ASPECTS OF THE THEORY OF DEEP INELASTIC SCATTERING

F.M. Lev

Laboratory of Nuclear Problems, Joint Institute for Nuclear Research, Dubna, Moscow region 141980, Russia

Abstract

The (electromagnetic or weak) current operator responsible for deep inelastic scattering (DIS) should be local and satisfy the well-known commutation relations with the representation operators of the Poincare group. The problem whether these conditions are compatible with the factorization theorem and operator product expansion is investigated in detail. We argue that the current operator contains a nontrivial nonperturbative part which contributes to DIS even in the Bjorken limit. Nevertheless there exists a possibility that many results of the standard theory remain.

1 The statement of the problem

The recent discovery of diffractive deep inelastic scattering (DIS) at HERA [1, 2] poses the problem whether DIS processes with small values of the Bjorken variable $x_{Bj}$ can be described in the framework of perturbative QCD (see e.g. ref. [3] and references therein). At the same time the applicability of perturbative QCD in the Bjorken limit (i.e. in the case when the momentum transfer $q$ is such that $Q = |q^2|^{1/2}$ is large and $x_{Bj}$ is not too close to 0 and 1) follows from the factorization theorem (FT) considered by several authors (see e.g. refs. [4, 5, 6, 7]).

The formulations of the FT given by different authors differ each other but the idea of all the formulations is that
• 1\(_{FT}\) The amplitudes of lepton-parton interactions can be calculated in the framework of perturbative QCD.

• 2\(_{FT}\) The parton distribution functions (PDFs) are essentially nonperturbative objects, but they are universal in the sense that they are the same in all hard-scattering processes.

The condition 1\(_{FT}\) follows from the property that in amplitudes of the lepton-parton interactions entering into diagrams dominating in DIS the off-shellness of quarks interacting with the virtual photon, W or Z bosons is very large. This implies that the (electromagnetic or weak) current operator \(\hat{J}^\mu(x)\) (where \(\mu = 0, 1, 2, 3\) and \(x\) is a point in Minkowski space) responsible for the transition nucleon \(\rightarrow\) hadrons in DIS can be considered in the framework of perturbation theory.

If \(|N\rangle\) is the state of the initial nucleon then the DIS cross-section is fully defined by the hadronic tensor

\[
W^{\mu\nu} = \frac{1}{4\pi} \int e^{ix\cdot q} \langle N | \hat{J}^\mu(x) \hat{J}^\nu(-x) |N\rangle d^4x \tag{1}
\]

The usual motivation of the FT is that since the main contribution to the integral (1) is given by small distances (more precisely by the values of \(x\) close to the light cone), the possibility to consider the operator \(\hat{J}^\mu(x)\) in the framework of perturbation theory is a consequence of asymptotic freedom (which implies that \(\alpha_s(Q^2) \rightarrow 0\) when \(Q \rightarrow \infty\), where \(\alpha_s\) is the QCD running coupling constant). However a rigorous proof of the FT encounters serious difficulties (see the discussion in ref. \[6\]). As noted in ref. \[5\], ”it is fair to say that a rigorous treatment of factorization has yet to be provided”. 

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In any relativistic quantum theory the system under consideration is described by some unitary representation \((a, l) \rightarrow \hat{U}(a, l)\) of the Poincare group. Here \(\hat{U}(a, l)\) is the representation operator corresponding to the element \((a, l)\) of the Poincare group, \(a\) is the four-vector describing the displacement of the origin in spacetime translation of Minkowski space and \(l \in SL(2, C)\). We use \(\hat{U}(a)\) to denote \(\hat{U}(a, 1)\) and \(\hat{U}(l)\) to denote \(\hat{U}(0, l)\). Let \(\hat{P} = (\hat{P}^0, \hat{\mathbf{P}})\) be the four-momentum operator, where \(\hat{P}^0\) is the Hamiltonian and \(\hat{\mathbf{P}}\) is the operator of ordinary momentum. Then \(\hat{U}(a) = \exp(i\hat{P}_\mu a^\mu)\) where a sum over repeated indices is assumed. Analogously it is well-known that \(\hat{U}(l)\) can be written in terms of the operators \(\hat{M}^{\mu\nu}\) (\(\hat{M}^{\mu\nu} = -\hat{M}^{\nu\mu}\)) which are the representation generators of the Lorentz group.

We shall always assume that the commutation relations for the representation generators of the Poincare group are realized in the form

\[
[\hat{P}^\mu, \hat{P}^\nu] = 0, \quad [\hat{M}^{\mu\nu}, \hat{P}^\rho] = -i(\eta^{\mu\rho} \hat{P}^\nu - \eta^{\nu\rho} \hat{P}^\mu),
\]

\[
[\hat{M}^{\mu\nu}, \hat{M}^{\rho\sigma}] = -i(\eta^{\mu\rho} \hat{M}^{\nu\sigma} + \eta^{\nu\sigma} \hat{M}^{\mu\rho} - \eta^{\mu\sigma} \hat{M}^{\nu\rho} - \eta^{\nu\rho} \hat{M}^{\mu\sigma})
\]

where \(\mu, \nu, \rho, \sigma = 0, 1, 2, 3\), the metric tensor in Minkowski space has the nonzero components \(\eta^{00} = -\eta^{11} = -\eta^{22} = -\eta^{33} = 1\), and we use the system of units with \(\hbar = c = 1\).

As explained in the well-known textbooks and monographs (see e.g. refs. [8, 9]), matrix elements of field operators over some states have correct transformation properties relative transformations from the Poincare group only if transformations of the operators are compatible with the transformations of the states. This implies in particular that the hadronic tensor \(W^{\mu\nu}\) in Eq. (1) will be the true tensor only if the current operator is
the relativistic vector operator satisfying the conditions

\[ \hat{U}(a)^{-1}\hat{J}^\mu(x)\hat{U}(a) = \hat{J}^\mu(x-a), \]  

(3)

\[ \hat{U}(l)^{-1}\hat{J}^\mu(x)\hat{U}(l) = L(l)^\mu_\nu\hat{J}^\nu(L(l)^{-1}x) \]  

(4)

where \( L(l) \) is the element of the Lorentz group corresponding to \( l \). The second expression implies that, as a consequence of Lorentz invariance,

\[ [\hat{M}^\mu\nu, \hat{J}^\rho(x)] = -i\{ (x^\mu \partial^\nu - x^\nu \partial^\mu)\hat{J}^\rho(x) + \eta^\mu\rho\hat{J}^\nu(x) - \eta^\nu\rho\hat{J}^\mu(x) \} \]

(5)

In addition, the current operator should be local in the sense that \([\hat{J}^\mu(x/2), \hat{J}^\nu(−x/2)] = 0 \) if \( x^2 < 0 \).

Let us now consider the difference between relativistic invariance and covariance although different authors understand these notions differently. The latter often means that each stage of calculations involves only objects belonging to finite-dimensional representations of the group \( SL(2,\mathbb{C}) \) — Dirac spinors, four-vectors, their scalar products in Minkowski metric etc. The former means that all operators corresponding to physical observables satisfy proper commutation relations with the Poincare group representation operators for the system under consideration. It is clear that in the general case relativistic invariance is the necessary physical condition while covariance is not.

In QED the electrons, positrons and photons are the fundamental particles, and the scattering space is the space of these almost free particles. The operators \( \hat{P}^\mu, \hat{M}^{\mu\nu} \) act in the scattering space of the system under consideration, and therefore, when the S-matrix is calculated in perturbation theory, they can be replaced by the corresponding free operators \( P^\mu, M^{\mu\nu} \). Then,
as shown in the standard textbooks (see e.g. ref. [8]), a covariant way of calculating the S-matrix guarantees that it will be relativistically invariant.

However in QCD the scattering space by no means can be considered as a space of almost free fundamental particles — quarks and gluons. For example, even if the scattering space consists of one particle (say the nucleon), this particle is the bound state of quarks and gluons, and the operators $\hat{P}^\mu, \hat{M}^{\mu\nu}$ considerably differ from $P^\mu, M^{\mu\nu}$. It is well-known (even in the nonrelativistic quantum mechanics) that in the presence of bound states the sets ($\hat{P}^\mu, \hat{M}^{\mu\nu}$) and ($P^\mu, M^{\mu\nu}$) cannot be unitarily equivalent, and, since perturbation theory does not apply to bound states, the former operators cannot be determined in the framework of perturbation theory. In this case covariance does not guarantee that the results will be relativistically invariant since the operators in question should properly commute with $\hat{P}^\mu, \hat{M}^{\mu\nu}$ and not with $P^\mu, M^{\mu\nu}$. The usual assumption in covariant calculations involving a bound state is that any matrix element describing a transition of this state can be (at least in principle) obtained by calculating matrix elements involving only transitions of fundamental particles into each other and taking the residue corresponding to the bound state. However in view of confinement it is not clear how to substantiate the scattering problem for free quarks and gluons. For these reasons we will be interested in cases when the representation operators in Eqs. (3) and (4) correspond to the full generators $\hat{P}^\mu, \hat{M}^{\mu\nu}$.

