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QED corrections in deep-inelastic scattering from tensor polarized deuteron target

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The leading-log model-independent radiative corrections in deep-inelastic scattering of unpolarized electron beam off the tensor polarized deuteron target have been considered. The calculation is based on the covariant parametrization of the deuteron quadrupole polarization tensor and use of the Drell–Yan like representation in electrodynamics to describe the radiation of real and virtual particles by the initial and scattered electron.

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1. Processes with polarized particles are a rich source of a new information on the structure of the nucleon and its fragmentation. Polarized deuterons and nuclei of ^{3}He are used to extract information on the neutron spin-dependent structure function $g_1(x)$ [1]. However, the polarized deuteron is interesting in its own right, because it has spin one. Therefore, due to the deuteron electric-quadrupole structure, other spin-dependent structure functions (as compared with one-half spin particles) appear [2]. The 15 GeV ELFE project provides a good opportunity for the measurement of some hadron tensor structure functions, which could give clues to physics of non-nucleonic components in spin-one nuclei and study the tensor structure on the quark-gluon level [3]. The use of the tensor polarized deuteron target at HERMES allows to investigate the nuclear binding effects and nuclear gluon components [4].

Current experiments at modern accelerators reached a new level of precision and this circumstance requires a new approach to data analysis and inclusion of all possible systematic uncertainties. One of the important source of such uncertainties are the electromagnetic radiative effects caused by physical processes which take place in higher orders of the perturbation theory with respect to the electromagnetic interaction. In present paper we give the covariant description of the cross-section of the deep–inelastic scattering of unpolarized electron beam off the tensor polarized deuteron target

$$e^-(k_1) + d_T(p_1) \to e^-(k_2) + X(p_x) \qquad (1)$$

and we use it to calculate the QED radiative corrections (RC) by means of the electron structure function method [5].

The corresponding approach is based on the covariant parametrization of the deuteron quadrupole polarization tensor in terms of the 4-momenta of the particles in process (1) [6] and use of the Drell–Yan like representation [7] in electrodynamics, which allows to sum the leading–log model–independent RC in all orders. Some applications of this representation for the calculation of RC to the polarization–independent and polarization–dependent contributions to the cross–section in different processes one can find in Ref. [8].

2. To begin with, we define the DIS cross–section of the process (1), with taking into account RC, in terms of the leptonic $L_{\mu\nu}$ and hadronic $H_{\mu\nu}$ tensors contraction

$$\frac{d\sigma}{dQ^2dy} = \frac{\pi\alpha^2}{Vq^4}L_{\mu\nu}H_{\mu\nu},\ y = \frac{2p_1(k_1-k_2)}{V},\ V = 2k_1p_1, \qquad (2)$$

where q is the 4-momentum of the intermediate heavy photon that probes the deuteron structure. Note that only in the Born approximation (without taking into account RC) $q = k_1 - k_2$.

The model–independent RC exhibits themselves by means of the corrections to the leptonic tensor. In the framework of the leading accuracy this tensor can be written as a convolution of two electron structure functions D and the Born form of the leptonic tensor $L^B_{\mu\nu}$ that depends on the scaled electron momenta

$$L_{\mu\nu}(k_1,k_2) = \int\int\frac{dx_1dx_2}{x_1x_2^2}D(x_1,L)D(x_2,L)L^B_{\mu\nu}(\hat{k}_1,\hat{k}_2),$$

$$L = \ln\frac{Q^2}{m^2}, \qquad (3)$$

$$\hat{k}_1 = x_1k_1,\ \hat{k}_2 = \frac{k_2}{x_2},\ \hat{k}_1 - \hat{k}_2 = q,\ Q^2 = -(k_1-k_2)^2,$$

where m is the electron mass.

