SU(2)_L \times SU(2)_R and U(1)_A restorations high in the hadron spectrum and what it tells us about.

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Abstract

Recent data for highly excited mesons suggest that not only the chiral SU(2)_L \times SU(2)_R symmetry of QCD is restored high in the spectrum but also the U(1)_A symmetry. This means that it is not a confining interaction in QCD which triggers the spontaneous breaking of chiral symmetry. The restoration of the U(2)_L \times U(2)_R symmetry of the QCD Lagrangian implies the appearance of multiplets of this group high in the hadron spectra. Such type of multiplets is naturally explained within the string picture of confinement. It also supports the scenario that the U(1)_A breaking is related to instantons and not to the gluonic interaction responsible for confinement.

I. INTRODUCTION

It has recently been suggested that high in the hadron spectra the spontaneously broken chiral symmetry of the QCD Lagrangian is effectively restored [1]. This phenomenon can be understood in very general terms from the well established concepts of quark-hadron duality, the validity of the operator product expansion in QCD at large space-like momenta, and the validity of the dispersion relation for the two-point correlator, which connects the space-like and time-like regions [2,3]. The phenomenological manifestation of the effective chiral symmetry restoration is that the high-lying hadrons in the (u, d) quark sector must fall into multiplets of the parity-chiral group; they manifest themselves as parity doublets or higher multiplets containing degenerate states of opposite parity [2-4]. This phenomenon does not mean that the spontaneous breaking of chiral symmetry in the QCD vacuum disappears, but rather that it becomes irrelevant once we are sufficiently high in the spectrum. While the chiral symmetry breaking condensates are crucially important for the physics of the low-lying states, the physics of the high-lying hadrons is such as if there were no chiral symmetry breaking in the vacuum. A very natural picture for the highly excited hadrons then is that they represent rotating strings (with the color-electric field in the string) with
practically massless bare quarks of definite chirality at the ends of the string and these valence quarks are combined into the parity-chiral multiplets \[4\].

The general phenomenon of chiral \(SU(2)_L \times SU(2)_R\) symmetry restoration high in the spectrum as well as the physical picture of excited hadrons as strings is also well compatible with the restoration of the higher symmetry \(U(2)_L \times U(2)_R\). For the latter it is necessary that not only the chiral symmetry is restored but also the \(U(1)_A\) symmetry. For the restoration of the \(U(1)_A\) symmetry both the explicit \(U(1)_A\) breaking through the axial anomaly and the spontaneous \(U(1)_V \times U(1)_A \rightarrow U(1)_V\) breaking must become unimportant \[2\]. While the effective chiral symmetry restoration high in the spectrum implies that the spontaneous \(U(1)_V \times U(1)_A\) breaking must also become irrelevant (because both are broken by the same quark condensates in the QCD vacuum), it does by no means cause the effects of the explicit \(U(1)_A\) breaking to disappear high in the spectra. The latter could only happen if the gluonic interactions that couple to the flavor-singlet current via the axial anomaly became unimportant due to some reason.

The main purpose of this Letter is to show that recent data on highly excited mesons obtained from the partial wave analysis of proton-antiproton annihilation at LEAR \[5,6\] indicate that not only the chiral symmetry but also the full \(SU(2)_L \times SU(2)_R\) symmetry of the QCD Lagrangian get restored.

The second purpose is to argue which possible picture could be behind both the chiral and \(U(1)_A\) symmetries restoration high in the hadron spectra. Historically the \(U(1)_A\) problem began with the observation that the three-flavor singlet state \(\eta'\) is too heavy to be considered as a ninth (pseudo)Goldstone boson, which would be required if the pattern of spontaneous symmetry breaking were \(U(3)_L \times U(3)_R \rightarrow U(3)_V\) \[8\]. While the Gell-Mann - Oakes - Renner relations (GOR) \[9\] are very successful for the octet states, \(\pi, K, \eta\), and the masses of these particles do vanish with the bare quark masses, this is apparently not the case for the flavor-singlet state \(\eta'\) in which case the GOR fails.

