

# QCD SYMMETRIES IN EXCITED HADRONS

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## Abstract

Recent developments for chiral and  $U(1)_A$  restorations in excited hadrons are reviewed. We emphasize predictions of the chiral symmetry restoration scenario for axial charges and couplings to Goldstone bosons. Using very general chiral symmetry arguments it is shown that strict chiral restoration in a given excited nucleon forbids its decay into the  $N\pi$  channel. We confront this prediction with the  $N^*N\pi$  coupling constants extracted from the decay widths and observe a 100 % correlation of these data with the spectroscopic parity doublet patterns. These results suggest that the lowest approximate chiral parity doublet is the  $N(1440) - N(1535)$  pair. In the meson sector we discuss predictions of the chiral symmetry restoration for still missing states and a signature of the higher symmetry observed in new  $\bar{p}p$  data. We conclude with the exactly solvable chirally symmetric and confining model that can be considered as a generalization of the 1+1 dimensional 't Hooft model to 4 dimensions. Complete spectra of  $\bar{q}q$  mesons demonstrate a fast chiral restoration with increasing  $J$  and a slow one with increasing  $n$ .

## 1 Introduction

In hadrons consisting of  $u$  and  $d$  quarks there are two crucially important properties of QCD - chiral symmetry and confinement. Their interrelations and mechanisms are not yet understood. What we do know theoretically is that at zero temperature and density in the confining phase chiral symmetry must be necessarily spontaneously broken in the vacuum [1]. Another conceptual and closely related issue is the generation of hadron mass in the light quark sector. It was considered almost self-evident that such a mass generation proceeds via the chiral symmetry breaking in the vacuum and the most important characteristics that determines the hadron mass is the quark

condensate of the vacuum. Indeed, it is firmly established both phenomenologically and on the lattice that to leading order the pion mass squared is proportional to the bare quark mass and the quark condensate [2]. In the baryon sector the very absence of the chiral partner to the nucleon implies that its mass is at least mostly related to the spontaneous breaking of chiral symmetry in the vacuum. This fact is supported by the Ioffe formula [3] that connects, though not rigorously, the nucleon mass with the quark condensate. Another obvious sign of the strong dynamical chiral symmetry breaking effects in the nucleon is the large pion-nucleon coupling constant. Indeed, it is well understood that the coupling of the Goldstone bosons to the nucleon is a direct consequence of the spontaneous chiral symmetry breaking and is a basis for nucleon chiral perturbation theory [4]. One more strong evidence for the chiral symmetry breaking in the nucleon is its large axial charge,  $g_A = 1.26$ .

A main message of this talk is that the mass generation mechanism in excited hadrons is essentially different - the quark condensate of the vacuum becomes less and less important with the excitation and the chiral as well as the  $U(1)_A$  symmetries get eventually approximately restored in the given hadron, even though they are strongly broken in the vacuum. This is referred to as effective restoration of chiral symmetry, for a review see ref. [10].

It is important to precisely characterize what is implied under effective restoration of chiral and  $U(1)_A$  symmetry in excited hadrons. A mode of symmetry is defined only by the properties of the vacuum. If a symmetry is spontaneously broken in the vacuum, then it is the Nambu-Goldstone mode and the whole spectrum of excitations on the top of the vacuum is in the Nambu-Goldstone mode. However, it may happen that the role of the chiral symmetry breaking condensates becomes progressively less important higher in the spectrum, because the valence quarks decouple from the quark condensates. This means that the chiral symmetry breaking effects become less and less important in the highly excited states and asymptotically the states approach the regime where their properties are determined by the underlying unbroken chiral symmetry (i.e. by the symmetry in the Wigner-Weyl mode). This effective restoration in excited hadrons should not be confused with the chiral symmetry restoration in the vacuum at high temperature/density. In the latter case the quark vacuum becomes trivial and the system is in the Wigner-Weyl mode. In the former case the symmetry is always broken in the vacuum, however this symmetry breaking in the vacuum gets irrelevant in the highly excited states.

