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Nonvanishing tensor polarization of sea quarks in polarized deuterons

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Abstract

We show how the dependence of the diffractive nuclear shadowing and of the nuclear excess of pions on the deuteron spin alignment leads to a substantial tensor polarization of sea partons in the deuteron. The corresponding tensor structure function $b_2(x,Q^2)$ rises towards small x and we predict the about one per cent tensor asymmetry $A_2(x,Q^2) = b_2(x,Q^2)/F_{2d}(x,Q^2)$ which by almost two orders in magnitude exceeds the effect evaluated earlier in the impulse approximation. We show that the integral $\int_0^1 dx\, b_1(x,Q^2)$ diverges, which implies that the sum rule $\int_0^1 dx\, b_1(x,Q^2) = 0$ suggested by Close and Kumano does not exist. We comment on the impact of tensor polarization on the determination of the vector spin structure function $g_{1d}(x,Q^2)$ for the deuteron. © 1997 Elsevier Science B.V.

The theoretical prejudice supported to a certain extent by experiment is that at high energies which in deep inelastic scattering (DIS) is $x \to 0$ total cross sections cease to depend on the beam and target polarizations (hereafter x, Q^2 are the DIS variables). For instance, the SLAC E143 data [1] suggest, and QCD evolution analyses of polarized DIS support [2], the vanishing vector polarization of sea quarks in the deuteron. In this paper we show that the tensor polarization of sea quarks in the deuteron is substantial and even rises at $x \to 0$. [1]

If $\mathbf{S}$ is a spin-1 operator then $T_{ik} = \frac{1}{2}(S_iS_k + S_kS_i - \frac{2}{3}\mathbf{S}^2\delta_{ik})$ is the T-even and P-even observable. Consequently, for DIS of *unpolarized* leptons the photoabsorption cross sections $\sigma_T^{(\lambda)}, \sigma_L^{(\lambda)}$ and structure functions $F_2^{(\lambda)}(x,Q^2) = \frac{Q^2}{4\pi^2\alpha_{em}}(\sigma_T^{(\lambda)} + \sigma_L^{(\lambda)})$ can depend on the spin projection $\mathbf{Sn} = \lambda$ onto the target spin quantization axis $\mathbf{n}$: $\sigma^{(\lambda)} = \sigma_0 + \sigma_2 T_{ik}n_in_k + \ldots$ [4–6]. Hereafter we focus on the γ^*-target collision axis (the z-axis) chosen for $\mathbf{n}$. The tensor spin structure function b_2 can be defined as

$$\begin{aligned} b_2(x,Q^2) &= \tfrac{1}{2}[F_2^{(+)}(x,Q^2) + F_2^{(-)}(x,Q^2) - 2F_2^{(0)}(x,Q^2)] \\ &= \tfrac{1}{2}\sum_f e_f^2 x[q_f^{(+)}(x,Q^2) + q_f^{(-)}(x,Q^2) \\ &\quad - 2q_f^{(0)}(x,Q^2)]\,. \end{aligned} \tag{1}$$

By parity conservation $\sigma^{(+)} = \sigma^{(-)}$ and $F_2^{(+)} = F_2^{(-)}$,

[1] A brief report on these results was presented at the Workshop on Future Physics at HERA [3].

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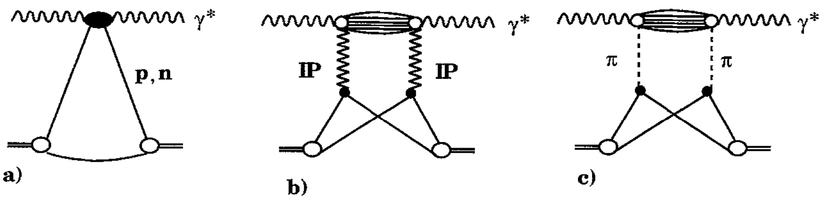

Fig. 1. Impulse approximation (a) and Gribov's inelastic shadowing (b), (c) diagrams for DIS on the deuteron.

$q_f^{(+)} = q_f^{(-)}$. Similarly, $\sigma_T^{(\lambda)}$ and $\sigma_L^{(\lambda)}$ define the transverse, $b_T = 2xb_1$, and the longitudinal, $b_L = b_2 - b_T$, tensor structure functions. The precursor of the structure function b_1 has been discussed by Pais in 1967 [7].

