The final state hadrons in polarised deep inelastic scattering

S. D. Bass

Cavendish Laboratory, University of Cambridge, Madingley Road, Cambridge CB3 0HE, U.K.

and

Institut für Kernphysik, KFA-Jülich, D-52425 Jülich, Germany

Abstract

General arguments suggest that the non-perturbative background field in QCD may have a non-trivial spin structure. We discuss how this effect may be manifest in semi-inclusive measurements of fast pions in polarised deep inelastic scattering.

1 Introduction

The discovery of the EMC spin effect [1] has inspired a vigorous programme to understand the internal spin structure of the nucleon – for reviews see [2,3]. In the absence of a solution to non-perturbative QCD we construct QCD inspired models of the nucleon where the valence quarks move in some background field that is defined self consistently as the deformation of the non-perturbative vacuum that is induced by the valence quarks. The background field includes the effects of confinement and the dynamical breaking of chiral symmetry. It is modelled by the bag in relativistic quark models of nucleon structure. One of the key questions in QCD spin physics is does this background field possess a non-trivial spin structure? This problem has been addressed by a number of authors [4-9]. (It is important to distinguish polarised gluons inside the bag and a possible spin structure of the long-range gluon induced fields which are described by the bag itself.)

Nachtmann and collaborators [4] have proposed that the quark sea may develop a nett transverse polarisation as the quarks move through the non-perturbative colour-magnetic fields in the QCD vacuum similar to the build up of electron polarisation in storage rings. This effect

Present address
could lead to a breakdown of factorisation in polarised Drell-Yan experiments due to initial state correlations where the anti-quark of one hadron moves through the non-perturbative vacuum of the second. This effect is consistent with the NA-10 data [10] and offers a possible explanation of the K-factor in Drell-Yan experiments [4].

The question of background field effects in polarised deep inelastic scattering has been addressed by Jaffe and Manohar [5,9], Fritzsch [6] and the present author [7,8]. Jaffe and Manohar [9] showed that it is necessary to consider non-perturbative gluon fields in order to fully understand the role of the anomaly in the first moment of $g_1$. This approach was also taken by Fritzsch [6] who proposed a model where the anomaly induces two gluonic contributions to $g_1$: the partonic contribution from the polarised gluon distribution that was initially found by Efremov and Teryaev and by Altarelli and Ross [11], and also a non-perturbative contribution which is associated with the internal structure of the constituent quark. Bass [7,8] generalised the arguments of Jaffe and Manohar to the higher moments of $g_1$ and suggested that the anomaly has the potential to contribute an OZI violation to $g_1$ at large $x$ ($x$ greater than 0.2) in the form of a local interaction between the hard photon and the background field.

2 Three reasons to look for a large $x$ anomaly contribution to $g_1$

Before discussing the phenomenology of a possible large $x$ anomaly in semi-inclusive reactions it is helpful to keep in mind the reasons why one might expect to see such an effect which does not appear in perturbation theory. We briefly discuss three reasons why one might look for a large $x$ anomaly in $g_1$. In this section we collect together ideas that are scattered throughout the literature and discuss the relationship between a possible large $x$ anomaly and quark model calculations of $g_1$. We start by defining the spin dependent quark distribution $\Delta q(x, Q^2)$ via the light-cone expansion so that:

$$2M s_+(p_+)^{2n} \int_0^1 dx x^{2n} \Delta q_k(x, Q^2) = \langle p, s | [\bar{q}(0) \gamma^+ \gamma_5 (i D_+)^{2n} \frac{\lambda^k}{2} q(0)]^{\text{GI}} | p, s > c,$$

where $D_\mu = \partial_\mu + igA_\mu$ is the QCD-gauge covariant derivative and the superscript $\text{GI}$ emphasises that we are using gauge invariant operators.