't Hooft has suggested a very elegant solution for this problem \[10\]. He realized that the coupling of quarks to the localized topological solutions of the pure gluonic field (instantons) \[11\] via the axial anomaly produces an effective interaction between quarks which is still \(SU(N_f)_L \times SU(N_f)_R\) invariant but breaks explicitly the \(U(1)_A\) symmetry and is repulsive in the flavor-singlet pseudoscalar channel. Hence the instantons break the \(U(N_f)_L \times U(N_f)_R\) symmetry of the QCD Lagrangian in the chiral limit to the lower symmetry \(SU(N_f)_L \times SU(N_f)_R \times U(1)_V\) and thus the lowest flavor-singlet pseudoscalar meson cannot be considered any longer a (pseudo)Goldstone boson.

Later on Witten \[12\] and Veneziano \[13\] argued that the solution of the \(U(1)_A\) problem should not necessarily come from instantons but rather from some other type of gluonic configurations in QCD (e.g. the ones related to confinement), which also couple to the flavor-singlet quark state via the axial anomaly. This resulted in the famous Witten-Veneziano formula which relates the flavor-singlet mass to the topological susceptibility in pure gauge theory and also shows that in the large \(N_c\) limit, where the axial anomaly vanishes, the
$U(N_f)_L \times U(N_f)_R$ symmetry is restored and the flavor-singlet pseudoscalar state becomes the (pseudo)Goldstone boson of the spontaneously broken $U(N_f)_L \times U(N_f)_R$ symmetry like the octet states.

A related problem may be the origin of chiral symmetry breaking in QCD. At approximately the same time Caldi as well as Callan, Dashen and Gross and others [14] have suggested that instantons under some conditions could also provide the spontaneous chiral symmetry breaking in the QCD vacuum. The mechanism of chiral symmetry breaking is rather obvious – the ’t Hooft interaction contains the Nambu and Jona-Lasinio [15] 4-fermion vertex and hence can trigger the chiral symmetry breaking. The phenomenological arguments for the instanton liquid structure of the vacuum as well as the typical size and separation between instantons have been given by Shuryak [16]. Diakonov and Petrov have derived a microscopical theory of chiral symmetry breaking by instantons [17]. Since then much work has been done by different groups and the present state of the art is summarized in review [18]. It should be stressed that there are both strong theoretical as well as lattice arguments that the instanton medium alone cannot generate the confinement mechanism in QCD and hence some additional gluonic configurations are required in the QCD vacuum in order to explain confinement, chiral symmetry breaking, and $U(1)_A$ symmetry breaking, if instantons are indeed important for the $U(1)_A$ and chiral symmetry breaking. So it is evident that the instanton liquid structure of the QCD vacuum cannot be a complete picture of the vacuum. Indeed the instantons provide only a very small part of the full action of QCD.

These issues, and in particular, to which extent instantons contribute (or not) to the structure of the QCD vacuum, to chiral symmetry breaking, and to the structure of the low-lying hadrons, are vividly discussed topics in the lattice community. Starting first with the technique of cooling of the gluonic part of the action [19], these studies have moved to the point to which extent fermions see the instantons (via the would-be-zero modes). There is still a controversy as one of the lattice groups argues against the instantons [20]. However, their results have been questioned and cross-checked by the other groups, which see evidence in favour of instantons [21]. For example, the results [22] show in particular a direct correlation of the would-be-zero modes of fermions and the self-dual or anti-self-dual lumps of gluonic field (see the second paper in the reference above). While these studies do evidence that instantons are very important for chiral symmetry breaking, it is still unclear to which extent they can explain it, whether it is the effect of instantons only or a combined effect of instantons and something else.

We want to present an empirical argument in favour of the point of view that it is instantons that are responsible for $U(1)_A$ breaking in QCD. The experimental data suggest that high in the hadron spectrum both the $SU(2)_L \times SU(2)_R$ and $U(1)_A$ symmetries are approximately restored. Since the physics of the highly-excited states is most probably due to the confinement in QCD, one can conclude (at least preliminary) that it is not a confining interaction in QCD which is responsible for both the $SU(2)_L \times SU(2)_R$ and $U(1)_A$ breakings. While it suggests that instantons are indeed important for the $U(1)_A$ breaking, these data cannot shed any light on whether only instantons or instantons and something else provide
the chiral symmetry breaking in the QCD vacuum.

