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THE PROTON "SPIN CONTENTS": Current Status & Perspectives

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Abstract

The present status of the phenomenological and theoretical interpretation of the EMC result on the polarized deep inelastic scattering is reviewed. We focus our discussion on the possibility of a significant gluonic contribution to the proton spin *via* the axial anomaly. We contrast the variant perspectives on this question: the viewpoint that stresses the interpretation in terms of the parton distributions *vs.* the one that concentrates on the matrix elements of local operators. Some remarks concerning the validity of OZI rule for the strange quark are also included.

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In this report we shall discuss the various issues raised by the EMC result concerning the intrinsic structure of the proton spin. In Sec. I we review the basis for interpreting the EMC data as indicating a significant proton matrix element of the strange quark axial vector current, which almost cancel that of the u, d quark contributions. In Sec. II we discuss the suggestion that these matrix elements $\Delta q'$ may contain significant gluonic fraction Δg because of axial anomaly:

$$\Delta q' = \Delta q - \frac{\alpha_s}{2\pi}\Delta g \tag{1}$$

where Δq & Δg are the spin-dependent distributions of the quark (of a particular flavor q) and the gluon, respectively. Arguments for the anomaly as local probe of the gluonic distribution, and some specific features of infrared collinear regularization of the perturbative QCD calculations will be reviewed. In Sec. III we emphasize the advantage, for the problem at hand, of working directly with the current divergence instead of the axial vector current itself. By taking the forward proton matrix element of the familiar axial anomaly equation, we obtain

$$\Delta q' = \Delta Q - \frac{\alpha_s}{2\pi}\Delta G \tag{1a}$$

where ΔQ & ΔG are defined directly in terms of the matrix elements of local operators $2im\bar{q}i\gamma_5 q$ and $TrG\tilde{G}$, respectively. The second viewpoint allows us to see more clearly the way how the decoupling theorem is enforced for the heavy quark contribution $\Delta q' \to 0$. This approach also naturally suggests a useful comparison with the anomaly equation for the trace of the energy momentum tensor. The proton matrix element of this equation yields a sum rule for the proton mass. In this connection the problem of pion-nucleon sigma term with its possible interpretation as indicating a significant proton matrix element of the strange quark scalar-density will be briefly recalled. In Sec. IV we show how ΔG can be estimated using the current divergence equation, and its peculiar feature of large isospin violation will be analyzed. In Sec. V, we summarize and contrast the various perspectives of this "spin content" problem. Finally we emphasize the relevance of the problem of the strange quark content of the nucleon as an important probe of the non-perturbative QCD. Our approach suggests that the quark model OZI rule is not generally applicable for the strange quark, and future investigation as prompted by such a suggestion is briefly outlined.

While concentrating on the question of an anomalous gluonic contribution, we will not review any work on low energy models of the proton spin. Work concerning models of parton distribution functions also will not be covered in this presentation. Thus our references represent only an incomplete set of the large number of papers published on this subject in the last couple of years.

I. The EMC result & its interpretations

Let us recall that the EMC collaboration measured the longitudinal spin-spin asymmetry in the deep inelastic muon-proton scattering [1]. This quantity is directly related to the spin-asymmetry of the virtual Compton scattering of polarized proton by the (circularly) polarized virtual photon. The EMC measurements extended the old polarized experiment performed at SLAC to considerable higher Q^2 (up to 70 GeV^2) and much smaller x (down to $x \sim 0.01$). In the region they overlap the two sets of data seem to agree. The asymmetry measurements directly yield information on the spin-dependent structure function $g_1(x,Q^2)$. Through the operator product expansion one can show that the first moment of g_1 is then related to the proton matrix element of the axial vector current (weighted by the quark charge-squared) [2]. Including the leading QCD radiative correction to the Wilson coefficient function [3], we have

$$\int_0^1 dx\, g_1{}^p(x) = \frac{1}{2}\left[\frac{4}{9}\Delta u' + \frac{1}{9}\Delta d' + \frac{1}{9}\Delta s'\right]\left[1 - \frac{\alpha_s}{\pi}\right] \tag{2}$$

with

$$\langle p,s|\,\bar{q}\gamma_\mu\gamma_5 q\,|p,s\rangle = \Delta q' s_\mu \tag{3}$$

where s_μ is the covariant spin vector of the proton. In the simple parton model one would naturally identify $\Delta q'$ with the polarized quark density.

