Instantons in the Large Q^2, x Region at HERA

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Abstract

A new mechanism for the excess of observed events above expectation in the large Q^2, x region at HERA is suggested. This mechanism results from a new type of instanton-induced quark-quark interaction, which is related to non-zero quark modes in the instanton field. We estimate the contribution of this interaction to the quark structure function F_2 using the gas instanton model approximation for the QCD vacuum. It is shown that this interaction can give a large contribution to F_2, especially at large values of Q^2. The strong dependence of this contribution on the quark masses and Q^2 is discussed, and it is concluded that a large charm quark contribution to events at high Q^2 can result from charm pair creation in the instanton field. We also obtain a sizeable contribution to the quark longitudinal structure function F_L a high Q^2. Experimental observables which can check this instanton mechanism in the high Q^2 region at HERA are suggested.

1 Introduction

Recently[1] the H1 and ZEUS collaborations at the electron proton collider HERA announced an intriguing excess of events above Standard Model expectation in the deep inelastic scattering (DIS) region of large Q^2 ($> 10000 GeV^2$) and large x. When confirmed by future data, this excess is likely to be the most exciting result in particle physics of the last years. Many interesting explanations of the this excess, mostly exploring physics scenarios beyond Standard Model, have been suggested (see for example [2]).

Within the Standard Model the largest uncertainty on the cross section prediction results from the assumed parton densities in the region of large x and large Q^2. These are based on two main ingredients. Firstly, the parton densities are parametrized with a relatively small amount of parameters at some low Q^2 scale, using a fit based on pQCD evolution equations of experimental data on lepton–parton and parton–parton cross sections at smaller values of Q^2. Secondly, to obtain the prediction for the parton densities at $Q^2 > 10000 GeV^2$, leading twist (DGLAP) evolution equations [4] are used.

A weakness of the first point was pointed out in [3] and [5]. In [3] it was shown that, not unexpectedly, the perturbative QCD evolution generates a feed down mechanism, i.e. effects at very large x and small Q^2 values propagate to smaller x at large Q^2 values. Hence e.g. an unexpected hard component to the partonic distributions at $x > 0.75$ and $Q^2 \sim 20 GeV^2$ has a visible effect in the region where HERA reports an excess. In [5] the possible importance of an intrinsic charm component in the proton wave function [6] to explain the excess was pointed out.

However, from our point of view, an important aspect of QCD has not been taken into account so far, namely the role of the complicated structure of the QCD vacuum related to the presence of strong nonperturbative fluctuations of gluon fields, called instantons [7]. In [8] a discussion was presented on a possible instanton contribution to the high Q^2 region at HERA. It was found that instanton induced interactions can enhance the expected event rate at high Q^2, resulting from multiple emission of gluons at the instanton vertex [9].

Instantons describe the tunnelling between different gauge rotated classical vacua in QCD field theory, and reflect the nonabelian character of the theory of strong interactions [1] . In QCD instantons play an important role in chiral symmetry breaking and in the origin of the masses of constituent quarks and hadrons. In recent reviews [11], [12] it was shown that instantons can reproduce fundamental quantities of the QCD vacuum, such as the values of the different quark and gluon condensates, and can also give a good description of the hadron spectrum.

Particularly interesting are the interactions which are connected with the so called quark–quark t'Hooft [13] interactions, which are induced by instantons. Four important features of these interactions should be stressed.

One feature is the flavor dependence of the interaction, which leads to large contributions in flavour singlet channels. For example, the quark–antiquark repulsion in the t'Hooft interaction for $SU(3)_f$ singlet $q\bar{q}$ states leads to a large enhancement of the η' mass and offers a solution to the $U_A(1)$ problem [13]. On the other hand, the t'Hooft interaction for quark-quark attraction in a flavor singlet state leads to the formation of an isoscalar diquark state inside the nucleon [14]. This has direct relevance to DIS: it offers a natural way to explain the experimental data on the ratio of proton to neutron structure functions $F_2^n(x, Q^2)/F_2^p(x, Q^2)$ at large $x > 0.1$, when assuming of the presence of a rather large isoscalar diquark component in the nucleon, which does not disappear for increasing Q^2 [15].

[1]For an introduction to the instanton physics see e.g. [10].

A second feature of the t'Hooft interaction is it's strong helicity dependence: contrary to the perturbative quark–gluon vertex, which conserves the quark helicity, instanton–induced quark–quark vertices exhibit a spin flip and consequently lead to quark helicity violation effects. In [16] it was shown that this contribution may explain the observed large violation of the Ellis–Jaffe sum rule for the first moment of spin–dependent structure function $g_1(x,Q^2)$ [17].

