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Partons and The EMC Spin Effect

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Abstract

We focus on the parton model and the role of the axial anomaly in polarised deep inelastic scattering. We show that the axial anomaly is relevant to each of the higher moments of the spin dependent structure function $g_1(x)$ and not just the first moment. This result implies that the factorisation of mass singularities is not sufficient to define the parton model in spin dependent QCD. (It is certainly a necessary condition.) We also need to consider the locality of the photon parton interaction. The anomaly is observed over all x in the EMC $g_1(x)$ data.

1. Introduction

In recent times there has been much excitement in the high energy physics community about polarised deep inelastic scattering (pDIS) from hadronic targets, particularly after the announcement of the EMC Spin Effect [1]. EMC measured the polarised proton structure function $g_1^P(x)$. After a Regge extrapolation of their data to $x = 0$ they found

$$\int_0^1 dx \; g_1^P(x) = 0.126 \pm 0.010(stat.) \pm 0.015(syst.) \tag{1}$$

The flavour singlet contribution to this quantity is the flavour singlet axial charge Δq_0. In the old parton model of Feynman Δq_0 determines the fraction of the proton's spin which is carried by its quarks. Hence the surprise when EMC found

$$\Delta q_0 = 0.120 \pm 0.094(stat.) \pm 0.138(syst.) \tag{2}$$

viz. consistent with zero.

In an interacting gauge theory like QCD Δq_0 does not measure quark spin after all because of the axial anomaly. As Veneziano has stressed [2], the EMC Spin Effect should be interpreted as a violation of OZI (in this case the Ellis-Jaffe hypothesis [3]) which is catalysed by the strong U(1) axial anomaly.

The partonic interpretation of the axial anomaly has been the source of much debate in the literature [4-17]. This work has been motivated by the need to explain the "spin crisis" and hence has concentrated on the EMC Spin Effect as a first moment problem. In this paper we argue that the EMC Spin Effect is really an all moment problem - viz. that the axial anomaly is relevant to each of the higher moments of $g_1(x)$. Rather than complicating things, this result simplifies the partonic interpretation of $g_1(x)$ considerably. However, we have to give up the idea that factorisation of mass singularities is sufficient to define the parton model in QCD. (It is certainly a necessary condition.) We need to introduce an extra locality condition so that the quark parton distribution is defined to include all partons that make a local interaction with the hard photon. The anomaly is observed over a complete range of x - even in the "valence region" $x \geq 0.2$. This is in contrast with the usual spin dependent gluon distribution which is observed only at small x - essentially outside the range of the present data.

We also consider the parity odd structure function $g_3(x)$ which occurs with pDIS using an (anti-)neutrino beam and a polarised proton target. We argue that $g_3(x)$ is anomaly free. Hence, $g_1(x)$ and $g_3(x)$ may be significantly different at large x. The polarised version of the Gross-Llewellyn Smith sum rule does measure a valence quark spin component in the proton. We conclude the paper by explaining how these results relate to semi-exclusive DIS where, for example, we separate out fast moving pions or kaons from the final state hadrons.

In section two we briefly review what is known about spin dependent parton distributions in QCD. We focus on the flavour singlet part of $g_1(x)$. The new results are presented in section 3, where we discuss the relationship between gauge invariance and the locality of photon parton scattering.

2. Parton Distributions in QCD

Formally, deep inelastic scattering is described in QCD by the operator product expansion (OPE) and the renormalisation group equations (RGE) (see eg. ref. [18]). We may write the flavour singlet contribution to $g_1(x)$ as a sum over spin dependent quark and gluon distributions each of which is convoluted with the relevant Wilson coefficient distribution, viz.

$$g_1(x,Q^2)|_S = \frac{1}{3}\sqrt{\frac{2}{3}}\int_x^1 \frac{dz}{z}\Delta q_0(z,Q^2)C_S^q(\frac{x}{z},\alpha_s(Q^2)) + \frac{1}{9}\int_x^1 \frac{dz}{z}\Delta g(z,Q^2)C_S^g(\frac{x}{z},\alpha_s(Q^2)) \tag{3}$$

In leading twist approximation, the quark $\Delta q_k(x,Q^2)$ and gluonic $\Delta g(x,Q^2)$ distributions are defined by relation to the OPE so that their odd moments satisfy

$$2Ms_+(p_+)^{2n}\int_0^1 dx\ x^{2n}\Delta q_k(x,Q^2) = < p,s|[\overline{q}(0)\gamma_+\gamma_5(iD_+)^{2n}\frac{\lambda_k}{2}q(0)]_{Q^2}|p,s >$$

$$2Ms_+(p_+)^{2n}\int_0^1 dx\ x^{2n}\Delta g(x,Q^2) = \tag{4}$$

$$< p,s|\mathrm{Tr}\ [G_{+\alpha}(0)(iD_+)^{2n-1}\tilde{G}^{\alpha}{}_{+}(0)]_{Q^2}|p,s > \quad (n\geq 1)$$

