I overview the physical picture of QCD at low energies that has led to the prediction of a narrow exotic baryon $\Theta^+$ which cannot be made of three quarks. The very narrow width of the $\Theta^+$ and a possible reason why it is seen in low- but not in high-energy experiments are briefly discussed.

PACS numbers: 12.38.-t, 12.39.-x, 12.39.Dc, 14.20-c

I would like to stress that the situation here is radically different from all other bound states we know. Be it atoms or nuclei, their masses are always less than the appropriate "$\sigma$-term" or the sum of their constituents’ masses. This paradox has to be explained first of all.

### II. ORIGIN OF HADRON MASSES

The most important happening in QCD from the point of view of the light hadron structure is the Spontaneous Chiral Symmetry Breaking: as its result, almost massless $u, d, s$ quarks get the dynamical momentum-dependent masses $M_{u,d,s}(p)$, and the pseudoscalar mesons $\pi, K, \eta$ become light (pseudo) Goldstone bosons.

This story starts from the gluon vacuum. In the absence of any matter and sources, i.e. in the vacuum, the gluon field experiences zero-point quantum fluctuations whose snapshot can be visualized in lattice simulations (Fig. 1, top). These are the normal quantum fluctuations, experienced also by the electromagnetic field. What is peculiar for the non-Abelian gluon field, is that beneath zero-point oscillations there are specific large fluctuations called instantons (for a recent review see Ref. [8]). One can measure their average size and the average separation between the peaks, which appear to be roughly 1/3 fm and 1 fm, respectively. These key numbers had been suggested by Shuryak [8] and obtained from lattice simulations long before lattice measurements became available.
FIG. 1: A snapshot of quantum fluctuations of the gluon field in the vacuum [12]. “Cooling” the normal zero-point oscillations (top) reveals that they lie on top of large gluon fluctuations, which are identified with instantons and anti-instantons with random positions and sizes (bottom). The left column shows the action density and the right column shows the topological charge density for the same snapshot, where instantons are peaks and anti-instantons are holes.

Now we switch in light quarks and ask how quarks propagate through this fluctuating gluon vacuum (Fig. 2). The result [11] is that the nearly massless quarks gain a large dynamical mass \( M(p) \) plotted in Fig. 3. This is in fact a beautiful and a non-trivial phenomenon. The point is, massless quarks generally conserve their helicity when interacting with gluons. Therefore, the mass is not generated in any order in perturbation theory. In other words, normal zero-point oscillations of the gluon field depicted on the top of Fig. 1 do not generate the quark mass. Only the large instanton fluctuations shown in the bottom of Fig. 1 have the power of flipping quark helicity. It is a non-perturbative effect. Quarks “hop” from one randomly situated instanton fluctuation to another, each time flipping the helicity. As a result they get a dynamical mass. It goes to zero at large momenta since quarks with very high momenta are not affected by the background, even if it is strong gluon field as in the case of instantons.

Another striking confirmation of the instanton mechanism of the spontaneous chiral symmetry breaking [11] came recently from another lattice study [14]. In Fig. 4 the ratio of the helicity-flip to non-flip correlation functions of flavor non-singlet currents is plotted. Random instantons explain very precisely the non-trivial shape of this ratio as function of the distance.

The dynamical mass plotted in Fig. 3 is key to understanding light hadrons’ properties. Its value at zero momentum,

\[
M(0) \simeq 350 \text{ MeV},
\]  

is what is usually called the “constituent” quark mass (to be distinguished from the small input “current” quark masses of the QCD lagrangian.) The half-width of \( M(p) \) corresponds roughly to the size of the constituent quarks,

\[
\rho \simeq \frac{1}{600 \text{ MeV}} = \frac{1}{3} \text{ fm},
\]  

In the instanton approach, these numbers are calculated from the average separation between and the average size
FIG. 5: The same instantons that bring in $M(p)$ are responsible for very strong binding of quarks into a nearly massless pion, which is guaranteed by the Goldstone theorem. Summation of diagrams demonstrating it: see Ref. [11].

of the instanton fluctuations which, in their turn, can be computed from the value of $\Lambda_{\overline{MS}}$. By the way, there is no real solution of the equation $p^2 = M^2(-p^2)$, meaning that quarks cannot be observable, only their bound states!

From the above numbers one can immediately estimate hadron sizes and masses.

Hadron sizes are determined, via the uncertainty principle, by the inverse $M(0)$, or are somewhat larger if constituent quark binding is loose:

$$R \geq \frac{1}{M} \sim 0.8 \text{ fm.} \quad (5)$$

It is crucial that hadron sizes are significantly larger than quark sizes, otherwise the constituent quark ideas would never have worked.

Hadron masses are roughly twice the constituent quark mass for the typical vector mesons, and thrice $M$ for baryons:

$$\text{Vector mesons: } m_V \approx 2M \approx 700 \text{ MeV,} \quad (6)$$

$$\text{3–quark baryons: } m_B \approx 3M \approx 1000 \text{ MeV.} \quad (7)$$

In pseudoscalar mesons, however, two constituent quark masses are completely ‘eaten up’ by strong interaction, which is guaranteed by the Goldstone theorem stating that once chiral symmetry of strong interactions is spontaneously broken, pseudoscalar mesons must be very light. Their masses are fully determined not by the constituent but by the current quark masses, in accordance with the Gell-Mann–Oakes–Renner formula:

$$m^2_{\pi} = (M + M - 2M)^2 + \frac{m_q \langle \bar{q}q \rangle}{F^2_{\pi}}. \quad (8)$$

The very strong interactions between quarks in the pseudoscalar channel is due to the same forces that are responsible for the originally nearly massless quarks obtaining a dynamical mass $M(p)$: they are due to instantons, see Fig. 5.

III. WHAT FORCES BIND CONSTITUENT QUARKS IN BARYONS?

The standard answer is: the linear confining potential, chromoelectric flux tubes. However, in the real world with very light pions it cannot be correct because it is energetically favorable to break an extended string, and produce pions, see Figs. 6,7.

In light baryons, it is sufficient to exceed the mass of one additional pion to make the flux tube energetically unfavorable. Pion production makes the color string break. Since pions are very light one can estimate that it happens at a very early $\sim 0.3 \text{ fm}$ separation between quarks, which is the size of the constituent quark itself! Meanwhile, the average separation of quarks inside a baryon is much bigger – typically about 0.8 fm. Therefore, color strings or color-electric flux tubes can hardly be responsible for quark binding. All what is left, are constituent quarks interacting with pions, but they themselves are made of constituent quarks.

