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INSTANTONS AND THE PROTON'S AXIAL CHARGE*

S. FORTE

I.N.F.N., Sezione di Torino

AND

Dipartimento di Fisica Teorica, Università di Torino
via P. Giuria 1, I-10125 Torino, Italy

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We show that non-perturbative contributions to the nucleon matrix elements of quark and gluon operators may explain the surprising experimental results of the EMC collaboration on the nucleon's axial charge. We discuss the phenomenological consequences of this way of understanding the data, and we argue that recent experimental results on the Gottfried sum rule may be understood in the same way.

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1. Introduction

In the past few years a series of deep-inelastic scattering experiments performed by the EMC collaboration (now the NMC collaboration) has produced results in contradiction with naive parton model expectations. In particular, the experimental determinations of the first moment of the polarized structure function $g_1(x)$ of the proton [1] has lead to the rather unexpected conclusion that the axial charge of a polarized nucleon is compatible with zero, the so-called "proton spin problem". This result has spawned a very ample theoretical literature, and it is now understood that the data are by no means in contradiction with QCD, contrary to the naive expectation.

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However, very few dynamical explanations of the experimental numbers are available. Here we shall discuss one such explanation [3], based on non-perturbative QCD. We shall argue that the EMC data [1] are giving us a hint of the influence of the topological structure of the QCD vacuum on hadron structure. We shall also propose that some very recent data on the flavor content of the nucleon [4] may be understood in a similar guise.

Indeed, these results have superficially several common features: in both cases the data deviate strongly from naive parton model expectations, and in both cases the quantity which is measured, although having a simple and straightforward parton model interpretation, cannot be expressed as the matrix element of a conserved current, and thus in neither case is the naive parton model a consequence of QCD; finally, in both cases theoretical and experimental arguments suggest that the source of trouble may reside in nonperturbative QCD effects. Here we shall sharpen this superficial analogy: we shall argue that if the matrix elements relevant to both sets of data are calculated in an instanton vacuum then the nonperturbative induced quark-instanton coupling provides contributions which are at least qualitatively in agreement with the observed effects.

First, we shall recall the so-called spin problem (comprehensive reviews are in Ref. [2]), related to the polarized scattering data [1], and its explanation in terms of instanton effects [3]; finally, we shall discuss the phenomenological consequences of an instanton explanation, and we shall argue that it may be relevant to the understanding of the Gottfried sum rule data [4].

2. The spin problem and instantons

The experimental data [1] on the proton structure function g_1 allow — in a rather indirect way (see Ref.[2]) — a determination of the matrix element

$$\Delta\tilde{q}s^{\mu} = \langle p, s|j_5^{\mu}|p, s\rangle \tag{1}$$

of the isosinglet axial current $j_5^{\mu} = \sum_i i\bar{\psi}_i\gamma^{\mu}\gamma_5\psi_i$ in a proton state with momentum p and spin s. Physically, $\Delta\tilde{q}$ is a measure of the constituent's axial charge:

$$h\Delta\tilde{q} \equiv \langle p, h|Q_5|p, h\rangle\,, \tag{2}$$

$$Q_5 \equiv \int d^3x \sum_{i=1}^{N_f} i\left(\bar{\psi}_i\gamma^0\gamma_5\psi_i\right) \tag{3}$$

in a polarized proton state $|p, h\rangle$ with helicity $h \equiv \hat{s}\cdot\hat{p}$. If $|p, h\rangle$ were a superposition of free, massless fermions (as in the most naive parton model), then $\Delta\tilde{q} = 1$, whereas the experimental value of $\Delta\tilde{q}$ is compatible with zero.

However, both perturbative [6] and nonperturbative [7,8] arguments show that in general the matrix element of the axial charge Q_5 is the sum of the quark chirality operator Q_5^q and a gluonic operator Q_5^{an} :

$$Q_5 = Q_5^q + Q_5^{an}. \tag{4}$$

The operator Q_5^{an} , whose mixing with the quark chirality operator is induced by the axial anomaly[1], is a functional of the gluon fields, which can be fully determined only through nonperturbative methods [8]

$$Q_5^{an} = \frac{g^2}{4\pi^2} \int d^3x \,\mathrm{tr}\, \epsilon^{ijk}(A_i \partial_j A_k + \frac{2}{3} g A_i A_j A_k) \tag{5}$$

in terms of the gluon operators A^i and the fermion-gluon coupling g^2.