Because vector polarized deuterons are always tensor polarized, in DIS of polarized leptons on a 100 % polarized deuteron the vector spin asymmetry equals $A_{1d} = g_{1d}/F_{1d}^{(\pm)}$. Consequently, the determination of the vector spin structure function g_{1d} only is possible if $F_1^{\pm}$ were known:

$$g_{1d}(x,Q^2) = A_{1d}F_1^{(\pm)}(x,Q^2)\,. \tag{2}$$

Hitherto $F_1^{(\pm)}, F_1^{(0)}$ are completely unknown and all the deuteron data have been analyzed under the unwarranted assumption that $F_{1d}(x,Q^2)$ for the unpolarized deuteron can be used instead of $F_1^{(\pm)}(x,Q^2)$ in (2). In order to justify this assumption and use the polarized deuteron data for tests of the Bjorken sum rule one needs the direct measurements, and the benchmark theoretical evaluations, of the tensor structure functions of the deuteron.

In the impulse approximation (IA) of Fig. 1a discussed ever since [5]

$$F_2^{(\lambda)}(x,Q^2) = \int \frac{dy}{y} f^{(\lambda)}(y) \times \left[F_{2p}\left(\frac{x}{y},Q^2\right) + F_{2n}\left(\frac{x}{y},Q^2\right)\right] \tag{3}$$

subject to the sum rule for the Doppler-Fermi smearing function

$$\int \frac{dy}{y} f^{(\lambda)}(y) = 1\,, \tag{4}$$

where y is a fraction of the lightcone momentum of the deuteron carried by the nucleon. For the pure S-wave deuteron $f^{(\pm)}(y) = f^{(0)}(y)$ and $b_2(x,Q^2) = 0$, the D-wave admixture makes $f^{(\pm)}(y) \neq f^{(0)}(y)$ but they differ markedly only when $f^{(\pm}(y), f^{(0)}(y) \ll 1$ [8]. As a result, in the IA the tensor polarization of quarks $A_2(x,Q^2) = b_2(x,Q^2)/F_{2d}(x,Q^2)$ is at most several units of 10^{-4} for $x \lesssim 0.5$ [5,9,10]. Because of the sum rule (4), in the IA $A_2(x,Q^2) \to 0$ at $x \to 0$. Furthermore, the widely discussed sum rule

$$\int dx b_1(x,Q^2) = 0 \tag{5}$$

holds in the IA. Close and Kumano conjectured that if the sea is unpolarized then this sum rule holds beyond the IA ([11], for earlier discussions see also [12,13]). For the positivity constraints on $b_{1,2}(x,Q^2)$ see [14], higher twist effects are discussed in [15].

Here we report the beyond-the-IA evaluation of the tensor polarization of sea in the deuteron from Gribov's inelastic shadowing diagrams of Fig. 1b and 1c [16] which lead to a nuclear eclipse effect: $\sigma_{\rm tot}(\gamma^* d) = \sigma_{\rm tot}(\gamma^* p) + \sigma_{\rm tot}(\gamma^* n) - \Delta\sigma_{\rm sh}(\gamma^* d)$. In diffractive nuclear shadowing (NSH) the intermediate state X is excited by pomeron exchange (Fig. 1b), in the second mechanism the state X is excited by pion exchange (Fig. 1c). The both contributions to the eclipse effect are sensitive to the alignment of nucleons in the polarized deuteron. Besides the sea they make gluons too tensor polarized (for an early discussion on gluon densities for spin axis **n** normal to the collision axis see [6,17,18]). Ever since [16] the both contributions to the eclipse effect in the deuteron have been discussed repeatedly [19-24], but the dependence on the tensor polarization has not been calculated before. The only exception is Ref. [25], but there the dominant effect of the D-wave has not been considered and the result for $b_1(x)$ evaluated in the reggeon exchange model vanishes at $x \to 0$, whereas we find $b_1(x)$ which rises faster than $\frac{1}{x}$, what makes the sum rule integral (5) divergent rather than receiving a finite and small correction, as was discussed in [25].

