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Chiral symmetry restoration in excited hadrons and dense matter[*]

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Abstract We overview two interconnected topics: possible effective restoration of chiral symmetry in highly excited hadrons and possible existence of confined but chirally symmetric matter at low temperatures and high densities.

Key words chiral symmetry, confinement, hadrons, QCD phase diagram.

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1 Introduction

The question of mass generation and the related question of interconnections of confinement and chiral symmetry breaking are central for QCD. In order to answer these questions we have to understand the gross structure of the hadron spectrum in the light quark sector and correlate it with interactions of hadrons with the Nambu-Goldstone bosons of broken chiral symmetry. Given symmetry patterns observed in the hadron spectra [1, 2], their correlations with the couplings to pions [3], or, more generally, with their axial properties, one can obtain the insight into the principal mechanism responsible for the mass generation, whether or not the mass of hadrons is directly related to the quark condensate of the vacuum. There are strong hints that in the highly excited hadrons the physics of mass generation is mostly unrelated to spontaneous breaking of chiral symmetry in the vacuum, i.e. most part of the hadron mass is not due to the quark condensate of the vacuum. This phenomenon, if correct, is referred to as effective restoration of chiral symmetry. This is just in contrast to physics of the lowest lying hadrons like nucleon, where mass is mostly induced by the quark condensate.

The issues of mass generation and interconnections of confinement and chiral symmetry are critically important for our view of the QCD phase diagram. For many years it was believed that in the con-

fining mode chiral symmetry should be strongly broken. This is certainly true in the vacuum, as it follows from the model-independent 't Hooft anomaly matching conditions [4]. Extrapolating these constraints to the nonzero temperature and density regions, one naively concludes that there cannot be a phase in the QCD phase diagram that is confining but with vanishing quark condensate. This picture was supported by simple models of confinement and chiral symmetry breaking. It is this argument which was a basis for a belief that the deconfinement and chiral restoration phase transitions coincide in the temperature - chemical potential plain. Then, given this apriori belief, the QCD phase diagram was modeled after the Nambu and Jona-Lasinio model phase diagram (or variants of thereof). This model is nonconfining.

In the large N_c limit QCD is confining at low temperatures up to arbitrary large density and such a matter was called quarkyonic [5]. Then, at some reasonably large density one can obtain a confining but chirally symmetric phase at low temperatures [6–8]. In such a phase the standard quark-antiquark condensate of the QCD vacuum vanishes and chiral symmetry is restored (or it can be slightly broken via the chiral breaking phenomena near the Fermi surface - the chiral density waves [9]). The important point is that the standard quark condensate of the vacuum vanishes at high density and the bulk mass in the confining mode has mostly the chirally symmetric origin.

2 Chiral symmetry breaking and its implications

The $SU(2)_L \times SU(2)_R$ axial symmetry of the QCD Lagrangian in the chiral limit is dynamically broken in the vacuum. Then there appear massless Goldstone bosons associated with broken axial symmetry.

Another direct evidence of dynamical chiral symmetry breaking in the vacuum is absence of the chiral parity doublets in the observed low-lying spectrum. If there is not a chiral partner to the nucleon, then its mass is generated through the quark condensate of the vacuum. Such a behavior can be modeled within the QCD sum rule approach, within the linear sigma-model, within the NJL model, variants of the bag model, constituent quark model with very massive constituent quarks, or within the Skyrme model. One cannot exclude, however, that some small part of nucleon mass is not related to the quark condensate. The (partial) axial vector current conservation translates this mass, via the Goldberger-Treiman relation, to the pion-nucleon coupling constant. Hence, the large pion-nucleon coupling constant encodes the physical origin of the nucleon mass as due to chiral symmetry breaking in the vacuum. It can be used as a natural measure for chiral symmetry breaking effect in a hadron.

3 Effective chiral restoration in baryon spectra

The nucleon excitation spectrum exhibits obvious patterns of parity doublets, see Fig. 1. Similar patterns are seen in the Delta spectrum. The linear axial transformation in the isospin space (i.e. the chiral $SU(2)_L \times SU(2)_R$ transformation) mixes the nucleon (or delta) states of a given parity with the nucleon or delta states of opposite parity. Then unbroken chiral symmetry requires existence of parity doublets in the nucleon and delta spectra that are not connected to each other, or quartets, i.e. parity doublets in the nucleon and delta spectra that are members of same higher representation. The absence of parity doublets in the low-lying spectrum is an evidence of chiral symmetry breaking in the vacuum. Appearance of parity doublets in highly excited nucleons and deltas was taken as evidence of effective chiral symmetry restoration [10, 11]. Of course, to claim a general pattern one needs a discovery of the still missing partners to the well established states with $7/2^-$ and $11/2^-$.