When evaluated on gluonic partons, Q_5^{an} is proportional to the gluon helicity operator Δg: $Q_5^{an} = (\alpha_s/2\pi)N_f \Delta g$, a result that may be obtained by evaluating [6] the contributions to $\Delta\tilde{q}$ (1) at order α_s in perturbative QCD. The Altarelli-Parisi equation indicates that the leading order evolution of Δg is $O(1/a_s)$, implying that the perturbative gluon contribution to (1) is not necessarily small [3]. This may reconcile the experimental result with the parton model expectation that the quark parton polarization $h\Delta q \equiv \langle p, h|Q_5^q|p, h\rangle$ should be sizable, provided Δg has the required magnitude. Whether this is actually the case rests with future experiments; however, bounds from the unpolarized gluon structure function (the total gluon helicity cannot exceed the total number of gluons), as well as phenomenological parameterization of the polarized structure function suggest [10] that it might be hard to reproduce the large value of Δg required, and hence perturbative QCD may not be the entire story.

Whatever the nature of the fields that contribute to the expectation value of Q_5 (3) — gluonic partons, sea quarks, or nonperturbative gluon configurations — the reason for the rather miraculous cancellation which leads to a vanishing (or almost vanishing) value of $\Delta\tilde{q}$ must be of nonperturbative nature. Let us now show that such a cancellation occurs in a simplified model. Namely, we show that in QCD with only one quark flavor, assumed to be massless, and with a topologically nontrivial vacuum structure, the axial charge of a free polarized quark vanishes because it is entirely screened by the vacuum gluon field fluctuations. In parton language, we consider a polarized massless quark, say right-handed, and we compute

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[1] For a review see Ref. [9].

[2] Eq.(5) holds in the gauge $A_0 = 0$ after fixing the gauge with respect to homotopically nontrivial gauge transformations, otherwise an additional nonlocal term is present on the r.h.s. in order to guarantee its gauge invariance [8].

its left-handed quark and antiquark content induced by the vacuum gluon field fluctuations: we find that, in the one flavor case, it is equal to one, *i.e.*, that the total chirality (which for massless quarks is equal to the helicity) of the given ("valence") quark, plus its sea quark and antiquark content vanishes.

We compute

$$h\Delta\tilde{q}_0 \equiv \langle 0|a(p,h)Q_5 a^\dagger(p,h)|0\rangle\,, \tag{6}$$

where $a^\dagger(p,h)$ is the creation operator for a quark with momentum p and helicity (or equivalently chirality) h, and $|0\rangle$ is the QCD vacuum state. The latter is a gauge invariant superposition of vacuum functionals peaked around the topologically nontrivial gauge copy of the n-th homotopy class $|n\rangle$ of the perturbative (trivial) vacuum; it is assumed to be approximated by a superposition of field configurations that tunnel between state $|n\rangle$ and state $|n\pm 1\rangle$ (see Refs [9, 11]). An instanton is an explicit example of a tunneling field: in the sequel, however, we shall not need its explicit form and indeed although we will speak of instantons for definiteness we shall only assume the tunneling picture of the QCD vacuum.

In order to compute $\Delta\tilde{q}_0$ (6) we use the decomposition (4) of the axial charge in a quark and a gluon contribution. Because Eq.(4) holds when quantizing quarks first in the gluon background [8], we ought to average over quarks first, then calculate the expectation value of the remaining gluon operators in the effective gluon states thus obtained. The quark contribution is trivial to compute: because $a^\dagger(p,h)|0\rangle$ is a helicity eigenstate it is also an eigenstate of Q_5^q with eigenvalue equal to h. Thus the contribution of Q_5^q to $\Delta\tilde{q}_0$ is equal to one. The expectation value of Q_5^{an} is also easily evaluated [8] for a tunneling configuration: it is a time-dependent function that varies between 0 and $2n$ for a configuration that tunnels between the k-th and the $k+n$-th field sector. In the case of a one-instanton (anti-instanton) $n=1$ ($n=-1$).

