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STRUCTURE FUNCTIONS OF SEA QUARKS, KOGUT-SUSSKIND POLE AND DYNAMICS of pp- AND $p\bar{p}$-INTERACTIONS AT HIGH ENERGIES

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INTRODUCTION.

Large violation of the **Ellis-Jaffe** parton sum rule[1] for the spin-dependent proton structure function $g_1^p(x)$ observed recently by the **EMC**[2] has led to the "spin crisis" (see[3]). Almost zero value of helicity carried by quarks inside the proton determined from an analysis of the results[2] means that, contrary to naive expectation, the helicity carried by sea quarks is so large that it almost completely compensates for the helicity of valence quarks. The anomalous large helicity of sea quarks has turned out extremely difficult to explain within the known parametrizations of the distributions of sea quarks[4].

A solution was suggested in our work[5] where it was noted that introducing a new Regge trajectory with a high intercept $\alpha(0) \approx 1$ caused by the Adler-Bell-Jackiw (ABJ) anomaly[6] and the Kogut-Susskind pole [7] (called further "anomalon") provides at least qualitative explanation of the experimental data on the structure functions of sea quarks.

In the present work we discuss in detail the nature of this trajectory and show that this trajectory defines not only the behavior of the spin-dependent sea quark structure functions in the region of small x but also some peculiarities of $pp-$ and $p\bar{p}-$ interactions at high energies.

STRUCTURE FUNCTIONS OF SEA QUARKS IN THE RANGE $X \to 1$.

The **EMC** has measured the proton structure function $g_1^p(x)$. Within the parton model this function is expressed via the quark distributions over helicity within a proton :

$$g_1^p(x) = \frac{1}{2}\sum_i e_i^2\ [q_+^i(x) - q_-^i(x)],$$

where $q_{+(-)}^i(x)$ is the probability that a quark of a flavour i has the helicity parallel (antiparallel) to the helicity of the proton. The integrals of the difference of the distribution functions:

$$\Delta q^i = \int_0^1 dx\ [q_+^i(x) - q_-^i(x)] \tag{1}$$

are connected with the matrix elements of the axial vector current between the baryon octet states which can be measured in the nucleon and hyperon $\beta-$decays and also in the elastic scattering of a neutrino off a proton. Joint analysis of these data produces the magnitude of the helicity carried by quarks[2]:

$$\Delta\Sigma = \Delta u + \Delta d + \Delta s = 0.12 \pm 0.24. \tag{2}$$

So, under generally accepted assumptions about $\mathbf{SU(3)}_f$ breaking, one can estimate the helicity of sea quarks[2]:

$$\Delta\Sigma^s = \Delta u^s + \Delta d^s + \Delta s^s = -0.95 \pm 0.16 \pm 0.23 \tag{3}$$

At the same time, it is known that the momentum fraction of the proton carried by sea quarks is small[8]:

$$p^s = \sum_q \int_0^1 dx\ x(q_+^s(x) + q_-^s(x)) \approx 0.074 \tag{4}$$

The difference of an order between the magnitudes of the integrals, Eq. (3) and Eq. (4), may be treated as a quantitative measure of the so-called "spin crisis".

In work[5], within the nonperturbative **QCD** based on the model of the **QCD** vacuum as an instanton liquid[9] the x dependence of the distribution functions of sea quarks inside the proton has been considered. Here we present a more detailed than in[5] derivation of basic relations and also point out some new moments concerning the distribution functions of sea quarks inside the proton.

To find expressions of distributions, one uses the noncovariant perturbative theory (NCPT) in the system of infinite momentum (see[10]). The distribution functions are connected with the light cone proton wave function expanded over free quark and gluon Fock states ($3q$, $3qq\bar{q}$, $3qg$, ...) by the relation[11]

$$q_{f/p}(x) \propto \sum_n \int [dk_{\perp i}][dx_i]\ \delta(x - x_q)\cdot \mid \Psi_{(n)}(k_{\perp i}, x_i) \mid^2, \tag{5}$$

where $\Psi_{(n)}(k_{\perp i}, x_i)$ is the contribution of an n-particle component of the Fock state , $x_i = (k^0 + k^3)_i/(p_0 + p_3)$ is the fractional (light-cone) momentum of the proton carried by the $i-$th parton ($\sum_{i=1}^n x_i = 1$) and $k_{\perp i}$ is its transverse momentum ($\sum_{i=1}^n \vec{k}_{\perp i} = 0$).

Within the NCPT at every vertex the momentum is conserved but the energy is not. Therefore, the contribution of an $n-$particle Fock state is

$$\Psi_n(k_{\perp i}, x_i) \propto \frac{\Gamma(k_{\perp i}, x_i)}{\frac{1}{2P}\left[M_p^2 - \sum_{i=1}^{n} \frac{m_i^2 + k_{\perp i}^2}{x_i}\right]}, \tag{6}$$

where m_i are the quark masses, M_p is the proton mass, $P \to \infty$ is the proton momentum in the infinite momentum frame, $\Gamma(k_{\perp i}, x_i)$ is a vertex function defined by the dynamics of the $n-$particle state production from the initial three-quark proton wave function, and the denominator is the energy difference of the initial and intermediate states $E_0 - E_N$.