How to present this queer situation mathematically? There is actually not much freedom here: the interaction of pseudoscalar mesons with constituent quarks is dictated by chiral symmetry. It can be written in the
from instantons in Ref. [11].

stituent quarks. In this form, eq. (9) has been derived are not 'elementary' but a composite field, made of con-

FIG. 8: Pseudoscalar mesons are themselves bound states of constituent quarks: they propagate and interact via virtual quark-antiquark pairs. The first diagram is pion's propagation, the second one describes (correctly) the s- and d-wave pion scattering at low energies, the third gives the 5-prong process \( K^+K^- \rightarrow \pi^+\pi^0\pi^- \) which is theoretically known from the Wess–Zumino term. The sum of all diagrams with any number of external legs is called the effective chiral lagrangian.

The following compact form [11]:

\[
\mathcal{L}_{\text{eff}} = \bar{q} \left[ i\slashed{\partial} - M \exp(i \gamma_5 \pi^A \lambda^A / F_{\pi}) \right] q, \quad \pi^A = \pi, K, \eta, (9)
\]

Because of the necessary exponent of the pseudoscalar field, it is an essentially non-linear interaction. The exponent here is a 3 × 3 flavor matrix acting on the 3-vector \((u, d, s)\). It is also a 4 × 4 matrix acting on the Dirac 4-spinor indices of quarks.

[A note for experts: Since eq. (9) is an effective low-energy theory, one expects formfactors in the constituent quark – pion interaction; in particular, as we know already \( M(p) \) is momentum-dependent. Eq. (9) is written in the limit of zero momenta. A possible wave-function renormalization factor \( Z(p) \) can be also admitted but it can be absorbed into the definition of the quark field. The low-energy theory similar to eq. (9) has been suggested in 1984 by several groups [17, 18, 19]. However, in these influential papers an additional kinetic energy term for pions has been added to eq. (9). It is neither necessary nor even seem to be correct: the kinetic energy term (and higher derivatives) for pions appears from integrating out quarks, or, in other words, from quark loops, see Fig. 8. It is in accordance with the fact that pions are not 'elementary' but a composite field, made of constituent quarks. In this form, eq. (9) has been derived from instantons in Ref. [11].

Constituent \( u, d, s \) quarks necessarily have to interact with the \( \pi, K, \eta \) fields according to eq. (9), and the dimensionless coupling constant is actually very large:

\[
g_{\pi qq}(0) = \frac{M(0)}{F_\pi} \simeq 4. \quad (10)
\]

Alternatively, one can estimate this coupling as the pion-nucleon constant \( g_{\pi NN} \simeq 13.3 \) divided by three quarks, which gives approximately the same.

IV. WHAT IS THE NUCLEON?

The chiral interactions of constituent quarks in baryons, following from eq. (9), are schematically shown

in Fig. 9. Antiquarks are necessarily present in the nucleon as pions propagate through quark loops. The non-linear effects in the pion field are essential since the coupling is strong. I would like to stress that this picture is a model-independent consequence of the spontaneous chiral symmetry breaking. One cannot say that quarks get a constituent mass but throw away their strong interaction with the pion field. In principle, one has to add perturbative gluon exchange on top of Fig. 9. However, \( \alpha_s \) is never really strong, such that gluon exchange can be disregarded in the first approximation. The large value of the pion-quark coupling (10) suggests that Fig. 9 may well represent the most essential forces inside baryons.

The low-momenta effective theory defined by eq. (9) is a big step forward, as compared to the original formulation of QCD: it operates with the adequate degrees of freedom, namely the dynamically-massive quarks and the light pseudoscalar meson fields, relevant at low energies. The transition to these new degrees of freedom is similar to the transition from QED (= the microscopic theory of the atoms) to the electrons in a material, whose mass is not the original 0.511 MeV but a heavier effective one, and whose most important interaction at the atomic “low energies” is not the Coulomb (read: gluon) but rather the phonon (read: pion) exchange. Phonons are collective excitations of atomic lattices, and they are Goldstone bosons associated with the spontaneous breaking of the translational symmetry by the lattice. Pions are collective excitations of the QCD ground state, and they are Goldstone bosons associated with the spontaneous breaking of chiral symmetry. Their interaction with quarks which obtained the dynamical mass also owing to the same spontaneous chiral symmetry breaking, is dictated by symmetry. Only the parameters of the low-energy effective action (9) have to be determined from a microscopic theory – in this case instantons seem to do this job very well.

Continuing the analogy, the Cooper pairing of electrons in a superconductor is due to phonon exchange which is much stronger than the Coulomb force between electrons (being actually a repulsion). As we shall see in a moment, the binding of quarks into a nucleon can be explained as due to their interaction with pions. Although the low-momenta effective theory (9) is a great simplification as compared to the microscopic QCD, it is still a strong-coupling relativistic quantum field theory. Summing up all interactions inside the nucleon of the kind shown in Fig. 9 is a difficult task. Nevertheless, it can be
solved exactly in the limit of large number of colors \( N_c \). It is widely believed \([20, 21]\) that the real world with \( N_c = 3 \) is not qualitatively different from the imaginary world at \( N_c \to \infty \). For \( N_c \) colors the number of constituent quarks in a baryon is \( N_c \), and all quark loop contributions in Fig. 9 are also proportional to \( N_c \). Therefore at large \( N_c \), quarks inside the nucleon produce a large, nearly classical pion field: quantum fluctuations about the mean field are suppressed as \( 1/N_c \). The same field binds the quarks; therefore it is called the self-consistent field. [A similar idea is exploited in the shell model for nuclei.] The problem of summing up all diagrams of the type shown in Fig. 9 is thus reduced to finding a classical self-consistent pion field. As long as \( 1/N_c \) corrections to the mean field results are under control, one can use the large-\( N_c \) logic and put \( N_c \) to its real-world value 3 at the end of the calculations.

The model of baryons based on these approximations has been named the Chiral Quark Soliton Model (CQSM) \([22]\). The “soliton” is another word for the self-consistent pion field in the nucleon. However, the model operates with explicit quark degrees of freedom, which enables one to compute any type of observables, e.g. relativistic quark (and antiquark!) distributions inside nucleons \([23]\), and the quark light-cone wave functions \([24]\). In contrast to the naive quark models, the CQSM is relativistic-invariant. Being such, it necessarily incorporates quark-antiquark admixtures to the nucleon. Quark-antiquark pairs appear in the nucleon on top of the three valence quarks either as particle-hole excitations of the Dirac sea (read: mesons) or as collective excitations of the mean chiral field.