Because the vacuum contains the same number of instantons and anti-instantons the vacuum expectation value of Q_5^{an} vanishes (more precisely, since Q_5^{an} is a CP-odd operator $\langle Q_5^{an}\rangle \propto \sin\theta$ [8] where θ is the vacuum angle). The matrix element of Q_5^{an} in a one-quark state, *i.e.*, the contribution of Q_5^{an} to the matrix element (6), however, need not vanish because of the correlation between instantons and the quark created by $a^\dagger$: the forward (in momentum space) matrix element of Q_5^{an} which appears in (6) is just the time average of this time-dependent function, evaluated over all the configurations of instantons and anti-instantons which are contained in the polarized one-quark state:

$$\langle 0|a(p,h)\, Q_5^{\rm an}\, a^\dagger(p,h)|0\rangle = \lim_{T\to\infty}\frac{1}{T}\int_{-T/2}^{T/2} dt\, \langle \Psi_1[A]|Q_5^{\rm an}(t)|\Psi_1[A]\rangle. \quad (7)$$

The expectation value of the gluonic operator $Q_5^{\rm an}$ in Eq.(7) is computed in the gluon state $|\Psi_1[A]\rangle$ obtained from the one-quark state after averaging over quark fields, *i.e.*,

$$\langle \Psi_1[A]|Q_5^{\rm an}(t)|\Psi_1[A]\rangle \equiv \int DA\, \exp\Big(i\int \Big(d^4x\, \mathcal{L}_0[A] + \mathcal{L}_{\rm eff}[A]\Big)\Big) Q_5^{\rm an}(t)$$

$$\exp\left(i\mathcal{L}_{\rm eff}\right) = \bar{u}(p,h)\not p \frac{\delta^2}{\delta J \delta \bar{J}}\Big[\int D\psi\, D\bar{\psi}$$

$$\times \exp\left(i\int d^4x\, \left(\bar{\psi} i\not D[A]\psi + \bar{J}\psi + \bar{\psi}J\right)\right)\Big]\Big|_{J,\bar{J}=0} \not p u(p,h)$$

$$= \bar{u}(p,h)\not p\, [i\not D[A]]^{-1} \not p u(p,h)\,, \quad (8)$$

where $\mathcal{L}_0$ is the usual (Maxwell) pure gauge gluon Lagrangian, and $u(p,h)$ is the Dirac spinor wave function for the given massless quark state.

Thus, $\mathcal{L}_{\rm eff}$ is just the quark propagation amplitude in the given gluon background, *i.e.*, the vacuum gluon fields which we approximate with instantons. Now, in the zero mass limit the propagator in an instanton background is dominated by the zero mode of $i\not D[A]$, which has definite chirality, and the propagator in an instanton background reduces to a projector on the zero mode wave function. Viewed as a coupling between a quark-antiquark pair and an instanton this is a helicity-flipping interaction: a right-handed (left-handed) incoming quark is turned by an instanton (anti-instanton) in an outgoing left-handed (right-handed) one. This restrict the possible arrangements of instantons and anti-instantons that appear in the state functional $\Psi_1[A]$ to those depicted in Fig. 1 alternating sequences of instantons and anti-instantons, with an even number of couplings, and with the ordering of the instanton-anti-instanton pair fixed by the initial quark chirality. Any other arrangement receives weight zero by $\mathcal{L}_{\rm eff}$ and does not contribute to the expectation value (8).

In order to compute the time average (7) the Feynman diagrams of Fig. 1 must be evaluated at zero momentum transfer; this enforces time ordering along the quark line. The function $Q_5^{\rm an}(t)$ in an instanton background was computed exactly in Ref.[12]; however, we need only to know that if the tunneling takes place from t_1 to t_2 (notice that in the gauge $A_0 = 0$

Fig. 1. Arrangements of instantons and anti-instantons with nonvanishing coupling to a massless quark propagation amplitude. ⊕ indicates coupling to an instanton and ⊖ to an anti-instanton *via* the 't Hooft vertex (15). L and R denote the chirality of the fermion.