1. Perhaps the most illuminating derivation of the anomaly is due to Schwinger [12] and uses point splitting regularisation (see also [13]). One finds that the chirality of a quark propagating in some background gauge field is not conserved with the result that the background field screens the spin of the quark.

We can write $\Delta q(x, Q^2)$ in terms of a light-cone correlation function in $A_+ = 0$ gauge as the Fourier transform of the axial vector current point-split along the light-cone:

$$\Delta q(x, Q^2) = \frac{1}{2\pi} \int dz_- \cos(xz_-p_+) < p, s | [\bar{q}(0) \gamma^+ \gamma_5 q(0)]_{Q^2} | p, s > c$$

The Fourier transform in equ.(2) means that small $z_-$ effects have the potential to contribute over a complete range of $x$ and that large $z_-$ effects contribute only at small $x$.

Light-cone correlation functions need to be treated with care in an interacting theory like QCD because of ultra-violet divergences [14]. (Each moment of equ.(2) has to be renormalised separately.) Quark and lattice models are endowed with an implicit or
explicit ultra-violet cut-off so that the correlation functions are formally well defined within such models. This result has led to suggestions that one can use the correlation functions to calculate the leading twist part of the structure function within one’s favourite model at some low scale [15,16]. (The model distributions are then evolved to higher $Q^2$, convoluted with the QCD radiative coefficients and compared with data.) There is one important practical difference between the correlation function for $\Delta q(x, Q^2)$ in interacting QCD and in these models: the treatment of the $z_- \rightarrow 0$ limit of the point-split matrix element in equ.(2). This limit is where one finds the axial anomaly in Schwinger’s derivation of the anomaly in the axial-vector current [12] – see also [8]. The way that we treat the $z_- \rightarrow 0$ limit determines the symmetry which is inherent in the polarised quark distribution. At this point one needs to consider QCD directly. The ultra-violet cut-off in the models means that the model correlation functions do not include this zero light-cone correlation length effect which, via the Fourier transform, has the potential to be important at all $x$.

2. In classical field theory the axial vector current is conserved and defines a good spin operator. In renormalised QCD one finds that the conserved axial vector current differs from the physical gauge-invariant operator by a gauge dependent counterterm, which is commonly denoted $k_\mu$. In general gauges, eg. the covariant Feynman gauge, one finds that the forward matrix element of this $k_\mu$ are invariant under “small” gauge transformations of perturbation theory but are not invariant under “large” gauge transformations with non-zero winding number [9]. The parton model is nearly always formulated in the light-cone gauge $A_+ = 0$. In this gauge one finds that the forward matrix elements of the $k_\mu$ are invariant under both small and large gauge transformations and also that $k_\mu$ corresponds with the gluon spin operator (evaluated in this gauge). For this reason, $k_\mu$ is commonly associated with the polarised gluon distribution $\Delta g(x, Q^2)$ in the parton model. ($\Delta g(x, Q^2)$ is a small $x$ effect.) When one considers the effect of the anomaly in the whole distribution rather than just the first moment, one finds a gauge dependent gluonic counterterm for each of the $C = +1$ axial tensor operators in equ.(1). These higher moment anomalous contributions are similarly associated with $\Delta g(x, Q^2)$ in the parton model. The interesting question is then: what is the physics content of the remaining quark distribution? Naively, it is tempting to say that since this distribution corresponds to a conserved axial vector current it should therefore correspond to the same canonical physics that we expect from semi-classical quark models. However, this is not clear in view of the general problem of invariance under large gauge transformations. In a covariant gauge one can continue to make large gauge transformations, going around the circle enough times, so that the anomalous “$k_\mu$ contribution” to equ.(1) becomes arbitrarily big at any given $x$. Given our present understanding of non-perturbative QCD, there is no reason to assume that the nett anomalous contribution to the physical $\Delta q(x, Q^2)$ is zero at large $x$ (say at $x$ greater than 0.2). If there is a nett anomalous contribution to $g_1$ at large $x$ it would be manifest as an OZI violation and would be wrapped up in what we call a constituent quark in the light-cone gauge. In this parton model gauge we write for the first moment of $\Delta q(x, Q^2)$

$$\Delta q_{GI}(Q^2) = (\Delta q_S + \Gamma) - \left(\frac{\alpha_s}{2\pi}\Delta g\right)(Q^2) \tag{3}$$