The baryon matrix elements of the axial vector current also appear in the neutron and hyperon semileptonic decays. They are related by flavor SU(3), and are parameterized in terms of the F and D coefficients:

$$\begin{aligned} \Delta u' - \Delta d' &= g_A = F + D \\ \Delta u' + \Delta d' - 2\Delta s' &= 3F - D \end{aligned} \tag{4}$$

Thus, using the isospin relations $(\Delta u')^p = (\Delta d')^n$ & $(\Delta d')^p = (\Delta u')^n$, one gets the

Bjorken sum rule [2] for the difference of the proton and the neutron first moments:

$$\int_0^1 dx\, (g_1^p(x) - g_1^n(x)) = \frac{1}{6} g_A \tag{5}$$

For the proton moment alone, because there is an SU(3) singlet piece, the r.h.s. cannot be completely expressed in terms of F and D parameters. Here we choose to write the remainder in terms of $\Delta s'$,

$$\int_0^1 dx\, g_1^p(x) = \frac{1}{18}(9F - D + 6\Delta s')\left[1 - \frac{\alpha_s}{\pi}\right] \tag{6}$$

Since the proton does not have strange valence quark, it seems reasonable to conjecture that $\Delta s' \cong 0$, expressing a kind of OZI rule [4]. This yields the Ellis-Jaffe sum rule [5]:

$$\int_0^1 dx\, g_1(x) = 0.171 \tag{7}$$

This has become the benchmark for experimenters and phenomenologists alike. What is considered so surprising about the EMC result is that it deviates significantly from this baseline expectation.

$$\int_0^1 dx\, g_1(x) = 0.114 \pm 0.012 \pm 0.026$$

$$= 0.126 \pm 0.010 \pm 0.015 \tag{8}$$

Here we have given both the published EMC numbers [1] as well as the result of a more recent reanalysis [6]. (This should give us a better idea as to the range of variability of the result.)

Using SU(3) to express the EMC result in terms of the Δq's, we have

$$\Delta u' = 0.74 \qquad\qquad 0.77$$
$$\Delta d' = -0.53 \qquad\qquad -0.49$$
$$\Delta s' = -0.20 \qquad\qquad -0.16$$

which sum up to

$$\Delta\Sigma' = 0.01 \qquad\qquad 0.12 \tag{9}$$

The first column is the published result, and the second from the 'new' analysis. Here we have used $F = 0.47$ & $D = 0.81$, which are the best-fit values of all the neutron and hyperon β decay measurements [7, 8]. This represents a significant downward revision, from 0.63 down to 0.58, of the $\frac{F}{D}$ ratio widely used up until recently [9]. We also note that, because not all measured F & D parameters are consistent with each other, by emphasizing any particular set of data the resultant Δq's can vary somewhat [10]. However we expect that the uncertainties not to exceed twenty percent, and the conclusions, that $\Delta s'$ is non-negligible and $\Delta\Sigma'$ is suppressed, will stand.

It has also been pointed out [7] that the EMC result is consistent with the existing data on low energy elastic $\nu p \rightarrow \nu p$ scattering [11], which is sensitive to a different combination of $\Delta u'$, $\Delta d'$, and $\Delta s'$.

The results of a significant $\Delta s'$ and a suppressed $\Delta\Sigma'$ seem to deviate from the naive quark model expectation, which would predict a negligible contribution from the strange quark term $\Delta s' \cong 0$ and it would have the proton spin given completely by the valence quarks $\Delta\Sigma' \cong \Delta u' + \Delta d' = 1$. Of course such an expectation is perhaps too simple, since we know from the deep inelastic scattering (DIS) momentum sum rule that the gluons carry about half of the proton momentum. Therefore it is likely that the situation will change considerably when the gluon degrees of freedom are introduced: while gluons do not contribute to the additive quantum numbers like electric charge, and strangeness, etc. there is no reason to expect the gluon spin and orbital angular momentum not to be an important factor in the discussion of the proton spin. Still, because of the simplicilty of the nonrelativistic quark model (and many of its successes), the EMC result of $\Delta s'$ & $\Delta\Sigma'$ had people calling it *a proton spin crisis*. Several different interpretations of the EMC data have been advanced. Even a breakdown of perturbative QCD has been suggested [12]. (For arguments against such an interpretation, see Refs. [10, 13].) We list below some of the less radical explanations:

(1) Data not yet in the scaling region ?

It has been suggested that perhaps the Q^2 values reached in the EMC experiment are still not large enough to be in the scaling region [14]: one observes that the $Q^2 = 0$ limit of the first moment is fixed by the Drell-Hearn-Gerasimov sum rule, and this implies a very different value for the r.h.s. of the integral (in fact a sign-change). However, this is not a viable option, since data over a sizable range of Q^2 exist and no appreciable Q^2-dependence is visible. Actually it is remarkable that the scaling curve represents, in the sense of Bloom and Gilman [15], a very good average over local fluctuations down to $Q^2 \cong 0.5\ GeV^2$. (And, one can even understand the change of sign in the DHG sum rule as reflecting the N^* dominance in the low Q^2 region [16].)

