RADIATIVE EVENTS IN DEEP-INELASTIC SCATTERING OF UNPOLARIZED ELECTRON BY TENSOR–POLARIZED DEUTERON. RADIATIVE CORRECTIONS

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UDC 530.145; 530.12

Deep-inelastic scattering (DIS) of unpolarized electron by tensor-polarized deuteron with tagged collinear photon, radiated from the initial electron, is considered. The cross section is derived in the Born approximation. The model-independent QED corrections to the Born cross section are also calculated using an approach based on the account of all essential Feynman diagrams.

Introduction

One of the main objectives of the HERA experiments is the determination of both spin-independent and spin-dependent structure functions of the proton and neutron over a broad range of kinematical variables. For the solution of some problems (for example, the separation of the F_L and F_2 structure functions), it is necessary to measure the DIS cross section at different cms energies. However, running of the collider at reduced beam energies increases some systematic errors, (e.g., luminosity uncertainties), in the experimental analysis [1]. To solve this problem, it was suggested [2] to use radiative events. This method was already used in a measurement of the structure function F_2 [3]. The photon detector (PD) necessary for the detection of the hard photon was a part of the luminosity monitoring system of the H1 and ZEUS experiments. PD was placed in the very forward direction of the incoming electron beam.

The feasibility of using tagged photons at high luminosity electron-positron storage rings, like the ϕ-factory DAPHNE or at B-factories, to measure σ_{had} has been proposed and studied in detail [4]. Preliminary experimental results using this method have been presented recently by the KLOE [5] and BaBar [6] Collaborations. Photon radiation from the initial e^+e^--state, in the events with missing energy, has been successfully used at LEP for the measurement of the number of light neutrinos and for search for the new physics signals.

Polarized deuterons and nuclei of ^{3}He are used to extract information on the neutron spin-dependent structure function $g_1(x)$ [7]. However, the polarized deuteron is of interest in its own right, because it has spin one. Therefore, other spin-dependent structure functions (as compared with one-half spin particles) appear [8]. 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 [9]. The use of the tensor polarized deuteron target at $HERMES$ allows one to investigate the nuclear binding effects and nuclear gluon components [10].

The spin-independent part of the DIS cross section with tagged photon has been investigated recently in detail [11]. The spin-dependent part (caused by the polarized spin-$\frac{1}{2}$ target and longitudinally polarized electron beam) of this cross section that is described by means of the nucleon structure functions g_1 and g_2 was considered in [12]. Now we consider the spin-dependent part of the tagged-photon DIS of an unpolarized electron beam by a tensor-polarized deuteron target

$$e^-(p_1) + d(p) \to e^-(p_2) + \gamma(k) + X. \quad (1)$$

We suggest, as in [11], that the hard photon is emitted very closely to the direction of the incoming electron beam ($\theta_\gamma = \widehat{\mathbf{p_1 k}} \le \theta_0,\ \theta_0 \ll 1$), and PD measures the energy of all photons inside the narrow cone with the opening angle $2\theta_0$ around the electron beam. Simultaneously, the scattered-electron 3-momentum is also measured.

1. Born Approximation

The hadronic tensor of process (1) is parametrized in terms of the tensor structure functions $b_i (i = 1 - 4)$ [8]. As the opening angle of the forward PD is very small, and we consider only the cross section where the tagged photon is integrated over the solid angle covered by PD, we can apply the quasi-real electron method [13] and parametrize these radiative events using the standard Bjorken variables $x = Q^2/2p(p_1 - p_2)$, $y = 2p(p_1 - p_2)/V$, $V = 2pp_1$, and the energy fraction of the electron after the initial-state radiation of a collinear photon $z = 2p(p_1 - k)/V = (\varepsilon_1 - \omega)/\varepsilon_1$, where ε_1 is the initial-electron energy and ω is the energy deposited in PD.

An alternative set of the kinematic variables, that is specially adapted to the case of the collinear-photon radiation, is given by the shifted Bjorken variables [2]: $\widehat{Q}^2 = -(p_1 - p_2 - k)^2$, $\hat{x} = \widehat{Q}^2/2p(p_1 - p_2 - k)$, $\hat{y} = 2p(p_1 - p_2 - k)/2p(p_1 - k)$, $\widehat{V} = 2p(p_1 - k)$. The relation between the shifted and standard Bjorken variables reads $\widehat{Q}^2 = zQ^2$, $\hat{x} = xyz/(z + y - 1)$, $\hat{y} = (z + y - 1)/z$, $\widehat{V} = zV$. At fixed values of x and y, the lower limit of z can be derived from the constraint on the shifted variable $\hat{x}$: $\hat{x} < 1 \ \rightarrow \ z > (1 - y)/(1 - xy)$.