Now the question arises whether perturbative consideration of the current operator in the Bjorken limit is compatible with the correct transformation properties of the quantity $W^{\mu\nu}$. Indeed, if the current operator is considered in perturbation theory
then the momentum and angular momentum operators in Eqs. (3) and (4) are not $\hat{P}^\mu, \hat{M}^{\mu\nu}$ but $P^\mu, M^{\mu\nu}$. At the same time, the Poincare transformations of the state $|N\rangle$ are described by the operators $\hat{P}^\mu, \hat{M}^{\mu\nu}$.

Another possible approach to DIS is as follows. Instead of the FT we use the operator product expansion (OPE) [10] according to which

- $1_{OPE}$) The product of the operators in Eq. (1) can be written in the form

$$\hat{J}(\frac{x}{2})\hat{J}(\frac{-x}{2}) = \sum_i C_i(x^2)x_{\mu_1}\cdots x_{\mu_n}\hat{O}_i^{\mu_1\cdots\mu_n}$$

where $C_i(x^2)$ are the $c$-number Wilson coefficients and the operators $\hat{O}_i^{\mu_1\cdots\mu_n}$ are regular at $x = 0$.

- $2_{OPE}$) These operators depend only on field operators and their covariant derivatives at the origin of Minkowski space and have the same form as in perturbation theory.

For example, the basis for twist two operators contains in particular

$$\hat{O}^\mu_V = \mathcal{N}\{\hat{\psi}(0)\gamma^\mu\hat{\psi}(0)\} \quad \hat{O}_A^\mu = \mathcal{N}\{\hat{\psi}(0)\gamma^\mu\gamma^5\hat{\psi}(0)\}$$

where $\mathcal{N}$ stands for the normal product, $\hat{\psi}(x)$ is the Heisenberg operator of the Dirac field and for simplicity we do not write flavor operators and color and flavor indices.

However the OPE has been proved only in perturbation theory of renormalized interactions [11] while in view of Eqs. (3) and (4) we have to use the OPE beyond perturbation theory. The problem of the validity of the OPE beyond perturbation
theory is rather difficult and so far concrete results in this field have been obtained only in two-dimensional models (see ref. [12] and references therein).

Although there exists a vast literature devoted to the theory of DIS, the restrictions imposed on the current operator by Eqs. (3) and (4) have not been considered. In view of the above discussion it seems reasonable to investigate whether these restrictions can add something to the understanding of the validity of the properties $1_{FT}$, $2_{FT}$, $1_{OPE}$ and $2_{OPE}$).

The paper is organized as follows. In Sects. 2 and 3 we discuss the general properties of the current operator in quantum field theory and the properties derived in the framework of canonical formalism. In addition to the results of many authors it is shown in Sect. 4, that the latter properties are not reliable since in some cases they are incompatible with Lorentz invariance. The current operator in DIS is considered in Sect. 5 and it is argued that the nonperturbative part of this operator contributes to DIS even in the Bjorken limit. In Sect. 6 we describe a model which, in our opinion, is important for understanding the problems considered in the paper. Finally Sect. 7 is discussion.

2 Problems with constructing the current operator

Strictly speaking, the notion of current is not necessary if the theory is complete. For example, in QED there exist unambiguous prescriptions for calculating the elements of the S-matrix to any desired order of perturbation theory and this is all we
need. It is believed that this notion is useful for describing the electromagnetic or weak properties of strongly interacted systems. It is sufficient to know the matrix elements $\langle \beta | \hat{J}^\mu (x) | \alpha \rangle$ of the operator $\hat{J}^\mu (x)$ between the (generalized) eigenstates of the operator $\hat{P}^\mu$ such that $\hat{P}^\mu | \alpha \rangle = P^\mu_\alpha | \alpha \rangle$, $\hat{P}^\mu | \beta \rangle = P^\mu_\beta | \beta \rangle$. It is usually assumed that as a consequence of Eq. (3)

$$
\langle \beta | \hat{J}^\mu (x) | \alpha \rangle = \exp \left[ i (P^\nu_\beta - P^\nu_\alpha) x_\nu \right] \langle \beta | \hat{J}^\mu | \alpha \rangle
$$

where formally $\hat{J}^\mu \equiv \hat{J}^\mu (0)$. Therefore in the absence of a complete theory we can consider the less fundamental problem of investigating the properties of the operator $\hat{J}^\mu$. From the mathematical point of view this implies that we treat $\hat{J}^\mu (x)$ not as a four-dimensional operator distribution, but as a usual operator function satisfying the condition

$$
\hat{J}^\mu (x) = \exp (i \hat{P} x) \hat{J}^\mu \exp (-i \hat{P} x)
$$

The standpoint that the current operator should not be treated on the same footing as the fundamental local fields is advocated by several authors in their investigations on current algebra (see e.g. ref. [13]). One of the arguments is that, for example, the canonical current operator in QED is given by [8]

$$
\hat{J}^\mu (x) = \mathcal{N} \{ \hat{\psi} (x) \gamma^\mu \hat{\psi} (x) \} = \frac{1}{2} [\hat{\psi} (x), \gamma^\mu \hat{\psi} (x)]
$$

but this expression is not a well-definition of a local operator. Indeed, as explained in several well-known monographs (see e.g. refs. [14, 9]), the interacting field operators can be treated only as operator valued distributions and therefore the product of two local field operators at coinciding points is not well defined. The problem of the correct definition of such products is known as
that of constructing composite operators (see e.g. ref. [15]). So far this problem has been solved only in the framework of perturbation theory for special models. When perturbation theory does not apply the usual prescriptions are to separate the arguments of the operators in question and to define the composite operator as a limit of nonlocal operators when the separation goes to zero (see e.g. ref. [16] and references therein). Since we do not know how to work with quantum field theory beyond perturbation theory, we do not know what is the correct prescription.

An additional difficulty with the current operator given by Eq. (10) is as follows. It is well-known that this operator is unitarily equivalent to the free current operator (to the current operator in interaction picture). At the same time, as noted above, in the general case the operators $(\hat{P}^\mu, \hat{M}^{\mu\nu})$ are not unitarily equivalent to $(P^\mu, M^{\mu\nu})$. Therefore the problem arises whether the operator given by Eq. (10) will satisfy proper commutation relations with the operators $(\hat{P}^\mu, \hat{M}^{\mu\nu})$ in the presence of bound states.

It is well-known (see e.g. ref. [16]) that it is possible to add to the current operator the term $\partial_\nu X^{\mu\nu}(x)$ where $X^{\mu\nu}(x)$ is some operator antisymmetric in $\mu$ and $\nu$. However it is usually believed [16] that the electromagnetic and weak current operators of strongly interacted systems are given by the canonical quark currents the form of which is similar to that in Eq. (10).

If the operator $\hat{J}^\mu$ can be correctly defined then, as follows from Eqs. (9) and (9),

$$[\hat{M}^{\mu\nu}, \hat{J}^\rho] = -i(\eta^{\mu\rho} \hat{J}^\nu - \eta^{\nu\rho} \hat{J}^\mu)$$

(11)
3 Canonical quantization and the forms of relativistic dynamics

In the standard formulation of quantum field theory the operators $\hat{P}_\mu, \hat{M}_{\mu\nu}$ are given by

$$\hat{P}_\mu = \int \hat{T}^\nu_\mu(x)d\sigma_\nu(x), \quad \hat{M}_{\mu\nu} = \int \hat{M}^\rho_{\mu\nu}(x)d\sigma_\rho(x)$$ \hspace{1cm} (12)

where $\hat{T}^\nu_\mu(x)$ and $\hat{M}^\rho_{\mu\nu}(x)$ are the energy-momentum and angular momentum tensors and $d\sigma_\mu(x) = \lambda_\mu\delta(\lambda x - \tau)d^4x$ is the volume element of the space-like hypersurface defined by the time-like vector $\lambda$ ($\lambda^2 = 1$) and the evolution parameter $\tau$. In turn, these tensors are fully defined by the classical Lagrangian and the canonical commutation relations on the hypersurface $\sigma_\mu(x)$. In this connection we note that in canonical formalism the quantum fields are supposed to be distributions only relative the three-dimensional variable characterizing the points of $\sigma_\mu(x)$ while the dependence on the variable describing the distance from $\sigma_\mu(x)$ is usual [9].