(2) Unreliable extrapolation to x = 0 ?

Although EMC measurements had reached, for significant Q^2's, very small values of x, to obtain the moment integral one still has to extrapolate down to $x = 0$. The small $x \equiv \frac{Q^2}{2M\nu}$ limit corresponds to the high energy, $\nu \to \infty$, limit of the virtual γp Compton scattering. The EMC extrapolation is quite consistent with the standard Regge expectation [13]:

$$g_1(x) \cong \sum_i \beta_i^{\gamma} \beta_i^{p} x^{-\alpha_i(0)} \tag{10}$$

Here one sums over the Regge trajectory contributions from $A_1(1270)$, $f_1(1285)$, $f_1(1420)$... Their intercepts $\alpha_i(0)$ are expected to be small and thus $g_1(x)$ should be regular in the $x = 0$ limit. Clearly if this expectation turns out to be incorrect and there is some sort of singular behavior, then the resulting moment integral will be significantly changed [17]. However there is not a hint of such a singular behavior in the data, and theoretically the Pomeron-Pomeron-cut $\frac{1}{x \ln x}$ behavior is not expected here either.

(3) Large SU(3) breaking ?

Unlike the vector charges which are protected from large symmetry breaking corrections by the Ademollo-Gatto theorem [18], the axial vector charges may not obey SU(3) relations all that well. Again there is no such way out. Several hyperon beta decays have been measured, and one finds that SU(3) symmetry actually works remarkably well. In fact the uncertainties in the *F & D* matrix elements result mostly from experimental systematic errors [8]. Also one can explicitly show that the baryon axial charges are not sensitive to large SU(3) breakings in the baryon wave functions [19].

(4) Gluonic contribution via axial anomaly ?

Perhaps the most interesting suggestion for revising the straightforward interpretation of the EMC result is that, because of the presence of axial anomaly [20, 21], $\Delta q'$ not only represents the quark contribution to the proton spin but has a significant fraction coming from the gluon polarization as well,

$$\Delta q' = \Delta q - \frac{\alpha_S}{2\pi}\Delta g \tag{11}$$

where Δq is the 'true' quark contribution, and is idendified with the spin-dependent distribution as

$$\Delta q = \int_0^1 dx\ [q^+(x) - q^-(x) + \bar{q}^+(x) - \bar{q}^-(x)] \tag{11a}$$

and, correspondingly, Δg is the polarized gluonic distribution:

$$\Delta g = \int_0^1 dx[g^+(x) - g^-(x)]. \tag{11b}$$

(The superscripts ± indicate whether the parton helicity having the same or opposite signs to the target proton helicity.) The remarkable claim is that one can in a perturbative QCD calculation single out a subset of gluon reducible diagrams with a small coefficient proportional to $\alpha_S(Q^2)$. This possibility was first raised by Efremov and Teryaev [22], and then Altarelli and Ross [23] correctly derived the above spin density relation. Thorough discussions have also been given by Carlitz, Collins, and Mueller [24, 25]. Further discussion as to the validity of this relation has been given in [8, 26]. We shall suggest [27] that current divergence equations and decoupling theorem can provide another perspective for clarifying some of the perplexities in this problem.

II. Gluonic contribution via axial anomaly

Let us recall some elementary aspects about the QCD axial anomaly. The gauged color-SU(3) is of course anomaly-free. The anomaly under discussion is the one associated with the global axial U(1) symmetry. Namely, because the triangle diagram of the (quark) axial vector current with two gluon legs is linearly divergent and any cut-off scheme will either break the vector or the axial-vector symmetry, the regularization scheme that keeps the vector current conservation necessarily leads to the anomaly term in the divergence of the axial vector current [20]:

$$\partial^{\mu}\left[\sum_{i}\bar{q}\gamma_{\mu}\gamma_{5}q\right] = \left[\sum_{i}2m_{i}\bar{q}i\gamma_{5}q\right] + n_{f}\frac{\alpha_{s}}{2\pi}TrG^{\mu\nu}\tilde{G}_{\mu\nu} \tag{12}$$

where $G^{\mu\nu}$ is the gluon field tensor, $\tilde{G}_{\mu\nu}$ its dual. $n_f = 3$ is the number of excited flavors, and $i = u, d, s$. For our purpose it is more convenient to think all this in terms of each flavor separately. Thus, we have

$$\partial^{\mu}\bar{q}\gamma_{\mu}q = 0$$

$$\partial^{\mu}\bar{q}\gamma_{\mu}\gamma_{5}q = 2m_{q}\bar{q}i\gamma_{5}q + \frac{\alpha_{s}}{2\pi}TrG\tilde{G} \tag{13}$$

The key point to keep in mind is that anomaly comes from regulating the ultraviolet divergence.