A further feature is the strong dependence of the t'Hooft interaction on the mass of the quarks involved. For the so called zero quark modes (zero energy states of quarks in the instanton field) the contribution behaves as $1/m_q^{*2} exp(-2m_q^* \rho_c)$, where m_q^* is the effective quark mass in the instanton vacuum and $\rho_c \approx 1.6 GeV^{-1}$ is the average instanton size in vacuum[11]. This leads to a large suppression of the contribution of zero quark modes to the quark distribution function for heavy flavours. A simple estimation in the framework of the instanton liquid model [11] shows that sea charm quarks are suppressed in F_2 by a factor 10^{-4} compared to light sea quarks, and can therefore be neglected.

A final feature is related to the strong dependence of the t'Hooft interaction on the virtuality of the interacting quark. The t'Hooft interaction results from the existence of zero quark modes in instanton fields, hence they are strongly localized in the field, which leads to an exponential dependence of their Green function on the quark virtuality.

One can calculate the instanton contribution to the sea quark distribution function by considering the interaction of a quark from a virtual photon with a valence quark in the nucleon through an instanton fluctuation (see Fig. 1). In this case all features of the instanton contribution are determined by the dependence of the instanton–induced quark–quark cross section on the virtuality of the incoming sea quark. Using the model of Landshoff [18] one can show that a fast reduction of the strength of this interaction with increasing sea quark virtuality results into a leading twist contribution in the sea quark distribution function [16] (see also [20]).

The exponential dependence of the quark-instanton vertex for zero quark modes on the incoming quark virtuality is expected to lead to a rather soft instanton–induced quark distribution as function of Bjorken x. This makes it difficult to explain a substantial excess at large values of x in the HERA region from quark zero modes in the instanton field (see [8]) alone.

The main goal of this study is to consider the contribution of so called *non-zero quark modes* in the instanton field (i.e. non-zero energy states of quarks in the instanton field) to the $F_2(x,Q^2)$ and $F_L(x,Q^2)$ structure functions, with special emphasis to the large Q^2 and x region at HERA. In most of the instanton calculations performed so far the contribution of non-zero quark modes has been neglected. It has been argued that this contribution is proportional to the product $m_q \rho_c$, which is indeed small for light u and d quarks. However, e.g. for strange quarks, this value amounts to $m_s \rho_c \approx 0.25$. Evidently for the heavy quarks c, b and t, the contribution of non-zero quark modes in the instanton field will dominate over the contribution from zero modes.

We show that the contribution of instanton interactions to unpolarized structure functions for these modes become increasingly important at large Q^2. Hence they can provide a QCD mechanism which leads to an any excess seen above calculations using parton distributions which result from 'standard' perturbative QCD evolution. For a quantitative prediction of these contributions we will use gas instanton model approximation for the QCD vacuum. In section 2 we calculate the instanton induced contribution to the quark structure functions, using the Green functions formalism. In section 3 we study the structure of the interaction using the effective Lagrangian formalism. The calculation of the non-zero modes contribution in framework of the liquid instanton model [11], [12] will be the subject of a future study.

2 Instanton induced contribution to the quark structure functions

To estimate the instanton induced contribution from non-zero quark mode states to the nucleon structure functions $F_2(x,Q^2)$, $F_L(x,Q^2)$ we consider the contribution to valence quark structure function which is given by the diagrams in Fig.1.

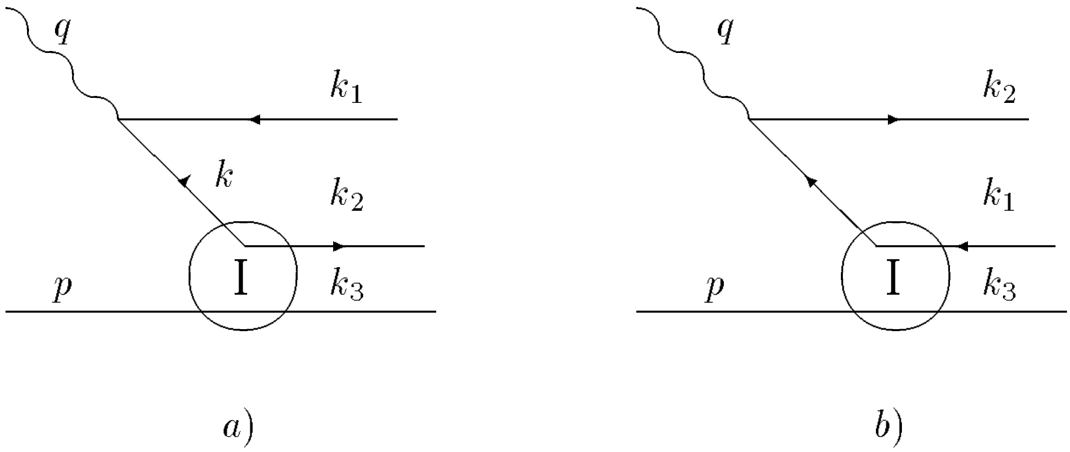

Figure 1: *The contribution from sea quarks created by instantons to the valence quark structure functions, a) the interaction of the sea quark with the valence quark and b) the interaction of the sea antiquark with the valence quark.*