In spinor QED we define $V(x) = -L_{\text{int}}(x) = e\hat{J}^\mu(x)\hat{A}_\mu(x)$, where $L_{\text{int}}(x)$ is the quantum interaction Lagrangian, $e$ is the (bare) electron charge and $\hat{A}_\mu(x)$ is the operator of the Maxwell field (let us note that if $\hat{J}^\mu(x)$ is treated as a composite operator then the product of the operators entering into $V(x)$ should be correctly defined).

At this stage it is not necessary to require that $\hat{J}^\mu(x)$ is given by Eq. (10), but the key assumption in the canonical formulation of QED is that $\hat{J}^\mu(x)$ is constructed only from $\hat{\psi}(x)$ (i.e. there is no dependence on $\hat{A}_\mu(x)$ and the derivatives of the fields $\hat{A}_\mu(x)$ and $\hat{\psi}(x)$). Then the canonical result derived in several
well-known textbooks and monographs (see e.g. ref. [8]) is

\[ \hat{P}^\mu = P^\mu + \lambda^\mu \int V(x)\delta(\lambda x - \tau)d^4x \]  

(13)

\[ \hat{M}^{\mu\nu} = M^{\mu\nu} + \int V(x)(x^\nu \lambda^\mu - x^\mu \lambda^\nu)\delta(\lambda x - \tau)d^4x \]  

(14)

It is important to note that if \( A^\mu(x) \), \( J^\mu(x) \) and \( \psi(x) \) are the corresponding free operators then \( \hat{A}^\mu(x) = A^\mu(x) \), \( \hat{J}^\mu(x) = J^\mu(x) \) and \( \hat{\psi}(x) = \psi(x) \) if \( x \in \sigma_\mu(x) \).

As pointed out by Dirac [17], any physical system can be described in different forms of relativistic dynamics. By definition, the description in the point form implies that the operators \( \hat{U}(l) \) are the same as for noninteracting particles, i.e. \( \hat{U}(l) = U(l) \) and \( \hat{M}^{\mu\nu} = M^{\mu\nu} \), and thus interaction terms can be present only in the four-momentum operators \( \hat{P} \) (i.e. in the general case \( \hat{P}^\mu \neq P^\mu \) for all \( \mu \)). The description in the instant form implies that the operators of ordinary momentum and angular momentum do not depend on interactions, i.e. \( \hat{P} = P \), \( \hat{M} = M \) (\( \hat{M} = (\hat{M}^{23}, \hat{M}^{31}, \hat{M}^{12}) \)), and therefore interaction terms may be present only in \( \hat{P}^0 \) and the generators of the Lorentz boosts \( \hat{N} = (\hat{M}^{01}, \hat{M}^{02}, \hat{M}^{03}) \). In the front form with the marked z axis we introduce the + and - components of the four-vectors as \( x^+ = (x^0 + x^2)/\sqrt{2}, x^- = (x^0 - x^2)/\sqrt{2} \). Then we require that the operators \( \hat{P}^+, \hat{P}^j, \hat{M}^{12}, \hat{M}^{+-}, \hat{M}^{+j} \) (\( j = 1, 2 \)) are the same as the corresponding free operators, and therefore interaction terms may be present only in the operators \( \hat{M}^{-j} \) and \( \hat{P}^- \).

In quantum field theory the form of dynamics depends on the choice of the hypersurface \( \sigma_\mu(x) \). The representation generators of the subgroup which leaves this hypersurface invariant are free since transformations from this subgroup do not involve dynamics. Therefore it is reasonable to expect that Eqs. (13)
and (14) give the most general form of the Poincare group rep-
representation generators in quantum field theory if the fields are
quantized on the hypersurface \( \sigma_\mu(x) \), but in the general case the
relation between \( V(x) \) and \( L_{\text{int}}(x) \) is not so simple as in QED.
The fact that the operators \( V(x) \) in Eqs. (13) and (14) are the
same follows from Eq. (2).

The most often considered case is \( \tau = 0, \lambda = (1, 0, 0, 0) \). Then
\( \delta(\lambda x - \tau) d^4x = d^3x \) and the integration in Eqs. (13) and (14)
is taken over the hyperplane \( x^0 = 0 \). Therefore, as follows from
these expressions, \( \hat{P} = P \) and \( \hat{M} = M \). Hence such a choice of
\( \sigma_\mu(x) \) leads to the instant form [17].

The front form can be formally obtained from Eqs. (13) and
(14) as follows. Consider the vector \( \lambda \) with the components
\( \lambda^0 = \frac{1}{(1 - v^2)^{1/2}}, \quad \lambda^j = 0, \quad \lambda^3 = -\frac{v}{(1 - v^2)^{1/2}} \quad (j = 1, 2) \)
(15)

Then taking the limit \( v \to 1 \) in Eqs. (13) and (14) we get
\[
\hat{P}^\mu = P^\mu + \omega^\mu \int V(x) \delta(x^+) d^4x,
\]
\[
\hat{M}^{\mu\nu} = M^{\mu\nu} + \int V(x)(x^\nu \omega^\mu - x^\mu \omega^\nu) \delta(x^+) d^4x
\]
(16)
where the vector \( \omega \) has the components \( \omega^- = 1, \omega^+ = \omega^j = 0 \).
It is obvious that the generators (16) are given in the front form
and that’s why Dirac [17] related this form to the choice of the
light cone \( x^+ = 0 \).

In ref. [17] the point form was related to the hypersurface
\( t^2 - x^2 > 0, \ t > 0 \), but as argued by Sokolov [18], the point
form should be related to the hyperplane orthogonal to the four-
velocity of the system under consideration. We shall not discuss
this question in the present paper.
In the case of systems with a finite number of particles all the forms are unitarily equivalent \[19\] and therefore the choice of the form is only the matter of convenience but not the matter of principle. However in quantum field theory it is not clear whether there exist forms in which the Poincare group representation operators are correctly defined \[14, 9\].

It is clear that when a form of dynamic is chosen, the expressions for the representation operators and the wave functions become noncovariant. Nevertheless, in the Feynman diagram technique each term of the perturbative series for the S-matrix is covariant. In QCD it is not clear whether covariance can be preserved in the presence of bound states (see the discussion in Sect. [1]) but anyway, as noted above, covariance is not a necessary condition.

4 Incompatibility of canonical formalism with Lorentz invariance for spinor fields

It has been shown in a vast literature (see e.g. refs. \[20, 21, 22, 23, 24, 16, 25\]) that the canonical treatment of commutation relations between the current operators can often lead to incorrect results. The results of these references show that it is often premature to trust assumptions which may seem physical or natural. The results of the present section can be considered as an additional argument in favor of this point of view. Namely, the purpose of this section is to show that the relation \(\hat{J}^\mu(x) = J^\mu(x)\) if \(x \in \sigma_{\mu}(x)\), which is the key property of the current operator for the spinor field in canonical formalism, is not correct since it is incompatible with the correct commuta-
tion relations between the operators \( \hat{J}_\mu(x) \) and \( \hat{M}^{\mu\nu} \) (i.e. with Lorentz invariance of the current operator).

A possible objection against the derivation of Eqs. (13) and (14) is that the product of local operators at one and the same value of \( x \) is not a well-defined object. For example, if \( x^0 = 0 \) then following Schwinger [20], instead of Eq. (10), one can define \( J_\mu(x) \) as the limit of the operator

\[
J_\mu(x) = \frac{1}{2} \left[ \bar{\psi}(x + \frac{1}{2}), \gamma^\mu \exp (ie \int_{x-\frac{1}{2}}^{x+\frac{1}{2}} A(x') dx') \psi(x - \frac{1}{2}) \right]
\]  

(17)

when \( l \to 0 \), the limit should be taken only at the final stage of calculations and in the general case the time components of the arguments of \( \hat{\psi} \) and \( \hat{\psi} \) also differ each other (the contour integral in this expression is needed to conserve gauge invariance). Therefore there is a ”hidden” dependence of \( \hat{J}_\mu(x) \) on \( \hat{A}_\mu(x) \) and hence Eqs. (13) and (14) are incorrect.