Here $G_{\mu\nu}$ is the gluon field tensor, $\tilde{G}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\alpha\beta}G^{\alpha\beta}$ is the corresponding dual tensor and $D_\mu = \partial_\mu + igA_\mu$ is the gauge covariant derivative in QCD. The kinematic variables p_μ and s_μ denote the proton momentum and spin vectors respectively (M is the proton mass). These gauge invariant flavour singlet quark and gluonic distributions mix under the renormalisation group according to Altarelli Parisi evolution [19].

The coefficient distributions $C^q(x,\alpha_s) = \delta(x-1) + O(\alpha_s)$ and $C^g(x,\alpha_s) \sim O(\alpha_s)$ are defined likewise. Their even moments reproduce the Wilson coefficients which accompany the quark and gluonic operators in the OPE. We have chosen to define the composite operators at the renormalisation scale $\mu^2 = Q^2$. The coefficient distributions are target independent and are calculated in perturbation theory. The parton distributions contain all of the target dependent information.

In this OPE language, the first moment of $g_1(x)|_S$ is given entirely by the target matrix element of the flavour singlet axial vector current multiplied by the relevant Wilson coefficient. There is no twist two, spin one, local gauge invariant gluonic operator which can contribute to the first moment of $g_1(x)$ [13]. This implies that there is no OPE $\Delta g = \int_0^1 dx\ \Delta g(x)$ term in the first moment of $g_1(x)$. One has

$$\Delta g\int_0^1 dx\ C^g_S(x,\alpha_s(Q^2)) = 0 \tag{5}$$

for an arbitrary choice of target. Although the first moment of $\Delta g(x)$ is not given by any single local gluonic operator there is no reason to suppose that it is identically zero - especially for an ideal free gluon target. The Mellin theorem uniquely determines Δq_0 in terms of its moments. However, the lack of any gauge invariant local gluonic operator with which to define the first moment of $\Delta g(x, Q^2)$ means that $\Delta g(x, Q^2)$ is uniquely defined up to a delta function at the origin. This result and the general target dependence of Δg leads us to the conclusion that the target independent coefficient in equ.(5) must vanish to all orders in perturbation theory. This result has been argued strongly by Bodwin and Qiu [12] for any choice of gauge invariant renormalisation scheme (see also [14]). It has the status of a theorem in QCD.

Whilst the OPE description of DIS is formally correct it is easy to lose track of the physics under towers of operator matrix elements. We would like a more intuitive description of what is going on. This leads us to the QCD Improved Parton Model (IPM) [20-24]. It is constructed following the OPE result for DIS. We factor the DIS cross section into a convolution of the IPM parton distributions of the target hadron with a hard parton scattering cross section.

In the parton model one usually defines the quark and gluon distributions at a scale Λ^2 to include all partons with transverse momentum squared $k_T^2 \leq \Lambda^2$. Thus, the IPM gluon distribution $\Delta g(x, \Lambda^2)_{IPM}$ is observed via two-quark-jet events with large $k_T^2 \geq \Lambda^2$ [10, 23]. It contributes to $g_1(x)$ as a convolution with the spin dependent asymmetry for producing $q\overline{q}$ jets from photon-gluon fusion with transverse momentum squared $k_T^2 \geq \Lambda^2$ (see Fig. 1). Once we impose this cut-off on the k_T^2 the hard photon-gluon scattering cross section is free of any mass singularity. All quarks with $k_T^2 \leq \Lambda^2$ are understood to be factored into the quark distribution. The picture that we have just described is the simplest form of factorisation. It dates from the work of Gribov and Lipatov [24]. The k_T^2 cut-off parton model has been developed to all orders of perturbation theory (for a review see [22]) and has application to a host of different hadronic processes. The

IPM parton distributions which are measured in one process (eg. DIS) can be used to make testable predictions in other hadron interactions (eg. Drell Yan and exclusive jet production). There is excellent agreement between the calculations and experiments without polarisation.

