Renormalization of Twist-4 Operators in QCD Bjorken and Ellis-Jaffe Sum Rules

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Abstract

The QCD effects of twist-4 operators on the first moment of nucleon spin-dependent structure function $g_1(x, Q^2)$ are studied in the framework of operator product expansion and renormalization group method. We investigate the operator mixing through renormalization of the twist-4 operators including those proportional to the equation of motion by evaluating off-shell Green’s functions in the usual covariant gauge as well as in the background gauge. Through this procedure we extract the one-loop anomalous dimension of the spin 1 and twist-4 operator which determines the logarithmic correction to the $1/Q^2$ behavior of the contribution from the twist-4 operators to the first moment of $g_1(x, Q^2)$.

*JSPS Research Fellow
†Supported in part by the Monbusho Grant-in-Aid for Scientific Research No. C-06640392
In the last several years there has been much interest in nucleon’s spin structure functions $g_1(x, Q^2)$ and $g_2(x, Q^2)$, which can be measured by deep inelastic scattering of polarized leptons on polarized targets. Recent experiments on the nucleon spin structure functions carried out at CERN [1, 2] and SLAC [3, 4], have stimulated intensive theoretical studies on the nucleon spin structure functions [5].

In the deep inelastic scattering, the perturbative QCD has been tested so far for the effects of the leading twist operators, namely twist-2 operators, for which the QCD parton picture holds. Now the spin structure functions would provide us with a good place to investigate higher-twist effects. Our purpose in this paper is to study the renormalization of higher-twist operators, especially the twist-4 operators, which are relevant for the first moment of $g_1(x, Q^2)$, that corresponds to the Bjorken and Ellis-Jaffe sum rules [6, 7]. The anomalous dimension of the twist-4 operators determines the logarithmic correction to the $1/Q^2$ behavior of the twist-4 operator’s contribution to the first moment of $g_1$.

The first moment of the $g_1(x, Q^2)$ structure functions for proton and neutron turns out to be up to the power correction of order $1/Q^2$:

$$
\Gamma_1^{p,n}(Q^2) \equiv \int_0^1 g_1^{p,n}(x, Q^2)dx = (\pm \frac{1}{12} g_A + \frac{1}{36} a_8)(1 - \frac{\alpha_s}{\pi} + \mathcal{O}(\alpha_s^2)) + \frac{1}{9} \Delta \Sigma (1 - \frac{33 - 8N_f}{33 - 2N_f} \frac{\alpha_s}{\pi} + \mathcal{O}(\alpha_s^2)),
$$

where $g_1^{p(n)}(x, Q^2)$ is the spin structure function of the proton (neutron) and the plus (minus) sign is for proton (neutron). On the right-hand side, $g_A \equiv G_A/G_V$ is the ratio of the axial-vector to vector coupling constants. Here we assume that the number of active flavors in the current $Q^2$ region is $N_f = 3$. Denoting $\langle p, s|\bar{\psi}\gamma_\mu\gamma_5\psi|p, s\rangle = \Delta q s_\mu$, the flavor-$SU(3)$ octet and singlet part, $a_8$ and $a_0 = \Delta \Sigma$ are given by

$$
a_8 \equiv \Delta u + \Delta d - 2\Delta s, \quad \Delta \Sigma \equiv \Delta u + \Delta d + \Delta s,
$$

(2)
and $\Delta \Sigma$ is related to the scale-dependent density $\Delta \Sigma(Q^2)$ which evolves as

$$\Delta \Sigma(Q^2) = \Delta \Sigma \left( 1 + \frac{6N_f}{33 - 2N_f} \frac{\alpha_s(Q^2)}{\pi} \right),$$

(hence $\Delta \Sigma$ is the density at $Q^2 = \infty$). Here we have suppressed the target mass effects, which can be taken into account by the Nachtmann moments [8]. Note that taking the difference between $\Gamma_1^p$ and $\Gamma_1^n$ leads to the QCD Bjorken sum rule, the first order QCD correction of which was calculated in [9, 10, 11] and the higher order corrections were given in [12, 13, 14, 15].

