Deep Inelastic Scattering on the Deuteron in the Bethe-Salpeter Formalism

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ABSTRACT

The nuclear effects in the spin structure functions of the deuteron g_1 and b_2 are estimated in a fully covariant approach of the Bethe-Salpeter formalism. The construction of the relativistic wave function of the deuteron is discussed in detail. Numerical results for the spin structure functions g_1 and b_2 are compared with nonrelativistic results and relativistic corrections are discussed.

1. Introduction

In this talk we present results of an investigation of the spin structure functions (SF) of the deuteron in the Bethe-Salpeter (BS) formalism and the operator product expansion (OPE) method. The approach was elaborated in Ref.[1, 2]. All calculations are made in the following well-defined approximations:

1. the ladder approximation for the BS equation;
2. twist-2 approximation in the OPE method.

The deuteron is considered as an "exactly" solvable model in the effective meson-nucleon theory[3]. Our consideration of the SF is limited with the impulse approximation.

This research is motivated by a number of experiments on deep inelastic scattering (DIS) of polarized (unpolarized) leptons on light nuclei being carried out in SLAC, CERN, DESY and CEBAF.

2. Basic Formalism

In the one photon approximation the DIS cross section can be calculated in terms of the hadronic tensor $W_{\mu\nu}$ which is the imaginary part of the amplitude for forward, virtual Compton scattering off the deuteron, $2\pi W_{\mu\nu}(P,q) = \operatorname{Im} T_{\mu\nu}(P,q)$,

$$T_{\mu\nu}(P,q) \equiv i \int d^4\xi \, e^{iq\cdot\xi} \langle P,D| T\left(J_\mu(\xi)J_\nu(0)\right) |P,D\rangle_c. \qquad (1)$$

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We work in the rest frame of the target, $P^\mu = (M_d, \mathbf{0})$, and orient the momentum of the virtual photon opposite to the z-axis, $q^\mu = (\nu, \mathbf{0}_\perp, -|\mathbf{q}|)$. As usual, we define the invariant variables: the square of the transferred momentum, $Q^2 \equiv -q^2 > 0$, the virtual photon energy, $\nu = P \cdot q/M_d$, and the Bjorken scaling variable, $x = Q^2/2M_d\nu$.

The hadronic tensor can be decomposed into independent SF which in the Bjorken limit ($\nu \to \infty$, $Q^2 \to \infty$, $x \in (0,1)$) depend on x only, resulting in scaling SF. Our further considerations are performed in the Bjorken limit. There are eight independent SF for the deuteron, as a spin one target, altogether[4]. It is convenient to classify the scaling SF in terms of helicity amplitudes for the forward Compton scattering, $\gamma_\lambda + \text{deuteron}_\mathcal{M} \to \gamma^*_\lambda + \text{deuteron}_\mathcal{M}$:

$$h_{\lambda\mathcal{M},\lambda\mathcal{M}} = \varepsilon^{\mu *}_\lambda W_{\mu\nu}(P,q)\varepsilon^\nu_\lambda, \tag{2}$$

where λ and $\mathcal{M}$ are the spin components of the virtual photon and the deuteron along the quantization axis, and ε_λ is the polarization vector of a helicity λ photon: $\varepsilon^\mu_\pm = \mp\frac{1}{\sqrt{2}}(0,1,\pm\imath,0)$, $\varepsilon^\mu_0 = \frac{1}{\sqrt{Q^2}}(-|\mathbf{q}|,0,0,\nu)$.

