Chiral multiplets of excited mesons.

L. Ya. Glozman

Institute for Theoretical Physics, University of Graz, Universitätsplatz 5, A-8010 Graz, Austria

Abstract

It is shown that experimental meson states with spins \( J=0,1,2,3 \) in the energy range 1.9 - 2.4 GeV obtained in a recent partial wave analysis of proton-antiproton annihilation at LEAR remarkably confirm all predictions of chiral symmetry restoration. Classification of excited \( \bar{q}q \) mesons according to the representations of chiral \( U(2)_L \times U(2)_R \) group is performed. There are two important predictions of chiral symmetry restoration in highly excited mesons: (i) physical states must fill out approximately degenerate parity-chiral multiplets; (ii) some of the physical states with the given \( I,J^{PC} \) are members of one parity-chiral multiplet, while the other states with the same \( I,J^{PC} \) are members of the other parity-chiral multiplet. For example, while some of the excited \( \rho(1,1^{-+}) \) states are systematically degenerate with \( a_1(1,1^{++}) \) states forming \((0,1)+(1,0)\) chiral multiplets, the other excited \( \rho(1,1^{-+}) \) states are degenerate with \( h_1(0,1^{++}) \) states \((1/2,1/2)\) chiral multiplets). Hence, one of the predictions of chiral symmetry restoration is that the combined amount of \( a_1(1,1^{++}) \) and \( h_1(0,1^{++}) \) states must coincide with the amount of \( \rho(1,1^{-+}) \) states in the chirally restored regime. It is shown that the same rule applies (and experimentally confirmed) to many other meson states.

I. INTRODUCTION

A recent partial wave analysis [1–4] of \( \bar{p}p \) annihilation at LEAR in the energy range 1.9 - 2.4 GeV has revealed a lot of new meson states in this mass region. Very many of them do not fit into the traditional potential description which is based on the \( ^2S+1L_J \) classification scheme. For instance, they cannot be accommodated by the constituent quark model [6]. Some of the particular examples are discussed in refs. [7,8]. If these experimental results are correct, then it means that a "harmony" of the potential description is certainly not adequate for highly excited states (where the valence quarks should be expected to be ultra-relativistic) and some other physical picture and classification scheme must be looked for.

*e-mail: leonid.glozman@uni-graz.at
FIG. 1. Pion and $n\bar{n} = \frac{u\bar{u} + d\bar{d}}{\sqrt{2}} f_0$ spectra. Since these $f_0$ states are obtained in $p\bar{p}$ and they decay predominantly into $\pi\pi$ channel, they are considered as $n\bar{n}$ states.

In ref. [9] it has been suggested that spontaneously broken chiral symmetry of QCD must be restored in the upper part of hadron spectra, which is evident by the (almost) systematical parity doubling of $N$ and $\Delta$ states at $M \geq 1.7$ GeV. This idea has been substantiated on the basis of the operator product expansion in QCD and analyticity of the two-point function (dispersion relation) [10,11]. This effective chiral symmetry restoration has been referred to as chiral symmetry restoration of the second kind [12] in order to distinguish this phenomenon from the other phenomena of chiral symmetry restoration in the vacuum state at high temperature and/or density. The essence of the present phenomenon is that the quark condensates which break chiral symmetry in the vacuum state (and hence in the low-lying excitations over vacuum) become simply irrelevant (unimportant) for the physics of the highly excited states and the physics here is such as if there were no chiral symmetry breaking in the vacuum. The valence quarks simply decouple from the quark condensates and consequently the notion of the constituent quarks with dynamical mass induced by chiral symmetry breaking becomes irrelevant in highly excited hadrons [9,13]. Instead, the string picture with quarks of definite chirality at the end points of the string should be invoked [8]. In recent lattice calculations DeGrand has demonstrated that indeed in the highly excited mesons valence quarks decouple from the low-lying eigenmodes of the Dirac operator (which determine the quark condensate via Banks-Casher relation) and hence decouple from the quark condensate of the QCD vacuum [14].

If it is true that chiral symmetry is effectively restored, then classification of hadrons should be performed according to the chiral group of QCD. This group strongly constrains the amount of certain mesons and their relative energies. In ref. [7] we have classified all excited mesons with spin $J = 0$ according to the chiral group and have shown that indeed the pattern of these mesons above 1.7 GeV region certainly favours chiral symmetry restoration.
of the second kind. This is illustrated in Fig. 1 where $\pi(I = 1, 0^+)$ and $\bar{q}q\ f_0(I = 0, 0^{++})$ states [7] are shown (which must be chiral partners in the chiral symmetry restored regime).

The symmetry restoration implies that it must be seen also in all other $\bar{q}q$ mesonic states. The purpose of this paper is to classify observed $\bar{q}q$ mesons according to chiral representations and demonstrate that the experimental patterns of $J = 1, 2, 3$ mesons do confirm this phenomenon. Experimental data on higher spin states are scarce and it is not yet possible to provide any systematic description though predictions can be made.