Now, the twist-4 operator gives rise to $O(1/Q^2)$ corrections [16, 17] to the first moment of $g_1(x,Q^2)$. As can be seen from the dimensional counting, there is no contribution from the four-fermi type twist-4 operators to the first moment of $g_1(x,Q^2)$. The only relevant twist-4 operators are of the form bilinear in quark fields and linear in the gluon field strength. This is in contrast to the unpolarized case [18], where both types of twist-4 operators contribute. The common feature for the renormalization of higher-twist operator is that there appear a set of operators proportional to equations of motion, which we call EOM operators [19, 20]. And there exists the operator mixing among twist-4 operators which can be studied in the off-shell Green’s functions where the EOM operators are inevitable. It should be emphasized that we have to keep the EOM operators to extract the physical observables like anomalous dimensions, which will be discussed later.

The relevant operators in our case has the following properties: The dimension of the operators is 5 and the spin is 1. Its parity is odd and it has to satisfy the charge conjugation invariance. The flavor non-singlet operators are bilinear in fermion fields. Here we have to consider gauge variant EOM operators as well.

Thus we have the following six operators which satisfy the above conditions:

$$R_1^\sigma = -\bar{\psi}\gamma_5\gamma^\sigma D^2\psi, \quad R_2^\sigma = g\bar{\psi}\tilde{G}^{\mu\nu}\gamma_\mu\gamma_\nu\psi,$$
\[ E_1^\sigma = \overline{\psi} \gamma_5 D^\sigma \psi - \overline{\psi} \gamma_5 D_\sigma \psi - \overline{\psi} \gamma_5 \gamma^\sigma D \psi, \]
\[ E_2^\sigma = \overline{\psi} \gamma_5 D^\sigma \psi + \overline{\psi} \gamma_5 D_\sigma \psi + \overline{\psi} \gamma_5 \gamma^\sigma D \psi, \]
\[ E_3^\sigma = \overline{\psi} \gamma_5 D^\sigma \psi + \overline{\psi} \gamma_5 \partial^\sigma \psi, \quad E_4^\sigma = \overline{\psi} \gamma_5 \gamma^\sigma \partial \psi + \overline{\psi} \gamma_5 \partial \gamma^\sigma \psi, \quad (4) \]

where \( D_\mu = \partial_\mu - igA_\mu T^a \) is the covariant derivative and \( \tilde{G}_{\mu\nu} = \frac{1}{2} \varepsilon_{\mu
u\alpha\beta} G^{\alpha\beta} \) is the dual field strength. And we work with massless quarks for simplicity of the argument.

Here one should note that not all of the above operators are independent, as in the case of twist-3 operators in \( g_2(x,Q^2) \) [21], and they are subject to the following constraint:
\[ R_1^\sigma = R_2^\sigma + E_1^\sigma, \quad (5) \]
where we have used the identities, \( D_\mu = \frac{1}{2} \{ \gamma_\mu, \partial \} \) and \([D_\mu,D_\nu] = -igG_{\mu\nu} \). Therefore any five operators out of (4) are independent and they mix through renormalization.

Here we take \((R_2,E_1,E_2,E_3,E_4)\) to be the basis of the independent operators. The only operator which actually contributes to the physical matrix element responsible for the Bjorken sum rule is \( R_2 \). This twist-4 operator corresponds to the trace-part of twist-3 operator, \((R_{\tau=3})_{\sigma_\mu_1\mu_2} = g \overline{\psi} \tilde{G}_\sigma \{ \gamma_\mu_1, \gamma_\mu_2 \} \psi - \text{traces}, \) but there is no relation between the basis for the twist-4 and that for the twist-3 operators.

