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Letter

Why the high-lying glueball does not mix with the neighbouring f_0

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Abstract. Chiral symmetry restoration in high-lying hadron spectra implies that hadrons which belong to different irreducible representations of the parity-chiral group cannot mix. This explains why the $f_0(2102 \pm 13)$, which was suggested to be a glueball, and hence must belong to the scalar (0,0) representation of the chiral group, cannot mix with the neighbouring $f_0(2040 \pm 38)$, which was interpreted as a $n\bar{n}$ state, and that belongs to the $(1/2, 1/2)$ representation of the chiral group. If confirmed, then we have an access to a "true" glueball of QCD.

PACS. 12.39.Mk Glueball and nonstandard multi-quark/gluon states – 11.30.Rd Chiral symmetries

1 Introduction

The puzzle of the 0^{++} f_0-resonances has attracted a significant attention for the last decade (for recent reviews and references see [1–3]). The reason is that there have been discovered more f_0-mesons than can be accomodated by the quark model. This also fits with the expectations that there should be a glueball state with the same quantum numbers around the 1.5 GeV region. And indeed, an analysis of numerous available experimental data suggests that $f_0(1370)$ is mostly a $n\bar{n} = \frac{u\bar{u}+d\bar{d}}{\sqrt{2}}$ state, $f_0(1710)$ is mostly a $s\bar{s}$ state and $f_0(1500)$ is dominantly a glueball [4]. Other alternatives have also been discussed [5]. The decay modes of these mesons suggest that the physical mesons above are some mixtures of the pure quark-antiquark and glueball states.

Recently, there have appeared results of the partial-wave analysis of the mesonic resonances obtained in the $p\bar{p}$ annihilation at LEAR in the region 1.8–2.4 GeV [6–8]. In particular, four high-lying f_0 have been reported:

$$f_0(1770 \pm 12), \quad f_0(2040 \pm 38)\,,$$
$$f_0(2102 \pm 13), \quad f_0(2337 \pm 14)\,.$$

It has been suggested that $f_0(2102 \pm 13)$ should be a glueball, while $f_0(1770 \pm 12)$, $f_0(2040 \pm 38)$ and $f_0(2337 \pm 14)$ are $n\bar{n}$ states. The motivation for such an interpretation was the following: i) all these states are observed in $p\bar{p}$, hence according to OZI rule they cannot be $s\bar{s}$ states; ii) $f_0(2102 \pm 13)$ decays more strongly to $\eta\eta$ than to $\pi\pi$. It requires a large mixing angle, in contrast to the other three f_0 states. It can be interpreted naturally only if $f_0(2102 \pm 13)$ is not a $n\bar{n}$ state, but a glueball.

If $f_0(2102 \pm 13)$ is a glueball, a question then arises as why does it not mix with the neighbouring broad $f_0(2040{\pm}38)$, which is supposed to be a $n\bar{n}$ state? *"When a particle is observed experimentally it does not come labeled "quark-antiquark" or "glueball", nor is there any strict objective criterion to distinguish them. The whole distinction is tied up with the naive quark model, which has no firm basis in quantum field theory, and ignores the inevitability of mixing."* [9].

The aim of the present letter is to offer a natural explanation of the absence of a strong mixing. It has been argued recently that high in the hadron spectrum the spontaneously broken chiral symmetry of QCD is effectively restored [10–14][1]. A possible fundamental physical ground is that, in hadrons with large n (radial quantum number) or large L, the semiclassical approximation must be valid and hence the effects of quantum fluctuations of the quark and gluon fields must be suppressed [16]. If chiral symmetry is approximately restored, then the index of the representation of the chiral group, to which the given hadron belongs, becomes a good quantum number. Hence, a mixing of hadrons that belong to different representations is

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[1] This phenomenon has been referred to [15] as *chiral symmetry restoration of the second kind* in order to distinguish it from the familiar phenomenon of chiral symmetry restoration in the QCD vacuum at high temperature and/or density.

forbidden, *even though all other quantum numbers are the same and hadrons are close in energy.* This then explains the absence of a strong mixing between the $f_0(2040 \pm 38)$ state, which is a member of the (1/2, 1/2) representation of the chiral group [13], with a glueball, which belongs to a scalar representation (0,0).