As $x \to 1$, the contribution of a five-quark state $\Psi_5(k_{\perp i}, x_i)$ dominates in the sea quark distributions. The diagrams corresponding to this contribution within the instanton model of the **QCD** vacuum are presented in Fig. 1.

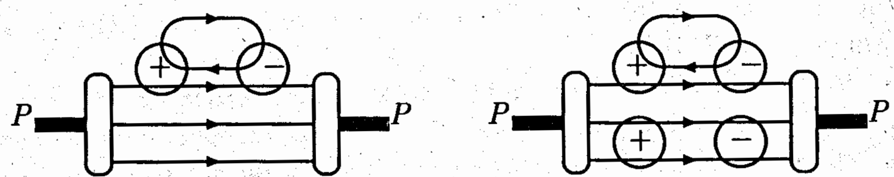

Fig. 1 The instanton contribution to the five-quark component of the nucleon wave function (+ (-) - instanton (anti-instanton))

Instanton vertices are defined by the 't Hooft effective interaction[12]:

$$\begin{aligned}\mathcal{L}_{eff}^{inst} &= \frac{4\pi^2\rho_c^2}{3}\{\sum_{i\neq j}\int \frac{d^4k_1\, d^4k_2\, d^4k_3\, d^4k_4}{(2\pi)^{12}}\,\delta^{(4)}(k_1 + k_2 - k_3 - k_4) \\ &\cdot \exp[-\rho_c \sum_{n=1}^{n=4} \mid k_n \mid]\, \bar{q}_{iR}(k_1) q_{iL}(k_3) \bar{q}_{jR}(k_2) q_{jL}(k_4) \\ &\cdot [1 + \frac{3}{32}(1 - \frac{3}{4}\sigma_{\mu\nu}^i \sigma_{\mu\nu}^j)\lambda_i^a \lambda_j^a] + (R \leftrightarrow L)\},\end{aligned} \tag{7}$$

where the coupling constant is obtained by factorization of the N_f-fermion 't Hooft Lagrangian within the instanton liquid model[13], $q_{R,L} = [(1 \pm \gamma_5)/2]q$, $i, j-$ quark flavours, $\rho_c \approx 2\ Gev^{-1}$ is an average size of an instanton in the **QCD** vacuum[9].

It should be stressed that the vertex, Eq. (7), differs principally from the perturbative quark-quark vertex caused by the one-gluon exchange. First, opposite to the quark-gluon vertex the instanton-induced vertex flips the quark helicity, so that for the N_f- fermion vertex the helicity change is equal to $2N_f$. Second, the vertex, Eq. (7), is nonzero only in the case of quarks of different flavours. Thus, the quark sea produced by instantons inside the proton should not be exactly $SU_f(2)$-symmetrical. Third, in Eq. (7) the exponential factor depending on an instanton size provides a natural cut-off parameter over the transverse momentum of intermediate state quarks (see Fig. 1). The last is easily proved if we remind that in fact Eq. (7) is obtained from the Furier-transform of the zero fermion modes of quarks in the instanton field being in the origin[9]. Lorenz invariance leads to a trivial change in Eq. (7): $\sum_{n=1}^{4} \mid k_n \mid \to \sum_{n\neq m} \mid k_n - k_m \mid$.

Thus, every instanton vertex in Fig. 1 produces the factor $\exp[-\rho_c(E_n - E_0)]$, where E_0 (E_n) is the energy sum of incoming (outgoing) fermion lines, respectively. This factor obviously takes into account that the life time of an instanton configuration in Euclidian time is of an order of $\tau \approx 1/\rho_c$.

So it follows that within the model of instanton vacuum the vertex function $\Gamma(x_i, k_{\perp i})$ in Eq. (6) produced from the diagrams of Fig. 1 has the exponential form

$$\Gamma(x_i, k_{\perp i}) \propto \exp\{-\frac{\rho_c}{2P}\left(\frac{m_i^2 + k_{\perp i}^2}{x_i} - M_p^2\right)\}.$$

In the integral, Eq. (5), the region

$$\sum_{i=1}^{5} \frac{m_i^2 + k_{\perp i}^2}{x_i} \approx M_p^2$$

dominates. So one can put:

$$q_{f/p}(x) \propto A \int \frac{[dx_i]\ \delta(x - x_q)}{\left(M_p^2 - \sum_{i=1}^{5} \frac{m_{\perp i}^2}{x_i}\right)^2}, \tag{8}$$

where $m_{\perp q} = \sqrt{m_q^2 + k_{\perp q}^2}$ is the transverse mass of q-quark inside the proton.

In addition, since the instanton vertex, Eq. (7), changes the quark helicity by $2N_f$, then the total angular momentum conservation requires nonzero angular momentum of quarks in the intermediate state in Fig. 1. Then, it is obvious that the angular momentum projection should be equal to the helicity change at an instanton. This allows ones to estimate possible transverse momentum in a five-quark configuration by the relation:

$$\rho_c k_\perp \approx 2N_f. \tag{9}$$

So for $\rho_c \approx 2\ Gev^{-1}$, $N_f = 2$ one obtains $k_\perp \approx 2\ Gev$ which is essentially larger than average momentum of the valence quark inside the proton $k_\perp \approx 0.35\ Gev$.