II. CLASSIFICATION OF THE $\bar{Q}Q$ MESONS ACCORDING TO THE CHIRAL GROUP.

Mesons reported in ref. [1–4] are obtained in $\bar{p}p$ annihilations, hence according to OZI rule we have to expect them to be $\bar{q}q$ states with $u$ and $d$ valence quark content. Hence we will consider

$$U(2)_L \times U(2)_R = SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A,$$

the full chiral group of the QCD Lagrangian. In the following chiral symmetry will refer to specifically the $SU(2)_L \times SU(2)_R$ symmetry, which is perfect due to very small masses of $u, d$ quarks as compared to the only dimensional parameter of QCD, $\Lambda_{QCD}$, and the typical hadronic scale of 1 GeV.

$SU(2)_L \times SU(2)_R$ is a symmetry with respect to two independent rotations of the left-handed and right-handed quarks in the isospin space. Hence the irreducible representations of this group can be specified by the isospins of the left and right quarks, $(I_L, I_R)$. The total isospin of the state can be obtained from the left and right isospins according to the standard angular momentum addition rules

$$I = |I_L - I_R|, ..., I_L + I_R.$$  

All hadronic states are characterised by a definite parity. However, not all irreducible representations of the chiral group are invariant under parity. Indeed, parity transforms the left quarks into the right ones and vice versa. Hence while representations with $I_L = I_R$ are invariant under parity (i.e. under parity operation every state in the representation transforms into the state of opposite parity within the same representation), this is not true for the case $I_L \neq I_R$. In the latter case parity transforms every state in the representation $(I_L, I_R)$ into the state in the representation $(I_R, I_L)$. We can construct definite parity states only combining basis vectors from both these irreducible representations. Hence it is only the direct sum of these two representations

$$(I_L, I_R) \oplus (I_R, I_L), \quad I_L \neq I_R.$$  

that is invariant under parity. This reducible representation of the chiral group is an irreducible representation of the larger group, parity-chiral group [11].
When we consider mesons of isospin $I = 0, 1$, only three independent irreducible representations of the parity-chiral group exist.

(i) $(0,0)$. Mesons in this representation must have isospin $I = 0$. At the same time $I_R = I_L = 0$. This can be achieved when either there are no valence quarks in the meson, or both valence quark and antiquark are right or left. If we denote $R = (u_R, d_R)$ and $L = (u_L, d_L)$, then the basis states of both parities can be written as

$$ |(0,0); \pm; J \rangle = \frac{1}{\sqrt{2}}(\bar{R}R \pm \bar{L}L)_J. $$

Note that such a system can have spin $J \geq 1$. Indeed, valence quark and antiquark in the state (4) have definite helicities, because generically helicity = +chirality for quarks and helicity = -chirality for antiquarks. Hence the total spin projection of the quark-antiquark system onto the momentum direction of the quark is $\pm 1$. The parity transformation property of the quark-antiquark state is then regulated by the total spin of the system [16]

$$ \hat{P}|(0,0); \pm; J \rangle = \pm(-1)^J |(0,0); \pm; J \rangle. $$

(ii) $(1/2,1/2)$. In this case the quark must be right and the antiquark must be left, and vice versa. These representations combine states with $I=0$ and $I=1$, which must be of opposite parity. The basis states within the two distinct representations of this type are

$$ |(1/2,1/2); +; I = 0; J \rangle = \frac{1}{\sqrt{2}}(\bar{R}L + \bar{L}R)_J, $$

$$ |(1/2,1/2); -; I = 1; J \rangle = \frac{1}{\sqrt{2}}(\bar{R}\tau L - \bar{L}\tau R)_J, $$

and

$$ |(1/2,1/2); -; I = 0; J \rangle = \frac{1}{\sqrt{2}}(\bar{R}L - \bar{L}R)_J, $$

$$ |(1/2,1/2); +; I = 1; J \rangle = \frac{1}{\sqrt{2}}(\bar{R}\tau L + \bar{L}\tau R)_J. $$

In these expressions $\tau$ are isospin Pauli matrices. The parity of every state in the representation is determined as

$$ \hat{P}|(1/2,1/2); \pm; I; J \rangle = \pm(-1)^J |(1/2,1/2); \pm; I; J \rangle. $$

1Hence glueballs must be classified according to this representation [15]; with no quark content this representation contains the state of only one parity.
Note that the two distinct \((1/2,1/2)\) irreducible representations of \(SU(2)_L \times SU(2)_R\) form one irreducible representation of \(U(2)_L \times U(2)_R\).

