Light meson spectrum and classical symmetries of QCD

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Abstract. Modern spectroscopic data on the light non-strange meson spectrum is analyzed. It is argued that the observed regularities of the experimental spectrum for highly excited states favour a partial restoration of all approximate classical symmetries of a QCD Lagrangian (conformal, chiral and axial) broken by quantum corrections. The rate of restoration of classical symmetries is estimated. The dependence of the resonance widths on the corresponding masses is systematically checked. On average, it turns out to be universal for high excitations as predicted by the effective string description.

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1 Introduction

The study of hadron resonances is of great importance for a deeper understanding of strong interactions. As the stable hadronic matter consists of *up*- and *down*-quarks, the resonances built up of these quarks are of a special interest. It is well known that masses of *up*- and *down*-quarks are very light (of the order of 5 MeV) in comparison with typical hadron masses (of the order 1000 MeV). Thus, with a good accuracy one can neglect them. In this massless limit strong interactions are chirally invariant in the two-flavor sector. The chiral $SU(2)_L \times SU(2)_R$ invariance is not a symmetry of the physical vacuum. This results in the spontaneous Chiral Symmetry Breaking (CSB) to the vector isospin subgroup $SU(2)_V$ and in the appearance of massless Goldstone bosons, the π-mesons. For this reason, the chiral symmetry is not seen in the hadronic spectrum. The vector $\rho(770)$- and axial $a_1(1230)$-mesons represent a typical textbook example. Another example is the axial $U(1)_A$ symmetry broken by chiral anomaly. This phenomenon is known to enhance significantly the mass of the η'-meson. However, all such examples refer to ground states only, whereas the higher radial and orbital excitations are usually avoided in QCD textbooks.

The classical QCD action in the chiral limit has also a symmetry with respect to the scale transformations stemming from the absence of dimensionful constants. The scale invariance is part of a larger conformal group. At the quantum level this symmetry is broken by scale anomaly. At high energies the conformal symmetry of QCD is, how-

ever, restored due to the asymptotic freedom. In particular, the scaling laws of the parton model can be derived directly from the conformal symmetry of the classical QCD Lagrangian. In recent years, it was understood that this fact provides one with powerful tools in practical calculations: The structure of perturbative predictions for light-cone–dominated processes reveals the underlying conformal symmetry of the QCD Lagrangian, see [1] for a review. This property turned out to be crucial in connection with the conjecture about AdS/CFT correspondence [2] in application to QCD, which attracted a lot of interest recently.

The conformal symmetry is incompatible with the existence of resonances at certain energies because the spectrum has to be scale invariant in this case. However, between the low-energy region, where the scale symmetry is badly violated, and the scale-invariant high-energy continuum there is an intermediate-energy region, where the resonances still exist but the conformal invariance should be partially restored. This raises an interesting problem: How does the partial restoration of conformal symmetry influence the meson spectrum?

It has been suspected for long ago that QCD is dual to some string theory. The term "duality" is commonly used when some phenomena can be described by two theories and the strong-coupling regime of one of them corresponds to the weak-coupling regime of the other one. There are examples of such duality in two-dimensional field theories. For four dimensions duality is usually only a hypothesis, which gives, however, a powerful tool for deriving various predictions. The large-N_c limit of QCD [3] provides, in a sense, a particular realization of certain duality: The the-

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ory of strongly interacting quarks and gluons is rewritten as a theory of weakly interacting mesons and glueballs, with baryons being the solitons in this dual theory. The two-point correlators can be then rewritten as a sum over meson contributions,

$$\langle J(p)J(-p)\rangle = \sum_n \frac{f_n^2}{p^2 - m_n^2} \sim N_c \log p^2, \qquad (1)$$

where $f_n = \langle 0|J|n\rangle \sim \mathcal{O}(\sqrt{N_c})$ are meson couplings. The logarithm in the r.h.s. of eq. (1) comes from the perturbation theory leading order (the so-called parton model logarithm) and it is related with the conformal invariance of classical QCD. Obviously, to reproduce this logarithm one needs an infinite number of states, provided the existence of confinement in the large-N_c limit. Hence, an infinite number of narrow ($\Gamma = \mathcal{O}(1/N_c)$) meson states is dual, at least, to the leading order of the perturbation theory, which is governed by the underlying conformal symmetry of QCD. On the other hand, there are many arguments that QCD in the large-N_c limit is dual to some (still unknown) string theory. A reason for such a belief is, for instance, the fact that the planar expansion in powers of $1/N_c$ is much reminiscent of perturbative expansion in the string theory; both expansions have a topological nature.

