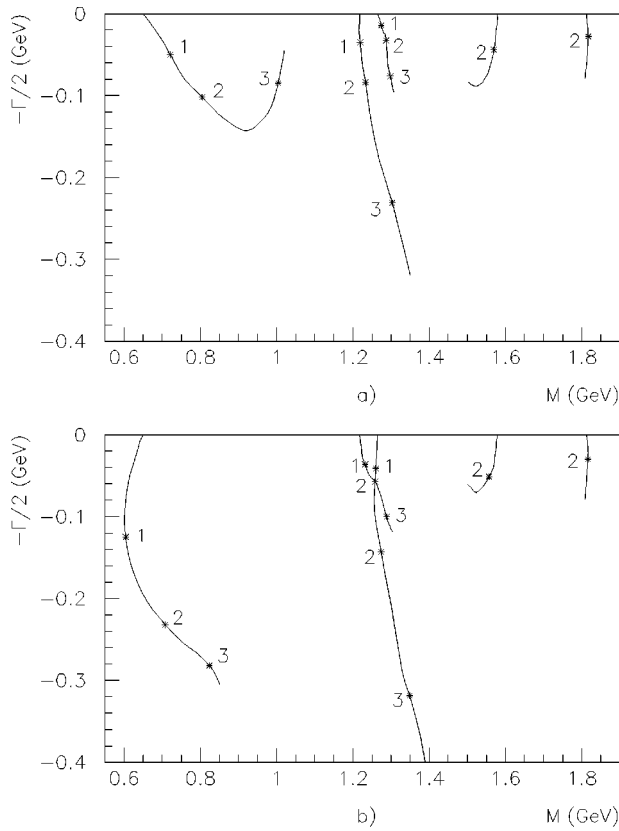


FIG. 2. The pole positions of the  $(00^{++})$  amplitude in the complex  $\sqrt{s}$  plane ( $\sqrt{s}=M-i\Gamma/2$ ), varying  $\xi$ . (a) the sheet under  $\pi\pi$  and  $\pi\pi\pi\pi$  cuts, (b) the sheet under  $\pi\pi$ ,  $\pi\pi\pi\pi$ , and  $K\bar{K}$  cuts. The crosses 1,2,3 indicate the values of  $\xi$ : 0.4, 0.6, and 0.9.

this wave has been pursued in a series of papers using a variety of technical approaches, the  $T$  matrix [6] and multi-channel  $K$ -matrix approaches [3] and the dispersion relation  $N/D$  method [4,5]. The most recent  $K$ -matrix analysis was done using data from GAMS [7], the Crystal Barrel Collaboration [8], and the BNL group [9]. It allows us to determine the resonance structure in this wave in the mass range up to 1900 MeV. There are five states in this mass region. Four of them may be identified as  $q\bar{q}$  states, members of  $1^3P_0q\bar{q}$  and  $2^3P_0q\bar{q}$  nonets; the fifth state in the mass region 1250–1650 MeV is an extra one not accommodated by  $q\bar{q}$  systematics. The restored coupling constants for decays into channels  $\pi\pi$ ,  $K\bar{K}$ ,  $\eta\eta$ , and  $\eta\eta'$  show that the extra state has at least a 40%–50% admixture from the lightest scalar glueball.

Parameters found in [3] for the  $K$  matrix  $0^{++}$  amplitude allow a direct comparison of the mixing dynamics with the simple model of Fig. 1. We introduce a scaling parameter  $\xi$  into the  $K$ -matrix elements

$$K_{ab} = \sum_{\alpha} \frac{g_a^{(\alpha)} g_b^{(\alpha)}}{M_{\alpha}^2 - s} + f_{ab} \rightarrow \sum_{\alpha} \xi \frac{g_a^{(\alpha)} g_b^{(\alpha)}}{M_{\alpha}^2 - s} + \xi f_{ab}, \quad (2)$$

and vary  $\xi$  in the interval 0–1. Figures 2(a) and 2(b) show the movement of the poles for two solutions, I and II, found in Ref. [3]. As  $\xi \rightarrow 0$ , the decay processes and the corre-