(iii) \((0,1) \oplus (1,0)\). The total isospin is 1 and the quark and antiquark must both be right or left. This representation is possible only for \(J \geq 1\). The basis states are

\[
| (0,1) + (1,0); \pm; J \rangle = \frac{1}{\sqrt{2}} (\bar{R} \tau R \pm \bar{L} \tau L)_J
\]

with parities

\[
\hat{P} | (0,1) + (1,0); \pm; J \rangle = \pm (-1)^J | (0,1) + (1,0); \pm; J \rangle.
\]

In the chirally restored regime the physical states must fill out completely some or all of these representations. We have to stress that usual quantum numbers \(I, J^{PC}\) are not enough to specify the chiral representation. It happens that some of the physical particles with the given \(I, J^{PC}\) belong to one chiral representation (multiplet), while the other particles with the same \(I, J^{PC}\) belong to the other multiplet. Classification of the particles according to \(I, J^{PC}\) is simply not complete in the chirally restored regime. This property will have very important implications as far as the amount of the states with the given \(I, J^{PC}\) is concerned.

In order to make this point clear, we will discuss some of the examples. Consider \(\rho(1,1^{--})\) mesons. Particles of this kind can be created from the vacuum by the vector current, \(\bar{\psi} \gamma^\mu \tau^i \psi\). Its chiral partner is the axial vector current, \(\bar{\psi} \gamma^\mu \gamma^5 \tau^i \psi\), which creates from the vacuum the axial vector mesons, \(a_1(1,1^{++})\). Both these currents belong to the representation \((0,1)+(1,0)\) and have the right-right \(\pm\) left-left quark content. Clearly, in the chirally restored regime the mesons created by these currents must be degenerate level by level and fill out the \((0,1)+(1,0)\) representations. Hence, naively the amount of \(\rho\) and \(a_1\) mesons high in the spectrum should be equal. This is not correct, however. \(\rho\)-Mesons can be also created from the vacuum by other type(s) of current(s), \(\bar{\psi} \sigma^i \tau^j \psi\) (or by \(\bar{\psi} \partial^\mu \tau^i \psi\)). These interpolators belong to the \((1/2,1/2)\) representation and have the left-right \(\pm\) right-left quark content. In the regime where chiral symmetry is strongly broken (as in the low-lying states) the physical states are mixtures of different representations. Hence these low-lying states are well coupled to both \((0,1)+(1,0)\) and \((1/2,1/2)\) interpolators. However, when chiral symmetry is (approximately) restored, then each physical state must be strongly dominated by the given representation and hence will couple only to the interpolator which belongs to the same representation. This means that \(\rho\)-mesons created by two distinct currents in the chirally restored regime represent physically different particles. The chiral partner of the \(\bar{\psi} \sigma^i \tau^j \psi\) (or \(\bar{\psi} \partial^\mu \tau^i \psi\)) current is \(\varepsilon^{ijk} \bar{\psi} \sigma^j \psi \) (or \(\bar{\psi} \gamma^5 \partial^\mu \psi\), respectively) \(^2\). The latter interpolators create from the vacuum \(h_1(0,1^{+-})\) states. Hence in the chirally restored regime, some of the \(\rho\)-mesons must be degenerate with the \(a_1\) mesons ((0,1)+(1,0) multiplets), but the others - with the \(^2\)Chiral transformation properties of some interpolators can be found in ref. [17].
h_1 mesons ((1/2,1/2) multiplets). Consequently, high in the spectra the combined amount of ρ_1 and h_1 mesons must coincide with the amount of ρ-mesons. This is a highly nontrivial prediction of chiral symmetry.

Actually it is a very typical situation. Consider f_2(0,2++) mesons as another example. They can be interpolated by the tensor field \( \bar{\psi}\gamma^\mu\partial^\nu\psi \) (properly symmetrised, of course), which belongs to the (0,0) representation. Their chiral partners are \( \omega_2(0,2^-) \) mesons, which are created by the \( \bar{\psi}\gamma^5\gamma^\mu\partial^\nu\psi \) interpolator. On the other hand, \( f_2(0,2++) \) mesons can also be created from the vacuum by the \( \bar{\psi}\partial^\mu\partial^\nu\psi \) type of interpolator, which belongs to the (1/2,1/2) representation. Its chiral partner is \( \bar{\psi}\gamma^5\partial^\mu\bar{\tau}\psi \), which creates \( \pi_2(1,2^-) \) mesons. Hence in the chirally restored regime we have to expect \( \omega_2(0,2^-) \) mesons to be degenerate systematically with some of the \( f_2(0,2++) \) mesons ((0,0) representations) while \( \pi_2(1,2^-) \) mesons must be degenerate with other \( f_2(0,2++) \) mesons (forming (1/2,1/2) multiplets). Hence the total number of \( \omega_2(0,2^-) \) and \( \pi_2(1,2^-) \) mesons in the chirally restored regime must coincide with the amount of \( f_2(0,2++) \) mesons.