It should be stressed thus that there is nothing odd or dramatic in calling the nucleon a chiral soliton. It is no more soliton than an atom which, at large \( Z \), can be well described by the Thomas–Fermi method where one considers \( Z \) electrons in the self-consistent field, in that case it is the electrostatic field. One can view an atom as a soliton of the electrostatic field, if one likes. Fortunately for the chiral soliton, corrections to the model go as \( 1/N_c \) or even as \( 1/N_c^2 \) and can be computed for many observables. In the Thomas–Fermi model of atoms corrections to the mean field are of the order of \( 1/\sqrt{Z} \) and for that reason are large unless atoms are heavy.

There are two instructive limiting cases in the CQSM:

- **Weak \( \pi(x) \) field.** In this case the Dirac sea is weakly distorted as compared to the no-field and thus carries small energy, \( E_{\text{sea}} \approx 0 \). Few antiquarks. The valence-quark level is shallow and hence the three valence quarks are non-relativistic. In this limit the CQSM becomes very similar to the constituent quark model remaining, however, relativistic-invariant and well defined.

- **Large \( \pi(x) \) field.** In this case the bound-state level with valence quarks is so deep that it joins the Dirac sea. The whole nucleon mass is given by \( E_{\text{sea}} \) which in its turn can be expanded in the derivatives of the mean field, the first terms being close to the Skyrme lagrangian. Therefore, in the limit of large and broad pion field, the model formally reduces to the Skyrme model.

The truth is in between these two limiting cases. The self-consistent pion field in the nucleon turns out to be strong enough to produce a deep relativistic bound state for valence quarks and a sufficient number of antiquarks, so that the departure from the non-relativistic constituent quark model is considerable. At the same time the self-consistent pion field is spatially not broad enough to justify the use of the Skyrme model which is just a crude approximation to the reality, although shares with reality some qualitative features. The CQSM demystifies the main paradox of the Skyrme model: how can one make a fermion out of a boson-field soliton. Since the “soliton” is nothing but the self-consistent pion field that binds quarks, the baryon and fermion number of the whole construction is equal to the number of quarks one puts on the valence level created by that field; it is three in the real world with three colors. See Ref. \([25]\) for a review.
resolution 1 fm
momentum transfer
q < 300 MeV

resolution 1/3 fm
300 < q < 1000 MeV

q > 1 GeV
perturbative cascade

FIG. 12: Nucleons under a microscope with increasing resolution.

V. NUCLEONS UNDER A MICROSCOPE WITH INCREASING RESOLUTION

Inelastic scattering of electrons off nucleons is a microscope with which we look into its interior. The higher the momentum transfer \( Q \), the better is the resolution of this microscope, see Fig. 12.

At \( Q < 300 \text{ MeV} \) one does not actually discern the internal structure; it is the domain of nuclear physics. At \( 300 < Q < 1000 \text{ MeV} \) we see three constituent quarks inside the nucleon, but also additional quark-antiquark pairs; mathematically, they come out from the distor- tion of the Dirac sea in Fig. 10. The appropriate quark and antiquark distributions have been found in Ref. [22]. In addition, the non-perturbative gluon distribution appears for the first time at this resolution. First and foremost, it is the glue in the interior of the constituent quarks that has been responsible for rendering them the mass, i.e. glue from the instanton fluctuations. It is amusing that these non-perturbative gluons are emitted not by the vector (chromoelectric) quark current but rather by the quarks’ large chromomagnetic moment, and their distribution has been found by Maxim Polyakov and myself to be given by a universal function \((1-x)/x\), see section 7 in Ref. [8].

At large \( Q > 1 \text{ GeV} \) one gets deep inside constituent quarks and starts to see normal perturbative gluons and more quark-antiquark pairs arising from bremsstrahlung. This part of the story is well-known: the perturbative evolution of the parton cascade gives rise to a small violation of the Bjorken scaling as one goes from moderate to very large momentum transfers \( Q \), but the basic shape of parton distributions serving as the initial condition for perturbative evolution, is determined at moderate \( Q \) by non-perturbative physics described above.

What is computed in the Chiral Quark Soliton Model

- ‘Static’ characteristics of baryons: masses, magnetic moments, formfactors, coupling constants, \( g_A \), etc., see Ref. [20]. Important: the large value of the nucleon \( \sigma \)-term explained [22], as due to the additional \( QQ \) pairs in the nucleon to which the \( \sigma \)-term is particularly sensitive.
- Parton distributions at low virtuality, automatically satisfying the known QCD sum rules [23]. Explained: a relatively large fraction of antiquarks, large flavor asymmetry of the sea \( u(x) - d(x) \) [28], small fraction of the nucleon spin carried by the quarks’ spin [20].
- Distribution amplitudes or light-cone quark wave functions in \( N \) and \( \Delta \). One recovers approximately the old non-relativistic \( SU(6) \) wave function but with relativistic corrections and the additional \( QQ \) pairs. It is important for understanding hard exclusive reactions [30].
- Prediction: large flavor asymmetry in the polarized antiquark parton distributions \( \Delta \bar{u}(x) - \Delta d(x) \) [31].
- Prediction: peculiar oscillatory behavior of generalized parton distributions [32].

VI. BARYON EXCITATIONS

There are excitations related to the fluctuations of the chiral field about its mean value in the baryons. In the context of the Skyrme model many resonances were found and identified with the existing ones in Ref. [33, 34] and quite recently in Ref. [35]. As I said before, the Skyrme model is too crude, and one expects only qualitative agreement with the Particle Data. The same work has to be repeated in the CQSM but it has not been done so far.

There are also low-lying collective excitations related to slow rotation of the self-consistent chiral field as a whole both in ordinary and flavor spaces. The result of the quantization of such rotations was first given by Witten [36]. The following \( SU(3) \) multiplets arise as rotational states of a chiral soliton: \( 8, \frac{1}{2}^{+} \), \( (10, \frac{1}{2}^{+}) \), \( (10, \frac{3}{2}^{-}) \), \( (27, \frac{1}{2}^{+}) \), \( (27, \frac{3}{2}^{-}) \)...
They are ordered by increasing mass, see Fig. 13. The first two (the octet and the decuplet) are indeed the lowest baryons states in nature. They are also the only two that can be composed of three quarks, according to the quantum numbers. However, the fact that one can manage to obtain the correct quantum numbers of the
octet and the decuplet combining only three quarks, does not necessarily mean that they are made of three quarks only. The difficulties of such an interpretation have been mentioned in the beginning. In fact we know for thirty years that the ‘three-quark’ view is simplistic: even baryons from the lowest $\left(8, \frac{1}{2}^+\right)$, $\left(10, \frac{3}{2}^+\right)$ multiplets have always an admixture of $Q\bar{Q}$ pairs, see Fig. 12.