Now let us see how does the axial anomaly enter into the discussion of partonic contributions to the proton spin. The current correlation function as measured in the inelastic eletroproduction in the scaling limit (specifically the first moment of the spin dependent structure function) is related, through the operator product expansion, to the proton matrix element of the local axial vector current. In the parton model, this can be evaluated by taking the axial vector current between the quark states and multiplying it by the probability of finding the quark in the target proton. Or, equivalently, this is the matrix element $\langle q|j_5|q\rangle$ for the quark state with helicity aligned to the proton helicity times the spin-dependent quark distribution Δq as given in Eq. (11a). What AR have shown [23], and is corroborated later by more detailed discussion by CCM [24], is that there is another potentially important contribution to this proton matrix element of the axial vector current: it is the spin-dependent gluon distribution Δg times the gluonic matrix element of the axial current $\langle g|j_5|g\rangle$. The last factor is just the anomaly triangle diagram which, with $\langle q|j_5|q\rangle$ normalized to one, yields a coefficient of $-\frac{\alpha_S}{2\pi}$. Thus, we have schematically:

$$\langle p|j_5|p\rangle = \langle q|j_5|q\rangle\Delta q + \langle g|j_5|g\rangle\Delta g \tag{14}$$

or

$$\Delta q' = \Delta q - \frac{\alpha_S}{2\pi}\Delta g$$

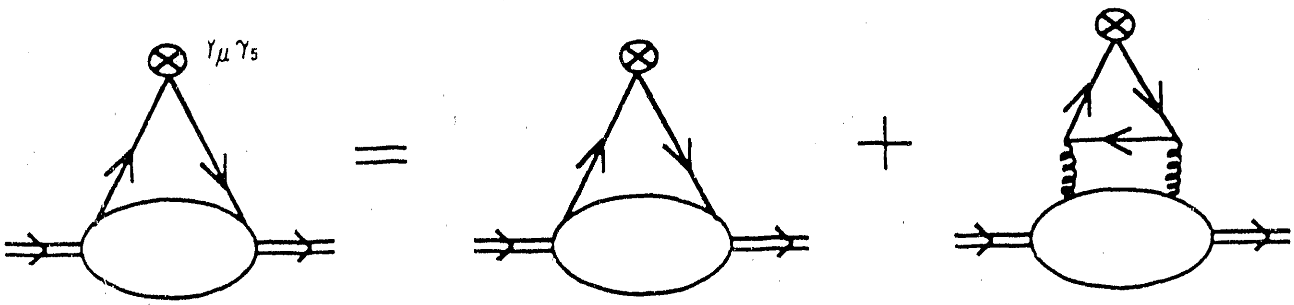

Fig. 1

Let us emphasize that the operator product expansion approach to DIS itself is not under challenge. What is under discussion is the contribution to the proton matrix element of the local operator of quark axial vector current. We note that in the naive parton model calculations the first term on the r.h.s. corresponds to the tree diagram and the second to the box diagram.

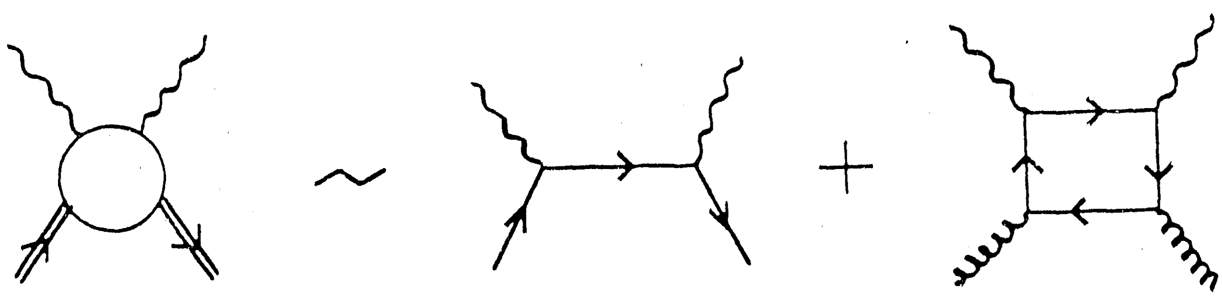

Fig. 2

Immediately one has to face the following two questions:

(1) Since this is a higher order contribution, why should it be important?

(2) If it comes from the gluonic partons, one needs to show that it's a point-like contribution. Namely, one is obliged to show that, in the matrix element of the local axial vector current operator (which is generally a long-distance-physics quantity), it is possible to identify a hard & perturbatively calculable contribution from the gluon.