These examples can be generalized to mesons of any spin \( J \geq 1 \). Those interpolators which contain only derivatives \( \bar{\psi}\partial^\mu\partial^\nu\psi \) ( \( \bar{\psi}\gamma^\mu\partial^\nu\psi \), respectively) have quantum numbers \( I = 0, P = (-1)^J, C = (-1)^J \) ( \( I = 1, P = (-1)^J, C = (-1)^J \) ) and transform as (1/2,1/2). Their chiral partners are \( \bar{\psi}\gamma^5\partial^\mu\partial^\nu\psi \) ( \( \bar{\psi}\gamma^5\partial^\mu\bar{\tau}\psi \), respectively) with \( I = 1, P = (-1)^{J+1}, C = (-1)^J \) ( \( I = 0, P = (-1)^{J+1}, C = (-1)^J \), respectively). However, interpolators with the same \( I, J^{PC} \) can be also obtained with one \( \gamma^9 \) matrix instead one of the derivatives, \( \partial^\mu \). \( \bar{\psi}\gamma^5\partial^\mu\partial^\nu\gamma^9\psi \) ( \( \bar{\psi}\gamma^5\partial^\mu\bar{\tau}\psi \) ). These latter interpolators belong to (0,0) ((0,1)+(1,0)) representation. Their chiral partners are \( \bar{\psi}\gamma^5\partial^\mu\partial^\nu\gamma^9\psi \) ( \( \bar{\psi}\gamma^5\partial^\mu\bar{\tau}\psi \) ) which have \( I = 0, P = (-1)^{J+1}, C = (-1)^J \) ( \( I = 1, P = (-1)^{J+1}, C = (-1)^J \) ). Hence in the chirally restored regime the physical states created by these different types of interpolators will belong to different representations and will be distinct particles while having the same \( I, J^{PC} \). One needs to indicate chiral representation in addition to usual quantum numbers \( I, J^{PC} \) in order to uniquely specify physical states in the chirally restored regime.

In the next section we will present experimental data and show that all these predictions of chiral symmetry are indeed verified.

### III. CHIRAL MULTIPLETS OF \( J = 1, 2, 3 \) MESONS.

A detailed analysis of all chiral multiplets with \( J = 0 \) has been performed in ref. [7] and we do not repeat it here. In Table 1 we present all mesons with \( J = 1, 2, 3 \) in the energy range 1.9-2.4 GeV obtained in refs. [1-4]. So it is important to see whether these experimental results fall into chiral multiplets and whether the data set is complete. Below we will show that indeed these data remarkably confirm the predictions of chiral symmetry.

---

3Those \( \rho(1,1^{--}) \) and \( \omega(0,1^{--}) \) mesons which belong to (1/2,1/2) cannot be seen in \( e^+e^- \to \text{hadrons} \).
We will start with chiral multiplets having $J = 2$ because the data seems to be indeed complete for the $J = 2$ states.

$$
\begin{array}{ll}
(0,0) & \\
\omega_2(0, 2^{--}) & f_2(0, 2^{++}) \\
1975 \pm 20 & 1934 \pm 20 \\
2195 \pm 30 & 2240 \pm 15 \\

(1/2, 1/2) & \\
\pi_2(1, 2^{--}) & f_2(0, 2^{++}) \\
2005 \pm 15 & 2001 \pm 10 \\
2245 \pm 60 & 2293 \pm 13 \\

(1/2, 1/2) & \\
a_2(1, 2^{++}) & \eta_2(0, 2^{--}) \\
2030 \pm 20 & 2030 \pm ? \\
2255 \pm 20 & 2267 \pm 14 \\

(0, 1) + (1, 0) & \\
a_2(1, 2^{++}) & \rho_2(1, 2^{--}) \\
1950^{+30}_{-70} & 1940 \pm 40 \\
2175 \pm 40 & 2225 \pm 35 \\
\end{array}
$$

We see systematic patterns of chiral symmetry restoration. In particular, the amount of $f_2(0, 2^{++})$ mesons coincides with the combined amount of $\omega_2(0, 2^{--})$ and $\pi_2(1, 2^{++})$ states. Similarly, number of $a_2(1, 2^{++})$ states is the same as number of $\eta_2(0, 2^{--})$ and $\rho_2(1, 2^{--})$ together. All chiral multiplets are complete. While masses of some of the states will be definitely corrected in the future experiments, if new states might be discovered in this energy region in other types of experiments, they should be either $\bar{s}s$ states or glueballs.

Consider now the $J = 1$ multiplets$^4$.

$$
\begin{array}{ll}
(0,0) & \\
\omega(0, 1^{--}) & f_1(0, 1^{++}) \\
? & 1971 \pm 15 \\
? & 2310 \pm 60 \\

(1/2, 1/2) & \\
\end{array}
$$

---

$^4$The state $\rho(1, 1^{--})$ mentioned in the Table at 2110$\pm$35 is given below at the mass 2150, because it is this value for this state which is given in PDG [18]. The $\rho(1, 1^{--})$ state at 1900, used in the classification, is seen in $e^+e^-$ (see PDG) and not seen in $\bar{p}p$. 