Therefore, one should not be a priori confused by the fact that higher-lying multiplets cannot be made of three quarks: even the lowest ones are not. A more important question is where to stop in this list of multiplets. Apparently for sufficiently high rotational states the rotations become too fast: the centrifugal forces will rip the baryon apart. Also the radiation of pions and kaons by a fast-rotating body is so strong that the widths of the corresponding resonances blow up \[37\]. Which precisely rotational excitation is the last to be observed in nature, is a quantitative question. Actually one needs to compute their widths in order to make a judgement. If the width turns out to be in the hundreds of MeV, one can say that this is where the rotational sequence ceases to exist.

An estimate of the width of the lightest member of the antidecuplet, shown at the top of the right diagram in Fig. 13, the $\Theta^+$, gave a surprisingly small result: $\Gamma_\Theta < 15$ MeV \[38\]. This result obtained in the CQSM, immediately gave credibility to the existence of the antidecuplet. It should be noted that there is no way to obtain a small width in the oversimplified Skyrme model.

Let us recall our bookkeeping for baryon masses.

Baryon masses again ($M_Q \approx 350$ MeV)

- mainly 3–quark baryons: $\mathcal{M} \approx 3M$ (+ strangeness), $5$–quark baryons, naively: $\mathcal{M} \approx 5M$ (+ strangeness)
  $\approx 1900$ MeV for $\Theta^+$ ??

- $5$–quark baryons, correct: $\mathcal{M} \approx 3M + \frac{1}{\text{baryon size}}$ (+ strangeness) $\approx 1550$ MeV for $\Theta^+$.

In pentaquarks forming the antidecuplet shown on the right of Fig. 13, the additional $Q\bar{Q}$ pair is added in the form of the excitation of the (nearly massless) chiral field. Energy penalty would be zero, had not the chiral field been restricted to the baryon volume. This is clearly seen from the formulae. The splitting between the average masses of octet, decuplet and antidecuplet baryons is given by the two moments of inertia $I_{1,2}$:

$$M_{\left(10, \frac{3}{2}\right)} - M_{\left(8, \frac{1}{2}\right)} = \frac{3}{2} \frac{1}{I_1},$$

$$M_{\left(1\pi\pi, \frac{3}{2}\right)} - M_{\left(8, \frac{1}{2}\right)} = \frac{3}{2} \frac{1}{I_2}.$$

The first, $I_1$, is the moment of inertia with respect to the usual rotations of a baryon; it coincides with the moment of inertia with respect to rotations of the chiral field in the isospin space. Numerically, it turns out to be about (1 fm)$^{-1}$, a very reasonable value for a baryon. The second, $I_2$, is the moment of inertia with respect to rotations mixing ‘strange’ directions of the flavor space. [In the chiral limit when the current mass $m_s$ is neglected such rotations are required by the $SU(3)$ symmetry, and are as good as the isospin rotations.] Numerically, it turns out to be about twice less than $I_1$, leading to approximately twice larger splitting of the antidecuplet from the ground-state octet than the decuplet-octet splitting.

Important, the splitting \[12\] is not twice the constituent mass $2M$ but less. One can imagine a large-size baryon: in such a case the moment of inertia, roughly $I \sim m < r^2>$, blows up, meaning that it costs little energy to excite the antidecuplet!

**VII. PREDICTION OF THE $\Theta^+$**

In 1997 Victor Petrov, Maxim Polyakov and I summarized what we knew about the antidecuplet. There were two striking features: it had to be a) relatively light and b) surprisingly narrow for strongly decaying resonances with masses above 1.5 GeV.

On the theoretical side, there were predictions for pentaquarks in the constituent quark models \[39\] with the lightest one appearing in the range $1700 - 1900$ MeV. These models predicted negative-parity baryons to be lighter than positive-parity ones, meaning that they could decay in the s-wave and had to be extremely broad. In contrast to those ad hoc models, the Skyrme model correctly emphasized the importance of the chiral degrees of freedom in baryons. It therefore predicted more light pentaquarks, the lightest being in the $1400$ MeV range \[40\]. However, the Skyrme model was just a qualitative approximation to the reality, and one faced a difficult task of getting the correct observables for ordinary baryons, not to mention the exotic ones. To the best of my knowledge, none of the authors attempted to estimate the widths of the antidecuplet baryons, which was the crucial thing (had such an attempt been made,
it would have yielded a factor of 5-15 larger width than in the CQSM). Since there were too many inherent inconsistencies inside the models considered, none of the authors seemed to have taken the antidecuplet for real.

On the experimental side, from 1960’s till early 80’s there have been intensive searches for the exotic $S = +1$ baryons (called $Z$ baryons in those days) with no convincing results. There have been several one- and two-star resonances summarized by the Particle Data Group in the 1986 edition, all in the 1700 – 1900 MeV mass range. After 1986 the Particle Data Group ceased to mention those resonances as insufficiently convincing. The 1992 partial wave analysis of the $KN$ scattering data [11] concluded that there might be broad resonances but, if any, they ought to be in a high-mass range, see Fig. 14.

Despite a heavy pressure from the unsuccessful attempts to find exotic baryons in the past and despite a general trend to dogmatize the “three quarks only” view on baryons, we decided to publish our findings:

**Exotic Anti-Decuplet of Baryons: Prediction from Chiral Solitons** [Z. Phys. A 359, 305 (1997)]

**Abstract:** We predict an exotic baryon (having spin 1/2, isospin 0 and strangeness +1) with a relatively low mass of about 1530 MeV and total width of less than 15 MeV. It seems that this region of masses has avoided thorough searches in the past.

In February 2000 at a conference in Adelaide, Australia, I managed to convince Takashi Nakano, the spokesman of the LEPS collaboration at SPring-8 near Osaka, that it is worthwhile to search for this particle in reactions induced by high-energy gamma quanta. In October 2002 Takashi Nakano reported on the first evidence of the new baryon from the $\gamma p$ reaction at the PANIC 2002 conference in Osaka [42]. Independently, the DIANA collaboration lead by Anatoly Dolgolenko at ITEP, Moscow, looked into the $K^+Xe$ bubble chamber data taken at ITEP back in 1986. They learned about our prediction from a talk given by Polyakov. Rescanning old films took more than two years, and in December 2002 the group reported on the observation of a very narrow $\Theta^+$ [43]. The SPring-8 and ITEP groups did not know about each other’s work, used completely different reactions, and observed the $\Theta^+$ peak using different techniques and final states, see Figs. 15,16. I believe that both groups can claim credit for the first independent evidence of the $\Theta^+$. The next groups knew about these two experiments, and gave very important confirmation using various reactions and final states, with higher statistical significance [44, 45, 46, 47, 48, 49, 50, 51, 52]. It should be noted, however, that there have been several experiments where $\Theta^+$ has not been seen [43, 54] (I briefly comment on their non-sighting at the end). Therefore, one is now looking forward to the next tour of dedicated experiments with higher statistics, for the issue to be finally resolved.