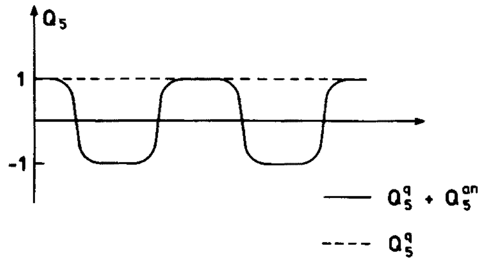

Fig. 2. The behavior of the axial charge $Q_5 = Q_5^q + Q_5^{an}$ as a function of time for a chain of instanton–anti-instanton pairs coupled to a right-handed source. The dashed line indicates the canonical source polarization Q_5^q. Notice that the time average $\overline{Q_5^{an}} = -Q_5^q$, implying $\overline{Q_5} = 0$.

tunneling always takes place in time) then

$$Q_5^{an}\left(\left(\frac{t_f - t_i}{2}\right) + t\right) - Q_5^{an}\left(\frac{t_f - t_i}{2}\right) = -\left[Q_5^{an}\left(\left(\frac{t_f - t_i}{2}\right) - t\right) - Q_5^{an}\left(\frac{t_f - t_i}{2}\right)\right]. \tag{9}$$

This follows from the O(4) symmetry of the instanton, which must be also a symmetry of any tunneling configuration which does not violate the Lorentz invariance (isotropy) of the vacuum state. Eq.(9) implies

$$\int_{t_i}^{t_f} dt\, Q_5^{an}(t) = (t_f - t_i) Q_5^{an}\left(\frac{t_f - t_i}{2}\right). \tag{10}$$

The shape of $Q_5 = Q^q + Q_5^{an}$ as a function of time along the diagrams of Fig. 1 is thus as given in Fig. 2 (for a right-handed source). Using Eq.(10) it follows immediately that Q_5 when evaluated along the diagrams of Fig.1 averages to zero

$$\frac{1}{T}\int_{-\frac{T}{2}}^{\frac{T}{2}} dt\, Q_5(t) = 0, \tag{11}$$

implying

$$\Delta\tilde{q}_0 = 0 \tag{12}$$

which is the claimed result.

The cancellation between the time average of Q_5^{an} and the constant canonical value of Q_5^q which leads to the vanishing result Eq.(12) holds diagram by diagram for all the contributions to the expectation value (7)–(8) where at least one instanton–anti-instanton pair couples to the quark line. However, since the matrix element is evaluated in the forward direction (in momentum space) in position space it should be computed in the limit $T \to \infty$; then, however small the instanton density, instantons contribute necessarily to the expectation value, *i.e.*, in this limit the diagram with no instanton couplings gives a vanishing contribution. Therefore, in the one-flavor case, for one single massless quark the axial charge form factor $\Delta\tilde{q}_0$ in the forward direction vanishes exactly.

Although everything so far has been done by quantizing quarks in the gauge background — which naturally provides the two-component interpretation (4) of the axial charge, also found in perturbative QCD [6] and current algebra [7] computations — one may also proceed in the usual way, and quantize gluon first. A determination of the cancellation (12) in this case requires computation of the full three-point function $\langle a(p,h) j_5^\mu a^\dagger(p,h)\rangle$ [3]. Without entering into the details of this computation, let us observe that in this picture one first path-integrates over instantons, then one computes the axial charge of the quark in presence of the effective instanton-induced fermion interaction ('t Hooft interaction). This leads again to the diagrams of Fig.1. Then, the oscillating charge shown in Fig.2 is actually the quark's axial charge, rather than the sum of a constant quark charge and an anomalous charge carried by gluons. The oscillations are due to the fact that the 't Hooft interaction flips the quark's chirality by creation of polarized sea quark-antiquark pairs (see, *e.g.*, [11]).

Although the vanishing of the axial charge persists, in this picture it is attributed to a cancellation between valence and sea quark polarization. In parton model terms, the diagrams of Fig. 1 lead to a left-handed quark content of a right-handed quark. Notice that since the vector charge is conserved, this opposite chirality content of a polarized quark is actually due to (anomalous) production of sea quark-antiquark pairs. The latter are the Fock states associated to the zero mode of the Dirac operator in the instanton background [12]. They cannot be written as a superposition of perturbative partons [8], in that their quantum numbers cannot be expressed as a linear superposition of quantum numbers carried by perturbative hard partons (see below).

Similar ambiguities in separating the axial charge in a quark and a gluon part are found in the perturbative approach [2].

3. Phenomenology of the instanton contribution

What is the relevance of the computation of the previous Section to the EMC result, and if any, how can instanton effects be tested experimentally? In order to relate the result of Section 2 to the physical case one must relax three simplifying assumptions: take massive quarks, let the number of flavors be greater than one, and compute the axial charge of a three-quark bound state, rather than a single free quark. Unfortunately, only the first approximation is quantitatively under control, while for the remaining one we can only give heuristic arguments.