Hereafter we focus on small x and do not consider negligible IA terms. We evaluate diffractive NSH using Gribov's theory [16] extended to DIS in [27]:

$$\Delta F_{\rm sh}^{(\lambda)}(x,Q^2) = \frac{2}{\pi}\frac{Q^2}{4\pi^2\alpha_{\rm em}}\int d^2\mathbf{k}_\perp \times \int dM^2 S_D^{(\lambda)}(4\mathbf{k}^2)\left.\frac{d\sigma^D(\gamma^*\to X)}{dtdM^2}\right|_{t=-\mathbf{k}_\perp^2}. \quad (6)$$

Here M is the mass of the state X, $\mathbf{k}=(\mathbf{k}_\perp,k_z)$ is the momentum of the pomeron,

$$\begin{aligned} S_D^{(\lambda)}(4\mathbf{k}^2) &= \int\frac{d^3\mathbf{p}}{(2\pi)^3}\mathrm{Tr}(\Psi_\lambda^\dagger(\mathbf{p}+\mathbf{k})\Psi_\lambda(\mathbf{p}))\\ &= 4\pi\int\frac{d^3\mathbf{p}}{(2\pi)^3}\Big\{u_0(p')u_0(p)(\mathbf{d}^*\mathbf{d})\\ &+\frac{u_0(p')w_2(p)}{\sqrt{2}}\left[3(\mathbf{d}^*\hat{\mathbf{p}})(\mathbf{d}\hat{\mathbf{p}})-(\mathbf{d}^*\mathbf{d})\right]\\ &+\frac{u_0(p)w_2(p')}{\sqrt{2}}\left[3(\mathbf{d}^*\hat{\mathbf{p}}')(\mathbf{d}\hat{\mathbf{p}}')-(\mathbf{d}^*\mathbf{d})\right]\\ &+\frac{w_2(p)w_2(p')}{2}\left[9(\mathbf{d}^*\hat{\mathbf{p}}')(\mathbf{d}\hat{\mathbf{p}})(\hat{\mathbf{p}}\hat{\mathbf{p}}')\right.\\ &\left.-3(\mathbf{d}^*\hat{\mathbf{p}})(\mathbf{d}\hat{\mathbf{p}})-3(\mathbf{d}^*\hat{\mathbf{p}}')(\mathbf{d}\hat{\mathbf{p}}')+(\mathbf{d}^*\mathbf{d})\right]\Big\}, \end{aligned} \quad (7)$$

$\Psi_\lambda(\mathbf{p}) = \sqrt{\pi}(\sqrt{2}u_0(p)(\boldsymbol{\sigma}\mathbf{d}) + w_2(p)(3(\boldsymbol{\sigma}\hat{\mathbf{p}})(\mathbf{d}\hat{\mathbf{p}}) -(\boldsymbol{\sigma}\mathbf{d})))$ [26], $u_0(p), w_2(p)$ are the S and D-wave functions normalized by the condition $S_D^{(\lambda)}(0)=1$, $\boldsymbol{\sigma}$ are the Pauli spin matrices, $\hat{\mathbf{p}}=\mathbf{p}/|\mathbf{p}|$ and $\mathbf{p}$ is the proton's momentum in the deuteron, $\mathbf{d}(\lambda)$ is the deuteron polarization vector. The polarization dependence of the form factor $S_D^{(\lambda)}(4\mathbf{k}^2)$ and the nonvanishing $b_2(x,Q^2)$ come only from the S-D-wave interference and D-wave terms $\propto u_0(p)w_2(p') + u_0(p')w_2(p)$ and $\propto w_2(p)w_2(p')$.

Decompose the diffractive DIS cross section into the pomeron valence and sea structure functions $F_{2\mathbf{P}}^i(\beta)$ (i = val, sea) and the pomeron fluxes in the proton $\phi_\mathbf{P}^i(x_\mathbf{P})/x_\mathbf{P}$ [29,30],

$$(M^2+Q^2)\frac{d\sigma_T^D(\gamma^*\to X)}{dtdM^2} = \frac{\sigma_{\rm tot}(pp)}{16\pi}\frac{4\pi^2\alpha_{\rm em}}{Q^2} \times\sum_{i={\rm val,sea}}\phi_\mathbf{P}^i(x_\mathbf{P})F_{2\mathbf{P}}^i(\beta,Q^2)e^{-B_i|t|}, \quad (8)$$

where $\beta = Q^2/(Q^2+M^2)$ is the Bjorken variable for DIS on pomerons and $x_\mathbf{P} = x/\beta$ is a fraction of the nucleon's momentum taken away by the pomeron, in the target rest frame $k_z = x_\mathbf{P}m_N$. Then diffractive NSH can be cast in the convolution form