Thus, the only five-quark configuration of the proton Fock state satisfies both the angular momentum and momentum conservation which contains at least two quarks with large transverse momentum. For this configuration from Eq. (8) one can easily obtain the asymptotics of the distribution functions:

$$q^s_{x\to 1} \propto \begin{cases} (1-x)^5 & m^2_{q\perp} << M_p^2, \\ (1-x)^3 & m^2_{q\perp} >> M_p^2, \end{cases} \tag{10}$$

So for the quark sea from light $u-$, $d-$, $s-$ quarks there are two regimes, Eq. (10), and for heavy $c-$, $b-$, $t-$ quarks there is one:

$$q^s_{x\to 1} \propto (1-x)^3.$$

Note that in fact the hardness of the sea is caused by kinematical reasons. The matter is that the dominant contribution to Eq. (8) is due to the configuration for which the denominator is minimal. It is easy to show (see f. i. [11]) that the minimum exists at $\frac{m^2_{\perp i}}{x_i} \approx \frac{m^2_{\perp j}}{x_j} \approx const$, i. e. the quark with larger transverse mass has larger value of x. This condition means that all quarks in an $n-$particle state have the same rapidity, and thus in this configuration a proton does not "decompose". Thus, we have got a hard nonperturbative sea inside the proton which for large transverse mass of a sea quark has the same form (as $x \to 1$) of distributions as that of a valence quark has.

The hard component of the quark sea produces very interesting consequences. Indeed, the experiments on charm production in the hadron interactions[14] produce the hard spectrum of charmed particles. At the same time, this spectrum is practically independent of the type of a hadron into which the charm is fragmented. The experiments on cumulative particle production off nuclei[15] also unambiguously indicate the hardness and similarity of spectra of all cumulative particles. So in work[16] it was noted that to explain the results[15] it is necessary to suppose the distribution of a quark sea in nuclei of the same hardness as the distribution function of valence quarks in nuclei.

In our approach, these effects are easily explained by the fact that a quark with a large transverse mass provides a dominant contribution to the momentum of a hadron produced. Therefore, the spectra of secondary particles are almost completely defined by the structure functions of a hard quark and are independent of the fragmentation process. Thus, we find the form of x-dependence as $x \to 1$ of the sea inside the proton

$$q^s_+(x)_{x\to 1} = P(x)(1-x)^n + N_1(x)(1-x)^5 + N_2(x)(1-x)^3, \tag{11a}$$
$$q^s_-(x)_{x\to 1} = P(x)(1-x)^n + 2N_1(x)(1-x)^5 + 2N_2(x)(1-x)^3, \tag{11b}$$

where the first terms of these relations describe the contribution of a quark sea due to a perturbative gluon ($n \approx 7$ within the quark-counting rule (see [17])), $P(x)$, $N_1(x)$, $N_2(x) \to const$ as $x \to 1$.

The difference between the coefficients in Eq. (11a) and Eq. (11b) comes from the fact that the sea quark helicity is antiparallel to the helicity of the valence quark off which the former is produced. Similarly, in our model it could be the substantial breakdown of $\mathbf{SU}(2)_f$ in the sea quark distribution functions as $x \to 1$: $\bar{d}_I(x) \approx 2\bar{u}_I(x)$ (I is the instanton part of Eq. (11a) and Eq. (11b)). Recently, the direct experimental evidence of the d sea excess has appeared[18] *):

$$\int_0^1 dx\ (\bar{d}(x) - \bar{u}(x)) = 0.140 \pm\ 0.024.$$

*)This result is apparently to be considered as a preliminary one because nuclear effects inside the deuteron can be essential (L.P. Kaptari, A.Yu. Umnikov private communication)

In fact, the negative helicity of sea quarks and $\mathbf{SU}(2)_f$–breakdown of the sea is caused by the properties of zero fermion modes in the instanton field from which the Lagrangian, Eq. (7), is constructed.

STRUCTURE FUNCTIONS OF SEA QUARKS IN THE RANGE $X \to 0$ AND THE KOGUT-SUSSKIND POLE.

As $x \to 0$ all quark and gluon configurations of the proton wave function are valuable and behavior of the distribution functions in this region is specified by the Regge asymptotics. Usually one assumes (see [17]) that the Pomeron exchange with intercept $\alpha_p(0) \approx 1$ dominates in the sum of the distributions $q^s(x) = q^s_+(x) + q^s_-(x)$, and hence:

$$\lim_{x\to 0} q^s(x) \propto 1/x, \tag{12}$$

whereas the difference $\Delta q^s(x) = q^s_+(x) - q^s_-(x)$ is specified by the A_1-meson trajectory with $\alpha_{A_1} \approx 0$, and therefore

$$\lim_{x\to 0} \Delta q^s(x) \propto const. \tag{13}$$

In fact, Eq. (13), which has been used in the analysis of experimental data by the EMC[2], is derived from the selection rule

$$\sigma(-1)^I G = -1 \tag{14}$$

for the Regge trajectories contributing to the structure function $g_1^p(x)$[17] (σ is signature). The well-known A_1–trajectory with $I=1$, $\sigma=-1$, $G=-1$ satisfies this selection rule. However, it is obvious that this trajectory cannot contribute to the isosinglet anomalous combination $\Delta\Sigma(x) = \Delta u(x) + \Delta d(x) + \Delta s(x)$. In accordance with the rule, Eq. (14), the only Regge singularity capable to contribute to $\Delta\Sigma$ is that with quantum numbers $I=0$, $\sigma=-1$, $G=1$, $C=1$.