The ITEP experiment [43] being so far the only one where the resonance is formed directly in the $K^+p$ reaction, provides an additional valuable information. First, our prediction from a talk given by Polyakov. Rescanning old films took more than two years, and in December 2002 the group reported on the observation of a very narrow $\Theta^+$ [43]. The SPring-8 and ITEP groups did not know about each other’s work, used completely different reactions, and observed the $\Theta^+$ peak using different techniques and final states, see Figs. 15,16. I believe that both groups can claim credit for the first independent evidence of the $\Theta^+$. The next groups knew about these two experiments, and gave very important confirmation using various reactions and final states, with higher statistical significance [44, 45, 46, 47, 48, 49, 50, 51, 52]. It should be noted, however, that there have been several experiments where $\Theta^+$ has not been seen [43, 54] (I briefly comment on their non-sighting at the end). Therefore, one is now looking forward to the next tour of dedicated experiments with higher statistics, for the issue to be finally resolved.

The ITEP experiment [43] being so far the only one where the resonance is formed directly in the $K^+p$ reaction, provides an additional valuable information. First,
it gives the most stringent direct restriction on the $\Theta^+$ width, less than 9 MeV. Moreover, since the formation cross section is proportional to the width it can be estimated indirectly from the number of the observed events in the peak, with the result $\Gamma_\Theta = 0.9 \pm 0.3$ MeV \cite{55} (a conservative estimate is less than 3.6 MeV). Such a narrow width is compatible with other indirect estimates based on the fact that $\Theta$ has not been seen in previous experiments \cite{54}. In the $K^+n$ reaction, $\Theta^+$ is formed with the lab momentum 440 MeV and the speed $v_\Theta = 0.27c$. Taking the 1 MeV width, one finds that it travels a distance of 55 fm before it decays, far beyond the heavy nucleus where it has been produced! Therefore, the second piece of information one can extract from the ITEP data is that the $\Theta^+$ must be a compact object with weak interaction with the nuclear medium, otherwise it would have been rescattered or destroyed on its way through the large Xe nucleus. In particular, the weakly bound molecule-type $KN$ bound state can be probably ruled out.

If it indeed exists, $\Theta^+$ must be not only exotic in quantum numbers but a very unusual baryon: extremely narrow, compact and weakly interacting.

**Why is it named $\Theta^+$?**

After the two groups have announced the evidence of a new exotic baryon, I sent out a letter to all parties involved suggesting to name it “the $\Theta^+$ baryon”, with the following arguments \cite{57}:

- The old name $Z$ used sometimes in the past for exotic candidates is unfortunate – first, because it can be confused with the electroweak boson and second, because the old one- and two-star candidates were broad and in a much higher mass range
- It must be named by an upper-case Greek character, according to the tradition of naming baryons
- It must be distinct from anything used before and carry no associations with bosons
- The character must exist in LaTeX

The only capital Greek character satisfying all criteria is $\Theta$. It is a nice “round” character hinting that it is an isosinglet, like the $\Omega^-$ hyperon also sitting at the vertex of a big triangle (Fig. 13). However, in distinction from the $\Omega^-$ which has been predicted by Gell-Mann to close the decuplet of known baryons, $\Theta^+$ (if confirmed) opens the antidecuplet of unknown exotic baryons; some of them might have already been seen but other still await observation.

**VIII. WHY IS $\Theta^+$ SO NARROW?**

All experiments giving evidence of the $\Theta$ see that it is narrow, the most stringent bound being $\Gamma < 9$ MeV \cite{58}. The indirect estimates \cite{58, 59} show that it can be actually as small as 1 MeV or even less. If correct, $\Theta$ would be the most narrow strongly decaying particle made of light quarks. Any theoretical model of $\Theta$ has to explain the unusually small width first of all.

The Chiral Quark Soliton Model from where the story started, gave an estimate $\Gamma < 15$ MeV \cite{55} or, to be more precise, $3.6$ MeV $< \Gamma < 11.2$ MeV \cite{58}. A more careful analysis of the semileptonic hyperon decays (serving as an input for the estimate) performed in Ref. \cite{58} led to $\Gamma < 5$ MeV. This is not bad but as for now there are big theoretical uncertainties in the estimate.

Recently, there have been several attempts to explain the narrow width qualitatively from various 5-quark non-relativistic wave functions of the $\Theta$ or from ideas how such wave functions could look like. I think such attempts are futile, for several reasons, and a different technique has to be used to answer this intriguing question.

First, already in the nucleon the constituent quarks are relativistic. Had they been non-relativistic, one would face the paradoxes mentioned in the beginning of the paper, and in fact many other. One may say: OK, let the three constituent quarks in the nucleon be relativistic. But it makes matters even worse: the more relativistic are the quarks, the less is their “upper component” of the Dirac bispinor, meaning there are less than three quarks in the nucleon. Since the number of quarks minus the number of antiquarks is the conserved baryon number, it automatically means that one gets a negative distribution of antiquarks from the relativistic Dirac equation \cite{23}. This is nonsense. People who have attempted to get parton distributions from the constituent quark models have not focused on this paradox but it is an inevitable mathematical consequence of the Dirac equation. It is cured by adding the Dirac sea to valence quarks; only then the antiquark distribution becomes positive definite and satisfies the general sum rules \cite{23}. Thus, a fully relativistic description of nucleons is a must if one does not wish to violate general theorems.

Second, when one passes from the nucleons to the $\Theta^+$, the situation becomes even more dramatic. As stressed above, the $\Theta^+$ is relatively light because the additional $QQ$ pair is added in the form of the excitation of a light and strongly bound chiral field. If there have been any doubts that the three constituent quarks in the nucleon are relativistic, there is no place for such doubts for pseudoscalars as the quark masses there have to be ‘eaten up’ to zero, since they are Goldstone particles. Therefore, quarks in the $\Theta^+$ are essentially relativistic.

Third, and most important. It has been known since the work of Landau and Peierls (1931) that the quantum-mechanical wave function description, be it non-relativistic or relativistic, fails at the distances of the order of the Compton wave length of the particle. Measuring the electron position with an accuracy better than $10^{-11}$ cm produces a new electron-positron pair, by the uncertainty principle. One observes it in the Lamb shift and other radiative corrections. Fortunately, the atom
size is $10^{-8}$ cm, therefore there is a gap of three orders of magnitude where we can successfully apply the Dirac or even the Schrödinger equation. In baryons, we do not have this luxury. Measuring the quark position with an accuracy higher than the pion Compton wave length of 1 fm produces a pion, i.e. a new $QQ$ pair. Hence, the quantum-mechanical description of baryons with a fixed number of quarks is senseless.

