Chiral symmetry and the string description of excited hadrons

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A large symmetry group is perhaps experimentally observed in excited hadrons which includes the chiral group $U(2)_L \times U(2)_R$ as a subgroup. To possess this large symmetry a dynamical model for excited hadrons, presumably a string model, should explain formation of chiral multiplets and, at the same time, predict coinciding slopes of the angular and radial Regge trajectories. This is possible only if both the dynamics of the string and the chirality of the quarks at the ends of the string are considered together. We construct a model-independent unitary transformation from the relativistic chiral basis to the $\{2S+1L_J\}$ basis, commonly used in hadronic phenomenology as well as in the string models, and demonstrate that a hadron belonging to the given chiral representation is a fixed superposition of the basis vectors with different $L$'s. Thus the description of highly excited hadron in terms of a fixed $L$ is not compatible with chiral symmetry and has to be disregarded in favour of the description in terms of the total hadron spin $J$. Therefore, dynamics of the string must deliver the principal quantum number $\sim n + J$, in order chiral multiplets with different spins to become degenerate, as required by the large symmetry group.

PACS numbers: 11.30.Rd, 12.38.Aw, 14.40.-n

Recent remarkable experimental results on highly excited mesons $^1$ from the $p \bar{p}$ annihilation at LEAR (CERN) reveal different kinds of symmetries of highly excited $n\bar{n}$ mesons$^1$: (i) equality of slopes of the angular and radial Regge trajectories $^2$, (ii) effective restoration of the chiral $SU(2)_L \times SU(2)_R$ and $U(1)_A$ symmetries $^3$, $^4$ (see Ref. $^6$ for a review), (iii) clustering of many states with different spins, parities, and isospins around definite energies $^5$, $^6$, $^7$, $^8$ (the latter property is obvious from Fig. 2 of Ref. $^6$). This clustering implies that perhaps a large symmetry group is observed which includes chiral symmetries as subgroups$^2$. Similar symmetries are also seen in excited baryon spectra $^6$, $^9$, $^{10}$, $^{11}$. These results suggest that physics, and in particular mass generation mechanisms, are essentially different in the low- and high-lying states. The fundamental underlying reason for this difference comes from the fact that quantum loop effects, which govern the properties of the low-lying states, are suppressed in high-lying hadrons, so that the dynamics becomes semi-classical $^6$, $^{10}$, $^{11}$. Certainly understanding this large symmetry group is one of the key problems to approach QCD in the infrared, in particular in what concerns confinement and its interrelation with chiral symmetry breaking.

Both the chiral restoration and asymptotically linear angular and radial Regge trajectories are reproduced $^{12}$ within the only known exactly solvable chirally symmetric and confining model (Generalised Nambu–Jona-Lasinio (GNJL) model) $^{13}$, $^{14}$. In this model the only gluonic interaction is the linear Coulomb-like confining potential. Then chiral symmetry breaking is obtained from the Schwinger–Dyson (mass–gap) equation, and the meson spectrum results from a Bethe–Salpeter equation $^{15}$, $^{14}$, $^{16}$, $^{17}$. While chiral symmetry restoration in excited hadrons is naturally explained in the framework of this model $^{17}$ and all possible highly excited states with the same $J$ and $n$ fall into $[(0,1/2) + (1/2,0)] \times [(0,1/2) + (1/2,0)]$ chiral representation (which combines all possible chiral multiplets with the same $J$) and become approximately degenerate $^{12}$, the degeneracy of states with different $J$’s is absent, because the slopes of the radial and angular Regge trajectories are different. It is not possible to obtain simultaneously chiral symmetry restoration and equal slopes of angular and radial Regge trajectories within any equal–time relativistic potential description $^{18}$.

Equal slopes of the radial and angular Regge trajectories is a generic property of the Veneziano dual amplitude $^{19}$ and of the open bosonic Nambu–Goto string $^{20}$. Quarks and the notion of chiral symmetry are absent in this pre–QCD approach, however.

There are modern attempts to model QCD within the AdS/QCD approach $^{21}$, based on the ideas of Maldacena on duality between the conformal supersymmetric field theory and the string theory in AdS. QCD in the infrared is a highly nonconformal theory, however, (which is evidenced already by the presence of the scale — the

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$^1$ These results have to be confirmed by an independent experiment. In particular, some of still missing states may be found.