Thus, QCD is believed to have some string dual and the large-N_c limit strongly supports this belief. Any self-consistent string theory possesses conformal invariance. The approximate conformal symmetry of QCD at the tree level gives a hope to find this string dual, at least in some kinematic regime. But what is then a signature of approaching to this regime for the light meson spectrum? Probably, the reply is that the spectrum should get reminiscent of that given by the string dual. The conformal symmetry is crucial in making this correspondence. This gives an idea for a mechanism of how the partial restoration of scale invariance influences the spectrum. The quasi-classical string approaches typically give the following law for the light meson spectrum: $m^2(n, J) \sim n + J$, where n is the principal (radial) quantum number and J is the spin. This spectrum is provided by the Veneziano type of strings. Hence, the QCD string dual, probably, represents a modification of *à la* Veneziano string with a similar spectrum. A special feature of this spectrum is that it predicts the clustering of states with different n and J near certain equidistant values of masses squared defined by the sum $n + J$. Thus, if one experimentally observes a tendency to such a clustering, this phenomenon could be interpreted as a manifestation of the partial restoration of conformal invariance of the underlying fundamental theory.

In principle, the corresponding physics can be figured out without exploiting the string ideas. If some "rest" of conformal symmetry is indeed realized in the meson spectrum, then the physical states must fill out the corresponding group representations with degenerate masses inside one multiplet. Experimentally, these multiplets should be then observed as clusters of states near some values of energy. Unfortunately, it seems that nothing is known on this subject.

Motivated by these discussions, in the present paper we address to the problem of relations between the experimental spectrum of light non-strange mesons and approximate classical symmetries of QCD broken at the quantum level.

The paper is organized as follows. The details of phenomenological analysis are given in sect. 2. Sections 3 and 4 are devoted to the interpretation of observed regularities for the masses and decay widths correspondingly. We conclude in sect. 5.

2 Experimental spectrum

The radial and orbital excitations were only poorly known in the time of fast development of QCD in the 70s. Since that time the experimental data has been accumulating and now the Particle Data Group (PDG) [4] lists a certain number of well-established higher excitations in the light non-strange meson sector (we will denote these states $\bar{n}n$) up to the energy of 1.9 GeV. At higher energies PDG enumerates only a few confirmed mesons and many unconfirmed states. At present it is difficult to draw any direct conclusions about general properties of meson excitations based on the well-confirmed states of PDG only. To reveal these properties we propose an indirect way: together with well-confirmed states one can analyse many unconfirmed (more precisely, not well-confirmed) states. As usually happens in a large statistical ensemble, one can hope that possible errors in different channels smooth each other providing finally a stable general picture, which can be described by some mean characteristics.

Since PDG cites so many unconfirmed states it is easy to go astray in searching for regularities. To avoid this, one inevitably should stick to some reasonable principles. Let us explain how we choose the unconfirmed resonances for the analysis. First, for reliability we will take only those states which were observed at least in two different reactions. Thus, we will deal with the "not well-confirmed" mesons rather than with the "unconfirmed" ones. Second, at energy above 1.9 GeV we will use the data of the Crystal Barrel Collaboration on proton-antiproton ($\bar{p}p$) annihilation in flight. The latest review of this data is contained in ref. [5]. The reasons for this choice are as follows:

1. It is the only experiment which performed a systematic study of the mass range 1.9–2.4 GeV. The coverage of this mass range from other experiments is very limited.
2. As a rule the states were independently observed in different channels, *i.e.* they are quite reliable. The reason why most of them are listed by PDG in a section "Other Light Unflavoured Mesons" is that PDG requires confirmation from a separate experiment. The appearance of other states in this section usually has a rather sporadic character.
3. As was realized long ago [6], meson resonances are strongly coupled to the $\bar{N}N$ reactions because mesons have the quantum numbers of the $\bar{N}N$ system. The dominant role of this system in the dynamics of meson states makes the data extracted from the $\bar{N}N$ reactions quite reliable.

4. A possible admixture of strange quarks is a serious problem for any classification of light states. A feature of $\bar{p}p$ annihilation is that the production of the $\bar{s}s$ component is strongly suppressed. Consequently, it is quite reliable that the discovered states are, except some rare cases, genuine $\bar{n}n$-mesons.

5. The obtained spectrum (first systematized in ref. [7]) turned out to be in full agreement with the old theoretical expectations from the hadron string models and the low-energy amplitudes [8]. Namely,

   a) linearity of Regge trajectories;

   b) equidistance of daughter Regge trajectories (linearity of radial Regge trajectories), as a consequence, approximate universality of slopes of trajectories;

   c) the intercept of the pion Regge trajectory is approximately equal to 0 and the intercept of the ρ-meson one is 0.5;

   d) the slope of the radial Regge trajectories is about $2m_\rho^2$, which is also consistent with the lattice calculations (see, e.g., [9]).

6. To reveal the general properties of the spectrum it is preferable to use the data of an individual systematic experiment. Only after that the overall picture should be compared with the one given by another systematic experiment. This type of comparison can lead to some global shifts but does not spoil the picture qualitatively. If one first performs the data averaging (as it is done by PDG), the errors accumulate rapidly and the final picture can be completely obscured. The case of light non-strange baryons is a good example: If one separately uses the data of individual systematic experiments (say, going under the names "Cutkosky" or "Hoehler" in PDG) the multispin-parity clusters of states are unambiguously seen, but if one takes the averaged data of PDG the clustering gets rather controversial.