The Chiral Quark Soliton Model provides such a quantum field-theoretic description of baryons, and to my knowledge, it is the only one today. This is why it is capable to explain the paradoxes of the constituent quark models, and of getting reasonable baryon properties, both from the phenomenological and theoretical points of views. However, it has a shortcoming: a reference to the large number of colors to justify the use of the mean field in baryons. Some day an accurate calculational scheme for baryons will be developed without referring to the mean field, as it has been developed for the atoms, but today we have to use the “Thomas–Fermi approximation” to understand baryons.

There is a well known alternative to the field-theoretic description when one returns to the more intuitive quantum mechanics. It is the use of the light cone quantization or the infinite momentum frame (IMF) \[62, 61\]. In that frame, and only in it, there is no production and annihilation of particles, so that the light-cone wave function makes sense, it does not contradict the Landau–Peierls limitation. There exists a prejudice that one uses the IMF only for high energy processes. This is not so: one can use it to find the static characteristics of baryons (like the magnetic moments or the axial constants). It is just a method where the language of wave functions makes sense despite that the number of constituents is not fixed. In the baryon rest frame the use of the wave function is contradictory in terms – it is like using coordinates and velocities in describing the hydrogen atom: no respect to the uncertainty principle.

In the IMF, the baryon wave function falls into separate sectors of the Fock space: three quarks, five quarks, etc. The general baryon light-cone wave function corresponding to the QCSM has been recently derived by Petrov and Polyakov \[24\]. The difference between the ordinary nucleon and the $\Theta^+$ is that the nucleon has a three-quark component (but necessarily has also a five-quark component) while $\Theta$’s Fock space starts from the five-quark component. In fact, the $\Theta$’s wave function is not qualitatively different from the five-quark component of the nucleon wave function (Fig. 18). They have similar distributions in longitudinal and transverse momenta, although their spin-flavor parts are, of course, different \[62\].

$\Theta^+ \rightarrow nK^+$ decay

Now we finally come to the question how to evaluate the $\Theta$ width. As I stressed above, in the constituent quark models where quarks are described by some wave function in the baryon rest frame, there is no possibility even to ask this question in a consistent way, leave alone to answer it. With the light-cone quantization, it is a legitimate question that can be answered.

Let us consider the decay amplitude of the $\Theta$ (or $\Xi_A/2$) into an octet baryon and a pseudoscalar meson. It is determined by the transition pseudoscalar coupling $g_{K\Theta}$ (similar to the diagonal $\pi N$ coupling $g_{\pi NN}$) which is related to the transition axial coupling $g_{A}^{\Theta \rightarrow NK}$ (similar to the nucleon axial constant $g_A$) by the approximate Goldberger–Treiman relation

$$g_{K\Theta} \approx \frac{g_{A}^{\Theta \rightarrow NK}(m_N + m_{\Theta})}{2F_K}$$

where $F_K \approx F_{\pi}$ is the kaon decay constant. Hence, it is sufficient to evaluate $g_{A}^{\Theta \rightarrow NK}$ as a transition matrix element of the axial charge between the $\Theta$ and the nucleon states. In the infinite-momentum frame there can be no production and annihilation of quarks, therefore the operator of the axial charge does not create or annihilate quarks but only measures the axial charge of the existing quarks \[24\]. (The same is true for the vector current as well.) Thus, the matrix element in question is non-zero only between the pentaquark and the five-quark component of the nucleon! \[63\].

FIG. 17: Uncertainty principle at work: When one attempts to measure the quark position in the nucleon to an accuracy better than the pion Compton wave length of 1 fm one produces a pion, i.e. a new $QQ$ pair. Hence, the quantum-mechanical description of baryons with a varying number of quarks, is senseless.

FIG. 18: In the infinite-momentum frame (and only there!) baryon wave functions are well defined and fall into separate sectors of the Fock space: 3 quarks, 5 quarks,... $\Theta$’s wave function is not qualitatively different from the five-quark component of the nucleon wave function.
To be concrete, let us consider the $\Theta^+ \rightarrow nK^+$ decay. In this case, the axial charge has the quantum numbers of the $K^+$ meson, $J^{PC}_{\Theta^+}=s\gamma_5u$. It annihilates the $u$ quark creating the $s$ quark and annihilates $\bar{s}$ quarks creating $\bar{u}$ ones. Correspondingly, there are two processes determining the $\Theta^+ \rightarrow nK^+$ decay, shown in Fig. 19. The transition axial constant is given by the normalized matrix element of this axial charge:

$$g_{A}^{\Theta^+ \rightarrow nK^+} = \frac{<n(5)|J_{05}^{K^+}|\Theta^+>}{\sqrt{N_n^{(3)}+N_n^{(5)}+...}}$$

where $N_n^{(3,5)}$ is the normalization of the 3- and 5-quark components of the neutron wave function in the infinite-momentum frame $n^{(3)}$ and $n^{(5)}$, respectively; $N_{\Theta}^{(5)}$ is the same for the $\Theta$: “...” stand for the omitted 7- and higher-quark Fock components. In the numerator, $n^{(5)}$ is normalized to $N_n^{(5)}$ whereas in the denominator one has to use the full neutron’s normalization. One expects that $N_n^{(5)} \ll N_n^{(3)}$, otherwise the neutron would be a mainly 5-quark baryon.

Therefore, the transition axial constant $g_{A}^{\Theta^+ \rightarrow nK^+}$ is first of all suppressed by the factor $(N_n^{(5)}/N_n^{(3)})^{1/2}$, that is suppressed to the extent the 5-quark component of the neutron is less than its 3-quark component. Additional suppression comes from the peculiar flavor structure of the neutron’s 5-quark component where the antiquark is in the flavor-singlet combination with one of the four quarks. A detailed calculation is on the way [62] but a very crude preliminary estimate shows that the $\Theta$ width can be as small as 0.7 MeV.

It is clear that this suppression of the pseudoscalar transition coupling of the $\Theta$ is quite general, it applies also to scalar, vector, magnetic etc. couplings. It has been already noticed in the CQSM [64]. In fact, all means of exciting the $\Theta^+$ from the nucleon seem to be suppressed, be it via $K$, $K^+$ or $K^*(1430)$ exchanges. From this point of view, exciting first some high nucleon resonance with a presumably large 5-quark component decaying then into the $\Theta$ can be a promising way to produce it, as suggested by one of the CLAS experiments [45].