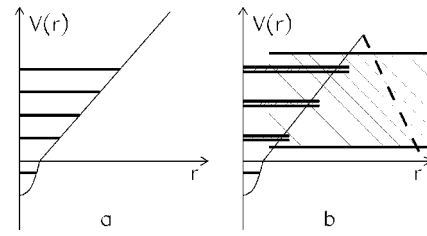


FIG. 3. Conventional pictures for potential model levels (a) without decay processes taken into account, (b) with decay processes taken into account.

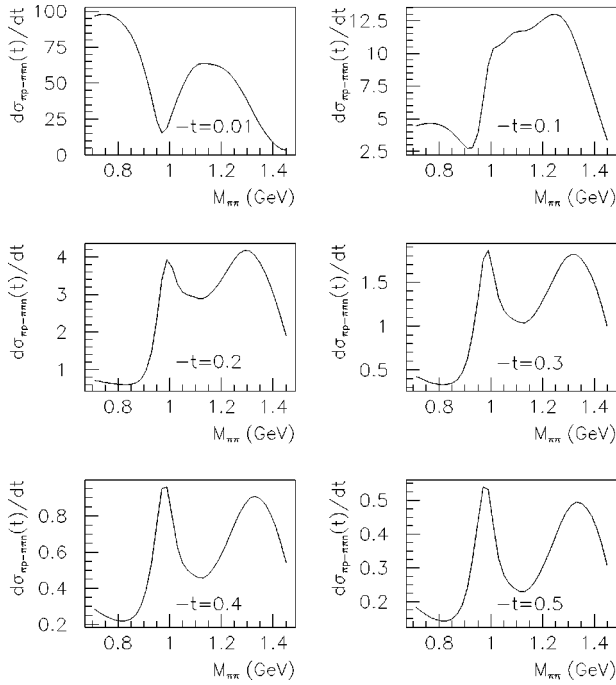
sponding mixing are switched off: The positions of  $K$ -matrix poles show the masses of the bare states  $f_0^{\text{bare}}(720 \pm 10)$ ,  $f_0^{\text{bare}}(1230 \pm 50)$ ,  $f_0^{\text{bare}}(1260 \pm 30)$ ,  $f_0^{\text{bare}}(1600 \pm 50)$ , and  $f_0^{\text{bare}}(1810 \pm 30)$ . In solution I, the gluonium state is associated with  $f_0^{\text{bare}}(1230 \pm 50)$ , while it is  $f_0^{\text{bare}}(1600 \pm 50)$  for solution II. The case  $\xi=1$  shows the pole positions for the physical states  $f_0(980)$ ,  $f_0(1300)$ ,  $f_0(1500)$ ,  $f_0(1530_{-250}^{+90})$ , and  $f_0(1750)$ . In both solutions, according to [4,5], the broad resonance  $f_0(1530_{-250}^{+90})$  carries roughly 40%–50% of the gluonium component. This should not come as a surprise. Stationary  $q\bar{q}$  states are orthogonal to one another (though the loop diagrams involving their decay cause some mixing), while the glueball may mix freely with  $q\bar{q}$  states; the rules of the  $1/N$  expansion [1] tell us that  $q\bar{q}$ /gluonium mixing is not suppressed, see [5,10]. That is the reason for an accumulation of the widths of the neighboring  $q\bar{q}$  states by the gluonium.

### BROAD RESONANCE AS A LOCKING STATE

We now discuss the role of the broad resonance in the formation of the confinement barrier. Let us consider as a guide the meson spectrum in the standard quark model with decay processes switched off [Fig. 3(a)]. Here we have a set of stable levels. If the decay processes are incorporated into highly excited states it would be naive to think that the decay processes result only in the broadening of levels. Because of the processes *bare state*  $\rightarrow$  *real mesons*, the resonances mix and one of them may transform into a very broad state. This creates a trap for the states with which it overlaps.