After these general arguments let us pass to the analysis. As said above, we do not consider the states with a large admixture of strange quarks (usually, it is clear from the analysis of corresponding decay channels) and we omit all states which were observed in one channel only (although many of them fill well the missing states on meson trajectories). In review [5] the latter states are: $\omega(2205)$, $a_1(1930)$, $a_1(2270)$, $a_2(1950)$, $a_2(2175)$, $\omega_4(2250)$, $b_5(2500)$ and $f_6(2485)$. For the same reasons, we omit $h_1(1595)$ and $b_1(1620)$ (see [5] for references). We do not use $\pi_2(1880)$ and $\eta_2(1870)$ which were cited in [5] and were shown to be inconsistent with the $\bar{n}n$ state. Similarly, we omit $f_0(2100)$ from [5], which is either a glueball or a $\bar{s}s$ state strongly mixed with $\bar{n}n$. The very narrow $\rho(1900)$ cited by PDG [4] (in the list of unconfirmed states) is also not considered since there are many doubts that it is a real resonance. The only well-confirmed states of PDG which are exotic for the quark model are π_1-mesons, namely $\pi_1(1400)$ and $\pi_1(1600)$. The state $\pi_1(2015)$ was seen in two reactions. We decided to include them into analysis because at least the first two of them are generally recognized observable resonances. We also include $f_0(980)$ and $a_0(980)$ although the nature of these states is still controversial, presumably they have a large admixture of strange quark (see, e.g., the note on the scalar mesons in ref. [4]). The reason will be explained below. The state $\eta(547)$ has a large admixture of strange component. Nevertheless, this admixture does not seem to be dominant in the corresponding radial excitations. For this reason, the η-meson is also considered.

Let us explain how we display the data. First, in the relativistic theories one deals with (masses)^2 which appear in the multiplets, Regge and string theory etc., and only these quantities are of theoretical interest. Second, it is better to normalize all masses to some typical hadron mass. In our opinion, the best candidate for the normalization is the mass of the ρ-meson.

The final picture of the meson spectrum resulting from our analysis is displayed in fig. 1. The corresponding experimental data is given in table 1.

A well-pronounced feature of the spectrum is that the observed states cluster at about some values of energy [5,10]. A similar phenomenon exists in the light non-strange baryons [11]. The clustering occurs at approximately 1.33, 1.70, 2.00 and 2.27 GeV. Some channels have additional states denoted by open circles or strips in fig. 1. They appear because the states in these channels can be created by different orbital momenta which result in doubling of the corresponding radial Regge trajectories [7]. For ρ and f_2 mesons there is polarization data which separates the S-wave from the D-wave and the P-wave from the F-wave states correspondingly [5]. For other channels such separation is tentative and new experiments are called for.

For the clusters we display in table 1 the mean mass $\bar{M}$ and the mean full decay width $\bar{\Gamma}$, which are defined as follows,

$$\bar{M} \equiv \sqrt{\frac{1}{k}\sum_k m_k^2}\,, \qquad \bar{\Gamma} \equiv \frac{1}{k}\sum_k \Gamma_k\,, \qquad (2)$$

where the index k enumerates the states in a cluster. The rules for the averaging are natural: The observable quantities are m_k^2 (as discussed above) and Γ_k. The data for $\bar{M}$ and $\bar{\Gamma}$ is presented in the form ("m." denotes "mean")

$$\bar{M}, \bar{\Gamma} = \text{m. value} \pm \text{m. square deviation} \pm \text{m. exper. error}.$$

It must be emphasized that the positions of clusters are very stable due to many states involved. For instance, above 1.9 GeV one can consider only those states from [5] which have the maximal star rating (rating 4* according to the classification in [5]). These resonances require observation of 3 or more strong, unmistakable peaks and a very good mass determination. Their reliability is practically equivalent to that of states in PDG. There exist 6 such states in the third cluster and 8 in the fourth one. One can check that if we consider only these states in the clusters, the positions of the clusters will not change (i.e. the change will be less than 0.01 GeV within our accuracy).