This contribution is determined by the deviation of the quark Green function in the instanton field (i.e. the quark propagator in the background colour instanton field) from a free quark Green function $S_0(x,y)$.

$$S^{int}(x,y) = S_I(x,y) - S_0(x,y), \tag{1}$$

where $S_I(x,y)$ is quark Green function in instanton field and

$$S_0(x,y) = -\frac{\gamma\cdot(x-y)}{2\pi^2(x-y)^4}. \tag{2}$$

The function $S_I(x,y)$ for a quark with current mass m can be presented as the sum of the zero quark modes and non-zero quark modes contribution (see [11]) [2]

$$S_I(x,y) = S_z(x,y) + S_{nz}(x,y) = \frac{\Psi_0(x)\Psi_0^+(y)}{im} + \sum_{\lambda\neq 0}\frac{\Psi_\lambda(x)\Psi_\lambda^+(y)}{\lambda + im}, \tag{3}$$

[2] We will present most formulae in Euclidean space. Only in the final step of the calculation the analytical continuation to the Minkowskii space will be performed.

where

$$\Psi_{0i,\alpha}(x)\Psi^{+}_{0\,j,\beta}(y) = \frac{1}{8}\phi(x)\phi(y)(\hat{x}\gamma_\mu\gamma_\nu\hat{y}\gamma_-)_{ij} \otimes (U\tau^-_\mu\tau^+_\nu U)_{\alpha\beta}, \tag{4}$$

$$\phi(x) = \frac{\rho}{\pi}\frac{1}{\sqrt{x^2}(x^2+\rho^2)^{3/2}} \tag{5}$$

and

$$\begin{aligned} S_{nz}(x,y) &= \frac{1}{\sqrt{1+\rho^2/x^2}}\frac{1}{\sqrt{1+\rho^2/y^2}}(S_0(x,y)(1+\frac{\rho^2 U\tau^-\cdot x\tau^+\cdot yU^+}{x^2y^2}) \\ &- D_0(x,y)\frac{\rho^2}{x^2y^2}(\frac{U\tau^- x\tau^+\cdot\gamma\tau^-\cdot(x-y)\tau^+\cdot yU^+}{\rho^2+x^2}\gamma_+ + \\ &\frac{U\tau^-\cdot x\tau^+\cdot(x-y)\tau^-\cdot\gamma\tau^+\cdot yU^+}{\rho^2+x^2}\gamma_-)), \end{aligned} \tag{6}$$

where i, j (α, β) are the spinor (colour) indices, $\gamma_\pm = (1\pm\gamma_5)/2$, $\tau^\pm_\mu = (\vec{\tau}, \mp)$ is the colour matrix, U is the orientation matrix of instanton in the colour space, ρ is the instanton size, and

$$D_0(x,y) = \frac{1}{4\pi^2(x-y)^2} \tag{7}$$

is the propagator of a scalar quark.

The contribution of zero quark modes to the polarized structure function $g_1(x)$ and the unpolarized structure function F_2 has been estimated in [16] and [8]. It was shown that this contribution is rather large, specially at low x, can explain the violation of the Ellis-Jaffe sum rule and gives a contribution to the high Q^2 region at HERA, on top of the perturbative QCD contribution.

However, as discussed in section 1, in the large x and Q^2 region one can expect also large contributions from non-zero quark modes, which are not very strongly localized in the instanton field and therefore their contribution is expected to have a different x and Q^2 dependence compared to the one from zero quark modes.

We will estimate this contribution from non-zero quark modes for the case of two flavor quark $N_f = 2$, for the case where one quark will be heavy. In this case the dominated contribution is related to the one shown in the diagrams of Fig.1, where the sea quark is in a non-zero mode state and the valence quark is in a zero mode state.

The contribution to the valence quark part of the structure functions $F_2(x, Q^2)$ and $F_L(x, Q^2)$ is given by

$$\begin{aligned} F_2(x,Q^2) &= (-g_{\mu\nu} + 6x\frac{p_\mu p_\nu}{p\cdot q})xW_{\mu\nu}(q,p), \\ F_L(x,Q^2) &= 4x^2\frac{p_\mu p_\nu}{p\cdot q}W_{\mu\nu}(q,p), \end{aligned} \tag{8}$$

where p is the momentum of the valence quark, q is the virtual photon momentum, $x = Q^2/2p\cdot q$ and

$$W_{\mu\nu}(p,q) = \frac{1}{4\pi}\prod_{i=1}^{3}\int\frac{d^4k_i}{(2\pi)^3}(2\pi)^4\delta(k_i^2-m_i^2)\delta(q+p-k_1-k_2-k_3)T_\mu(p,q)T^*_\nu(p,q), \tag{9}$$

where $T_\mu(p,q)$ is the matrix element of the virtual photon-valence quark scattering.