However, the excitation of a 5-quark nucleon resonance cannot be large either.

As to the high energy experiments, here the situation with respect to the $\Theta^+$ production is probably even worse. At high energies all exchanges with the non-vacuum quantum numbers die out and only the gluonic pomeron survives. It may be therefore even more difficult to excite the 5-quark $\Theta$ by touching a mainly 3-quark nucleon by soft gluons, than via the meson exchange which is itself small. This could explain the present day nonsighting of the $\Theta^+$ in high energy experiments. Probably certain kinematical cuts and/or association production should be imposed on the high energy data to help disclosing the $\Theta^+$.

**IX. SUMMARY**

1. 93% of the light baryon masses are due to the Spontaneous Chiral Symmetry Breaking well explained by instantons. It implies that quarks get a large dynamically-generated mass, which inevitably leads to their strong coupling to the chiral fields ($\pi, K, \eta$).

2. There is nothing queer in calling baryons solitons – they are no more solitons than atoms, which one may like to call “the solitons of the self-consistent electrostatic field”. It is a technical aspect, not a crucial one. The important question is what forces are responsible for binding quarks.

3. Assuming that the chiral forces are essential in binding quarks together in baryons, one gets the lowest multiplets $(8, \frac{8}{3})$, $(10, \frac{5}{3})$ and $(\overline{10}, \frac{-5}{3})$. The predicted $\Theta^+ \approx uudd\bar{s}$ is light because it is not a sum of constituent quark masses but rather a collective excitation of the mean chiral field inside baryons.

4. The non-relativistic wave-function description of an atom is valid at distances $10^{-8}$ cm, but fails at $1/m_c = 10^{-11}$ cm. For baryons, “$10^{-11}$ cm” is 1 fm. Relativistic field-theoretic description of baryons is a must.

5. The very narrow width of the $\Theta^+$ can be probably explained.

6. If confirmed, $\Theta^+$ will not only be a new kind of subatomic particle but will seriously influence our understanding of how do ordinary nucleons “tick” and what are they “made of”.

Borrowing John Collins’ joke: it would be as if a new chemical element between hydrogen and helium is discovered.
Acknowledgments

I thank Victor Petrov, Pavel Pobylitsa and Maxim Polyakov for a long-time collaboration during which the views presented here were formulated, and Yakov Azimov and Mark Strikman for their comments on the manuscript and useful discussions. I am grateful to the Physics Department of the Pennsylvania State University where I have spent a part of the sabbatical year, for hospitality and support.