## 2 Empirical evidence for chiral restoration in excited nucleons

The nucleon excitation spectrum is shown in Fig. 1. Only well-established states (i.e. without stars in boxes) should be seriously considered. It is well seen that there is no chiral partner to the nucleon. This necessarily implies that chiral symmetry is strongly broken in the nucleon and consequently is realized nonlinearly [5]. Obvious approximate parity doublets are observed in the region 1.7 GeV and higher. An absence of parity doublets for the lowest-lying states and their apparent appearance for (highly) excited states was taken in refs. [6–9] as evidence for chiral restoration in excited baryons, for a review see ref. [10]. The parity doublets in the 1.7 GeV region have been assigned to the  $(0, 1/2) + (1/2, 0)$  representation of the parity-chiral group because there are no approximately degenerate doublets in the same mass region in the spectrum of the delta-resonance [8, 10]. A clear testable prediction of the chiral symmetry restoration scenario is an existence of chiral partners of the well established high-lying resonances  $N(2190)$  and  $N(2600)$ . A dedicated experimental search of these missing states can be undertaken [11]. Similar situation takes place in the Delta-spectrum.

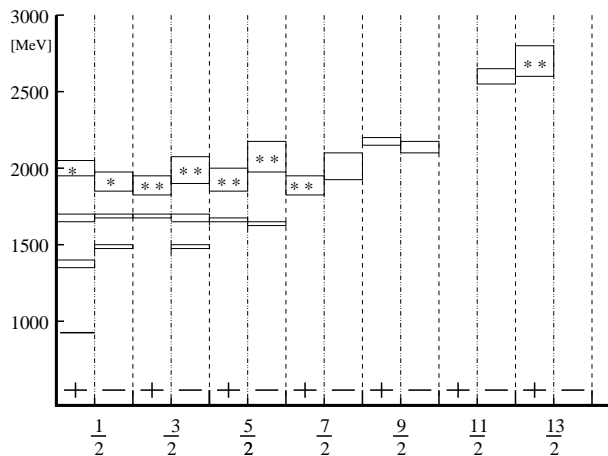


Figure 1: Low- and high-lying nucleons. Those states which are not yet established are marked by \*\* or \* according to the PDG classification.

While these parity doubling patterns are impressive, they are still only suggestive, because so far no other complementary experimental data would independently tell us that these parity doublets are due to effective chiral symmetry restoration. Strict chiral restoration in a given baryon would imply that its diagonal axial charge is zero and hence the diagonal coupling

constant to the pion must vanish [10, 12–15]. This is one of the most important implications of the chiral symmetry restoration and is reviewed below.

Assume that we have a free  $I = 1/2$  chiral doublet  $B$  in the  $(0, 1/2) + (1/2, 0)$  representation and there are no chiral symmetry breaking terms. This doublet is a column [16]

$$B = \begin{pmatrix} B_+ \\ B_- \end{pmatrix}, \quad (1)$$

where the bispinors  $B_+$  and  $B_-$  have positive and negative parity, respectively. The chiral transformation law under the  $(0, 1/2) \oplus (1/2, 0)$  representation provides a mixing of two fields  $B_+$  and  $B_-$ <sup>1</sup>

$$B \rightarrow \exp\left(i\frac{\theta_V^a \tau^a}{2}\right) B; \quad B \rightarrow \exp\left(i\frac{\theta_A^a \tau^a}{2}\sigma_1\right) B. \quad (2)$$

Here  $\sigma_i$  is a Pauli matrix that acts in the  $2 \times 2$  space of the parity doublet. Then the chiral-invariant Lagrangian of the free parity doublet is given as

$$\mathcal{L}_0 = i\bar{B}\gamma^\mu\partial_\mu B - m_0\bar{B}B = i\bar{B}_+\gamma^\mu\partial_\mu B_+ + i\bar{B}_-\gamma^\mu\partial_\mu B_- - m_0\bar{B}_+B_+ - m_0\bar{B}_-B_-.$$

Alternative forms for this Lagrangian can be found in refs. [17, 18].

