A novel multiquark approach to hadron resonances

S. S. Afonin

Saint Petersburg State University, 7/9 Universitetskaya nab., St.Petersburg, 199034, Russia
E-mail: s.afonin@spbu.ru

Abstract

A new approach to description of hadron spectroscopy is proposed. By assumption, the form of spectrum is dictated by the trace of energy momentum tensor in QCD. This provides the relativistic and renormalization invariance of hadron masses. The constructed scheme is applied to the light mesons for which two complementary interpretations are worked out. The first one represents an "atomic" structure of resonances above 1 GeV in which the quanta of non-perturbative gluon contributions are quantified via an effective formation of quasiparticles representing likely gluon analogues of positronium. This picture allows to build a "periodic table of the hadronic elements", i.e. to classify hadrons, in some sense, in analogy with Mendeleev table in chemistry. The classification does not require introduction of the orbital angular momentum associated with hadron constituents. The Regge and radial trajectories emerge in a natural way. The second interpretation is based on a "collisional" nature of some (or many) hadrons, specifically of scalar resonances below 1 GeV. Here the role of quasiparticles is played by another hadron. This picture in particular leads to a simple explanation of the puzzling scalar sector below 1 GeV with correct predictions for masses and dominant decay modes.
1 Introduction

The strong interactions between light $u$ and $d$ quarks are known to generate almost 99% of mass of the visible universe. In spite of many successes of Quantum Chromodynamics, the mechanism of hadron mass generation remains unclear. The problem is the most pronounced in the sector of light quarks where the underlying physics is highly non-perturbative and ultrarelativistic. The production of a large number of hadron resonances observed in the hadron collisions and leptonic processes [1] provides a striking manifestation of this physics. It is widely believed that the mass generation mechanism is enciphered in the spectroscopy of resonances. A wish to understand the rich hadron spectroscopy has driven an active model building since the 1960s. And although a great deal of efforts was invested in this enterprise in the last fifty years, a generally accepted analytical approach that would describe systematically the whole hadron spectrum has still not been elaborated. Such a situation could mean that some basic concepts which are being used for interpretation of hadron resonances and for model building may happen to be somewhat incomplete or even partly wrong. The proliferating observations of unconventional resonances in the heavy quark sector strengthens this impression [2]. It is not unlikely that for further advancing in the field of hadron spectroscopy one needs essentially new ideas and approaches.

Our research was inspired by the following problems. The standard interpretation of observable hadrons includes the non-relativistic notion of orbital angular momentum $L$ associated with quarks. In particular, $L$ dictates the spatial and charge parities for the quark-antiquark systems, $P = (-1)^{L+1}$ and $C = (-1)^{L+S}$, where $S$ is the quark spin. This picture does not explain the scalar mesons below 1 GeV for which a tetraquark structure is often assumed [3]. A question arises why do not we see many light multiquark hadrons above 1 GeV? Also the existence of $\pi_1$-mesons (among them the $\pi_1(1400)$ and $\pi_1(1600)$ are well established [1]) is not compatible with the standard quark model. Usually the $\pi_1$-mesons are interpreted as some hybrid quark-gluon states. A question then appears why do not we see other hybrid states with exotic quantum numbers, e.g. with $J^{PC} = 0^{--}$ or $2^{--}$, among light mesons? There are some deep theoretical questions as well. The light hadrons represent highly relativistic quantum systems in which $L$ is not a conserved quantity. The quark angular momentum $L$ is nevertheless a standard ingredient in constructing dynamical models for light hadrons. Another question concerns the observable quantities like hadron masses — they must be renorminvariant in the field-theoretical sense. Stated simply, they must represent some constants independent of energy scale. The quark masses, for instance, are not such constants since they have anomalous di-
mension in QCD. A relation of calculated observables to the renormalization invariance is obscure in almost all phenomenological models of hadrons.

In the present work, we put forward a principally new approach to the light hadrons in which the masses are relativistic and renorminvariant by construction. The approach allows to classify the light mesons without use of any angular momentum associated with hadron constituents. The states $\pi_1$ emerge in a natural way while the other exotic quantum numbers remain forbidden. The constructed mass counting scheme permits to obtain hadron masses from very simple relations with a typical accuracy comparable to numerical calculations in complicated dynamical models.

2 Motivation and idea of the approach

Consider the Gell-Mann–Oakes–Renner relation for pion mass [4],

$$m_\pi^2 = -\frac{\langle \bar{q}q \rangle}{f_\pi^2} (m_u + m_d) = \Lambda \cdot 2m_q,$$  \hspace{1cm} (1)

where we set $m_u = m_d = m_q$ and $\Lambda = -\frac{\langle \bar{q}q \rangle}{f_\pi^2}$. We will use the standard values for the quark condensate and masses of current quarks at the scale of the pion mass from the QCD sum rules [5] and Chiral Perturbation Theory (ChPT) [6], $\langle \bar{q}q \rangle = -(250 \text{ MeV})^3$, $m_u + m_d = 11 \text{ MeV}$. Together with $f_\pi = 92.4 \text{ MeV}$ (the value of the weak pion decay constant in the normalization used in [1]), the relation (1) yields $m_\pi = 140 \text{ MeV}$ and $\Lambda = 1830 \text{ MeV}$. The famous relation (1) can be derived in various ways as a consequence of the Spontaneous Chiral Symmetry Breaking (SCSB) in strong interactions. The interpretation of pion as the pseudogoldstone boson arising due to the SCSB in QCD became a common lore in particle physics.

Let us look at the relation (1) from a different angle. According to Quantum Mechanics, when a fermion-antifermion pair has the lowest energy it carries the quantum numbers of a pseudoscalar particle. This is the standard description of ground pseudoscalar particles in the hadron potential models where the notion of SCSB is absent. They have problems, however, with small pion mass and ensuing low-energy physics. We believe that the quantum-mechanical intuition about the pion is right and the problem with small pion mass appears because the relativistic and renormalization invariance of underlying theory are not properly reflected even in ”relativized” potential models. If this viewpoint is correct, there should exist an alternative way for derivation of relation (1), a way which does not use the assumption of SCSB. We do not have such a derivation in a rigorous way but we can outline a possible direction one should pursue to get it.
The mass of a hadron state $|h\rangle$ can be related to the trace of energy momentum tensor in QCD via the following Ward identity known in deep inelastic phenomenology [7],

$$2m_h^2 = \langle h|\Theta^{\mu}_\mu|h\rangle,$$

(2)

where $\Theta^{\mu}_\mu$ is given by the scale anomaly,

$$\Theta^{\mu}_\mu = \frac{\beta}{2g_s}G^2_{\mu\nu} + \sum_{\psi=u,d,...} m_\psi \bar{\psi}\psi.$$

(3)

Here $\beta$ denotes the QCD Beta-function. As follows from derivation of the identity (2), the squared mass (not the linear one!) in the l.h.s. appears due to the relativistic invariance and the general factor 2 stems from the symmetry between particles and antiparticles in relativistic Hamiltonians. The r.h.s. of (2) represents a renormalization invariant quantity. One can build two substantially different renormalization invariant operators in QCD, the both are present in (3). Applying the identity (2) to $\pi$-meson, one should prove that $\langle \pi|G^2_{\mu\nu}|\pi\rangle \neq 0$ and $\langle \pi|\bar{q}q|\pi\rangle = \frac{2f_\pi^2}{F_L}\langle 0|\bar{q}q|0\rangle$, $q = u, d$. This leads to the relation (1).

