Spin effects in instanton models

A.E. DOROKHOV

Joint Institute for Nucelar Researchs, Bogoliubov Laboratory of Theoretical Physics, 141980 Dubna, Russia

Nonperturbative nolocal structure of QCD vacuum is well described by instanton model. Specific helicity and flavor structure of zero modes of quarks in instanton field allows simultaneously to explain some important features of low- and high- energy hadron phenomemology. The basic characteristics of hadron spectrum, partonic sum rules, heavy-quark potential etc within the instanton liquid model are briefly discussed.

1 Introduction

Main features of *strong interaction* was defined long before the QCD has been discovered. They are quarks and strangeness; symmetries and color; spontaneous breaking of chiral symmetry; and, finally, confinement. QCD is the non-abelian gauge theory of quarks and gluons called for solving these problems. In the region of small distances (large energies) perturbative theory of QCD (pQCD) works well with very high precision and describes the scaling and its logarithmic violation. In particular, pQCD describes the evolution of parton distributions (experimental data rescaled to one scale) with Q^2 at high Q^2.

Still the main phenomena of strong interaction are not explicitly explained. There are problems with description of hadron spectrum, hadron distribution functions of quarks and gluons at intermediate energies, *etc.* Thus, the hadron processes at (asymptotically) high energies are described by pQCD and very low Q^2 region is fixed by symmetries of strong interaction (low energy theorems, chiral quark models, etc.)

The transition region between very low and high energies is very interesting aspect of modern experimental and theoretical searches. It is this region where different nonperturbative approaches was suggested: Lattice QCD, QCD sum rules, effective field models, etc. We are going to consider the instanton model of QCD vacuum and some of it applications.

2 The instanton model of QCD vacuum

To describe the strong interactions not only at short distances but also at medium and long distances one needs to understand how to compute the fundamental elements of the theory, its Green functions. This task may be done within the lattice QCD simulations [1]–[5] or effective QCD models like instanton model of QCD vacuum [6]. Lorentz invariance allows us to decompose the full quark propagator into Dirac vector and scalar pieces

$$S^{-1}(p^2) = Z^{-1}(p^2)[i\gamma \cdot p + M(p^2)]. \tag{1}$$

Asymptotic freedom means that, as $p^2 \to \infty$, $S^{-1}(p^2) \to i\gamma \cdot p + m$, (the free propagator) where m is the bare quark mass. The gluon propagator in the Landau gauge is given by

$$D^{ab}_{\mu\nu}\left(p^2\right) = \delta^{ab}\left(\delta^{\mu\nu} - \frac{p_\mu p_\nu}{p^2}\right) D\left(p^2\right) \tag{2}$$

with asymptotic large p^2 behavior $D\left(p^2\right) \to p^{-2}$. The explicit form of nonperturbative functions $M(p^2)$, $D\left(p^2\right)$ etc. may be found within the instanton vacuum model.

The vacuum gauge field is taken to be sum of the individual instanton fields with their centra at positions z_j. Each instanton is in the singular gauge

$$A^a_\mu(x)\frac{\tau^a}{2} = \frac{1}{g}\tau^a \eta^a_{\mu\nu}(x - z_j)_\nu \phi(x - z_j), \tag{3}$$

where $\eta^a_{\mu\nu}$ is the 't Hooft symbols. The ansatz for the profile function for the constrained instanton field[7] in the singular gauge is

$$\phi_I(x) = \frac{\overline{\rho}(x^2)}{x^2\left(x^2 + \overline{\rho}(x^2)\right)}, \tag{4}$$