In this picture, which is based entirely on factorisation in k_T^2, one finds that the IPM gluon distribution does contribute to the first moment of $g_1(x)$ - in apparent contradiction with the OPE. If we define $\Delta q = \int_0^1 dx \Delta q(x)$ then one finds that (for each flavour q)

$$\Delta q_{OPE} \to (\Delta q - \frac{\alpha_s}{2\pi}\Delta g)_{IPM} \tag{6}$$

in the expression for the first moment of $g_1(x)$. This is the Efremov-Teryaev, Altarelli-Ross result [4, 5]. The gluon term in the first moment of $g_1(x)$ is induced by a local photon gluon interaction, which appears to generate two quark jet events with k_T^2 of order Q^2 [6]. There is no contribution to the first moment from jets with a range of k_T^2 between the factorisation scale Λ^2 and Q^2, which is in contrast with the results in unpolarised DIS and the higher moments of $g_1(x)$.

The apparent contradiction between the OPE and IPM results for the first moment of $g_1(x)$ has a simple resolution when one considers the axial anomaly in QCD. The axial anomaly [25] is related to the definition of γ_5 and hence is relevant to each moment of $g_1(x)$. The simplest way to see this is at one loop in perturbation theory. We now evaluate the OPE quark distribution of a gluon at $O(\alpha_s)$. How we treat the axial anomaly determines the gluonic Wilson coefficients for each moment.

3. The Axial Anomaly and pDIS

The spin dependent asymmetry $\Delta\sigma(x, Q^2, P^2, m^2) = \sigma(\gamma \uparrow \; g \uparrow \to q\bar{q}) - \sigma(\gamma \uparrow \; g \downarrow \to q\bar{q})$ for photon-gluon fusion at $O(\alpha_s)$ is calculated from the box graphs in Fig.

1. This asymmetry is just $g_1(x)$ for an ideal gluon target at one loop. We may apply the OPE and write $\Delta\sigma(x, Q^2, P^2, m^2)$ as the sum of the target independent gluonic Wilson coefficient distribution $C^g(x, \frac{Q^2}{\mu^2}, \alpha_s(\mu^2))$ and the OPE defined spin dependent quark distribution of the gluon $\Delta q^{gluon}(x, \mu^2)$. Here μ^2 is the renormalisation scale in the problem. The quark distribution is defined by the absorptive part of the forward vector, vector, axial-vector (VVA) amplitudes, which are obtained by replacing the axial vector current by the general spin-odd axial-tensor operator $\overline{q}(0)\gamma_+\gamma_5(D_+)^{2n}q(0)$ in the Adler Bell Jackiw triangle amplitude. In $A_+ = 0$ gauge, $\Delta q^{gluon}(x)$ is given by the x dependent triangle amplitude (see Fig. 2) [6]

$$\begin{aligned} 2p_+\Delta q^{gluon}(x) &= -ig^2 2T \\ \int \frac{d^4k}{(2\pi)^4}\delta(x - \frac{k_+}{p_+})\delta((k-p)^2 - m^2)&\frac{\mathrm{Tr}[(\hat{k}+m)\hat{\epsilon}^*(\hat{k}-\hat{p}+m)\hat{\epsilon}(\hat{k}+m)\gamma_+\gamma_5]}{(k^2-m^2)^2} \end{aligned} \tag{7}$$

Here we consider one flavour of quark and $T = \frac{1}{2}$ is a group factor; m and p_μ denote the quark mass and gluon momentum respectively. We shall evaluate $\Delta q^{gluon}(x, \mu^2)$ via minimal subtraction. This calculation provides the most transparent insight into the role of the axial anomaly in pDIS. With dimensional regularisation the axial anomaly becomes a problem of how to continue the γ_5 into the regulator dimensions. The correct procedure was established by 't Hooft and Veltman [26]. Gauge invariant regularisation is equivalent to the continuation

$$\begin{aligned} \{\gamma_\mu, \gamma_5\}_+ &= 0 \qquad \mu = 0,1,2,3 \\ [\gamma_\mu, \gamma_5]_- &= 0 \qquad \mu = \text{regulator dimensions} \end{aligned} \tag{8}$$

where $\gamma_5 = i\gamma_0\gamma_1\gamma_2\gamma_3$. This 't Hooft Veltman prescription for the γ_5 reproduces the axial anomaly. Hence, it should be used consistently for the renormalisation of each of the axial tensor operators $\overline{q}(0)\gamma_+\gamma_5(D_+)^{2n}q(0)$. (The choice $\{\gamma_\mu, \gamma_5\}_+ = 0 \ \forall\mu$ defines the gauge dependent Noether symmetry current A^S_μ, which satisfies the divergence equation $\partial^\mu A^S_\mu = 2m\bar{q}i\gamma_5 q$. We shall refer to it as the S continuation.)