We assume thus that the relation (1) is a result of the Ward identity (2) applied to the lightest hadron. The final mass formula has a form enabling the interpretation based on the SCSB. The identity (2) suggests in this case that masses of other hadrons composed of $u$ and $d$ quarks can be given by extensions of (1) in which $\langle h|G^2_{\mu\nu}|h\rangle \neq 0$. Namely,

$$m_h^2 = \Lambda (E_h + 2m_q) = \Lambda E_h + m_\pi^2.$$

(4)

The quantity $E_h \sim \langle h|G^2_{\mu\nu}|h\rangle \neq 0$ depends on a hadron under consideration, with the combination $\Lambda E_h$ being renormalization invariant. As hadron masses are inseparably connected with properties of non-perturbative QCD vacuum, there should be a universal way to relate $E_h$ to the gluon condensate $\langle 0|G^2_{\mu\nu}|0\rangle \neq 0$, at least in the chiral limit. In the ansatz (4), we take into account that ”switching off” the explicit contribution from gluons we must come to the pion mass.

The simultaneous restrictions imposed by the relativistic and renormalization invariance are likely so strong that we do not have a large choice in constructing expressions for hadron masses. Below we develop a simple

\footnote{We plan to address this issue in a follow-up paper.}

\footnote{The vast majority of hadron resonances have no well defined gravitational (and inertial) mass as they do not propagate in space: Their typical lifetime of the order of}
model based on the ansatz (4) that describes the whole spectrum of light mesons.

The first excitation of pion is the $\rho$-meson. In hadron reactions with pions, this excitation can be interpreted as a quark spin flip caused by absorption of a long-distance gluon. Within the QCD sum rules, such non-perturbative processes are parametrized via a contribution of the gluon condensate
\[ \langle G_{\mu\nu}^2 \rangle = \langle G_{\mu\nu}^2 \rangle \] [5]. The given spin flip ”costs” energy $E_\rho \approx 310 \text{ MeV}$ in our model. The ansatz (4) yields then\(^3\) $m_\rho \approx 766.5 \text{ MeV}$. The crucial point here is that the non-relativistic logic (or rather the ”non-renorminvariant” one) $m_\rho = E_\rho + m_\pi$ does not work\(^4\) because $E_\rho$ is not renorminvariant: The value of 310 MeV refers to the scale where $\Lambda = 1830 \text{ MeV}$. The relativistic and renorminvariant relation between $m_\rho$ and $E_\rho$ is given by our ansatz (4). Under some assumptions, $E_\rho$ can be expressed via the values of vacuum quark and gluon condensate\(^5\)
\[ E_\rho = \frac{\beta_0 (\alpha_s / \pi) \langle G_{\mu\nu}^2 \rangle}{\langle \bar{q}q \rangle} \] , where $\beta_0 = 11 - \frac{2}{3} n_f$ and we take $\frac{\alpha_s}{\pi} \langle G_{\mu\nu}^2 \rangle = 0.012(3) \text{ GeV}^4$ [5]. For $n_f = 2$, the value $E_\rho \approx 310 \text{ MeV}$ is reproduced for the center value of the gluon condensate. In the present phenomenological model, however, we will treat $E_\rho$ as an input parameter, this is enough for our purposes. It is interesting to observe that the value of $E_\rho$ resembles a typical value of constituent quark mass\(^6\), $M_c \approx m_p / 3$, where $m_p$ is the proton mass.

10\(^{-23}\)–10\(^{-24}\)s does not allow to leave the reaction zone of the order of 1 fm which is their size simultaneously. They show up only as some structures in the physical observables at certain energy intervals. The resonance mass is commonly associated with the real part of an $S$-matrix pole on the second (unphysical) sheet. How this definition is related with the gravitational mass is by far not obvious in a theory with confinement. The Ward identity (2) can be considered as a definition of gravitational hadron mass that extends naturally to the masses of unstable resonances (assuming the existence of a limit where their gravitational masses are well defined).

\(^3\)The Particle Data reports $m_\rho = 775.11 \pm 0.34 \text{ MeV}$ for the charged $\rho$-meson [1]. This averaged mass is seen in leptonic processes [1]. In the hadronic processes, however, the averaged mass of charged $\rho$ is $m_\rho = 766.5 \pm 1.1 \text{ MeV}$ [1]. The difference can be easily explained: One needs an extra energy to create two current quarks in the leptonic processes and usually does not need this in the hadronic reactions. The latter situation is our case.

\(^4\)A remote analogy: The non-relativistic addition law for velocities, $v = v_1 + v_2$, does not work in the relativistic case and must be replaced by a non-linear law.

\(^5\)See the footnote 1.

\(^6\)Strictly speaking, the value of constituent quark mass is very model dependent — as far as we know, it ranges from 220 to 450 MeV in various models. The value $M_c(p^2) \approx 310 \text{ MeV}$ at small momentum $p$, however, proves to be seen in unquenched lattice simulations in the chiral limit (see, e.g., Ref. [8]). We may indicate a ”pocket” way for reproducing this value. When one incorporates the vector and axial mesons into the low-energy models describing the SCSB in QCD and related physics, e.g. into the Nambu–Jona-Lasinio model [9] or the linear sigma-model, a model-independent relation
3 Meson excitations above 1 GeV

We assume that the masses of higher spin and radial excitations are relativistic and renorminvariant quantities and can be given by the ansatz (4). The description with correct quantum numbers can be realized within the following scheme. We conjecture that gluodynamics leads to formation of isoscalar quasiparticles $A_0$ and $A_1$ which carry the pseudoscalar, $J^{PC} = 0^{--}$, and vector, $J^{PC} = 1^{--}$, quantum numbers, respectively. They represent gluonic counterparts of para- and ortho-positronium — the massive states that can be formed in photonic processes and "dematerializing" into pure photons. The quasiparticles $A_0$ and $A_1$, generally speaking, are not necessarily colorless. The underlying nature of these objects is not important for our further purposes. They reside likely under the valent $q\bar{q}$ shell and we will call them "underquarkonia" in what follows. The creation of underquarkonia requires energy $E_0$ for $A_0$ and $E_1$ for $A_1$. By assumption, these energies are proportional to the gluon condensate $\langle G^2_{\mu\nu} \rangle$, so that the resulting products $\Lambda E_0$ and $\Lambda E_1$ are renorminvariant. The energies $E_0$ and $E_1$ are "building blocks" which constitute the total excitation energy $E_h$ in (4). Within this scheme, the meson mass depends on the total number of underquarkonia inside the given meson. It is convenient to divide the non-strange resonances above 1 GeV into excited pions and excited $\rho(\omega)$ mesons. The excited $\rho$ mesons containing $n$ underquarkonia $A_0$ and $l$ underquarkonia $A_1$ will be labeled $J^{\rho A_0 A_1}$, where $J$ means the total spin. The mass of these excitations is given by

$$m_{\pi A_0 A_1}^2 = \Lambda (nE_0 + lE_1) + m_{\pi}^2.$$  \hspace{1cm} (5)