The limits of integration on the right side of Eq. (3) can be derived from the condition that the DIS process

(1) takes place. It is possible if the final undetected hadron system consists, at least, of a deuteron and a pion. In this case

$$x_1x_2+y-1-x_1xy\geq x_2\delta,\, x=\frac{Q^2}{2p_1(k_1-k_2)},\quad \delta=\frac{(M+m_\pi)^2-M^2}{V}, \tag{4}$$

where M (m_π) is the deuteron (pion) mass. This inequality defines the integration limits as follows

$$1\geq x_1\geq\frac{1+\delta-y}{1-xy},\ 1\geq x_2\geq\frac{1-y+x_1xy}{x_1-\delta}\ . \tag{5}$$

In fact, the representation (3) for the leptonic tensor contains all mass collinear singularities arising due to the radiation of real and virtual photons and electron–positron pairs by the initial and scattered electrons. It reflects the essence of the quasi–real electron method [9] suitable to describe the real collinear particle emission by means of the contribution of the so–called Θ-term into the electron structure function (for the details of the D-function see, for example, Ref. [10]).

The Born leptonic tensor is

$$L^B_{\mu\nu}(k_1,k_2)=-2k_1k_2g_{\mu\nu}+2(k_{1\mu}k_{2\nu}+k_{1\nu}k_{2\mu}). \tag{6}$$

To write the hadron tensor we define first the tensor polarized deuteron density matrix (here we do not consider the effect caused by the vector polarization of the deuteron)

$$\rho_{\mu\nu}=-\frac{1}{3}(g_{\mu\nu}-\frac{p_{1\mu}p_{1\nu}}{M^2})+Q_{\mu\nu},\ Q_{\mu\nu}=Q_{\nu\mu},\quad Q_{\mu\mu}=0,\ p_{1\mu}Q_{\mu\nu}=0, \tag{7}$$

where $Q_{\mu\nu}$ is the deuteron quadrupole–polarization tensor. The corresponding hadron tensor has polarization–independent and polarization–dependent parts

$$\begin{gathered}H_{\mu\nu}=H^{(u)}_{\mu\nu}+H^{(T)}_{\mu\nu},\ H^{(u)}_{\mu\nu}=-W_1\tilde g_{\mu\nu}+\frac{W_2}{M^2}\tilde p_{1\mu}\tilde p_{1\nu},\\ \tilde g_{\mu\nu}=g_{\mu\nu}-\frac{q_\mu q_\nu}{q^2},\ \tilde p_{1\mu}=p_{1\mu}-\frac{p_1q}{q^2}q_\mu,\\ H^{(T)}_{\mu\nu}=aB_1\tilde g_{\mu\nu}+\frac{aB_2}{p_1q}\tilde p_{1\mu}\tilde p_{1\nu}+\\ +\frac{M^2}{(p_1q)^2}B_3q_\alpha(\tilde p_{1\mu}Q_{\tilde\nu\alpha}+\tilde p_{1\nu}Q_{\tilde\mu\alpha})+\frac{M^2}{p_1q}B_4\widetilde Q_{\mu\nu},\\ a=\frac{M^2}{(p_1q)^2}Q_{\alpha\beta}q_\alpha q_\beta.\end{gathered} \tag{8}$$

In general all the hadron structure functions $W_i (i=1,2)$ and $B_j (j=1,2,3,4)$ depend on two independent variables: q^2 and $x'=-q^2/2p_1q$ (within the chosen accuracy $x'=\hat x=x_1xy/(x_1x_2+y-1)$). We used the following notation on the right side of Eq. (8)

$$\begin{gathered}Q_{\mu\tilde\nu}=Q_{\mu\nu}-\frac{q_\nu q_\alpha}{q^2}Q_{\mu\alpha},\ Q_{\mu\tilde\nu}q_\nu=0,\\ \widetilde Q_{\mu\nu}=Q_{\mu\nu}+\frac{q_\mu q_\nu}{q^4}Q_{\alpha\beta}q_\alpha q_\beta-\frac{q_\nu q_\alpha}{q^2}Q_{\mu\alpha}-\\ -\frac{q_\mu q_\alpha}{q^2}Q_{\nu\alpha},\ \widetilde Q_{\mu\nu}q_\nu=0.\end{gathered} \tag{9}$$

Note also that the hadron structure functions B_j are related with the structure functions b_j, introduced in Ref. [2] (HJM), in the following way

$$B_1=-b_1,\ B_2=\frac{b_2}{3}+b_3+b_4,$$
$$B_3=\frac{b_2}{6}-\frac{b_4}{2},\ B_4=\frac{b_2}{3}-b_3.$$