Here $-\frac{\alpha_s}{2\pi}\Delta q$ is the partonic polarised gluon contribution [11] and $\Gamma$ is a background field contribution which includes the remaining OZI violation. The $\Delta q_S$ are the quark “spin fractions” with good OZI and correspond to the quantities that one would calculate in
a semi-classical quark model with no anomaly. The background field contribution is not present in perturbation theory. It is present in the spin dependent quark distribution, which is a non-perturbative quantity, so that the sum \((\Delta q_S + \Gamma)\) is the polarised “quark contribution” to the first moment of \(g_1\) in the parton model. Indeed, the background field contribution has the same \(Q^2\) evolution equation and the same local coupling to the hard photon as we expect of a “quark” [7]. Equ.(3) is implicit in the model of Fritzsch [6].

3. It is well known that the current \(k_\mu\) couples to a massless (Kogut-Susskind) ghost pole [17], which is exactly cancelled by the massless pole that couples to the conserved (gauge dependent) axial-vector current. Let us temporarily consider the effect of \(k_\mu\) in isolation to the physics of the gauge invariant current – that is, in the spirit of the parton model neglecting a possible finite non-perturbative background field contribution. In this case, it is natural to ask what is the contribution of the (unphysical) ghost to the shape of \(g_1\). Here one employs the Sullivan mechanism [18] to calculate mesonic contributions to structure functions. The zero mass ghost would give a contribution to \(g_1\) at much larger \(x\) than the light mass pion. Of course, this contribution is unphysical and is cancelled by the ghost coupling to the conserved current. However, it does illustrate the point that one cannot self-consistently treat the physics of “the anomaly” (as described by \(k_\mu\)) in isolation to the conserved current as a small \(x\) effect.

The correct way to describe the \(x\) dependence of the anomaly in \(g_1\) would be an all moment generalisation of the analysis of Shore and Veneziano [19], where the OZI violations are isolated explicitly. However, it is not clear (at least to the author) how one would estimate the size of the effect. At the present time one can only say that non-perturbative QCD appears to permit an anomalous contribution to \(g_1\) at large \(x\). One should therefore try to observe this effect in the laboratory. This may be possible in semi-inclusive measurements of fast pions in polarised deep inelastic scattering – experiments which are being carried out as part of the HERMES [20] and SMC [21] programme. Before we discuss how the anomaly may show up in these experiments, it is worth spending time to discuss whether this effect consistent with – or indeed is ruled out by – existing data. We consider only the inclusive spin asymmetries. Measurements of semi-inclusive asymmetries are just beginning [21] and the errors are presently much too large to make definite conclusions.

It is sometimes emphasised that valence quark models predicted the inclusive longitudinal spin asymmetry (and hence \(g_1\)) at \(x \geq 0.2\) within the present errors [22-24] – see also [25]. Is a large \(x\) anomaly consistent with these calculations? The answer to this question is yes. In the picture we are presenting, the anomaly emerges as a non-trivial spin structure in the transition from the constituent to current quark. The valence quark model calculations of \(g_1\) employ either a Melosh transformation or phenomenological “spin dilution factors” to transform from constituent quark to current quarks in the transition from constituent quark to current quark degrees of freedom. Suppose that we compare \(\Delta q(x, Q^2)\) and the C-odd, anomaly free, polarised valence distribution

\[
\Delta q_V(x, Q^2) = (q - \bar{q})_{GI}^\uparrow - (q - \bar{q})_{GI}^\downarrow = (q - \bar{q})_S^\uparrow - (q - \bar{q})_S^\downarrow.
\]