We are interested in the spin SF of the deuteron which in the Bjorken limit are characterized with $g_1^D(x)$ and $b_2^D(x)$[4]. These functions are two features of the spin effects in the deuteron. While $g_1(x)$ measures the spin distribution of nucleons, $b_2(x)$ reflects the dependence of the nucleon momentum distributions on the spin projection of the deuteron. In all it is sufficient to calculate three helicity amplitudes, $h_{++,++}$, $h_{+-,+-}$ and $h_{+0,+0}$, and we find

$$g_1(x) = -\frac{1}{4\pi}\mathrm{Im}\left(h_{++,++} - h_{+-,+-}\right), \tag{3}$$

$$b_2(x) = -\frac{1}{8\pi}\mathrm{Im}\left(2h_{+0,+0} - h_{++,++} - h_{+-,+-}\right). \tag{4}$$

Since in the DIS kinematics the behavior of the $T_{\mu\nu}$ is controlled by the behavior of the product of the currents in Eq. (1) in the vicinity of the light cone, $\xi^2 = 0$, it can be derived from Wilson's operator product expansion. The product of two operators is expanded in sets of local operators with the increasing order of their twist. The lowest twist accounts for the leading contributions in the Wilson's series and therefore gives rise to the leading twist SF, i. e. g_1 and b_2. Structure function $g_1(x)$ can be determined by measuring the DIS cross section for a polarized beam to scatter from a polarized target, whereas $b_2(x)$ can be determined for an unpolarized beam, we need to calculate both the symmetrical, $T_{\{\mu\nu\}} = (T_{\mu\nu} + T_{\nu\mu})/2$, and antisymmetrical, $T_{[\mu\nu]} = (T_{\mu\nu} - T_{\nu\mu})/2$, parts of the Compton amplitude. In the OPE of the helicity amplitude $h_{+\mathcal{M},+\mathcal{M}}$ in the leading twist order receives contributions from vector (V) and axial-vector (A) operators with twist 2

$$h^{V(A)}_{+,\mathcal{M},+\mathcal{M}} \propto \sum_{n=0}^{\infty}\sum_{a=\text{fields}} E^{V(A)}_{a,n}(Q^2)\left(\frac{2}{Q^2}\right)^n q_{\mu_1}\dots q_{\mu_n}\langle P|O^{\{\mu_1\dots\mu_n\}}_{V(A),a}|P\rangle, \tag{5}$$

where a stands for fundamental fields of a theory under consideration, $E^i_{a,n}$ are the coefficient function (being c-numbers) and $\mathcal{O}^{\{\mu_1...\mu_n\}}_{i,a}$ is a set of twist two operators totally symmetrical over Lorentz indices. The explicit form of the operators is to be constructed from the field entering the Lagrangian of the theory.

At present time it is impossible to calculate both pieces appearing in the OPE (5). If either is calculated in a self-consistent way, the other remains to be the unknown element. We are interested in the investigation of the nuclear structure in DIS processes, hence we examine the matrix elements in Eq. (5). To this end one needs a field theory within that the deuteron bound state is well described and the explicit form of the operators are then found. We apply the effective meson-nucleon theory with renormalizable interactions to the OPE method. For example, the nucleon operators of the second twist are:

$$\mathcal{O}^{\{\mu_1...\mu_n\}}_{N,V} = \left(\frac{i}{2}\right)^{n-1} \{:\bar{N}(0)\gamma^{\mu_1} \overleftrightarrow{\partial}^{\mu_2} \ldots \overleftrightarrow{\partial}^{\mu_n} N(0):\}, \tag{6}$$

$$\mathcal{O}^{\{\mu_1...\mu_n\}}_{N,A} = \left(\frac{i}{2}\right)^{n-1} \{:\bar{N}(0)\gamma^{\mu_1}\gamma_5 \overleftrightarrow{\partial}^{\mu_2} \ldots \overleftrightarrow{\partial}^{\mu_n} N(0):\}, \tag{7}$$

where N is the spinor nucleon field. The matrix elements of the operators in Eq. (6) and (7) with the BS amplitudes for the deuteron can be calculated explicitly. On the other hand, the coefficients $E^i_{a,n}$ **cannot be** calculated in such approach and they are new constants to be determined from an independent experiment, e.g. from DIS off free nucleons. In the impulse approximation $E^i_{a,n}$ are identical to moments of the SF of the physical nucleon, $E^V_{N,n} = \mathsf{M}_n(F^N_2)$ and $E^A_{N,n} = \mathsf{M}_n(g^N_1)$ with $\mathsf{M}_n(f) = \int\limits_0^1 dx x^{n-1} f(x)$. The corresponding moments of the g^D_1 and b^D_2 in the ladder approximation have the form