At chosen normalization the elastic limit ($p_x^2=M^2$) can be reached by simple substitution in hadronic tensor

$$\begin{gathered}W_i(x',q^2)\to-\frac{1}{q^2}\delta(1-x')W_i^{(el)},\\ B_j(x',q^2)\to-\frac{1}{q^2}\delta(1-x')B_j^{(el)},\end{gathered} \tag{10}$$

where

$$W_1^{(el)}=-\frac{2}{3}q^2(1-\frac{\rho}{4})G_2^2,$$

$$\begin{gathered}W_2^{(el)}=M^2\Big\{4[1-\frac{\rho}{3}(1-\frac{\rho}{4})]G_1^2+\rho^2G_3^2-\\ -\frac{2}{3}\rho(1-\frac{\rho}{2})G_2^2+\frac{4}{3}\rho(1+\frac{\rho}{2})G_1G_3-\\ -\frac{4}{3}\rho(1-\frac{\rho}{2})G_1G_2+\frac{2}{3}\rho^2G_2G_3\Big\},\end{gathered} \tag{11}$$

$$\rho=\frac{q^2}{M^2},\ G_i=G_i(q^2),$$

$$\begin{gathered}B_1^{(el)}=-\frac{\rho q^2}{4}G_2^2,\\ B_2^{(el)}=-\frac{\rho^2q^2}{8}[(2G_1+G_2)^2+8G_1G_3],\\ B_3^{(el)}=-\frac{\rho^2q^2}{8}(G_2^2+2G_1G_2+4G_2G_3),\\ B_4^{(el)}=\frac{\rho q^2}{2}(1-\frac{\rho}{4})G_2^2.\end{gathered} \tag{12}$$

The elastic deuteron electromagnetic formfactors G_E (electric), G_M (magnetic) and G_Q (quadrupole) can be expressed in terms of the formfactors $G_i (i = 1, 2, 3)$ [11] by means of the following relations

$$G_E = G_1 + \frac{\rho}{2}G_3, \quad G_M = G_2,$$
$$G_Q = G_1 + 2G_2 + \left(2 - \frac{\rho}{4}\right)G_3.$$

3. Because the polarization–independent part of the hadronic tensor depends on the scaled electron momenta only (by means of $q = \hat{k}_1 - \hat{k}_2$) we can write the respective contribution to the cross–section in the form of the Drell–Yan representation in the electrodynamics that takes into account the leading part of RC

$$\frac{d\sigma^{(u)}(k_1, k_2)}{dQ^2 dy} = = \int\int \frac{dx_1 dx_2}{x_2^2} D(x_1, L) D(x_2, L) \frac{d\sigma_B^{(u)}(\hat{k}_1, \hat{k}_2)}{d\widehat{Q}^2 d\hat{y}}, \quad (13)$$

where $\widehat{Q}^2 = x_1 Q^2 / x_2$, $\hat{y} = (x_1 x_2 + y - 1)/x_1 x_2$ and unpolarized Born cross–section reads

$$\frac{d\sigma_B^{(u)}(k_1, k_2)}{dQ^2 dy} = = \frac{2\pi\alpha^2}{VQ^2}\Big[W_1(x, Q^2) + \frac{1 - y - xy\tau}{2xy\tau} W_2(x, Q^2)\Big], \quad (14)$$

$$\tau = \frac{M^2}{V}.$$

As concerns the polarization–dependent contribution to the cross–section $d\sigma^{(T)}$, the situation is somewhat different. In general we cannot use for it the representation (13) with simple substitution

$$\frac{d\sigma^{(u)}}{dQ^2 dy} \to \frac{d\sigma^{(T)}}{dQ^2 dy} \quad (15)$$

in both sides of Eq. (13). The reason is that the axes, respect to which the components of the deuteron quadrupole polarization tensor $Q_{\alpha\beta}$ are defined, can change their directions at the scale transformation of the electron momenta: $k_{1,2} \to \hat{k}_{1,2}$. For example, this occurs obviously when one of the axes is directed along $(\boldsymbol{k}_1 - \boldsymbol{k}_2)$ direction. But substitution (15) can be useful and applicable if all axes remain stabilized under this transformation.