A crucial property of this Lagrangian is that the fermions  $B_+$  and  $B_-$  are exactly degenerate and have a nonzero chiral-invariant mass  $m_0$ . In contrast, for usual (Dirac) fermions chiral symmetry in the Wigner-Weyl mode restricts particles to be massless, hence they acquire their mass only in the Nambu-Goldstone mode of chiral symmetry due to chiral symmetry breaking in the vacuum (i.e. via the coupling with the quark condensate of the vacuum). The chiral parity doublets have their chiral-invariant mass term already in the Wigner-Weyl mode and this mass term has no relation with the quark condensate.

From the axial transformation law (2) one can read off the axial charge matrix, which is  $\gamma_5\sigma_1$ . Hence the diagonal axial charges of the opposite parity baryons are exactly 0,  $g_+^A = g_-^A = 0$ , while the off-diagonal axial charge is 1,  $|g_{+-}^A| = |g_{-+}^A| = 1$ . This is another crucial property that distinguishes the parity doublets from the Dirac fermions where  $g^A = 1$ . The axial vector current conservation,  $q^\mu\langle B_\pm|A_\mu|B_\pm\rangle = 0$ , translates this axial charge matrix via the Goldberger-Treiman relation into the  $\pi B_\pm B_\pm$  coupling constants which are zero. Hence a small (vanishing) value of the pion-baryon coupling constant taken together with the large baryon mass would tell us that the

<sup>1</sup>Note that the axial transformation given in [16] is incorrect as it breaks chiral symmetry of the kinetic term. The correct axial transformation is given in ref. [10].

origin of this mass is not due to chiral symmetry breaking in the vacuum. An experimental verification of the smallness of the diagonal axial charges or smallness of the pion-baryon coupling constants would be a direct verification of the chiral symmetry restoration scenario in excited nucleons. It is unclear, however, how to measure experimentally these quantities.

There is rich experimental data on strong decays of excited hadrons. It turns out that the chiral restoration implies a very strong selection rule [19]. Namely, it predicts that if chiral symmetry is completely restored in a given excited nucleon ( $B$ ), then it cannot decay into the  $\pi N$  channel, i.e. the coupling constant  $f_{BN\pi}$  must vanish. This selection rule is based exclusively on general properties of chiral symmetry and hence is model-independent.

Let us prove this selection rule. Assume that a  $\pi N$  decay of an exact chiral doublet is possible. Then there must be a self-energy contribution  $B_{\pm} \rightarrow \pi N \rightarrow B_{\pm}$  into its mass. Then the axial rotation (2) would require that the S-wave  $\pi N$  state transforms into the P-wave  $\pi N$  state. However, in the Nambu-Goldstone mode the axial rotations of the pion and nucleon states are fixed - these are the nonlinear axial transformations [14,15]. Given these well known axial transformation properties of the Goldstone boson and nucleon [5] it is not possible to rotate the S-wave  $\pi N$  state into the P-wave  $\pi N$  state. Therefore, there cannot be any  $\pi N$  self-energy component in  $B_{\pm}$ . Hence a decay  $B_{\pm} \rightarrow \pi N$  is forbidden. However, a decay of the exact chiral doublet into e.g.  $N\rho$  or  $N\pi\pi$  is not forbidden. Hence, if a state is a member of an approximate chiral multiplet, then its decay into  $N\pi$  must be strongly suppressed,  $(f_{BN\pi}/f_{NN\pi})^2 \ll 1$ .

If, in contrast, the excited baryon has no chiral partner, then its mass, like in the nucleon case is exclusively due to chiral symmetry breaking in the vacuum. Its axial charge should be comparable with the nucleon axial charge. Then nothing forbids its strong decay into  $N\pi$ . One then expects that the decay coupling constant should be of the same order as the pion-nucleon coupling constant. These two extreme cases suggest that a magnitude of the  $BN\pi$  decay constant can be used as an indicator of the mass origin.