$^2$ In this respect it is important to find still missing chiral multiplets for the high–spin states at the levels $M \sim 2.3$ GeV, $M \sim 2$ GeV and, possibly, at $M \sim 1.7$ GeV.
Regge slope), and such a modeling is based in fact on ad hoc and uncontrolled deformations in the infrared which simulate confinement and break conformal symmetry. The interrelations between chiral symmetry breaking and confinement, as demanded in QCD by the 't Hooft anomaly matching conditions \cite{23} and by the Coleman–Witten theorem \cite{24}, are not clear in these models and it is possible to achieve, at least in some AdS/QCD models, confinement without manifest chiral symmetry breaking in the vacuum. It is possible to choose the infrared boundary in this approach in such a way that the radial and angular Regge slopes coincide \cite{22}, however chiral multiplets of excited hadrons do not show up.

In this note we address the question whether it is possible or not to reconcile a view of an excited hadron as an open bosonic string (electric flux tube) with quarks at the ends with the approximately restored chiral symmetry in this hadron. It is generally assumed within such a picture that quarks serve mainly as sources of the colour–electric field. The energy of the system is stored in the gluonic electric string (flux tube) which generates an effective interquark interaction traditionally described in terms of the string orbital angular momentum $L$ (for example, in a specific approach of this kind \cite{25} the resulting interquark interaction is strongly angular momentum ($L$) dependent, is nonlocal and does not amount to a plain potential \cite{26}). The spins of the quarks, irrelevant for the dynamics of the string, are used then to obtain the total spin of the hadron within the LS-coupling scheme, $\vec{J} = \vec{L} + \vec{S}$, i.e. $J = L - 1, L, L + 1$ (the dependence on the quark spins orientation enters for low–lying states only through corrections due to spin–dependent terms in the effective interquark potential). Then the energy of the hadron is determined, like in the Nambu–Goto string, by the orbital angular momentum $L$ of the flux tube, that describes its rotational motion, and by the radial quantum number $n$, that describes the vibrational motion of the string (we ignore here possible transverse excitations of the flux tube and hence do not consider hybrids in this paper).

Such a picture has obtained a certain support in lattice simulations with the static sources \cite{27}, where at large separation between the sources only electric flux tube is seen, that can be approximated by a linear potential. The observed flux tube is highly nondynamical, however, because its ends are static. Both the Nambu–Goto string as well as the static lattice simulations do not appeal to chiral degrees of freedom which are present in the light quark hadrons.

There are notorious experimental and theoretical reasons to believe that chiral symmetry is highly relevant and is actually approximately manifest in the highly excited hadrons. Then the question arises whether the popular nonrelativistic $\{2S+1L_J\}$–inspired classification scheme and dynamical pictures based upon it are adequate for highly excited hadrons? In other words, we put in question the very possibility to describe highly excited hadrons in terms of the quantum numbers $n$ and $L$, for the latter do not comply with chiral symmetry. Hence, in a dynamically generated hadron consisting of the quarks at the ends of the string, the role of the quarks and in particular of their spin orientations is highly nontrivial, contrary to the traditional belief. Below we prove this statement. As a byproduct it is then clear that the nonrelativistic $\{2S+1L_J\}$ classification, traditionally used in the quark picture and somewhat justified in the heavy quark systems, makes no sense in the highly excited light–quark hadrons with the approximately restored chiral symmetry, because $L$ is not a separately conserved quantity. The chiral basis is to be used instead.

Before diving into technicalities we briefly outline the content of the proof. Since both the basis of chiral multiplets and the $\{2S+1L_J\}$ basis are complete ones, then we can construct a unitary transformation from one basis to the other, similar to the transformation from the two–body relativistic helicity basis to the nonrelativistic one \cite{28}. This transformation is a pure mathematical enterprise and is completely model–independent. While the derivation of this unitary transformation is rather straightforward and simple, it has quite interesting physical implications directly relevant to hadron modelling. It turns out that a given chiral state that belongs to some of the chiral multiplets is a fixed superposition of two $\{2S+1L_J\}$ basis states with different $L$‘s. Hence if chiral symmetry is approximately restored in a given excited hadron, its description in terms of a fixed $L$ is impossible. The energy of the dynamical flux tube cannot be determined solely by its orbital angular momentum $L$ and can depend generally only on $J$ and on the chiral index. Hence the role of the quark spin orientations is highly nontrivial for the energy of the dynamical flux tube.