In the Born approximation, we use the following determination of the DIS cross section of process (1) in terms of the contraction of the leptonic and hadronic tensors (further, we will be interested in the spin-dependent part of the cross section only)

$$\frac{d\sigma}{\hat{y}d\hat{x}d\hat{y}} = \frac{4\pi\alpha^2(\widehat{Q}^2)}{\widehat{Q}^4} zL^B_{\mu\nu}H_{\mu\nu} \ , \tag{2}$$

where $\alpha(\widehat{Q}^2)$ is the running electromagnetic coupling constant that takes into account the effects of the vacuum polarization, and the Born leptonic current tensor in the considered case reads [14]

$$L^B_{\mu\nu} = \frac{\alpha}{4\pi^2}\int_\Omega \frac{d^3k}{\omega}\big[B_g\tilde{g}_{\mu\nu} + B_{11}\tilde{p}_{1\mu}\tilde{p}_{1\nu} + $$
$$+ B_{22}\tilde{p}_{2\mu}\tilde{p}_{2\nu}\big], \tag{3}$$

where Ω covers the solid angle of PD. The hadronic tensor is determined by the standard way [15]. The quantities B_g, B_{11}, and B_{22}, for the case of the initial-state collinear radiation, were calculated in [16]. The trivial integration of (3) over the collinear-photon angular variables gives, in accordance with the quasi-real electron approximation [13],

$$L^B_{\mu\nu} = \frac{\alpha}{4\pi}P(z, L_0)dz\big(-Q^2\tilde{g}_{\mu\nu} + 4z\tilde{p}_{1\mu}\tilde{p}_{1\nu}\big),$$

$$L_0 = \ln\frac{\varepsilon_1^2\theta_0^2}{m^2} \ , \tag{4}$$

$$P(z, L_0) = \frac{1 + z^2}{1 - z}L_0 - \frac{2z}{1 - z} \ ,$$

where m is the electron mass.

To write the hadron tensor, we define first the deuteron spin-density matrix

$$\rho_{\mu\nu} = -\frac{1}{3}\big(g_{\mu\nu} - \frac{p_\mu p_\nu}{M^2}\big) - \frac{i}{2M}\epsilon_{\mu\nu\rho\sigma}s_\rho p_\sigma + Q_{\mu\nu},$$

$$Q_{\mu\nu} = Q_{\nu\mu}, \quad Q_{\mu\mu} = 0 \ , \quad p_\mu Q_{\mu\nu} = 0 \ ,$$

where s_ρ is the 4-vector of the deuteron vector polarization, which satisfies the following conditions: $s^2 = -1, sp = 0$; $Q_{\mu\nu}$ is the deuteron quadrupole-polarization tensor.

The corresponding hadron tensor has polarization-independent and polarization-dependent parts

$$H_{\mu\nu} = H^{(u)}_{\mu\nu} + H^{(T)}_{\mu\nu},$$

$$H^{(u)}_{\mu\nu} = -W_1\tilde{g}_{\mu\nu} + W_2M^{-2}\tilde{p}_\mu\tilde{p}_\nu \ ,$$

$$\tilde{g}_{\mu\nu} = g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2} \ , \tilde{p}_\mu = p_\mu - \frac{pq}{q^2}q_\mu \ ,$$

$$H^{(T)}_{\mu\nu} = aB_1\tilde{g}_{\mu\nu} + \frac{aB_2}{pq}\tilde{p}_\mu\tilde{p}_\nu +$$

$$+\frac{M^2}{(pq)^2}B_3q_\alpha(\tilde{p}_\mu Q_{\tilde{\nu}\alpha} + \tilde{p}_\nu Q_{\tilde{\mu}\alpha}) + \frac{M^2}{pq}B_4\widetilde{Q}_{\mu\nu} \ ,$$

$$a = \frac{M^2}{(pq)^2}Q_{\alpha\beta}q_\alpha q_\beta. \tag{5}$$

The structure functions W_i and B_j depend on two independent invariant variables: q^2 and $\hat{x}$. We used the following notation on the right side of Eq. (5):