[1] Based on the invited talks at the meeting of the American Physical Society (Denver, May 1, 2004) and at the international meeting Continuous Advances in QCD (Minneapolis, May 14, 2004).
[2] Yu. Dokshitzer, D. Diakonov and S. Troian, Phys. Rep. 147, 269 (1986).
[3] M. Glick, E. Reya and A. Vogt, Zeit. Phys. C 67, 433 (1995).
[4] B.W. Filippine and X. Ji, Adv. Nucl. Phys. 26, 1 (2001), hep-ph/0101224.
[5] M.M. Pavan, R.A. Arndt, I.I. Strakovsky and R.L. Workman, in Proceedings of 9th International Symposium on Meson-Nucleon Physics and the Structure of the Nucleon (MENUS2001), Washington, DC, USA, July 26-31, 2001, H. Haberzettl and W.J. Briscoe (eds.); Nucl. Phys. B 611, 110 (2001), hep-ph/0111066.
[6] P. Schweitzer, hep-ph/0312376.
[7] H. Leutwyler, Nucl. Phys. Proc. Suppl. 94, 108 (2001), hep-ph/0011049.
[8] D. Diakonov, Nucl. Phys. Proc. Suppl. 51, 173 (2003), hep-ph/0212026.
[9] E. Shuryak, Nucl. Phys. B 203, 93 (1982).
[10] D. Diakonov and V. Petrov, Nucl. Phys. B 245, 259 (1984); D. Diakonov, M. Polyakov and C. Weiss, Nucl. Phys. B 461, 539 (1996), hep-ph/9510232.
[11] D. Diakonov and V. Petrov, Lett. B 147, 351 (1984); Nucl. Phys. B 272, 457 (1986).
[12] M.-C. Chu, J. Grandy, S. Huang and J. Negele, Phys. Rev. D 49, 6039 (1994); J. Negele, Nucl. Phys. Proc. Suppl. 73, 92 (1999), hep-lat/9810053.
[13] P. Bowman, U. Heller, D. Leinweber, A. Williams and J. Zhang, Nucl. Phys.Proc. Suppl. 128, 23 (2004), hep-lat/0403002.
[14] P. Faccioli and T. DeGrand, Phys. Rev. Lett. 91, 182001 (2003), hep-ph/0304219.
[15] F. Karsch, E. Laermann and A. Peikert, Nucl. Phys. B 605, 579 (2001), hep-lat/0012023.
[16] C.W. Bernard et al., Phys. Rev. D 64, 054506 (2001), hep-lat/0104002.
[17] H. Georgi and A. Manohar, Nucl. Phys. B 234, 189 (1984).
[18] S. Kahana, G. Ripka and V. Soni, Nucl. Phys. A 415, 351 (1984); S. Kahana and G. Ripka, Nucl. Phys. A 429, 402 (1984).
[19] M.S. Birse and M.K. Banerjee, Phys. Lett. B 136, 284 (1984).
[20] G. ’t Hooft, Nucl. Phys. B 72, 461 (1974).
[21] E. Witten, Nucl. Phys. B 160, 57 (1979).
[22] D. Diakonov and V. Petrov, JETP Lett. 43, 75 (1986) [Pisma Zh. Eksp. Teor. Fiz. 43, 57 (1986)]; D. Diakonov, V. Petrov and P.V. Pobylitsa, Nucl. Phys. B 306, 809 (1988).
[23] D. Diakonov, V. Petrov, P. Pobylitsa, M. Polyakov and C. Weiss, Nucl. Phys. B 480, 341 (1996), hep-ph/9606314, hep-ph/9703420.
[24] V. Petrov and M. Polyakov, hep-ph/0307077.
[25] D. Diakonov and V. Petrov, in Handbook of QCD, M. Shifman, ed., World Scientific, Singapore (2001), vol. 1, p. 359, hep-ph/0009006.
[26] C. Christov, A. Blotz, H.-C. Kim, P. Pobylitsa, T. Watabe, Th. Meissner, E. Ruiz Arriola and K. Goeko, Prog. Part. Nucl. Phys. 37, 91 (1996), hep-ph/9604441.
[27] D. Diakonov, V. Petrov and M. Praszalowicz, Nucl. Phys. B 323, 53 (1989).
[28] P.V. Pobylitsa, M.V. Polyakov, K. Goeko, T. Watabe and C. Weiss, Phys. Rev. D 59, 034024 (1999), hep-ph/9804436.
[29] M. Wakamatsu and H. Yoshih, Nucl. Phys. A 524, 561 (1991). The first qualitative explanation of the ‘spin crisis’ from the Skyrme model point of view, namely zero fraction of proton spin carried by quarks’ spin, was given in: S.J. Brodsky, J.R. Ellis and M. Karliner, Phys. Lett. B 206, 309 (1988).
[30] K. Goeko, M. Polyakov and M. Vanderhaeghen, Prog. Part. Nucl. Phys. 47, 401 (2001), hep-ph/0106012.
[31] B. Dressler, K. Goeko, M. Polyakov and C. Weiss, Eur. Phys. J. C 14, 147 (2000), hep-ph/9909541.
[32] V. Petrov, P. Pobylitsa, M. Polyakov, I. Börnig, K. Goeko and C. Weiss, Phys. Rev. D 57, 4325 (1998), hep-ph/9710270.
[33] A. Hayashi, G. Eckart, G. Holzwarth and H. Walliser, Phys. Lett. B 147, 5 (1984).
[34] M. Karliner and M.P. Mattis, Phys. Rev. D 31, 2833 (1985); ibid. 34, 1991 (1986).
[35] N. Itzhaki, I.R. Klebanov, P. Ouyang and L. Rastelli, Nucl. Phys. B 684, 264 (2004), hep-ph/0309305.
[36] E. Witten, Nucl. Phys. B 160, 433 (1983).
[37] D. Diakonov and V. Petrov, hep-ph/0312144.
[38] D. Diakonov, V. Petrov and M. Polyakov, Z. Phys. A 359, 305 (1997), hep-ph/9703373.
[39] R.L. Jaffe, SLAC-PUB-1774 (1976); H. Hogaasen, P. Sorba, Nucl. Phys. B 145, 119 (1978); J.J. de Swart, P.J. Mulders, L.J. Somers, in Baryon 1980, N. Isgur, ed., Toronto University (1980); hep-ph/9703373.
[40] L.C. Biedenharn and Y. Dothan, in From SU(3) to Gravity (Ne’eman Festschrift), E. Costman and G. Tauber, eds., Cambridge University Press (1986), p. 15 [Duke University preprint (1984)]; M. Praszalowicz, in Skyrmi ons and Analogies, M. Jezabek and M. Praszalowicz, eds., World Scientific, Singapore (1987), p. 112, see the commented reprint in Phys. Lett. B 575, 234 (2003), hep-ph/0309305, H. Walliser, Nucl. Phys. A 548, 649 (1992), see also H. Walliser and V.B. Kopeliovich, J. Phys. G: Theor. Phys. 29, 433 (2003) [Zh. Eksp. Teor. Fiz. 124, 483 (2003)], hep-ph/0304055.
[41] J.S. Hyslop, R.A. Arndt, L.D. Roper and R.L. Workman, Phys. Rev. D 46, 961 (1992).
[42] T. Nakano (LEPS Collaboration), Talk at the PANIC 2002 (Oct. 3, 2002, Osaka); T. Nakano et al., Phys. Rev. Lett. 91, 012002 (2003), hep-ex/0301020.
[43] V.A. Shebanov (DIANA Collaboration), Talk at the Session of the Nuclear Physics Division of the Russian Academy of Sciences (Dec. 3, 2002, Moscow); V.V. Barmin, A.G. Dolgolenko et al., Phys. Atom. Nucl. 66, 1715 (2003) [Yad. Fiz. 66, 1763 (2003)], hep-ex/0304040.
[44] S. Stepanyan, K. Hicks et al. (CLAS Collaboration), Phys. Rev. Lett. 91, 252001 (2003), hep-ex/0307018.
[45] V. Kubarovsky et al. (CLAS Collaboration), Phys. Rev. Lett. 92, 032001 (2004), hep-ex/0311046.
[46] J. Barth et al. (SAPHIR Collaboration), Phys. Lett B 572, 127 (2003), hep-ex/0307083.
[47] A.E. Asratyan, A.G. Dolgolenko and M.A. Kubantsev, submitted to Yad. Fiz., hep-ex/0309042.
[48] A. Airapetian et al. (HERMES Collaboration), Phys. Lett. B 585, 213 (2004), hep-ex/0312044.
[49] A. Aleev et al. (SVD Collaboration), hep-ex/0401024.
[50] M. Abdel-Bary et al. (COSY-TOF Collaboration), hep-ex/0403044.
[51] F.Zh. Aslanyan, V.N. Emelyanenko, G.G. Rikhvkitzkaya, hep-ex/0403085.
[52] S. Chekanov et al. (ZEUS Collaboration), hep-ex/0403051.
[53] J.Z. Bai et al. (BES Collaboration), hep-ex/0402012.
[54] K.T. Knopfle, M. Zavertyae and T. Zivko (HERA-B Collaboration), hep-ex/0403020.
[55] R.N. Cahn and G.H. Trilling, Phys. Rev. D 69, 011501 (2004), hep-ph/0311245.
[56] S. Nussinov, hep-ph/0307367 R.A. Arndt, I.I. Strakovsky and R.L. Workman, Phys. Rev. C 68, 042201 (2003), nucl-th/0308012 J. Haidenbauer and G. Krein, Phys. Rev. C 68, 052201 (2003), hep-ph/0309243.
[57] D. Diakonov, e-letter to parties involved (April 12, 2003), see http://thd.pnpi.spb.ru/news/Theta.html.
[58] D. Diakonov, V. Petrov and M. Polyakov, hep-ph/0404212.
[59] A. Rathke, Diploma thesis, Bochum University (1998).
[60] S.J. Brodsky and G.P. Lepage, in Perturbative Quantum Chromodynamics, A.H. Mueller (ed.), World Scientific, Singapore (1989).
[61] R. Perry, A. Harindranath and K. Wilson, Phys. Rev. Lett. 65, 2959 (1990).
[62] D. Diakonov, V. Petrov and M. Polyakov, in preparation.
[63] D. Diakonov and V. Petrov, Phys. Rev. D 69, 094011 (2004), hep-ph/0310212.
[64] M. Polyakov and A. Rathke, Eur. Phys. J. A 18, 691 (2003), hep-ph/0303138 M. Polyakov, private communication.