2 The evidence and theoretical justification of chiral symmetry restoration in high-lying hadrons

We overview first a QCD-based justification of chiral symmetry restoration in high-lying spectra [11,12]. Consider two local currents (interpolating fields) $J_1(x)$ and $J_2(x)$ which are connected by chiral transformation, $J_1(x) = UJ_2(x)U^\dagger$, where $U \in SU(2)_{\rm L} \times SU(2)_{\rm R}$. These currents, when acting on the QCD vacuum $|0\rangle$, create hadron states with quantum numbers "1" and "2", respectively. All these hadrons are the intermediate states in the two-point correlators

$$\Pi_{J_\alpha}(q) = i \int {\rm d}^4x \; e^{-iqx} \langle 0|T\left\{J_\alpha(x)J_\alpha(0)\right\}|0\rangle, \quad (1)$$

where all possible Lorentz indices (which are specific for a given interpolating field) have been omitted, for simplicity. At large space-like momenta $Q^2 = -q^2 > 0$ the correlator can be adequately represented by the operator product expansion, where all nonperturbative effects reside in different condensates [17]. The only effect that spontaneous breaking of chiral symmetry can have on the correlator is via the quark condensate of the vacuum, $\langle q\bar{q}\rangle$, and higher dimensional condensates that are not invariant under chiral transformation U. However, the contributions of all these condensates are suppressed by inverse powers of momenta Q^2. This shows that at large space-like momenta the correlation function becomes chirally symmetric. In other words

$$\Pi_{J_1}(Q) \to \Pi_{J_2}(Q) \qquad \text{at} \qquad Q^2 \to \infty. \quad (2)$$

The dispersion relation provides a connection between the space-like and time-like domains for the Lorentz scalar (or pseudoscalar) parts of the correlator. In particular, the large-Q^2 correlator is completely dominated by the large-s spectral density $\rho(s)$, which is an observable. Hence the large-s spectral density should be insensitive to the chiral symmetry breaking in the vacuum and must satisfy

$$\rho_1(s) \to \rho_2(s) \qquad \text{at} \qquad s \to \infty. \quad (3)$$

This is in contrast to the low-s spectral densities $\rho_1(s)$ and $\rho_2(s)$, which are very different because of the chiral symmetry breaking in the vacuum. This manifests a smooth chiral symmetry restoration from the low-lying spectrum to the high-lying spectrum (chiral symmetry restoration of the second kind)[2].

Since the inclusive data indicate that the quark-hadron duality picture starts to work in the resonance region, we have to anticipate in this region a nontrivial implication of chiral symmetry. Indeed, if chiral symmetry restoration happens in the regime where the spectrum is still quasidiscrete (*i.e.* it is dominated by resonances and the successive resonances with the given spin are well separated), then these resonances must fill out representations of the parity-chiral group. There are evidences both in baryon [10–12] and in meson spectra [13,14] that light hadrons above $m \sim 1.7$ GeV fill out representations of the parity-chiral group, which are manifest as parity doublets or higher chiral representations. Nevertheless, a systematic experimental exploration of the high-lying hadrons is required in order to make definitive statements.

While the asymptotic prediction (3) is rather robust (it is based only on the asymptotic freedom of QCD in the deep space-like domain and on the analyticity of the two-point correlator; the earliest application of asymptotic freedom to $e^+e^- \to hadrons$ is ref. [22]), it does not tell us which physics could be associated with the chiral symmetry restoration in the isolated hadron. This question has been addressed in ref. [16]. One of the possibilities is that the chiral symmetry breaking (*i.e.* dynamical quark mass generation) is due to quantum fluctuations of the quark and gluon fields. This can be seen from the fact that the chiral symmetry breaking can be formulated via the Schwinger-Dyson equation. For the present context it is not important at all which specific gluonic interactions are the most important ones in the kernel of the Schwinger-Dyson equation, instantons, gluon exchanges or anything else. If the effects of quantum fluctuations of the quark and gluon fields are suppressed, then the dynamical mass of quarks must vanish. At large n (radial quantum number) or at large angular momentum L we know that in quantum systems the effects of quantum fluctuations become indeed suppressed and the semiclassical approximation (WKB) must work[3]. Physically, this approximation applies in these cases because the de Broglie wavelength of the valence quarks in the hadron is small in comparison with the size of the hadron. If so, the chiral symmetry