$$Q_{\mu\tilde{\nu}} = Q_{\mu\nu} - \frac{q_\nu q_\alpha}{q^2}Q_{\mu\alpha} \ , \quad Q_{\mu\tilde{\nu}}q_\nu = 0 \ ,$$

$$\widetilde{Q}_{\mu\nu} = Q_{\mu\nu} + \frac{q_\mu q_\nu}{q^4}Q_{\alpha\beta}q_\alpha q_\beta - \frac{q_\nu q_\alpha}{q^2}Q_{\mu\alpha} - \frac{q_\mu q_\alpha}{q^2}Q_{\nu\alpha} \ ,$$

$$\widetilde{Q}_{\mu\nu}q_\nu = 0 \ . \tag{6}$$

The structure functions B_j are related to the structure functions b_j, introduced in [8], in the following way: $B_1 = -b_1$, $B_2 = b_2/3 + b_3 + b_4$, $B_3 = b_2/6 - b_4/2$, $B_4 = b_2/3 - b_3$.

In our problem, it is convenient to parametrize the polarization state of the deuteron target in terms of the 4-momenta of the particles participating in the reaction under consideration. Therefore, first, we have to find the set of axes and write them in covariant form in terms of 4-momenta. If we choose the longitudinal direction **l** along the electron beam and the transverse one $\mathbf{t}$ in the plane $(\mathbf{p_1}, \mathbf{p_2})$ normally to $\mathbf{l}$, then

$$S_\mu^{(l)} = \frac{2\tau p_{1\mu} - p_\mu}{M},$$

$$S_\mu^{(t)} = \frac{p_{2\mu} - (1-y-2xy\tau)p_{1\mu} - xyp_\mu}{d},$$

$$S_\mu^{(n)} = \frac{2\varepsilon_{\mu\lambda\rho\sigma}p_\lambda p_{1\rho}p_{2\sigma}}{Vd}\ ,\ \tau = \frac{M^2}{V}, \tag{7}$$

where $d = \sqrt{Vxyb}$, $b = 1 - y - xy\tau$. The set of the 4-vectors $S_\mu^{(l,t,n)}$ has the following properties: $S_\mu^{(\alpha)} S_\mu^{(\beta)} = -\delta_{\alpha\beta}$, $S_\mu^{(\alpha)} p_\mu = 0$, $\alpha, \beta = l, t, n$. In the deuteron rest frame, we have

$$S_\mu^{(l)} = (0, \mathbf{l}),\quad S_\mu^{(t)} = (0, \mathbf{t}),\quad S_\mu^{(n)} = (0, \mathbf{n})\ ,$$

$$\mathbf{l} = \mathbf{n_1},\quad \mathbf{t} = \frac{\mathbf{n_2} - (\mathbf{n_1 n_2})\mathbf{n_1}}{\sqrt{1-(\mathbf{n_1 n_2})^2}},$$

$$\mathbf{n} = \frac{\mathbf{n_1} \times \mathbf{n_2}}{\sqrt{1-(\mathbf{n_1 n_2})^2}},\quad \mathbf{n_{1,2}} = \frac{\mathbf{p_{1,2}}}{|\mathbf{p_{1,2}}|}\ .$$

If to add one more 4-vector $S_\mu^{(0)} = p_\mu/M$ to the set (7), we receive the complete set of orthogonal 4-vectors with the following properties:

$$S_\mu^{(m)} S_\nu^{(m)} = g_{\mu\nu},\quad S_\mu^{(m)} S_\mu^{(n)} = g_{mn},$$

$$m, n = 0, l, t, n. \tag{8}$$

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

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

$$R_{\alpha\beta} = R_{\beta\alpha}, R_{\alpha\alpha} = 0 \tag{9}$$

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

Using the expressions for the leptonic (4) and hadronic (5) tensors and parametrization (9), we derive the spin-dependent part of the cross section of process (1) in the following form:

$$\frac{d\sigma}{d\hat{x}d\hat{y}dz} = \frac{2\pi\alpha^2(\hat{Q}^2)}{\hat{y}\hat{Q}^4}\frac{\alpha}{2\pi}P(z, L_0)\big[\hat{S}_{ll}R_{ll}+$$