The instanton–induced non-zero quark mode contribution in leading order in the mass expansion, for sea quarks with mass m_q, is

$$
\begin{aligned}
T_\mu(p,q) &= -ie_q \int_0^{\rho_{cut}} \frac{d(\rho)d\rho}{\rho^5} \int d^4y e^{-iqy} \bar{s}(k_2)\hat{k}_2(\bar{S}^{int}_{nz}(y,k_2)\gamma_\mu S_{nz}(k_1,y) + \\
&\bar{S}_{nz}(y,k_2)\gamma_\mu S^{int}_{nz}(k_1,y))\hat{k}_1 s(k_1)\bar{u}(k_3)\hat{k}_3 S_z(k_3,p)\hat{p}u(p),
\end{aligned} \tag{10}
$$

where $s(k_i)$ and $u(p)$ are Dirac spinors for sea quarks and valence quark respectively, and ρ_{cut} is a parameter which cuts the contribution from large instantons. The density of the instantons is

$$
d(\rho) = \frac{C_1}{2} e^{-3C_2+N_f C_3} \left(\frac{2\pi}{\alpha_s(\mu)}\right)^6 e^{-\frac{2\pi}{\alpha_s(\mu)}} (\rho\mu)^b \prod_{q=1}^{Nf} (m_q \rho), \tag{11}
$$

where

$$
b = b_0 + \frac{\alpha_s(\mu)}{4\pi}(b_1 - 12b_0), \tag{12}
$$

$$
\alpha_s(\mu) = \frac{4\pi}{b_0 \ln(\frac{\mu^2}{\Lambda^2})} \left\{ 1 - \frac{b_1}{b_0^2} \frac{\ln\ln(\frac{\mu^2}{\Lambda^2})}{\ln\frac{\mu^2}{\Lambda^2}} \right\},
$$

and in the $\overline{MS}$-scheme we have $C_1 = 0.466$, $C_2 = 1.54$ and $C_3 = 0.153$, where μ is the renormalization scale, and

$$
b_0 = 11 - \frac{2}{3}N_f \quad b_1 = 102 - \frac{38}{3}N_f. \tag{13}
$$

For numerical estimations below the value $\Lambda = 230MeV$ was taken.

The Fourier transform of the Green function for non-zero mode states for quarks is

$$
S_{nz}(y,k_2) = \int d^4x e^{ik_2x} S_{nz}(y,x) \tag{14}
$$

and antiquarks is $\bar{S}_{nz}(k_1,y)$ for the case $k_1^2 \to 0$, $k_2^2 \to 0$ has been calculated in [20]

$$
\begin{aligned}
S_{nz}(y,k_2)\hat{k}_2 &= -\frac{iy}{\sqrt{y^2+\rho^2}} e^{ik_2y} \left[1 + \frac{\rho^2}{2y^2} \frac{U\tau^- \cdot y\tau^+ \cdot k_2 U^+}{k_2 \cdot y}(1 - e^{-ik_2y})\right], \\
\hat{k}_1 \bar{S}_{nz}(k_1,y) &= \frac{iy}{\sqrt{y^2+\rho^2}} e^{ik_1y} \left[1 + \frac{\rho^2}{2y^2} \frac{U\tau^- \cdot k_1\tau^+ \cdot yU^+}{k_1 \cdot y}(1 - e^{-ik_1y})\right].
\end{aligned} \tag{15}
$$

The Fourier transformed Green function for zero quark modes in an instanton field equals

$$
\hat{k}_3 S_z(k_3,p)\hat{p} = \frac{\rho^2\pi^2}{2}(\gamma_\mu\gamma_\nu\gamma_-) \otimes (U\tau^-_\mu\tau^+_\nu U^+)F(p^2)F(k_3^2), \tag{16}
$$

where

$$
F(k^2) \approx e^{-\rho|k|} \to 1 \; at \; k^2 \to 0. \tag{17}
$$

The result of the calculation of (10), using (15), (16) at $k_i^2 \to 0$ is

$$
\begin{aligned}
T_\mu(p,q) &= -e_q\pi^4 \int_0^{\rho_{cut}} \frac{d(\rho)d\rho}{\rho} \bar{s}(k_2)\gamma_\mu s(k_1) F_I(k,k_1,k_2) \cdot \\
&\frac{\bar{u}(k_3)(\gamma_\mu\gamma_\nu\gamma_-) \otimes (U\tau^-_\mu\tau^+_\nu U^+)u(p)}{m_u},
\end{aligned} \tag{18}
$$