By assumption, the nucleus formed by underquarkonia is in the ground $S$-wave state. This leads to the maximal spin $J = l + 1$, the $P$ and $C$ parities obey to the standard multiplicative law: $(P, C) = ((-1)^{n+l+1}, (-1)^{l+1})$. The excited pions follow the same principle: The states $J^{\pi A_0 A_1}$ have mass

$$m_{\pi A_0 A_1}^2 = \Lambda (nE_0 + lE_1) + m_{\pi}^2.$$  \hspace{1cm} (6)

maximal spin $J = l$, and parities $(P, C) = ((-1)^{n+l+1}, (-1)^{l})$.

Since the pair $A_0 A_0$ does not change neither spin nor parities, a meson containing even number $2k$ of $A_0$ and no one $A_1$ will be referred to as the $k$-th radial excitation of the ground state, by definition. The radial excitations appear thus with a "period" $2E_0\Lambda$. The value of $E_0$ can be fixed from masses}

emerges: $m_{\rho A_0}^2 = m_{\rho}^2 + 6M_{c}^2$. On the other hand, the idea of SCSB was also exploited in the derivation of famous Weinberg relation, $m_{\rho A_0}^2 = 2m_{\rho}^2$ [10]. Combining both relations, one has $M_c = m_{\rho}/\sqrt{6}$ that gives the aforementioned value.
of pion radial excitations $\pi(1300)$ and $\pi(1800)$ which are $^0\pi_{A_0}$ and $^0\pi_{A_4}$ states in our notations. This fixes $E_0 \approx 440 - 450$ MeV. $E_1$ is fixed in the next Section.

It is easy to see that the excitations of the kind $^J\rho_{A^k}$, $k = 0, 1, 2, \ldots$, will give rise to a degenerate family of resonances with $J^{PC} = (1, 3, \ldots, 2k+1)^{--}$ (the even spins do not appear as explained in the next Section). The states on the main $\rho$-meson Regge trajectory follow thus with a "period" $2E_1\Lambda$. These excitations, in addition, generate daughter trajectories. For instance, the excited $\rho$ mesons emerge with $J^{PC} = 1^{--}$ which follow with the same "period" $2E_1\Lambda$. They are different from the "radial" $\rho$ mesons following with another period $2E_0\Lambda$.

Let us summarize the constructed picture for the excited light non-strange mesons above 1 GeV. A meson looks like an atom. Strictly speaking, it looks like a unusual positronium atom containing usual atomic nucleus. The role of electron and positron is played by valent quark and antiquark in the spin-singlet ("para") or spin-triplet ("ortho") states. The role of neutron and proton is played by the para- and ortho-underquarkonia. The mass of the usual atom, in the first approximation, is just sum of masses of all nucleons and electrons. In the case of hadronic atom, this mass counting rule contradicts both the relativistic and renormalization invariance. It must be replaced by the rule [3] for the ortho- and by [4] for the para-states. The spins of the ortho-underquarkonia must be added as spins of massless particles (see the next Section). Following these prescriptions, one can construct a "periodic table of the hadronic elements" describing the masses and quantum numbers of light non-strange mesons. In real situations, however, this idealistic picture is contaminated by contributions of the strange quark. A careful systematization requires a thorough phenomenological study of available data which we leave for a future work.

If underquarkonia are interpreted as real quark-antiquark pairs then the hadron excitations containing $A_0$ and $A_1$ represent multiquark states.

4 Some examples of meson excitations

Below we dwell on some meson excitations to demonstrate our scheme.

The state $^0\pi_{A_0}$ has the quantum numbers of scalar particle, $J^{PC} = 0^{++}$. Its mass is $m = \sqrt{\Lambda E_0 + m_{\pi}^2} \approx 910$ MeV. It could be $a_0(980)$ in which contributions from the strange quarks were not taken into account.

\footnote{We used $\pi(1800)$ as the uncertainty in its mass is much smaller than in the case of $\pi(1300)$.}
$^1\rho_{A_0}$ and $^1\omega_{A_0}$ have quantum numbers $J^{PC} = 1^{++}$ and mass about 1180 MeV. The natural candidates are $h_1(1170)$ and $b_1(1230)$ [1].

$^2J\rho_{A_1}$ and $^2J\omega_{A_1}$ have $J^{PC} = (0, 1, 2)^{++}$. They should be the series of states $a_1(1260), f_1(1285), f_2(1270), a_2(1320), \text{and likely } f_0(1370)$ [1]. Fitting to the masses of well measured resonances $f_1(1285)$ and $f_2(1270)$, we obtain the estimate $E_1 \approx 570$ MeV which will be used in what follows.

$^3J\rho_{A_2}$ and $^3J\omega_{A_2}$ of mass near 1490 MeV are first radial excitations of $\rho$ and $\omega$ mesons — the resonances $\rho(1540)$ and $\omega(1420)$ [1].

$^4J\rho_{A_0A_1}$ and $^4J\omega_{A_0A_1}$ have quantum numbers $J^{PC} = (0, 1, 2)^{+-}$, i.e. 3 possible spins and pion parities. These states describe resonances $\pi_1, \pi_1, \pi_2$ above 1 GeV. In contrast to the quark model, the $\pi_1$-meson is not exotic in our scheme! The relation (6) predicts mass about 1570 MeV. The possible candidates are $\pi_1(1600)$ and $\pi_2(1670)$ [1]. In the real hadron reactions, however, it seems that the pseudoscalar underquarkonium $A_0$ is sometimes replaced by the lighter $\pi^0$ meson. This may happen when photons come into play and participate in formation of underquarkonia along with gluons. In the case under consideration, a lighter state $^4J\rho_{A_0A_1}$ with mass near 1370 MeV can be formed instead of $^4J\rho_{A_0A_1}$. The resonance $^4J\rho_{A_0A_1}$ in particular is a candidate for $\pi_1(1400)$ (experimentally, $m_{\pi_1(1400)} = 1354 \pm 25$ MeV [1]). The state $^0J\rho_{A_0A_1}$ is another candidate for the radially excited pion. This may mean that $\pi(1300)$ has two possible structures — $^0J\pi_{A_0}^0$ and $^4J\rho_{A_0A_1}$ — that could be a reason for a large uncertainty in the mass of $\pi(1300)$, $m_{\pi(1300)} = 1300 \pm 100$ MeV [1]. In general, multiple possibilities for formation of a hadron resonance become rather common for highly excited states. This could be among reasons of the observed tendency of growing total width as a function of mass.