$$\mathsf{M}_n(g^D_1) = \mathsf{M}_n(g^N_1)\Theta^{N/D}_{A,n}, \tag{8}$$

$$2\mathsf{M}_n(b^D_2) = \mathsf{M}_n(F^D_2)\Theta^{N/D}_{V,n+1} + \mathsf{M}_n(F^B_2)\Theta^{B/D}_{n+1}, \qquad n = 2, 4, \ldots \tag{9}$$

where $\Theta^{a/D}_{V(A),n}$ are the reduced matrix elements: $\langle P|\mathcal{O}_{\mu_1...\mu_n}(0)|P\rangle = \Theta_n\{P_{\mu_1}\ldots P_{\mu_n}\}$. The last term in the r. h. s. of Eq. (9) is the contribution of meson exchange currents which will not be considered here.

3. The bound state wave function

Our next point concerning the description of the state vector $|P\rangle$ is the homogeneous BS equation. We use the matrix representation of the BS amplitude for the deuteron[2] which is denoted by $\chi(p; P)$. In the momentum space the BS equation for the bound state in the ladder approximation reads

$$K(\epsilon, \mathbf{p})\chi(p; P) + \sum_{B=\sigma,\pi,\omega,\rho,...} \frac{\lambda_B}{4i\pi^3} \int d^4p' \frac{\Lambda(p_1)\Gamma_B\chi(p'; P)\Gamma_B\Lambda(p_2)}{(p-p')^2 - \mu^2_B} = 0, \tag{10}$$

$$K(\epsilon,\mathbf{p}) \equiv \left(E_{\mathbf{p}}^2 - \epsilon^2 - M_d^2/4\right)^2 - \epsilon^2 M_d^2. \tag{11}$$

where $\Lambda(p) = \hat{p} - m$ and $p_i = (\epsilon, \mathbf{p})$ is the 4-momentum of the ith nucleon expressed in terms of relative 4-momenta p, p' and the center-of-mass (c. m.) momentum $P = (M_d, \mathbf{0})$: $p_1 = P/2 + p$ and $p_2 = P/2 - p$; μ_α is the mass of the relevant meson and Γ_α stands for the interaction vertex between the nucleon and relevant boson, $\lambda_B \equiv g_B^2/(4\pi)^2$ (g_B denoting NNB-coupling constant).

The BS amplitude χ and its conjugate $\bar{\chi}$ satisfy the normalization condition given by the matrix element of the electromagnetic current in the impulse approximation

$$\int \frac{d^4p}{i(2\pi)^4} Tr\left\{\bar{\chi}(p;P)\gamma_\mu \chi(p;P)(m - \hat{p}_1)\right\} = 2P_\mu. \tag{12}$$

The BS amplitude is parametrized using the decomposition in terms of the complete set of Dirac matrices

$$\chi(p;P) = \gamma_5 \mathsf{P} + \gamma^5\gamma^0 \mathsf{A}^0 - (\boldsymbol{\gamma}\cdot\vec{\mathsf{V}}) - \gamma_5(\boldsymbol{\gamma}\cdot\vec{\mathsf{A}}) - 2i\gamma^0(\boldsymbol{\gamma}\cdot\vec{\mathsf{T}}_0) - 2\gamma^0\gamma_5(\boldsymbol{\gamma}\cdot\vec{\mathsf{T}}), \tag{13}$$

where P, A^0 are the scalar and $\vec{\mathsf{T}}^0$, $\vec{\mathsf{T}}$ and $\vec{\mathsf{V}}$ are vector functions depending upon the relative 4-momentum p in the c. m. frame. The angular dependence for an angular momentum $J = 1$ and its projection $\mathcal{M}$ owing to the rotational invariance of Eq. (10) is expressed in terms of the spherical and vector spherical harmonics. For example,

$$\mathsf{P}(\epsilon,\mathbf{p}) = \mathsf{P}_1(\epsilon,|\mathbf{p}|)Y_{1\mathcal{M}}(\Omega_p), \quad \vec{\mathsf{T}}(\epsilon,\mathbf{p}) = \sum_{L=0;2} \mathsf{T}_L(\epsilon,|\mathbf{p}|)\mathbf{Y}^L_{1\mathcal{M}}(\Omega_p). \tag{14}$$