where $k = q - k_1 - k_2$ and

$$F_I(k,k_1,k_2) = I_1(k^2) + \frac{U\tau^- \cdot k_2\tau^+ \cdot kU^+}{k_2 \cdot k} I_2(k,k_1,k_2) + \frac{U\tau^- \cdot k_2\tau^+ \cdot k_1U^+}{k_2 \cdot k_1} I_3(k,k_1,k_2) + \frac{U\tau^- \cdot k\tau^+ \cdot k_1U^+}{k_1 \cdot k} I_4(k,k_1,k_2) + k_1 \leftrightarrow k_2,$$

with

$$\begin{aligned}
I_1(k^2) &= \frac{2}{\rho^2 k^4}\int_0^\infty \frac{dz z^3(z - \sqrt{z^2 - \rho^2 k^2})}{z^2 - \rho^2 k^2} J_1(z), \\
I_2(k,k_1,k_2) &= \rho\left(\frac{K_1(\sqrt{-(k+k_2)^2}\rho)}{\sqrt{-(k+k_2)^2}} - \frac{K_1(\sqrt{-k^2}\rho)}{\sqrt{-k^2}}\right), \\
I_3(k,k_1,k_2) &= \rho^2 k_2 \cdot k_1 \int_0^1 d\beta \frac{K_0(\sqrt{-(k+k_1+\beta k_2)^2}\rho) - K_0(\sqrt{-(k+\beta k_2)^2}\rho)}{2(k \cdot k_1 + \beta k_1 \cdot k_2)}, \\
I_4(k,k_1,k_2) &= \int_0^\infty \left(\frac{z}{-(k+k_1)^2(z^2 - \rho^2(k+k_1)^2)} + \frac{1}{k^2(z^2 - \rho^2 k^2)}\right) J_1(z).
\end{aligned} \tag{19}$$

To get the final result we have to integrate over the colour orientation dU which is given by

$$\begin{aligned}
\int dU U_{ij} U^+_{kl} &= \frac{1}{N_c}\delta_{jk}\delta_{li} \\
\int dU U_{ij} U^+_{kl} U_{mn} U^+_{op} &= \frac{1}{N_c^2}\delta_{jk}\delta_{li}\delta_{no}\delta_{mp} + \frac{1}{4(N_c^2-1)}(\lambda^a)_{jk}(\lambda^a)_{li}(\lambda^b)_{no}(\lambda^b)_{mp}.
\end{aligned} \tag{20}$$

In this paper only the colour singlet exchange between quarks induced by instantons is taken into account. The contribution of colour octet exchange is suppressed by a factor $1/N_c$ and will be studied in a forthcoming paper. For the colour singlet case the final result for the non-zero quark modes contribution to the quark structure functions F_2 and F_L, after averaging over the initial quark polarization and colour, equals

$$\begin{aligned}
F_2(x,Q^2) &= -\frac{2^9\pi^7 e_q^2 x}{27 m_u}\int dPS^3 \int_0^{\rho_{cut}} d\rho \frac{d(\rho)}{\rho} \bar{F}_I^2(k,k_1,k_2) t^2 \left(k_1 \cdot k_2 + \frac{12x^2 k_1 \cdot p k_2 \cdot p}{Q^2}\right), \\
F_L(x,Q^2) &= -\frac{2^{12}\pi^7 e_q^2 x^3}{27 m_u}\int dPS^3 \int_0^{\rho_{cut}} d\rho \frac{d(\rho)}{\rho} \bar{F}_I^2(k,k_1,k_2) t^2 \frac{k_1 \cdot p k_2 \cdot p}{Q^2},
\end{aligned} \tag{21}$$

where dPS^3 is the phase space integration, and $t^2 = (p-k_3)^2$, and

$$\bar{F}_I(k,k_1,k_2) = I_1(k^2) + I_2(k,k_1,k_2) + I_3(k,k_1,k_2) + I_4(k,k_1,k_2) + k_1 \leftrightarrow k_2, \tag{22}$$

The integration is performed in the center of mass frame of the sea quarks. In this frame the momenta of particles are (see for example [21])

$$\begin{aligned}
q &= (Q_0, 0, 0, E_q) \\
t &= (t_0, 0, 0, -E_q) \\
p &= E_p(1, \sin\beta, 0, \cos\beta) \\
k_{1,2} &= (E_k, \mp q_k \sin\theta\cos\Phi, \mp q_k \sin\theta\sin\Phi, \mp q_k\cos\theta),
\end{aligned} \tag{23}$$

where

$$\begin{aligned}
E_k &= \sqrt{\hat{s}}/2,\ q_k = \sqrt{\hat{s} - 4m_q^2}/2 \\
Q_0 &= (\hat{s} - Q^2 - t^2)/4E_k,\ t_0 = 2E_k - Q_0, \\
E_q &= \sqrt{(\hat{s} + Q^2 + t^2)/4\hat{s} - t^2},\ S_G = (q+p)^2, \\
E_p &= (S_G + Q^2 + t^2)/4E_k,\ \hat{s} = (t+q)^2, \\
\cos\beta &= (-S_G - Q^2 + (S_G + Q^2 + t^2)Q_0/2E_k)/(2E_pE_q).
\end{aligned} \tag{24}$$