However, any attempt to separate the arguments of the \( \hat{\psi} \) operators in \( \hat{J}_\mu(x) \) immediately results in breaking of locality. In particular, at any \( l \neq 0 \) in Eq. (17) the Lagrangian is non-local and the whole edifice of local quantum field theory (including canonical formalism) becomes useless. Meanwhile the only known way of constructing the generators \( \hat{P}_\mu, \hat{M}^{\mu\nu} \) in local quantum field theory is canonical formalism. For these reason we first consider the results which formally follow from canonical formalism and then show that they are inconsistent.

In addition to the properties discussed above, the current operator should also satisfy the continuity equation \( \partial \hat{J}_\mu(x) / \partial x^\mu = 0 \). As follows from this equation and Eq. (3), \([\hat{J}_\mu(x), \hat{P}_\mu] = 0\). The canonical formalism in the instant form implies that if \( x^0 = 0 \) then \( \hat{J}_\mu(x) = J_\mu(x) \). Since \( J_\mu(x) \) satisfies the condition
\[ [J^\mu(x), P^\nu] = 0, \] it follows from Eq. (13) that if \( \hat{P}^\mu = P^\mu + V^\mu \)
then the continuity equation is satisfied only if
\[ [V^0, J^0(x)] = 0 \] (18)
where
\[ V^0 = \int V(x) d^3x, \quad V(x) = -eA(x)J(x) \] (19)
We take into account the fact that the canonical quantization on the hypersurface \( x^0 = 0 \) implies that \( A^0(x) = 0 \).

As follows from Eqs. (3) and (5), the commutation relation between the operators \( \hat{M}^{0i} (i = 1, 2, 3) \) and \( J^0(x) \) should have the form
\[ [\hat{M}^{0i}, J^0(x)] = -x^i[\hat{P}^0, J^0(x)] - \nu J^0(x) \] (20)
Since
\[ [M^{0i}, J^0(x)] = -x^i[P^0, J^0(x)] - \nu J^0(x) \] (21)
it follows from Eqs. (14), (18) and (19) that Eq. (20) is satisfied if
\[ \int y^iA(y)[J(y), J^0(x)]d^3y = 0 \] (22)
It is well-known that if the standard equal-time commutation relations are used naively then the commutator in Eq. (22) vanishes and therefore this equation is satisfied. However when \( x \to y \) this commutator involves the product of four Dirac fields at \( x = y \). The famous Schwinger result [20] is that if the current operators in question are given by Eq. (17) then
\[ [J^i(y), J^0(x)] = C \frac{\partial}{\partial x^i} \delta(x - y) \] (23)
where \( C \) is some (infinite) constant. Therefore Eq. (22) is not satisfied and the current operator \( \hat{J}^\mu(x) \) constructed in the
framework of canonical formalism does not satisfy Lorentz invariance.

At the same time, Eq. (23) is compatible with Eqs. (18) and (19) since div(A(x)) = 0. One can also expect that the commutator [M^{0i}, J^k(x)] is compatible with Eq. (1). This follows from the fact [26] that if Eq. (23) is satisfied then the commutator [J^i(x), J^k(y)] does not contain derivatives of the delta function.

While the arguments given in ref. [20] prove that the commutator in Eq. (23) cannot vanish, one might doubt whether the singularity of the commutator is indeed given by the right hand side of this expression. Of course, at present any method of calculating such a commutator is model dependent, but the incompatibility of canonical formalism with Lorentz invariance (see Eq. (20)) follows in fact only from algebraic considerations. Indeed, Eqs. (18), (20) and (21) imply that if M^{\mu\nu} = M^{\mu\nu} + V^{\mu\nu} then

\[ [V^{0i}, J^0(x)] = 0 \]  

(24)

Since V^{0i} in the instant form is a nontrivial interaction dependent operator, there is no reason to expect that it commutes with the free operator J^0(x). Moreover for the analogous reason Eq. (18) will not be satisfied in the general case.

To better understand the situation in spinor QED it is useful to consider scalar QED [27]. The formulation of this theory can be found, for example, in ref. [28]. In contrast with spinor QED, the Schwinger term in scalar QED emerges canonically [20, 16]. We use \( \varphi(x) \) to denote the operator of the scalar complex field at \( x^0 = 0 \). The canonical calculation yields

\[ \hat{J}^0(x) = J^0(x) = i[\varphi^*(x)\pi^*(x) - \pi(x)\varphi(x)], \]
\[ \hat{J}^i(x) = J^i(x) - 2eA^i(x)\varphi^*(x)\varphi(x), \]
\[ J^i(x) = i[\varphi^*(x) \cdot \partial^i \varphi(x) - \partial^i \varphi^*(x) \cdot \varphi(x)] \] (25)

where \( \pi(x) \) and \( \pi^*(x) \) are the operators canonically conjugated with \( \varphi(x) \) and \( \varphi^*(x) \) respectively. In contrast with Eq. (19), the operator \( V(x) \) in scalar QED is given by

\[ V(x) = -eA(x)J(x) + e^2A(x)^2\varphi^*(x)\varphi(x) \] (26)

However the last term in this expression does not contribute to the commutator (20). It is easy to demonstrate that as pointed out in ref. [27], the commutation relations (3) in scalar QED are satisfied in the framework of the canonical formalism. Therefore the naive treatment of the product of local operators at coinciding points in this theory is not in conflict with the canonical commutation relations. The key difference between spinor QED and scalar QED is that, in contrast with spinor QED, the spatial component of the canonical current operator is not free if \( x_0 = 0 \) (see Eq. (25)). Just for this reason the commutator \([\hat{M}_0^i, J^0(x)]\) in scalar QED agrees with Eq. (3) since the Schwinger term in this commutator gives the interaction term in \( \hat{J}^i(x) \).

Now let us return to spinor QED. As noted above, the canonical formalism cannot be used if the current operator is considered as a limit of the expression similar to that in Eq. (17). In addition, the problem exists what is the correct definition of \( V(x) \) as a composite operator. One might expect that the correct definition of \( J^\mu(x) \) and \( V(x) \) will result in appearance of some additional terms in \( V(x) \) (and hence in \( V^0 \) and \( V^{0i} \)). However it is unlikely that this is the main reason of the violation of Lorentz invariance. Indeed, as noted above, for only algebraic reasons it is unlikely that both conditions (18) and (24) can be simultaneously satisfied. Therefore, taking into account the situation in scalar QED, it is natural to think that the main reason...
of the failure of canonical formalism is that either the limit of \( \hat{\mathcal{J}}^{\mu}(x^0, \mathbf{x}) \) when \( x^0 \to 0 \) does not exist or this limit is not equal to \( \mathcal{J}^{\mu}(\mathbf{x}) \) (i.e. the relation \( \hat{\mathcal{J}}^{\mu}(\mathbf{x}) = \mathcal{J}^{\mu}(\mathbf{x}) \) is incorrect).

The fact that the relation \( \hat{\mathcal{J}}^{\mu}(\mathbf{x}) = \mathcal{J}^{\mu}(\mathbf{x}) \) cannot be correct follows from simpler considerations. Indeed, assume first that this relation is valid. Then we can use canonical formalism in the framework of which the generator of the gauge transformations is \( \text{div} \mathbf{E}(\mathbf{y}) - J^0(\mathbf{y}) \), and if \( \mathbf{J}(\mathbf{x}) \) is gauge invariant then \( [\text{div} \mathbf{E}(\mathbf{y}) - J^0(\mathbf{y}), \mathbf{J}(\mathbf{x})] = 0 \). The commutator \( [J^0(\mathbf{y}), \mathbf{J}(\mathbf{x})] \) cannot be equal to zero \([20]\) and therefore \( \mathbf{J}(\mathbf{x}) \) does not commute with \( \text{div} \mathbf{E}(\mathbf{y}) \) while the free operator \( \mathbf{J}(\mathbf{x}) \) commutes with \( \text{div} \mathbf{E}(\mathbf{y}) \). The relation \( \hat{\mathcal{J}}^{\mu}(\mathbf{x}) = \mathcal{J}^{\mu}(\mathbf{x}) \) also does not take place in explicitly solvable two-dimensional models \([9]\). In addition, once we assume that the field operators on the hypersurface \( \sigma_{\mu}(\mathbf{x}) \) are free we immediately are in conflict with the Haag theorem \([29, 14, 9]\). However for our analysis of the current operator in DIS in Sect. \([9]\) it is important that \( \hat{\mathcal{J}}^{\mu}(\mathbf{x}) \neq \mathcal{J}^{\mu}(\mathbf{x}) \) as a consequence of Lorentz invariance.