A large \(x\) anomaly in \(\Delta q(x, Q^2)\) would be manifest within these models as a different choice of the parameters in the spin dilution factors for each of \(\Delta q_V(x, Q^2)\) and \(\Delta q(x, Q^2)\). Given the present experimental error on the large \(x\) data points, there is already plenty of room to vary these parameters and still provide a good fit to the measured asymmetries [25]. In our picture the “spin” carried by the \(q_S\) is conserved. If one applies Schwinger’s derivation of the anomaly to each of the valence quarks in the model wavefunction, it follows that the screening effect is proportional to the valence quarks’ spin. It follows that the leading large \(x\) behaviour of
\( g_1 \sim (1 - x)^3 \) that follows from the counting rules [26] should not be affected by the presence (or otherwise) of \( \Gamma \). (It is interesting to note that the phenomenological parametrisation of the anomaly obtained in [27] has the correct leading \((1 - x)^3\) behaviour.)

3 How to look for a large \( x \) anomaly in semi-inclusive measurements of polarised deep inelastic scattering

In an ideal world one would like to measure the C-odd spin structure function directly and compare the C-even and C-odd polarised quark distributions at large \( x \). This experiment would involve a neutrino beam and a polarised target and is clearly impracticable at the present time. The best available tool is to reconstruct the valence distributions from semi-inclusive measurements of fast pions in polarised deep inelastic scattering. In the rest of this paper we discuss how a large \( x \) anomaly should show up in these experiments. We shall concentrate only on the region \( x > 0.2 \) where polarised gluons, the Dirac sea and mesonic effects are not important.

If we assume a large \( x \) anomaly then

\[
g_{1|x>0.2} = \frac{2}{9}(u^\uparrow - u^\downarrow)_{S} + \frac{1}{18}(d^\uparrow - d^\downarrow)_{S} + \frac{1}{3}\Gamma
\]

where we work within flavour SU(3). Given that the \( q_S \) are sufficient to describe the physics of the C-odd structure function and also the local coupling of \( \Gamma \) to the hard photon [7], we shall treat \( \Gamma \) as a new “parton”, unique to \( g_1 \), where “parton” is defined according to Feynman [28].

Given our “new parton”, we can modify the naive parton model analysis of semi-inclusive deep inelastic scattering [29,30] to include the background field contribution. We shall follow the notation of Frankfurt et al. [29]. The \( q_S \) fragment to a fast pion in the same way as the naive parton model quark – independent of its helicity. We let \( z \) denote the fraction of the hard photon’s energy which is taken by the fast pion. Following [29], we use \( D_1(z) \) to denote the favoured fragmentation

\[
D_1(z) \equiv D_{u_S}^+(z) = D_{d_S}^-(z) = D_{\bar{u}_S}^+(z) = D_{\bar{d}_S}^-(z)
\]

and \( D_2(z) \) to denote the unfavoured fragmentation function

\[
D_2(z) \equiv D_{\bar{u}_S}^+(z) = D_{d_S}^-(z) = D_{\bar{d}_S}^+(z) = D_{u_S}^- (z).
\]

The strange quark fragmentation function

\[
D_3(z) \equiv D_{s_S}^+(z) = D_{\bar{s}_S}^-(z) = D_{\bar{s}_S}^+(z) = D_{s_S}^- (z)
\]

is not relevant to our analysis since we are considering only the large \( x \) region \( x > 0.2 \) which is not sensitive to the strange quark components in the Fock expansion of the nucleon wavefunction.