Chiral multiplets of the $SU(2)_L \times SU(2)_R$ group for $q\bar{q}$ mesons were classified in Refs. \cite{3, 6}. They can be naturally described in terms of the chiral basis with the basis vectors depending on the index of the chiral representation $R$ ($R = (0,0)$, $(1/2,1/2)_a$, $(1/2,1/2)_b$, or $(0,1) + (1,0)$), on the spatial parity $P$, on the total spin $J$ and its projection $M$, as well as on the isospin $I$ and its projection $i$. To simplify notations, we omit the spin

| $R$ | $J = 0$ | $J = 1, 3, ...$ | $J = 2, 4, ...$ |
|-----|---------|-----------------|-----------------|
| $(0,0)$ | $0J^+$ ↔ $0J^-$ | $0J^+$ ↔ $0J^+$ | $0J^+$ ↔ $0J^+$ |
| $(1/2,1/2)_a$ | $1J^+ ↔ 0J^{++}$ | $0J^+ ↔ 0J^- ↔ 1J^+$ | $0J^+ ↔ 0J^+$ |
| $(1/2,1/2)_b$ | $1J^{++} ↔ 0J^{++}$ | $1J^- ↔ 0J^+ ↔ 1J^+$ | $0J^- ↔ 1J^+ ↔ 0J^+$ |
| $(0,1) + (1,0)$ | $1J^- ↔ 1J^+ ↔ 1J^−$ | $1J^- ↔ 1J^+ ↔ 1J^−$ | $1J^- ↔ 1J^+ ↔ 1J^−$ |

### Table I: The complete set of $q\bar{q}$ states classified according to the chiral basis.
and isospin projections in the notation of the chiral basis vectors, referring to them as to $|R; IJ^{PC}\rangle$, where for a neutral $q\bar{q}$ system the $C$–parity is related to the other quantum numbers in the standard way $^{28}$. The chiral basis $\{R; IJ^{PC}\}$ is obviously consistent with the Poincaré invariance. Highly excited hadrons belonging to the same chiral representation $R$ but possessing opposite parities $P$ are approximately degenerate. Hence the chiral basis is extremely convenient for describing hadron wave functions. In Table II we give the complete set of $q\bar{q}$ states classified according to this chiral basis.

The chiral basis vectors can be constructed as 

$$
|R; IJ^{PC}\rangle = \sum_{\lambda_q \lambda_{\bar{q}}} \sum_{i_q} \chi^{R_{\lambda_q \lambda_{\bar{q}}}} \langle l_q | q\bar{q} \rangle |J \lambda_q \lambda_{\bar{q}}\rangle , \tag{1}
$$

where $i_q$ and $\lambda_q$ ($i_{\bar{q}}$ and $\lambda_{\bar{q}}$) are the quark (antiquark) isospin and helicity, respectively. The coefficients $\chi^{R_{\lambda_q \lambda_{\bar{q}}}}$ can be extracted from the explicit form of the basis vectors given in Refs. $^{5, 6}$. They form a unitary transformation from the quark helicity basis to the chiral basis with definite parity in the state. In Table II we give these coefficients for various chiral representations and quantum numbers. Notice that the chiral basis is closely related to the notion of the quark helicity as helicity of the quark coincides with its chirality while helicity of the antiquark is just opposite to its chirality. In other words, one can refer to the states $|J \lambda_q \lambda_{\bar{q}}\rangle$ as to the states with the given chirality of the quarks. Vectors $|J \lambda_q \lambda_{\bar{q}}\rangle$ do not possess definite parity — states with a definite parity can be constructed only as linear combinations of such vectors with opposite helicity.

To proceed we are to build an explicit form of the vectors $|J \lambda_q \lambda_{\bar{q}}\rangle$ in terms of the single–quark states with the given helicity and to establish the relation between these states and the basis vectors of the $^{2S+1}L_J$ scheme, that is to find the matrix $\langle J \lambda_q \lambda_{\bar{q}} |^{2S+1}L_J \rangle$, where $L$ is the orbital angular momentum and $S$ is the total spin of two particles. Then, following Ref. $^{28}$, we use$^3$

$$
|J \lambda_q \lambda_{\bar{q}}\rangle = D_{J_{\lambda_q \lambda_{\bar{q}}, L}}(\vec{n}) \sqrt{\frac{2L+1}{4\pi}} |l_q\rangle , \tag{2}
$$