The clusters describe the behavior of the spectrum as a whole (the relevant discussions for the light non-strange baryons can be found in ref. [12]). With a good accuracy they are equidistant, hence, one can parametrize them by
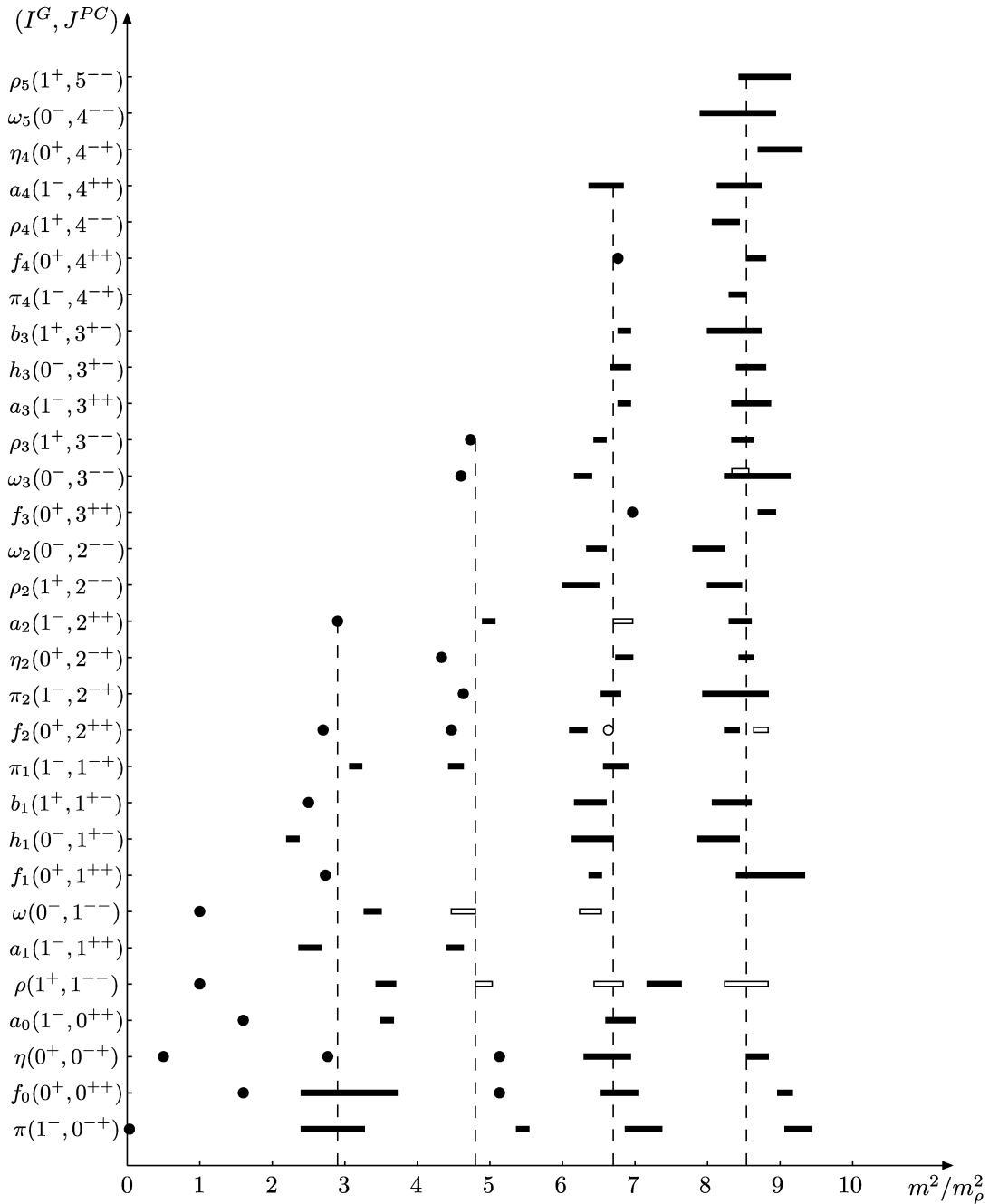
Fig. 1. The spectrum of light non-strange mesons from refs. [4] and [5] (for the last two clusters) in units of $m^2_{\rho(770)}$. Experimental errors are indicated. Circles stay when errors are negligible. Open circles and strips denote the additional states (see text). The dashed lines mark the mean (mass)^2 in each cluster. The absolute values of the masses are given in table 1. The masses of the lightest states not displayed in table 1 are (in MeV): π: 140; f_0: 980 ± 10; η: 547.75 ± 0.12; a_0: 984.7 ± 1.2; ρ: 775.8 ± 0.5; ω: 782.59 ± 0.11.

a linear function. For the data in fig. 1 the result of the fit is

$$M^2(n) = an + b, \quad n = 1, 2, 3, 4; \ a \approx 1.13, \ b \approx 0.63, \quad (3)$$

where $M^2(n)$ is the position of the n-th cluster in GeV^2. The slope a in the cluster spectrum (3) is nothing but the mean slope of radial Regge trajectories. Its numerical value is within the interval found in [5]: $a = 1.14 \pm 0.013$.

The parameter b is the mean intercept of the radial Regge trajectories. In [5,7] this quantity was not estimated, but for us it is of importance as will be seen below.

Finally, we would like to estimate to what extent the cluster spectrum in eq. (3) is vitiated if one excludes the Crystal Barrel data (the last two clusters). Then we have only two clear-cut clusters. The parametrization of two points by the linear function can look doubtful, so we will

Table 1. The masses and widths (in MeV) of the states in fig. 1. Experimental errors are indicated.