$$+\hat{S}_{tt}(R_{tt} - R_{nn}) + \hat{S}_{lt}R_{lt}\big], \tag{10}$$

where

$$\hat{S}_{ll} = \hat{V}\hat{y}\big([2\hat{x}\hat{b}\hat{\tau} - \hat{y}(1+2\hat{x}\hat{\tau})^2]\hat{G}+$$

$$+2\hat{b}(1+3\hat{x}\hat{\tau})B_3 + (\hat{b} - \hat{a})B_4\big),$$

$$\hat{S}_{lt} = 2\hat{V}\sqrt{\hat{x}\hat{y}\hat{b}\hat{\tau}}\big[2\hat{y}(1+2\hat{x}\hat{\tau})\hat{G}+$$

$$+(2-\hat{y}-4\hat{b})B_3 + \hat{y}B_4\big], \hat{S}_{tt} = 2q^2\hat{b}\hat{\tau}(\hat{G} + B_3),$$

$$\hat{G} = \hat{x}\hat{y}B_1 - \frac{\hat{b}}{\hat{y}}B_2\ ,\ \hat{b} = 1 - \hat{y} - \hat{a},$$

$$\hat{\tau} = \frac{M^2}{\hat{V}},\ \hat{a} = \hat{x}\hat{y}\hat{\tau},\ \hat{V} = zV.$$

2. QED Corrections

We restrict ourselves to the model-independent QED radiative corrections (RC) related to the radiation of the real and virtual photons by leptons. The remaining sources of RC in the same order of perturbation theory, such as virtual corrections with the double photon exchange mechanism and bremsstrahlung off the deuteron and partons, are more involved and model-dependent. They are not considered here. Our approach to the calculation of RC is based on the account of all essential Feynman diagrams that describe process (1) in framework of the used approximation. To get rid of cumbersome expressions, we will retain the terms in RC that are accompanied at least by one power of large logarithms. In our case, three different types of such logarithms appear: L_0, $L_Q = \ln(Q^2/m^2)$, and $L_\theta = \ln(\theta_0^2/4)$. Besides, in the chosen approximation, we neglect the terms of the order of θ_0^2, $m^2/\varepsilon_1^2\theta_0^2$, and m^2/Q^2 in the cross section.

To calculate the contribution of the virtual- and soft-photon emission, we use the expression for the one-loop corrected Compton tensor with a heavy photon [16]. Then it is necessary to account for the hard collinear initial-state radiation which is considered here. The obtained expression contains the fictitious photon mass (introduced for the regularization of the infrared divergence). In order to eliminate it, we have to add the contribution due to additional soft-photon emission with the energy less than $\Delta\varepsilon_1$, $\Delta \ll 1$ [17] (for details, see [16]). Using this, we derive the contribution due to the emission of virtual and soft photons to the Born cross section (10) as

$$\frac{d\sigma^{V+S}}{d\hat{x}d\hat{y}dz} = \frac{\alpha^2}{4\pi^2}[\tilde{\rho}P(z, L_0) - T]\Sigma(\hat{x}, \hat{y}, \widehat{Q}^2),$$

$$\Sigma(\hat{x},\hat{y},\widehat{Q}^2)=\frac{2\pi\alpha^2(\hat{Q}^2)}{\hat{y}\hat{Q}^4}\big[\hat{S}_{ll}R_{ll}+$$

$$+\hat{S}_{tt}(R_{tt}-R_{nn})+\hat{S}_{lt}R_{lt}\big], \qquad (11)$$

$$\tilde{\rho}=2(L_Q-1)\ln\frac{\Delta^2}{Y}+3L_Q+3\ln z-\ln^2 Y-$$

$$-\frac{\pi^2}{3}-\frac{9}{2}+2Li_2(\cos^2\frac{\theta}{2})\ ,$$

$$T=\frac{1+z^2}{1-z}L_0[2L_Q\ln z-\ln z(L_0+2\ln(1-z))+$$

$$+\ln^2 z-2Li_2(1-z)]-\frac{4z}{1-z}L_Q\ln z-\frac{2-(1-z)^2}{2(1-z)}L_0,$$

$$Y=\frac{\varepsilon_2}{\varepsilon_1}\ ,\quad \mathrm{Li}_2(z)=-\int\limits_0^z\frac{dx}{x}\ln(1-x),$$

where ε_2 is the scattered-electron energy and θ is the electron scattering angle ($\theta=\widehat{\mathbf{p_1p_2}}$).