We first integrate over k_+ and k_- to evaluate the the momentum conserving delta functions in equ.(7). The Dirac trace is evaluated in $4+\epsilon$ regulator dimensions. We find

$$\Delta q^{gluon}(x,\mu^2) = \pm \frac{\alpha_s}{2\pi^2}\mu^{-2\epsilon}\int d^{2+\epsilon}k_T \left[\frac{(k_T^2+m^2)(1-2x)-2m^2(1-x)-2\frac{\epsilon}{2+\epsilon}k_T^2(1-x)}{(k_T^2+m^2+P^2x(1-x))^2}\right] \tag{9}$$

Here μ^2 is the renormalisation scale and $P^2=-p^2$ is the virtuality of the gluon. The overall positive (negative) sign indicates a left (right) handed gluon polarisation. The $\frac{\epsilon}{2+\epsilon}k_T^2(1-x)$ term comes from the continuation of γ_5 into the regulator dimensions and is absent when we use the S continuation. If no ultraviolet cut-off is imposed the k_T^2 integral develops a $\frac{1}{\epsilon}$ pole. This pole cancels with $\frac{\epsilon}{2+\epsilon}k_T^2(1-x)$ to give a finite contribution in the limit $\epsilon \rightarrow 0$, which is the axial anomaly [6, 9].

Our one loop parton distribution and hence the gluonic coefficient distribution depends upon how we choose to continue the γ_5 into the regulator dimensions. We find the result

$$\Delta q^{gluon}(x,\mu^2)|_{\text{'tHV}} - \Delta q^{gluon}(x,\mu^2)|_{\text{S}} = \mp\frac{\alpha_s}{\pi}(1-x) \tag{10}$$

which comes from the maximum possible k_T^2 and describes a contact interaction between the hard photon and the gluon target. This result is important for our understanding of the parton model in spin dependent QCD.

Each moment of $\Delta q^{gluon}(x,\mu^2)$ and, hence, each gluonic Wilson coefficient is dependent upon how we treat the anomaly. The one loop gluon matrix elements of the axial tensor operators $\bar{q}(0)\gamma_+\gamma_5(D_+)^{2n}q(0)$ change by a polynomial in the external gluon momentum when if we swap from the 't Hooft Veltman to the S continuation of γ_5. In other words, there is a generalised anomalous gluonic (gauge dependent) counterterm associated with each of the spin-odd axial tensor operators. At order $O(\alpha_s)$ these local counterterms generate the (gauge independent) anomalous contribution to the gauge invariant parton distribution $\Delta q^{gluon}(x, \mu^2)$, viz. equ.(10).

Our calculation at $O(\alpha_s)$ is really QED up to the group factor $T = \frac{1}{2}$, which comes from summing over colour indices at the quark gluon vertex. From equ.(10) we can deduce the generalised (spin 2n+1) gauge dependent counterterm in QED. In $A_+ = 0$ gauge it reads

$$k_{+\ldots+(2n+1)} = \frac{2\alpha}{\pi} \xi_{2n} \epsilon_{+\lambda\alpha\beta} A^\alpha \partial^\lambda (\partial_+)^{2n} A^\beta \tag{11}$$

where $\xi_{2n} = \int_0^1 dx x^{2n}(1-x)$. These gauge dependent counterterms could also be isolated in QED by calculating the surface term which arises when we make the momentum shift $k_\mu \to k_\mu + a_\mu$ in the general VVAxial tensor amplitude. There will be a contribution to the surface terms from the gauge invariant photon operators $F_{+\alpha}(i\partial_+)^{2n-1}\tilde{F}^\alpha_{\ +}$ which enter due to operator mixing under renormalisation. These terms are subtracted out to isolate the gauge dependent counterterm.

The gauge-invariant axial tensor operators appear in the OPE. These operators are defined following the 't Hooft and Veltman prescription for γ_5. Thus, the OPE spin dependent quark distribution includes the local photon gluon interaction. It can only be shifted into the gluonic Wilson coefficient distribution (as the k_T^2 IPM would suggest) if we use the S continuation of γ_5 into the regulator dimensions. However, the axial tensor operators which are constructed via the S continuation are not gauge invariant. They differ from the gauge invariant axial tensor operators by a generalised gauge dependent gluonic counterterm. Hence, the quark distribution which is defined following the S continuation of γ_5 is not gauge invariant. It is defined in a gauge non-invariant world where chiral symmetry is exact in the Noether sense. In summary, the parton model which is based entirely upon factorisation in k_T^2 is inconsistent with gauge invariance when we consider spin dependent processes.