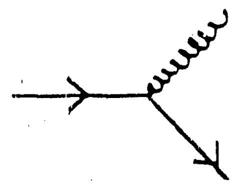

Fig. 3

(1) $\alpha_S \Delta g$ can be large

The QCD correction term $\alpha_S \Delta g$ in $\Delta q'$ can be large because Δg grows with Q^2. It is easy to see that helicity is not conserved in a tree diagram representing gluon bremsstrahlung by a quark. While quark helicity does not change, there is a gain of the gluon helicity. Of course what is conserved is the total angular momentum, for which one has to include the orbital contribution [28]. As it turns out for large Q^2, Δg goes as lnQ^2 [29], and, since the strong coupling $\alpha_S \propto \frac{1}{lnQ^2}$, the product $\alpha_S \Delta g$ is Q^2-independent at the leading order. In effect, one can regard $\alpha_S \Delta g$ as being of the zeroth order, and thus potentially it can be substantial.

We can state this more precisely in terms of the QCD evolution equation of Altarelli & Parisi. The evolution of the moment of the structure function is controlled by the moment of the 'splitting function' which can be calculated from QCD vertices. For $t \equiv lnQ^2$, we have [23]

$$\frac{d}{dt}\begin{bmatrix}\Delta\Sigma \\ \Delta g\end{bmatrix} = \alpha_S(t)\begin{bmatrix}0 & 0 \\ A_{qg} & b\end{bmatrix}\begin{bmatrix}\Delta\Sigma \\ \Delta g\end{bmatrix} \qquad (15)$$

The zero entries reflect the Q^2 independent nature of the quark helicity distribution $\Delta\Sigma \equiv \sum_q \Delta q$. The key point is that the (g,g) entry in the evolution matrix is proportional to the leading coefficient in the QCD β-function: $\frac{d\alpha_S}{dt} = -b\alpha_S{}^2$. Because to this order $\Delta\Sigma$ is t-independent, thus we can ignore the $\alpha_S(t)\Delta\Sigma$ term for large Q^2, and obtain

$$\frac{d}{dt}\Delta g(t) \cong b\alpha_s(t)\Delta g(t)$$

This implies that
$$\alpha_s \frac{d\Delta g}{dt} + \frac{d\alpha_s}{dt}\Delta g = \frac{d}{dt}(\alpha_s \Delta g) \cong 0 \tag{16}$$

(2) Anomaly as a local probe of the gluon distribution

In the parton model language, one can represent the probing of the *quark* distribution by the tree diagram. The phenomenological manifestation is the 1-jet process corresponding to a hard quark. As for the gluon, we have box diagrams, *Fig. 2.* Since this involves a integration over all values of momentum it generally does not represent a pointlike probe. However the finite result $-\frac{\alpha_s}{2\pi}$ extracted above in the perturbation calculation comes from the high $k_T^2 \cong Q^2$ region of the loop integration [24]. One way to see this is as follows: For the massless quark ($m = 0$) & off-shell gluon ($p^2 \neq 0$), the box-graph contribution to the moment sum of $g_1(x)$ is given in Eq. (17). The x-integral on the r.h.s. vanishes for small values of k_T because in that limit the integrand is odd under the interchange of $x \to 1-x$.

$$\int_0^1 dx g_1(x,Q^2) = -\frac{\alpha_s}{2\pi}\int_0^1 dx \int_0^{K^2} \frac{d(k_T^2)}{\sqrt{1-\frac{(k_T^2)}{K^2}}} \left\{ \frac{(1-2x)\frac{(k_T^2)}{2}}{[k_T^2-p^2x(1-x)]^2} - \frac{x(1-2x)}{Q^2(1-x)} \right\} \tag{17}$$

where $K^2 \equiv \frac{Q^2(1-x)}{4x}$. Because only the high k_T region contributes, this loop graph is effectively a pointlike coupling. Furthermore because the quark lines are hard the phenomenological signal of gluon parton contribution will be 2-jet processes.