The states $^1J\rho_{A_1}$ and $^1J\omega_{A_1}$ have $J^{PC} = J^{--}$. A new feature appears in these states: We must add spins of two ortho-underquarkonia $A_1$. The standard spin counting rule would give $J = 0, 1, 2, 3$ for the total hadron spin. But our spin addition law will be different. To motivate this new rule consider first two photons. The photon is massless, hence, its spin has two projections only — two possible helicities. The total helicity of two photons can be 0 or 2. In other words, the addition law for spins of spin-1 massless particles is the same as for massive spin-$\frac{1}{2}$ particles, just the final result must be multiplied by 2. Consider now an ortho-positronium which decays into 3 photons (decay into 2 photons is forbidden). If 2 ortho-positronia are simultaneously created by 6 photons, their total spin can be 0 or 2 because the total helicity of original photons is even$^8$. We come thus to the conclusion that although ortho-positronia are massive spin-1 particles,

$^8$The total spin of ortho-positronia can be "rotated" to the value of 1 by some subsequent manipulations.
their spins must be added (immediately after formation) as if they were massless. The same rule we apply to gluons and underquarkonia. The total spin of the pair $A_1 A_1$ can be 0 or 2, hence, the total spin of $J^{PC} A_1^2$ can be $J = 1, 3$. A natural candidate is the couple of $J^{PC} = (1, 3)^{--}$ states $\rho(1700)$ and $\rho_3(1690)$ (and $\omega(1650)$ with $\omega_3(1670)$ for $\omega$) \cite{1}. Our results do not contradict the quark model predictions. First, the exotic states with quantum numbers $J^{PC} = (0, 2)^{--}$ are absent. Second, $\rho(1700)$ is not the second radial excitation of $\rho(770)$.

The recipe for constructing higher radial and spin excitations is straightforward. An interesting question emerge: Can the underquarkonia $A_0$ and $A_1$ materialize as real hadrons with pseudoscalar and vector quantum numbers? If this happens, the masses of pure underquarkonia should be estimated as $M_{A_i} = \sqrt{\Lambda E_i}$, $i = 0, 1$, according to our ansatz. The mass of pseudoscalar "materialized" underquarkonium, $M_{A_0} \approx 910$ MeV, is suspiciously close to the mass of $\eta'$-meson, $m_{\eta'} = 958$ MeV \cite{1}. The vector underquarkonium would materialize with the mass $M_{A_1} \approx 1020$ MeV which coincides with the $\varphi$-meson mass. These observations suggest that the idea of underquarkonia might be able to give an alternative non-instanton solution for the old $U(1)_A$ problem, i.e. to explain why $\eta'$ is much heavier than it is prescribed by the quark model, and to clarify why the strange component in $\varphi$ is almost unmixed with the lighter quarks.

5 Regge trajectories

The relations (5) and (6) lead to linear Regge, equidistant daughter Regge and radial trajectories. Below we give examples for some of them.

The states $0\pi A_2^0$, $n = 0, 1, 2, \ldots$, form linear trajectory for the radial excitations of pion,

$$m_\pi^2(n) = 2\Lambda E_0 n + m_\pi^2.$$  \hfill (7)

The radial excitations of $\rho(770)$ — the resonances $^1\rho A_2^0$ — lie on the first radial $\rho$-trajectory,

$$m_\rho^2(n)_I = 2\Lambda E_0 \left(n + \frac{E_\rho}{2E_0}\right) + m_\pi^2.$$  \hfill (8)

The second radial $\rho$-trajectory is composed of the states $^1\rho A_2 A_0^0$,

$$m_\rho^2(n)_H = 2\Lambda E_0 \left(n + \frac{E_1}{E_0} + \frac{E_\rho}{2E_0}\right) + m_\pi^2.$$  \hfill (9)

\footnote{In the quark potential models, $\rho(1700)$ (and $\omega(1650)$) is a $D$-wave state, while the radial excitations of $\rho(770)$ should be $S$-wave ones.}
It is evident that the states $^1\rho_{A_{1}^{2S-0}}$ formally give rise to the $k$-th radial $\rho$-trajectory. The resonances having structure $^1\rho_{A_{1}A_{2}^0}$ form the first radial $a_1$-trajectory,

$$m_{a_1}^2(n)_I = 2\Lambda E_0 \left(n + \frac{E_1}{2E_0} + \frac{E_\rho}{2E_0}\right) + m_\pi^2.$$  \hspace{1cm} (10)

The further axial radial trajectories can be easily written. It should be remarked that the standard quark model allows of only two radial $\rho$-trajectories ($S$- and $D$-wave ones) and one $a_1$-trajectory (a $P$-wave one). Our scheme is richer.

The spin $\rho$-trajectory is composed from the states $^J\rho_{A_{1}J-1}$ with $J = 1, 3, 5, \ldots$. The corresponding masses are

$$m_{\rho_J}^2 = \Lambda E_1 \left(J - 1 + \frac{E_\rho}{E_1}\right) + m_\pi^2.$$  \hspace{1cm} (11)

For the even spins, $J = 2, 4, 6, \ldots$, the trajectory (11) describes $a_J$ ($f_J$) mesons. The Regge trajectory (11) describes thus states with alternating parities and this agrees with the phenomenology.

The slopes of radial trajectories (7)–(10) are universal. This property was observed in the phenomenology of light mesons [11]. The slope $2\Lambda E_1$ of spin Regge trajectory (11) is different from the radial one as it is determined by the energy $E_1$ of different underquarkonium. If we consider all spins in (11) then the slope is $\Lambda E_1 \approx 1.83 \cdot 0.57 \approx 1.04$ GeV$^2$. This value is in accord with the phenomenology [11]. The phenomenological value of Regge slope can be taken, in principle, from the experimental data to fix $E_1$.

The principal $\rho$-meson Regge trajectory (11) is accompanied by the daughter trajectories following with the step $2\Lambda E_1$. The spin-1 $\rho$-mesons of the kind $^1\rho_{A_{1}J-1}$ are the lowest states on the daughters. For example, the lowest state in (9) is the lowest state on the first daughter. The spectrum of $^1\rho_{A_{1}J-1}$ excitations reads

$$m_{\rho_{A_{1}J-1}}^2 = 2\Lambda E_1 \left(k + \frac{E_\rho}{2E_1}\right) + m_\pi^2,$$  \hspace{1cm} (12)

where $k = 0, 1, 2, \ldots$ enumerates the daughters. Similar relations can be written for the axial and other mesons. An obvious consequence of the emerging spectrum is the degeneracy of spin and daughter radial excitations of the type $m^2(J, k) \sim J + k$, which is typical for the Veneziano dual amplitudes [12] and the Nambu–Goto open strings.