In order to understand how an explicit background field contribution should fragment into fast pions in the final state it is important to note that \( \Gamma \) does not itself have any Fock components so that its fragmentation is constrained by the fragmentation of the valence quarks. (For the same reason, since the background field is a property of the vacuum – or bag quarks – its contribution to \( g_1 \) is independent of the number of accessible flavours.) We give \( \Gamma \) its own fragmentation function \( D_4(z) \). (We make the assumption here that \( q_S \) and \( \Gamma \) can separately be treated in impulse approximation and that the fragmentation of our \( \Gamma \) parton into fast pions...
factorises.) Since $\Gamma$ is a flavour singlet effect it will fragment equally into fast $\pi^+$ and $\pi^-$ in the final state, whence $D_4$ must be proportional to the sum of the favoured and unfavoured light-quark fragmentation functions $D_1$ and $D_2$. (The fragmentation properties of the anomaly were also discussed in ref. [31], where it was assumed that the anomaly can be associated entirely with the polarised gluon distribution $\Delta g(x,Q^2)$ – that is, a purely small $x$ effect.)

We let $N_{\uparrow \downarrow}(x, z)$ denote the number of $\pi^+$ produced in a bin characterised by Bjorken $x$ ($> 0.2$) and $z$, where the virtual photon helicity is $\uparrow$ and the target proton helicity is $\downarrow$. The spin averaged pion production rates are independent of the anomaly. It follows that:

\[
N_{\uparrow \downarrow}^{\pi^+} \sim \frac{4}{9} u_S(x) D_1(z) + \frac{1}{9} d_S(x) D_2(z) + \frac{1}{12} \Gamma(x) D_4(z)
\]

\[
N_{\downarrow \uparrow}^{\pi^+} \sim \frac{4}{9} u_S(x) D_1(z) + \frac{1}{9} d_S(x) D_2(z) - \frac{1}{12} \Gamma(x) D_4(z)
\]

\[
N_{\uparrow \downarrow}^{\pi^-} \sim \frac{4}{9} u_S(x) D_2(z) + \frac{1}{9} d_S(x) D_1(z) + \frac{1}{12} \Gamma(x) D_4(z)
\]

\[
N_{\downarrow \uparrow}^{\pi^-} \sim \frac{4}{9} u_S(x) D_2(z) + \frac{1}{9} d_S(x) D_1(z) - \frac{1}{12} \Gamma(x) D_4(z)
\]

whence

\[
N_{\uparrow \downarrow}^{\pi^+ - \pi^-} \sim \frac{4}{9} (u_S^\uparrow - u_S^\downarrow) - \frac{1}{9} (d_S^\downarrow - d_S^\uparrow) (D_1 - D_2)(z)
\]

\[
N_{\downarrow \uparrow}^{\pi^+ - \pi^-} \sim \frac{4}{9} (u_S^\downarrow + u_S^\uparrow) - \frac{1}{9} (d_S^\uparrow + d_S^\downarrow) (D_1 - D_2)(z)
\]

\[
N_{\uparrow \downarrow}^{\pi^+ + \pi^-} \sim \frac{4}{9} (u_S^\uparrow - u_S^\downarrow) + \frac{1}{9} (d_S^\downarrow - d_S^\uparrow) (D_1 + D_2)(z) + \frac{1}{3} \Gamma(x) D_4(z)
\]

\[
N_{\downarrow \uparrow}^{\pi^+ + \pi^-} \sim \frac{4}{9} (u_S^\downarrow + u_S^\uparrow) + \frac{1}{9} (d_S^\uparrow + d_S^\downarrow) (D_1 + D_2)(z)
\]

It follows from equs.(10) that the spin asymmetry $A_{\pi^+ - \pi^-}$ measures the C-odd valence part of $g_1^N$ in QCD; viz.