For the phase space integration we have

$$dPS^3 = \frac{x}{128\pi^3 Q^2}\int_{S_0}^{S_G} d\hat{s} \int_0^{t^2{}_{max}} dt^2 \int_{t^2_{1min}}^{t^2_{1max}} \frac{dt_1^2}{\hat{s} - t^2 + Q^2}\frac{d\Phi}{2\pi}, \tag{25}$$

where $t_1^2 = (q - k_1)^2$, $S_0 = (2m_q)^2$, $S_G = Q^2(1-x)/x$, $t^2_{max} = -(S_G - \hat{s})/(1-x)$, and

$$t_1^{max,min} = -\frac{\hat{s} + Q^2 - t^2}{2}\left\{1 \pm \sqrt{(1 - \frac{4m_q^2}{\hat{s}})(1 + \frac{4Q^2t^2}{(\hat{s} + Q^2 - t^2)^2})}\right\}. \tag{26}$$

In the calculation we take for the renormalization scale $\mu = 1/\rho_{cut}$. There is a very strong dependence on the cutoff scale ρ_{cut} of the instanton density for most calculations on instanton induced contributions to physical quantities. From our point of view a natural way to derive an estimate for this parameter is by considering the confinement of the valence quarks inside the proton. For a zero mode valence quark in the instanton field we expect a cutoff factor $exp(-2|p|\rho)$, where $|p| \approx 300MeV$ represents the average virtuality of a valence quark in the proton. This leads to a value of $\rho_{cut} \approx 1.6GeV^{-1}$, which happens to coincide with the average size of the instantons in QCD vacuum (see [11]).

In Fig.2 we present the result of the calculation of the instanton induced contribution from the non-zero charm quark modes to the structure functions $F_2(x, Q^2)$ and $F_L(x, Q^2)$ at $x = 0.5$. This is approximately the x region where the HERA experiments have reported an excess at high Q^2. In the gas instanton model and for case $N_f = 2$, as studied here, the strange quark contribution to the F_2 and F_L structure functions from non-zero quark modes is expected to be about a factor 10^{-2} suppressed compared to the one of the charm quark, due to the smaller current quark mass. However, in the framework of the perhaps more realistic liquid instanton model for the QCD vacuum [11] an additional enhancement of strange quark production is expected, due to chiral symmetry breaking which leads to an increase of the effective quark mass.

In Fig.3 the result of the calculation of instanton contribution to the $F_2(x, Q^2)$ structure function is compared with the perturbative QCD contribution to the same structure function, due to gluon emission

$$F_2^{pert}(x, Q^2) = e_u^2 x \frac{\alpha_s(\mu)}{2\pi} P_{qq}(x) \ln(\frac{Q^2}{\mu^2}), \tag{27}$$

where

$$P_{qq}(x) = \frac{4}{3}\frac{1+x^2}{1-x}. \tag{28}$$

We see that the instanton contribution gives the same order correction to the parton distributions as the perturbative gluon correction. Therefore they should be taken into account in the Q^2 evolution of a proton $F_2^p(x, Q^2)$ structure function at a large x and Q^2.

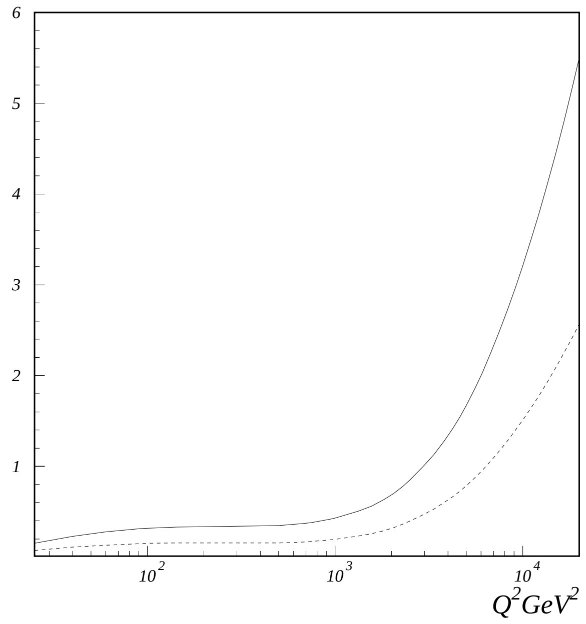

Figure 2: *The instanton–induced charm quark pairs contribution to valence quark structure function $F_2(x,Q^2)$ at $x = 0.5$ (full line) and $F_L(x,Q^2)$ (dashed line). The mass of charm quark is $m_c = 1.5GeV$.*