The standard instanton solution is obtained with the constant size of instanton $\overline{\rho}^2(x^2) = \rho^2$. For the constrained instantons one uses exponentially-decreasing functions $\overline{\rho}^2(x)$, normalized as $\overline{\rho}^2(0) = \rho^2$. The constrained quark zero mode [8] is given by

$$\psi^\pm_{\text{sing}}(x) = \sqrt{2}\varphi_{\text{sing}}(x)\frac{\widehat{x}}{|x|}\chi^\pm, \qquad \varphi_{\text{sing}}(x) = \frac{\overline{\rho}(x^2)}{\pi(x^2 + \overline{\rho}^2(x^2))^{3/2}}, \tag{5}$$

where χ is a color Dirac spinor given by $\chi^\pm\overline{\chi}^\pm = (\gamma_\mu\gamma_\nu/16)\left(1 \pm \gamma_5\right)/2\tau^\pm_\mu\tau^\mp_\nu$ and $\tau^\pm_\nu = (\mp i, \vec{\tau})$, with the upper (lower) signs corresponding to solutions in the instanton (anti-instanton) field. We shell use the form [7]

$$\overline{\rho}^2(x^2) = \frac{2}{\Gamma(1/3)3^{1/3}}\left(\frac{\rho}{R}\right)^2 x^2 K_{4/3}\left(\frac{2}{3}\left(\frac{x^2}{R^2}\right)^{3/4}\right), \tag{6}$$

where $K_\nu\left(z\right)$ is the modified Bessel function. The specific feature of the constrained instanton is that at small distances it is close to the standard instanton profile of size ρ, and at large distances it has exponentially-decreasing asymptotics governed by a large-scale parameter R, such as $\rho < R$. These shapes are motivated by considering modifications of the instanton in the background field of large-scale vacuum fluctuations [7]. The constrained instanton profile, as opposed to the unconstrained one, provides the correct large-distance asymptotics of the quark and gluon field correlators [7].

The dynamical quark mass generated by the instanton background is expressed through the four-dimensional Fourier transform of the quark zero mode profile $\widetilde{\varphi}_{\text{sing}}(p)$ [9]

$$M(p^2) = Cp^2\widetilde{\varphi}^2(p^2), \tag{7}$$

where the constant $C > 0$ is determined from the gap equation [9, 10]

$$\int \frac{d^4k}{(2\pi)^4} \frac{M^2(k)}{k^2 + M^2(k)} = \frac{n}{4N_c}, \tag{8}$$

with $n = 0.0016$ GeV^4 [6] denoting the instanton density. The single instanton contribution to the gluon propagator

$$G^{CI,ab}_{\mu\nu}(p) \equiv \int d^4x e^{ipx} \left\langle 0 \left| A^{a,I}_{\mu}(x) A^{b,I}_{\nu}(0) \right| 0 \right\rangle_I = \delta^{ab} \left(\delta^{\mu\nu} - \frac{p^{\mu}p^{\nu}}{p^2} \right) G^{CI}(p), \tag{9}$$
$$G^{CI}(p) = -\frac{n_c}{N_c^2 - 1} p^2 \widetilde{\phi}^2(p^2),$$

is expressed through the Fourier transform of the instanton profile function $\widetilde{\phi}(p^2)$. Summing the contribution to the propagator of any number of instantons, which is analog of a self-energy resummation in perturbative theory, we get the dressed in the instanton vacuum gluon propagator in the Landau gauge

$$G^{ab}_{\mu\nu}(p) = \frac{\delta^{ab}}{p^2 + M^2_{CI}(p^2)} \left(\delta^{\mu\nu} - \frac{p^{\mu}p^{\nu}}{p^2} \right), \tag{10}$$

with a dynamical gluon mass

$$M^2_{CI}(p^2) = p^4 G^{CI}(p).$$

Important that the Green function of the effective model have correct large p^2 behavior and models nonperturbative dynamics at large distances. Also, the result of instanton model (7) and (10) are close to the results of lattice calculations [1]–[5] in the infrared region.

In general instanton induced interaction leads to effective fermion interaction which is for two flavors becomes

$$L_2 = \frac{2N_c^2}{n} \int \frac{d^4k_1 d^4k_2 d^4l_1 d^4l_2}{(2\pi)^{12}} \sqrt{M(k_1)M(k_2)M(l_1)M(l_2)}$$
$$\cdot \frac{\epsilon^{f_1 f_2} \epsilon_{g_1 g_2}}{2(N_c^2 - 1)} \left[\frac{2N_c - 1}{2N_c} (\psi^{\dagger}_{Lf_1}(k_1) \psi^{g_1}_L(l_1)) (\psi^{\dagger}_{Lf_2}(k_2) \psi^{g_2}_L(l_2)) \right.$$
$$\left. + \frac{1}{8N_c} (\psi^{\dagger}_{Lf_1}(k_1) \sigma_{\mu\nu} \psi^{g_1}_L(l_1)) (\psi^{\dagger}_{Lf_2}(k_2) \sigma_{\mu\nu} \psi^{g_2}_L(l_2)) + (L \to R) \right]. \tag{11}$$

This is so called effective 't Hooft vertex.