II. EMPIRICAL PATTERN OF THE $SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A$ SYMMETRY BREAKING

In this Letter we limit ourselves to the two-flavor version of QCD. There are two reasons for doing this. First of all, the $u$ and $d$ quark masses are very small as compared to $\Lambda_{QCD}$ and the typical hadronic scale of 1 GeV. Thus the chiral $SU(2)_L \times SU(2)_R$ and more generally the $U(2)_L \times U(2)_R$ symmetries of the QCD Lagrangian are nearly perfect. This is not the case if the $s$ quark is included, and a-priori it is not clear whether one should regard this quark as light or "heavy". The second reason is a practical one – there are good new data on highly excited $u, d$ mesons, but such data are still missing for the strange mesons. Certainly it would be very interesting and important to extend the analysis of the present paper to the $U(3)_L \times U(3)_R$ case. We hope that the present results will stimulate the experimental and theoretical work in this direction.

The lowest pseudoscalar and scalar mesons give an idea of how strongly the $SU(2)_L \times SU(2)_R$ and $U(1)_V \times U(1)_A$ symmetries are broken. In QCD the meson masses are extracted from the two-point correlator,

$$\langle 0|T \{j_\alpha(x)j_\alpha(0)\}|0\rangle,$$

where $\alpha$ specifies the set of quantum numbers of the current (interpolating field), $j_\alpha(x)$; it coincides with the set of quantum numbers of the meson of interest. All the information about the hadron spectrum is encoded in the complicated structure of the QCD vacuum and the physical hadrons with the quantum numbers $\alpha$ represent a response of the vacuum to the external probe $j_\alpha(x)$. For the pseudoscalar and scalar mesons $\pi, f_0, a_0$ and $\bar{\eta}$ the interpolating fields are given as

$$j_\pi(x) = \bar{q}(x)\frac{i}{2}\gamma_5q(x), \quad (2)$$

$$j_{f_0}(x) = \frac{1}{2}\bar{q}(x)q(x), \quad (3)$$

$$j_{\bar{\eta}}(x) = \frac{1}{2}\bar{q}(x)i\gamma_5q(x), \quad (4)$$

1 The $\bar{\eta}$ represents the singlet state in two-flavor QCD which is analogous to the flavor-singlet state $\eta'$ in three-flavor QCD; its mass can be approximately extracted from the masses of physical $\eta$ and $\eta'$ mesons by unmixing the $(u\bar{u} + d\bar{d})/\sqrt{2}$ and $s\bar{s}$ components - see Appendix.
\[ j_{a_0}(x) = \bar{q}(x) \frac{\tau^2}{2} q(x). \] (5)

These four currents belong to the irreducible representation of the \( U(2)_L \times U(2)_R = SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A \) group. It is instructive to see how these currents transform under different subgroups of the group above. The irreducible representations of \( SU(2)_L \times SU(2)_R \) can be labeled as \((I_L, I_R)\) with \( I_L \) and \( I_R \) being the isospins of the left and right subgroups. However, generally the states that belong to the given irreducible representation of the chiral group cannot be ascribed a definite parity because under parity transformation the left-handed quarks transform into the right-handed ones (and vice versa). Therefore under a parity operation the irreducible representation \((I_L, I_R)\) transforms into \((I_R, I_L)\). Hence, in general, the state (or current) of definite parity can be constructed as a direct sum of two irreducible representations \((I_L, I_R) \oplus (I_R, I_L)\), which is an irreducible representation of the parity-chiral group \([3]\).

The \( SU(2)_L \times SU(2)_R \) transformations consist of vectorial and axial transformations in the isospin space. The axial transformations mix the currents of opposite parity:

\[ j_\pi(x) \leftrightarrow j_{f_0}(x) \] (6)
as well as

\[ j_{a_0}(x) \leftrightarrow j_{\bar{\eta}}(x). \] (7)

Each pair of currents belongs to the \((1/2, 1/2)\) representation of the parity-chiral group, which contains both \( I = 0 \) as well as \( I = 1 \) states.

The \( U(1)_A \) transformation mixes the currents of the same isospin but opposite parity:

\[ j_\pi(x) \leftrightarrow j_{a_0}(x) \] (8)
as well as

\[ j_{f_0}(x) \leftrightarrow j_{\bar{\eta}}(x). \] (9)

All four currents together belong to the irreducible representation \((1/2, 1/2) \oplus (1/2, 1/2)\) of the \( U(2)_L \times U(2)_R \) group.