Furthermore, the gauge dependent S distributions do not satisfy Altarelli-Parisi evolution. The gauge invariant operators in the OPE are renormalised independently of any gauge non-invariant operators. In other words, it is the gauge invariant spin dependent quark distribution which mixes with the OPE defined gluon distribution under QCD evolution. When we write this gauge invariant quark distribution as the sum of anomalous and non-anomalous contributions then it follows that each behaves identically so far as mixing with gluon distribution is concerned. The S distribution *by itself* does not contribute to the renormalisation of the OPE $\Delta g(x)$. The formal basis for Altarelli-Parisi evolution is that the moments of the splitting functions reproduce the anomalous dimensions of the gauge invariant operators in the OPE. Thus, the Altarelli Parisi equations describe the evolution of the gauge invariant quark and gluon distributions which are defined via the OPE in equs. (3) and (4).

The IPM parton distributions can be made gauge invariant if we *define* the quark distribution to *include* all partons which make a local interaction with a hard photon in DIS. This definition was previously proposed by Gribov in a little noticed remark at the SLAC Lepton Photon Symposium [27]. It means that we need to factor the large $k_T^2 \sim Q^2$ part of the total phase space for photon-gluon fusion into the quark distribution $\Delta q(x, \Lambda^2)$ as well as the soft part $k_T^2 \leq \Lambda^2$. That is, we also have to worry about the top endpoint of the k_T^2 integration. The local photon-gluon interaction is a renormalisation effect associated with the gluonic dressing of the quark partons.

At this stage, one might think of splitting the OPE defined quark distribution $\Delta q(x, Q^2)$ into an intrinsic quark and a second (anomalous) gluon distribution, both of which are gauge invariant. The topological charge density and its higher spin generalisation $\frac{\alpha_s}{2\pi} G_{\mu\nu} D_{\mu_1} ... D_{\mu_{2n}} \tilde{G}^{\mu\nu}$. are not good candidate operators with which to define the second gluon distribution. Firstly, the hadronic matrix elements of the topological charge density have large isospin violations in them [28]. Furthermore, when we add any number of gauge covariant derivatives the resultant operator is not a topological invariant. This can be verified by varying with respect to the gluon field. Hence, $\frac{\alpha_s}{2\pi} G_{\mu\nu} D_{\mu_1} ... D_{\mu_{2n}} \tilde{G}^{\mu\nu}$ is not a total derivative for $n \geq 1$. The only sensible way to separate out a second spin dependent gluon distribution would be to generalise the renormalisation group analysis of Shore and Veneziano [29] to all of the higher moments of $g_1(x)$. However, this requires a prior knowledge of the scale dependence of the new anomalous distributions beyond $O(\alpha_s)$. The separation of $\Delta q_0(x)$ into intrinsic quark and gluon distributions is seemingly arbitrary unless the the two terms have a different experimental signature (see below). It is better to talk in terms of the distribution $\Delta q_0(x)$ which has nothing to do with spin in the interacting theory due to the axial anomaly.

Since the axial anomaly contributes to $g_1(x)$ through the OPE quark distribution $\Delta q(x, Q^2)$ it is convoluted with the quark coefficient distribution $C_S^q(x) = \delta(x-1) + O(\alpha_s)$ (see equ.(3)). This means that it can appear over a complete range of x - even in the "valence region" $x \geq 0.2$. One can see that the anomaly contributes at large x by the following physical argument. The intrinsic glue carries about 50% of the proton's momentum. Hence, a non-perturbative gluon can exist at x close to one. In the local photon gluon interaction this gluon is unable to first radiate away momentum and is seen in its bare state (at large x) by the hard photon. This is in contrast to the usual partonic gluon distribution $\Delta g(x, Q^2)$ which is observed only at very small x - essentially outside the range of the present data [15, 16, 17]. In the latter case, the gluon interacts with the hard photon by first radiating into a $q\overline{q}$ pair, which dissipates the momentum.