7
\[ \omega(0, 1^{--}) \quad b_1(1, 1^{--}) \]
\[
\begin{align*}
1960 \pm 25 & \quad 1960 \pm 35 \\
2205 \pm 30 & \quad 2240 \pm 35 \\
\end{align*}
\]
\[
(1/2, 1/2)
\]
\[
\begin{align*}
\omega_1(0, 1^{--}) & \quad \rho(1, 1^{--}) \\
1965 \pm 45 & \quad 1970 \pm 30 \\
2215 \pm 40 & \quad 2150 \pm ? \\
\end{align*}
\]
\[
(0, 1)+(1, 0)
\]
\[
\begin{align*}
a_1(1, 1^{++}) & \quad \rho(1, 1^{--}) \\
1930 \pm 30 & \quad 1900 \pm ? \\
2270 \pm 30 & \quad 2265 \pm 40 \\
\end{align*}
\]

Here, like for the \( J = 0, 2 \) states, we again observe patterns of chiral symmetry restoration. Two missing \( \omega(0, 1^{--}) \) states are yet missing.

Below are the multiplets for \( J = 3 \).

\[
(0, 0)
\]
\[
\begin{align*}
\omega_3(0, 3^{--}) & \quad f_3(0, 3^{++}) \\
? & \quad 2048 \pm 8 \\
2285 \pm 60 & \quad 2303 \pm 15 \\
\end{align*}
\]
\[
(1/2, 1/2)
\]
\[
\begin{align*}
\omega_3(0, 3^{--}) & \quad b_3(1, 3^{--}) \\
1945 \pm 20 & \quad 2032 \pm 12 \\
2255 \pm 15 & \quad 2245 \pm ? \\
\end{align*}
\]
\[
(1/2, 1/2)
\]
\[
\begin{align*}
h_3(0, 3^{++}) & \quad \rho_3(1, 3^{--}) \\
2025 \pm 20 & \quad 1982 \pm 14 \\
2275 \pm 25 & \quad 2260 \pm 20 \\
\end{align*}
\]
\[
(0, 1)+(1, 0)
\]
\[
\begin{align*}
a_3(1, 3^{++}) & \quad \rho_3(1, 3^{--}) \\
2031 \pm 12 & \quad 2013 \pm 30 \\
2275 \pm 35 & \quad 2300 \pm 50 \\
\end{align*}
\]

Data on \( J \geq 4 \) is scarce. While some of the multiplets are well seen, it is not yet possible to provide any systematic analysis. The prediction is that for \( J = 4 \) the pattern should be the same as for \( J = 2 \), while for \( J = 5 \) it should be similar to \( J = 1, 3 \) cases.
IV. EVIDENCE FOR $U(1)_A$ RESTORATION.

It is important to see whether there are also signatures of the $U(1)_A$ restoration. This can happen if two conditions are fulfilled [10]: (i) unimportance of the axial anomaly in excited states, (ii) chiral $SU(2)_L \times SU(2)_R$ restoration (i.e. unimportance of the quark condensates which break simultaneously both types of symmetries in the vacuum state). Some evidence for the $U(1)_A$ restoration has been reported in ref. [7] on the basis of $J = 0$ data. Yet missing $a_0$ and $\eta$ states have to be discovered to complete the $U(1)_A$ multiplets in the $J = 0$ spectra. In this section we will demonstrate that the data on the $J = 1, 2, 3$ present convincing evidence on $U(1)_A$ restoration.

First, we have to consider which mesonic states can be expected to be $U(1)_A$ partners. The $U(1)_A$ transformation connects interpolators of the same isospin but opposite parity. But not all such interpolators can be connected by the $U(1)_A$ transformation. For instance, the vector currents $\bar{\psi}\gamma^\mu\psi$ and $\bar{\psi}\tau\gamma^\mu\psi$ are invariant under $U(1)_A$. Similarly, the axial vector interpolators $\bar{\psi}\gamma^5\gamma^\mu\psi$ and $\bar{\psi}\tau\gamma^5\gamma^\mu\psi$ are also invariant under $U(1)_A$. Hence those interpolators (states) that are members of the $(0, 0)$ and $(0, 1) + (1, 0)$ representations of $SU(2)_L \times SU(2)_R$ are invariant with respect to $U(1)_A$. However, interpolators (states) from the distinct $(1/2, 1/2)$ representations which have the same isospin but opposite parity transform into each other under $U(1)_A$. For example, $\bar{\psi}\psi \leftrightarrow \bar{\psi}\gamma^5\psi$, $\bar{\psi}\tau\psi \leftrightarrow \bar{\psi}\tau\gamma^5\psi$, and those with derivatives: $\bar{\psi}\partial^\mu\psi \leftrightarrow \bar{\psi}\gamma^5\partial^\mu\psi$, $\bar{\psi}\tau\partial^\mu\psi \leftrightarrow \bar{\psi}\tau\gamma^5\partial^\mu\psi$, etc. If the corresponding states are systematically degenerate, then it is a signal that $U(1)_A$ is restored. In what follows we show that it is indeed the case.