Let us consider the emission of an additional hard photon (radiated at arbitrary angle) with 4-momentum $\tilde{k}$ and the energy more than $\Delta\varepsilon_1$

$$e^-(p_1)+d(p)\to e^-(p_2)+\gamma(k)+\gamma(\tilde{k})+X(p_x). \qquad (12)$$

To calculate the contribution due to the real hard bremsstrahlung, which corresponds in our case to double hard photon emission, with at least one photon seen in the forward PD, we divide the phase space of additional hard photon into three kinematical regions (for details, see [12]). Below we use the notation of this paper.

The contribution from kinematic region i), when both hard photons hit PD and every one has the energy more than $\Delta\varepsilon_1$, can be written as follows:

$$\frac{d\sigma_{\gamma\gamma}^{i)}}{d\hat{x}d\hat{y}dz}=\Big(\frac{\alpha}{2\pi}\Big)^2L_0(\frac{1}{2}L_0A+B)\Sigma(\hat{x},\hat{y},\widehat{Q}^2), \qquad (13)$$

$$A=4\frac{1+z^2}{1-z}\ln\frac{1-z}{\Delta}+(1+z)\ln z-2(1-z),$$

$$B=3(1-z)+\frac{3+z^2}{2(1-z)}\ln^2 z-2\frac{(1+z)^2}{1-z}\ln\frac{1-z}{\Delta}.$$

In the case of kinematic region ii), there are two possibilities for the experimental setup depending on the detection method for the scattered electron and the hard photon that is collinear to it: exclusive and calorimetric ones.

For the exclusive event selection, when only the scattered electron is detected, while the hard photon, that is emitted almost collinear to it (i.e. within the opening angle $2\theta_0'$ around the momentum of the scattered electron), goes unnoticed or is not taken into account in the determination of the kinematic variables, we have, in accordance with [13],

$$\frac{d\sigma_{\gamma\gamma}^{ii)\mathrm{excl}}}{\hat{y}d\hat{y}d\hat{x}dz}=\frac{\alpha^2}{4\pi^2}P(z,L_0)\int\limits_{\Delta/Y}^{y_{1\,\max}}\frac{dy_1}{y_s(1+y_1)}\times$$

$$\times\Big[\frac{1+(1+y_1)^2}{y_1}(\widetilde{L}-1)+y_1\Big]\Sigma(x_s,y_s,Q_s^2)\ , \qquad (14)$$

where $y_1=\tilde{\omega}/\varepsilon_2$. The expressions for the quantities $\widetilde{L}$, x_s, y_s, Q_s^2, and y_{1max} can be founed in Ref. [12]. Here the parameter θ_0' is pure auxiliary and escapes the final result when the contribution of region iii) will be added.

From the experimental point of view, the calorimeter event selection is more realistic. Here the hard photon and the electron cannot be distinguished inside a narrow cone with the opening angle $2\theta_0'$ along the outgoing-electron momentum direction. Therefore, only the sum of the hard photon and electron energies can be measured if the photon belongs to this cone. In this case, we obtain

$$\frac{d\sigma_{\gamma\gamma}^{(ii)\mathrm{cal}}}{d\hat{x}d\hat{y}dz}=\frac{\alpha^2}{4\pi^2}P(z,L_0)\int\limits_{\Delta/Y}^{\infty}\frac{dy_1}{(1+y_1)^3}\times$$

$$\times\Big[\frac{1+(1+y_1)^2}{y_1}(\widetilde{L}-1)+y_1\Big]\Sigma(\hat{x},\hat{y},\widehat{Q}^2)=$$

$$=\frac{\alpha^2}{4\pi^2}P(z,L_0)\Big[(\widetilde{L}-1)\Big(2\ln\frac{Y}{\Delta}-$$

$$-\frac{3}{2}\Big)+\frac{1}{2}\Big]\Sigma(\hat{x},\hat{y},\widehat{Q}^2)\ . \qquad (15)$$

Here the parameter θ_0' is the physical one, and the final result depends on it (see below).