The corresponding equations for the radial functions can be found by the partial wave decomposition of the kernel in Eq. (10). The basis states (13) employed so far are not traditional to discuss the properties of the deuteron. It is more convenient to introduce the basis labeled by J, L, S and ρ. In spectroscopic notation of Ref.[5] we have the eight coupled states

$$Y^T = (v_s^o, v_t^e, v_s^e, v_t^o, u^+, u^-, w^+, w^-), \tag{15}$$

which are connected with the adopted states

$$\Psi^T = (\mathsf{A}_1^0, \mathsf{T}_1^0, \mathsf{P}_1, \mathsf{V}_1, \mathsf{X}_0^+, \mathsf{X}_0^-, \mathsf{X}_2^+, \mathsf{X}_2^-), \tag{16}$$

where $\vec{\mathsf{X}}^\pm \equiv \sqrt{2}(\vec{\mathsf{T}} \pm \vec{\mathsf{A}}/2)$, with the orthogonal transformation. While the odd states in relative energy are not mixed

$$v_s^o = \mathsf{A}_1^0, \quad v_t^e = 2i\mathsf{T}_1^0, \tag{17}$$

the even states are connected with the matrix

$$U = \frac{\eta}{2\sqrt{1+\eta^2}} \tag{18}$$

$$\times \begin{pmatrix} -\frac{2}{\eta} & 0 & \sqrt{\frac{2}{3}} & -\sqrt{\frac{2}{3}} & -\frac{2}{\sqrt{3}} & \frac{2}{\sqrt{3}} \\ 0 & \frac{2}{\eta} & \frac{2}{\sqrt{3}} & \frac{2}{\sqrt{3}} & \sqrt{\frac{2}{3}} & \sqrt{\frac{2}{3}} \\ \sqrt{\frac{2}{3}} & -\frac{2}{\sqrt{3}} & \frac{1+\sqrt{1+\eta^2}}{\eta} & \frac{1-\sqrt{1+\eta^2}}{3\eta} & 0 & \frac{2\sqrt{2}}{3}\frac{1-\sqrt{1+\eta^2}}{\eta} \\ -\sqrt{\frac{2}{3}} & -\frac{2}{\sqrt{3}} & \frac{1-\sqrt{1+\eta^2}}{3\eta} & \frac{1+\sqrt{1+\eta^2}}{\eta} & \frac{2\sqrt{2}}{3}\frac{1-\sqrt{1+\eta^2}}{\eta} & 0 \\ \frac{2}{\sqrt{3}} & \sqrt{\frac{2}{3}} & 0 & -\frac{2\sqrt{2}}{3}\frac{1-\sqrt{1+\eta^2}}{\eta} & -\frac{1+\sqrt{1+\eta^2}}{\eta} & \frac{1-\sqrt{1+\eta^2}}{3\eta} \\ -\frac{2}{\sqrt{3}} & \sqrt{\frac{2}{3}} & -\frac{2\sqrt{2}}{3}\frac{1-\sqrt{1+\eta^2}}{\eta} & 0 & \frac{1-\sqrt{1+\eta^2}}{3\eta} & -\frac{1+\sqrt{1+\eta^2}}{\eta} \end{pmatrix},$$

where $\eta \equiv |\mathbf{p}|/m$.

The BS Eq. (10) is solved numerically using standard methods. The kernel of the equation is the one-boson-exchange potential, consisting of the exchange of π, η, δ, σ, ω and ρ mesons. The pion-nucleon interaction is taken to be of the axial-vector type. The cutoff Λ at high momenta is introduced. This is done by inserting a form-factor $F_B(t) = (\mu_B^2 - \Lambda^2)/(t - \Lambda^2)$ at the meson-nucleon vertex. The Wick rotation is applied, yielding the nonsingular two-dimensional integral equation. The eigenvalue problem for the deuteron mass in Eq. (10) can be reduced to the eigenvalue problem in the space of the coupling constants. The input coupling constants λ_B and the masses of the exchange bosons μ_B are taken to be the same as in Ref.[5]. The parameters reproduce phase shifts for the elastic NN-scattering up to $E_{\rm Lab} \leq 250$ MeV and they are the unique ones known in the literature.