The decay constants  $f_{BN\pi}$  can be extracted from the  $B \rightarrow N + \pi$  decay widths, see e.g. [20,21]. The pion-nucleon coupling constant is well-known,  $f_{NN\pi} = 1.0$ . In Table 1 we show ratios  $(f_{BN\pi}/f_{NN\pi})^2$  for all well-established states. It is well seen that this ratio is  $\sim 0.1$  or smaller for approximate  $J = 1/2, 3/2, 5/2$  parity doublets. For the high-spin states this ratio is practically vanishing. This is consistent with the recent demonstration of the large J-rate of chiral restoration within the only known exactly solvable confining and chirally-symmetric model [22].

From Fig. 1 one can see that the only well established excited state which has no obvious chiral partner is  $3/2^-, N(1520)$ . It decays very strongly into

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  $\Delta$  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

$N\pi$ , indeed. This implies that a nature of mass of this state is rather different compared to approximate parity doublets. One observes a 100% correlation of the spectroscopic patterns with the  $N\pi$  decays, as predicted by the chiral symmetry restoration.

The Fig. 1 and the Table 1 suggest that the lowest approximate chiral doublet is  $N(1440) - N(1535)$ . If correct, the diagonal axial charges of these states must be small. While it is impossible to measure these charges experimentally, this can be done on the lattice. The axial charge of  $N(1535)$  has just been measured by Takahashi and Kunihiro and they report it to be surprisingly small, smaller than 0.2 [23]. Certainly lattice studies of other states are welcome.

### 3 Symmetries in excited mesons

Fig. 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 [24, 25] of  $\bar{p}p$  annihilation at LEAR (CERN). Obvious high symmetry of the high-lying  $\bar{n}n$  states is seen. These data have been analysed in ref. [26] and it turned out that the high-lying  $\bar{n}n$  mesons perfectly fit all possible linear chiral multiplets of both  $SU(2)_L \times SU(2)_R$  and  $U(1)_A$  groups with a few still missing states. In particular, the chiral symmetry predicts a duplication of some of the  $J > 0$  states with the given quantum numbers, which is indeed observed in data. If the chiral symmetry is indeed responsible for positive-negative

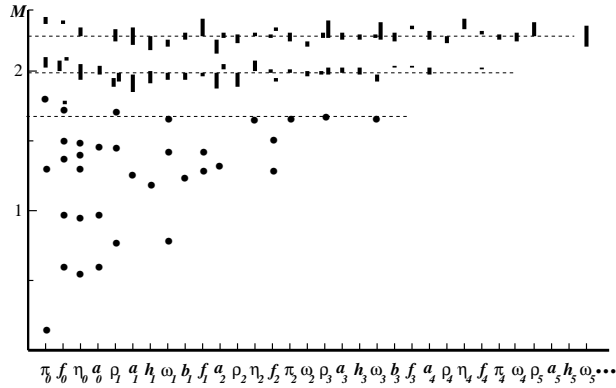


Figure 2: Masses (in GeV) of the well established states from PDG (circles) and new  $\bar{n}n$  states from the proton-antiproton annihilation (strips). Note that the well-established states include  $f_0(1500)$ ,  $f_0(1710)$ , which are the glueball and  $\bar{s}s$  states with some mixing and hence 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.

parity degeneracy of the states, then there should be chiral multiplets for the high-spin states at the levels  $M \sim 2$  GeV,  $M \sim 2.3$  GeV and, possibly, at  $M \sim 1.7$  GeV. These states are presently missing in refs. [24,25] and it would be extraordinary important to find them or to reliably exclude them. Note that such high-spin parity doublets are well seen in the nucleon spectrum - see Fig. 1.

The chiral and  $U(1)_A$  symmetries can connect only states with the same spin. Certainly we observe larger degeneracy, the states with different spins are also degenerate. The large degeneracy might be understood if, on top of chiral and  $U(1)_A$  restorations, a principal quantum number  $N = n + J$  existed.