Corrections due a nonzero quark mass would appear in the effective Lagrangian (8). The new terms proportional to quark masses have different helicity structure: for example, the linear term in m is helicity preserving. However, any correction of this kind will appear weighted by a power of the (current) quark mass divided by the quark energy, which in the present case is for light quarks of order 10^{-3}. Nevertheless, strange quark corrections may be non negliglible.

The most interesting complication appears when one lets the number of flavors $N_f > 1$. Then, the effective Lagrangian (8) is given by a complicated $2N_f$-fermion vertex [11], and the result depends critically on the structure of the instanton vacuum [13]. First, consider the case of a single quark propagating in an instanton medium which is so dilute that the vacuum is chirally symmetric. Then, the extra $2N_f - 2$ quark lines can only be contracted with each other. The helicity structure of the interaction prevents tadpole contraction of lines emanating from the same vertex; however, lines that come from an instanton vertex may be contracted with those from an anti-instanton vertex, leading to diagrams as those in Fig. 3(a). This generates strong correlations between instantons. Because each fermion line in Fig. 3(a) connecting two vertices depends on the instanton–anti-instanton separation x as $1/x^3$, quasiparticles tend to bind into instanton–anti-instanton "molecules" [14]. In this regime the time dependence of the axial charge no longer has the symmetric form of Fig. 2, rather, it is equal to the charge of the source quark when this propagates between molecules, and it tunnels to the opposite value only when the quark propagates inside an instanton molecule (see Fig. 4). On the other hand, the variation of $Q_5^{\rm an}$ in an instanton field is now $\pm 2N_f$, rather than 2. It follows that there are two competing effects: on the one hand, the axial charge created by tunneling is larger, but on the other hand tunneling is weighted by a small number proportional to the density of molecules.

If instead the instanton ensemble is very dense and random, so that the pseudoparticle positions are completely uncorrelated, then the multi-fermion 't Hooft interaction effectively reduces to the two-fermion one [13,14], leading back to the $N_f = 1$ case. This is because if the instanton density is large enough, the vacuum quark condensate acquires a nonvanishing value. Then, the extra fermion lines may go into this condensate, leading to the diagrams of Fig. 3(b). Now, although the physical instanton vacuum is not chirally symmetric, numerical simulations [14] suggest that the physical case is somewhat intermediate between these two limiting pictures. Although the chiral symmetry is spontaneously broken, the correlations between positions of pseudoparticle are rather strong. While the diagrams of Fig.3(b) lead to the contribution calculated in Section 2, the diagrams of Fig. 3(a) are present, too. Unfortunately, we cannot provide a firm estimate of the average Q_5^{an} for the diagrams of Fig. 3(a), which as we argued would depend on the details of the instanton molecules, and *a fortiori* we cannot determine the relative importance of diagrams Fig. 3(a) and Fig. 3(b).

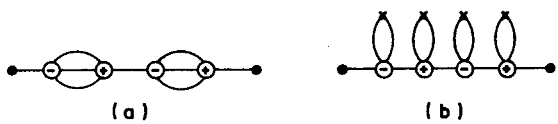

Fig. 3. Some of the diagrams that contribute to polarized quark propagation in an instanton background in the three-flavor case. (a) Chirally symmetric vacuum. (b) Chirally asymmetric vacuum. × indicates contraction with the vacuum chiral condensate.

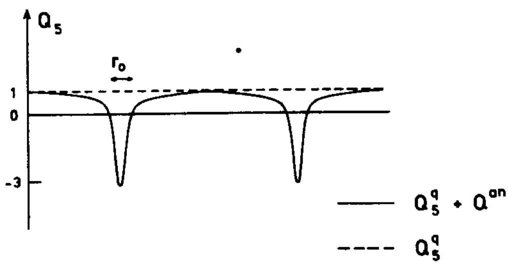

Fig. 4. Shape of Q_5^{an} for quark propagation in the background of small instanton molecules when $N_f = 2$; r_0 indicates the size of the molecule. For generic values of N_f the dip extends to $-(2N_f - 1)$, rather than to -1.

Finally, if we consider propagation of a multiquark state, then quarks may interact with each other through the same instantons which contribute to the charge, thereby making an analytic computation of Q_5^{an} impossible. However, we may assume that in the deep-inelastic regime each quark is independently probed. Although physically plausible, this picture cannot be supported by the usual perturbative arguments since the instanton fields are nonperturbative in the first place.