The result for the instanton induced non-zero quark modes contribution is proportional to the square of the mass of the quark which is in the non-zero mode state. In our calculation we have taken into account only terms of order m_c^2 but not explicitely the mass of the sea quark in it's Green function [3]. At small $Q^2 \approx 20GeV^2$ this contribution is rather small due to phase space restriction. But for large Q^2 the contribution to both $F_2(x,Q^2)$ and $F_L(x,Q^2)$ becomes very large due to the strong Q^2 dependence of this contribution. It is interesting to note that the non-zero quark modes lead to a sizable contribution to the $F_L(x,Q^2)$ structure function at large Q^2. This results from the structure of the Green function for quarks in the instanton field which includes a part that comes from to the Green function of scalar quarks (see (6)).

Therefore the charm quark contribution induced by instantons to structure functions is large. We also expect some contribution from another heavy quarks b, t due to the same mechanism. Experimentally this effect could be verified by studying the evolution of the charm component of the structure function. A considerable increase with increasing Q^2 is predicted.

While this study was obviously stimulated by the excess of the high Q^2 events seen at HERA, it should be noted that independently of it's confirmation in future, this type of contribution can play an important role in QCD analyses of very high Q^2 data, and should therefore be considered. In a future paper we plan to make more complete study of its implications for QCD studies.

[3]Threshold effect due to the quark masses are taken into account in the phase space integration.

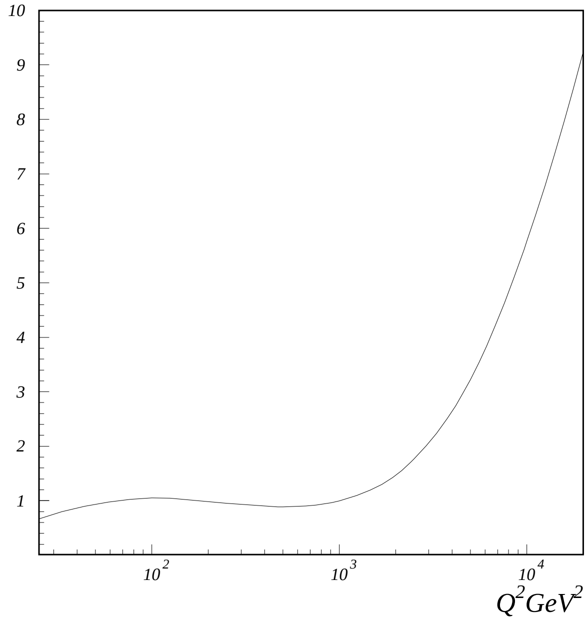

Figure 3: *The ratio of the instanton induced charm quark pairs contribution to the perturbative QCD contribution, to valence quark structure function* $F_2(x, Q^2)$ *at* $x = 0.5$.

3 Effective quark-quark interaction induced by non-zero quark modes in instanton field

In the previous section a striking Q^2 dependence of non-zero quark modes contribution to the structure functions has been found. This Q^2 dependence results from the structure of quark-quark interaction vertex induced by non-zero quark modes in the instanton field. Let us consider an instanton induced two quark interaction for the case when one of the incoming quarks is off-shell, a situation which is directly related to the sea quark contribution to the structure function. One part of the interaction is related to the t'Hooft interaction induced by quark zero mode states. The effective vertex can obtained by amputating of the four-point zero modes quark Green function in instanton field

$$\mathcal{L}_z = \int \frac{d\rho d(\rho)}{\rho^5} \bar{u}(k_2)\hat{k}_2 S_z(k_2, k_1)\hat{k}_1 u(k_1)\bar{d}(k_3)\hat{k}_3 S_z(k_3, p)\hat{p} d(p). \tag{29}$$

After substitution the Fourier transformed zero modes (16) and performing the integration over the instanton orientation in colour space dU, one obtains the following effective Lagrangian for the t'Hooft interaction

$$\mathcal{L}_z = \int \frac{d\rho d(\rho)}{m_u m_d \rho^5} (\frac{4\pi^2\rho^2}{3})^2 (\bar{u}_R(k_2) u_L(k_1) \bar{d}_R(k_3) d_L(p)(1+\frac{3}{8}(1-\frac{3}{4}\sigma^u_{\mu\nu}\sigma^d_{\mu\nu})t^a_u t^a_d + (R \longleftrightarrow L))) F(|k_1|\rho), \tag{30}$$

where $F(|k_1|\rho) \approx exp(-|k_1|\rho)$ for $\rho \to 0$. The density of the instantons is proportional to the product of the quark masses (11) and therefore, in spite of the quark masses in the denominator of Eq. (30), the final result is finite in the limit $m_u, m_d \to 0$.