3 Instanton forces and spectroscopy.

The QCD vacuum has quite complicated structure. Conventionally the nonperturbative fields can be divided by two parts: short wave component which provides the interaction of quarks at small distances and long wave one which respects for the confinement. In the framework of instanton liquid model the first part is connected with short distance vacuum correlations where the single-instanton contribution with effective size $\rho_c \propto 1.5$–$2\,\mathrm{GeV}^{-1}$ dominates. Second component is related with long wave collective excitations of instanton liquid with wave length $\lambda \approx R \approx R_{\mathrm{conf}}$, where $R \approx 3\rho_c$ is average distance between instantons and $R_{\mathrm{conf}} \approx 5$–$6\,\mathrm{GeV}^{-1}$ is a confinement radius. The hadron model based on these assumptions was considered in [11, 12].

Within this model the hadron energy is

$$E_H = \frac{1}{R} \sum_{i=\mathrm{flavor}} N_i \omega\left(m_i R\right) + E_{\mathrm{vac}}, \tag{12}$$

where the first term is a sum of kinetic energy of quarks confined in the hadron bag with energy $\omega\left(m_i R\right)/R$ and the second term takes into the quark interaction with physical vacuum. The interaction with long-ranged vacuum fluctuations (condensates) has form of power corrections

$$E_{\mathrm{cond}} = -R^2 \sum_{i=\mathrm{flavor}} N_i C^{QQ}\left(m_i R\right) \left\langle 0 \left|\overline{q}_i q_i\right| 0\right\rangle + \sim \left\langle 0 \left|G^a_{\mu\nu} G^a_{\mu\nu}\right| 0\right\rangle R^3 + \ldots \tag{13}$$

and stabilizes the bag balancing the internal and external pressures $\partial E_H/\partial R = 0$. The kinetic energy and E_{cond} with stabilization condition generates the hadron mass scale

$$E_H = \frac{3}{2} N_q \left(\frac{2\pi\omega^2}{24\left(\omega - 1\right)}\right)^{1/3} \left|\left\langle 0 \left|\overline{q} q\right| 0\right\rangle\right| \approx \begin{cases} 750 \quad \mathrm{MeV} & \text{for mesons} \\ 1100 \quad \mathrm{MeV} & \text{for baryons} \end{cases}. \tag{14}$$

This estimate for mass scale of the nonstrange hadrons is consistent with QCD sum rule estimates [13].

The interaction of quarks through exchange by small size instantons due to specific structure of effective 't Hooft interaction (11)

$$\Delta E_{\mathrm{inst}} = \frac{\beta \rho_c}{R^3} \sum_{i \neq j} \frac{N_{ij} I\left(m_i, m_j\right)}{m_i^* m_j^*} \left(1 - \vec{\sigma}_i \vec{\sigma}_j\right) \tag{15}$$

produces spin-spin forces and solves the $U_A(1)$ problem

$$\Delta E^{\mathrm{p}}_{\mathrm{inst}} = -\frac{3}{2}\frac{\lambda_0}{R^3}, \qquad \Delta E^{\pi}_{\mathrm{inst}} = -\frac{\lambda_0}{R^3}, \qquad \Delta E^{\eta}_{\mathrm{inst}} = \frac{\lambda_0 - 4\lambda_s}{3R^3},$$

$$\Delta E^{\Delta}_{\mathrm{inst}} = 0, \qquad \Delta E^{\rho}_{\mathrm{inst}} = 0, \qquad \Delta E^{\eta'}_{\mathrm{inst}} = \frac{\lambda_0 + 2\lambda_s}{3R^3}.$$

By using standard vacuum parameters $\left\langle 0 \left|\overline{q}_i q_i\right| 0\right\rangle \approx -\left(250 \quad \mathrm{MeV}\right)^3$ and $\rho_c \approx 2$ GeV^{-1} the satisfactory results for ground state hadrons have been obtained [11].