We now study the renormalization of the operators. The composite operators, \( O_i \), are renormalized by introducing the renormalization constants \( Z_{ij} \) as
\[ (O_i)_R = \sum_j Z_{ij} (O_j)_B, \quad (6) \]
where the suffix \( R \) (\( B \)) denotes renormalized (bare) quantities. For the present basis we have the following renormalization mixing matrix:
\[ \begin{pmatrix} R_2 \\ E_1 \\ E_2 \\ E_3 \\ E_4 \end{pmatrix}_R = \begin{pmatrix} Z_{11} & Z_{12} & Z_{13} & Z_{14} & Z_{15} \\ 0 & Z_{22} & Z_{23} & Z_{24} & Z_{25} \\ 0 & Z_{32} & Z_{33} & Z_{34} & Z_{35} \\ 0 & 0 & 0 & Z_{44} & 0 \\ 0 & 0 & 0 & 0 & Z_{55} \end{pmatrix} \begin{pmatrix} R_2 \\ E_1 \\ E_2 \\ E_3 \\ E_4 \end{pmatrix}_B. \quad (7) \]
The general features for the mixing matrix are the following \cite{19, 20, 22}: (1) The counter terms for the EOM operators are supplied by the EOM operators themselves. This is because the on-shell matrix elements of the EOM operators ought to vanish. (2) A certain type of operators do not get renormalized. And if we take those operators as one of the independent base, the calculation becomes much simpler. (3) The gauge variant operators also contribute to the mixing.

We compute $Z_{ij}$ by evaluating the off-shell Green’s function of twist-4 composite operators keeping the EOM operators as independent operators. Thus we can avoid the subtle infrared divergence which may appear in the on-shell amplitude with massless particle in the external lines. Another advantage to study the off-shell Green’s function is that we can keep the information on the operator mixing problem. And further, the calculation is much more straightforward than the one using the on-shell conditions.

At the tree level, $R_2$ operator contributes to the 3-point functions with quarks $\psi$, $\bar{\psi}$ and a gluon, $A_\mu$ in the external lines. So we consider the following one-particle irreducible (1PI) Green’s function:

$$\Gamma_{O_\sigma}^{\psi\bar{\psi}A} \equiv \langle 0 | T(O_\sigma(p')A_\rho^a(l)\bar{\psi}(p)) | 0 \rangle_{1PI}^\text{1PI},$$

where the fields and the coupling constant involved represent the bare quantities. Here we employ the dimensional regularization ($D = 4 - 2\varepsilon$) and take the minimal subtraction scheme. The Green’s functions are renormalized as follows:

$$(\Gamma_{O_i})_R = \sum_j Z_2 \sqrt{Z_3} Z_{ij} (\Gamma_{O_j})_B,$$

where $Z_2$ and $Z_3$ are wave function renormalization constants for quarks and gluon fields. We first present the evaluation in the usual covariant gauge. The one-loop radiative corrections arising from eight diagrams for $R_2$ are represented as:

$$(\Gamma_{R_2}^{\psi\bar{\psi}A})_{\text{1-loop}} = \left\{ 1 + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} \left[ -\frac{5}{3} C_2(R) + C_2(G) \right] \right\} (\Gamma_{R_2}^{\psi\bar{\psi}A})_{\text{tree}}$$
\[ + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} \left[ -\frac{3}{2} C_2(R) + \frac{3}{8} C_2(G) \right] (\Gamma_{E_1}^{\psi \bar{\psi} A})_{\text{tree}} \]
\[ + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} \left[ -\frac{1}{6} C_2(R) + \frac{1}{8} C_2(G) \right] (\Gamma_{E_2}^{\psi \bar{\psi} A})_{\text{tree}} \]
\[ + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} \left[ -\frac{1}{4} C_2(G) \right] (\Gamma_{E_3}^{\psi \bar{\psi} A})_{\text{tree}} \]
\[ + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} \frac{1}{4} C_2(G) (\Gamma_{E_4}^{\psi \bar{\psi} A})_{\text{tree}}, \]

where the quadratic Casimir operators are \( C_2(R) = 4/3 \) and \( C_2(G) = 3 \) for QCD. In (10), the tree-level Green’s functions are given by