The effective Lagrangian connected with non-zero modes contribution to leading order in m_q, and for a small instanton size $\rho \to 0$ can be obtained in a similar way

$$\mathcal{L}_{nz} = \int \frac{d\rho d(\rho)}{\rho^5} \bar{s}(k_2)\hat{k}_2 \left\{S_{nz}(k_2,k_1) - S_0(k_2,k_1)\right\} \hat{k}_1 s(k_1)\bar{d}(k_3)\hat{k}_3 S_z(k_3,p)\hat{p}d(p). \quad (31)$$

The final result is

$$\begin{aligned}\mathcal{L}_{nz} &= \int \frac{d\rho d(\rho)}{m_d \rho^5}(\frac{4\pi^2\rho^2}{3})^2 \Big\{(1+\frac{3}{32}t_s^a t_d^a)(\overline{F_1^n}(t^2)+\overline{F_2^n}(k^2,t^2)) \\ &- \frac{9}{32}\frac{t_s^a t_d^a k_2^\tau k_1^\sigma \sigma_{\tau\sigma}{}^d}{k_2\cdot k_1}\overline{F_2^n}(k^2,t^2)\Big\} \otimes \bar{s}(k_2)\hat{k}_1 s(k_1)\bar{d}(k_3)d(p), \quad (32)\end{aligned}$$

where $t^2 = (k_1 - k_3)^2$ and the form factors

$$\begin{aligned}\overline{F_1^n}(t^2) &= \frac{2}{\rho^2 t^4}\int_0^\infty dz \frac{z^2(z-\sqrt{z^2-\rho^2 t^2})}{\sqrt{z^2-\rho^2t^2}} J_1(z), \\ \overline{F_2^n}(k_1^2,t^2) &= \int_0^\infty dz z(\frac{1}{k_1^2\sqrt{z^2-k_1^2\rho^2}} - \frac{1}{t^2\sqrt{z^2-t^2\rho^2}})J_1(z). \quad (33)\end{aligned}$$

A very interesting feature of this interaction (32) is that it is zero for on-shell massless quarks but it is large for the case when one of the quarks is off-shell. This is the origin of the striking Q^2 dependence of the contribution of this interaction to the DIS structure functions, as found in the previous section. Another feature of the non-zero mode interaction, as compared to the zero mode interaction (31) is the helicity flip. The t'Hooft interaction leads to a double spin flip of quarks and therefore this interaction can give a large contribution to the double spin asymmetry in the polarized DIS (see [16]). The contrary is true for non-zero quark mode interactions, which do not have a definite helicity and can lead only to a single spin flip; their contribution to the double spin asymmetry is *zero* . Hence these instanton induced interactions do not contribute to the spin-dependent structure function $g_1(x,Q^2)$. The zero quark mode interactions on the other hand do contribute to $g_1(x,Q^2)$, and this contribution is *negative* [8]. Thus the total contribution of instantons to $g_1(x,Q^2)$ is negative, contrary to the perturbative QCD expectation which yield a *positive* value for $g_1(x,Q^2)$ at large x and Q^2. Hence the instanton mechanism for DIS events excess at HERA can be checked by a measurement of the double spin asymmetry for these events at high Q^2, when a polarized proton beam at HERA will be available. The zero quark mode component will contribute with a negative asymmetry, while the non-zero quark mode is predicted to reduce the asymmetry, compared to the positive asymmetry predicted from perturbative QCD.

4 Summary

In summary, a new contribution to the proton structure at high Q^2 and high x is suggested. This contribution is connected with complicated structure of the QCD vacuum, related to the existence of strong fluctuations of the vacuum gluon fields: instantons. It is shown that a new type of quark-quark interaction in QCD, which is due to non-zero quark mode states in instanton field, gives a large heavy quark contribution to the quark structure functions in the large Q^2, x region. This effect could already be visible at HERA.

The large Q^2 dependence of this interaction comes from the dependence of the effective instanton–induced vertex on the quark virtuality. We also predict anomalous spin properties of high Q^2 events induced by this new interaction. Due to the strong dependence of this interaction on the quark virtuality it should show up also in other lepton-hadron and hadron-hadron processes with large transfer momentum. Note that this contribution is missing in standard QCD analyses. Independently of a possible excess at large Q^2 and x of the HERA data this contribution probably should be considered when evolving to large scales. Finally we note that enhancement of the charm contribution to the structure function at large x can lead also to an explicit QCD model for intrinsic charm [6] in the proton.

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