While these parity doublets are impressive, by themselves they are only suggestive, because they could be accidental. Given a statistical analysis of Ref. [12] the latter is unlikely. However, there could other reason, not related to chiral symmetry, responsible for the doubling. Then we need other observables that are sensitive specifically to chiral symmetry and that would correlate with the observed degeneracies. Such observables are axial properties of states: their axial charges and couplings to pions.
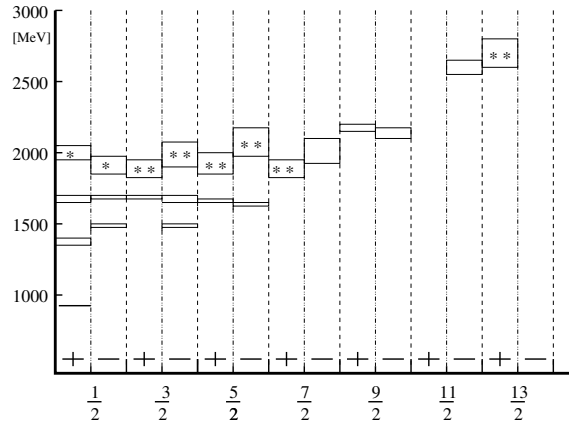

Fig. 1. Low- and high-lying nucleons. Those states which are not yet established are marked by ** or * according to the PDG classification.

Assume that the parity doublets are accidental and they are degenerate due to some other reason, not related to chiral symmetry. This means that these states are not chiral partners. Then their mass is induced by the quark condensate of the vacuum, like nucleon mass. The axial charges of these states should be expected to be of order 1. The Goldberger-Treiman relation then tells that these states must be very strongly coupled to pions and that the coupling constant to pion should be comparable with the pion-nucleon coupling constant. Such states should have a large decay coupling to the πN channel. In contrast, the effective chiral restoration requires that the axial charges of these states should be small (as compared to the nucleon axial charge) and they must have small decay coupling constants into the πN channel [3, 13]. The diagonal axial charges of excited states as well as their diagonal couplings to pions is difficult to extract from experiment. The decay coupling constants can be obtained from the known decay widths, however. These decay coupling constants in units of the well-known pion-nucleon coupling constant are shown in Table 1. One clearly observes that all those excited nucleons that are assumed from the spectroscopic patterns to be in approximate chiral multiplets have a very small decay coupling constant into the πN channel. In contrast, the only well established excited

nucleon, $N_{3/2^-}(1520)$, in which case a chiral partner cannot be identified from the spectrum, has a very large πN decay coupling constant. Consequently its mass origin, should be the quark condensate of the vacuum, similar to nucleon. It is a very interesting question what dynamics makes this state so peculiar. One observes a 100% correlation of the spectroscopic

patterns with the πN decays as predicted by effective chiral restoration. It is noteworthy that a specific small value of the decay coupling constant cannot be predicted by chiral symmetry and depends on the microscopic structure of the state. The approximate chiral symmetry only requires that these decay constants must be small.

Table 1. Chiral multiplets of excited nucleons. Comment: There are two possibilities to assign the chiral representation: $(1/2,0)\oplus(0,1/2)$ or $(1/2,1)\oplus(1,1/2)$ because there is a possible chiral pair in the Δ spectrum with the same spin with similar mass.

Spin	Chiral multiplet	Representation	$(f_{B_+N\pi}/f_{NN\pi})^2 - (f_{B_-N\pi}/f_{NN\pi})^2$
1/2	$N_+(1440) - N_-(1535)$	$(1/2,0)\oplus(0,1/2)$	$0.15 - 0.026$
1/2	$N_+(1710) - N_-(1650)$	$(1/2,0)\oplus(0,1/2)$	$0.0030 - 0.026$
3/2	$N_+(1720) - N_-(1700)$	$(1/2,0)\oplus(0,1/2)$	$0.023 - 0.13$
5/2	$N_+(1680) - N_-(1675)$	$(1/2,0)\oplus(0,1/2)$	$0.18 - 0.012$
7/2	$N_+(?) - N_-(2190)$	see comment	$? - 0.00053$
9/2	$N_+(2220) - N_-(2250)$	see comment	$0.000022 - 0.0000020$
11/2	$N_+(?) - N_-(2600)$	see comment	$? - 0.000000064$
3/2	$N_-(1520)$	no chiral partner	2.5