The approach of [11] allows us to extract the leading contributions (proportional to $\ln\theta_0$ and $\ln\theta_0'$) in the cross section as well as to separate the infrared singularities. The contribution of kinematic region iii) may be presented as

$$\frac{d\sigma_{\gamma\gamma}^{iii)}}{d\hat{x}d\hat{y}dz}=\frac{\alpha^2}{4\pi^2}P(z,L_0)\Big\{\Big[\hat{y}\int\limits_{\Delta}^{x_{1\,\max}}\{dx_1[z^2+$$

$$+(z-x_1)^2]\}\{x_1z(z-x_1)\}^{-1}\ln\frac{2(1-c)}{\theta_0^2}y_t^{-1}\times$$

$$\times\Sigma(x_t,y_t,Q_t^2)+\hat{y}\int\limits_{\Delta/Y}^{y_{1\,\max}}\frac{dy_1[1+(1+y_1)^2]}{y_1(1+y_1)}\times$$

$$\times\ln\frac{2(1-c)}{\theta_0'^2}y_s^{-1}\Sigma(x_s,y_s,Q_s^2)\Big]+Z\Big\}. \qquad (16)$$

where $x_1=\tilde{\omega}/\varepsilon_1$. The expressions for the quantities x_t, Q_t^2, y_t, and $x_{1\max}$ can be found in [12].

The dependence on the infrared auxiliary parameter Δ as well as on the angles θ_0 and θ_0' is contained in the first two terms on the right side of Eq. (16), whereas the quantity Z does not contain the infrared and collinear singularities. It can be written as

$$Z=\hat{y}\frac{4(1-c)}{zQ^2}\int\limits_0^\infty\frac{du}{1+u^2}\Big\{\int\limits_{-1}^1 dc_2(1-c_2)^{-1}\times$$

$$\times|c_2-c|^{-1}\int\limits_0^{x_m}\frac{dx_1}{x_1}\Big[\Phi(c_2,\tilde{c}_1)-$$

$$-\Phi(c,1)\Big]+\int\limits_c^1\frac{dc_2}{(1-c)(1-c_2)}\int\limits_0^{x_m}\frac{dx_1}{x_1}\times$$

$$\times\Big[\Phi(c,1)-\Phi(1,c)\Big]\Big\}\ , \qquad (17)$$

where we use the following notation: $\tilde{c}_1=c_-+(c_+-c_-)[1+(1-c_+)z^2/(1-c_-)]$, $c_\pm=cc_2\pm\sqrt{1-c^2-c_2^2+c^2c_2^2}$, $c_{1,2}=\cos\theta_{1,2}$, $\theta_{1,2}=\widehat{\mathbf{kp_{1,2}}}$. Quantity $\Phi(c_2,c_1)$ reads

$$\Phi(c_2,c_1)=\frac{\alpha^2(\tilde{q}^2)}{\tilde{q}^4}2\varepsilon_1^3\varepsilon_2 zx_1^2(1-c_1)(1-c_2)S. \qquad (18)$$

The upper limit of the integration x_m can be found in [12]. The term S can be represented as follows: $S=A_0Q_0+A_1Q_1+A_2Q_2+A_{11}Q_{11}+A_{22}Q_{22}$, where $Q_0=Q_{\alpha\beta}q_\alpha q_\beta$, $Q_i=Q_{\alpha\beta}q_\alpha p_{i\beta}$, $Q_{ii}=Q_{\alpha\beta}p_{i\alpha}p_{i\beta}$, $(i=1,2)$ and the functions A_i are the following:

$$A_0=\frac{1}{\tilde{s}\tilde{t}}\frac{M^2}{(p\tilde{q})^2}(2(\tilde{q}^4+\tilde{u}^2-2\tilde{s}\tilde{t})B_1+$$

$$+[\frac{M^2}{p\tilde{q}}(\tilde{q}^4+\tilde{u}^2-2\tilde{s}\tilde{t})-4zpp_1(\tilde{u}+\tilde{t})+$$

$$+4pp_2(\tilde{u}+\tilde{s})+\frac{4\tilde{q}^2}{p\tilde{q}}(z^2(pp_1)^2+(pp_2)^2)]B_2-$$

$$-2B_3[2zpp_1(\tilde{u}+\tilde{t})-2pp_2(\tilde{u}+\tilde{s})]),$$

$$A_1=\frac{4z}{\tilde{s}\tilde{t}}\frac{M^2\tilde{q}^2}{(p\tilde{q})^2}[2zpp_1B_3-\frac{p\tilde{q}}{\tilde{q}^2}(\tilde{u}+\tilde{t})(B_3+$$

$$+B_4)],\ A_{11}=\frac{4z^2}{\tilde{s}\tilde{t}}\frac{M^2\tilde{q}^2}{p\tilde{q}}B_4,$$