	$m(1)$	$\Gamma(1)$	$m(2)$	$\Gamma(2)$	$m(3)$	$\Gamma(3)$	$m(4)$	$\Gamma(4)$
π	1300 ± 100	200–600	1812 ± 14	207 ± 13	2070 ± 35	310^{+100}_{-50}	2360 ± 25	310^{+100}_{-50}
f_0	1200–1500	200–500	1770 ± 12	220 ± 40	2020 ± 38	405 ± 40	2337 ± 14	217 ± 33
η	1294 ± 4	55 ± 5	1760 ± 11	60 ± 16	2010^{+35}_{-60}	270 ± 60	2285 ± 20	325 ± 30
a_0	1474 ± 19	265 ± 13			2025 ± 30	300 ± 25		
ρ	1465 ± 25	400 ± 60	1720 ± 20	250 ± 100	2000 ± 30 2110 ± 35	260 ± 45 230 ± 50	2265 ± 40	325 ± 80
a_1	1230 ± 40	250–600	1647 ± 22	254 ± 27				
ω	1400–1450	180–250	1670 ± 30	315 ± 35	1960 ± 25	195 ± 60		
f_1	1281.8 ± 0.6	24.1 ± 1.1			1971 ± 15	240 ± 25	2310 ± 60	255 ± 70
h_1	1170 ± 20	360 ± 40			1965 ± 45	345 ± 75	2215 ± 40	325 ± 55
b_1	1229.5 ± 3.2	142 ± 9			1960 ± 35	230 ± 50	2240 ± 35	320 ± 85
π_1	1376 ± 17	300 ± 40	1653^{+18}_{-15}	225^{+45}_{-28}	2013 ± 25	230 ± 105		
f_2	1275 ± 1	$185.1^{+3.5}_{-2.6}$	1638 ± 6	99^{+28}_{-24}	1934 ± 20 2001 ± 10	271 ± 25 312 ± 32	2240 ± 15 2293 ± 13	241 ± 30 216 ± 37
π_2			1672 ± 3	259 ± 9	2005 ± 15	200 ± 40	2245 ± 60	320^{+100}_{-40}
η_2			1617 ± 5	181 ± 11	2030 ± 16	205 ± 18	2267 ± 14	290 ± 50
a_2	1318.3 ± 0.6	107 ± 5	1732 ± 16	194 ± 40	2030 ± 20	205 ± 30	2255 ± 20	230 ± 15
ρ_2					1940 ± 40	155 ± 40	2225 ± 35	335^{+100}_{-50}
ω_2					1975 ± 20	175 ± 25	2195 ± 30	225 ± 40
f_3					2048 ± 8	213 ± 34	2303 ± 15	214 ± 29
ω_3			1667 ± 4	168 ± 10	1945 ± 20	115 ± 22	2255 ± 15 2285 ± 60	175 ± 30 230 ± 40
ρ_3			1688 ± 2.1	161 ± 10	1982 ± 14	188 ± 24	2260 ± 20	160 ± 25
a_3					2031 ± 12	150 ± 18	2275 ± 35	350^{+100}_{-50}
h_3					2025 ± 20	145 ± 30	2275 ± 25	190 ± 45
b_3					2032 ± 12	117 ± 11	2245 ± 50	320 ± 70
π_4							2250 ± 15	215 ± 25
f_4					2018 ± 6	182 ± 7	2283 ± 17	310 ± 25
ρ_4							2230 ± 25	210 ± 30
a_4					2005^{+25}_{-45}	180 ± 30	2255 ± 40	330^{+110}_{-50}
η_4							2328 ± 38	240 ± 90
ω_5							2250 ± 70	320 ± 95
ρ_5							2300 ± 45	260 ± 75
$\bar{M}$	$1325 \pm 89 \pm 31$		$1697 \pm 56 \pm 12$		$2004 \pm 40 \pm 24$		$2269 \pm 37 \pm 32$	
$\bar{\Gamma}$		$248 \pm 132 \pm 57$		$199 \pm 66 \pm 29$		$224 \pm 69 \pm 38$		$266 \pm 56 \pm 53$
$\bar{\Gamma}/\bar{M}$		0.187		0.117		0.112		0.117

consider the ground ρ- and ω-mesons as two non-strange constituents of the lowest cluster near $0.78\,\mathrm{GeV}$. We note in passing that fit (3) predicts the lowest cluster for $n = 0$ near $0.79\,\mathrm{GeV}$, so our assumption is well justified. We have then

$$M_{\mathrm{PDG}}^2(n) = an + b, \quad n = 0, 1, 2; \quad a \approx 1.14, \, b \approx 0.61. \quad (4)$$

Both cluster spectra (3) and (4) turn out to be very close. Thus, PDG contains enough data to arrive at our conclusions. The data of Crystal Barrel provides a dramatic confirmation for the observed regularities.

Concluding this section we would like to make the following remark. The hypothesis that mesons should appear as towers of states (which we call clusters, fig. 1 is self-explanatory in the analogy with towers) was proposed before QCD [6] for explaining the absence of backward peaks

in $\pi^+\pi^-$, π^+K^-, K^+K^-, and $\bar{N}N$ elastic scattering in the framework of Regge theory. There was a hope that further $\bar{N}N$ studies (see reason 3 above) would provide crucial tests for the existence of these towers. The Crystal Barrel experiment on $\bar{p}p$ annihilation can be considered as such a test.