The values of $E_\rho$, $E_0$ and $E_1$ are of similar magnitude. This allows to consider a reasonable limit where they are equal. The notions of radial and
daughter radial trajectories coincide in the limit $E_0 = E_1$. In the limit $E_\rho = E_1$, the radial vector and axial trajectories are related by

$$m_{a_1}^2(n) = m_\rho^2(n) + m_\rho^2 - m_\pi^2. \quad (13)$$

This relation holds both for radial and for daughter radial trajectories. In the chiral limit, $m_\pi = 0$, the relation (13) for the ground states, $n = 0$, reduces to the old Weinberg relation, $m_{a_1}^2 = 2m_\rho^2 \ [10]$. In the most symmetric limit, $E_\rho = E_0 = E_1$, the vector and axial radial spectrum reduce to a simple form in the chiral limit,

$$m_\rho^2(n) = 2m_\rho^2 \left( n + \frac{1}{2} \right), \quad (14)$$

$$m_{a_1}^2(n) = 2m_\rho^2 (n + 1), \quad (15)$$

The relations (14) and (15) first appeared in the variants of Veneziano amplitude which incorporated the Adler self-consistency condition [12]. This condition (the amplitude of $\pi\pi$ scattering is zero at zero momentum) removes degeneracy between the $\rho$ and $a_1$ spectra. Within the QCD sum rules, the relations (13) and (15) may be interpreted as a large-$N_c$ generalization of the Weinberg relation [13].

We see thus that in certain limits the Regge phenomenology of our approach reproduces various known relations.

Let us clarify further how the excited resonances with identical quantum numbers can have different origin in the proposed scheme. They may represent the radial states, states on daughter Regge trajectories and various "mixed" ones. For instance, the second $\rho$-meson excitation with the same parities (i.e. containing 2 pairs of underquarkonia) appears in three forms: the second radial excitation $^1\rho_A^2$, the vector state on the second daughter trajectory $^1\rho_{A_2}$, and the mixed one $^1\rho_{A_2A_2}$. They are degenerate only in the limit $E_0 = E_1$. It is likely difficult to detect such a splitting experimentally because of overlapping widths. In reality, one would observe rather a "broad resonance region".

We note also that placing the observed "radial" states on a certain trajectory should be made with care — an incorrect interpretation of states leads to a false (or more precisely, introduced by hands) non-linearity of the trajectory. Take again the $\rho$ meson as a typical example. The first 3 radial excitations of $\rho$ are the states $^1\rho_{A_0}$, $^1\rho_{A_2}$, and $^1\rho_{A_4}$. They are accompanied by the following states with the quantum numbers of $\rho$: $^1\rho_{A_2}$, $^1\rho_{A_2A_2}$, and $^1\rho_{A_4}$. Since $\frac{3}{2}E_0 > E_1 > E_0$ in our fits, the sequence of first 7 $\rho$-mesons is: $\rho$, $^1\rho_{A_2}$, $^1\rho_{A_2}$, $^1\rho_{A_4}$, $^1\rho_{A_2A_2}$, $^1\rho_{A_2}$, $^1\rho_{A_4}$. They likely correspond to the vector resonances $\rho(770)$, $\rho(1450)$, $\rho(1700)$, $\rho(1900)$, $\rho(2000)^*$, $\rho(2150)$, $\rho(2270)^*$. 

11
Here the states marked by asterisk are taken from ”Further States” in the Particle Data [1].

6 σ-meson as a collisional excitation

The valent electron in atom can be excited in two ways: (a) via absorption of photon (with ensuing spin flip due to momentum conservation); (b) via a direct collision with some particle. Looking at an excited atom at some moment in time $t_0$ we cannot say how it was excited without any information got prior to the moment $t_0$. There exists the third way for changing energy of atomic electron: (c) via a change of nucleus charge caused, e.g., by absorption of α-particle.

Let us apply these electromagnetic analogies to the pion excitations. Within our approach, the case (a) corresponds to the ρ-meson. The case (c), in some sense, corresponds to the formation of underquarkonia. But what about the option (b)? In practical situations of collisions of hadronic beams, the ”collisional” excitations should play an important or even very important role.

Consider a low-energy $\pi\pi$ scattering. We may have a situation when a pion collides as a whole with one of quarks of another pion. This may happen when the first pion is faster, hence, has smaller de Broglie wavelength, i.e. when one pion wave packet penetrates another one[10]. The collision lasts a very short time $\Delta t$ which determines the lifetime of the formed coherent state. During the time $\Delta t$, the second quark (quark-spectator) ”feels” the first one as a particle with unchanged color charge and spin (because pion is colorless and spinless) but with different proper mass: $m_q \rightarrow m_q + m_{\pi}$.

According to our basic principle (4), this coherent state, let us denote it $\sigma$ beforehand, obeys to the mass relation

$$m_{\sigma}^2 = \Lambda m_{\pi} + m_{\pi}^2,$$

that yields $m_\sigma \approx 525$ MeV. This state must have the quantum numbers of scalar meson, decay into two pions, and should be an extremely short-living particle. The formation of such a coherent state is favored by Coulomb attraction between $\pi^+$ and $\pi^-$ which leads to the dominance of isosinglet channel. All these observations strongly suggest that we must interpret this resonance as the $f_0(500)$-meson [1], widely known as $\sigma$-meson.

We may propose a quantum-mechanical interpretation for the relation (16). Consider a stationary state of some quantum system with energy $E$ which

\[\text{\footnote{A picture of this process depends on a reference frame.}}\]
is described by the Schrödinger equation, \( H |\psi\rangle = E |\psi\rangle \). Consider now a
sudden perturbation of this system, e.g. by a very fast particle. A sudden
perturbation in Quantum Mechanics is a perturbation lasting so short time
\( \Delta t \) that the wave function \( \psi \) does not change during \( \Delta t \) (such a change always
requires a finite time). The perturbed system is described by the equation
\( H' |\psi\rangle = E' |\psi\rangle \). But \( \psi \) is not an eigenfunction of the new Hamiltonian \( H' \)
and the perturbed system becomes unstable. An important point for us is
that the form of \( \psi \) determines the functional dependence of energy \( E \) from
the parameters \( a_i \) of the Hamiltonian \( H, E = E(a_i) \). Since \( \psi \) is unchanged
during the time \( \Delta t \), this dependence remains the same after the perturbation.
It means that if \( H' \) is determined by a set of perturbed parameters
\( a'_i \) then \( E' = E(a'_i) \). Returning to the \( \pi\pi \) scattering, one of pion can be
considered as a stationary system in which \( m_q \) plays the role of a parameter
determining its energy \( m_\pi \). A collision causes a sudden perturbation that
changes the parameter, \( m_q \to m_q + m_\pi \), for a short time \( \Delta t \). This change
must be substituted to (1) (for one of quarks) to obtain the energy (16) of
unstable state \( \sigma \).