Figure 3(a) shows the conventional potential model with a linear potential rising indefinitely. That is unrealistic, since it does not allow decays. In Fig. 3(b) we show instead a confinement barrier through which states may decay, thus creating a broad locking state. The broad resonance prevents decay of other states, which are left in the small- $r$  region, the broad state plays the role of a dynamical barrier. This is a familiar phenomenon: an absorption process acts as a reflecting barrier. It means that comparatively narrow locked states lie inside the confinement well, while the broad locking state appears mainly outside the well. Experimental data confirm this idea.

Direct experimental evidence that the broad  $0^{++}$  state is associated with large  $r$  is given by the GAMS data on  $\pi^- \pi^+(t) \rightarrow \pi^0 \pi^0$  [7]. The broad state is clearly visible at  $|t| < 0.2 \text{ GeV}^2$ , but disappears at large  $|t|$ , leaving peaks re-

FIG. 4. Mass spectra at different  $t$  ( $\text{GeV}^2$ ).

lated to narrower states clearly visible. This effect is demonstrated in Fig. 4, which depicts  $|A_{\pi\pi(t)\rightarrow\pi\pi}(M)|^2$  found in [3] for different  $t$ .

By fitting these data we determine the ratios of  $t$  distributions for narrow resonances and background. We find the ratios

$$\frac{d\sigma_{\pi\rightarrow f_0(980)}(t)}{dt} \bigg/ \frac{d\sigma_{\pi\rightarrow f_0(1530_{-250}^{+90})}(t)}{dt},$$

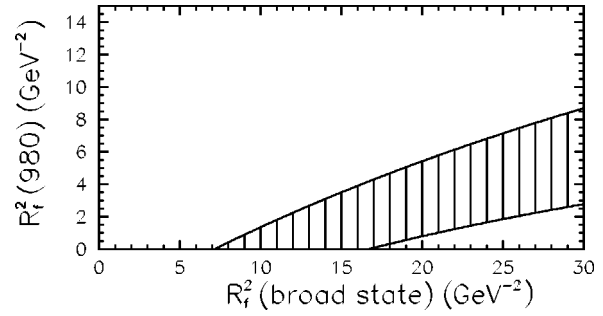
$$\frac{d\sigma_{\pi\rightarrow f_0(1300)}(t)}{dt} \bigg/ \frac{d\sigma_{\pi\rightarrow f_0(1530_{-250}^{+90})}(t)}{dt} \quad (3)$$

as functions of  $t$ . Assuming an universal  $t$ -exchange mechanism for all  $f_0$  mesons, these ratios are given by ratios of the transition form factors squared:

$$F_{\pi\rightarrow f_0(980)}^2(t)/F_{\pi\rightarrow f_0(1530_{-250}^{+90})}^2(t),$$

$$F_{\pi\rightarrow f_0(1300)}^2(t)/F_{\pi\rightarrow f_0(1530_{-250}^{+90})}^2(t). \quad (4)$$

A method for calculating transition form factors at small and moderate momentum transfers is given in [11]. Following it we estimate mean radii squared of the  $f_0$  mesons using a simple parametrization of the wave functions in an exponential form (exponential approximation works sufficiently well for wave functions of mesons in the mass below 1500 MeV). Another simplification follows: in the triangle diagram, which is responsible for the transition form factor at moderately small  $|t|$ , we approximate the light cone variable energy squared as  $s = m^2 + k_{\perp}^2/x(1-x) \approx 4(m^2 + k_{\perp}^2)$  (at small  $|t|$  the region  $x \approx \frac{1}{2}$  gives the main contribution). Then

FIG. 5. The correlation between radius squared of  $f_0(980)$  and broad state  $f_0(1530_{-250}^{+90})$ .