In [5] it was noted that to explain the difference of an order between the momentum fraction carried by sea quarks, Eq. (4) , and their helicity , Eq. (3), it is necessary to suppose the singular behavior at small x of the flavour singlet spin-dependent distribution function of sea quarks. We remind that earlier[19] an attempt was made to explain the EMC effect by anomalous dependence of the structure function as $x \to 0$

$$g_1^p(x)_{x\to 0} \to C/(x \ln^2 x). \tag{15}$$

The assumption, Eq. (15) , with $C = 0.135$ essentially enhanced the value of the quark helicity extracted from the EMC data

$$\Delta\Sigma \approx 0.5,$$

which matches the EMC data with the predictions of the constituent quark model $\Delta\Sigma \approx 1$ ($\Delta\Sigma \approx 0.65$ in the bag model). In[19] the behavior, Eq. (15) , was connected with the contribution of the Pomeron-Pomeron cut $P \otimes P$. However, in [20] it was argued that the $P \otimes P$-cut cannot contribute to $g_1^p(x)$ due to the selection rule, Eq. (14).

In [5] we have suggested a new Regge trajectory with a high intercept $\alpha(0) \approx 1$ which allows us to obtain the singularity of the distribution functions as $x \to 0$. There it was also pointed out that the trajectory was probably related with the Kogut-Susskind ghost in QCD[7]. The pole in the correlator of the pseudo-vectors

$$K_\mu = \frac{\alpha_s}{4\pi}\epsilon_{\mu\rho\sigma\alpha}A^a_\rho(\partial_\alpha A^a_\sigma - \frac{1}{3}g f_{abc}A^b_\alpha A^c_\sigma),$$
$$\int e^{iqx} < K_\mu(x)K_\nu(0) >_{q\to 0} \to \frac{g_{\mu\nu}}{q^2}\lambda^4 \tag{16}$$

was introduced in the work[7] in order to resolve the puzzle of the anomalously large mass of the η'-meson ($U_A(1)$–problem). Further, it has been shown [21, 22] that this pole is due to the nonperturbative properties of the **QCD** vacuum, namely, with the periodical structure of the **QCD** vacuum over a collective coordinate

$$Q = \int d^4x\, \partial_\mu K_\mu,$$

which is the topological charge of the large gauge transformation. At an instanton (anti-instanton) the topological charge change is $\Delta Q = 1\ (-1)$. Thus, there are many energy degenerate states with different values of the topological charge Q separated by penetrable (due to instantons) wells. The vacuum state is the superposition of the states with definite Q:

$$\Psi_{vac} = \sum_Q e^{iQ\theta}\Psi_Q,$$

where θ is the quasimomentum of the "θ-vacuum"[21, 22].

This problem is completely similar to the problem of calculating the Green function of an electron in the periodic lattice (see [22]). It is well known that here the gapless excitation appears with the Green function of a free electron on the periodic lattice:

$$\int dt\ e^{i\omega t} i < TX(t)X(0) >|_{\omega\to 0} \to -\frac{1}{\omega^2 m^*},$$

where m^* is the effective mass of an electron determined by the barrier penetrability. Thus, the pole of Eq. (16) also corresponds to free motion of the system but now along the variable of the topological charge.

Further, if we put a quark inside the "$\theta-$vacuum", it naturally would not have definite helicity. This follows from that the divergence of the axial-vector current is related through the ABJ anomaly to the topological charge density:

$$\partial_\nu j_\nu^5 = 2N_f \partial_\nu K_\nu \equiv 2N_f \frac{\alpha_s}{8\pi} G^a_{\mu\nu} \tilde{G}^a_{\mu\nu}, \tag{17}$$

where the terms proportional to the quark current masses are omitted.

From Eq. (17) it is easy to show that the change of the axial charge is connected with variation of the topological charge by the relation[23, 24]:

$$\Delta Q_5 = -2N_f \Delta Q. \tag{18}$$

Therefore, the free motion along Q means unconservation of the quark helicity in the "$\theta-$vacuum". For example, at an instanton $\Delta Q = 1$ and the helicity change is:

$$\Delta\Sigma_q = \Delta Q_5 = -2N_f.$$

Thus, in order to find an average value of the quark helicity inside the proton one should take into account the interaction of quarks with the ghost pole.

This problem was considered by Veneziano[25] who within the effective QCD Lagrangian approach derived the Goldberger-Treiman relation for the singlet axial-vector coupling constant

$$G_A^{(1)} = \Delta\Sigma = G_A^{\eta'} + G_A^{\eta'_0}, \tag{19}$$

where the first term of the right part is the contribution of a direct interaction of the η'-meson with the proton and the second is the contribution of the interaction via the Kogut-Susskind pole (Fig 2). Note that to explain the EMC result, Eq. (2) , one needs a large negative contribution of the ghost to the divergence[25]: $G_A^{\eta'_0} \approx -G_A^{\eta'} \approx -1$.