[2] A theoretical expectation that chiral symmetry must be restored high in the spectra is supported by the recent data on the difference between the vector and axial-vector spectral densities. This difference has been extracted from the weak decays of the τ-lepton by the ALEPH and OPAL Collaborations [18, 19]. The nonzero difference is entirely from the spontaneous breaking of chiral symmetry. It is well seen from the results that while the difference is large at the masses of $\rho(770)$ and $a_1(1260)$, it becomes very strongly reduced towards $m = \sqrt{s} \sim 1.7$ GeV. This is also seen from $e^+e^- \to hadrons$, where starting approximately from the same energy the spectral density oscillates around perturbative QCD prediction [20]. Similarly, recent data of JLab on inclusive electroproduction of baryonic resonances in the mass region $3.1 \leq M^2 \leq 3.9$ GeV^2 are perfectly dual to deep inelastic data [21].

[3] For example, the Lamb shift in the hydrogen atom is entirely due to the quantum fluctuations of the electromagnetic and electron fields, which are the vertex correction, electron self-energy and the vacuum polarization diagrams. The Lamb shift vanishes very fast with n, $\sim 1/n^3$, and also very fast with L.

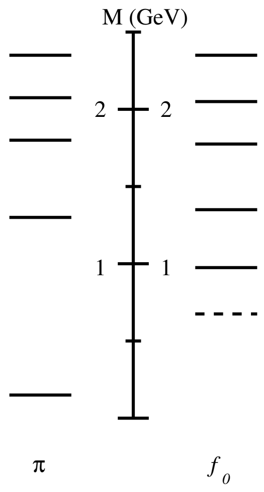

Fig. 1. Pion and $n\bar{n}$ f_0 spectra. The three highest states in both pion and f_0 spectra are taken from [6–8]. Since these f_0 states are obtained in $p\bar{p}$ and they decay predominantly into the $\pi\pi$ channel, they are considered in [6–8] as $n\bar{n}$ states.

must be effectively restored in high-lying hadrons. So a very natural picture for highly excited hadrons is a string with bare quarks of definite chirality at the end-points [14].

3 Chiral quantum number as a good quantum number

In the low-lying hadrons, where the effects of chiral symmetry breaking in the vacuum $SU(2)_\mathrm{L} \times SU(2)_\mathrm{R} \to SU(2)_I$ are strong, only isospin is a good quantum number among those that are potentially supplied by the chiral group. In the regime where chiral symmetry is (approximately) restored, a new good quantum number appears, that characterizes a parity-chiral multiplet. This quantum number is an index of the irreducible representation of the parity-chiral group[4]. For example, the vector (ρ) and axial-vector (a_1) mesons, in the chirally restored regime become chiral partners and fill out in pairs $(0,1)\oplus(1,0)$ irreducible representations of the parity-chiral group. Their valence content is given as [14]

$$\rho: \quad \frac{1}{\sqrt{2}}\left(\bar{R}\vec{\tau}\gamma^\mu R + \bar{L}\vec{\tau}\gamma^\mu L\right), \tag{4}$$

$$a_1: \quad \frac{1}{\sqrt{2}}\left(\bar{R}\vec{\tau}\gamma^\mu R - \bar{L}\vec{\tau}\gamma^\mu L\right). \tag{5}$$

The valence composition of $n\bar{n}$ f_0- and π-mesons in the chirally restored regime is

$$f_0: \quad \frac{1}{\sqrt{2}}\left(\bar{R}L + \bar{L}R\right), \tag{6}$$

$$\pi: \quad \frac{1}{\sqrt{2}}\left(\bar{R}\vec{\tau}L - \bar{L}\vec{\tau}R\right). \tag{7}$$