4 The axial anomaly, the nucleon structure and the parton sum rules

The deep inelastic lepton-nucleon scattering processes (DIS) occurring at small distances characterize the internal structure of the elementary particles. In the past decade new experimental data with high precision and in large kinematic region has become available.

This is primarily a result of the SLAC–EMC–SMC–HERMES measurements [14, 15, 16, 17, 18] of helicities of the charged constituents of the proton and neutron. The EMC data analysis[15] has resulted in a striking conclusion: the sum of the helicities of the quarks inside a proton, $\Delta\Sigma$, was found to be extremely small, ($\Delta\Sigma \ll 1$), and the Ellis-Jaffe (EJ) [20] sum rule (SR) is strongly violated.

Then, it was concluded from NMC data analysis of the unpolarized structure function of nucleon $F_2(x, Q^2)$ [21, 19] that u-quark sea in the proton is suppressed w.r.t. d-quark sea, that is the Gottfried (G) SR [22] is violated, too. All this is in dramatical contradiction with the expectation of the naive parton model where all these sum rules are fulfilled. From the other side, the polarization experiments on neutron-contained targets confirmed that fundamental Bjorken (Bj) SR [23] is valid [17, 18].

In this part we want to argue [24] that the observed inconsistency in parton sum rules are a manifestation of nonperturbative structure of the QCD vacuum. Within the framework of this approach the breaking mechanism of QCD partonic SR is connected with a mixture of sea quarks with large transverse momentum in the nucleon wave function. This quark sea results from scattering valence quark off nonperturbative vacuum fluctuation, instanton.

This interaction in the limit of small size instanton is defined by the effective 't Hooft Lagrangian:

$$\mathcal{L}^{\text{inst}}(x) = (2n_c k^2)\Re\det(q_R q_L) \tag{16}$$

with the anomaly equation $\partial_\mu j^5_\mu = -2N_f(2n_c k^2)\Im\det(q_R q_L)$, where $n_c \approx 0.8\cdot 10^{-3}$ GeV^4 is effective instanton density, and $k = \frac{4\pi\rho_C^3}{3}\frac{\pi}{(m_*\rho_C)}$ is effective instanton-quark coupling. The value $\rho_c \approx 1$–$2\,\text{GeV}^{-1}$ defines the constituent quark radius, and $m_* = m - 2/3\rho_c^2 < 0\,|\bar{Q}Q|0>$ is the effective quark mass in the physical vacuum.

It's important that instanton induced interaction 16 changes the chirality of a quark by the value $\Delta Q = -2N_f$ and acts only for differently flavored quarks. From this it immediately follows that **sea quarks have negative helicity** and screen the helicity of a valence quark on which they are produced. On instanton the sea pair in the state with *Right chirality* is created and on the anti-instanton the quark pair with *Left chirality* is appeared.

Another thing is that on u-(d-)quark only $d\bar{d}$-($u\bar{u}$-) and $s\bar{s}$-quark sea is possible. Therefore there is **more d-sea quarks in the proton**. As a result it turns out that in the framework of the instanton mechanism the spin and flavor structure of nucleon quark sea is strongly correlated with the spin-flavor of the valence nucleon wave function.

Thus, specific helicity and flavor structure of quark zero mode interaction in the instanton field allows us simultaneously to explain the breaking of both Ellis-Jaffe SR related to significant breaking of quark helicity conservation and the Gottfried SR caused by the violation of the $SU_f(2)$-symmetry of quark sea.

I should note that a perturbative quark-gluon vertex does not flip the helicity neither does feel the flavor of the valence quark. Thus, within the perturbative QCD it is not possible, in principle, to explain the experimentally observed significant violation of both sum rules.