$$\Delta F_{\rm sh}^{(\lambda)}(x,Q^2) = \int_x^1\frac{dx_\mathbf{P}}{x_\mathbf{P}}\left(\Delta n_{\rm val}^{(\lambda)}(x_\mathbf{P})F_{2\mathbf{P}}^{\rm val}(\frac{x}{x_\mathbf{P}},Q^2) + \Delta n_{\rm sea}^{(\lambda)}(x_\mathbf{P})F_{2\mathbf{P}}^{\rm sea}(\frac{x}{x_\mathbf{P}},Q^2)\right), \quad (9)$$

where $\Delta n_i^{(\lambda)}(x_\mathbf{P})$ is a nuclear modification of the pomeron flux functions (i = val, sea):

$$\Delta n_i^{(\lambda)}(x_\mathbf{P}) = \frac{2}{\pi}\frac{\sigma_{\rm tot}(pp)}{16\pi}\phi_\mathbf{P}^i(x_\mathbf{P}) \times\int d^2\mathbf{k}_\perp S_D^{(\lambda)}(4\mathbf{k}^2)e^{-B_i|t|}. \quad (10)$$

Because NSH is a leading twist effect [27,28] we find a manifestly leading twist $b_2(x,Q^2)$. In the practical calculations we use the predictions [29] for $F_{2\mathbf{P}}^i$ and $\phi_\mathbf{P}^i(x_\mathbf{P})$ from the color dipole approach to diffractive DIS [28,31] which is known to agree well with the H1 and ZEUS data [32,33]. Following [28] we take for diffraction slopes $B_{\rm sea} = B_{3\mathbf{P}} = 6$ GeV^{-2} and $B_{\rm val} = 2B_{\rm sea}$, the results for NSH only weakly depend on $B_{\rm val}, B_{\rm sea}$.

The pion exchange dominates $\gamma^*p\to nX$, contributes substantially to $\gamma^*p\to pX$ at a moderately small $x_\mathbf{P}$ [34,35] and is a well established source of the $\bar{u}$-$\bar{d}$ asymmetry in DIS on the nucleon [36–38]. The relevant input is $d\sigma(\gamma^*p\to Xn)/dtdM^2$ times the isospin factors. The diagram of Fig. 1c can be interpreted as DIS on nuclear excess pions in the deuteron and its contribution to $\Delta F_{\rm sh}$ can be cast in the convolution form

$$\Delta F_{{\rm sh},\pi}^{(\lambda)}(x,Q^2) = -\int_x^1\frac{dz}{z}\Delta n_\pi^{(\lambda)}(z) \times\left[\frac{2}{3}F_2^{\pi^\pm}(\frac{x}{z},Q^2)+\frac{1}{3}F_2^{\pi^0}(\frac{x}{z},Q^2)\right], \quad (11)$$

where F_2^π is the pion structure function and the flux of excess pions equals

$$
\begin{aligned}
\frac{\Delta n_\pi^{(\lambda)}(z)}{z} &= -6\frac{g_{\pi NN}^2}{4\pi} z \int \frac{d^2\mathbf{k}_\perp}{(2\pi)^2}\frac{G_{\pi NN}^2(z,t)}{(t-\mu_\pi^2)^2} \\
&\times \int \frac{d^3\mathbf{p}}{(2\pi)^3} 4\pi \Big\{ u_0(p)u_0(p')\big(2(\mathbf{kd})(\mathbf{kd}^*) - \mathbf{k}^2\big) \\
&+ \frac{u_0(p)w_2(p')}{\sqrt{2}}\big[3(\mathbf{d}^*\hat{\mathbf{p}}')\big(2(\hat{\mathbf{p}}'\mathbf{k})(\mathbf{kd}) - \mathbf{k}^2(\hat{\mathbf{p}}'\mathbf{d})\big) \\
&- \big(2(\mathbf{kd})(\mathbf{kd}^*) - \mathbf{k}^2\big)\big] \\
&+ \frac{u_0(p')w_2(p)}{\sqrt{2}}\big[3(\mathbf{d}^*\hat{\mathbf{p}})\big(2(\hat{\mathbf{p}}\mathbf{k})(\mathbf{kd}) - \mathbf{k}^2(\hat{\mathbf{p}}\mathbf{d})\big) \\
&- \big(2(\mathbf{kd})(\mathbf{kd}^*) - \mathbf{k}^2\big)\big] \\
&+ \frac{w_2(p)w_2(p')}{2} \\
&\times \big[9(\mathbf{d}\hat{\mathbf{p}})(\mathbf{d}^*\hat{\mathbf{p}}')\big(2(\mathbf{k}\hat{\mathbf{p}})(\mathbf{k}\hat{\mathbf{p}}') - \mathbf{k}^2(\hat{\mathbf{p}}\hat{\mathbf{p}}')\big) \\
&- 3(\mathbf{d}^*\hat{\mathbf{p}}')\big(2(\hat{\mathbf{p}}'\mathbf{k})(\mathbf{kd}) - \mathbf{k}^2(\hat{\mathbf{p}}'\mathbf{d})\big) \\
&- 3(\mathbf{d}^*\hat{\mathbf{p}})\big(2(\hat{\mathbf{p}}\mathbf{k})(\mathbf{kd}) - \mathbf{k}^2(\hat{\mathbf{p}}\mathbf{d})\big) \\
&+ 2(\mathbf{kd})(\mathbf{kd}^*) - \mathbf{k}^2\big]\Big\}, \qquad (12)
\end{aligned}
$$