If the vacuum were invariant with respect to \( U(2)_L \times U(2)_R \) transformations, then all four mesons, \( \pi, f_0, a_0 \) and \( \bar{\eta} \) would be degenerate (as well as all their excited states). Once the \( U(1)_A \) symmetry is broken explicitly through the axial anomaly, but the chiral \( SU(2)_L \times SU(2)_R \) symmetry is still intact in the vacuum, then the spectrum would consist of degenerate \((\pi, f_0)\) and \((a_0, \bar{\eta})\) pairs. If in addition the chiral \( SU(2)_L \times SU(2)_R \) symmetry is spontaneously broken in the vacuum, the degeneracy is also lifted in the pairs above and the pion becomes a (pseudo)Goldstone boson. Indeed, the masses of the lowest mesons are \([23]\)

\[ m_\pi \simeq 140 \text{MeV}, \quad m_{f_0} \simeq 400 - 1200 \text{MeV}, \quad m_{a_0} \simeq 985 \text{MeV}, \quad m_{\bar{\eta}} \simeq 782 \text{MeV}. \]
This immediately tells that both $SU(2)_L \times SU(2)_R$ and $U(1)_V \times U(1)_A$ are broken in the QCD vacuum to $SU(2)_I$ and $U(1)_V$, respectively.

There are a few possible scenarios that are consistent with the given pattern.
(i) both the $SU(2)_L \times SU(2)_R$ and $U(1)_V \times U(1)_A$ breakings come from the particular gluodynamics in QCD that is responsible for confinement;
(ii) the $SU(2)_L \times SU(2)_R$ breaking is due to the gluodynamics that is responsible for confinement and the $U(1)_V \times U(1)_A$ breaking comes from instantons;
(iii) the $U(1)_V \times U(1)_A$ breaking is due to the confinement while the $SU(2)_L \times SU(2)_R$ breaking is from other sources;
(iv) the $U(1)_V \times U(1)_A$ breaking is provided by instantons while the $SU(2)_L \times SU(2)_R$ breaking is related to instantons alone or to a combination of instantons and some other possible gluonic interactions that are not related directly to confinement.

In the following we will show that the scenarios (i), (ii) and (iii) are very unlikely in view of the new empirical data on highly excited mesons.

### III. WHAT DO WE LEARN FROM THE HIGHLY EXCITED HADRONS?

Systematic data on highly excited mesons are still missing in the PDG tables. We will use the recent results of the partial wave analysis of mesonic resonances obtained in $p\bar{p}$ annihilation at LEAR \[5,6\]. For the scalar and pseudoscalar mesons in the mass range from 1.8 GeV to 2.4 GeV the corresponding results are summarized in the Table below. We note that the $f_0$ state at $2102 \pm 13$ MeV is not considered by the authors as a $q\bar{q}$ state (but rather as a candidate for glueball) because of its very unusual decay properties and very large mixing angle. This is in contrast to all other $f_0$ mesons in the Table, for which the mixing angles are small. Therefore these mesons are regarded as predominantly $u, d = n$ states. Hence, in the following we will exclude the $f_0$ state at $2102 \pm 13$ from our analysis which applies only in the chiral symmetry broken regime the use of effective degrees of freedom in the low-lying hadrons is certainly fruitful. The chiral symmetry breaking implies that practically massless quarks acquire a quasiparticle (dynamical or constituent) mass through their coupling to the quark condensates of the vacuum. How it happens is well seen from the schematical Nambu and Jona-Lasinio model. Pions are also well understood from the Nambu and Jona-Lasinio picture of chiral symmetry breaking and the formation of the lowest excitation over the vacuum is analogous to the Anderson mode in superconductors. In this picture the pion is a relativistic bound state of two quasi-particles $Q\bar{Q}$. The quasiparticle $Q$ itself is a result of chiral symmetry breaking in the vacuum. The ”residual” attraction of these quasiparticles in the isovector-pseudoscalar channel is unambiguously fixed by chiral symmetry and once it is taken into account within the Bethe-Salpeter approach it necessarily leads to the zero mass of pions in the chiral limit. The pion is a highly collective mode, but not a simple $q\bar{q}$ excitation, because the quasiparticle $Q$ itself is a highly collective coherent excitation of bare quarks and antiquarks.