Both these mesons fill out in pairs the $(1/2, 1/2)$ representations of the parity-chiral group [13]. The experimental data in the range 1.8–2.4 GeV are summarized in the following table:

Meson	I	J^P	Mass (MeV)	Width (MeV)	Reference
f_0	0	0^+	1770 ± 12	220 ± 40	[6]
f_0	0	0^+	2040 ± 38	405 ± 40	[7]
f_0	0	0^+	2102 ± 13	211 ± 29	[7]
f_0	0	0^+	2337 ± 14	217 ± 33	[7]
π	1	0^-	1801 ± 13	210 ± 15	[1]
π	1	0^-	2070 ± 35	310^{+100}_{-50}	[8]
π	1	0^-	2360 ± 25	300^{+100}_{-50}	[8]

It is well seen that while the chiral symmetry is strongly broken low in the spectrum, the high-lying $n\bar{n}$ f_0- and π-mesons indeed form chiral pairs (see also fig. 1):

$$\pi(1300 \pm 100)\text{-}f_0(1370^{+130}_{-170}), \tag{8}$$

$$\pi(1801 \pm 13)\text{-}f_0(1770 \pm 12), \tag{9}$$

$$\pi(2070 \pm 35)\text{-}f_0(2040 \pm 38), \tag{10}$$

$$\pi(2360 \pm 25)\text{-}f_0(2337 \pm 14). \tag{11}$$

A true glueball (G) has no valence quark content and hence must belong to the scalar representation of the parity-chiral group, $G \sim (0,0)$.

If chiral symmetry is a good symmetry then the index of the corresponding irreducible representation of the parity-chiral group becomes a good quantum number. Hence, quantum mechanics forbids a mixing of the states that belong to different representations, even though all other quantum numbers coincide and states are close in energy. Since the true glueball by definition belongs to the $(0,0)$ representation and the approximately degenerate f_0- and π-mesons form the $(1/2,1/2)$ representation, they cannot be mixed. This explains why $f_0(2102)$, which is presumably a glueball and has no partner in the pion spectrum, and which is very close to $f_0(2040)$, that is a $n\bar{n}$ state, does not mix with the latter. If confirmed by a detailed study of decay modes, it would be better to rename $f_0(2102)$ as $G(2102)$.

Clearly, the chiral restoration is not exact at masses of 2 GeV, so a small amount of mixing is still possible. Hence, it is a very important experimental task to study in detail

[4] The irreducible representations of $SU(2)_\mathrm{L} \times SU(2)_\mathrm{R}$ can be labeled as $(I_\mathrm{L}, I_\mathrm{R})$ with I_L and I_R being the isospins of the left and right subgroups. However, generally the states that belong to the given irreducible representation of the chiral group cannot be ascribed a definite parity because under parity transformation the left-handed quarks transform into the right-handed ones (and *vice versa*). Therefore under a parity operation the irreducible representation $(I_\mathrm{L}, I_\mathrm{R})$ transforms into $(I_\mathrm{R}, I_\mathrm{L})$. Hence, in general, the state (or current) of definite parity can be constructed as a direct sum of two irreducible representations $(I_\mathrm{L}, I_\mathrm{R}) \oplus (I_\mathrm{R}, I_\mathrm{L})$, which is an irreducible representation of the parity-chiral group [12].

decay modes and mixings of these high-lying states. The prize for these efforts would be that we identify and study a "true" glueball.

We can also look at the same problem from the other perspective. If a detailed study of decay modes confirmes that $f_0(2102)$ is a glueball and that $f_0(2040)$ is a $n\bar{n}$ state, then it would be an independent confirmation of chiral symmetry restoration.

As a conclusion, chiral symmetry restoration high in the hadron spectra provides a natural basis for the absence of a strong mixing between those scalar mesons that are chiral partners of pions, and those scalar mesons, that represent a glueball. Then we have an opportunity to study a "true" glueball of QCD.

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