Within our approach, the main difference of a collisional resonance from
the resonances considered above consists in the absence of explicit renormin-
variance of its mass: The product \( \Lambda m_\pi \) in (16) is not renorminvariant and \( \Lambda \)
should be taken at the scale of \( m_\pi \). If we replace pion in (16) by some other
particle, we should take \( \Lambda \) at the scale of that particle.

The recent lattice simulation of Ref. [14] reported an effect of evolving
\( \sigma \)-meson into a stable bound state lying below the \( \pi\pi \) threshold as \( m_\pi \) is
increased. This observation follows directly from the mass relation (16):
Imposing \( m_\sigma \geq 2m_\pi \) we obtain \( m_\pi \leq \Lambda/3 \). This restriction is nontrivial as
\( \Lambda \) depends on the mass scale. If we normalize \( \Lambda \) to the numerical results of
Ref. [14], \( m_\sigma = 758 \text{ MeV} \) when \( m_\pi = 391 \text{ MeV} \), we get \( \Lambda = 1078 \text{ MeV} \) that
gives the estimate \( m_\pi \leq 359 \text{ MeV} \). This restriction agrees with the lattice
results of Ref. [14]: \( \sigma \) represents a bound state at \( m_\pi = 391 \text{ MeV} \) and a
broad resonance at \( m_\pi = 236 \text{ MeV} \).

7 \( h_1 \) and \( b_1 \) as collisional excitations

In the potential quark models, the lightest scalar and axial-vector mesons
represent \( P \)-wave states and, as a result, their masses are rather close since
the difference stemming from the spin-orbital interaction is relatively small.
Within the picture of "hadronic atom", we have identified the lightest axial
resonances \( h_1 \) and \( b_1 \) as \( ^1P_{A_0} \) or \( ^1\omega_{A_0} \) states. But we can give also a collisional
interpretation for theses states. Imagine that a pion collided with a \( \rho \)-meson
and this \( \rho \), having a smaller de Broglie wavelength, excited one of quarks inside the pion in the collisional way. Let us denote the formed coherent state \( \pi_\rho \) (means "\( \rho \) inside \( \pi \)"). The \( \sigma \)-meson represents the \( \pi_\sigma \) collisional resonance in this notation. According to our prescriptions, the mass of a collisional state of the kind \( \pi_\rho \) is given by the following extension of relation (16),

\[
m^2_{\pi_\rho} = \Lambda m_\rho + m^2_{\pi},
\]

(17)

where \( \Lambda \) should be taken at the scale \( m_\rho \). For making estimates in the first approximation we will consider \( \Lambda \) as a universal constant and set as before \( \Lambda = 1830 \text{ MeV} \). Taking in (17) \( h = \rho \), we obtain \( m_{h_1} \approx 1190 \text{ MeV} \) that is very close to our previous result. As in the \( \sigma \)-case, the formation of the coherent state \( \pi_\rho \) should be favored by the Coulomb attraction, i.e. the favored channel is \( \pi^+_{\rho^-} \) or \( \pi^-_{\rho^+} \) that entails zero isospin. A natural consequence of \( \pi_\rho \) structure of \( h_1 \) is the absolute dominance of the decay mode \( h_1 \rightarrow \rho \pi \). The isotriplet partner of \( h_1(1170) \) — the \( b_1(1230) \) meson — represents the collisional resonance \( \pi_\omega \) that determines its isospin 1 (it inherits the pion isospin) and dominant decay \( b_1 \rightarrow \omega \pi \). \( b_1 \) is expected to be heavier than \( h_1 \) because \( m_\omega > m_\rho \) and narrower than \( h_1 \) because \( \Gamma_\omega < \Gamma_\rho \). These expectations agree with the experimental data [1], at least qualitatively.

It should be noted that the state \( ^1\pi_{A_1} \) has the quantum numbers of \( b_1 \) as well. The mass of this state is, however, about \( 1030 \text{ MeV} \). We should hence explain the suppression of this state within the atomic picture — a problem that does not arise in the collisional one. The suppression seems to occur dynamically: As the energy \( E_1 \) is accumulating inside a pion, the system prefers first to turn into the \( \rho \) meson since \( E_\rho < E_1 \) and the rest of energy \( E_1 - E_\rho \) is not enough for further excitation of \( \rho \).

We see thus that the axial \( h_1 \) and \( b_1 \) mesons can be interpreted both as "hadronic atoms" and as collisional resonances. This duality takes place due to an interesting fine-tuning of parameters, \( m_\rho \approx E_\rho + E_0 \), that provides an approximate coincidence of predicted masses in both interpretations. Some (may be even all?) higher excitations allow the dual interpretation as well. The collisional interpretation is convenient for explaining dominant decay modes and isospin while the atomic one suits better for deriving various mass relations between different states, the Regge trajectories are a good example.

\[ ^{\text{11}} \text{The state } ^1\eta_{A_1} \text{ nevertheless predicts the correct mass of 1170 MeV for the } h_1(1170) \text{ meson (see the next Section).} \]
8 Towards inclusion of strange quark

The ChPT predicts that all masses in the octet of pseudogoldstone bosons follow from (1) if one replaces \( m_q \) by \( m_s \) as prescribed by the quark model [6]. An important point here is that \( \Lambda \) remains universal. We will use this approximation to extend our mass counting scheme to light mesons containing the strange quarks.

The quark content \( q \bar{s} \) and \((u \bar{u} + d \bar{d} - 2s \bar{s})/\sqrt{6}\) of pseudoscalar \( K \) and \( \eta \) mesons leads to the ChPT prediction for their masses:

\[
m^2_{K} = \Lambda(m_q + m_s) \quad \text{and} \quad m^2_{\eta} = \Lambda(2m_q + 4m_s)/3.\]

The vector \( K^* \) and \( \varphi \) mesons have the quark content \( q \bar{s} \) and \( s \bar{s} \), respectively. According to our ansatz (4), the corresponding masses are

\[
m^2_{K^*} = \Lambda(E_\rho + m_q + m_s) = \Lambda E_\rho + m^2_K, \quad \text{and} \quad m^2_{\varphi} = \Lambda(E_\rho + 2m_s). \tag{18} \tag{19}
\]

The value of \( m_s \) at the scale 1 GeV is about \( m_s \approx 130 \text{ MeV} \) [1]. Accepting this value we obtain \( m_{K^*} \approx 900 \text{ MeV} \) and \( m_\varphi \approx 1020 \text{ MeV} \) in a very good agreement with their experimental values [1]. Combining Eqs. (18) and (19) with (5) for the ground \( \rho \)-meson, we arrive at the Gell-Mann–Okubo mass relation,

\[
m^2_\rho + m^2_\varphi = 2m^2_{K^*}. \tag{20}
\]

In the standard quark model, the relation (20) holds for the linear masses. The relativistic nature of our approach leads to masses squared in (20) automatically.