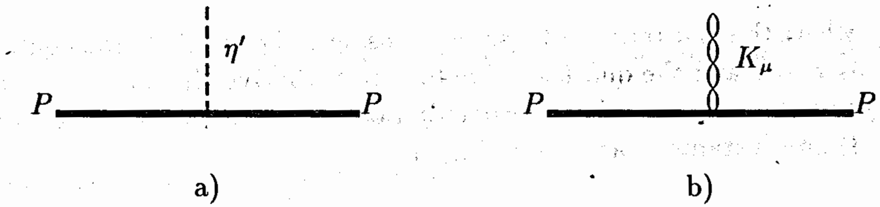

Fig. 2 The contribution to the divergence of axial-vector current of a) direct interaction of $\eta'-$meson, b) via the Kogut-Susskind pole

Further, the Kogut-Susskind pole cannot contribute to the structure function $g_1^p(x)$ as it causes the flip of the chirality of the proton. However, the double-pole exchange can contribute to the Compton amplitude of the forward scattering off the nucleon, as in Fig 3, which is produced by quadrating the diagram of Fig 2b.

In works[21, 22], it was argued that the Kogut-Susskind pole effectively takes into account the contribution of heavy gluonic states to the correlator of hadron currents. Hence, we suppose that this pole is reggezied and its trajectory ("anomalon") contains gluonic states with odd spin. It is easy to see that the "anomalon" has the quantum numbers $\sigma = -1$, $P = +1$, $C = -1$, $I = 0$. The intercept of the new trajectory should obviously be equal to unity, $\alpha_A(0) = 1$, as $t \to 0$ we have the massless pseudo-vector pole, Eq. (16).

From this it is easy to obtain x-dependence of the contribution of the "anomalon" to the structure function $g_1^p(x)$. Namely, the $A \otimes A$-cut corresponds to the diagram of Fig 3 and its contribution is

$$g_1^p(x)_{x\to 0} \to\ -a/(x \ln^2 x). \tag{20}$$

The minus sign in Eq. (20) has the principle meaning since it leads to lowering of the quark helicity inside the proton. Joining Eq. (11a), Eq. (11b) and Eq. (20) we obtain the expression for the singlet distribution functions of sea quarks over the helicity inside the proton

$$q_+^s(x) = \frac{A_P}{x}(1-x)^7 + \frac{1}{x\ln^2 x}[B(1-x)^5 + C(1-x)^3], \quad (21a)$$

$$q_-^s(x) = \frac{A_P}{x}(1-x)^7 + \frac{2}{x\ln^2 x}[B(1-x)^5 + C(1-x)^3], \quad (21b)$$

where the first terms of these relations describe the Pomeron contribution as $x \to 0$ and the quark sea due to a perturbative gluon as $x \to 1$ and the second terms describe the contribution of the $A \otimes A$-cut as $x \to 0$ (Fig. 3) and instantons as $x \to 1$ (Fig. 1).

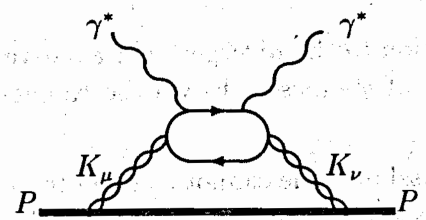

Fig. 3 The Kogut-Susskind pole contribution to the Compton effect off the proton

It should be noted that the powers of $1/x$ and $(1-x)$ in Eq. (21a) and Eq. (21b) reflect only the asymptotics as $x \to 0,\ 1$. In the intermediate region it is obviously necessary to take into account more complicated configurations of the proton wave function.

From Eq. (21a) and Eq. (21b) we obtain the contribution of the ABJ anomaly to the spin-dependent distribution function in the form

$$g_1^p(x) = -\frac{1}{x\ln^2 x}[B_1(1-x)^5 + C_1(1-x)^3] \quad (22)$$

and to the proton momentum

$$\Delta p^s = \int_0^1 dx\, \frac{3}{\ln^2 x}[B_1(1-x)^5 + C_1(1-x)^3]. \quad (23)$$

Unfortunately, the EMC data at low x : $x < 0.05$ have large errors; so it is impossible to determine the constants B_1 and C_1 separately from the data[2]. In this case we do the following. It is obvious that the anomaly dominates at small x and essentially affects the value of the integral of $g_1^p(x)$:

$$I_p = \int_0^1 dx\, g_1^p(x). \quad (24)$$

If the anomaly were absent, then the Ellis-Jaffe sum rule would be correct[3]:

$$I_p^{EJ} = 0.175 \pm 0.018 \quad (25)$$

which corresponds to the quark helicity:

$$\Delta\Sigma^{EJ} = 0.60 \pm 0.12. \quad (26)$$

Recently, the work[26] has appeared where the estimation was derived :

$$\Delta\Sigma^{KZ} = -\frac{2N_f}{11N_c - 2N_f}. \quad (27)$$

In [26] the circumstance has been used that the change of the axial charge is related with the motion of the Dirac sea levels in the field of nonperturbative fluctuations (for example instantons). Then, the magnitude of the proton axial charge should not depend on the manner of regularization of the Dirac vacuum. By using this independence and some natural assumptions in[26] the model independent value, Eq. (27), has been derived. For $N_f = 3$ one has

$$\Delta\Sigma^{KZ} = -0.22, \quad (28)$$

which does not agree even in sign neither with Eq. (26) nor with the EMC number, Eq. (2). We think that this discrepancy is due to incorrect extrapolation of the EMC data into the small x region.