$$A_2=\frac{4}{\tilde{s}\tilde{t}}\frac{M^2\tilde{q}^2}{(p\tilde{q})^2}[2pp_2B_3+\frac{p\tilde{q}}{\tilde{q}^2}(\tilde{u}+\tilde{s})(B_3+$$

$$+B_4)],\quad A_{22}=\frac{4}{\tilde{s}\tilde{t}}\frac{M^2\tilde{q}^2}{p\tilde{q}}B_4, \qquad (19)$$

where $\tilde{q}=zp_1-p_2-\tilde{k}$, $\tilde{u}=-2zp_1p_2$, $\tilde{s}=2\tilde{k}p_2$, $\tilde{t}=-2z\tilde{k}p_1$, $\tilde{x}=-\tilde{q}^2/2p\tilde{q}$, $B_i=B_i(\tilde{x},\tilde{q}^2),(i=1-4)$.

If we use the set mentioned above of the four-vectors $S_\mu^{(i)}(i=l,t,n)$ we can write the coefficients $Q_i(i=0,1,2,11,22)$ in the form

$$Q_0=(a_1^2-\frac{1}{2}b_1^2-\frac{1}{2}e_1^2)R_{ll}+2a_1b_1R_{lt}+$$

$$+\frac{1}{2}(b_1^2-e_1^2)(R_{tt}-R_{nn})+2a_1e_1R_{ln}+2b_1e_1R_{tn},$$

$$Q_1=a_1a_2R_{ll}+b_1a_2R_{lt}+e_1a_2R_{ln},$$

$$Q_2=(a_1a_3-\frac{1}{2}b_1b_3)R_{ll}+(a_1b_3+b_1a_3)R_{lt}+$$

$$+\frac{1}{2}b_1b_3(R_{tt}-R_{nn})+e_1a_3R_{ln}+b_3e_1R_{tn},$$

$$Q_{11}=a_2^2R_{ll},\ Q_{22}=(a_3^2-\frac{1}{2}b_3^2)R_{ll}+$$

$$+2a_3b_3R_{lt}+\frac{1}{2}b_3^2(R_{tt}-R_{nn}),$$

where $a_1=[\tau(\tilde{u}+\tilde{t})-zp\tilde{q}]/Mz$, $a_2=-V/2M$, $a_3=-V(b-xy\tau)/2M$,

$$b_1=-\frac{1}{2d}[2xyp\tilde{q}+(b-xy\tau)\frac{\tilde{u}+\tilde{t}}{z}+\tilde{u}+\tilde{s}],$$

$$b_2=e_2=e_3=0,\quad b_3=-\frac{Q^2}{d}b,$$

$$e_1^2=\frac{1}{4z^2d^2}\Big(-[z\tilde{s}+(1-y)\tilde{t}]^2+4xyz[\tau\tilde{s}\tilde{t}+$$

$$+(z\tilde{s}-(1-y)\tilde{t})p\tilde{k}-xyz(p\tilde{k})^2]\Big).$$

3. Total RC

The total RC to the Born cross section is the sum of the virtual and soft corrections and hard-photon emission contribution. In the considered approximation, it is convenient to write this RC as

$$\frac{d\sigma^{\rm RC}}{d\hat{x}d\hat{y}dz} = \frac{\alpha^2}{4\pi^2}\left(\Sigma_i + \Sigma_f\right) . \tag{20}$$

The term Σ_i is independent of the experimental selection rules for the scattered electron and reads

$$\Sigma_i = L_0\Big\{\frac{1}{2}L_0 P_\theta^{(2)}(z) + \frac{1+z^2}{1-z}\Big[3\ln z - -\frac{1}{2}\ln^2 z + \ln^2 Y - 2\ln z\ln Y + +2\ln z\ln(1-z) - 2\ln(1-z) - \frac{\pi^2}{3} + +2{\rm Li}_2(1-z) + 2{\rm Li}_2\Big(\frac{1+c}{2}\Big)\Big] + +\frac{1}{(1-z)}\ln^2 z + \frac{4z}{1-z}\ln(\frac{z}{1-z}) + +\frac{1-16z-z^2}{2(1-z)}\Big\}\Sigma(\hat{x},\hat{y},\widehat{Q}^2) + (z,L_0)\hat{y}\ln\frac{2(1-c)}{\theta_0^2}\times \times\int\limits_0^{u_0}\frac{du}{1-u}P^{(1)}(1-u)y_t^{-1}\Sigma(x_t,y_t,Q_t^2) + +\frac{1+z^2}{1-z}L_0 Z, \quad u_0 = \frac{x_{1\,\max}}{z} . \tag{21}$$

Here we use the notation $P_\theta^{(2)}(z)$ for the Θ-part of the second-order electron structure function $D(z,L)$ [12]. The quantities x_t, y_t, Q_t^2 depend on $u = x_1/z$.