Therefore, first we have to find the set of stabilized axes and write them in covariant form in terms of 4–momenta of the particles participating in the reaction. If we choose the longitudinal direction $\boldsymbol{l}$ along the electron beam and the transverse one $\boldsymbol{t}$ in the plane $(\boldsymbol{k}_1, \boldsymbol{k}_2)$ and perpendicular to $\boldsymbol{l}$, then

$$S_\mu^{(l)} = \frac{2\tau k_{1\mu} - p_{1\mu}}{M},$$
$$S_\mu^{(t)} = \frac{k_{2\mu} - (1 - y - 2xy\tau)k_{1\mu} - xyp_{1\mu}}{d}, \quad (16)$$
$$S_\mu^{(n)} = \frac{2\varepsilon_{\mu\lambda\rho\sigma} p_{1\lambda} k_{1\rho} k_{2\sigma}}{Vd},$$
$$d = \sqrt{Vxyb}, \quad b = 1 - y - xy\tau.$$

One can verify that the set $S_\mu^{(l,t,n)}$ remains stabilized under the scale transformation and

$$S_\mu^{(\alpha)} S_\mu^{(\beta)} = -\delta_{\alpha\beta}, \; S_\mu^{(\alpha)} p_{1\mu} = 0, \; \alpha, \beta = l, t, n.$$

One can make sure also that in the rest frame of deuteron

$$S_\mu^{(l)} = (0, \boldsymbol{l}), \; S_\mu^{(t)} = (0, \boldsymbol{t}), \; S_\mu^{(n)} = (0, \boldsymbol{n}),$$

$$\boldsymbol{l} = \boldsymbol{n}_1, \; \boldsymbol{t} = \frac{\boldsymbol{n}_2 - (\boldsymbol{n}_1\boldsymbol{n}_2)\boldsymbol{n}_1}{\sqrt{1 - (\boldsymbol{n}_1\boldsymbol{n}_2)^2}},$$
$$\boldsymbol{n} = \frac{\boldsymbol{n}_1 \times \boldsymbol{n}_2}{\sqrt{1 - (\boldsymbol{n}_1\boldsymbol{n}_2)^2}}, \; \boldsymbol{n}_{1,2} = \frac{\boldsymbol{k}_{1,2}}{|\boldsymbol{k}_{1,2}|}. \quad (17)$$

If to add one more 4-vector $S_\mu^{(0)} = p_{1\mu}/M$ to the set (16), we receive the complete set of orthogonal 4-vectors with the following properties

$$S_\mu^{(m)} S_\nu^{(m)} = g_{\mu\nu}, \; S_\mu^{(m)} S_\mu^{(n)} = g_{mn}, \; m, n = 0, l, t, n. \quad (18)$$

This allows to express the deuteron quadrupole polarization tensor in general case as follows

$$Q_{\mu\nu} = S_\mu^{(m)} S_\nu^{(n)} R_{mn} \equiv S_\mu^{(\alpha)} S_\nu^{(\beta)} R_{\alpha\beta},$$
$$R_{\alpha\beta} = R_{\beta\alpha}, \; R_{\alpha\alpha} = 0 \quad (19)$$

because the components R_{00}, $R_{0\alpha}$ and $R_{\alpha 0}$ identically equal to zero due to the condition $Q_{\mu\nu} p_{1\nu} = 0$.

Within the leading accuracy, when the radiation of non–collinear particles by the initial and the scattered electrons is not considered, the components R_{nl} and R_{nt} do not contribute and the expansion (19) can be rewritten in the following standard form

$$Q_{\mu\nu} = \big[S_\mu^{(l)} S_\nu^{(l)} - \frac{1}{2} S_\mu^{(t)} S_\nu^{(t)}\big] R_{ll} +$$
$$+ \frac{1}{2} S_\mu^{(t)} S_\nu^{(t)} (R_{tt} - R_{nn}) + \big(S_\mu^{(l)} S_\nu^{(t)} + S_\mu^{(t)} S_\nu^{(l)}\big) R_{lt}, \quad (20)$$

where we took into account that $R_{ll} + R_{tt} + R_{nn} = 0$.