\[
(\Gamma_{R_2}^{\psi \bar{\psi} A})_{\text{tree}} = ig\varepsilon_{\sigma \rho \alpha \beta} l^\alpha \gamma^\beta T^a,
\]
\[
(\Gamma_{E_1}^{\psi \bar{\psi} A})_{\text{tree}} = g\gamma_5 \gamma_\rho (p + p')_\rho T^a - ig\varepsilon_{\sigma \rho \alpha \beta} l^\alpha \gamma^\beta T^a,
\]
\[
(\Gamma_{E_2}^{\psi \bar{\psi} A})_{\text{tree}} = -2g\gamma_5 g_{\sigma \rho} (p + p')^\sigma T^a - 2g\gamma_5 \gamma_\rho (p + p')_\rho T^a
\]
\[+ g\gamma_5 \gamma_\rho (p + p')_\rho T^a - ig\varepsilon_{\sigma \rho \alpha \beta} l^\alpha \gamma^\beta T^a
\]
\[
(\Gamma_{E_3}^{\psi \bar{\psi} A})_{\text{tree}} = -g\gamma_5 \gamma_\rho (p + p')_\rho T^a,
\]
\[
(\Gamma_{E_4}^{\psi \bar{\psi} A})_{\text{tree}} = g\gamma_5 g_{\sigma \rho} (p + p')^\sigma T^a - 2g\gamma_5 \gamma_\rho (p + p')_\rho T^a
\]
\[ - g\gamma_5 \gamma_\rho (p + p')_\rho T^a + ig\varepsilon_{\sigma \rho \alpha \beta} l^\alpha \gamma^\beta T^a \]

Here one can easily see that these five operators have their tree-level 3-point functions as linear combinations of four independent tensor structures. So in order to identify the counter terms properly as given in (10) we need to make use of the conditions for \( Z_{ij} \) extracted from the 2-point functions with \( \psi, \bar{\psi} \) in the external lines.

Note that the tree-level tensor structure for \( R_2, ig\varepsilon_{\sigma \rho \alpha \beta} l^\alpha \gamma^\beta T^a \), appears also in those for \( E_1, E_2 \) and \( E_4 \). Therefore, in order to extract the correct mixing-matrix element, it is crucial to keep the EOM operators. This feature is quite in contrast to the case of twist-2 operators, where we do not have to consider EOM operators at all.

For the Green’s functions of the EOM operators, we have additional Feynman diagrams due to the presence of the two-point vertices at the tree level. Further,
the EOM operators like $E_3$ and $E_4$ which are of the form $E = \overline{\psi} B \frac{\delta S}{\delta \psi}$, where $B$ is independent of fields, do not get renormalized: $Z_{44} = Z_{55} = 1$.

To summarize we get the following result for the renormalization constants. (The detailed calculation will be discussed elsewhere [23]):

$$
\begin{align*}
z_{11} &= \frac{8}{3} C_2(R), & z_{12} &= \frac{3}{7} C_2(R) - \frac{3}{8} C_2(G), \\
z_{13} &= \frac{1}{6} C_2(R) - \frac{1}{8} C_2(G), & z_{14} &= \frac{1}{8} C_2(G), \\
z_{15} &= -\frac{1}{8} C_2(G), & z_{22} &= \frac{1}{3} C_2(R) + \frac{3}{8} C_2(G), \\
z_{23} &= -\frac{1}{3} C_2(R) - \frac{1}{8} C_2(G), & z_{24} &= \frac{1}{4} C_2(G), \\
z_{25} &= \frac{1}{8} C_2(G), & z_{32} &= -\frac{3}{7} C_2(R) - \frac{3}{8} C_2(G), \\
z_{33} &= -\frac{1}{8} C_2(R) + \frac{1}{8} C_2(G), & z_{34} &= -\frac{1}{4} C_2(G), \\
z_{35} &= -\frac{1}{8} C_2(G), & z_{44} &= z_{55} = 0.
\end{align*}
$$

where we have written the renormalization constants as

$$
Z_{ij} \equiv \delta_{ij} + \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} z_{ij}.
$$

This result is in agreement with the general theorem on the renormalization mixing matrix discussed above.