By analogy with ref. \([20]\) it is easy to show that if \( x^+ = 0 \) then the canonical current operator in the front form \( \mathcal{J}^+(x^-, \mathbf{x}_\perp) \) (we use the subscript \( \perp \) to denote the projection of the three-dimensional vector onto the plane 12) cannot commute with all the operators \( \mathcal{J}^i(x^-, \mathbf{x}_\perp) \) \( (i = -, 1, 2) \). As easily follows from the continuity equation and Lorentz invariance \([2]\), the canonical operator \( \mathcal{J}^+(x^-, \mathbf{x}_\perp) \) should satisfy the relations

\[
[V^-, \mathcal{J}^+(x^-, \mathbf{x}_\perp)] = [V^-, \mathcal{J}^+(x^-, \mathbf{x}_\perp)] = 0 \quad (j = 1, 2) \quad (27)
\]

By analogy with the above consideration it is natural to think that these relations cannot be simultaneously satisfied and therefore either the limit of \( \hat{\mathcal{J}}^{\mu}(x^+, x^-, \mathbf{x}_\perp) \) when \( x^+ \to 0 \) does
not exist or this limit is not equal to \( J^\mu(x^-, \mathbf{x}_\perp) \). Therefore the canonical light cone quantization does not render a Lorentz invariant current operator for spinor fields.

Let us also note that if the theory should be invariant under the space reflection or time reversal, the corresponding representation operators in the front form \( \hat{U}_P \) and \( \hat{U}_T \) are necessarily interaction dependent (this is clear, for example, from the relations \( \hat{U}_P \hat{P}^+ \hat{U}_P^{-1} = \hat{U}_T \hat{P}^+ \hat{U}_T^{-1} = \hat{P}^- \)). In terms of the operator \( \hat{J}^\mu \) one can say that this operator should satisfy the conditions

\[
\hat{U}_P(\hat{J}^0, \hat{J}) \hat{U}_P^{-1} = \hat{U}_T(\hat{J}^0, \hat{J}) \hat{U}_T^{-1} = (\hat{J}^0, -\hat{J}) \tag{28}
\]

Therefore it is not clear whether these conditions are compatible with the relation \( \hat{J}^\mu = J^\mu \). However this difficulty is a consequence of the difficulty with Eq. (11) since, as noted by Coester [30], the interaction dependence of the operators \( \hat{U}_P \) and \( \hat{U}_T \) in the front form does not mean that there are discrete dynamical symmetries in addition to the rotations about transverse axes. Indeed, the discrete transformation \( P_2 \) such that \( P_2 \mathbf{x} := \{x^0, x_1, -x_2, x_3\} \) leaves the light front \( x^+ = 0 \) invariant. The full space reflection \( P \) is the product of \( P_2 \) and a rotation about the 2-axis by \( \pi \). Thus it is not an independent dynamical transformation in addition to the rotations about transverse axes. Similarly the transformation \( TP \) leaves \( x^+ = 0 \) invariant and \( T = (TP)P_2R_2(\pi) \).

In terms of the operator \( \hat{J}^\mu \) the results of this section are obvious. Indeed, since at \( x = 0 \) the Heisenberg and Schroedinger pictures coincide then in view of Eq. (10) one might think that the operator \( \hat{J}^\mu \) is free, i.e. \( \hat{J}^\mu = J^\mu \). However there is no reason for the interaction terms in \( \hat{M}^{\mu\nu} \) to commute with all the operators \( J^\mu \) (see Eq. (11)). Therefore the results of this section
show that the algebraic reasons based on Eq. (11) are more solid than the reasons based on formal manipulations with local operators and in the instant and front forms $\hat{J}^\mu \neq J^\mu$ (moreover, $\hat{J}^\mu(x) \neq J^\mu(x)$ if $x \in \sigma_\mu(x)$). In particular, if some interaction operators are present in $\hat{M}^{\mu\nu}$, one has to expect that they are also present in some of the operators $\hat{J}^\mu$.

Let us note that although the model considered in this section is spinor QED, the above results are not very important for QED itself and other standard theories where the S-matrix can be calculated in perturbation theory. The matter is that, as pointed out in Sect. [1], in this case it is sufficient to consider only commutators involving $P^\mu$ and $M^{\mu\nu}$. For example, the problem important for calculating the S-matrix elements in such the ories is that of constructing the covariant T-product $T^*(\hat{J}^\mu(x)\hat{J}^\nu(y))$ (see e.g. refs. [21, 22, 8, 28]). It has been shown that in the presence of Schwinger terms the standard T-product $T(\hat{J}^\mu(x)\hat{J}^\nu(y))$ is not covariant but it is possible to add to $T(\hat{J}^\mu(x)\hat{J}^\nu(y))$ a contact term (which is not equal to zero only if $x^0 = y^0$) such that the resulting $T^*$-product will be covariant. In standard theories covariance is equivalent to Lorentz invariance since covariance is satisfied when the operators in question properly commute with the $M^{\mu\nu}$. However in QCD Lorentz invariance is satisfied only if the operators in question properly commute with the full operators $\hat{M}^{\mu\nu}$. Therefore in this case Lorentz invariance and covariance differ each other and the former is of course more important.

In the next section we argue that the algebraic consideration based on Eq. (11) is important for investigating the properties of the current operator for strongly interacting particles.
5 Current operator in DIS

The DIS cross-section is fully defined by the hadronic tensor given by Eq. (1). The nucleon state $|N\rangle$ is the eigenstate of the operator $\hat{P}$ with the eigenvalue $P'$ and the eigenstate of the spin operators $\hat{S}^2$ and $\hat{S}^z$ which are constructed from $\hat{M}^{\mu\nu}$. In particular, $\hat{P}^2|N\rangle = m^2|N\rangle$ where $m$ is the nucleon mass.

The structure of the four-momentum operator $\hat{P}$ in QCD is rather complicated (see e.g. ref. [31]) but anyway some of the components of $\hat{P}$ necessarily contain a part which describes the interaction of quarks and gluons at large distances where $\alpha_s$ is large and perturbation theory does not apply. In view of the last relation, this part is responsible for binding of quarks and gluons in the nucleon. We will call this part the nonperturbative one.

Suppose that the Hamiltonian $\hat{P}^0$ contains the nonperturbative part and consider the relation $[\hat{M}^{0i}, \hat{P}^k] = -i\delta_{ik}\hat{P}^0 (i, k = 1, 2, 3)$ which follows from Eq. (2). Then it is obvious that if $\hat{P}^k = P^k$ then all the operators $\hat{M}^{0i}$ inevitably depend on the nonperturbative part and vice versa, if $\hat{M}^{0i} = M^{0i}$ then all the operators $\hat{P}^k$ inevitably depend on this part. Therefore in the instant form all the operators $\hat{M}^{0i}$ inevitably depend on the nonperturbative part and in the point form all the operators $\hat{P}^k$ inevitably depend on this part. In the front form the fact that all the operators $\hat{M}^{-j}$ inevitably depend on the nonperturbative part follows from the relation $[\hat{M}^{-j}, \hat{P}^l] = -i\delta_{jl}\hat{P}^l (j, l = 1, 2)$ which also is a consequence of Eq. (2). Of course, the physical results should not depend on the choice of the form of dynamics and in the general case all ten generators can depend on the nonperturbative part.
In turn, Eq. (5) and the consideration in Sect. 4 give grounds to think that some of the operators \( \hat{J}^\mu(x) \) in the instant form and \( \hat{J}^\mu(x^-, x_\perp) \) in the front one inevitably depend on the non-perturbative part. If it is possible to define \( \hat{J}^\mu \) in the point form then as follows from Eq. (12), the relation \( \hat{J}^\mu = J^\mu \) does not contradict Lorentz invariance but, as follows from Eq. (9), the operator \( \hat{J}^\mu(x) \) in that form inevitably depend on the non-perturbative part. As noted in Sect. 1, the fact that the same operators \( (\hat{P}^\mu, \hat{M}^{\mu\nu}) \) describe the transformations of both the operator \( \hat{J}^\mu(x) \) and the state \( |N\rangle \) guaranties that \( W^{\mu\nu} \) has the correct transformation properties.

We see that the relation between the current operator and the state of the initial nucleon is highly nontrivial. Meanwhile, in the present theory they are considered separately. As noted in Sect. 1, the possibility of the separate consideration follows from the FT.

Let us now discuss the following question. If the current operator depends on the nonperturbative part then this operator depends on the integrals from the quark and gluon field operators over the region of large distances where \( \alpha_s \) is large. Is this property compatible with locality? In the framework of canonical formalism compatibility is obvious but, as shown in the preceding section, the results based on canonical formalism are not reliable. Therefore it is not clear whether in QCD it is possible to construct local electromagnetic and weak current operators beyond perturbation theory. We will now consider whether such a possibility can be substantiated in the framework of the OPE.

