New criteria of QCD-instantons search in deep inelastic scattering

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Abstract

QCD-instanton-induced deep inelastic scattering is considered. Arbitrary quark multiplicity is studied in contrast with previous investigations taking into account only minimal required number of quarks. Poisson distribution of number of quark pairs is obtained in addition to fixed quark multiplicity. Predictions for hadron spectrum done by means of Monte-Carlo method are discussed in connection with recent results on QCD-instantons search at HERA.

1 Introduction

Instantons [1, 2] playing an important role in QCD do not have direct experimental confirmation yet. A possibility to observe instantons in deep inelastic scattering (DIS) was first discussed in [3]. Systematic calculations of instanton-induced electron-proton DIS were initiated at HERA (Fig.1) [4]. At the parton level, these processes are characterized by high multiplicity (about 10 particles) in comparison with background processes [3, 4], presence of all light quark flavors and chirality violation [2], isotropy in instanton rest system [4]. Comparison of the instanton-induced and normal DIS final states was carried out with the help of Monte Carlo generator QCDINS [5]. Search strategy was based on isolating an instanton-enriched kinematic regions and observables by means of suitable cuts [6].

First results on QCD-instantons search at HERA are intriguing, but not conclusive. While the excess of events with instanton-like topology over the expectation of the standard DIS background is observed, it can not be claimed to be significant given the uncertainty of the Monte Carlo simulations of the standard DIS background [7]. Thus it is of importance to study QCS-instanton-induced DIS more deeply both experimentally and theoretically.

In this paper we calculate the corrections to the instanton processes coming from taking into account the arbitrary quark multiplicity. Note that the existing Monte Carlo package QCDINS [4] generating instanton-induced events implies Poisson distribution of gluons (with average value $< n_g >\sim 3$) and *fixed* quark (antiquark) multiplicity in final parton state:

$$q + g \to I \to < n_g > g + (2n_f - 1)q. \tag{1}$$

where only three quark flavors $n_f = 3$ in chiral limit are considered [1]. The presence of at least five quarks in final state is necessary to obtain a non-vanishing amplitude [2]. On

[1] production of heavy quarks are suppressed [4]

the other hand higher quark (antiquark) multiplicity is possible

$$q + g \to I \to < n_g > g + (2n_f - 1)q + additional\, quarks, \tag{2}$$

and, moreover, is not suppressed by small coupling constant (unlike perturbative corrections to gluon distribution).

Here we find the accurate distribution of quark multiplicities and analyze the consequences for hadron spectrum by means of slightly modernized QCDINS generator.

Let us consider Euclidean QCD [2] with only one massless quark flavor

$$S[A, \psi, \bar{\psi}] = \int dx \, \left(\frac{1}{2} Tr[F_{\mu\nu} F_{\mu\nu}] - \bar{\psi} i \not{D} \psi \right), \tag{3}$$

$$\not{D} = (\partial_\mu - igA_\mu)\, \gamma_\mu, \quad F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - ig[A_\mu, A_\nu], \quad A_\mu = A_\mu^a \frac{\lambda_a}{2}.$$

Group and spinor fermion indices are omitted for simplicity.

The process under consideration is instanton-induced quark-gluon collision with $2m-1$ quarks (antiquarks) in final state

$$q + g \to I \to (2m - 1)q. \tag{4}$$

Euclidean Green function for the process (4) reads

$$\int DAD\phi D\psi D\bar{\psi} e^{-S[A,\psi\,\bar{\psi}]} \prod_{i=1}^{m} \psi(x_i)\bar{\psi}(y_i) A_\mu(y) \approx \int d^4 z d\rho dU \mu(\rho) e^{-\frac{8\pi^2}{g^2}} \times$$

$$\times A_\mu^I (y - z) \int D\psi D\bar{\psi} e^{\int dx \bar{\psi} i \not{D}[A^{inst}]\psi} \prod_{i=1}^{m} \psi(x_i - z)\bar{\psi}(y_i - z), \tag{5}$$

where z_μ, ρ, U and $\mu(\rho)$ are instanton center, size, orientation and density correspondingly [2], $A_\mu^I(y)$ is instanton solution [1].

Here we focus attention on the fermion contribution

$$< \psi(x_1) ... \bar{\psi}(y_m) >^{inst} \equiv \int D\psi D\bar{\psi} e^{\int dx \bar{\psi} i \not{D}[A^{inst}]\psi} \prod_{i=1}^{m} \psi(x_i)\bar{\psi}(y_i) \tag{6}$$

which can be considered as a fermion Green function in the *external* instanton field.