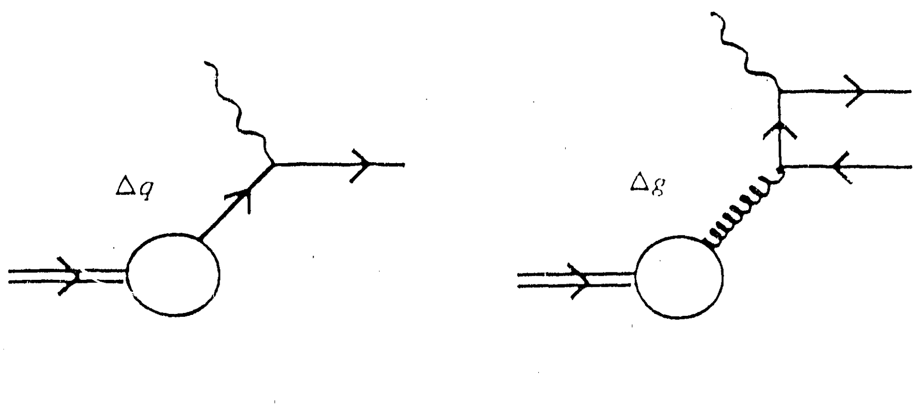

Fig. 4

Also, had the low momentum scales contributed, the running coupling would not be small and the perturbative approach could not be trusted. Fortunately since the present situation only involves hard quark lines, we do have a small coupling $\alpha_S(Q^2)$, as guaranteed by asymptotic freedom, and the relevant perturbation calculation is reliable.

In the operator product expansion language, the above box graph generates in the scaling limit the triangle diagram of the axial vector current matrix element. And the above discussion is compatible with our understanding that anomaly comes from the ultraviolet behavior of the loop integration. Again this means that one can view the hard quark triangle loop as effetively pointlike. Thus we have a rather subtle situation with respect to OPE in this problem: Usually the purpose of expressing the current product in terms of local operators is to have pointlike probes. In this problem, initially one would think that there was no pointlike gluon contribution to the first moment because there was no (color) gauge invariant twist-two gluon operator. But the above calculations suggest that, the gluon contribution via the anomaly being effectively pointlike, we can think of the local quark axial vector current operator $j_{5\mu}$, as also containing a gluon part:

$$j_{5\mu} = \tilde{j}_{5\mu} - \frac{\alpha_S}{2\pi} k_\mu \tag{18}$$

where k_μ is a twist-two local gluonic operator.

This possibility has long been discussed in the context of the anomaly equation. The anomalous divergence itself can be written as the four-divergence of a gluonic current:

$$\partial^\mu k_\mu = -Tr G\tilde{G},$$

with

$$k_\mu = -\epsilon_{\mu\nu\lambda\sigma} Tr\left[A^\nu G^{\lambda\sigma} - \frac{2}{3} A^\nu A^\lambda A^\sigma\right] \tag{19}$$

Thus the decomposition of Eq. (18) is possible, $\tilde{j}_{5\mu}$ being conserved in the chiral limit. Indeed it is tempting to suggests that Eq. (1) can be trivially obtained from the current decompostion (18) if the forward matrix elements of the currents k_μ & $\tilde{j}_{5\mu}$ are to be idendified with Δg & Δq , just as the forward matrix element of the axial vector current $j_{5\mu}$ defines $\Delta q'$ in Eq. (3). If this could be done, we would have then simple operator matrix element expressions of the spin-dependent parton density functions Δq & Δg. But the topological current k_μ is not (color) gauge invariant. (That is why it does not appear explicitly as an independent local operator in OPE and it must enter together with $\tilde{j}_{5\mu}$ to maintain gauge invariance.) Related to this, the matrix elements of k_μ & $\tilde{j}_{5\mu}$ have massless gnost-pole terms and the corresponding zero momentum transfer limits are singular. Thus it seems that we cannot go beyond the partonic definitions of Δq & Δg as given in Eqs. (14, 11a,b).

To be more specific we write out the nonforward matrix elements, with the momentum difference of the final and initial proton states being $t_\mu \equiv p'_\mu - p_\mu$, as follows:

$$<p',s|\, j_{5\mu}\,|\, p,s> = \bar{u}(p')[\gamma_\mu\gamma_5 G_1(t^2) - t_\mu\gamma_5 G_2(t^2)]u(p) \tag{21a}$$

$$<p',s|\, \tilde{j}_{5\mu}\,|\, p,s> = \bar{u}(p')\left[\gamma_\mu\gamma_5 \tilde{G}_1(t^2) - t_\mu\gamma_5 \tilde{G}_2(t^2)\right]u(p) \tag{21b}$$

$$<p',s|\, k_\mu\,|\, p,s> = \bar{u}(p')[\gamma_\mu\gamma_5 H_1(t^2) - t_\mu\gamma_5 H_2(t^2)]u(p) \tag{21c}$$

In the forward limit of $t_\mu = 0$ there are massless ghost-pole terms in the gauge variant current matrix elements $\tilde{G}_2(t^2)$ and $H_2(t^2)$:

$$\tilde{G}_2(t^2) = \frac{2MR_q}{t^2} \quad and \quad H_2(t^2) = \frac{2MR_g}{t^2} \tag{22}$$