In order to overcome this difficulty one should obviously take into account the anomaly contribution , Eq. (22). In addition, its contribution to the integral, Eq. (24), should be such that the value, Eq. (28) , of the quark helicity be reproduced. Further, it is easy to show that the value of the integral , Eq. (24) , equal to

$$I_p^{KZ} \approx 0.086 \quad (29)$$

corresponds to Eq. (28). Then, from the difference of Eq. (25) and Eq. (29) we find the contribution of the anomaly to the integral, Eq. (24),

$$I_p^A \approx -0.089. \tag{30}$$

In Fig 4 we present our prediction of $g_1^p(x)$ satisfying Eq. (22), Eq. (29), and Eq. (30) with the parametrization[†)]:

$$g_1^p(x) = -\frac{0.151}{x \ln^2 x}(1-x)^5 + 0.625(1-x)^{2.57}. \tag{31}$$

Here, the first term is the contribution of the anomaly; the second, which is regular over x, was determined from the fit of the EMC data in accordance with Eq. (31). This term corresponds to the valence quark contribution to $g_1^p(x)$.

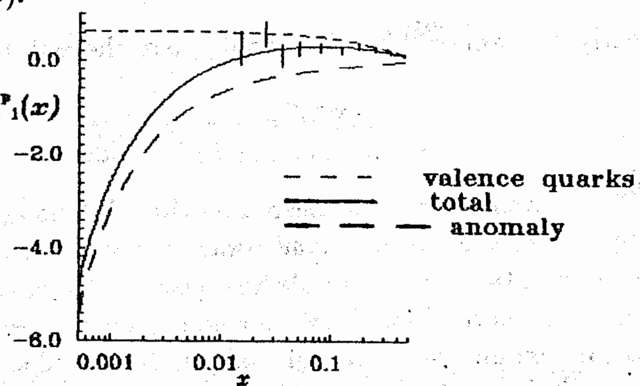

Fig. 4 Spin-dependent structure function $g_1^p(x)$. Continues line is the fit of EMC data by Eq. (31).

Under the parametrization, Eq. (31), the contribution of the anomaly to the proton momentum, Eq. (23), is equal to 2.7 % which is smaller than half a value of the proton momentum carried by sea quarks, Eq. (4), (the rest 4.7 % is the contribution of the perturbative sea). Here, we should stress that as $x \to 0$ only the anomaly contribution to Eq. (31) allows us to match two numbers Eq. (28) and Eq. (4).

[†)] The parametrization with the asymptotic form as $x \to 1$ $(1-x)^3$ gives essentially the same behavior.

Thus, the analysis of the distribution functions of sea quarks performed within the nonperturbative QCD points out the necessity of introducing new Regge trajectory related with the Adler-Bell-Jackiw anomaly and the Kogut-Susskind pole ("anomalon") in the QCD.

ANOMALON AND DYNAMICS OF $PP-$ AND $P\bar{P}-$ INTERACTIONS AT HIGH ENERGIES.

In the previous section we have shown that to explain the EMC data it is necessary to introduce new Regge trajectory with the intercept $\alpha_A(0) \approx 1$. It is natural that this trajectory should manifest itself in the processes of hadron-hadron scattering at high energies too. Now, the necessity of introducing the Regge trajectory with a high intercept additional to the Pomeron is intensively discussed[27, 28]. This is primarily connected with the necessity to explain the form of differential cross sections of elastic $pp-$ and $p\bar{p}-$ interactions at high transfers $|t| \geq 1\ Gev^2$. Namely, the experiment points out two salient facts: absence of a second diffractive minimum in pp interactions and existence of the shoulder in $p\bar{p}$ interactions at ISR energies and, second, independence of the differential cross section of energy at $|t| \geq 2\ Gev^2$.

Usually, these facts are explained by the contribution of the Odderon trajectory [28] with quantum numbers $\sigma = -1,\ P = -1,\ C = -1$. Within perturbative **QCD** the Odderon resembles the three-gluon exchange between hadrons. However, there arise some problems if the Odderon is applied to the data. First, if the Odderon intercept becomes higher than unity as in[28], then it increases with energy the difference of the total cross sections $\sigma_{pp}^{tot} - \sigma_{p\bar{p}}^{tot}$. Therefore, one needs an additional mechanism to suppress the Odderon contribution as $t \to 0$. However, a direct calculation[29] of the three-gluon exchange does not provide this suppression and, moreover, it produces the sign of the real part of the Odderon amplitude at $t = 0$ which does not agree with that needed to explain a large magnitude of the real part of the forward amplitude of $p\bar{p}$ scattering measured by the **UA4** Collaboration[30]. Second, at large t the radiation corrections induce an essential dependence of the elastic pp and $p\bar{p}$ cross section on energy $\sqrt{s}$[31], which is also not supported by experiment.