As noted above, there are grounds to think that the operator \( \hat{J}^\mu(x) \) necessarily depends on the nonperturbative part while Eq. (5) has been proved only in perturbation theory. Therefore
if we use Eq. (6) in DIS we have to assume that either nonperturbative effects are not important to some orders in \(1/Q\) and then we can use Eq. (6) only to these orders (see e.g. ref. [32]) or it is possible to use Eq. (6) beyond perturbation theory. The question also arises whether Eq. (6) is valid in all admissible forms of dynamics (as it should be if it is an exact operator equality) or only in some forms.

In the point form all the components of \(\hat{P}\) depend on the nonperturbative part and therefore, in view of Eqs. (3) or (9), it is not clear why there is no nonperturbative part in the dependence of the right hand side of Eq. (6), or if (for some reasons) it is possible to include the nonperturbative part only into the operators \(\hat{O}_i\) then why they have the same form as in perturbation theory.

One might think that in the front form the \(C_i(x^2)\) will be the same as in perturbation theory, at least in the case when the process is considered in the infinite momentum frame (IMF) where the initial nucleon has a large positive momentum along the \(z\) axis. Then the value of \(q^-\) in DIS is very large and therefore only a small vicinity of the light cone \(x^+ = 0\) contributes to the integral (11). Therefore we indeed have the description in the front form where the only dynamical component of \(\hat{P}\) is \(\hat{P}^-\). Since the eigenvalues of \(\hat{P}^-\) in the IMF are small, the dependence of \(\hat{P}^-\) on the nonperturbative part is of no importance. Since there exist final hadrons (which are the bound states of quarks and gluons) moving in the negative direction of the \(z\) axis in the IMF, the problem arises whether the final state interaction (FSI) can be neglected. In addition, it is necessary to take into account that the operator \(\hat{J}^\mu(x^-, \mathbf{x}_\perp)\) in the front form depends on the nonperturbative part. Nevertheless we assume that Eq.
in the front form is valid (see also the consideration in the next section).

If we assume as usual that there is no problem with the convergence of the OPE series then experiment makes it possible to measure each matrix element $\langle N|\hat{O}^{\mu_1\ldots\mu_n}|N\rangle$ which should have correct transformation properties (if only the series as a whole has the correct properties then the decomposition (6) is of no practical importance). Let us consider, for example, the matrix element $\langle N|\hat{O}^\mu_V|N\rangle$. By analogy with the consideration in Sect. 1 we conclude that this matrix element transforms as a four-vector if the Lorentz transformations of $\hat{O}^\mu_V$ are described by the operators $\hat{M}^{\mu\nu}$ describing the transformations of $|N\rangle$, or in other words, by analogy with Eq. (11),

$$[\hat{M}^{\mu\nu}, \hat{O}^\rho_V] = -i(\eta^{\mu\rho}\hat{O}^\nu_V - \eta^{\nu\rho}\hat{O}^\mu_V) \quad (29)$$

It is also clear that Eq. (29) follows from Eqs. (2-7). Since the $\hat{M}^{-j}$ in the front form depend on the nonperturbative part, the results of Sect. 4 give grounds to think that at least some components $\hat{O}_V^\mu$, and analogously some components $\hat{O}_i^{\mu_1\ldots\mu_n}$, also depend on the nonperturbative part. Since Eq. (29) does not contain any $x$ or $q$ dependence, this conclusion has nothing to do with asymptotic freedom and is valid even in leading order in $1/Q$. Therefore the problem arises whether this fact is compatible with the FT [4, 5, 6, 7].

If the operators $\hat{O}_i^{\mu_1\ldots\mu_n}$ depend on the nonperturbative part, then by analogy with the above considerations we conclude that the operators in Eq. (7) are ill-defined and the correct expressions for them involve integrals from the field operators over large distances where the QCD coupling constant is large. Therefore it is not clear whether the operators $\hat{O}_i^{\mu_1\ldots\mu_n}$ are local.
and whether the Taylor expansion at $x = 0$ is correct, but even if it is, the expressions for $\hat{O}_{\mu_1...\mu_n}^i$ will depend on higher twist operators which contribute even in leading order in $1/Q$.

If (for some reasons) Eq. (6) is valid and no form of the operators $\hat{O}_{\mu_1...\mu_n}^i$ is prescribed then all the standard results concerning the $Q^2$ evolution of the structure functions remain. Indeed the only information about the operators $\hat{O}_{\mu_1...\mu_n}^i$ we need is their tensor structure since we should correctly parametrize the matrix elements $\langle N | \hat{O}_{\mu_1...\mu_n}^i | N \rangle$. However the form of the operators $\hat{O}_{\mu_1...\mu_n}^i$ is important for the derivation of sum rules in DIS. We will discuss this question below.

6 A model

The above discussion shows that the compatibility of Poincare invariance of the current operator with the FT and OPE is the difficult problem of the present theory. For this reason it is interesting to consider models in which the well-defined current operator and the representation generators of the Poincare group can be explicitly constructed and one can explicitly verify that they satisfy the commutation relations (2-5). In this section we consider a model the detailed description of which will be given elsewhere [33].

Consider the wave function of a system of $n$ particles with the four-momenta $p_i$, and masses $m_i$ ($i = 1, 2, ... n$). Since $p_i^2 = m_i^2$, only three components of $p_i$ are independent and we choose $p_i^\perp, p_i^+$ as these components. Instead of the individual variables $(p_{1\perp}, p_1^+, ... p_{n\perp}, p_n^+)$ we introduce the $\perp$ and $+$ components of...
the total momentum
\[ \mathbf{P}_\perp = \mathbf{p}_1\perp + ... + \mathbf{p}_n\perp, \quad P^+ = p^+_1 + ... + p^+_n \quad (30) \]

and the internal (Sudakov) variables
\[ \xi_i = \frac{p_i^+}{P^+}, \quad \mathbf{K}_\perp = \mathbf{p}_i\perp - \xi_i \mathbf{P}_\perp \quad (31) \]

Let us define the ”internal” space \( \mathcal{H}_{int} \) as the space of functions \( \chi(\mathbf{k}_1\perp, \xi_1, ... \mathbf{k}_n\perp, \xi_n) \) such that
\[ \|\|\chi\|\|^2 = \int |\chi(\mathbf{k}_1\perp, \xi_1, ... \mathbf{k}_n\perp, \xi_n)|^2 d\rho(int) < \infty \quad (32) \]
where \( d\rho(int) \) is the volume element in the internal momentum space and the spin variables are dropped for simplicity (the explicit expression for \( d\rho(int) \) is of no importance for us). Then the full Hilbert space \( \mathcal{H} \) can be realized as the space of functions \( \phi(\mathbf{P}_\perp, P^+) \) with the range in \( \mathcal{H}_{int} \) and such that
\[ \int \|\|\phi(\mathbf{P}_\perp, P^+)\|\|^2 \frac{d^2P_\perp dP^+}{2(2\pi)^3 P^+} < \infty \quad (33) \]

Suppose that the particles interact with each other, and their interactions are described in the front form. Then the \( \perp \) and \( + \) components of the four-momentum operator are the operators of multiplication by the corresponding variable, i.e. \( \hat{\mathbf{P}}_\perp = \mathbf{P}_\perp \), \( \hat{\mathbf{P}}^+ = P^+ \). As shown by several authors (see e.g. refs. [34]), there exists the representation where the remaining seven generators have the form
\[ \hat{\mathbf{P}}^- = \frac{\hat{M}^2 + \mathbf{P}_\perp^2}{2P^+}, \quad M^{+\perp} = iP^+ \frac{\partial}{\partial P^+}, \]
\[ M^{+j} = -iP^+ \frac{\partial}{\partial P^j}, \quad M^{xy} = -i\epsilon_{jkl} \frac{\partial}{\partial P^k} + S^z, \]
\[ \hat{M}^{-j} = -i(P^j \frac{\partial}{\partial P^+} + \hat{\mathbf{P}}^- \frac{\partial}{\partial P^j}) - \frac{\epsilon_{jkl}}{P^+}(\hat{M}\hat{S}^l + P^k S^z) \quad (34) \]
Here \( j = (x, y) \), \( \epsilon_{xx} = \epsilon_{yy} = 0 \), \( \epsilon_{xy} = -\epsilon_{yx} = 1 \), \( \hat{M} \) is the mass operator and \( \hat{S} \) is the spin operator. The latter act only through the variables of the space \( \mathcal{H}_{int} \).