The simplest non-vanishing Green function reads [2]

$$< \psi(x)\bar{\psi}(y) >^{inst} = \prod_n \int dc_n d\bar{c}_n e^{\lambda_n c_n \bar{c}_n} \sum_s c_s \psi_s(x) \sum_r \bar{c}_r \bar{\psi}_r(y) =$$

$$= det'(\rho)\psi_0(x)\bar{\psi}_0(y), \quad det'(\rho) = \prod_{n \neq 0} \lambda_n, \quad i\, \not{D}[A^{inst}]\psi_n(x) = \lambda_n \psi_n(x), \tag{7}$$

[2] we will present most formulas in Euclidean space. Analytical continuation to the Minkovsky space will be performed in the final step

where Grassman algebra is used. Explicit expressions for the zero mode $\psi_0(x)$ (eigenfunction belonging to $\lambda_0 = 0$) and $det'(\rho)$ are known [2].

Let us calculate 4-fermion Green function in the instanton field in the generalized approximation

$$< \psi(x_1)\bar{\psi}(y_1)\psi(x_2)\bar{\psi}(y_2) >^{inst} = \prod_n \int dc_n d\bar{c}_n e^{\lambda_n c_n \bar{c}_n} \psi(x_1)\bar{\psi}(y_1)\psi(x_2)\bar{\psi}(y_2) =$$

$$= \int dc_0 d\bar{c}_0 c_0 \psi_0(x_1) \bar{c}_0 \bar{\psi}_0(y_1) \prod_{n\neq 0} \int dc_n d\bar{c}_n (1+\lambda_n c_n \bar{c}_n)\psi(x_2)\bar{\psi}(y_2),$$

$$\prod_{n\neq 0} \int dc_n d\bar{c}_n (1+\lambda_n c_n \bar{c}_n)\psi(x_2)\bar{\psi}(y_2) = \sum_{k\neq 0}\Big(\int dc_k d\bar{c}_k c_k \psi_k(x_2) \bar{c}_k \bar{\psi}_k(y_2) \times$$

$$\times \prod_{n\neq 0,k} \int dc_n d\bar{c}_n \lambda_n c_n \bar{c}_n \Big) = \sum_{k\neq 0}\Big(\psi_k(x_2)\bar{\psi}_k(y_2) \prod_{n\neq 0,k}\lambda_n\Big) = S^{nz}(x_2,y_2)det'(\rho). \quad (8)$$

Here and below permutations of coordinates of the corresponding terms is implied. Propagator $S^{nz}(x,y)$ of non-zero fermion modes is known analytically [9]:

$$S^{nz}(x,y) \equiv \sum_{j\neq 0} \frac{\psi_j(x)\bar{\psi}_j(y)}{\lambda_j} = \frac{1}{\sqrt{1+\rho^2/x^2}} \frac{1}{\sqrt{1+\rho^2/y^2}} \Bigg[\frac{i(x-y)_\mu \gamma_\mu}{2\pi^2 (x-y)^4} \times$$

$$\times \left(1 + \rho^2 \frac{x_\nu \tau_\nu^- y_\kappa \tau_\kappa^+}{x^2 y^2}\right) - \frac{\rho^2}{4\pi^2 (x-y)^2 x^2 y^2} \times$$

$$\times \left(\frac{x_\nu \tau_\nu^- \gamma_\mu \tau_\mu^+ (x-y)_\lambda \tau_\lambda^- y_\omega \tau_\omega^+}{\rho^2 + x^2} \gamma_+ + \frac{x_\nu \tau_\nu^- (x-y)_\lambda \tau_\lambda^+ \gamma_\omega \tau_\omega^- y_\mu \tau_\mu^+}{\rho^2 + y^2} \gamma_- \right) \Bigg], \quad (9)$$

where $\gamma_\pm = (1 \pm \gamma_5)/2$, color matrices $\tau_\mu^\pm = (\vec{\tau}, \mp i)$ are defined in [9]. The propagator can be generalized to arbitrary instanton position and color orientation.

Finally, for the 4-fermion Green function we obtain

$$< \psi(x_1)\bar{\psi}(y_1)\psi(x_2)\bar{\psi}(y_2) >^{inst} = \psi_0(x_1)\bar{\psi}_0(y_1) S^{nz}(x_2,y_2) det'(\rho). \quad (10)$$

Higher Green functions are calculated similarly:

$$< \psi(x_1)\bar{\psi}(y_1)...\psi(x_m)\bar{\psi}(y_m) >^{inst} = \psi_0(x_1)\bar{\psi}_0(y_1) S^{nz}(x_2,y_2)...$$

$$...S^{nz}(x_m,y_m)det'(\rho). \quad (11)$$

The amplitude of the process (4) can be calculated by Lehman-Symanzik-Zimmermann technique on the basis of Green function (11):

$$A(k,k_1,p_1,...,k_m,p_m) = \int d\rho \int dU \mu(\rho) e^{-\frac{8\pi^2}{g^2}} \tilde{A}_\mu^{inst}(k) k^2 \bar{u}(k_1) k_1^2 \tilde{\psi}_0(k_1) \tilde{\bar{\psi}}_0(p_1)$$

$$\times p_1^2 u(p_1) det'(\rho) \prod_{j=2}^{m} \bar{u}(p) p_j^2 \tilde{S}^{nz}(p_j,k_j) k_j^2 u(k) \delta^{(4)}\Big(k + k_1 - p_1 - \sum_{n=2}^{m}(p_n+k_n)\Big), \qquad (12)$$

where k and k_1 denote 4-momenta of initial particles, $p_1, k_2, ..., k_m, p_m$ are 4-momenta of final fermions, tildes mean Fourier transformations of corresponding functions, δ-function appears due to the integration over instanton size z_μ. It is supposed that all momenta satisfy the mass-shell equation.