Using the "anomalon" we arrive at a more natural explanation of these data. First, due to the double-spin flip helicity amplitude induced by "anomalon" the latter does not contribute to the total cross sections and thus the difference $\sigma_{pp}^{tot} - \sigma_{p\bar{p}}^{tot} \to 0$ as $\sqrt{s} \to \infty$. Further, we expect a very small slope α'_A of the "anomalon". It is connected with that this slope is defined by the size of an instanton while the slope of the Pomeron is related to the confinement radius. Thus, we have simple estimation $\alpha'_A/\alpha'_P \approx (\rho_c/R_{conf})^2 \approx 10^{-1}$. If as usual $\alpha'_P = 0.2 - 0.3$, then we have $\alpha'_A = 0.02 - 0.03$. Therefore, we can neglet the slope of the "anomalon".

Thus, the "anomalon" contribution to the amplitude of pp ($p\bar{p}$) scattering may be written as

$$T_A(t,s) = \mp\, \gamma_A(t)(\frac{s}{s_0}), \qquad (32)$$

where the up (down) sign corresponds to pp ($p\bar{p}$) scattering. The residue in Eq. (32) should be related to the distribution of the axial charge inside the proton and therefore,

$$\gamma_A(t) \propto [G_A^{I=0}(t)]^2 \propto 1/[1 - t/M_A^2]^4, \qquad (33)$$

where $M_A^2 \approx 1.4\ Gev^2$[32].

The Pomeron amplitude corresponds to the expression:

$$T_P(t,s) = i\gamma_P(t)(\frac{s}{s_0})^{\alpha_P(t)} \exp[-i\frac{\pi}{2}(\alpha_P(t) - 1)] \qquad (34)$$

where the Pomeron residue is related with the electromagnetic form factor of a nucleon[27]:

$$\gamma_P(t) \propto [G_{em}(t)]^2 \propto 1/[1 - t/M_p^2]^4, \qquad (35)$$

where $M_p^2 \approx 0.71\ Gev^2$.

At large t the "anomalon" becomes dominant over the Pomeron ($M_A^2 > M_p^2$), which results in the observed change in the slope of the elastic cross sections of pp and $p\bar{p}$ scattering at $|\ t\ | \approx 1\ Gev^2$. Therefore, the absence of diffraction minima connected with multi-Pomeron exchanges is explained by the fact that they remain under the large contribution of the "anomalon" at $|\ t\ | \geq 2\ Gev^2$.

Owing to the absence of the double-spin flip amplitude of Pomeron it does not interfere with the "anomalon", and therefore, the differential

cross sections of $pp-$ and $p\bar{p}-$ interactions are

$$\frac{d\sigma}{dt} = \frac{1}{16\pi s^2}[|\ T_P\ |^2 + |\ T_A\ |^2]. \qquad (36)$$

In Fig. 5 fits of the total and differential cross sections at high energies with the Pomeron and "anomalon" trajectories are presented. As we see, there is a satisfactory agreement with experiment in the regions $|\ t\ | \leq 1\ Gev^2$ and $|\ t\ | \geq 2\ Gev^2$. Note here that a small slope of the "anomalon" trajectory, $\alpha'_A \approx 0$, is needed for explaining the energy independence of elastic cross sections of $pp-$ and $p\bar{p}-$ interactions at $|\ t\ | \geq 2\ Gev^2$. Essential deflection is seen only in the region of the dip in pp at ISR energies but it seems to be of a very complicated nature due to secondary Reggeons and their cut-offs[27]; and its dynamics is a subject of our forthcoming study.

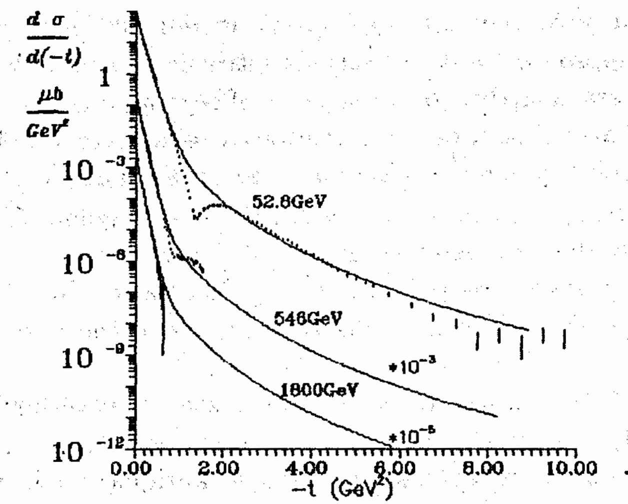

Fig. 5 Elastic pp and $p\bar{p}$ scattering at 52.8 GeV, 546 GeV, 1800 GeV. Data for $\sqrt{s} = 1800\ GeV$ are from[38]. The parameters are: $\alpha_P = 1.08,\ \alpha'_P = 0.3,\ M_P^2 = 0.97\ GeV^2,\ \gamma_P(0) = 21.8;$ $\alpha_A = 1,\ \alpha'_A = 0,\ M_A^2 = 1.08\ GeV^2,\ \gamma_A(0) = 0.57.$

In fact, at high energies two parts of structure functions of sea quarks, Eq. (21a) and Eq. (21b), correspond to two regions of transfers in the $pp-$ and $p\bar{p}-$ interactions. At $|\ t\ | \leq 1\ Gev^2$ the Pomeron dominates but at $| t | \geq 2\ Gev^2$ the "anomalon" does. The dominance of the "anomalon" at large transfers is clear since its contribution to the structure functions, Eq. (21a) and Eq. (21b), corresponds to the Fock state with large transfer momentum, Eq. (9).