3 Interpretation of data

It is well known that properties of any quantum system approach to its classical ones while the quantum numbers defining the stationary states of this system are large enough (see, *e.g.*, [13]). In our case these quantum numbers are the spin J and the radial excitation number n. The valent quarks in such hadrons on average have high

energies and, hence, practically do not "feel" the non-perturbative structure of QCD vacuum which is the underlying reason of CSB (see the relevant discussions in [14] and references therein). The quasiclassical description implies, in particular, the universal linear string-like behavior of the meson mass spectrum which can be compactly written as

$$m^2(I, G, P, C, L, J, n) \simeq a(n + J) + b. \qquad (5)$$

The fact that relation (3) is an experimental result and that the number of states is indeed growing in clusters seems to confirm the validity of the quasiclassical treatment. Needless to say that the manifest cluster structure of high meson excitations includes the full linearly realized approximate $U(2) \times U(2)$ chiral flavor symmetry of QCD as a particular case. The restoration of this symmetry at high energies leads to degeneracy inside the chiral multiplets. Different aspects of the relation between the chiral symmetry restoration in highly excited hadrons and the parity doubling were widely discussed in the literature [10,14–19]. It happens, however, that with the same accuracy the observed mass degeneracy is much higher than predicted by the restoration of chiral and axial symmetries of QCD Lagrangian. Even models of generalized chiral symmetry like in ref. [16] cannot explain such a high degeneracy because in the chiral multiplets one has states with equal spin only. This phenomenon should be a manifestation of some additional symmetry. If we believe that all regularities in the spectrum must be related to the symmetries of QCD, then we have only one possible candidate: the conformal symmetry. According to the general principle in the quantum theories discussed at the beginning of this section, one expects the restoration of *all* broken classical symmetries of the QCD Lagrangian in highly excited states. The conformal invariance is among the classical symmetries. Consequently, it should be effectively restored at high energies. Indeed, we know that QCD is nearly conformal in the ultraviolet region. As discussed briefly in the introduction, this could be intimately related with the existence of string dual for QCD. However, following the arguments given in the introduction, even without this duality, it is natural to suggest that *the observed degeneracy of the light non-strange mesons is a combined effect of a partial restoration of chiral and conformal symmetries at high energies.*

Unlike the case of chiral invariance, where the complete restoration is possible in the spectrum (*i.e.* the complete parity doubling), we cannot observe in the spectrum the complete restoration of conformal invariance, which could mean the ideal degeneracy of states with different spin inside a cluster, like in the string theories of Veneziano type. In QCD the existence of hadrons at discrete energies is incompatible with the absence of scale in the problem. What we can observe is only approaching to that regime. Finally, the resonances disappear and the scale-invariant continuum sets in. Thus, the approaching to the perturbative continuum and the grouping of resonances into clusters seem to be tightly related.

As seen qualitatively from fig. 1 and numerically from table 1 the higher are the resonances one considers the

more clear-cut clusters they form. Let us estimate the rate of clustering. In doing this certain care should be exercised. This procedure makes sense only if the deviations from the averaged values are substantially larger than the corresponding experimental errors. As seen from table 1 this is indeed the case for the first and second clusters, where the deviations are by a factor of 3–4 larger than the averaged experimental errors. For the third cluster the difference is by a factor of 2 only, while for the fourth one there is practically no difference at all. Thus, only the first two clusters (the PDG data) can serve for our purpose more or less reliably. The Crystal Barrel data will be used for a qualitative check.

The equidistant cluster spectrum with deviations can be written in the form

$$M(n) = \sqrt{an + b} \pm \delta(n). \qquad (6)$$

Now we should interpolate the deviation $\delta(n)$ by some smooth function. *A priori* we have no theoretical idea how this function should look like. In the literature there exist some arguments for the rate of chiral symmetry restoration only. Namely, in [17,19] the deviations were argued to be exponential, while in [14] a polynomial minimal rate was derived. We will consider both possibilities for $\delta(n)$,

$$\delta_{\mathrm{e}}(n) \sim \frac{e^{-\beta_{\mathrm{e}} n}}{\sqrt{n+1}}, \qquad \delta_{\mathrm{p}}(n) \sim (n+1)^{-\beta_{\mathrm{p}} n}. \qquad (7)$$

Here in the first ansatz we introduced the square root in order to have a purely exponential correction for the mass squared. Taking the corresponding values for the first two clusters in table 1, $\delta(1) \approx 89\,\mathrm{MeV}$, $\delta(2) \approx 56\,\mathrm{MeV}$, one arrives at the following estimates:

$$\beta_{\mathrm{e}} \approx 0.26, \qquad \beta_{\mathrm{p}} \approx 1.14. \qquad (8)$$

The resulting predictions for the exponential and polynomial deviations are (in MeV): $\delta_{\mathrm{e}}(3) \approx 37$, $\delta_{\mathrm{e}}(4) \approx 26$, $\delta_{\mathrm{p}}(3) \approx 40$, $\delta_{\mathrm{p}}(4) \approx 31$, while experimentally $\delta(3) \approx 40$, $\delta(4) \approx 37$. It is seen that the polynomial ansatz works slightly better. The modern level of experimental accuracy, however, does not allow to indicate convincingly which ansatz is really preferable.