From the vertex 16 we obtain the anomalous instanton contributions to different PSR and axial charges [25]:

Gottfried sum rule

$$\Delta S_G = -\frac{2}{3}\left(\bar{d}-\bar{u}\right) = -\frac{2}{3}a;$$

flavor triplet and octet axial constants

$$\Delta g_A^3 = \Delta u - \Delta d = -\frac{10}{3}a_s;$$

$$\Delta g_A^8 = \Delta u + \Delta d - 2\Delta s = -4a + 2a_s;$$

flavor singlet axial charge

$$\Delta\Sigma_{\text{inst}} \equiv \Delta g_A^0 = \Delta u + \Delta d + \Delta s = -4a - 4a_s; \tag{17}$$

Ellis-Jaffe integrals for proton and neutron

$$\Delta S_{EJ}^p = \sum_p e_q^2 \Delta q/2 = -\frac{1}{9}(5a + 6a_s);$$

$$\Delta S_{EJ}^n = \sum_n e_q^2 \Delta q/2 = -\frac{1}{9}(5a + a_s);$$

Bjorken integral

$$\Delta S_{Bj} = \Delta S_{EJ}^p - \Delta S_{EJ}^n = -\frac{5}{9}a_s = \frac{1}{6}\Delta g_A^3$$

where a ($a_s \lesssim a/2$) is the probability to create nonstrange (strange) sea quark pair in instanton field. If we attribute all $SU_f(2)$ asymmetry of sea measured by the E866 Collab. [19]

$$\int_0^1 dx\, [\bar{d}(x) - \bar{u}(x)] = 0.118 \pm 0.012, \tag{18}$$

to the instanton contribution then we obtain the value for the coupling $a = 0.118 \pm 0.012$.

To estimate the (nonanomalous) valence quark contribution we shall use the quark model where relativistic effect reduces the helicity of quarks with respect to the nonrelativistic quark model result and take as a conservative estimation the value

$$\Delta u_v + \Delta d_v = 0.8 \pm 0.15 \tag{19}$$

From 17, 18, 19 we obtain the final result for the singlet axial charge of the proton:

$$\Delta\Sigma = -0.04 \pm 0.30.$$

Our theoretical error is the sum of indefinetness in experimental number for GSR for sea quarks and relativistic effects for valence quarks. Physically this compensation for the helicity of initial quark means a transformation of the valence quark spin momentum into the angular momentum of quark pair (in O^{++} state) created by instanton.

It should be stressed that in spite of that the instanton induced interaction contributes to g_A^3 it does not violate the Bjorken sum rule and violation of GSR is strongly correlated with contribution of axial anomaly to EJSR.

Another predictions of the instanton model for parton distribution integrals are:

1) Asymmetries in flavor structure of quark sea

$$\overline{d} - \overline{s} = 2\left(1 - \frac{3}{2}\frac{a_s}{a}\right)a; \qquad \overline{u} - \overline{s} = 2\left(1 - 3\frac{a_s}{a}\right)a; \tag{20}$$

2) New sum rule that is saturated by nonperturbative asymmetries in the nucleon sea

$$\frac{9}{25}\int_0^1 dx\left[g_1^p(x) - 6g_1^n(x)\right] + \frac{3}{2}\int_0^1 \frac{dx}{x}\left[F_2^{\mu p}(x) - F_2^{\mu n}(x)\right] = \frac{3}{5} \tag{21}$$

By using the E866 data for the GSR integral and the theoretical prediction for the BjSR integral [18]

$$\int_0^1 dx\left[g_1^p(x) - g_1^n(x)\right] = 0.181 \pm 0.003$$

the model predict the neutron integral of $g_1^n(x)$ as

$$\int_0^1 dx\, g_1^n(x) = -0.085 \pm 0.007,$$

that may be compared with world average experimental value [18]

$$\int_0^1 dx\, g_1^n(x) = -0.075 \pm 0.031.$$

3) Sea quark spin asymmetries

$$\Delta\overline{u} - \Delta\overline{d} = \frac{5}{3}a > 0; \qquad \Delta\overline{s} = -a_s < 0. \tag{22}$$

Thus we demonstrate that the instanton motivated nonperturbative sea fluctuations of nucleon sea are well describe the current experimental situation[1]).

[1]) In the papers [24, 25] the model of the sea quark distributions induced by instanton interaction has been considered.