where $D_{J_{\lambda_q \lambda_{\bar{q}}, L}}(\vec{n})$ is the standard Wigner $D$–function describing rotation from the quantization axis to the quark momentum direction $\vec{n} = \vec{p}/p$. Finally, after simple algebraic transformations, we find:

$$
\langle J \lambda_q \lambda_{\bar{q}} |^{2S+1}L_J \rangle = \sqrt{\frac{2L+1}{2J+1}} C^{\text{SA}}_{\lambda_q \lambda_{\bar{q}}, \Lambda} C^{\text{JA}}_{L \overline{\Lambda} \Sigma}, \tag{3}
$$

where $\Lambda = \lambda_q - \lambda_{\bar{q}}$.

Combining Eqs. $^{1} – ^{3}$ together and using the property that $^{2S+1}L_J |J \lambda_q \lambda_{\bar{q}}\rangle = \langle J \lambda_q \lambda_{\bar{q}} |^{2S+1}L_J \rangle^*$, we arrive at:

$$
|R; IJ^{PC}\rangle = \sum_{LS} \chi^{R_{\lambda_q \lambda_{\bar{q}}}^{\text{PC}}} \sum_{i_q} \chi_{i_q}^{R_{\lambda_q \lambda_{\bar{q}}}^{\text{PC}}} \langle i_q | J \lambda_q \lambda_{\bar{q}} \rangle |J \lambda_q \lambda_{\bar{q}}\rangle \times \sqrt{\frac{2L+1}{2J+1}} C^{\text{SA}}_{\lambda_q \lambda_{\bar{q}}, \Lambda} C^{\text{JA}}_{L \overline{\Lambda} \Sigma} |^{2S+1}L_J \rangle. \tag{4}
$$

Consequently, one ends up with a unitary transformation from the basis vectors of the $SU(2)_L \times SU(2)_R$ chiral group with the fixed $IJ^{PC}$ quantum numbers to the $\{I;^{2S+1}L_J\}$ basis:

$$
|R; IJ^{PC}\rangle = \sum_{LS} \chi^{R_{\lambda_q \lambda_{\bar{q}}}^{\text{PC}}} \sum_{i_q} \chi_{i_q}^{R_{\lambda_q \lambda_{\bar{q}}}^{\text{PC}}} \langle i_q | J \lambda_q \lambda_{\bar{q}} \rangle |J \lambda_q \lambda_{\bar{q}}\rangle \times \sqrt{\frac{2L+1}{2J+1}} C^{\text{SA}}_{\lambda_q \lambda_{\bar{q}}, \Lambda} C^{\text{JA}}_{L \overline{\Lambda} \Sigma} |I;^{2S+1}L_J \rangle. \tag{5}
$$

Hence every state from the chiral basis is a fixed superposition of allowed states in the $\{I;^{2S+1}L_J\}$ basis. Note that in the sum above only those $L$ are allowed which satisfy $P = (-1)^L$.

The unitary transformation derived above is a model–independent result. It implies stringent limitations on description of physical states with approximate chiral symmetry. If chiral symmetry breaking in a given physical state is only a small perturbation, this state is described by one of the basis vectors $|R; IJ^{PC}\rangle$ to a good accuracy. For particular basis vectors the sum in Eq. $^{5}$ is restricted to only one fixed combination $\{L, S\}$ (still, there is a summation over the quark helicities). For example, for the axial–vector mesons $a_1$, described by the chiral vector $|(0, 1) + (1, 0); 1 \ 1^{+}\rangle$, the sum in $L$ and $S$ contains only one term, namely $|1;^3P_1\rangle$. Similarly, the $h_1$ meson corresponds to the state $|(1/2, 1/2); 0 \ 1^{–}\rangle$, in the chiral basis, and to the state $|0;^1P_1\rangle$, in the $\{I;^{2S+1}L_J\}$ basis. However, there are two kinds of the $\rho$–mesons, given by two different vectors, $|(0, 1) + (1, 0); 1 \ 1^{-}\rangle$ and $|(1/2, 1/2); 1 \ 1^{-}\rangle$, in the chiral basis. Each of these $\rho$–mesons is represented by the mutually orthogonal fixed superpositions of two different partial waves:

$$
|(0, 1) + (1, 0); 1 \ 1^{-}\rangle = \sqrt{\frac{2}{3}} |1;^3S_1\rangle + \sqrt{\frac{1}{3}} |1;^3D_1\rangle ,
$$

$$
|(1/2, 1/2); 1 \ 1^{-}\rangle = \sqrt{\frac{1}{3}} |1;^3S_1\rangle - \sqrt{\frac{2}{3}} |1;^3D_1\rangle .
$$