Any interacting system should satisfy cluster separability which has been discussed by several authors (see e.g. refs. [33, 30, 37, 38, 39]). We will use the algebraic version of cluster separability: for any partition of the system under consideration into the subsystems \( (\alpha_1, \ldots, \alpha_k) \) such that all interactions between these subsystems are turned off, the representation of the Poincare group for the system as a whole becomes the tensor product of the representations for the subsystems. This implies that the representation generators for the system as a whole become sums of the corresponding generators for the subsystems.

As shown in ref. [40], a consequence of the cluster separability property is that the spin operator is not equal to the free one; only the \( z \) component of \( \hat{S} \) is free while \( \hat{S}_\perp \neq S_\perp \). However in models where \( n \) is finite

\[
\hat{S} = \hat{A} S \hat{A}^{-1}
\]  

where the unitary operator \( \hat{A} \) is the front form analog of the Sokolov packing operators [36].

The explicit expression for \( \hat{A} \) is rather complicated and essentially model dependent [40]. In addition, the operator satisfying Eq. (35) is not unique since it is defined up to unitary operators commuting with \( S \).

The final step specifying our model is as follows. Since the expressions (33-35) do not explicitly depend on \( n \), we assume that field theory models also can be described in such a way. In this case, taking into account the experience with the finite values of \( n \), we have grounds to believe that the operator \( \hat{A} \) is
highly nontrivial and can be determined only beyond perturbation theory.

If $\chi$ is the internal nucleon wave function and $P'$ is its total four-momentum, then its wave function in the representation (34) has the form

$$|N\rangle = 2(2\pi)^3 P'^+ \delta^{(2)}(P_\perp - P'_\perp) \delta(P^+ - P'^+) \chi \quad (36)$$

As follows from Eq. (9), Eq. (1) can be written in the form

$$W^{\mu\nu} = \frac{1}{4\pi} \sum_X (2\pi)^4 \delta^{(4)}(P' + q - P_X) \langle N|\hat{J}^\mu|X\rangle \langle X|\hat{J}^\nu|N\rangle \quad (37)$$

where the sum is taken over all possible final states $X$ with the four-momenta $P_X$.

We will consider the matrix elements of the current operator in the reference frames where $P'_z$ is positive and very large. Then $P'^-$ is very small and therefore the value of $P'$ is approximately the same as in the case of the free momentum operator. In the Bjorken limit it is also possible to find such frames where $|P_X| \gg m$. We suppose as usual that the FSI of the struck quark with the remnants of the target is a higher twist effect, i.e. the effect which is suppressed as $(m/Q)^{2n} (n = 1, 2...).$ Then there exist reference frames where the values of $P_X$ also are approximately the same as for the free four-momentum operator.

As follows from Eq. (11)

$$[\hat{M}^{-j}, \hat{J}^+] = -i\hat{J}^j, \quad [\hat{M}^{-j}, \hat{J}^l] = -i\delta_{jl} \hat{J}^- \quad (38)$$

These expressions make it possible to obtain all the components of $\hat{J}^\mu$ if $\hat{J}^+$ is known.

Suppose that in the Bjorken limit the action of $\hat{J}^+$ in the reference frames under consideration can be replaced by the action of the free operator $J^+$ and consider first the case when
$A = 1$. The action of $\hat{M}$ on $|X\rangle$ can be replaced by the action of $M$ and therefore the interaction dependence of the operators $\hat{M}^{-j}$ manifests itself only when $\hat{M}$ acts on $|N\rangle$. However since $m/P_+^\prime$ is small, this dependence can be neglected. We conclude that, as a consequence of the first expression in Eq. (38), the matrix elements of the operator $\hat{J}^j$ in the reference frames under consideration can be calculated with $\hat{J}^j$. Analogously, as a consequence of the second expression, the matrix elements of $\hat{J}^-$ can be calculated with $J^-.$

We see that if $A = 1$ then in the reference frames described above the matrix elements of $\hat{J}^\mu$ can be calculated with the operator $J^\mu$ and the interaction dependence of the four-momentum operator can be neglected too. Therefore the tensor (1) can be calculated with the operator $J^\mu(x)$.

However, as explained above, the interaction dependence of the operators $\hat{M}^{-j}$ comes not only from the interaction dependence of the mass operator but also from the interaction dependence of the operator $S$. In this case the operators $\hat{M}^{-j}$ are unitarily equivalent to the operators with $A = 1$. Then the calculation of matrix elements in the reference frames where $P_\prime$ and $P_X$ are directed along the $z$ axis by using $J^\mu(x)$ does not contradict Poincare invariance, but it is not clear where the current operator as a whole is local [33]. At the same time, since $A$ becomes unity when all interactions in the system under consideration are turned off, it is obvious that the matrix elements of the operator $\hat{J}^\mu(x)$ in the reference frames under consideration can be calculated with the operator

$$\hat{J}^\mu(x) = AJ^\mu(x)A^{-1}$$

which is local since $J^\mu(x)$ is local.
As noted in Sect. 1, the operators considered in perturbation theory properly commute with the free representation generators of the Poincare group. Therefore the obvious generalization of Eq. (39) is

\[ \hat{J}_\mu(x) = AJ_\mu_{\text{pert}}(x)A^{-1} \]  

(40)

where the operator \( J_\mu_{\text{pert}} \) can be calculated in perturbative QCD.

Let us stress that Eqs. (39) and (40) are considered not as exact operator equalities but only as the relations which can be used if the matrix elements in Eq. (37) are calculated in some reference frames (for example, these relations cannot be used in the reference frames where \( P'_z \) is negative and \( |P'_z| \) is large).

As already mentioned, the OPE has been proved only in perturbation theory but several authors tried to prove the OPE beyond that theory in 1+1 dimensional models [12]. Let us note in this connection that in the \((t, z)\) space the operator \( \hat{S} \) is absent (moreover, the Lorentz group is one-dimensional and the only generator of this group is \( M^{+-} \)) and therefore we can choose \( A = 1 \). Hence the validity of the OPE beyond perturbation theory in 1+1 dimensions does not necessarily imply that the same takes place in 3+1 dimensions.

As follows from Eq. (40),

\[ \hat{J}_\mu(x/2) \hat{J}^\mu(-x/2) = AJ_\mu_{\text{pert}}(x/2)J^\mu_{\text{pert}}(-x/2)A^{-1} \]  

(41)

Now we apply the OPE to \( J_\mu_{\text{pert}}(x/2)J^\nu_{\text{pert}}(-x/2)\):

\[ J_\mu_{\text{pert}}(x/2)J^\nu_{\text{pert}}(-x/2) = \sum_i C_i(x^2)x_{\mu_1} \cdots x_{\mu_n} \tilde{O}_{\mu_1 \cdots \mu_n} \]  

(42)

where the operators \( \tilde{O}_{\mu_1 \cdots \mu_n} \) depend only on field operators and their covariant derivatives at the origin of Minkowski space. As
follows from Eqs. (1), (41) and (42), the $Q^2$ evolution of the structure functions is defined by the Wilson coefficients as well as in the standard theory.

Instead of $\chi$ we now introduce a new internal wave function $\tilde{\chi}$ such that $\chi = A \tilde{\chi}$. Let us also introduce the functions

$$
\rho_i(\xi_i) = \sum_{\sigma_i} \int |\chi(\xi_i, k_{i\perp}, \sigma_i, int_i)|^2 \frac{d^2 k_{i\perp} d\rho(int_i)}{2(2\pi)^3 \xi_i(1 - \xi_i)},
$$

$$
\tilde{\rho}_i(\xi_i) = \sum_{\sigma_i} \int |\tilde{\chi}(\xi_i, k_{i\perp}, \sigma_i, int_i)|^2 \frac{d^2 k_{i\perp} d\rho(int_i)}{2(2\pi)^3 \xi_i(1 - \xi_i)} \quad (43)
$$

where $int_i$ means the internal variables for the system $(1, \ldots i - 1, i + 1, \ldots)$ and the integration over these variables is assumed. As follows from Eqs. (31-33) and (36), the quantity $\rho_i(\xi_i) d\xi_i$ is the probability of the event that the light cone momentum fraction of quark $i$ is in the interval $(\xi_i, \xi_i + d\xi_i)$. For this reason the functions $\rho_i(\xi_i)$ are just the PDFs considered in the FT (streaktly speaking they coincide with the PDFs only in the light cone gauge for the gluon field [6, 7]).

In the parton model the current operator is replaced by the free one and Eqs. (36) and (37) make it possible to explicitly calculate the structure functions of DIS.