The contribution of zero quark modes as well as instanton gauge field are well known. Therefore, we restrict ourselves to non-zero mode contribution into the amplitude (12). Fourier transformation of propagator of non-zero modes is known [8]

$$\tilde{S}^{nz}(p,k) \equiv \int dxdy e^{-ipx-ipy} S^{nz}(x,y) = -2\pi^2\rho^2(p_\mu\gamma_\mu\gamma_+ + k_\nu\gamma_\nu\gamma_-)\frac{p_\lambda\tau^-_\lambda k_{\lambda'}\tau^+_{\lambda'}}{p^2k^2(pk)}, \qquad (13)$$

where we expand the expression in powers of ρ^2 and neglect terms of order $O(\rho^4)$, because only small-size instantons ($\rho \le 1/|Q'|$, $Q'^2 = -k_1^2$) can contribute significantly in DIS processes [4].

Calculating, squaring, integrating over momenta $k_1, \ldots, k_n, p_1, \ldots, p_n$ and normalizing (12) we obtain the *Poisson* distribution of number of quark pairs in a final state

$$P_n = \frac{<n>^n}{n!}e^{-<n>}, \qquad <n>= \frac{1}{2}\left(\frac{1-x'}{x'}\right)^2, \quad 0.5 < x' < 1, \qquad (14)$$

in addition to two quarks (one initial and one final) corresponding to zero modes, $x' = Q'^2/(2kk_1)$ is Bjorken variable of the instanton subprocess. Generalization into $n_f = 3$ leads to the distribution (14) for every light quark flavor in addition to usually considered six fermions.

In order to obtain quantitative contribution of additional quarks into hadron spectrum of instanton induced DIS we suppose $<n>= 1$ in the region $x' < 0.5$ for u, d, s flavors and correspondingly modify QCDINS2.0 generator [5]. The results of simulation which are given by Fig.1 and Fig.2 show the difference between some instanton DIS characteristics with and without account of additional quarks.

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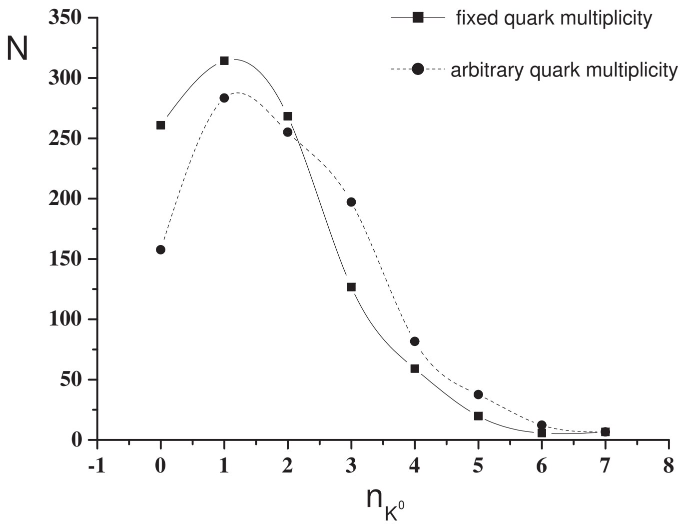

Figure 1: Number of instanton-induced events depending on final K^0 multiplicity (solid line corresponds to fixed quark multiplicity, dashed line corresponds to arbitrary quark multiplicity).

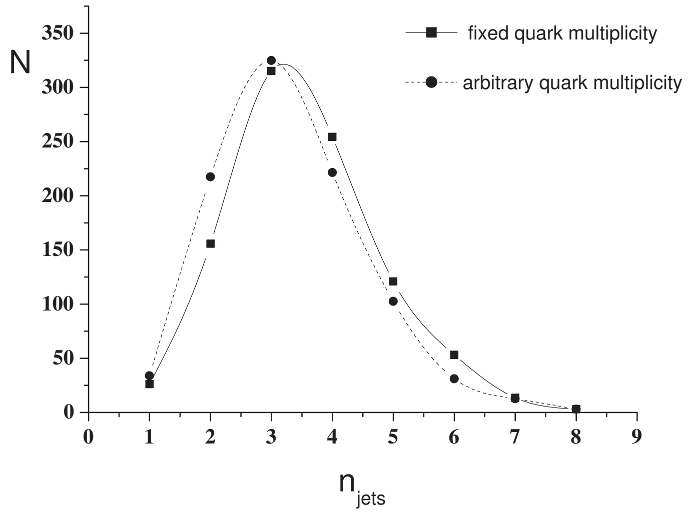

Figure 2: The same as in Fig.1, but for the dependence on number of hadron jets.