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2 In the chiral symmetry broken regime the use of effective degrees of freedom in the low-lying hadrons is certainly fruitful. The chiral symmetry breaking implies that practically massless quarks acquire a quasiparticle (dynamical or constituent) mass through their coupling to the quark condensates of the vacuum. How it happens is well seen from the schematical Nambu and Jona-Lasinio model. Pions are also well understood from the Nambu and Jona-Lasinio picture of chiral symmetry breaking and the formation of the lowest excitation over the vacuum is analogous to the Anderson mode in superconductors. In this picture the pion is a relativistic bound state of two quasi-particles $Q\bar{Q}$. The quasiparticle $Q$ itself is a result of chiral symmetry breaking in the vacuum. The ”residual” attraction of these quasiparticles in the isovector-pseudoscalar channel is unambiguously fixed by chiral symmetry and once it is taken into account within the Bethe-Salpeter approach it necessarily leads to the zero mass of pions in the chiral limit. The pion is a highly collective mode, but not a simple $q\bar{q}$ excitation, because the quasiparticle $Q$ itself is a highly collective coherent excitation of bare quarks and antiquarks.
to $n\bar{n}$ states.

| Meson | $I$ | $J^P$ | Mass (MeV) | Width (MeV) | Reference |
|-------|-----|-------|------------|-------------|-----------|
| $f_0$ | 0   | 0+    | 1770 ± 12  | 220 ± 40    |           |
| $f_0$ | 0   | 0+    | 2040 ± 38  | 405 ± 40    |           |
| $f_0$ | 0   | 0+    | 2102 ± 13  | 211 ± 29    |           |
| $f_0$ | 0   | 0+    | 2337 ± 14  | 217 ± 33    |           |
| $\eta$ | 0   | 0−    | 2010$^{+35}_{−60}$ | 270 ± 60 |           |
| $\eta$ | 0   | 0−    | 2285 ± 20  | 325 ± 30    |           |
| $\pi$ | 1   | 0−    | 1801 ± 13  | 210 ± 15    |           |
| $\pi$ | 1   | 0−    | 2070 ± 35  | 310$^{+100}_{−50}$ |           |
| $\pi$ | 1   | 0−    | 2360 ± 25  | 300$^{+100}_{−50}$ |           |
| $a_0$ | 1   | 0+    | 2025±?     | 320±?       |           |

The prominent feature of the data is an approximate degeneracy of the three highest states in the pion spectrum with the three highest states in the $f_0$ spectrum:

$$\pi(1801 \pm 13) - f_0(1770 \pm 12), \quad (10)$$

$$\pi(2070 \pm 35) - f_0(2040 \pm 38), \quad (11)$$

$$\pi(2360 \pm 25) - f_0(2337 \pm 14). \quad (12)$$

This can be considered as a manifestation of chiral symmetry restoration high in the spectra. The approximate degeneracy of these physical states indicates that the chiral $SU(2)_L \times SU(2)_R$ transformation properties of the corresponding currents (see section 2) are not violated by the vacuum. This means that the chiral symmetry breaking of the vacuum becomes irrelevant for the high-lying states and the physical states above form approximately the chiral pairs in the $(1/2, 1/2)$ representation of the chiral group. The physics of the highly excited hadrons is such as if there were no chiral symmetry breaking in the vacuum.

A similar behaviour is observed from a comparison of the $a_0$ and $\eta$ masses high in the spectra:

$$a_0(2025±?) - \eta(2010^{+35}_{−60}). \quad (13)$$

The authors of ref. 6 are confident of the existence of $a_0(2025)$, however it is difficult to extract the error bars for its mass from the existing data. Some of the missing states with these quantum numbers are still to be discovered; technically the identification of the $a_0$ and $\eta$ resonances is a rather difficult task.