(M is the proton mass.) However, because the matrix element of the gauge invariant current j_5 is nonsingular (even in the chiral limit there is no massless Goldstone boson in the singlet channel*), the singular parts on the r.h.s. of Eq. (18) must cancel:

$$R_q - \frac{\alpha_s}{2\pi} R_g = 0 \tag{23}$$

The forward matrix elements of Eq. (18) are then related:

$$\Delta q' = \tilde{G}_1(0) - \frac{\alpha_s}{2\pi} H_1(0) \tag{24}$$

where we have used $G_1(0) = \Delta q'$. Indeed it is tempting to recover Eq. (1) by making the idendifications of

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* The relevance of the 'axial U(1) problem' (η' gaining a mass through instanton interactions, *etc.*) has been discussed previously in Refs. [8, 30-34].

$$\Delta q = \tilde{G}_1(0) \quad and \quad \Delta g = H_1'(0). \tag{25}$$

However this is not warranted because we do not expect $H_1(0)$ & $\tilde{G}_1(0)$ to be invariant under large gauge transformations. Thus, for Δq & Δg, it is not clear how to go beyond the partonic definitions [as given in Eqs. (14, 11a,b)], to the local operator matrix element definitions. We shall return to this problem of expressing quark and gluon contributions in terms of operator matrix elements in Secs. III and V.

(3) Regulating the infrared collinear divergence

In the chiral limit ($m = 0$) with the gluon on-shell ($p^2 = 0$), the box- and triangle-graphs have infrared collinear divergences. The problem of a proper regularization has been studied in detail by CCM [24]. They note that the final result (symbolically represented by Δ_5 below) depends sensitively on the form of the IR regulation:

$$\Delta_5 = -\frac{\alpha_S}{2\pi} \quad for \quad \frac{m^2}{p^2} \to 0 \tag{26a}$$

$$\Delta_5 = 0 \quad for \quad \frac{p^2}{m^2} \to 0 \tag{26b}$$

CCM point out that the first limit $\frac{m^2}{p^2} \to 0$ is the appropriate one: gluons inside a confined hadron is expected to be off-shell by an amount characterized by the QCD confinement scale $p^2 \cong \Lambda^2$ of a few hundred MeV^2, which is large compared to the u, d, s quark masses. Furthermore the result $\Delta_5 = -\frac{\alpha_S}{2\pi}$ comes from high k_T region of the loop integration, in agreement with our expectation that the anomaly results from UV regularization of the triangle graph.

The zero of the second limit $\frac{p^2}{m^2} \to 0$ results from a cancellation of this high k_T term by an equal and opposite low k_T contribution. As we have already pointed out, the high k_T part corresponds to hard quark lines and thus a 2-jet process. It is compatible with our expectation of a scattering off the gluon parton. The low k_T contribution to the scatterings does not produce 2-jet events. Thus this factor should really be counted as (higher order) part of Δq. In short, the second form of IR regularization is not the physically relevant one.

For different viewpoints with regard to this question of infrared sensitivity, see Refs. [8, 26].

III. The perspective from current divergence equations & decoupling theorem

(1) Axial anomaly equation and the proton spin

The problems of gauge invariance and infrared regularization prompt us [27] to suggest that a useful approach to the problem of gluonic contribution via the axial anomaly is to work directly with current divergences. (For previous discussions of the axial U(1) Ward idenditiy in connection with the proton spin problem, see Refs. [8, 30-35].) The current divergences are expected to be less IR sensitive, and gauge invariance is also manifest.

For the gauge invariant axial vector current $j_{5\mu}$, becuase there is no massless Goldstone boson even in the chiral limit, it is trivial to go from the forward matrix element of the current to that of the divergence: Denoting the proton (axial) spin-vector and pseudoscalar density by $s_\mu \equiv \bar{u}\gamma_\mu\gamma_5 u$ *and* $v \equiv \bar{u}i\gamma_5 u$, respectively, we have

$$\langle p,s|\, j_{5\mu}\,|\, p,s\rangle = s_\mu \Delta q' \quad \rightarrow \quad \langle \partial j_5\rangle = 2Mv\Delta q' \tag{27}$$

For the gauge variant currents $\tilde{j}_{5\mu}$ and k_μ, even though the current matrix elements are singular, those of the divergences are regular and manifestly gauge invariant:

$$\langle \partial \tilde{j}_5\rangle = \langle p,s|\, 2m\bar{q}i\gamma_5 q\,|\, p,s\rangle \equiv 2Mv\Delta Q \tag{28a}$$

$$\langle \partial k\rangle = \langle p,s|\, -G\tilde{G}\,|\, p,s\rangle \equiv 2Mv\Delta G \tag{28b}$$

From the expressions in Eq. (21) and the definitions in Eq. (22), we have

$$\Delta Q = \tilde{G}_1(0) - R_q \quad \textit{and} \quad \Delta G = H_1(0) - R_g . \tag{29}$$