Short of a conclusive argument to show that the cancellation found in the one-quark, one-flavor case carries over to a physical nucleon (which, if it exists, will be presumably based on some elusive symmetry of the fermion-instanton effective Lagrangian), we may take a phenomenological approach, and see which would be the experimental signals of such an explanation.

First, it should be noticed that the anomalous instanton contribution to the axial charge is totally unrelated to angular momentum or spin. Indeed, since the instanton field is rotationally invariant it carries vanishing angular momentum; also, the gluon helicity operator averages to zero in a one-instanton state [8]. Thus the contribution to Q_5^{an} from the instanton is totally unrelated to the gluon helicity. Notice that this is true *only* for contributions to the charge that are due to tunneling in the space parametrized by the topological charge Q_5^{an}, which in turn are due to nonperturbative instanton-like gluon configurations. Other contributions, related to the perturbative mixing induced by the anomaly between quark and gluon operators [6], just correspond to conversion of gluon helicity into quark helicity. The same result is found if we look at the diagrams of Fig.1 in the pure quark picture discussed at the end of the previous section. Then, the vanishing of the total axial charge is due to cancellation of the valence quark's charge by that of the sea quark-antiquark pairs created by the anomalous instanton-induced 't Hooft interaction. This interaction is anomalous because it does not conserve chirality, *i.e.*, the axial charge. However, angular momentum *is* conserved, hence, the sea contribution to the axial charge *is not* a contribution to the angular momentum. Indeed, the fermion states created because of the anomaly are those associated to the zero modes of the Dirac operator, which carry vanishing angular momentum and helicity.[3]

This means that, at least in principle, an anomalous contribution due to tunneling would appear as a violation of the spin sum rule:

$$\tfrac{1}{2}\Delta\tilde{q} + \Delta g + L_z \neq \tfrac{1}{2}, \qquad (13)$$

(where Δg is the gluon helicity and L_z the orbital angular momentum of quarks and gluons) because the correct sum rule is

$$\tfrac{1}{2}\Delta q + \Delta g + L_z = \tfrac{1}{2}, \qquad (14)$$

[3] This is not in contradiction with the fact that the zero mode is a chirality eigenstate: the identification of helicity and chirality holds only for free Dirac spinors. The zero mode carries vanishing energy and momentum and it is a superposition of spin-$\frac{1}{2}$ states with opposite polarization [12]. It does show, however, that the instanton contribution cannot be understood in terms of perturbative partons, which are helicity and chirality eigenstates, and for which the identification of chirality and helicity holds.

where $\Delta q = \langle p, h|Q_5^q|p, h\rangle$ differs from $\Delta\tilde{q}$ by the addition [8] of the perturbative contribution proportional to the gluon helicity, and the nonperturbative tunneling contributions discussed above. For instance, in the extreme case in which $\Delta g = L_z = 0$ then the r.h.s. of Eq. (13) is just zero. Of course a precision test is totally out of the question; nevertheless, Δg is an experimentally accessible quantity [2]. A very small value of Δg would be an indication in favor of a nonperturbative explanation.

A second remarkable feature of the instanton contribution to the axial charge is that it is due to the classical gauge field configuration about which gluons are quantized. Thus, it is unaffected by quantum evolution and it is a renormalization group invariant. This is consistent with the fact that it might be reinterpreted in parton language as a quark contribution, since in the parton model the quark and gluon contributions are associated to the two eigenvalues of the Altarelli-Parisi kernel, respectively, as a renormalization group invariant quantity and a noninvariant one. Therefore, a further signal of the instanton-induced cancellation would be a scale invariance of the cancellation: if the vanishing of the axial charge were due instead to a perturbative mechanism, it would disappear at higher scales. Again, this is unfortunately hard to test experimentally, since the QCD evolution of the perturbative contribution to Q_5^{an} only appears at two loops and is thus very weak [2].

Finally, a remarkable aspect of the instanton mechanism is its universality: the same cancellation should occur if one measured the axial charge of other polarized baryons.

A somewhat different approach to the phenomenology of instanton-induced interactions consists of asking oneself whether there exist other signals of an important quark-instanton coupling. That is, we assume that the effective instanton-induced interaction, described by the diagrams of Fig. 3 should be added to the usual perturbative quark-gluon interactions for the sake of parton-model computations. This is posited as a phenomenological assumption, and we investigate its consequences.