The higher spin and radial excitations can be built following the identical scheme. The same applies to the collisional excitations. For instance, the coherent state \( K_{K^*} \) has mass \( m_{K_{K^*}} = \sqrt{\Lambda(m_q + m_s + m_{K^*})} \approx 1380 \text{ MeV} \), quantum numbers of \( h_1 \), and decays into \( K \bar{K}^* \) or \( K K^* \). This state is observed as the resonance \( h_1(1380) \) [1].

9 The scalar sector below and near 1 GeV

We have already interpreted \( \sigma \)-meson as the collisional \( \pi_s \) state and \( h_1(b_1) \) as the \( \pi_\rho(\pi_\omega) \) one. Adding now the \( K \) and \( \eta \) mesons we can construct other scalar collisional states which can be formed, e.g., in the \( K \pi \) scattering.

Consider the state \( \pi_\kappa \). Setting \( h = K \) in (17) we obtain its mass \( m_{\pi_\kappa} \approx 970 \text{ MeV} \). As in the case of \( \sigma \) and \( h_1 \), the expected isospin of \( \pi_\kappa \) is zero. Its natural decay mode would be \( \pi_\kappa \to K \pi \) but such a decay is forbidden by the isospin conservation if \( \pi_\kappa \) represents a genuine resonance. \( \pi_\kappa \) should be then relatively narrow and, as its mass lies slightly below the \( K \bar{K} \) threshold,
its dominant decay mode is expected to be $\pi_K \rightarrow \pi\pi$. The scalar resonance $f_0(980)$ meets all these expectations \[1\].

Let us include now the $\eta$ meson. Taking $h = \eta$ in (17) we predict the characteristics of $\pi\eta$ resonance: $m_{\pi\eta} \approx 1010$ MeV, $I_{\pi\eta} = 1$ (it inherits $I_\pi$), the dominant decay mode $\pi\eta \rightarrow \eta\pi$ and it should be broader than $\pi_K$ because this mode is not forbidden. The scalar resonance $a_0(980)$ satisfies these predictions \[1\].

Consider a hypothetical $K\pi$ collisional resonance. The formation of the coherent state $K\pi$ is much harder than $\pi K$ because $\pi$ has larger de Broglie wavelength, i.e. the pion wave packet is larger then the kaon one. This is also true for the measured mean sizes, $\langle r_\pi \rangle > \langle r_K \rangle$ \[1\]. But one might assume that the probability of $K\pi$ were non-zero due to some quantum effects. The Coulomb attraction would favor then the $K^+_\pi$ or $K^-_\pi$ channel, the $K^{\pm}_0$ or $K^{\pm}_0$ are much less plausible. In any case, the mass of $K\pi$ would be given by (we make the replacements $\pi \rightarrow K$ and $h \rightarrow \pi$ in (17)),

$$m^2_{K\pi} = \Lambda(m_\pi + m_\eta + m_s) = \Lambda m_\pi + m^2_K,$$

resulting in $m_{K\pi} \approx 710$ MeV. It is tempting to interpret $K\pi$ as the unconfirmed scalar resonance $K^{*}_0(800)$ called also $k$-meson \[1\]. The Particle Data reports the following mass of this elusive resonance: $682 \pm 29$ MeV \[1\]. Comparison of (21) with (16) shows that, as expected, $K\pi$ would be a partner of $\sigma$ in which one of $u$ or $d$ quarks is substituted by the $s$ quark. The observed isospin $I_k = \frac{1}{2}$, however, contradicts the favorable isospin zero predicted by the assumed mechanism of $K\pi$ formation. We see thus that the existence of $k$ meson, if confirmed, is not in conflict with our general principle for collisional resonances and one should look for a correct formation mechanism.

We propose the following explanation. Let us take a closer look at the production of $k$-resonance. The main source of information on $k$ are decays of $J/\psi$ meson into kaons and pions. The decays of vector charmonia are always accompanied by an abundant photon background. In all this ”mixture”, one can have situations when photons produce $\pi^0$ meson inside a kaon. The formed coherent state would then inherit the kaon isospin, i.e. would give rise to the scalar partners of the pseudoscalar $K^+, K^-, K^0$, and $K^{*0}$ mesons, with the mass being described by the relation (21).

Our approach predicts other collisional scalar resonances as well: $K_\eta$ with mass about 1120 MeV, $\eta_q$ with mass about 1150 MeV, almost unfeasible $\eta_s$ having mass near 760 MeV and a formation mechanism similar to that of $k$ meson, and resonances with $\eta'$ like the $\pi_{\eta'}$ state of mass about 1330 MeV. It is likely very hard to detect these resonances in the $\pi\pi$, $K\pi$ and $KK$ scattering (in experiments of direct $\eta\pi$ and $\eta K$ scattering this would be easier) but they may contribute to the strong background emerging in these reactions.
Within the framework of our collisional interpretation, the scalar resonances below and slightly above 1 GeV represent thus two-meson states. In terms of the quark degrees of freedom, they are tetraquarks, as is also suggested by many other models and observations [3].

10 Once more on inclusion of strange quark

The inclusion of strange quark considered above is not yet a full extension of our approach to the strange sector. We must switch on the third quark flavor in the trace of energy-momentum tensor [4]. The $\beta$-function changes slightly in the gluon part and this numerical change can be neglected in the first approximation. The main effect comes from the appearance of operator $m_s\bar{s}s$ in the quark part. The universality of $\Lambda$ used above is related with the assumption $\langle \bar{s}s \rangle = \langle \bar{q}q \rangle$ which is compatible both with the ChPT [6] and QCD sum rules [5].

We have already observed an automatic fine-tuning of parameters $m_{\rho} \approx E_{\rho} + E_0$ that enables a dual interpretation of the axial $h_1$ meson. Our approach contains more interesting "coincidences". Taking as before $m_s = 130$ MeV at the scale 1 GeV, one can observe the following relations between parameters: $E_0 \approx E_{\rho} + m_s$ (perhaps, $E_0 = E_{\rho} + m_s + 2m_q$) and $E_1 = E_{\rho} + 2m_s$. It is the second relation that leads to equality of masses of $\varphi$ meson [4] and "materialized" vector underquarkonium $A_1$. The phenomenological values of $E_0$ and $E_1$ seem to indicate that in quantum world we cannot build a spectroscopy of excited "pure non-strange" mesons: If there is available excitation energy $E_{\rho} \approx 310$ MeV then there is enough energy for creating an $s\bar{s}$ pair and the strange quarks intervene. The matter looks as if we added the third quark flavor to (4), $2m_q \rightarrow 2m_q + m_s$, and after that ascribed $m_s$ to a redefined $E_h$, $E_h + m_s \rightarrow E_h$. The "mass-symmetric" limit $E_{\rho} = E_0 = E_1$ considered in our discussion of Regge trajectories means then the limit $m_s = 0$. An explanation of observed "quantization" of $E_0$ and $E_1$ with respect to $m_s$ is challenging [12].