Let us speculate about the physical sense of these estimates. If the partial restoration of conformal invariance indeed takes place, then the obtained results can be considered as a rough estimate for the rate of this restoration. On the other hand, they may be regarded as an estimate for the minimal rate of the chiral symmetry restoration in excited hadrons. Obviously, in particular channels this effect can occur faster. For instance, the fits in [17] yield $\beta_{\mathrm{e}} \approx 1$, while for the polynomial ansatz there exist an estimate for the scalar channels [14], $\beta_{\mathrm{p}} \gtrsim 1.5$. In any case it should be noted that although the estimation of this rate is still a rather controversial problem, the effect of chiral symmetry restoration in highly excited states *per se* seems to be now well settled both experimentally and theoretically.

In principle, one can try to give some alternative explanations of experimental data, say, for parity doubling.

For instance, it may be that at high energies the influence of CSB ceases to depend on the quantum numbers of concrete channel, but the strong CSB persists at all energies, just higher excitations "feel" its presence equally. Or it may be that the mass degeneracy is just an effect of vanishing spin-orbit forces in quark interactions as it was proposed for light baryons in ref. [20]. Some independent tests are needed.

Actually, such a test can be provided by the results of some recent papers. Due to the observed universality of spectra expressed by the existence of clusters (3) it seems to be sufficient to check the situation for some channels only. The channels with $J = 1$ are good candidates for our purpose because the problem can be directly addressed in these cases within the QCD sum rules in the planar limit. Since the spectrum is linear with a good accuracy, one can saturate the sum rules by the linear ansatz for the masses. As it was noted in [17,18] and developed in [19], if the chiral symmetry is not broken (which is equivalent to the absence of the weak pion decay constant and the quark condensate) then the linear spectrum turns out to be: $M^2(n) = a(n + 1/2)$. For $J = 1$ mesons this relation is a consequence of the absence of the gauge-invariant local dimension-two gluon condensate. If the CSB is present at any energy then the intercept should be substantially larger, e.g., in ref. [18] it was obtained for this case $M^2(n) = a(n + 1)$. Experimentally, one has for $J = 1$ clusters: $m^2(n) \approx a(n + 0.51)$ with $a \approx 1.13 \, \text{GeV}^2$. Consequently, the experimental spectrum on average favors the chirally symmetric pattern, i.e. it reveals a strong supression of CSB effects for high excitations where the linear behavior sets in. For the experimental spectrum (3) one obtains $m^2(n) \approx a(n + 0.55)$ with the same slope. Thus, the universality works remarkably well providing a solid ground for the extension of the conclusion to the whole spectrum.

Finally, let us consider the states $f_0(980)$ and $a_0(980)$ the nature of which is a subject of many discussions in the literature. A large amount of phenomenological arguments (see ref. [21] and references therein) indicates that these mesons are genuine $\bar{n}n$ states with a large admixture of $\bar{s}s$ component which shifts their masses almost to the $\bar{K}K$ threshold. The analysis of different reactions (see [21] for references) yields the estimation of the strange component in $f_0(980)$ to be about 60–70%. In fact, this estimate can be easily obtained theoretically. After the CSB the ground vector-isovector meson and the ground scalar-isoscalar meson practically do not mix [22]. If the latter state is $f_0(980)$ then we should have $m_\rho^2 = m_{f_0}^2$. Since the ρ-meson is a pure $\bar{n}n$ state, the estimation of $\bar{s}s$ admixture in $f_0(980)$ follows immediately: $(m_{f_0}^2 - m_\rho^2)/m_\rho^2 \approx 0.6$ (we remind that in all relevant formulae one deals with (masses)^2). The situation with $a_0(980)$ happens to be similar. Such an estimation does not give any insight into the mechanism of this admixture and it may be that for $f_0(980)$ and $a_0(980)$ this mechanism is different. It only shows that numerically it is consistent with the hypothesis that these mesons are genuine quark-antiquark ground scalar states. Thus, if the strange quarks were "switched

off" the lowest cluster at about 0.78 GeV would consist of four mesons: $\rho(770)$, $\omega(782)$, $f_0(980)$ and $a_0(980)$ (that is why the last two particles have been included into our analysis). As noted above, this cluster is in good agreement with the spectrum (3) for $n = 0$.

4 Analysis of decay widths

The clusters of meson states have not only stable positions at some equidistant values of energy, but also stable mean decay widths. They are shown in table 1. What do they tell us about? In the given section we address ourselves to this question.

It is widely believed that the light mesons can be considered as an effective hadron string with relativistic quarks at the ends. The conjecture is that a flux tube of the chromoelectric field between a quark and an antiquark can be effectively described as a string. On the basis of such a simple qualitative picture the following behaviour for the full decay width was predicted [23]: $\Gamma(n) = Bm(n)$, with $B = \mathcal{O}(1/N_c)$ being a universal constant. Originally, this relation was derived for highly excited states, where a quasiclassical treatment can be applied. With this result at hand, let us consider the behaviour of the mean width in the clusters, namely introduce the number $B(N)$ defined as,

$$B(N) \equiv \frac{\bar{\Gamma}(N)}{\bar{M}(N)} \, . \tag{9}$$

For $N = 1, 2, 3, 4$ the corresponding values are given in table 1.