The positive-energy components of the wave function, $u^+(\epsilon, |\mathbf{p}|)$ and $w^+(\epsilon, |\mathbf{p}|)$, can be compared with those obtained from nonrelativistic (NR) calculations. To do this we consider the question of the magnitude of probabilities of the various components. By the use of the matrix (18) the integrand in Eq. (12) takes the form

$$\int \frac{d\epsilon_4 d|\mathbf{p}||\mathbf{p}|^2}{(2\pi)^4} (Y^+(\epsilon_4, |\mathbf{p}|), \omega Y(\epsilon_4, |\mathbf{p}|)) = 1, \tag{19}$$

where ω is the diagonal matrix

$$\omega = -diag(1, 1, 1, 1, 1 - \frac{2E_{\mathbf{p}}}{M_d}, \frac{2E_{\mathbf{p}}}{M_d} + 1, 1 - \frac{2E_{\mathbf{p}}}{M_d}, \frac{2E_{\mathbf{p}}}{M_d} + 1). \tag{20}$$

Each integral in Eq. (19) defines the probability of the relevant component. The numerical values are given in Table I. As a result we see that an admixture of the negative-energy amplitudes reinforces the contribution of the positive-energy states.

Table I. The probabilities of the components of the BS amplitude

State	$^3S_1^+$	$^3D_1^+$	$^3S_1^-$	$^3D_1^-$
$P_\alpha(\%)$	95.014	5.106	−0.002	−0.003
State	$^1P_1^e$	$^3P_1^o$	$^1P_1^o$	$^3P_1^e$
$P_\alpha(\%)$	−0.010	−0.082	−0.015	−0.008

The values of the probabilities for the relativistic components are small and close to those in Ref.[5]. The D-state probability of the BS wave function is compared with $P_D^{BS} = 4.8$ % of Ref.[5], $P_D = 4.3$ % of Bonn[6] and $P_D = 5.9$ % Paris[7] potentials.

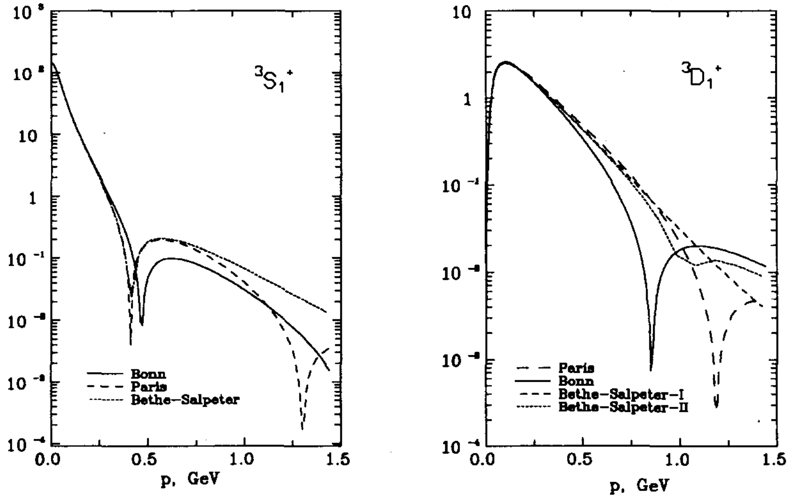

Figure 1: The module of $^3S_1^+$ and $^3D_1^+$ components of the deuteron wave function. For comparison the correspondent components of the NR wave functions of the Bonn and Paris potentials are plotted.

Additionally, the various static characteristics of the deuteron were calculated. The calculation of the quadrupole moment and the magnetic moments does not contrast to the values of the NR nuclear potential models for the same values of P_D. A thorough analysis will be reported elsewhere.