Hence the description of the $\rho$–meson with approximately restored chiral symmetry is impossible in terms of a fixed $L$. Obviously, this situation occurs for all mesons from Table II for which two different chiral representations can be assigned to the given $IJ^{PC}$ set. Consequently, for many mesons with $J > 0$ there must be a duplication of the Regge trajectories, each of them can be uniquely specified by a proper chiral index. Such a duplication of

$^3$We use the standard definition of the phase of the spherical functions, which differs by the factor $i^J$ from the definition in Ref. $^{28}$.
TABLE II: The complete set of nonzero coefficients $\chi_{nI}$.

| $R = (0,0)$ | $P = (-1)^J$ | $I = 0$ |
|---|---|---|
| $\chi_{1/2}^{RPI} = \chi_{1/2}^{RPI} = \frac{1}{\sqrt{2}}$ |

| $R = (0,0)$ | $P = (-1)^{J+1}$ | $I = 0$ |
|---|---|---|
| $\chi_{1/2}^{RPI} = -\chi_{1/2}^{RPI} = \frac{1}{\sqrt{2}}$ |

| $R = (1/2,1/2)$ | $P = (-1)^J$ | $I = 0$ |
|---|---|---|
| $\chi_{1/2}^{RPI} = \chi_{1/2}^{RPI} = \frac{1}{\sqrt{2}}$ |

| $R = (1/2,1/2)$ | $P = (-1)^{J+1}$ | $I = 1$ |
|---|---|---|
| $\chi_{1/2}^{RPI} = -\chi_{1/2}^{RPI} = \frac{1}{\sqrt{2}}$ |

| $R = (0,1) \oplus (1,0)$ | $P = (-1)^J$ | $I = 1$ |
|---|---|---|
| $\chi_{1/2}^{RPI} = \chi_{1/2}^{RPI} = \frac{1}{\sqrt{2}}$ |

| $R = (0,1) \oplus (1,0)$ | $P = (-1)^{J+1}$ | $I = 1$ |
|---|---|---|
| $\chi_{1/2}^{RPI} = -\chi_{1/2}^{RPI} = \frac{1}{\sqrt{2}}$ |

the Regge trajectories is indeed observed in the Crystal Barrel data [12].

One has to conclude that generally the description of the highly excited hadronic states with approximately restored chiral symmetry is impossible in terms of the fixed $L$. This simple result calls for a review of the thirty–year tradition in classification of the light–quark sector based on the naive quark model scheme[4]. A naive string picture with the fixed $L$ is to be revised either to include quark chiralities and thus to become compatible with chiral symmetry [6, 7]. A possible string description of the large experimental degeneracy of mesons with different spins, isospins, and parities would amount to demonstration that the string energy is determined only by the principal quantum number and hence is independent on $|R; JPC|$. In particular, the large degeneracy seen at Fig. 2 of Ref. [6] might be understood if, on top of the chiral symmetry restoration, a principal quantum number $N = n + J$ existed. Note that the chiral restoration requires chiral multiplets for the highest spin states, which are presently missing in Refs. [1, 2]. Hence it is a very important experimental task to find them or to reliably establish their absence[5].

L. Ya. G. is grateful to Tom Cohen for useful comments and acknowledges support of the Austrian Science Fund through grant P19168-N16. Work of A.N. was supported by the Federal Agency for Atomic Energy of Russian Federation, by the grants RFFI-05-02-04012-NNO, DFG-436 RUS 113/820/0-1(R), NSh-843.2006.2, and PTDC/FIS/70843/2006-Fisica, as well as by the Federal Programme of the Russian Ministry of Industry, Science, and Technology No. 40.052.1.1.1112.

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Note that the notion of the spin–orbit force is not defined for the states with a definite chirality since the chirality operator does not commute with the orbital angular momentum operator $L$. This is directly relevant to the widely discussed “problem” of missing spin–orbit force [6, 7].

Such high–spin parity doublets are well seen in the nucleon spectrum — see Fig. 1 of Ref. [6].
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