\[
A_{\pi^+ - \pi^-} = \frac{4 \Delta u_v - \Delta d_v}{4 u_v - d_v}
\]

for a proton target and (modulo nuclear effects [32])

\[
A_{\pi^+ - \pi^-} = \frac{\Delta u_v + \Delta d_v}{u_v + d_v}
\]

for a deuteron target. A measurement of these asymmetries will allow us to extract the valence spin distributions $\Delta u_v(x)$ and $\Delta d_v(x)$. Note that equs.(11,12) are the same expressions that one also obtains in the naive parton model with no background field contribution [29,30]. However, these C-odd distributions are not in general equal to the C-even distributions which describe the inclusive structure function $g_1$. It is a challenge for future experiments to check whether the parton distributions that one uses to describe $g_1$ [33] can also be used to describe these semi-inclusive C-odd asymmetries. This will require high quality data with much reduced errors in the large $x$ region. If different parton distributions are required to describe both $g_1$ and also the C-odd semi-inclusive asymmetries then we would evidence for a large $x$ anomaly. (Note that if one of our impulse or factorisation hypotheses were to fail, then it would be most unlikely that our parton distributions would describe both $g_1$ and equs.(11,12). In this respect, these assumptions are not necessary. They are necessary if one finds a large $x$ anomaly and then wishes to make predictions for other processes.)
In this analysis, the spin asymmetry which is obtained by summing over fast $\pi^+$ and $\pi^-$ in the final state can be used to deduce the fragmentation function for the large $x$ part of the anomaly. Given a high quality measurement of $\Gamma(x)$, we can use the asymmetry

$$A^{\pi^+\pi^-} = \frac{g_1^N(x)}{F_1(x)} - \left(1 - \frac{D_4(z)}{(D_1 + D_2)(z)}\right) \frac{1}{3} \frac{\Gamma(x)}{F_1(x)}$$

(13)

to extract the relative fragmentation of the anomaly viz. $\left(1 - \frac{D_4(z)}{(D_1 + D_2)(z)}\right)$. Since the final state hadrons are dominated by light mass pions it is reasonable to take

$$A^{\pi^+\pi^-} \approx A_1 \equiv \frac{g_1^N(x)}{F_1^N(x)}$$

(14)

whence

$$D_4(z) = D_1(z) + D_2(z)$$

(15)

describes the fragmentation of the background field contribution.

As a guide to the size of the effect that one is looking for we evaluate the asymmetry in equ.(12) using the phenomenological parametrisation of the anomaly in [27] (which was deduced via the MIT bag model). This is shown in Fig.1. Here the dashed curve is the semi-inclusive asymmetry assuming zero large $x$ anomaly; the bold curve is the model calculation with a large $x$ anomaly included. (Given the semi-classical nature of the model and the non-classical nature of the anomaly, it is important to regard Fig.1 more as a guide to the experimental accuracy that is required than a rigorous prediction of the asymmetry.)

In summary, semi-inclusive measurements of fast pions in polarised deep inelastic scattering offer a window where we may hope to make an explicit measurement of the background field in QCD. Given the programme of present [20,21] and proposed [34] spin experiments in this field together with data from future high-energy, polarised, hadron hadron collisions [4], we may soon begin to learn about this important non-perturbative physics.

Acknowledgements

I thank A.W. Thomas and W. Melnitchouk for helpful discussions. This work has been supported in part by the EC Programme “Human Capital and Mobility”, Network “Physics at High Energy Colliders”, contract CHRX-CT93-0357 (DG 12 COMA).
References

1. The EMC Collaboration, J. Ashman et al., Phys. Lett. B206 (1988) 364, Nucl. Phys. B328 (1990) 1

2. R. Windmolders, Int. J. Mod. Phys. A7 (1992) 639

3. S.D. Bass and A.W. Thomas, Prog. Part. Nucl. Phys. 33 (1994) 449
   G. Altarelli and G. Ridolfi, Nucl. Phys. B (Proc. Suppl.) 39B (1995) 106

4. O. Nachtmann and A. Reiter, Z Physik C24 (1984) 283
   A. Brandenburg, O. Nachtmann and E. Mirkes, Z Physik C60 (1993) 697
   O. Nachtmann, to appear in Proc. ELFE Summer School, Cambridge 1995 (Editions Frontieres 1995)