As was stressed before 23, the chiral symmetry restoration high in hadron spectra does not mean that the chiral symmetry breaking in the QCD vacuum disappears, but rather
that the chiral asymmetry of the vacuum becomes irrelevant once we are sufficiently high in the spectra. While the quark condensates of the QCD vacuum are crucially important for the physics of low-lying states and ”remove” the axial part of the chiral symmetry, thereby preventing a parity doubling low in the hadron spectra, their role high in the spectrum becomes progressively less important and eventually the chiral symmetry is restored.

It is quite natural to assume that the physics of the highly excited hadrons is due to confinement in QCD. If so, it follows that the confining gluodynamics is still important. On the other hand the chiral symmetry breaking effects in the vacuum become irrelevant. Then the scenarios (i) and (ii) are ruled out.

A very natural physical picture for the highly excited states is that these hadrons are relativistic strings (with the color-electric field in the string) with practically massless quarks at the ends; these massless quarks are combined into parity-chiral multiplets \([4]\). The string picture is compatible with the chiral symmetry restoration because there always exists a solution for the right-handed and left-handed quarks at the end of the string with exactly the same energy and total angular momentum. Since the nonperturbative field in the string is pure electric and the electric field is ”flavor-blind”, the string dynamics itself is not sensitive to the specific flavor of a light quark once the chiral limit is taken. This picture explains the empirical parity-doubling because for every intrinsic quantum state of the string there necessarily appears parity doubling of the states with the same total angular momentum of hadron. Hence the string picture is compatible not only with the \(SU(2)_L \times SU(2)_R\) restoration, but more generally with the \(U(2)_L \times U(2)_R\) one.

This picture should be contrasted with the nonrelativistic or (semi)relativistic potential description of hadrons. Within the potential description the parity of the state is unambiguously prescribed by the relative orbital angular momentum \(L\) of quarks. For example, all the states on the radial pion Regge trajectory are \(^1S_0\) \(q\bar{q}\) states, while the members of the \(f_0\) trajectory are \(^3P_0\) states. Clearly, such a picture cannot explain the systematical parity doubling as it would require that the stronger centrifugal repulsion in the case of \(^3P_0\) mesons (as compared to the \(^1S_0\) ones) as well as the strong and attractive spin-spin force in the case of \(^1S_0\) states (as compared to the weak spin-spin force in the \(^3P_0\) channel) must systematically lead to an approximate degeneracy for all radial states. This is very improbable. The potential picture also implies strong spin-orbit interactions between quarks while the spin-orbit splittings are absent or very small for excited mesons and baryons in the \(u, d\) sector. The strong spin-orbit interactions inevitably follow from the Thomas precession (once the confinement is described through a scalar confining potential)\(^3\), and this very strong spin-orbit force must be practically exactly compensated by other strong spin-orbit force from the one-gluon-exchange interaction in this picture. In principle such a cancellation could be provided by tuning the parameters for some specific (sub)families of mesons. However, in

\(^3\) Note also that a scalar potential explicitly breaks the chiral symmetry in contradiction to the requirement that the chiral symmetry must be restored high in the spectra.
this case the spin-orbit forces become very strong for other (sub)families. In contrast, in the string picture there is no spin-orbit force at all once the chiral symmetry is restored [4]. That the potential description fails high in the spectra also follows from a comparison of the prediction of, e.g., ref. [24] with the recent experimental data: the potential picture simply does not predict very many states in the region of 2 GeV. For example, while the tuning of parameters of the model provides an accurate description of the three lowest states in the pion spectrum, it does not predict at all the existence of \( \pi(2070) \) and \( \pi(2360) \); the forth and the fifth radial states of the pion do not appear in this picture up to 2.4 GeV (which means that they are predicted to be at least \( \sim 0.5 \) GeV heavier than in reality). A similar situation occurs also in other channels. The failure of the potential description is inherently related to the fact that it cannot incorporate chiral symmetry restoration high in the spectra.

The nonrelativistic or (semi)relativistic potential picture is justified, however, once the current quarks are heavy and move slowly (e.g. like in charmonium and bottomonium) or the (semi)relativistic description can be still justified to some extent once the proper effective degree of freedom is a rather heavy quasiparticle but not a bare quark (constituent quark in the low-lying nucleons and deltas is not yet ultrarelativistic). Here the situation is similar to atomic physics or to the physics of positronium. For heavy fermions the relativistic effects represent only small \( \beta^2/c^2 \) corrections to the nonrelativistic picture. However, once the quarks are ultrarelativistic, it is not justified at all. As a manifestation, the potential picture requires the \( f_0 \) mesons to be \( P \) states of quarks, contrary to the \( S \) states of quarks in pions. On the other hand the string picture attributes both \( \pi \) and \( f_0 \) states (in pairs) to the same intrinsic quantum state of the string with the same angular momentum [4]. The opposite parity of these high-lying mesons is provided by different right-left configurations of the quarks at the ends of the string.