**J=1**

| State | $J$ | $I^- | $I^+ |
|-------|----|------|------|
| $\omega(0, 1^-)$ | 1 | 0 | 1 |
| $1960 \pm 25$ | 1965 | ± 45 |
| $2205 \pm 30$ | 2215 | ± 40 |
| $b_1(1, 1^-)$ | 1 | 1 |
| $1960 \pm 35$ | 1970 | ± 30 |
| $2240 \pm 35$ | 2150 | ± ? |

**J=2**

| State | $J$ | $I^- | $I^+ |
|-------|----|------|------|
| $f_2(0, 2^+)$ | 2 | 0 |
| $2001 \pm 10$ | 2030 | ± ? |
| $2293 \pm 13$ | 2267 | ± 14 |
| $\pi_2(1, 2^-)$ | 2 | 1 |
| $2005 \pm 15$ | 2030 | ± 20 |
| $2245 \pm 60$ | 2255 | ± 20 |

**J=3**

| State | $J$ | $I^- | $I^+ |
|-------|----|------|------|
| $\omega_3(0, 3^-)$ | 3 | 0 |
| $h_3(0, 3^+)$ | 3 | 1 |
We see systematic doublets of $U(1)_A$ restoration. Hence two distinct $(1/2,1/2)$ multiplets of $SU(2)_L \times SU(2)_R$ can be combined into one multiplet of $U(2)_L \times U(2)_R$. So we conclude that the whole chiral symmetry of the QCD Lagrangian $U(2)_L \times U(2)_R$ gets approximately restored high in the hadron spectrum.

V. DISCUSSION AND CONCLUSIONS

We have classified $\bar{q}q$ meson states in the $u,d$ sector according to chiral symmetry of QCD. Then we have presented recent experimental data on highly excited meson states which provide evidence that in the high-lying mesons both $SU(2)_L \times SU(2)_R$ and $U(1)_A$ get restored. Hence the highly excited states should be classified according to $U(2)_L \times U(2)_R = SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A$ symmetry group of QCD Lagrangian. Usual quantum numbers such as $I,J^{PC}$ are not enough to uniquely specify the states.

It is useful to quantify the effect of chiral symmetry breaking (restoration). An obvious parameter that characterises effects of chiral symmetry breaking is a relative mass splitting within the chiral multiplet. Let us define the chiral asymmetry as

$$\chi = \frac{|M_1 - M_2|}{(M_1 + M_2)},$$

where $M_1$ and $M_2$ are masses of particles within the same multiplet. This parameter gives a quantitative measure of chiral symmetry breaking at the leading (linear) order and has the interpretation of the part of the hadron mass due to chiral symmetry breaking.

For the low-lying states the chiral asymmetry is typically 0.3 - 0.6 which can be seen e.g. from a comparison of the $\rho(770)$ and $a_1(1260)$ or the $\rho(770)$ and $h_1(1170)$ masses. If the chiral asymmetry is large as above, then it makes no sense to assign a given hadron to the chiral multiplet since its wave function is a strong mixture of different representations and we have to expect also large nonlinear symmetry breaking effects. However, at meson masses about 2 GeV the chiral asymmetry is typically within 0.01, as can be seen from the multiplets presented in this paper, and in this case the hadrons can be believed to be members of multiplets with a tiny admixture of other representations. Unfortunately there are no systematic data on mesons below 1.9 GeV and hence it is difficult to estimate the chiral asymmetry as a function of mass ($\sqrt{s}$). Such a function would be crucially important for a further progress of the theory. So a systematic experimental study of hadron spectra is difficult to overestimate. However, thanks to the $0^{++}$ glueball search for the last 20 years, there are such data for $\pi$ and $f_0$ states, as can be seen from Fig. 1 (for details we
refer to [7,15]). According to these data we can reconstruct $\chi(\sqrt{s} \sim 1.3 GeV) \sim 0.03 \div 0.1$, $\chi(\sqrt{s} \sim 1.8 GeV) \sim 0.008$, $\chi(\sqrt{s} \sim 2.3 GeV) \sim 0.005$. We have to also stress that there is no reason to expect the chiral asymmetry to be a universal function for all hadron channels. Hadrons with different quantum numbers feel chiral symmetry breaking effects differently, as can be deduced from the operator product expansions of two-point functions for different currents. A task of the theory is to derive these chiral asymmetries microscopically. A comment on the present theoretical state of the arts is in order.