5. R. L. Jaffe, MIT preprint MIT-CTP-2466 (1995), submitted to Phys. Lett. B

6. H. Fritzsch, Phys. Lett. B256 (1991) 75

7. S.D. Bass, Z Physik C55 (1992) 653, ibid C60 (1993) 343

8. S.D. Bass, Phys. Lett. B342 (1995) 233

9. R.L. Jaffe, Phys. Lett. B193 (1987) 101
   R.L. Jaffe and A. Manohar, Nucl. Phys. B337 (1990) 509.

10. The NA-10 Collaboration, S. Falciano et al., Z Physik C31 (1986) 513
    The NA-10 Collaboration, M. Guanziroli et al., Z Physik C37 (1988) 545

11. A.V. Efremov and O.V. Teryaev, Dubna Preprint E2-88-287 (1988)
    G. Altarelli and G.G. Ross, Phys. Lett. B212 (1988) 391

12. J. Schwinger, Phys. Rev. 82 (1951) 664

13. R. Jackiw, in “Current algebra and anomalies”, eds. S. B. Treiman, R. Jackiw, B. Zumino and E. Witten (World-Scientific, 1985)

14. C. H. Llewellyn Smith, Oxford preprint OX-89/88 (1988)
    C. H. Llewellyn Smith, Nucl. Phys. A434 (1985) 35c

15. R. L. Jaffe and G. G. Ross, Phys. Lett. B93 (1980) 313
    R. L. Jaffe, Nucl. Phys. B229 (1983) 205

16. A. W. Schreiber, A. W. Thomas and J. T. Londergan, Phys. Rev. D42 (1990) 2226
    A. W. Schreiber, A. I. Signal and A. W. Thomas, Phys. Rev. D44 (1991) 2653

17. J. Kogut and L. Susskind, Phys. Rev. D11 (1974) 3594

18. J. D. Sullivan, Phys. Rev. D5 (1972) 1732

19. G. M. Shore and G. Veneziano, Nucl. Phys. B381 (1992) 23
20. The HERMES Proposal, K. Coulter et al., DESY/PRC 90-1 (1990)  
M. Düren, Proc. Zuoz Summer School, PSI-Proceedings 94-01 (1994) 273

21. W. Wislicki, SMC report, [hep-ex 9405012] (1994)

22. J. Kuti and V. Weisskopf, Phys. Rev. D4 (1971) 3418

23. F. E. Close, Nucl. Phys. B80 (1974) 269

24. R. Carlitz and J. Kaur, Phys. Rev. Lett. 38 (1977) 673

25. A. Schäfer, Phys. Lett. B208 (1988) 175

26. S.J. Brodsky, M. Burkardt and I. Schmidt, Nucl. Phys. B441 (1995) 197

27. S.D. Bass and A.W. Thomas, Phys. Letts. B312 (1993) 345

28. R. P. Feynman, Phys. Rev. Lett. 23 (1969) 1415

29. L.L. Frankfurt et al., Phys. Lett. B320 (1989) 141

30. F.E. Close and R.G. Milner, Phys. Rev. D44 (1991) 3691

31. St. Gullenstern et al., Phys. Lett. B312 (1993) 166

32. W. Melnitchouk, G. Piller and A. W. Thomas, Phys. Lett. B346 (1995) 165  
S. A. Kulagin, W. Melnitchouk, G. Piller and W. Weise, Phys. Rev. C52 (1995) 932

33. T. Gerhmann and W. J. Stirling, Z Physik C65 (1995) 461

34. The HMC Collaboration, E. Nappi et al., Letter of Intent, CERN/SPSLC 95-27 (1995)

**Figure caption**

Fig.1.
The deuteron asymmetry equ.(12) evaluated in the model of ref.[27]. The bold curve is the spin asymmetry given a large $x$ anomaly contribution to $g_1$. The dashed curve is the asymmetry with no large $x$ anomaly contribution.