$$\frac{d}{dt} \ln \left[ \frac{d\sigma_{\pi\rightarrow f_0}(t)}{dt} \bigg/ \frac{d\sigma_{\pi\rightarrow f_0(1530_{-250}^{+90})}(t)}{dt} \right]$$

$$= \frac{d}{dt} \ln [F_{\pi\rightarrow f_0}^2(t)/F_{\pi\rightarrow f_0(1530_{-250}^{+90})}^2(t)]$$

$$= \frac{1}{3} (R_{\pi\rightarrow f_0}^2 - R_{\pi\rightarrow f_0(1530_{-250}^{+90})}^2), \quad (5)$$

where the transition radius squared  $R_{\pi\rightarrow f_0}^2$  is determined as

$$R_{\pi\rightarrow f_0}^2 = \frac{2R_{\pi}^2 R_{f_0}^2}{\frac{5}{3} R_{\pi}^2 + R_{f_0}^2}. \quad (6)$$

$R_{\pi}^2$ ,  $R_{f_0}^2$  are pion and  $f_0$ -meson radii squared. Constituent quarks of the pion and  $f_0$  meson are correspondingly in  $S$  and  $P$  waves: that results in the factor  $5/3$  in the denominator of Eq. (6).

A fit to data for the ratios of Eq. (3) gives

$$R_{\pi\rightarrow f_0(980)}^2 \approx R_{\pi\rightarrow f_0(1300)}^2, \quad \text{or} \quad R_{f_0(980)}^2 \approx R_{f_0(1300)}^2, \quad (7)$$

and

$$R_{\pi\rightarrow f_0(1530_{-250}^{+90})}^2 - R_{\pi\rightarrow f_0(980)}^2 \approx (8 \pm 2) \text{ GeV}^{-2}. \quad (8)$$

The last equality defines the correlation between  $R_{f_0(1530_{-250}^{+90})}^2$  and  $R_{f_0(980)}^2$ , provided the mean pion radius squared is fixed. Figure 5 demonstrates this correlation for  $R_{\pi}^2 = 0.41 \text{ fm}^2$ . It is clearly seen that  $R_{f_0(1530_{-250}^{+90})}^2$  is distinctly larger than  $R_{f_0(980)}^2$  or  $R_{f_0(1300)}^2$ ; it means that a broad state at large  $r$  can definitely play the role of the locking state.

The same effect may take place for other exotic hadrons. By this we mean glueballs and hybrids with different quantum numbers. One example is now known for  $J^P = 0^{-+}$  [12], where a very broad  $0^{-+}$  state around 1800–2100 MeV ( $\Gamma \approx 1 \text{ GeV}$ ) has been identified in radiative  $J/\psi$  decays. Branching ratios for this state to  $\rho\rho$ ,  $\omega\omega$ ,  $K^*K^*$ , and  $\phi\phi$  channels are in approximate agreement with flavor-blindness, suggesting once again a strong glueball component. The paper [12] also advances arguments for a broad

$2^{++}$  resonance at 2000–2400 MeV; we have been involved in the analyses of two sets of experimental data, to be published shortly, providing evidence for this broad  $2^{++}$  state.

### CONCLUSIONS

In the deconfinement of quarks of an excited  $q\bar{q}$  level, there are two stages.

(i) An inevitable creation of new quark-antiquark pairs that results in the production of white hadrons. This stage is the subject of QCD and is beyond our present discussion.

(ii) The outflow of the created white hadrons and their mixing results in the production of a very broad state. The broad resonance locks other  $q\bar{q}$  levels into the small- $r$  region, thus playing the role of a dynamical barrier; this is the

reason for calling the broad resonance a locking state.

The bare states are the subjects of the quark and gluon classification. Exotic hadrons like glueballs and hybrids mix readily with  $q\bar{q}$  states and are good candidates to generate locking states in all waves. Rich physics is hidden in the broad states, and an investigation of them is an important and unavoidable step in understanding the spectroscopy of highly excited states and their confinement.

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