So, if the components of the deuteron polarization tensor are defined in the coordinate system with the axes along the directions $\boldsymbol{l}, \boldsymbol{t}$ and $\boldsymbol{n}$, as given by Eqs. (17), the polarization–dependent contribution to the cross–section of the process (1), with taking into account the leading RC can be written in the same way as polarization–independent one

$$\frac{d\sigma^{(T_s)}(k_1,k_2)}{dQ^2dy} = $$
$$= \int\int \frac{dx_1dx_2}{x_2^2} D(x_1,L)D(x_2,L)\frac{d\sigma_B^{(T_s)}(\hat{k}_1,\hat{k}_2)}{d\widehat{Q}^2d\hat{y}}. \quad (21)$$

Symbol T_s indicate that components of the quadrupole polarization are defined with respect to stabilized set (16). The simple calculation, using (16), (20) and the definition of the Born leptonic (6) and hadronic (8) tensors, gives

$$\frac{d\sigma_B^{(T_s)}(k_1,k_2)}{dQ^2dy} = $$
$$= \frac{2\pi\alpha^2}{yQ^4}\big[S_{ll}R_{ll} + S_{tt}(R_{tt}-R_{nn}) + S_{lt}R_{lt}\big], \quad (22)$$

where

$$S_{ll} = \big[2xb\tau - y(1+2x\tau)^2\big]G +$$
$$+ 2b(1+3x\tau)B_3 + (b-xy\tau)B_4, \quad (23)$$

$$S_{lt} = 2\sqrt{\frac{xb\tau}{y}}\big[2y(1+2x\tau)G + (2-y-4b)B_3 + yB_4\big], \quad (24)$$

$$S_{tt} = 2xb\tau(-G-B_3), \quad G = xyB_1 - \frac{b}{y}B_2\ . \quad (25)$$

4. Consider now the case when components of the deuteron polarization tensor are defined in the coordinate system with the axes along directions $\boldsymbol{L}$, $\boldsymbol{T}$ and $\boldsymbol{N}$ in the rest frame of the deuteron, where

$$\boldsymbol{L} = \frac{\boldsymbol{k}_1-\boldsymbol{k}_2}{|\boldsymbol{k}_1-\boldsymbol{k}_2|}, \quad \boldsymbol{T} = \frac{\boldsymbol{n}_1 - (\boldsymbol{n}_1\boldsymbol{L})\boldsymbol{L}}{\sqrt{1-(\boldsymbol{n}_1\boldsymbol{L})^2}}, \quad \boldsymbol{N} = \boldsymbol{n}. \quad (26)$$

As one can see, directions $\boldsymbol{L}$ and $\boldsymbol{T}$ have became unstable under the scale transformation. The respective covariant form of the set (26) reads

$$S_\mu^{(L)} = \frac{2\tau(k_1-k_2)_\mu - yp_{1\mu}}{M\sqrt{yh}},$$
$$S_\mu^{(t)} = \frac{(1+2x\tau)k_{2\mu} - (1-y-2x\tau)k_{1\mu} - x(2-y)p_{1\mu}}{\sqrt{Vxbh}},$$
$$S_\mu^{(N)} = S_\mu^{(n)}, \quad h = y + 4x\tau, \quad (27)$$

and the expansion of the deuteron polarization tensor is defined in full analogy with (20)

$$Q_{\mu\nu} = \big[S_\mu^{(L)}S_\nu^{(L)} - \frac{1}{2}S_\mu^{(T)}S_\nu^{(T)}\big]R_{LL} +$$
$$+ \frac{1}{2}S_\mu^{(T)}S_\nu^{(T)}(R_{TT}-R_{NN}) +$$
$$+ \big(S_\mu^{(L)}S_\nu^{(T)} + S_\mu^{(T)}S_\nu^{(L)}\big)R_{LT}. \quad (28)$$