Carlitz et al. [6] have suggested that the axial anomaly should be characterised by two-quark-jet events with large $k_T^2 \sim Q^2$. If this is correct then we should expect to see these events over all x in both polarised and unpolarised DIS. (Recall that $g_1(x)$ is measured in the difference of two spin dependent deep inelastic muon-proton cross sections - the sum of which defines the unpolarised structure function. Cross sections are positive definite by definition. Any effect seen in the difference of two positive quantities is seen in their sum.) A significant two-quark-jet cross section at large x in unpolarised DIS has no place in our understanding of the unpolarised gluon distribution. For this reason, I believe that the cross section for two quark jet events with $k_T^2 \sim Q^2$ at large x will be found to be negligible if any experiment is made to measure it. These jets are probably suppressed by some soft effects in QCD, for example those which generate the isospin dependence of the hadronic matrix elements of the topological charge density $\frac{\alpha_s}{2\pi} G_{\mu\nu}\tilde{G}^{\mu\nu}$. (Non-perturbative effects have been discussed in ref. [30]). A small two-quark-jet cross section does not imply a small anomaly contribution to Δq_0. Altarelli and Stirling [7] have investigated the possibility of reconciling the EMC data with the Ellis-Jaffe hypothesis by invoking a large Δg in the proton. Following previous work [4, 5, 6], they tried to relate Δg to the anomaly in an IPM defined purely by factorisation of soft k_T^2. They convoluted $\Delta g(x)$ with $\delta(x-1)$ giving a gluon contribution that was much too inflated at large x. However, the preceding argument shows that they were on the right track - only that the axial anomaly has nothing to do with $\Delta g(x)$.

When polarised e^+e^- beams are available at LEP or at SLAC it will be possible to measure the polarised photon structure function $g_1^\gamma(x)$. Here we also have to consider the axial anomaly in QED [31]: the generalised anomalous currents in QED (see equ.(11)) describe a contact measurement of the soft target photon at leading order α. To $O(\alpha)$ there are two spin dependent photon distributions of the polarised photon: the usual partonic one $\Delta\gamma(x, Q^2)$, which is defined via the operators $F_{+\alpha}(i\partial_+)^{2n-1}\tilde{F}^{\alpha}{}_{+}$, and also a distribution which describes a contact hard-photon soft-photon interaction. This second photon distribution has the form given in equ.(10) where we make the replacement $\alpha_s \rightarrow 2\alpha$. It is gauge invariant at $O(\alpha)$ and can be made gauge invariant to all orders if we construct it from the derivative operators $F_{\mu\nu}(i\partial_+)^{2n}\tilde{F}^{\mu\nu}$ instead of the generalised anomalous currents. Since the axial anomaly is in the gauge invariant axial tensor operators, it follows that this anomalous distribution behaves exactly as the usual spin dependent quark distribution when it mixes with $\Delta\gamma(x, Q^2)$ under QED evolution. With the polarised photon structure function we would expect to see large $k_T^2 \sim Q^2$ jets (as described in ref. [6]) as a signature of the axial anomaly *per se* in QED. The important difference is that the axial anomaly in QED describes a contact measurement of the physical target photon, unlike the case with strong U(1) anomaly in QCD and hadronic targets. Since the contact hard-photon soft-photon interaction is observed over all x in the unpolarised photon structure function [32], a finite two-quark-jet cross section at large x would not violate the positivity condition which we used when discussing the polarised proton target.

So far we have concentrated on the anomalous (gauge dependent) counterterms associated with each of the spin-odd axial-tensor operators, which define $\Delta q(x)$. The spin dependent valence distribution

$$\Delta q_V(x) = (q^{\uparrow} - \overline{q}^{\uparrow})(x) - (q^{\downarrow} - \overline{q}^{\downarrow})(x) \tag{12}$$

could (in principle) be measured in polarised DIS with a polarised proton target and (anti-)neutrino beam in the parity odd structure function $g_3(x)$. (The cross section is infinitesimal in practice.) It is formally defined by the forward proton matrix elements of the spin-even axial tensor operators $\overline{q}(0)\gamma_+\gamma_5(D_+)^{2n-1}q(0)$. The spin-even axial-tensor operators have odd charge conjugation. The spin-even gluonic operators $G_{+\alpha}(0)(D_+)^{2n-2}\tilde{G}^{\alpha}{}_{+}(0)$ are even under charge conjugation and do not contribute to DIS. There is no operator mixing between quark and gluonic operators with $\Delta q_V(x)$. Similarly, we cannot construct a generalised spin-even, anomalous gluonic counterterm with odd charge conjugation. This means that the valence distribution is anomaly free.

The polarised analogy to the Gross-Llewellyn Smith sum-rule [33], viz.

$$\int_0^1 dx\ g_3^{\nu+\overline{\nu}}(x) = (\Delta q_V)_0 \tag{13}$$

(above the charm production threshold) does measure a valence quark spin component in the proton.