We found that the positive-energy states $^3S_1^+$ and $^3D_1^+$ correspond to the NR states. Not looking for the exact analytical correspondence, we can employ Eq. (19) and define the following quantity: $u^+(|\mathbf{p}|) = \sqrt{\int d\epsilon\, \omega^+ |u^+(\epsilon, |\mathbf{p}|)|^2}$ and $w^+(|\mathbf{p}|) = \sqrt{\int d\epsilon\, \omega^+ |w^+(\epsilon, |\mathbf{p}|)|^2}$ with $\omega^+ = 1 - 2E_{\mathbf{p}}/M_d$, for $^3S_1^+$ and $^3D_1^+$ components. The dependence of such functions on $|\mathbf{p}|$ is shown in Fig. 1. Then we tentatively introduce D-wave with inclusion of the negative-energy states: $w = \sqrt{w^{+2} + w^{-2} + u^{-2} + \ldots}$ (dotted line labeled as Bethe-Salpeter-II on Fig. 1). This amplitude changes its sign what is the feature of the D-components of a deuteron wave function. Numerically the corresponding curves are in a reasonable agreement in the NR region, up to $|\mathbf{p}| \sim m$.

4. The spin structure functions

To calculate the SF of the deuteron we use the OPE within the effective meson-nucleon theory. The matrix elements of the operators are calculated according to the Mandelstam prescription[8]

$$\langle P|\mathcal{O}^{\mu_1 \ldots \mu_n}(\xi)|P\rangle = \int d^4y d^4y' \bar{\chi}_{\alpha\gamma}(y,Y)\Lambda^{\mu_1 \ldots \mu_n}_{\alpha\beta;\gamma\delta}(\xi; y-y', Y-Y')\chi_{\gamma\beta}(y',Y'), \quad (21)$$

where $\Lambda^{\mu_1 \ldots \mu_n}$ is the Mandelstam vertex corresponding to the given operator. Neglecting "off-mass-shell" corrections[9, 10] and working in the impulse approximation, the Mandelstam vertex of the operators (6) and (7) can be written as follows

$$q^{\mu_1} \ldots q^{\mu_n} \Lambda^0_{\mu_1 \ldots \mu_n}(p,p',P)_{\alpha\beta;\gamma\delta} = \frac{1}{2}(2\pi)^4\delta^{(4)}(p-p')\nu^{n-1} \quad (22)$$
$$\times \left\{ \Gamma^+_{\alpha\gamma}(m-\hat{p}_2)_{\beta\delta} p_1^{+\,n-1} + \Gamma^+_{\beta\delta}(m-\hat{p}_1)_{\alpha\gamma} p_2^{+\,n-1} \right\},$$

where the kinematical variables are defined in the c. m. system and DIS kinematics is used, $p^+ = p_0 + p_z$ and $pq = \nu p^+$. The matrix Γ^+ corresponds to the spinor structure of the operators: $\Gamma^+_V = \gamma^0 + \gamma^3$ and $\Gamma^+_A = \gamma_5(\gamma^0 + \gamma^3)$. Then the reduced matrix elements can be obtained

$$\Theta^{N/D}_{V,n} = \frac{1}{2M_d^n} \int \frac{d^4p}{\imath(2\pi)^4} \quad (23)$$

$$\left\{ Tr\left(\bar{\chi}_{1\mathcal{M}}(p;P)\Gamma^+_V \chi_{1\mathcal{M}}(p;P)(m-\hat{p}_2)\right)\Big|_{\mathcal{M}=0} - Tr(\ldots)\Big|_{\mathcal{M}=1} \right\} p_1^{+\,n-1}, \quad (24)$$

$$\Theta^{N/D}_{A,n} = \frac{1}{2M_d^n} \int \frac{d^4p}{\imath(2\pi)^4} Tr\left(\bar{\chi}_{1\mathcal{M}}(p;P)\Gamma^+_A \chi_{1\mathcal{M}}(p;P)(m-\hat{p}_2)\right)\Big|_{\mathcal{M}=1} p_1^{+\,n-1}, \quad (25)$$