To use the Drell–Yan like representation in this case we have to express the unstable 4–vectors $S_\mu^{(A)}$ $(A = L, T)$ through the stabilized 4–vectors $S_\mu^{(\alpha)}$ $(\alpha = l, t)$. These sets are connected by means of orthogonal matrix which describes the rotation in the plane perpendicular to direction $\boldsymbol{n} = \boldsymbol{N}$

$$S_\mu^{(L)} = \cos\theta S_\mu^{(l)} + \sin\theta S_\mu^{(t)},$$
$$S_\mu^{(T)} = -\sin\theta S_\mu^{(l)} + \cos\theta S_\mu^{(t)}, \quad (29)$$
$$\cos\theta = \frac{y(1+2x\tau)}{\sqrt{yh}}, \quad \sin\theta = -2\sqrt{\frac{xb\tau}{h}}.$$

Now we must substitude the expressions (29) into expansion (28) and perform the contraction of the leptonic and hadronic tensors. This procedure leads to the following formula for the polarization dependent part of the cross–section in the considered case of unstable axes

$$\frac{d\sigma^{(T_u)}(k_1,k_2)}{dQ^2dy} = $$
$$= X_{ij}\int\int \frac{dx_1dx_2}{x_2^2} D(x_1,L)D(x_2,L)\frac{d\sigma^{(ij)}(\hat{k}_1,\hat{k}_2)}{d\widehat{Q}^2d\hat{y}}, \quad (30)$$

where we bear in mind the summation over i and j $(i,j = l,t)$. Symbol T_u indicate that the components of the deuteron polarization tensor are defined with respect to unstable axes (27).

The principal point of the representation (30) resides in that matrix X_{ij} depends on non–scaled electron momenta through the angle θ defined in Eq. (29). Its elements read

$$X_{ll} = \frac{1}{4}(1+3\cos 2\theta)R_{LL} +$$
$$+ \frac{1}{4}(1-\cos 2\theta)(R_{TT}-R_{NN}) - \sin 2\theta R_{LT},$$

$$X_{lt} = \frac{1}{4}\sin 2\theta[3R_{LL} - (R_{TT}-R_{NN})] + \cos 2\theta R_{LT},$$

$$X_{tt}=\frac{3}{4}(1-\cos 2\theta)R_{LL}+ \\ +\frac{1}{4}(3+\cos 2\theta)(R_{TT}-R_{NN})+\sin 2\theta R_{LT}, \tag{31}$$

where

$$\cos 2\theta=\frac{y+4x\tau(-1+2y+2xy\tau)}{h}, \\ \sin 2\theta=-\frac{4(1+2x\tau)}{h}\sqrt{xyb\tau}.$$

The partial cross–sections under integral sign on the right side of Eq. (30) depends just on the scaled electron momenta. They are

$$\frac{d\sigma^{(ij)}(k_1,k_2)}{dQ^2dy}=\frac{2\pi\alpha^2}{yQ^4}S_{ij}, \tag{32}$$

where elements of S_{ij} are defined by Eqs. (23)–(25).

To derive the cross–section on the left side of Eq. (30) in the Born approximation one can use the ordinary δ-functions instead of D-functions on the right side of this equation. This leads to the result very similar to the Eq.(22)

$$\frac{d\sigma_B^{(T_u)}}{dQ^2dy}=X_{ij}\frac{d\sigma^{(ij)}}{dQ^2dy}= \\ =\frac{2\pi\alpha^2}{yQ^4}\left[S_{LL}R_{LL}+S_{TT}(R_{TT}-R_{NN})+S_{LT}R_{LT}\right], \tag{33}$$

$$S_{LL}=-hG+2bB_3+ \\ +\frac{B_4}{h}[(1-y)(y-2x\tau)-2xy\tau(y+x\tau)], \\ S_{TT}=\frac{2xb\tau}{h}B_4,\ S_{LT}=2\sqrt{\frac{xb\tau}{y}}(2-y)\big(B_3+\frac{y}{h}B_4\big). \tag{34}$$

Note that relations (34) can be checked independently by the straightforward calculation using the expansion of the deuteron polarization tensor as given by Eq. (28) and covariant expressions $S_\mu^{(L)}$ and $S_\mu^{(T)}$ for the axes defined by Eq. (27).

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