It is desirable to have an estimate for $B(0)$ as well. Here one must exercise certain care because the averaging of widths for the ground states should not be the same as for the excited ones. First of all, we do not consider the mesons $f_0(980)$ and $a_0(980)$ because a large admixture of the strange quark in these states is expected to change dramatically their widths. Although above we have presented an argument why these states can be considered as the members of the lowest cluster in the analysis of the mass spectrum, this hardly can be done for the widths. Second, the full width of the $\omega(782)$-meson, $\Gamma = 8.49 \pm 0.08 \, \text{MeV}$, is almost by 18 times less than that of the $\rho(770)$-meson, $\Gamma = 150.3 \pm 1.6 \, \text{MeV}$. The flavour symmetry predicts an approximate mass degeneracy for these states, but this is emphatically not the case for their widths. The reason is that the decay $\omega \to \pi\pi$ is strongly suppressed for the flavour singlet and the dominant decay is $\omega \to \pi\pi\pi$. The latter has much less phase space. Its first radial excitation, the $\omega(1420)$-meson, avoids this by decaying into $\rho\pi$. This phenomenon has nothing to do with our subject and, hence, we have to exclude the ω-meson from the averaging. Finally, we have only the ρ-meson, which gives

$$B(0) = \Gamma_\rho/m_\rho = 0.194. \tag{10}$$

It is interesting to note that this number was proposed in [24] as an educated guess in order to estimate the constant B in the real world. Surprisingly enough, it turned

out to be very close to a numerical estimate for B in the 't Hooft model (QCD in two dimensions in the large-N_c limit, see [25]) performed in [26].

Finally, our analysis yields the following estimates (see table 1):

$$B(0) \approx B(1) \approx 0.2, \quad B(2) \approx B(3) \approx B(4) \approx 0.1. \quad (11)$$

Thus, the educated guess made in [24] turns out to be correct for the next cluster ($N = 1$) as well. However, then one observes a sudden jump down by about 2 times. The Crystal Barrel Data, which was used to estimate $B(3)$ and $B(4)$, dramatically confirms this jump. How can we interpret this phenomenon? Looking at the states in clusters more attentively, one can make the following observations: a) the $N = 1$ cluster mainly consists of the ground states; b) the mesons in the $N = 0, 1$ clusters prefer to decay into two particles [4]; c) the mesons in the $N > 1$ clusters prefer to decay into three or four particles [4]. The nature of the phenomenon c) is enigmatic, at least for the author. This very phenomenon leads to a suppression of the available phase space for the decays. But how can one understand this within the effective hadron string? In the simplest case of open string one could assume, for instance, that the string breaks in two points simultaneously, producing three final particles. This is an $\mathcal{O}(1/N_c^2)$ effect. The question arises, why this effect might become dominant? A more plausible assumption is that the string decay is a cascade process for the excited states. Experimentally, one cannot detect the intermediate stages of this process. The universal quantity B is somehow decreased in this case.

Thus, the stability of numbers in eq. (11) supports the possibility of the effective string description. However, the experiment seems to tell us that this description for the excited states should be different from that of the ground states, at least with respect to the issue of string decays.

Last but not least. If one considers the individual channels, the result (11) hardly can be detected. In this respect the t' Hooft model provides an instructive example. It has so little degrees of freedom for exciting the bound states that each cluster consists of one state only. As a result, the quantity $B(n)$ has seemingly random fluctuations around a constant value, which is clearly seen when one computes the widths for several hundreds of radial excitations [24, 26]. Dealing with the first several states, this asymptotic value cannot be guessed at all. In four dimensions we are more lucky. The multitude of states in each cluster smooth significantly these fluctuations after the averaging. Due to this effect already a few clusters are able to provide the asymptotic value for B.

5 Conclusions

Our analysis shows that the available experimental spectrum of light non-strange mesons reveals the universal string-like behavior expressed by eq. (5) and, on average, a strong suppression of chiral symmetry breaking effects for sufficiently high resonances. The observed degeneracy

of the spectrum, however, cannot be explained by effective restoration of the chiral and axial symmetries only. A possible explanation is that a partial restoration of the conformal invariance happens simultaneously.

Independently of interpretation, the modern experimental data seems to point out two remarkable facts [27], which hold on average for the high excitations of light non-strange mesons. First, the spectrum globally behaves as that of the Lovelace-Shapiro dual amplitude (the intercept is the half of the slope). Second, the full decay width is proportional to the mass of decaying particle, just as expected within various string models.

New systematic experiments for the search of light non-strange hadrons above $1.9\,\text{GeV}$ are indispensable. Unfortunately, at present such experiments are not very widespread because they are not expected to bring a new physics. The analysis carried out in the paper is, in a sense, an attempt to overcome this prejudice. The detailed knowledge of the experimental spectrum for high meson excitations can help significantly to answer some fundamental questions and, hence, to extend our understanding of QCD.

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27. A curious historical remark: Balmer guessed the spectrum of excitations for the hydrogen atom just scrutinizing the available experimental data. Subsequently, the derivation of this spectrum within the Bohr model marked a milestone in the foundation of Quantum Mechanics. At present, we may be close to Balmer's situation in the case of light meson spectrum. . .