The light non-strange excited mesons have usually some decay modes with the strange component. Within the standard picture, masses of these states do not depend on $m_s$ and the $s\bar{s}$ pair pops up from the vacuum triggering the

\[ E_{\rho} + E_1 + m_\pi = 2(m_s + E_{\rho}) + m_\pi = m_\varphi. \]

This relation allows of an obvious non-relativistic interpretation: The $\varphi$ meson is composed of two constituent strange quarks with mass $m_s + E_{\rho}$ and binding energy $m_\pi$. The "free" constituent strange quarks would have the mass $m_s + E_{\rho} + m_\pi/2 = 510$ MeV which is exactly the value usually extracted from the observed magnetic moment of $\Lambda$ baryon within a naive quark model [15].

12 We would indicate another one "coincidence": $E_{\rho} + E_1 + m_\pi = 2(m_s + E_{\rho}) + m_\pi = m_\varphi$. This relation allows of an obvious non-relativistic interpretation: The $\varphi$ meson is composed of two constituent strange quarks with mass $m_s + E_{\rho}$ and binding energy $m_\pi$. The "free" constituent strange quarks would have the mass $m_s + E_{\rho} + m_\pi/2 = 510$ MeV which is exactly the value usually extracted from the observed magnetic moment of $\Lambda$ baryon within a naive quark model [15].
corresponding decay modes. The observed relation $E_1 \simeq E_0 + m_s \simeq E_\rho + 2m_s$ may suggest that, in reality, the strange quark does contribute to their masses via the contribution of $m_s$ to $E_0$ and $E_1$.

11 Conclusions

We have proposed a novel approach to classification of hadrons and description of hadron spectroscopy. The approach is very simple, allows of a clear visual interpretation and gives possibility to estimate hadron masses by a trivial arithmetics on any available calculator.

In brief, we conjectured an ansatz for hadron masses that should follow from the trace of energy-momentum tensor in QCD, providing thereby the relativistic and renormalization invariance of masses. The Gell-Mann–Oakes–Renner relation for pion mass turns out to be a particular case of this ansatz. The heavier non-strange mesons appear via excitation of one of quarks inside pion. A quark-spectator ”feels” a sudden perturbation of energy-momentum of the excited quark and this affects the hadron mass in the same manner as a sudden increase of mass of one of quarks. The excitation is caused either by the gluon exchange or in the collisional way in hadron-hadron collisions when a hadron with smaller wave packet scatters as a whole on one of quarks of another hadron. For the case of gluon excitations, we developed a model of ”atomic” structure of excited mesons, where the nucleus is formed by gluon quasiparticles of two kinds — the spin-singlet and spin-triplet one. They are not necessarily colorless and their energies are not renorminvariant. But together with valent quarks they form a colorless hadron with renorminvariant mass. The given picture leads to a new systematics of light mesons above 1 GeV, including the exotic $\pi_1$-states as an integral part. This systematics does not use the non-relativistic notion of orbital angular momentum associated with hadron constituents, or more precisely, it is equal to zero. The proposed mass counting scheme reproduces various long-known relations for hadron masses. The ”atomic” picture happens to be smoothly connected with the collisional one near 1 GeV due to a miraculous fine-tuning of parameters occurring somehow automatically. The collisional picture explains the whole scalar sector below 1 GeV in a remarkably simple way — the sector, that is the hardest for traditional approaches, turns out to be the easiest one within the proposed scheme.

In the given paper, the approach was developed for the light mesons. The presented analysis is partly raw (and perhaps the constructed scheme is still incomplete) and does not explain all state-by-state properties of meson spectrum but we consider it as a possible step forward in a right direction.
The extension to the light baryons is straightforward and is planned for a future work, together with some theoretical justifications. It seems that many highly excited $N$ and $\Delta$ baryons can be described as collisional excitations of the kind $M_B$, where $M$ is a meson (typically the $\pi$ meson and resonances which are abundantly produced in reactions with $\pi$, like $\rho$ ($\omega$) and $f_J$ mesons) and $B$ is a baryon (typically the proton and $\Delta(1232)$). An extension to the heavy quarks is in progress. In the heavy sector, it is important to understand a transition from the relativistic to non-relativistic picture. The interpretation of many unconventional heavy resonances observed recently [2] is especially challenging.

It must be emphasized that the approach proposed in this work is broader than ”just another one model” as it gives a new language for discussion of hadron resonances, for interpretation of data in the hadroproduction and formation experiments, and a possible starting point for construction of essentially new dynamical models.

References

[1] C. Patrignani et al. (Particle Data Group), Chin. Phys. C 40, 100001 (2016).
[2] H. X. Chen, W. Chen, X. Liu and S. L. Zhu, Phys. Rept. 639, 1 (2016).
[3] J. R. Pelaez, Phys. Rept. 658, 1 (2016).
[4] M. Gell-Mann, R. J. Oakes and B. Renner, Phys. Rev. 175, (1968) 2195.
[5] M. A. Shifman, A. I. Vainstein and V. I. Zakharov, Nucl. Phys. B 147, 385, 448 (1979).
[6] J. Gasser and H. Leutwyler, Nucl. Phys. B 250, 465 (1985).
[7] R. L. Jaffe and A. Manohar, Nucl. Phys. B 337, 509 (1990).
[8] P. O. Bowman et al., Phys. Rev. D 71, 054507 (2005).
[9] S. P. Klevansky, Rev. Mod. Phys. 64, 649 (1992).
[10] S. Weinberg, Phys. Rev. Lett. 18, 507 (1967).
[11] A. V. Anisovich, V. V. Anisovich and A. V. Sarantsev, Phys. Rev. D 62, 051502(R) (2000); D. V. Bugg, Phys. Rept. 397, 257 (2004).
[12] P. D. B. Collins, *An Introduction to Regge Theory and High-Energy Physics* (Cambridge University Press, Cambridge, 1977).

[13] S. S. Afonin, Phys. Lett. B 576, 122 (2003).

[14] R. A. Briceno, J. J. Dudek, R. G. Edwards and D. J. Wilson, Phys. Rev. Lett. 118, 022002 (2017).

[15] F. Halzen, A. D. Martin, *Quarks and Leptons* (John Wiley & Sons, New York, 1984).