There are suggestions in the literature to explain this large degeneracy without resorting to chiral symmetry, assuming the  $\vec{J} = \vec{L} + \vec{S}$  coupling scheme and that there is a principal quantum number  $N = n + L$ , where  $L$  is the *conserved* orbital angular momentum in the quark-antiquark system [27–29]. This suggestion is hard to reconcile with the Lorentz and chiral symmetries, however [30].

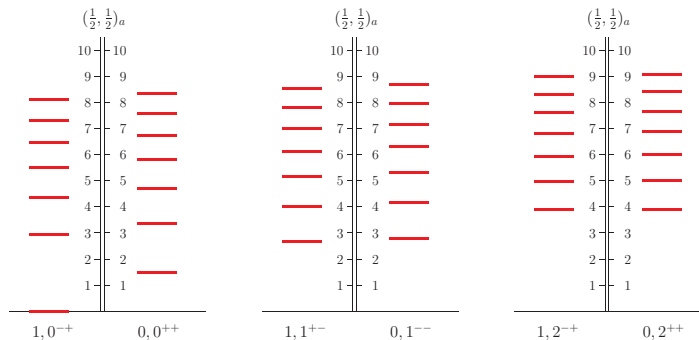


Figure 3: Spectra of the  $\bar{q}q$  mesons in the  $(1/2, 1/2)_a$  representations.

## 4 Chirally symmetric and confining solvable model

There exists only one known manifestly chirally-symmetric and confining model in four dimensions that is solvable [31], sometimes called Generalized Nambu and Jona-Lasinio model (GNJL). This model can be considered as a generalization of the 1+1 dimensional 't Hooft model, that is QCD in the large  $N_c$  limit [32]. It is postulated within the GNJL model that there exists a linear confining potential of the Coulomb type in four dimensions. The chiral symmetry breaking and the properties of the Goldstone bosons have been obtained from the solution of the Schwinger-Dyson and Bethe-Salpeter equations [33–38]. The complete spectrum of  $\bar{q}q$  mesons has been calculated only recently, in ref. [22], which exhibits restoration of the chiral symmetry.

Part of the spectra is shown in Fig. 3 and a fast chiral restoration with increasing of  $J$  is observed, while a slow rate is seen with respect to the radial quantum number  $n$ . It is possible to see directly a mechanism of the chiral restoration. The chiral symmetry breaking Lorentz-scalar dynamical mass of quarks  $M(q)$  arises via selfinteraction loops and vanishes fast at large momenta. When one increases the spin of the hadron  $J$ , or its radial quantum number  $n$ , one also increases the typical momentum of valence quarks. Consequently, the chiral symmetry violating dynamical mass of quarks becomes small and chiral symmetry gets approximately restored. This mechanism of chiral restoration is in accord with a general semiclassical analysis [9, 10, 39].

A higher degeneracy is recovered for  $J \rightarrow \infty$ . In this limit all states with the same  $J$  and  $n$  fall into reducible representation  $[(0, 1/2) \oplus (1/2, 0)] \times [(0, 1/2) \oplus (1/2, 0)]$ , hence the quantum loop effects become irrelevant and all

possible states with different quark chiralities become equivalent.