Naively one would expect that the operator product expansion of the two-point correlator, which is valid in the deep Euclidean domain [19], could help us. This is not so, however, for two reasons. First of all, we know phenomenologically only the lowest dimension quark condensate. Even though this condensate dominates as a chiral symmetry breaking measure at the very large space-like $Q^2$, at smaller $Q^2$ the higher dimensional condensates, which are suppressed by inverse powers of $Q^2$, are also important. These condensates are not known, unfortunately. But even if we knew all quark condensates up to a rather high dimension, it would not help us. This is because the OPE is only an asymptotic expansion [20]. While such kind of expansion is very useful in the space-like region, it does not define any analytical solution which could be continued to the time-like region at finite $s$. While convergence of the OPE can be improved by means of the Borel transform and it makes it useful for SVZ sum rules for the low-lying hadrons, this cannot be done for the higher states. So in order to estimate chiral symmetry restoration effects one indeed needs a microscopic theory that would incorporate at the same time chiral symmetry breaking and confinement.

The question is which physical picture is compatible with this symmetry restoration for highly excited mesons. It has been shown in ref. [8] that viewing excited hadrons as strings with quarks of definite chirality at the end points of the string is consistent with chiral symmetry of QCD. However, we can extend chiral symmetry to the larger symmetry. If one assumes that the spin degree of freedom of quarks (i.e. their helicity) is uncorrelated with the energy of the string, then we have to expect that all possible states of the string with all possible helicity orientations of quarks with the same spin must have the same energy. This means that all possible different $SU(2)_L \times SU(2)_R$ multiplets with the same $J$ must be degenerate. If we look carefully at the data presented in this paper we do see this fact. Does it mean that we observe a restoration of larger symmetry which includes chiral $U(2)_L \times U(2)_R$ group as a subgroup, i.e. symmetry which is higher than the symmetry of the QCD Lagrangian!? There is simpler solution to this problem. There is a reducible representation which combines $(0,0)$, $(1/2,1/2)$, $(1/2,2/2)$ and $(0,1) + (1,0)$ - it is $[(0,1/2) \oplus (1/2,0)] \times [(0,1/2) \oplus (1/2,0)]$, which is still a representation of the parity-chiral group, i.e. it is still a representation of the symmetry group of the QCD Lagrangian. This representation is nothing else but a product of the massless quark and antiquark fields. Hence a degeneracy within such a reducible representation is indeed compatible with the string with massless quarks with definite chirality at the end points of the string.
VI. ACKNOWLEDGEMENTS

I am grateful to D.V. Bugg for comments on the data [1–4]. I am also thankful to C. Lang for careful reading of the manuscript. The work was supported by the FWF project P16823-N08 of the Austrian Science Fund.
REFERENCES

[1] A. V. Anisovich et al, Phys. Lett. B491 (2000) 47.
[2] A. V. Anisovich et al, Phys. Lett. B517 (2001) 261.
[3] A. V. Anisovich et al, Phys. Lett. B542 (2002) 8.
[4] A. V. Anisovich et al, Phys. Lett. B542 (2002) 19.
[5] A. V. Anisovich et al, Phys. Lett. B513 (2002) 281.
[6] S. Godfrey, N. Isgur, Phys. Rev. D32 (1985) 189.
[7] L. Ya. Glozman, Phys. Lett. B539 (2002) 257.
[8] L. Ya. Glozman, Phys. Lett. B541 (2002) 115.
[9] L. Ya. Glozman, Phys. Lett. B475 (2000) 329.
[10] T. D. Cohen and L. Ya. Glozman, Phys. Rev. D65 (2002) 016006.
[11] T. D. Cohen and L. Ya. Glozman, Int. J. Mod. Phys. A17 (2002) 1327.
[12] L. Ya. Glozman, Progr. Part. Nucl. Phys., 50 (2003) 247.
[13] E. Swanson, hep-ph/0309296.
[14] T. DeGrand, hep-ph/0310303.
[15] L. Ya. Glozman, hep-ph/0301012, Eur. Phys. J, A19 (2004).
[16] L.D. Landau and E. M. Lifshitz, Course of Theoretical Physics, vol. IV, §69, Pergamon, 1982
[17] T. D. Cohen and X. Ji, Phys. Rev. D55 (1997) 6870.
[18] PDG, K. Hagiwara et al, Phys. Rev. D66 (2002) 010001.
[19] M. A. Shifman, A. I. Vainstein, and V. I. Zakharov, Nucl. Phys. B147 (1979) 385.
[20] A. R. Zhitnitsky, Phys. Rev. D53 (1996) 5821.