$\mathbf{k} = (\mathbf{k}_\perp, zm_N)$ is the pion momentum (notice the similarity between z and $x_\mathbf{P}$), $t = -(\mathbf{k}_\perp^2 + z^2m_N^2)/(1-z)$, $g_{\pi NN}$ is the πNN coupling and $G_{\pi NN}(z,t)$ is the πNN form factor. At $z \ll 1$ relevant to the present analysis the Regge model parameterization $G_{\pi NN}(z,t) = \exp(-(\Lambda_0^2 + \alpha'_\pi \log(\frac{1}{z}))(|t| + \mu_\pi^2))$ is appropriate (for instance, see [39]). This form factor is constrained by the pion induced $\bar{u}$-$\bar{d}$ asymmetry in the proton, with $\Lambda_0^2 = 0.3$ GeV^{-2}, the Regge slope $\alpha'_\pi = 0.9$ GeV^{-2} and the F_2^π from [40] we find the same results as in the phenomenologically successful lightcone approach ([36,37], for the review see [38]).

Brief comments on $b_L(x,Q^2)$ are in order. It does not vanish. Evidently, b_L from excess pions is found replacing F_2^π in (11) by $F_L^\pi(x,Q^2)$. The experimental data on $d\sigma_L^D(\gamma^* \to X)/dtdM^2$ are as yet lacking, we evaluate $b_L(x,Q^2)$ from NSH using the perturbative QCD approach [41] which gives $R = b_L/b_T$ similar to $R = \sigma_L/\sigma_T$ for DIS on protons.

We use the Bonn (Table 11 in Ref. [42]) and Paris [43] wave functions as representatives of the weak and strong D-wave in the deuteron. The deuteron wave functions in (7) and (12) make $\Delta n_\pi^{(\lambda)}(z)$, $\Delta n_{\rm val}^{(\lambda)}(x_\mathbf{P})$ and $\Delta n_{\rm sea}^{(\lambda)}(x_\mathbf{P})$ nonvanishing only at $x_\mathbf{P}, z \lesssim \langle z\rangle \sim (R_d m_N)^{-1} \sim 0.2$, where R_d is the deuteron radius.

Table 1
The S-D wave decomposition of the mean multiplicities of excess pions in the deuteron for different polarizations λ along the $\gamma^* d$-collision axis. The numerical results were obtained for the Bonn and Paris wave functions

	$\langle \Delta n_\pi^{(\lambda)}\rangle \cdot 10^2$			
	Bonn		Paris	
	$\lambda = 0$	$\lambda = \pm 1$	$\lambda = 0$	$\lambda = \pm 1$
S-wave	0.146	0.999	0.133	0.929
S-D interference	−0.537	−0.268	−0.676	−0.331
D-wave	0.035	0.085	0.070	0.153
total	−0.356	0.807	−0.466	0.750

The nonrelativistic treatment of the deuteron holds, and binding effects can be neglected, at such a small $x_\mathbf{P}, z$ ([22] and references therein). The arguments for using the on-mass shell pion structure function are reviewed in [38]. After the $d\sigma(\gamma^* p \to nX)$ will be measured at HERA [34], the pion effect will be calculated free of any uncertainties with the off-mass shell behaviour [44]. To have a quick impression on the effect of excess pions notice that at $x \lesssim \langle z\rangle$ the convolution (11) can be approximated by

$$
\Delta F_{{\rm sh},\pi}^{(\lambda)}(x,Q^2) = -\left(\frac{2}{3}F_2^{\pi^\pm}\left(\frac{x}{\langle z\rangle},Q^2\right) + \frac{1}{3}F_2^{\pi^0}\left(\frac{x}{\langle z\rangle},Q^2\right)\right)\langle \Delta n_\pi^{(\lambda)}\rangle . \qquad (13)
$$