A standard calculation for the unpolarized structure functions $F_1$ and $F_2$ (see e.g. ref. [41]) gives

$$
F_1(x_{Bj}) = \frac{1}{2} \sum_i e_i^2 \rho_i(x_{Bj}), \quad F_2(x_{Bj}) = x_{Bj} \sum_i e_i^2 \rho_i(x_{Bj}) \quad (44)
$$

where $e_i$ is the electric charge of particle $i$. This result shows that the Bjorken variable $x_{Bj}$ can be interpreted as the light cone momentum fraction of the struck quark. At the same time
if the current operator is given by Eq. (39) we have obviously

\[ F_1(x_{Bj}) = \frac{1}{2} \sum_i e_i^2 \tilde{\rho}_i(x_{Bj}), \quad F_2(x_{Bj}) = x_{Bj} \sum_i e_i^2 \tilde{\rho}_i(x_{Bj}) \]  \hspace{1cm} (45)

Since the wave function of the moving nucleon is defined by the function \( \chi \) and not \( \tilde{\chi} \) (see Eq. (36)), only the functions \( \rho_i \) are universal while the functions \( \tilde{\rho}_i \) are not. As follows from what has been said about the operator \( A \), the relations between them are rather complicated and the functions \( \tilde{\rho}_i \) are not defined uniquely.

As noted in the preceding section, the sum rules in DIS depend on the form of the operators \( \hat{O}_i^{\mu_1...\mu_n} \). If the matrix elements of the current operator are calculated in the reference frames discussed above, then, as follows from Eqs. (6) and (42), we can use the relation

\[ \hat{O}_i^{\mu_1...\mu_n} = A \hat{O}_i^{\mu_1...\mu_n} A^{-1} \]  \hspace{1cm} (46)

Therefore the operators \( \hat{O}_i^{\mu_1...\mu_n} \) in our model indeed contains the nonperturbative part but this part is present only in the operator \( A \).

The momentum carried by quarks is defined by sum rules for the functions \( \xi_i \rho_i(\xi_i) \) and in the general case

\[ \int_0^1 \xi_i \rho_i(\xi_i) d\xi_i \neq \int_0^1 \xi_i \tilde{\rho}_i(\xi_i) d\xi_i \]  \hspace{1cm} (47)

Therefore if the current operator is given by Eq. (39) or Eq. (40), then DIS experiments cannot directly determine the fraction of the nucleon momentum carried by quarks. Analogously these experiments cannot directly determine the fraction of the nucleon spin carried by quarks. It is also unclear how one can extract from DIS experiments the PDFs which should be used for the description of Drell-Yan pairs and hard \( pp \) collisions.
As follows from the above discussion we have to choose between the following possibilities

- If we want to work only with universal PDFs, we have to work in the representation where the current operator is given by Eqs. (39) or (40). Then the amplitudes of lepton-parton interactions cannot be calculated in perturbation theory. In other words, the condition $2_{FT}$ is satisfied but the condition $1_{FT}$ is not.

- On the other hand, we can exclude the $A$ dependence of the current operator by using a proper unitary transformation. Then we will have the situation described by Eq. (45) when the condition $1_{FT}$ is satisfied but the condition $2_{FT}$ is not.

We see that if the possibility (40) is realized in nature then one can satisfy either $1_{FT}$ or $2_{FT}$ but not the both conditions simultaneously.

Since $\chi$ is obtained from $\tilde{\chi}$ by a unitary transformation, it is clear from Eq. (43) that

$$\int_0^1 \rho_i(\xi_i)d\xi_i = \int_0^1 \tilde{\rho}_i(\xi_i)d\xi_i = 1 \quad (48)$$

It is easy to show that for this reason the Adler, Bjorken and Gross - Llewellyn Smith sum rules for unpolarized DIS [42, 43, 44] are satisfied also in the case when the current operator is given by Eq. (14) (according to the present theory, in leading approximation in QCD running coupling constant these sum rules are given by the parton model). Therefore the fact that the operators $\hat{O}_i^{\mu_1...\mu_n}$ and $\tilde{O}_i^{\mu_1...\mu_n}$ differ each other does not affect the sum rules [42, 43, 44].
At the same time the question arises whether other sum rules are satisfied. In the framework of the OPE the validity of the Bjorken sum rule for polarized DIS [21] is a consequence of the fact that $\hat{O}_A^\mu$ coincides with the axial current operator $\hat{J}_A^\mu$ (see Eqs. (7) and (10)). In our model this sum rule will be satisfied if the axial current operator in the $\beta$ decay satisfies Eq. (40) at $x = 0$. However such a choice is defined by low-energy physics and is an additional assumption. The fact that the Bjorken sum rule [21] involves an assumption about the relation between the current operators in DIS and in the $\beta$ decay has been pointed out by several authors (see e.g. ref. [45]).

7 Discussion

The results of the present paper give grounds to conclude that if the operators $\hat{J}^\mu(x)$ in the instant form and $\hat{J}^\mu(x^-, x_\perp)$ in the front one can be correctly defined then some of them inevitably depend on the nonperturbative part of the quark-gluon interaction which contributes to DIS even in the Bjorken limit. Then the question arises whether the canonical equal time or light cone commutation relations between the components of the current operators are satisfied and whether the deviation of these relations from canonical ones can be investigated in perturbation theory (as noted in ref. [16], at present the only known way of investigating the commutation relations is the renormalizable perturbation theory). The main problem is whether these results can add something to the understanding of the questions formulated in Sect. 1.

In Sect. 6 we have considered a model where the nonper-
turbative contribution to the current operator is contained only in the unitary operator $A$ (see Eq. (40)). If this possibility is realized in nature then one can satisfy either $1_{FT}$ or $2_{FT}$ but not the both conditions simultaneously.

In Sect. \[ we argue that difficulties in substantiating Eq. (6) beyond perturbation theory exist not only in the next-to-leading orders in $1/Q$ (see ref. [32] and references therein) but also in the Bjorken limit. Our results definitely show that the usual physical arguments based on asymptotic freedom at small distances are insufficient to substantiate the expansion (6) beyond perturbation theory, and, if it takes place, this is a consequence of deeper reasons.

The important mathematical fact which is demonstrated by the model considered in Sect. \[ is that it is not necessary to require the validity of Eq. (6) in all reference frames. In other words, it is not necessary to require for the operator product expansion to be a true operator relation. For the validity of the standard results about the $Q^2$ evolution of the structure functions it is quite sufficient to require the validity of Eq. (6) only in special reference frames.

If Eq. (6) is understood in such a way then, as shown in Sect. \[, there exists a possibility that this expression is valid, i.e. the property $1_{OPE}$ takes place. The validity of Eq. (6) is some reference frames will automatically ensure its validity in all reference frames only if the $\hat{O}_{\mu_1...\mu_n}$ are true tensor operators with respect to all transformations of the Lorentz group. However the arguments given in Sects. \[ show that the problem of constructing such operators is rather complicated and the property $2_{OPE}$ is problematic.

The model considered in Sect. \[ shows that there exists a
possibility that many results of the standard theory (including the $Q^2$ evolution of the structure functions) take place even if the current operator contains a nonperturbative contribution which survives in the Bjorken limit. At the same time one has to take into account that, although the present theory has many impressive successes in describing experimental data, several problems remain. For example, as noted in the recent review paper [7], "...On the other hand, these analyses also are calling into question, for the first time, the ultimate consistency of the existing theoretical framework with all existing experimental measurements! (This can be regarded as testimony to the progress made in both theory and experiment - considering the fact that contradictions come with precision, and they are a necessary condition for discovery overlooked shortcomings and/or harbingers of new physics.)". The existence of difficulties in the present approach has been also demonstrated in the recent experiments at Tevatron and HERA [46]. It is interesting to note that the most serious difficulties appear at very large values of the momentum transfer, i.e. when according to the usual philosophy there should be no doubt about the validity of the standard approach.

The possibility that the nonperturbative part of the current operator is important even in the Bjorken limit was considered by several authors (see e.g. ref. [47] and references therein). Our consideration differs from that of those authors since we tackle the problem of nonperturbative effects by considering the commutation relations between the current operator and the representation generators of the Poincare group.

In view of the above discussion it is very important to know what is the contribution of the nonperturbative effects to the current operator but at the present stage of QCD this contribu-
tion cannot be calculated. It is also important to know which experiments can shed light on our understanding of the structure of the current operator. In ref. [48] we argue that deuteron DIS at large $Q$ and $x_{Bj} \leq 0.01$ is such an experiment.

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