The anomaly is (in principle) distributed over all x in $\Delta q(x)$. Hence, the violation of the Ellis-Jaffe hypothesis suggests that $\Delta q_V(x)$ and $\Delta q(x)$ may be significantly different at large x. This is indicated in Fig. 3. It means that one has to be careful when trying to extract the valence distribution from semi-exclusive measurements of $g_1(x)$ (eg. fast pions in the final state). The method proposed in ref. [34] assumes a negligible sea contribution to $\Delta q(x)$ in the traditional "valence region" $x \geq 0.2$. However, the axial anomaly is a property of the sea. The experiment devised in ref. [35] is free from this assumption and would provide a measurement of the spin dependent valence distribution. It would be interesting to measure $\Delta q_V(x)$ in this way. One could then determine the x dependence of the sea contribution to $g_1(x)$ (which includes the anomaly) at large x.

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Figures :

Fig.1 : The box graphs relevant to photon-gluon fusion.

Fig.2 : The OPE defined quark parton distribution of a gluon is calculated at $O(\alpha_s)$ (and in $A_+ = 0$ gauge) from the forward x dependent triangle graph.

Fig.3 : Here we show the contributions to $g_1(x)$ from $\Delta g(x)$ and the axial anomaly. The anomaly is manifest over all x (see the dotted region) whilst the usual partonic spin dependent gluon distribution is manifest only at small x-essentially outside the range of the present data (the shaded region). The bold curve is the EMC fit to $g_1(x)$. The dashed curve is a possible shape for $g_3(x)$. The point is that the total and valence distributions may be significantly different at large x.