5 The instanton effects on the heavy quark potential.

The gauge invariant potential between a very heavy quark and antiquark in a color-singlet state is given by [9, 26]

$$V(R) = -\lim_{T\to\infty} \ln \langle 0 | W(T,R) | 0 \rangle , \tag{23}$$

where $W(T,R)$ is a Wilson-loop $Tr\left(P \exp ig \oint_{C(R,T)} dz^\mu A_\mu(z)\right)$, with loop $C(R,T)$ being the rectangular closed curve with a spatial length R and a temporal length T. The origin of this expression is in the fact that during scattering heavy quark off the instanton its wave function changes as

$$\Psi'_Q(x) = \left[O(\vec{x}, \vec{p}, \vec{\sigma})\, U^{-1}(\vec{x})\right] \Psi_Q(x)$$

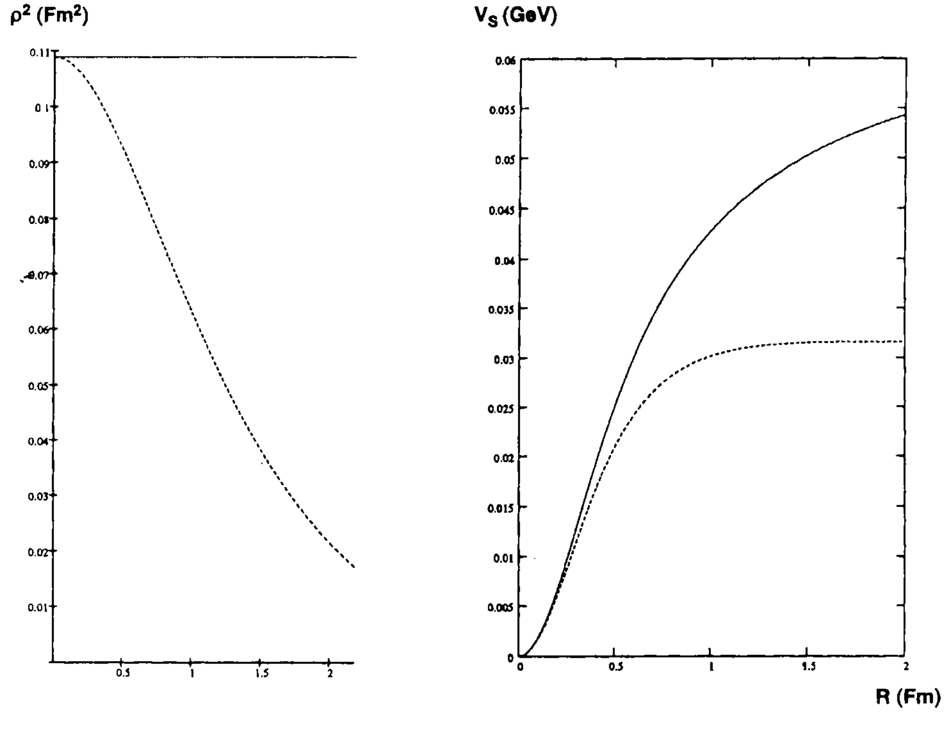

Fig. 1. Left: Instatnon size distribution. Right: Scalar interaction potential. The solid line is for instanton and the dashed one is for constrained instanton.

with phase factor

$$U^{-1}(\vec{x}) = P\exp\left[ig\int_{-\infty}^{\infty} dx_4 A_4(\vec{x},x_4)\right]$$

and spin-orbit operator

$$O = 1 - \frac{i}{m_Q}\vec{L}\cdot\vec{\nabla} - \frac{i}{2m_Q}\vec{\sigma}\cdot\vec{\nabla} - \frac{1}{2m_Q^2}\left(\vec{L}\cdot\vec{\nabla}\right)\left(\vec{\sigma}\cdot\vec{\nabla}\right) + \frac{1}{4m_Q^2}\vec{\sigma}\left[\vec{p}\cdot\vec{\nabla}\right].$$

Within the single-instanton approximation the phase factor may be reduced to

$$U^{-1}(\vec{x}) = \cos\alpha(\vec{x}) - i\tau^a n^a \sin\alpha(\vec{x}),$$

where

$$\alpha(\vec{x}) = \int_{-\infty}^{\infty} dx_4 \phi(\vec{x},x_4), \qquad n^a = \widehat{\vec{x}}.$$