Table 1 Mesons obtained in $\bar{p}p$ annihilation at 1.9 - 2.4 GeV. The extra broad $f_2(0, 2^{++})$ state with $M = 2010 \pm 25, \Gamma = 495 \pm 35$ seen in the $\eta\eta$ channel only is not included in the Table. If it is a real state, then it can be either glueball or $\bar{s}s$ state because it is not seen in the $\pi\pi$ channel.
| Meson | $J^{PC}$ | Mass (MeV) | Width (MeV) | Reference |
|-------|---------|------------|-------------|-----------|
| $f_1$ | 0 1++   | 1971 ± 15  | 240 ± 45    | 1         |
| $f_1$ | 0 1++   | 2310 ± 60  | 255 ± 70    | 1         |
| $f_2$ | 0 2++   | 1934 ± 20  | 271 ± 25    | 1         |
| $f_2$ | 0 2++   | 2001 ± 10  | 312 ± 32    | 1         |
| $f_2$ | 0 2++   | 2240 ± 15  | 241 ± 30    | 1         |
| $f_2$ | 0 2++   | 2293 ± 13  | 216 ± 37    | 1         |
| $\eta_2$ | 0 2−+ | 2030±?    | 205±?       | 1         |
| $\eta_2$ | 0 2−+ | 2267 ± 14 | 290 ± 50    | 1         |
| $f_3$  | 0 3++   | 2048 ± 8  | 213 ± 34    | 1         |
| $f_3$  | 0 3++   | 2303 ± 15 | 214 ± 29    | 1         |
| $a_1$  | 1 1++   | 1930±30   | 155 ± 45    | 2         |
| $a_1$  | 1 1++   | 2270±35   | 305±70      | 2         |
| $a_2$  | 1 2++   | 1950±30   | 180±30      | 2         |
| $a_2$  | 1 2++   | 2030 ± 20 | 205 ± 30    | 2         |
| $a_2$  | 1 2++   | 2175 ± 40 | 310±45      | 2         |
| $a_2$  | 1 2++   | 2255 ± 20 | 230 ± 15    | 2         |
| $\pi_2$ | 1 2−+ | 2005 ± 15 | 200 ± 40    | 2         |
| $\pi_2$ | 1 2−+ | 2245 ± 60 | 320±100     | 2         |
| $a_3$  | 1 3++   | 2031 ± 12 | 150 ± 18    | 2         |
| $a_3$  | 1 3++   | 2275 ± 35 | 350±100     | 2         |
| $\rho$ | 1 1−−   | 1970 ± 30 | 260 ± 45    | 3         |
| $\rho$ | 1 1−−   | 2110 ± 35 | 230 ± 50    | 3         |
| $\rho$ | 1 1−−   | 2265 ± 40 | 325 ± 80    | 3         |
| $b_1$  | 1 1−+   | 1960 ± 35 | 230 ± 50    | 3         |
| $b_1$  | 1 1−+   | 2240 ± 35 | 320 ± 85    | 3         |
| $\rho_2$ | 1 2−− | 1940 ± 40 | 155 ± 40    | 3         |
| $\rho_2$ | 1 2−− | 2225 ± 35 | 335±100     | 3         |
| $\rho_3$ | 1 3−−  | 1982 ± 14 | 188 ± 24    | 3         |
| $\rho_3$ | 1 3−−  | 2013 ± 30 | 165 ± 35    | 5         |
| $\rho_3$ | 1 3−−  | 2260 ± 20 | 160 ± 25    | 3         |
| $\rho_3$ | 1 3−−  | 2300±50  | 340 ± 50    | 1         |
| $b_3$  | 1 3++   | 2032 ± 12 | 117 ± 11    | 3         |
| $b_3$  | 1 3++   | 2245±?    | 320 ± 70    | 3         |
| $\omega$ | 0 1−−  | 1960 ± 25 | 195 ± 60    | 4         |
| $\omega$ | 0 1−−  | 2205 ± 30 | 350 ± 90    | 4         |
| $h_1$  | 0 1−+   | 1965 ± 45 | 345 ± 75    | 4         |
| $h_1$  | 0 1−+   | 2215 ± 40 | 325 ± 55    | 4         |
| $\omega_2$ | 0 2−− | 1975 ± 20 | 175 ± 25    | 4         |
| $\omega_2$ | 0 2−− | 2195 ± 30 | 225 ± 40    | 4         |
| $\omega_3$ | 0 3−−  | 1945 ± 20 | 115 ± 22    | 4         |
| $\omega_3$ | 0 3−−  | 2255 ± 15 | 175 ± 30    | 4         |
| $\omega_3$ | 0 3−−  | 2285 ± 60 | 230 ± 40    | 4         |
| $h_3$  | 0 3++   | 2025 ± 20 | 145 ± 30    | 4         |
| $h_3$  | 0 3++   | 2275 ± 25 | 190 ± 45    | 4         |