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## References

- [1] G. 't Hooft, in: Recent developments in Gauge Theories, Eds. G. 't Hooft et al, Plenum, New York, 1980; S. Coleman, E. Witten, Phys. Rev. Lett. **45**, 100 (1980).
- [2] M. Gell-Mann, R. J. Oakes, B. Renner, Phys. Rev. **175**, 2195 (1968).
- [3] B. L. Ioffe, Nucl. Phys. **B 188**, 317 (1981); E: **B 191**, 591 (1981).
- [4] U-G. Meissner, Chiral QCD: Baryon Dynamics, in: At the frontiers of Particle Physics. Handbook of QCD., edited by M. Shifman, v. 1, p. 417, World Sc. 2001.
- [5] S. Weinberg, Phys. Rev. **166**, 1568 (1968).
- [6] L. Ya. Glozman, Phys. Lett. B **475**, 329 (2000).
- [7] T. D. Cohen and L. Ya. Glozman, Phys. Rev. D **65**, 016006 (2002).
- [8] T. D. Cohen and L. Ya. Glozman, Int. J. Mod. Phys. A **17**, 1327 (2002).
- [9] L. Ya. Glozman, Int. J. Mod. Phys. A **21**, 475 (2006).
- [10] L. Ya. Glozman, Phys. Rep. **444**, 1 (2007).
- [11] A. Sibirtsev et al, arXiv: 0706.0183 [nucl-th].
- [12] T. D. Cohen and L. Ya. Glozman, Mod. Phys. Lett. **A 21**, 1939 (2006).
- [13] L. Ya. Glozman, A. V. Nefediev, Phys. Rev. **D 73**, 074018 (2006).
- [14] R. L. Jaffe, D. Pirjol, A. Scardicchio, Phys. Rev. **D 74**, 057901 (2006); Phys. Rev. Lett. **96**, 121601 (2006).
- [15] R. L. Jaffe, D. Pirjol, A. Scardicchio, Phys. Rep. **435**, 157 (2006).
- [16] B. W. Lee, Chiral Dynamics, Gordon and Breach, New York, 1972
- [17] C. DeTar, T. Kunihiro, Phys. Rev. **D 39**, 2805 (1989).

- [18] D. Jido, M. Oka, A. Hosaka, *Progr. Theor. Phys.* **106**, 873 (2001).
- [19] L. Ya. Glozman, arXiv:0706.3288 [hep-ph], *Phys. Rev. Lett.*, in print.
- [20] W. K. Cheng, C. W. Kim, *Phys. Rev.* **154**, 1525 (1967).
- [21] D. O. Riska and G. E. Brown, *Nucl. Phys. A* **679**, 577 (2001).
- [22] R. F. Wagenbrunn, L. Ya. Glozman, *Phys. Lett. B* **643**, 98 (2006); *Phys. Rev. D* **75**, 036007 (2007).
- [23] T. Takahashi and T. Kunihiro, talk at MENU2007.
- [24] A. V. Anisovich et al, *Phys. Lett. B* **491**, 47 (2000); *ibid.* B **517**, 261 (2001); *ibid.* B **542**, 8 (2002); *ibid.* B **542**, 19 (2002); *ibid.* B **513**, 281 (2001).
- [25] D. V. Bugg, *Phys. Rep.* **397**, 257 (2004).
- [26] L. Ya. Glozman, *Phys. Lett. B* **539**, 257 (2002); *ibid* **587**, 69 (2004).
- [27] S. S. Afonin, *Phys. Rev. C* **76**, 015202 (2007).
- [28] E. Klempt and A. Zaitsev, arXiv:0708.4016 [hep-ph].
- [29] M. Shifman and A. Vainshtein, arXiv:0710.0863 [hep-ph].
- [30] L. Ya. Glozman and A. V. Nefediev, arXiv:0704.2673 [hep-ph], to appear in *Phys. Rev. D*.
- [31] A. Le Yaouanc et al, *Phys. Rev. D* **29** (1984) 1233; **31** (1985) 137.
- [32] G. 't Hooft, *Nucl. Phys. B* **75** (1974) 461; I. Bars, M. B. Green, *Phys. Rev. D* **17** (1978) 537.
- [33] S. L. Adler and A. C. Davis, *Nucl. Phys. B* **244** (1984) 469.
- [34] R. Alkofer and P. A. Amundsen, *Nucl. Phys. B* **306** (1988) 305.
- [35] P. Bicudo and J. E. Ribeiro, *Phys. Rev. D* **42** (1990) 1611; 1625.
- [36] P. J. A. Bicudo and A. V. Nefediev, *Phys. Rev. D* **68** (2003) 065021.
- [37] F. J. Llanes-Estrada and S. R. Cotanch, *Phys. Rev. Lett.*, **84** (2000) 1102.
- [38] R. Alkofer et al, *Phys. Rev. Lett.*, **96**, 022001 (2006).
- [39] L. Ya. Glozman, A. V. Nefediev, J. Ribeiro, *Phys. Rev. D* **72** (2005) 094002.