Then the quark-quark scattering potential is given by potential

$$H_{QQ} = [-1 + O_{s-s} + O_{s-L}]\, V(\vec{x}_1 - \vec{x}_2)$$

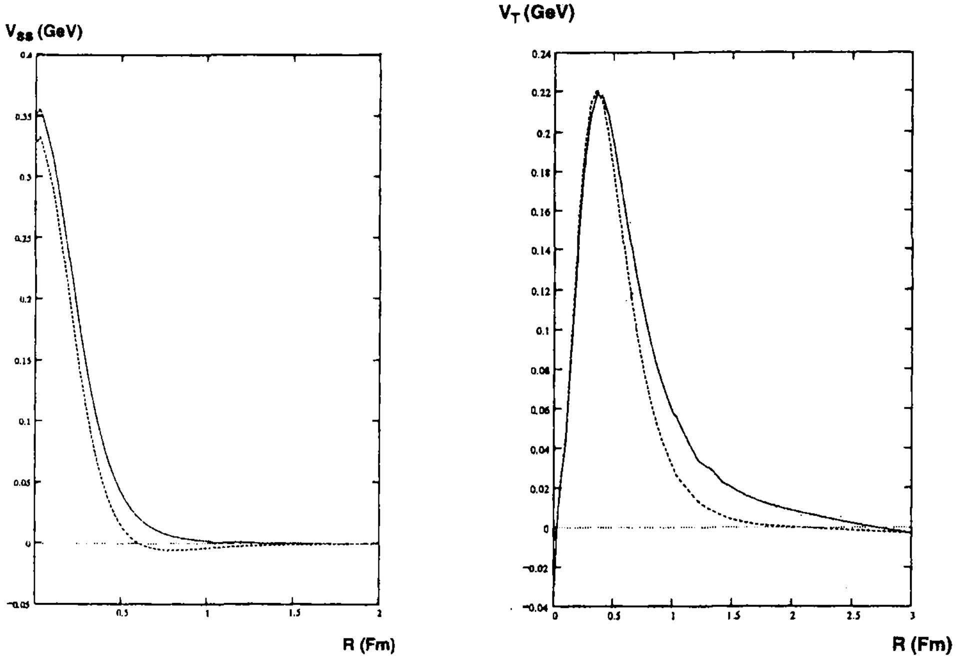

Fig. 2. Left: Spin-spin interaction potential. Right: Tensor interaction potential

with $V(R) = -2\int dn(\rho)\,\rho^3 W(R)$ and expansion in scalar part

$$W(R) = \frac{1}{N_c}\int d\vec{z}\,tr\left[1 - U\left(\frac{\vec{R}}{2} - \vec{z}\right)U^{-1}\left(\frac{\vec{R}}{2} + \vec{z}\right)\right],$$

spin-spin part proportional to $\vec{\sigma}\cdot\vec{\sigma}$

$$W_{s-s}(R) = \frac{1}{N_c}\left[2\frac{W'(R)}{R} + W''(R)\right]$$

and tensor part proportional to $\left[\vec{\sigma}_i\cdot\vec{\sigma}_j - \frac{\delta_{ij}}{3}\vec{\sigma}_i\times\vec{\sigma}_j\right]\vec{x}_i\vec{x}_j$

$$W_T(R) = \frac{1}{N_c}\left[\frac{W'(R)}{R} - W''(R)\right].$$

The instanton and constrained instanton model predictions are given in Fig. 1, 2.

6 Conclusion

In this talk we illustrate few examples of how instanton physics works in description of structure of QCD vacuum and light hadrons. In particular we demonstrate that instantons are responsible for nonlocal properties of vacuum condensates. Next, we show that the effective quark interaction due to instanton exchange gives to spin-spin splittings within the hadron multiplets. The same interaction leads to nontrivial spin-flavor structure of nucleon sea providing large violation of isovector symmetry of sea and essentially reducing the observed spin carried by charge constituents. Both results was confirmed experimentally. In the last part we considered the instanton effect on the potential between two infinitely heavy quarks. The instanton interaction gives only renormalization of heavy quark mass, but not to confinement of quarks. At the same time it provides small spin effects on heavy quark potential. All this show that the instanton physics is rather important part of nonperturbative aspects of QCD at intermediate energies and need further considerations.