The term Σ_f explicitly depends on the rule for the event selection. It includes the main effect of the scattered-electron radiation. In the case of exclusive event selection, this contribution is

$$\Sigma_f^{\rm excl} = \hat{y}P(z,L_0)\int\limits_0^{y_{1\,\max}} dy_1[(L_Q + \ln Y - 1) P^{(1)}(\frac{1}{1+y_1}) + \frac{y_1}{1+y_1}]y_s^{-1}\Sigma(x_s,y_s,Q_s^2) .$$

Here the parameter θ_0', that separates kinematic regions *ii*) and *iii*), is not physical, and we see that the final result does not contain it. But the mass singularity, that is connected with the scattered-electron radiation, exhibits itself through the L_Q term.

The situation is quite different for the calorimeter event selection. For such an experimental setup, we derive

$$\Sigma_f^{\rm cal} = P(z,L_0)\Big[\hat{y}\ln\frac{2(1-c)}{\theta_0'^2}\int\limits_0^{y_{1\,\max}} dy_1 \times \times P^{(1)}(\frac{1}{1+y_1})\frac{\Sigma(x_s,y_s,Q_s^2)}{y_s} + \frac{1}{2}\Sigma(\hat{x},\hat{y},\widehat{Q}^2)\Big] .$$

For the calorimeter setup, the parameter θ_0' defines the rule of the event selection and has, therefore, the physical sense. The final result depends on it. However, the mass singularity due to photon emission by the final electron is cancelled in accordance with the Kinoshita–Lee–Nauenberg theorem [18]. The absence of the mass singularity indicates clearly that the term containing $\ln\theta_0'$ arises due to the contribution of kinematic region *iii*), where the scattered electron and the photon radiated from the final-state are well separated. That is why no question appears to determine the quantity ε_2 that enters the expression for $y_{1\,\max}$.

Note that the correction to the usually measured asymmetry, which is the ratio of the spin-dependent part of the cross section to the spin-independent one, is not large because the main factorized contribution due to the virtual- and soft-photon emission trends to cancellation in this case. If the experimental information about spin observables is extracted directly from the spin-dependent part of the cross section (for the corresponding experimental method, see [19]) such cancellation does not take place and the factorized correction gives the basic contribution.

We acknowledge numerous useful discussions with N.P.Merenkov.

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Received 28.01.02

РАДІАЦІЙНІ ПОДІЇ У ГЛИБОКО НЕПРУЖНОМУ РОЗСІЯННІ НЕПОЛЯРИЗОВАНОГО ЕЛЕКТРОНА ТЕНЗОРНО ПОЛЯРИЗОВАНИМ ДЕЙТРОНОМ. РАДІАЦІЙНІ ПОПРАВКИ

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Р е з ю м е

Розглянуто глибоко непружне розсіяння неполяризованого електрона тензорно поляризованим дейтроном із міченим колінеарним фотоном, який випромінюється початковим електроном. Переріз обчислено у борнівському наближенні. Також обчислено модельно незалежні КЕД-поправки до борнівського перерізу з використанням підходу, який грунтується на врахуванні всіх істотних діаграм Фейнмана.

РАДИАЦИОННЫЕ СОБЫТИЯ В ГЛУБОКО НЕУПРУГОМ РАССЕЯНИИ НЕПОЛЯРИЗОВАННОГО ЭЛЕКТРОНА ТЕНЗОРНО ПОЛЯРИЗОВАННЫМ ДЕЙТРОНОМ. РАДИАЦИОННЫЕ ПОПРАВКИ

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Р е з ю м е

Рассмотрено глубоко неупругое рассеяние неполяризованного электрона тензорно поляризованным дейтроном с меченым коллинеарным фотоном, который излучается начальным электроном. Сечение вычислено в борновском приближении. Также вычислены модельно независимые КЭД-поправки к борновскому сечению с использованием подхода, основанного на учете всех существенных диаграмм Фейнмана.