Upon examining the experimental data more carefully one notices not only a degeneracy in the chiral pairs, but also an approximate degeneracy in \( U(1)_A \) pairs (\( \pi, a_0 \)) and (\( f_0, \eta \)) (in those cases where the states are established). If so, one can preliminary conclude that not only the chiral \( SU(2)_L \times SU(2)_R \) symmetry is restored high in the spectra, but the whole \( U(2)_L \times U(2)_R \) symmetry of the QCD Lagrangian. Then the approximate \( (1/2, 1/2) \oplus (1/2, 1/2) \) multiplets of this group are given by:

\[
\pi(1801 \pm 13) - f_0(1770 \pm 12) - a_0(?) - \eta(?); \quad (14)
\]

\[
\pi(2070 \pm 35) - f_0(2040 \pm 40) - a_0(2025?) - \eta(2010_{-60}^{+35}); \quad (15)
\]

\[
\pi(2360 \pm 25) - f_0(2337 \pm 14) - a_0(?) - \eta(2285 \pm 20). \quad (16)
\]

This preliminary conclusion would be strongly supported by a discovery of the missing \( a_0 \) meson in the mass region around 2.3 GeV as well as by the missing \( a_0 \) and \( \eta \) mesons in the 1.8 GeV region. This would also rule out the scenario (iii) and only the scenario (iv) would be viable.
We have to stress, that the $U(1)_A$ restoration high in the spectra does not mean that the axial anomaly of QCD vanishes, but rather that the specific gluodynamics (e.g. instantons) that are related to the anomaly become unimportant there.

It should also be emphasized that the only restoration of $U(1)_V \times U(1)_A$ symmetry (without the $SU(2)_L \times SU(2)_R$) is impossible. This was discussed in ref. [2]. The reason is that even if the effects of the explicit $U(1)_A$ symmetry breaking via the axial anomaly vanish, the $U(1)_V \times U(1)_A$ would still be spontaneously broken once the $SU(2)_L \times SU(2)_R$ were spontaneously broken. This is because the same quark condensates in the QCD vacuum that break $SU(2)_L \times SU(2)_R$ do also break $U(1)_V \times U(1)_A$.

IV. HOW DO THE INSTANTONS DO THE JOB?

The present analysis suggests that indeed instantons cause the $U(1)_A$ breaking. Then it is instructive to outline the possible scheme of how this happens and why instantons are not important high in the spectra.

The instanton-induced interaction between quarks in the two-flavor case is given as

$$H_{int} \sim -G \left\{ [\bar{q}(x)q(x)]^2 + [\bar{q}(x)i\gamma_5 q(x)]^2 - [\bar{q}(x)i\gamma_5 q(x)]^2 - [\bar{q}(x)v\gamma_5 q(x)]^2 \right\}. \quad (17)$$

Since this is a local 4-fermion vertex, the ultraviolet cut-off must be introduced to regularize the integrals. The physical interpretation of this cut-off is obvious: the instanton-induced interaction is operative only when the squared four-momenta of the quarks are small. The reason is that the effective interaction comes from the existence of the zero modes of quarks (i.e. the zero mass quark is bound by the instanton exactly with zero energy). If the three-momentum of the travelling quark is very high but its energy is small, it does not see instantons at all and the instanton-induced interaction between such quarks must vanish. The strength of the interaction $G$ as well as the ultraviolet cut-off are directly related to such parameters as an average size of instantons as well as an average separation between them [16,17]. This interaction is attractive in $f_0$ and $\pi$ channels (the first and the second terms) and repulsive in $a_0$ and $\eta$ channels (the third and the fourth terms). The repulsion in the latter channels must be contrasted with the attraction in these channels that is prescribed by perturbative gluon exchange. The repulsion in these channels explicitly breaks the $U(1)_A$. The interaction is $SU(2)_L \times SU(2)_R$ symmetric. The first two terms in eq. (17) represent the Nambu and Jona-Lasinio Hamiltonian. Hence if the interaction is strong enough it can also provide the spontaneous breaking of chiral symmetry.