We are thankful the organizing committee for invitation to make a talk and very nice condition for discussions.

References

[1] D. Becirevic et al.: Phys. Rev. **D 61** (2000) 114507.

[2] J.-I. Skullerud and A.G. Williams: Phys. Rev. **D 63** (2001) 054508.

[3] F.D.R. Bonnet et al.: Phys. Rev. **D 64** (2001) 034501.

[4] C. Alexandrou, Ph. de Forcrand and E. Follana: Phys. Rev. **D 63** (2001) 094504.

[5] A. Cucchieri, D. Zwanziger: hep-th/0012024.

[6] T. Schäfer and E.V. Shuryak: Rev. of Mod. Phys. **70** (1998) 323.

[7] A.E. Dorokhov, S.V. Esaibegyan, A.E. Maximov, S.V. Mikhailov: Eur. Phys. J. **C13** (2000) 331.

[8] I. Affleck: Nucl. Phys. **B191** (1981) 429; O. Espinosa: Nucl. Phys. **B343** (1990) 310.

[9] C. Callan, R. Dashen, D. Gross: Phys. Rev. **D17** (1978) 2717.

[10] D.I. Dyakonov and V. Yu. Petrov: Zh. Eksp. Teor. Fiz. **89** (1985) 751 [Sov. Phys. JETP **62** (1985) 431]; Nucl. Phys. **B272** (1986) 457.

[11] A.E. Dorokhov, N.I. Kochelev: Z. Phys. **C37** (1988) 377; A.E. Dorokhov, N.I. Kochelev, Yu. A. Zubov: Sov. J. Part. Nucl. **23** (1992) 522.

[12] E.V. Shuryak, J.L. Rosner: Phys. Lett. **B218** (1989) 72.

[13] V.M. Belyaev and B.L. Ioffe: Sov. Phys. JETP **56** (1982) 493 [Zh. Eksp. Teor. Fiz. **83** (1982) 876].

[14] SLAC–Yale Collaboration, G. Baum et al.: Phys. Rev. Lett. **51** (1983) 1135.

[15] The EMC Collaboration, J. Ashman et al.: Nucl. Phys. **B238** (1989) 1.

[16] The SMC Collaboration, B. Adeva et al.: Phys. Lett. **B302** (1993) 533; Phys. Lett. **B329** (1994) 399.

[17] The E142 Collaboration, P.L. Anthony et al.: Phys. Lett. Lett. **71** (1993) 959.

[18] SMC Collab.: Phys. Rev. **D58** (1998) 112002.

[19] E866 Collab., E.A. Hawker et al.: Phys. Rev. Lett.**80** (1998) 3715; G.T. Garvey, J.-C. Peng: nucl-ex/0109010.

[20] J. Ellis, R. L. Jaffe: Phys. Rev. **D9** (1974) 1444.

[21] P. Amaurdruz et al.: Phys. Rev. Lett. **66** (1991) 2712.

[22] K. Gottfried: Phys. Rev. Lett. **18** (1967) 1154.

[23] J.D. Bjorken: Phys. Rev. **148** (1966) 1467; *ibid* **D1** (1971) 1376.

[24] A.E. Dorokhov, N.I. Kochelev, Yu.A. Zubov: Int. Journ. Mod. Phys. **A8** (1993) 603; A.E. Dorokhov, N.I. Kochelev: Phys. Lett. **B304** 167 (1993).

[25] A.E. Dorokhov: in *Proc. of Quantum Field Theoretcal Aspects of High Energy Physics*, Kyffhauser, Germany (1993) 33; *Adelaide preprint* **ADP-93-224/T141** (1993).

[26] C. Callan, R. Dashen, D. Gross, F. Wilczek, A. Zee: Phys. Rev. **D18** (1978) 4684; E. Eichten, F. Feinberg: Phys. Rev. **D23** (1981) 2724.