We now determine the anomalous dimension of $R_2^\sigma$ operator. In physical matrix elements, the EOM operators do not contribute [18] and we have

$$
\langle \text{phys}|(R_2^\sigma)_B|\text{phys}\rangle = Z_{11}^{-1} \langle \text{phys}|(R_2^\sigma)_R|\text{phys}\rangle = (1 - \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} \frac{18}{3} C_2(R)) \langle \text{phys}|(R_2^\sigma)_R|\text{phys}\rangle.
$$

Therefore the anomalous dimension $\gamma_{R_2}$ turns out to be

$$
\gamma_{R_2}(g) \equiv Z_{11} \mu \frac{d}{d\mu} (Z_{11}^{-1}) = \frac{g^2}{16\pi^2} \gamma_{R_2}^0 + O(g^4), \quad \gamma_{R_2}^0 = 2z_{11} = \frac{16}{3} C_2(R),
$$

which coincides with the result obtained by Shuryak and Vainshtein [24] based on the background field method [25] in the coordinate space, where they discarded the contribution from the EOM operators by taking the on-shell quark external states using the equations of motion for massless quarks given by

$$
\overline{\mathcal{D}} \psi = \overline{\psi} \mathcal{D} \psi = 0.
$$
Here we also present our result for the renormalization mixing of the twist-4 operators in the background field method [26]. We shall work with the momentum space. In this method we decompose the gauge field into classical background field and the quantum field as:

\[ A_\mu^a = A_\mu^{a(\text{cl})} + a_\mu^a, \]

and set up the Feynman rule, where we have an additional term in the three-gluon vertex [26] contributing to this calculation. In the background field method, there appear only gauge invariant operators contributing the mixing through renormalization [27]. We take the independent operator basis to the three gauge invariant operators; \( R_2, E_1 \) and \( E_2 \). Here we calculated the Green’s function (8) with \( A_\mu^{\text{cl}} \) as the external gauge field. Taking into account the wave function renormalization constant of the background gauge field, we obtain the renormalization mixing matrix:

\[
\begin{pmatrix}
R_2 \\
E_1 \\
E_2
\end{pmatrix}_R =
\begin{pmatrix}
1 + \frac{8}{3} C_2(R) \hat{\alpha}/\varepsilon & \frac{3}{2} C_2(R) \hat{\alpha}/\varepsilon & \frac{1}{6} C_2(R) \hat{\alpha}/\varepsilon \\
0 & 1 + \frac{1}{2} C_2(R) \hat{\alpha}/\varepsilon & -\frac{1}{2} C_2(R) \hat{\alpha}/\varepsilon \\
0 & -\frac{3}{2} C_2(R) \hat{\alpha}/\varepsilon & 1 - \frac{1}{2} C_2(R) \hat{\alpha}/\varepsilon
\end{pmatrix}
\begin{pmatrix}
R_2 \\
E_1 \\
E_2
\end{pmatrix}_B
\]

where \( \hat{\alpha} = g^2/16\pi^2 \). This result leads to the same physically observable anomalous dimension of \( R_2 \) as given in (15).

Including the twist-4 effect the Bjorken sum rule becomes

\[
\int_0^1 dx \left[ g_1^p(x, Q^2) - g_1^n(x, Q^2) \right] = \frac{1}{6} \left\{ g_A \left( 1 - \frac{\alpha_s(Q^2)}{\pi} + \mathcal{O}(\alpha_s^2) \right) - \frac{8}{9 Q^2} f_3 \left\{ \frac{\alpha_s(Q_0^2)}{\alpha_s(Q^2)} \right\}^{-32/9 \beta_0} \right\},
\]

where \( f_3 \) is the reduced matrix element of \( R_{2\sigma}^3 \), renormalized at \( Q_0^2 \), which is defined for the general flavor indices, with \( t^i \) being the flavor matrices, as

\[
R_{2\sigma}^i = g \bar{\psi} \tilde{G}_{\sigma\nu} \gamma^\nu t^i \psi, \quad \langle p, s | R_{2\sigma}^i | p, s \rangle = f_i s_\sigma \quad (i = 0, \ldots, 8).
\]

So far we have considered the flavor non-singlet part. Now we turn to the flavor singlet component. Here we note that there is only one non-vanishing independent
gluon operator: $\tilde{G}^\alpha \sigma D^\mu G_{\mu \alpha}$. This operator is equal to the flavor-singlet operator $R_{2g}^0$ up to the gluon’s equation of motion:

$$\tilde{G}^\alpha \sigma D^\mu G_{\mu \alpha} = g\bar{\psi} \gamma^\alpha \tilde{G}^\sigma \alpha \psi.$$  \hspace{1cm} (19)