Summarizing, once the ’t Hooft determinant interaction is introduced between the valence quarks in a meson, it automatically solves the $U(1)_A$ problem. If this interaction is taken between the sea quarks in the vacuum, it can provide the spontaneous breaking of chiral symmetry. For the present context it is crucially important that this interaction is a low-momentum interaction. Hence, for the low-lying hadrons, where the typical momenta of
valence quarks are not high, the nonperturbative dynamics due to instantons is important.

When we are high in the spectra, the three-momenta of valence quarks increase (contrary to their energy)\(^4\), hence the 't Hooft interaction between the valence quarks vanishes and the \(U(1)_A\) symmetry is restored. Similarly, the fast moving valence quarks do not interact via instantons (or via some other gluonic interaction that is responsible for chiral symmetry breaking) with the sea quarks in the vacuum; thus they decouple from the quark condensates of the vacuum. Consequently the chiral \(SU(2)_L \times SU(2)_R\) symmetry is also effectively restored.

The QCD vacuum is a very complicated medium and has many facets (like cubistic paintings). It contains instantons, other possible topological configurations and mostly the quantum fluctuations around them. Different probes see different facets of the vacuum. For example, if one probes the vacuum by heavy quarks, those facets that are important for the breaking of chiral as well as \(U(1)_A\) symmetries become irrelevant. The heavy quarks simply do not see them. However the aspects of the vacuum that are important for confinement are relevant in this case. Indeed, it could be possible that predominantly the stochastic structure of the QCD vacuum [23], which nicely explains the area law of the Wilson loop and also the Casimir scaling [26], does underly the physics of the heavy quarks. However, once we probe the vacuum by light quarks, in addition other facets (like instantons) become important and the physics become reacher. The light quarks do see the instantons inspite their weight in the full QCD action is very small.

As a conclusion, the present results show that high in the spectrum the chiral \(SU(2)_L \times SU(2)_R\) symmetry is restored and probably also the \(U(1)_A\) one. It then follows that it is not a confining gluodynamics in QCD that is responsible for chiral and \(U(1)_A\) breakings.

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V. APPENDIX

In this appendix we show how the unmixing of the pure \(\bar{\eta} = (u\bar{u} + d\bar{d})/\sqrt{2}\) and \(\bar{\eta}_s = s\bar{s}\) states from the physical \(\eta\) and \(\eta'\) mesons is done.

The physical \(\eta\) and \(\eta'\) mesons can be written as

\[
\eta = \left( \frac{1}{\sqrt{3}} \cos 10 + \frac{\sqrt{2}}{\sqrt{3}} \sin 10 \right) \bar{\eta} + \left( \frac{1}{\sqrt{3}} \sin 10 - \frac{\sqrt{2}}{\sqrt{3}} \cos 10 \right) \bar{\eta}_s, \]

\(\text{This is natural in the string picture where practically the whole energy of the hadron is accumulated in the string while the quarks at the ends have a large three-momentum.}\]
\[ \eta' = \left( \frac{\sqrt{2}}{\sqrt{3}} \cos 10 - \frac{1}{\sqrt{3}} \sin 10 \right) \bar{\eta} + \left( \frac{\sqrt{2}}{\sqrt{3}} \sin 10 + \frac{1}{\sqrt{3}} \cos 10 \right) \bar{\eta}_s. \]

Assuming the mass-squared mixing matrix, the masses of physical \( \eta \) and \( \eta' \) can be found from the eigenvalue problem

\[
\begin{vmatrix}
  m_{\eta}^2 - m_{\eta,\eta'}^2 & V^2 \\
  V^2 & m_{\bar{\eta}_s}^2 - m_{\eta,\eta'}^2
\end{vmatrix} = 0.
\]

This equation together with the mixing relations above allow to determine the masses \( m_{\bar{\eta}} \simeq 782 \text{ MeV} \) and \( m_{\bar{\eta}_s} \simeq 778 \text{ MeV} \).
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