The diagonal axial charges of excited states cannot be measured experimentally, but in principle can be obtained in lattice simulations. On the lattice it is an intrinsically difficult problem to extract the highly excited states. Nevertheless a progress has been achieved in identification of the lowest negative parity excitations [14]. Recently first lattice results for the diagonal axial charges of the states $N_{1/2^-}(1535)$ and $N_{1/2^-}(1650)$ have appeared [15]. These first results are limited to rather large quark masses and also require confirmation by other groups. Assuming a naive extrapolation of the results to the physical point one concludes that the axial charge of the lowest negative parity excitation, $N_{1/2^-}(1535)$, is very small. This is consistent with a possible identification of the $N_{1/2^+}(1440)-N_{1/2^-}(1535)$ pair as the lowest chiral pair. The chiral symmetry breaking effects are still large in this case (because of rather large splitting of the states). May be this also explains a long-standing puzzle why the Roper state has so small mass. Given large error bars for the next excited state, $N_{1/2^-}(1650)$, it is difficult to extrapolate its axial charge to the physical point. One should mention that the obtained values of the axial charges are also consistent with the $SU(6)_{FS}\times O(3)$ quark model prediction assuming that there is not mixing of the $S = 1/2$ state $N_{1/2^-}(1535)$ with the $S = 3/2$ state $N_{1/2^-}(1650)$ via the tensor quark-quark force [16–18]. Within the Isgur-Karl type quark models such a mixing is very strong; it makes the axial charge of $N_{1/2^-}(1535)$ to be of the order 1. Such a strong mix-

ing is also very important to obtain a reasonable fit of strong baryon decays within the constituent quark model.

Here one more comment on the quark (or large N_c) description of baryon decays is relevant. These models operate with the nonrelativistic decay amplitudes with improper nonrelativistic phase space factors. They try to fit, with some free parameters, decay widths, rather than the coupling constants. Physics is contained in the coupling constants, however. Then a proper procedure would be to compare the quark model predictions with the coupling constants.

4   Highly excited mesons

Figure 2 shows the spectra of the well established mesons from the PDG and new, not yet confirmed $\bar{n}n$ states from the partial wave analysis [19, 20] of $\bar{p}p$ annihilation at LEAR (CERN). Obvious large degeneracy of the high-lying mesons is seen. How could we understand such a degeneracy? This data have been analyzed in Refs. [21, 22] and it was shown that the degeneracies of the high-lying states with $J = 0-3$ are consistent with the conjecture of effective $SU(2)_L\times SU(2)_R$ and $U(1)_A$ restorations. A prediction was made that the pattern for the high-lying $J = 4$ states should be similar to the pattern of $J = 2$ states and the pattern for the $J = 5$ mesons should be the same as the pattern for the $J = 3$ mesons. There are a_4 and f_4 positive parity states in the band
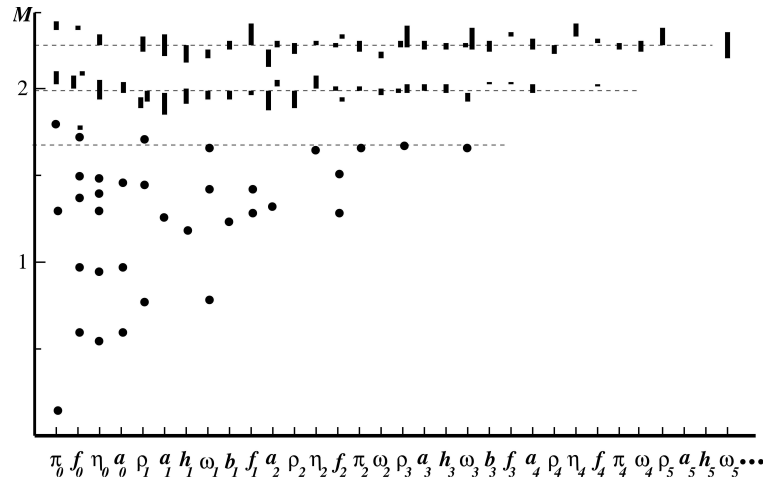
Fig. 2.   Masses (in GeV) of the well established states from PDG (circles) and new $\bar{n}n$ states from the proton-antiproton annihilation (strips). Note that the well-established states include $f_0(1500)$, $f_0(1710)$, which are the glueball and $\bar{s}s$ states with some mixing and hence are irrelevant from the chiral symmetry point of view. Similar, the $f_0(980)$, $a_0(980)$ mesons most probably are not $\bar{n}n$ states and also should be excluded from the consideration. The same is true for $\eta(1475)$, which is the $\bar{s}s$ state and $\eta(1405)$ with the unknown nature.

around 2 GeV and their possible partners of opposite parity are missing. Similar happens with the $J=5$ states in the 2.3 GeV band. The absence of the chiral partners for these highest spin mesons would be a difficulty for the chiral restoration scenario. Consequently a key question is whether these states do not exist or they cannot be seen in the $\bar{p}p$ annihilation, even if they exist. It turns out that the latter is correct [23].