Applying the inverse Mellin transform to Eq. (24) and Eq. (25), spin SF g_1^D and b_2^D, are cast in the convolution form

$$g_1^D(x) = \int_x^1 \frac{d\xi}{\xi} f^{N/D}_{\rm spin}(\xi) g_1^N(x/\xi), \quad b_2^D(x) = \int_x^1 d\xi \Delta f^{N/D}(\xi) F_2^N(x/\xi), \quad (26)$$

where g_1^N and F_2^N is the SF of the isoscalar nucleon. The effective distribution functions $f^{N/D}_{\rm spin}(\xi)$ and $\Delta f^{N/D}(\xi)$ are defined with the moments (24) and (25), respectively.

Their first moments measure the dependence of the momentum distributions on the nucleon and deuteron spin

$$\int_0^1 d\xi f_{\rm spin}^{N/D}(\xi) = \langle P|\bar{N}\gamma_5 N|P\rangle\big|_{\mathcal{M}=1}, \quad \int_0^1 d\xi \Delta f^{N/D}(\xi) = 0. \tag{27}$$

Using the distribution functions $f_{\rm spin}^{N/D}(\xi)$ and $\Delta f^{N/D}(\xi)$ and the realistic parametrization of the nucleon SF g_1[11] and F_2[1], we calculate the SF of the deuteron.

5. Results

The numerical results of the SF are shown in Fig. 2. The nuclear effects in b_2^D and g_1^D calculated within the NR approach of Ref.[12] are presented as well. The magnitude of $b_2^D(x)$ on Fig. 2a is small and about the same in different models, the order of magnitude is of $|\mathbf{p}|^2/m^2$. The source of the effects in b_2^D is the interference of the S- and D-waves in the deuteron wave function. The NR calculation for g_1^D includes

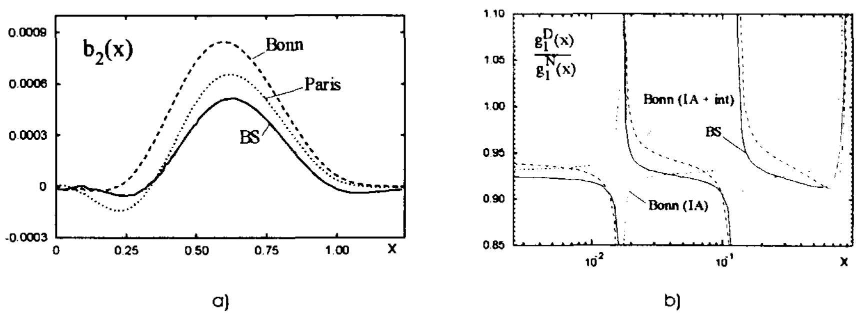

Figure 2: The spin SF of the deuteron: a) $b_2^D(x)$ calculated in the relativistic and NR approaches; b) the ratio of $g_1^D(x)/g_1^N(x)$ calculated in the relativistic and NR approaches.

the effects of Fermi-motion of the polarized nucleons (dotted curve on Fig. 2b), Fermi-motion of the polarized nucleons with the binding effects (dashed curve Fig. 2b). The ratio in the interval $0.2 < x < 0.7$ is governed by the distractive contribution of the D-wave admixture which generates a polarization of the deuteron along the z-axis even though the nucleons have their spin aligned in the opposite to the polarization. The "poles" of the ratio on Fig. 2b are by no means the nuclear effects and are the consequence of nodes of the parametrization of the $g_1^N(x)$[11]. All curves in Fig. 2b tend to the limit $1 - \frac{3}{2}P_D$ as $x \to 0$.

6. Summary

In summary, we have considered a relativistic description of DIS scattering on the deuteron within the BS formalism and the OPE method. The leading twist spin SF g_1^D and b_2^D are calculated in terms of the Mandelstam vertices and BS amplitude which is the solution of the spinor-spinor BS equation with a realistic meson exchange potential. We show the obtained amplitude provides an appropriate description of static properties of the deuteron. It is found that the results for the SF in the relativistic impulse approximation agree with the previous NR calculations motivated by the nuclear physics.

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