So now we have only to take into account the mixing between $R_{2g}^0 = g\bar{\psi} \gamma^\alpha \tilde{G}^\sigma \alpha \psi$ and

$$E_G^\sigma = \tilde{G}^\alpha \sigma D^\mu G_{\mu \alpha} - g\bar{\psi} \gamma^\alpha \tilde{G}^\sigma \alpha \psi,$$  \hspace{1cm} (20)
in addition to the previous results for the non-singlet part. The mixing between $R_2^0$ and $E_G$ can be studied by computing the Green’s function with two-gluon external lines, $\Gamma_{AA}^{\pi\pi}$, shown in Fig.1. Now we introduce the renormalization constant $Z_{16}$ as

$$(R_2)_R = Z_{11}(R_2)_B + Z_{12}(E_1)_B + Z_{13}(E_2)_B + Z_{14}(E_3)_B + Z_{15}(E_4)_B + Z_{16}(E_G)_B.$$  \hspace{1cm} (21)

From the diagrams of Fig.1, we get for the number of flavors $N_f$:

$$Z_{16} = \frac{1}{\varepsilon} \frac{g^2}{16\pi^2} \times \frac{2}{3} N_f,$$  \hspace{1cm} (22)

hence we obtain the exponent for the singlet part

$$-\frac{\gamma_0^N}{2\beta_0} = -\frac{\gamma_0^N}{3\beta_0} - \frac{2}{3} N_f = -\frac{1}{3\beta_0} \left( \frac{32}{9} + 2\frac{N_f}{3} \right).$$  \hspace{1cm} (23)

Including the twist-4 effects the first moment of $g_1^{p,n}(x, Q^2)$ becomes

$$\Gamma_1^{p,n}(Q^2) \equiv \int_0^1 g_1^{p,n}(x, Q^2)dx$$

$$= (\pm \frac{1}{12} g_A + \frac{1}{36} a_8)(1 - \frac{\alpha_s}{\pi} + O(\alpha_s^2)) + \frac{1}{9} \Delta \Sigma (1 - \frac{33}{33} - \frac{8}{2N_f} \frac{\alpha_s}{\pi} + O(\alpha_s^2))$$

$$- \frac{8}{9Q^2} \left\{ (\pm \frac{1}{12} f_3 + \frac{1}{36} f_8) \left( \frac{\alpha_s(Q_0^2)}{\alpha_s(Q^2)} \right) - \frac{\gamma_8^S}{2N_f} f_0 \left( \frac{\alpha_s(Q_0^2)}{\alpha_s(Q^2)} \right) \right\},$$  \hspace{1cm} (24)

where $f_0$, $f_3$ and $f_8$ are the twist-4 counter parts of $a_0$, $a_3$ and $a_8$. $f_i$’s are scale dependent and here they are those at $Q_0^2$.

If we take into account the ghost terms in our QCD lagrangian, we get extra terms for the gluon EOM operator, which are expressed in terms of the ghost fields and satisfy
the BRST invariance. In addition, there appears the so-called BRST exact operator \[19, 20, 22\] which participates in the operator mixing. However, it turns out that their contributions cancel with each other, and the final result does not change \[23\]. This can be more easily seen in the background gauge where we have only \(E_G\) for the additional independent operator and no ghost fields.

Finally it should be noted that the matrix elements of the twist-4 operators \(f_i\)'s are considered to have ambiguities due to the renormalon singularity as discussed in the literatures \[28\]. However, the exponents of logarithmic corrections to the \(1/Q^2\) behavior, which we computed in the present paper, have definite values without any ambiguity. In case the \(Q^2\) dependence of the moment \(\langle 24 \rangle\) could be measured with enough accuracy in future experiments, we would be able to examine the presence of the twist-4 effects.

We would like to thank K. Tanaka for valuable discussion.
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\textbf{Figure Caption}

Fig.1 The Feynman diagrams for $\Gamma_{R_2}^{AA}$ contributing to the mixing of $R_2^0$ and $E_G$. 

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Fig. 1