Consider, for example, the missing $J=4$ states of negative parity in the 2 GeV band. They all require the $L=4$ partial wave in the $\bar{p}p$ system. In contrast, the observed a_4 and f_4 mesons are produced in the $L=3$ partial wave. From Fig. 2 it follows that the missing chiral partners should be expected with mass 2000±50 MeV. At such energy the $L=4$ partial wave in the $\bar{p}p$ system is suppressed as compared to the $L=3$ partial wave by the factor 10^1–10^3 by the centrifugal repulsion in the $\bar{p}p$ system. A signal from this missing states is very weak as compared to the observed a_4 and f_4 states. The same suppression factor is valid with respect to the negative parity states seen in the 2.3 GeV band. With such a weak signal the χ^2 fit cannot reveal these missing states, even if they exist. Similar analysis can be done for the $J=5$ states in the 2.3 GeV band. One then concludes that the existing experimental data [19, 20] on highly excited mesons cannot answer a question about existence or non-existence of these missing states and new types of experiments with polarization are required to answer this conceptually important question.

The chiral and $U(1)_A$ symmetries cannot explain degeneracies of the states with different spins. Such a

degeneracy can be obtained, however, if one assumes a principal quantum number $n+J$ on top of chiral restoration [24].

There exists an alternative conjecture about nature of the large degeneracy. If these high-lying states behaved nonrelativistically (i.e., the valence quarks were nonrelativistic) the degeneracy could be obtained assuming the standard nonrelativistic LS coupling scheme and a principal quantum number $n+L$, where L is the orbital angular momentum of quarks [25–27]. In the nonrelativistic quantum mechanics L can indeed be a good quantum number in absence of the spin-orbit force (compare this, e.g., with the nonrelativistic Hydrogen atom).

Such a scenario is inconsistent with two basic facts. In QCD, which is a highly relativistic quantum field theory, there is only one conserved angular momentum, J. The principal quantum number $n+L$ would imply that there are three independent conserved angular momenta, L, S, J! Such an assumption is also inconsistent with the stringy picture, on which it is based. The ends of the rotating string move at the speed of light. Then the quarks at the ends of the string are ultrarelativistic and must have a definite chirality because only chiral quarks can move at the speed of light. The stringy picture does imply the *unbroken* chiral symmetry and would in fact predict the missing states in the 2 and 2.3 GeV bands. A consistent relativistic string model with quarks at the ends of the string is an open issue.

Another interesting question is a linearity of the Regge trajectories. The leading nucleon angular Regge trajectory is approximately linear, a well

known fact. The daughter trajectory is highly non-linear, however. While the high-spin states $J = 5/2$ and $J = 9/2$ of both positive and negative parity are approximately degenerate, there is not a degenerate state of the opposite parity for the nucleon. In the meson spectrum the leading angular Regge trajectory is also approximately linear. Should one expect a linearity of all daughter Regge trajectories in the meson spectrum?

5　Models

One cannot solve QCD, even in the large N_c limit. Hence at the moment the only useful tool to address the problem of highly excited hadrons is modeling. The model must contain all principal elements of QCD that are relevant to the present problem. It must be (i) relativistic and field-theoretical in nature, (ii) chirally symmetric, (iii) confining, (iv) it should provide dynamical breaking of chiral symmetry. It is highly nontrivial to meet all these requirements within one and the same model. For example the NJL model (or variants of thereof) is chirally symmetric and guarantees spontaneous breaking of chiral symmetry. It is not confining, however. In contrast, the potential constituent quark models do not respect points (i), (ii) and (iv).