It is immediately clear that the quark sea produced through the diagrams of Fig. 3 stands against that produced *via* perturbative evolution because of its peculiar structure in chirality *and flavor*. The chirality structure of the effective 't Hôoft interaction is at the origin of the spin effects discussed above, and appears as an anticorrelation between the chirality of the valence quark, and that of the outgoing sea quark-antiquark pairs. A similar anticorrelation exists in flavor space: the $2N_f$ lines emanating from the 't Hooft vertices depicted in Fig. 3 correspond to a particle-antiparticle pair per flavor. Thus, for instance, in the diagram of Fig. 3(a) (in the two-flavor case) if the initial and final "valence" quark state is, say, an up quark, then the intermediate three-quark states must be made of an up quark and a down quark-antiquark pair. If we define a flavor charge Q_f as the difference in numbers of up and down partons

$$Q_f \equiv n_u - n_d \,, \tag{15}$$

where i-flavor quarks and antiquarks contribute to n_i with the same sign, then the behavior of Q_f along the diagram in Fig. 3(a) is the same as that of Q_5 displayed in Fig. 2.

An important difference between the flavor charge (15) and the axial charge is that the diagrams of Fig. 3(b), while providing an anticorrelated contribution to the axial charge, do not contribute to the flavor charge — the flavor of the propagating quark remains unchanged. More in general, if we consider more complicated arrangements of instantons and anti-instantons, the chirality flip of each quark which takes part to the instanton induced interaction is a universal property, whereas the behavior of the flavor charge depends on the way the quark lines are contracted in the diagram. This is mirrored by the fact that the instanton-induced contribution to the axial charge may be written as a universal functional (3) of the gluon field, while no such expression may exist for the flavor charge, because the gluon fields are flavor singlets. Accordingly, a detailed computation of the instanton-induced effects on the flavor charge of a given valence quark is more involved. A detailed quantitative investigation is currently in progress [5].

In sum, although we cannot yet give a quantitative estimate of the instanton-induced effects on the flavor charge of a quark, qualitatively it is clear that the effect will correspond to a reduction of the charge from its valence value to a smaller one. In general, however, we do not expect complete cancellation, contrary to what happens for the axial charge.

It is interesting to note that the same problem has been studied in perturbative QCD [15]: there, one may calculate the flavor charge of the sea quark-antiquark pairs contained in a quark of given flavor. An analogous effect (*i.e.*, reduction of the valence value) is found. The effect, however, is of order $\sim 0.01\alpha_s^2$, which is much smaller than the magnitude we would expect for the nonperturbative effect, that should be comparable to that on the axial charge.

Now, the flavor charge (15) of the nucleon is an experimentally accessible quantity. More precisely, the difference in moments of the structure functions F_2 for the proton and neutron has been measured [4] and is proportional to the flavor charge of the proton:

$$\begin{aligned}\int_0^1 dx\, \frac{\left(F_2^p(x) - F_2^n(x)\right)}{x} &= \frac{4}{9}\left(n_u^p - n_u^n\right) + \frac{1}{9}\left(n_d^p - n_d^n\right) \\ &= \frac{1}{3} Q_f^p ,\end{aligned} \tag{16}$$

where in the last step we have used isospin symmetry. The valence value of Q_f^p is of course one, leading to the value of $\frac{1}{3}$ for the difference of moments of structure functions (Gottfried sum rule). Sea quarks produced by perturbative two-loop evolution should not modify this value significantly at realistic energies [15]. The experimental value reported by the NMC collaboration is 0.24 ± 0.016, implying that the flavor charge is of order $Q_f^p \sim \frac{2}{3}$, thus calling again for a nonperturbative explanation. The instanton effects described above produce an effect which is at least qualitatively in agreement with this.

4. Conclusion

Instanton effects have been suggested several times as an important ingredient in would-be nonperturbative QCD computations; however, most of the results obtained in instanton models so far are either qualitative, or very strongly model dependent. The precision data on structure functions which have become available recently are offering the possibility of changing this state of affairs. Perhaps, it will be possible to treat the instanton-induced interactions in a way which is analogous to the naive quark model: although the theoretical foundation of the model is only heuristic there exist a well-defined scheme to perform quantitative computations. This would open a new window on quantitative nonperturbative QCD.

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