The numbers of excess pions $\langle \Delta n_\pi^{(\lambda)}\rangle = \int_0^1 \frac{dz}{z} \times \Delta n_\pi^{(\lambda)}(z)$ are cited in Table 1. The pure S- and (negligibly small) D-wave terms give $\langle \Delta n^{(\lambda)}\rangle > 0$, i.e., the antishadowing effect in $\Delta F_{\rm sh}$. The tensor polarization of sea quarks is proportional to $\langle \Delta n_\pi^{(\pm)}\rangle - \langle \Delta n_\pi^{(0)}\rangle = 1.16\cdot 10^{-2}$, whereas NSH for unpolarized deuterons is proportional to $\frac{1}{3}(\langle 2\Delta n_\pi^{(\pm)}\rangle + \langle \Delta n_\pi^{(0)}\rangle) = 0.42 \cdot 10^{-2}$ which is strongly reduced by the S-D wave interference from the pure S-wave contribution. For a comparison, for the number of pions in a nucleon one has $\langle n_\pi\rangle_N \sim 0.2$ [36,37]. Table 1 shows the results are stable against the admissible variations of the deuteron wave function. The same is true of the diffractive NSH effects.

In Fig. 2 and Fig. 3 we show the contribution from excess pions to $\Delta F_{\rm sh}(x,Q^2)$ and $b_2(x,Q^2)$ calculated from Eqs. (11), (12) for $Q^2 = 10$ GeV^2 and for the

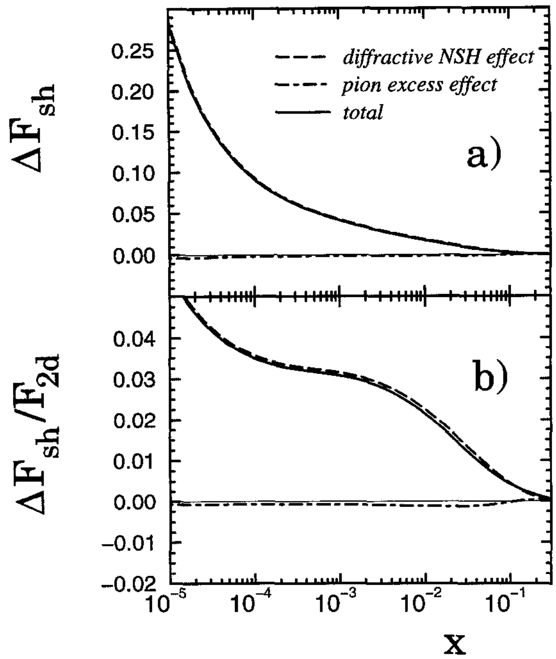

Fig. 2. (a) Nuclear shadowing correction in the deuteron $\Delta F_{sh}(x,Q^2)$ for $Q^2 = 10$ GeV^2 and its diffractive NSH and pion excess components for the Bonn wave function of the deuteron. (b) The fractional shadowing correction to the deuteron structure function $F_{2d}(x,Q^2)$ for $Q^2 = 10$ GeV^2.