There exists such a model, however [28, 29]. The model is a generalization of the 't Hooft model [30]. The 't Hooft model is QCD in large N_c limit in 1+1 dimensions. Due to its low-dimensional nature it is an *exactly* solvable field theory. One can analytically calculate all required quantities: the quark condensate, the meson spectrum, etc. It is useful to understand *how* this field theory is solved. In 1+1 dimensions the highly nonlinear gluodynamics is exactly reduced to the Coulomb interaction alone, which is an instantaneous linear potential of the Lorentz-vector type. To address the problem of dynamical chiral symmetry breaking one has to solve the gap (Schwinger-Dyson) equation. Given the quark Green function obtained from the gap equation, it is possible to solve the Bethe-Salpeter equation for mesons, etc. However, in 1+1 dimensions a rotational motion and angular momenta are absent, that are crucial for effective chiral restoration. The model [28, 29] is a straightforward generalization of the 't Hooft model to 3+1 dimensions. It is postulated within the model that there exists a linear instantaneous Coulomb-type potential in 3+1 dimensions. All other possible gluonic interactions are neglected. Given its "simplicity", the model can be solved numerically. The problem of chiral

restoration in excited mesons has been addressed in Ref. [31]. A complete spectrum of mesons has been calculated and a fast effective chiral restoration with increasing spin J has been demonstrated. It is instructive to outline a physical mechanism responsible for the phenomenon. When one increases the spin of a hadron J, one also increases a typical momentum of valence quarks. The chiral symmetry breaking dynamical mass of quarks is important only at low momenta. At large J all low momenta components are suppressed in the meson wave function by the centrifugal repulsion and consequently the chiral symmetry breaking dynamical mass of quarks gets irrelevant. Consequently one observes the effective chiral restoration. Due to its simplicity, the model cannot reproduce a degeneracy of chiral multiplets with different spins [32]. It is much more difficult to solve the model for baryons. Nevertheless some steps have been done [33, 34] and similar effective chiral restoration with increasing J has been observed.

The problem of highly excited hadrons has been addressed for the last years within many different holographic models. All existing holographic models of hadrons suffer an essential disease, however. The holographic models are assumed to satisfy the AdS/CFT matching conditions at the ultraviolet border, where chiral symmetry is not broken. It has been recently proven [35] that in reality they do not satisfy these matching conditions. As such, they are not suitable to address the problems related to interconnections of confinement and chiral symmetry in hadrons.

6　Chiral restoration in the quarkyonic matter

Quite recently McLerran and Pisarski suggested a new state of the matter - the quarkyonic matter [5]. Their crucial observation was that in the large N_c limit at low and moderate temperatures, confinement persists up to arbitrary high densities. There are no dynamical quark loops and hence nothing screens the confining gluon propagator.

At some critical density one should expect a chiral phase transition, namely one expects that the standard quark-antiquark condensate of the QCD vacuum should vanish. Then one arrives at a subphase within the quarkyonic matter that is manifestly confining and at the same time the quarks condensate vanishes. How could it be ?! This question was addressed in Refs. [6–8]. The same confining and chirally symmetric model was chosen that had been

used for study of effective chiral restoration in excited hadrons. The following mechanism for confining matter with vanishing quark-antiquark condensate was observed. The quark Green function, that is a solution of the gap equation, acquires not only the chiral symmetry breaking Lorentz-scalar part but also the chirally symmetric Lorentz-vector part. Both these parts are infrared-divergent, which guarantees that there are no single quarks in the spectrum. The infrared divergence cancels exactly in all color-singlet quantities, such us the quark condensate or hadronic excitations. At some critical density the Lorentz-scalar part of the quark self-energy as well as the quark condensate of the vacuum vanishes and one observes a chiral phase transition. This happens exclusively due to the Pauli blocking of the quark levels that are required for existence of the quark condensate. The Lorentz-vector part of the quark self-energy does not vanish, however, and is still infrared-divergent. This guarantees that a single quark does not exist even above the phase transition. This is just

in contrast to the nonconfining NJL model (or variants of thereof) where the Lorentz-vector part of the quark self-energy is absent and one inevitably has a massless free quark in the chirally symmetric regime. The infrared singularity cancels exactly in any possible color-singlet excitation of the matter and such an excitation has a finite well-defined energy. Consequently even in the absence of the chiral symmetry breaking via the quark-antiquark condensate one obtains a confining matter with color-singlet excitations only. The mass of a dense confining matter is not related to the chiral symmetry breaking via the quark condensate.

It does not mean, however, that such a matter with vanishing quark condensate will be exactly chirally symmetric. It can have a small amount of chiral symmetry breaking via the chiral density waves near the Fermi surface [9]. But in any case the bulk of the mass of this confining matter is not related to chiral symmetry breaking.

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