CONCLUSION.

Within the instanton liquid model of the **QCD** vacuum the parametrization of the distribution functions of sea quarks inside the proton is derived. It should be stressed that the instanton mechanism of resolution of "spin crisis" differs entirely from the perturbative explanation[33] based on the contribution of the **ABJ**-anomaly on polarized gluons. Within the perturbative mechanism a very large value of the angular momentum $\Delta G \approx 5$ carried by gluons is required, which is in turn difficult to understand within the well-working constituent three-quark model of the proton. In addition, in this approach there are serious intrinsic problems[3, 34]. First, the answer is very sensitive to the manner of regularization of the ABJ-anomaly and the contribution of perturbative gluons to the first moment of $g_1^p(x)$ is probably zero[34]. Second, isotopic dependence of the anomaly[35] excludes its interpretation as the contribution of gluons to the proton spin. Note also that the perturbative gluon contribution to $g_1^p(x)$ would mean also the Pomeron contribution , Eq. (21a) and Eq. (21b), to this structure function , which is in disagreement with the selection rule, Eq. (14).

In our approach there is a natural way to generate negative helicity of sea quarks. Namely, the quark helicity is flipped in the field of strong vacuum fluctuation - instanton. In this case, a quark-antiquark pair with a large relative angular momentum arises to compensate for the changes of helicity by $2N_f$ since the spins of sea and valence quarks are opposite to the spin of the original quark. Thus, our mechanism leads to a completely definite orientation of rotation of the quark-antiquark cloud within a constituent quark. However, the total angular momentum carried by a quark is unchanged, and this is the reason for the good results of the constituent quark model in the description of static properties of hadrons.

The analysis of the spin-dependent distributions of sea quarks indicates the necessity of introducing a new Regge trajectory related with the Kogut-Susskind pole. The specific features of dynamics of $pp-$ and $p\bar{p}-$ interactions at large transfers and high energies also require a trajectory additional to the Pomeron which does not die out with growing energy.

Note that probably there exists a direct experimental confirmation of the "anomalon".[‡)] For instance, the **OMEGA** Collaboration[36] has observed that the differential cross section of reaction $\gamma p \to b_1(1^{+-}, 1235)p$ at $E_\gamma = 40 \div 70\ Gev$ is almost energy independent although in quantum numbers the Pomeron cannot produce the contribution in this reaction. On the other hand, the "anomalon" can result in the constancy of the cross section in energy and in the slope over t about twice as small as the slope of the diffractive cone in elastic pp scattering at these energies, which has been observed in experiment [36].

From our point of view the "anomalon" is also needed in order to explain anomalously large polarization fenomena in hadron-hadron processes at large transfers[37]. Thus, e.g., the three-anomalon vertex (3A) leads to the amplitude with single flip of helicity Φ_5 which does not die out with growing energy. Within the perturbative QCD this amplitude behaves as $m/\sqrt{s}$ where m is the quark current mass and therefore the perturbative QCD in principle cannot explain the anomaly in the scattering of polarized particles at high energies[37].

To provide a complete answer to the question of the existence of the new trajectory, experimental efforts should be made along the following directions: first, precision measurement of the DIS spin-dependent proton structure function $g_1^p(x)$ in the region of small x; second, the measurement of the differential cross section of $\gamma p \to b_1(1^{+-}, 1235)p$ at $E_\gamma >\ 70\ Gev$.

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Дорохов А.Е., и др. E2-91-375
Структурные функции морских кварков,
полюс Когута–Сусскинда и динамика pp-
и $p\bar{p}$-взаимодействий при высоких энергиях

Показано, что аномальное поведение спин-зависимой структурной функции $g_1^p(x)$ определяется вкладом в эту структурную функцию при $x \to 0$ новой траектории Редже, связанной с полюсом Когута–Сусскинда. Отмечено, что вклад новой траектории в процессах pp- и $p\bar{p}$-взаимодействия позволяет объяснить некоторые особенности этих реакций при высоких энергиях.

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Dorokhov A.E., Kochelev N.I., Zubov Yu.A. E2-91-375
Structure Functions of Sea Quarks,
Kogut-Susskind Pole and Dynamics of pp-
and $p\bar{p}$-Interactions at High Energies

It is shown that anomalous behavior of the spin-dependent proton structure function $g_1^p(x)$ is defined by the contribution to this structure function as $x \to 0$ of a new Regge trajectory caused by Kogut-Susskind pole. It is pointed out that manifestation of a new trajectory in pp- and $p\bar{p}$-interactions allows us to explain some peculiarities of these reactions at high energies.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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