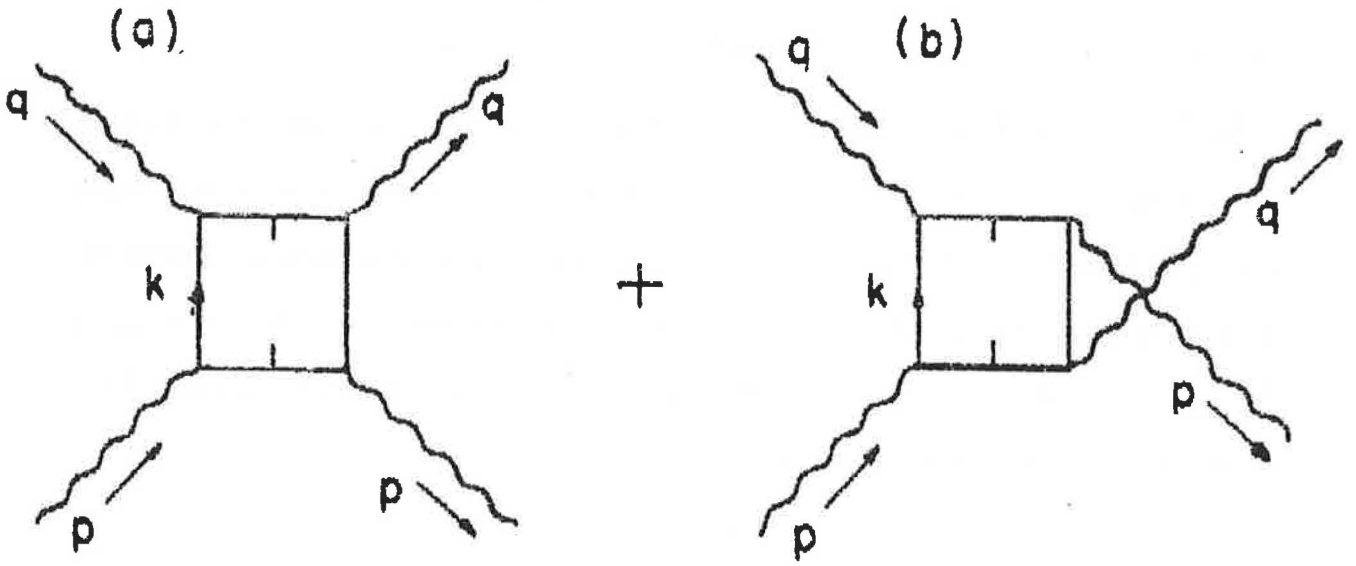

Fig 1

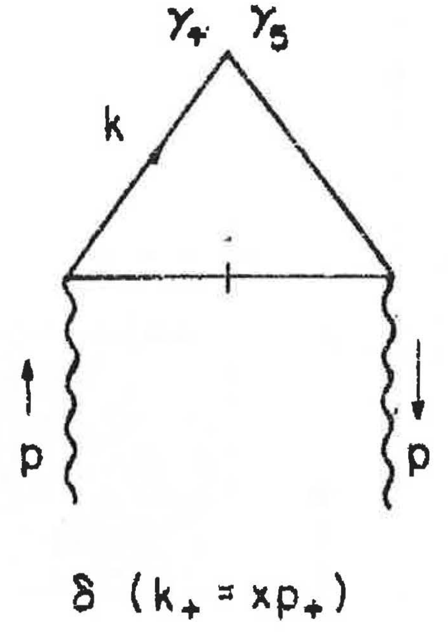

Fig 2

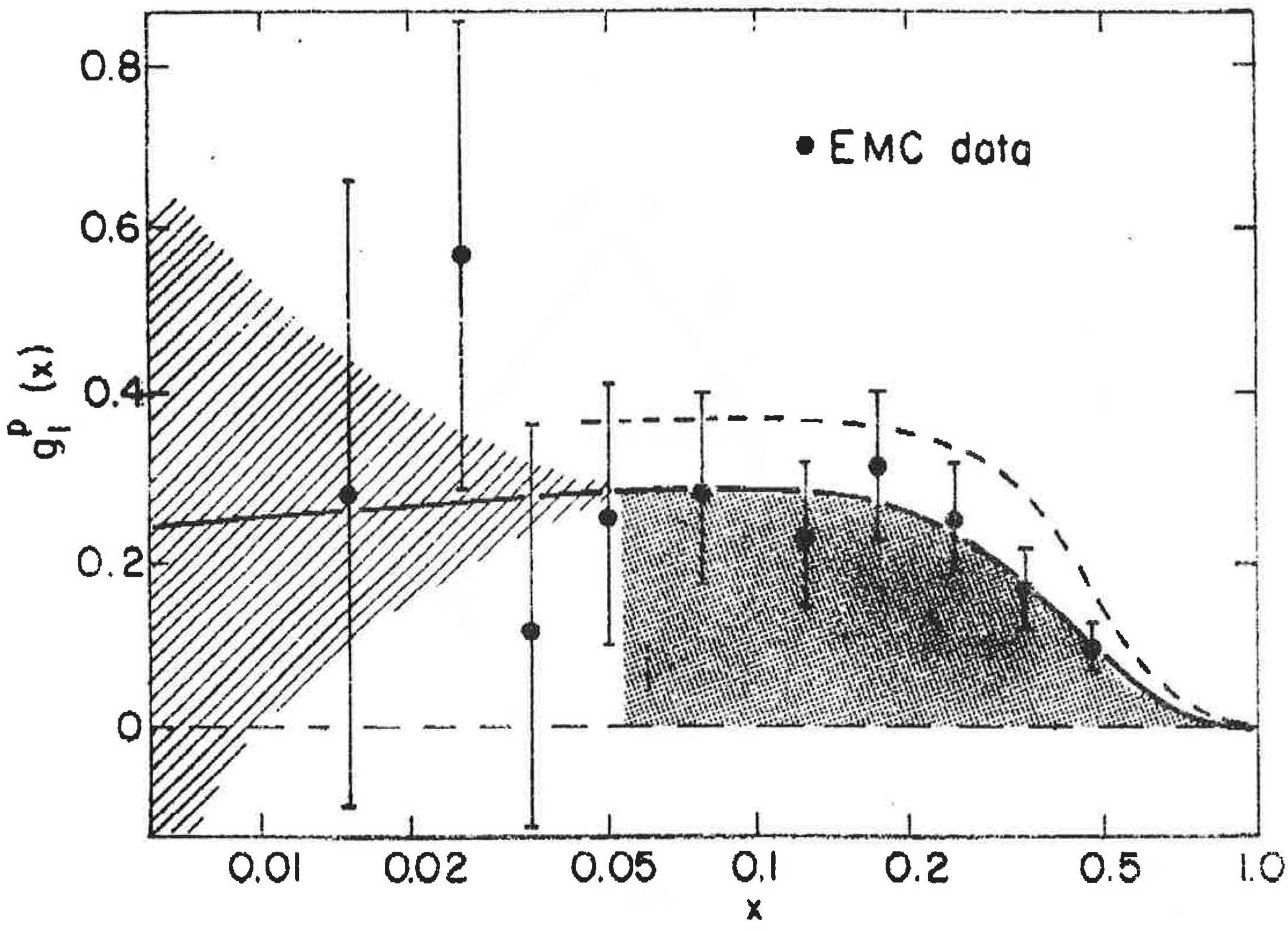

Fig 3