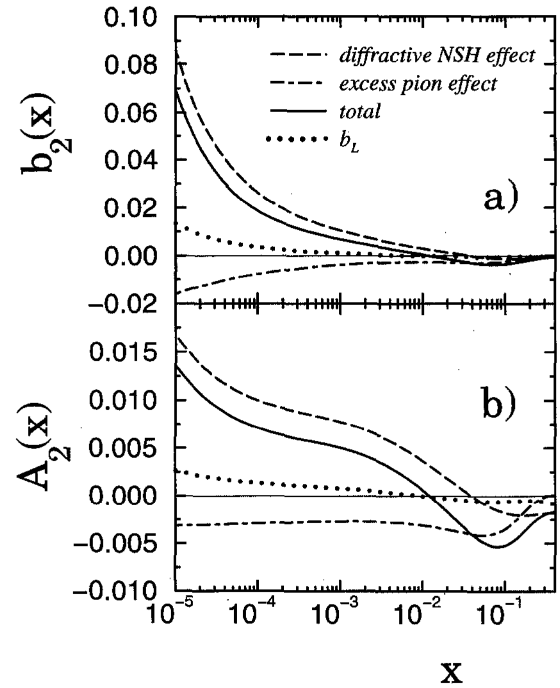

Fig. 3. (a) The tensor structure function of the deuteron $b_2(x,Q^2)$ at $Q^2 = 10$ GeV^2 and its diffractive NSH and pion excess components for the Bonn wave function of the deuteron. (b) The tensor polarization $A_2(x,Q^2)$ of quarks and antiquarks in the deuteron at $Q^2 = 10$ GeV^2. Also shown (by the dotted line) is the longitudinal tensor structure function $b_L(x,Q^2)$ (in (a)), and the corresponding asymmetry (in (b)).

GRV pion structure function [40] (for similar parameterizations see also [45]). Because the pion effect is small, our estimate of NSH for unpolarized deuterons in Fig. 2 is close to that of Ref. [24] in which only the pure diffractive NSH has been evaluated in the related model for diffractive DIS. A good description of NSH in heavy nuclei in the same approach [46] lends support to our numerical estimates. The results for $b_2(x,Q^2)$ for the deuteron are shown in Fig. 3. The excess pions play much larger rôle in $b_2(x,Q^2)$ than in nuclear shadowing and dominate $b_2(x,Q^2)$ at $x \gtrsim 0.01$. At smaller x the diffractive NSH mechanism takes over. In Fig. 2 and Fig. 3 we also show our estimate for $b_L(x,Q^2)$. Fig. 3b shows that the tensor polarization of sea quarks in the deuteron $A_2(x,Q^2)$ is substantial and rises towards small x.

Summary and conclusions. Regarding the numerical results, the uncertainties with the wave function of the deuteron are marginal. There are larger uncertainties with the pion and pomeron structure functions, which can be reduced with more data from HERA and do not call into question our finding of the striking spin dependence of DIS which persists, and even gets stronger, at small x. We regard our analysis as a benchmark calculation of the tensor structure function of the deuteron at small x, very similar results must follow from other models if they give similar description of the $\bar{u}$-$\bar{d}$ asymmetry in the proton and of diffractive DIS.

Regarding the measurements of g_{1d} in polarized $\vec{e}\vec{d}, \vec{\mu}\vec{d}$ scattering, our results justify the substitution of the unpolarized deuteron structure function for the more correct but still unknown $F_2^{(\pm)}$ in (2) at the present accuracy in the spin asymmetry A_{1d}, but corrections for the tensor polarization will eventually become important for the interpretation of high accuracy DIS off polarized deuterons in terms of the neutron spin structure function g_{1n} and for the precision tests of the Bjorken sum rule.

The both diffractive nuclear shadowing and pion excess mechanisms give a leading twist $b_2(x,Q^2)$. The pion contribution is obviously GLDAP evolving. The sea structure function of the pomeron also obeys the GLDAP evolution [31]. The detailed form of scaling violations in the valence structure function of the pomeron at $\beta \to 1$ remains an open issue [30,47], but the contribution from this region of β to $b_2(x,Q^2)$ is marginal. Consequently, QCD evolution of $b_2(x,Q^2)$ must not be any different from that of $F_2(x,Q^2)$, see also a related discussion in [6,18]. Furthermore, because of the similar origin of nuclear shadowing and of tensor asymmetry, we can invoke the theoretically [27,46] and experimentally (for the review see [48]) well established weak Q^2 dependence of nuclear shadowing. For instance, we expect a similarly large tensor asymmetry in $\sigma_{\rm tot}(\gamma D)$ for real photons ($Q^2=0$).

Our finding of the rising tensor polarization of sea quarks in the deuteron invalidates the assumptions invoked in the derivation of the Close-Kumano sum rule (5). Because we find $b_1(x) = \frac{1}{2x}b_2(x)$ which rises faster than $\frac{1}{x}$, the Close-Kumano integral simply diverges at small x.

Finally, Eqs. (9) and (12) are fully applicable to calculation of the tensor polarization of gluon densities in the deuteron polarized along the $\gamma^* d$ collision axis in terms of gluon structure functions of the pomeron and pion. The results for the gluon structure functions will be presented elsewhere.

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Note added. After the present paper has been submitted for publication, G. Piller informed us of a related analysis of the diffractive shadowing contribution to $b_1(x)$ [49].

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