Namely, the (large) gauge non-invariant parts on the r.h.s. cancel to yield gauge invariant ΔQ & ΔG. Taking the forward proton matrix elements of the familiar anomaly equation (13):

$$\partial j_5 = 2m\bar{q}i\gamma_5 q + \frac{\alpha_s}{2\pi} Tr G\tilde{G}$$

we obtain a relation very similar to that of Eq. (1):

$$\Delta q' = \Delta Q - \frac{\alpha_s}{2\pi}\Delta G. \qquad (1a)$$

Given the expressions (28) for ΔQ & ΔG in terms of the quark and gluon fields, we can naturally interprete Eq. (1a) as another way of separating the quark and gluon contributions to the matrix element of the axial vector current. In particular, the interpretation of ΔQ & Δq as the "intrinsic quark contributions" is justified as in the free-field limit we have:

$$(\Delta q')_0 \quad (\Delta Q)_0 = (\Delta q)_0$$

Because Eqs. (1) and (1a) have similar quark-gluon decompositions, we can view their relation as follows: Eq. (1a) is just the Eq. (1) with certain constant being shifted from the first to the second term on the r.h.s.*

$$\sigma_q - \frac{\alpha_s}{2\pi}\sigma_g = 0,$$

with
$$\sigma_q = \Delta q - \Delta Q, \qquad \sigma_g = \Delta g - \Delta G \qquad (30)$$

Since the σ's in the quark and gluon factors should always cancel, it really does not matter which set, $(\Delta q,\ \Delta g)$ or $(\Delta Q,\ \Delta G)$, one uses in the discussion of the $\Delta q'$ properties. (Of course, this is not the case when one wishes to discuss the separate quark or gluon contributions.) In this sense, one can look upon Eq. (1) as being the forward matrix element of the axial anomaly equation. The main advantage of Δq & Δg is their direct parton interpretations, while that of ΔQ & ΔG is their simple expression in terms of the (manifestly gauge invariant) field operators.

One of the advantages of having the explicit operator representations of (28) is that it allows us to see more clearly the source of 'IR regulation ambiguity' discussed in Sec. II.(3). Because we now view the relation Eq. (1) as the forward proton matrix elements of the anomaly equation Eq. (13), it is straightforward to see how a heavy quark ($q = c,\ b,\ t$) contribution to the proton spin 'decouples'. According to the Appelquist-Carrazzone theorem [36] $\Delta q'$ for heavy quarks should be suppressed by powers of the quark masses. However from our Eq. (1a) for $\Delta q'$ we see that ΔQ does *not* vanish in the $m \to \infty$ limit, because when we integrate out the heavy quark

* G. Altarelli, private communication.

field the pseudoscalar density for the heavy quark gives rise to a dimension-four operator:

$$m\bar{q}i\gamma_5 q \rightarrow c\ TrG\tilde{G} + O(m^{-2}) \tag{31}$$

The relevant diagram is of course the pseudoscalar triangle diagram with two gluon legs. This yields the coefficient $c = -\frac{\alpha_s}{4\pi}$. Consequently, to the leading order of this 'heavy quark expansion' [37], the quark and gluon pseudoscalar density terms cancel. Thus, the gluonic contribution via anomaly is needed to enforce the physical requirement of decoupling of heavy quarks.

It is easy to see how this cancellation comes about. Recall that anomaly term results from regulating the UV divergence. We can think that this is accomplished by introducing a 'regulating fermion' with a large mass. This of course breaks the chiral symmetry, and it gives rise to a mass term in the divergence of the axial vector current. Integrating out this heavy fermion we can then transform this naive divergence term into the anomaly factor. In every step it contributes the same as our original heavy quark. The difference is that for the regulating fermion loop there is an extra minus sign. This accounts for the cancellation.

Given our previous comment on the equivalence of $(\Delta Q, \Delta G)$ and $(\Delta q, \Delta g)$ in any calculation of $\Delta q'$, this cancellation due to decoupling is the same one encountered above in the discussion of the IR collinear regularization, Eq. (26). This shows once again that the $\frac{p^2}{m^2} \rightarrow 0$ limit is not the appropriate one for the u, d, s quarks. (They are not heavy quarks on the QCD confinement scale.)

Viewed in this perspective, it should also be clear that there is no reason to expect that the strange quark fraction $\Delta s'$ be suppressed. Although $m_s >> m_u, m_d$ it is still not a heavy quark, and there is no reason to expect that the approximation of Eq. (31) is valid for the strange quark. (After all, treating the s quark as a light quark in chiral SU(3) and kaon PCAC does lead to reasonablly good results.) In this sense the EMC result finding $\Delta s'$ to be comparable to $\Delta u'$, $\Delta d'$ should not really be looked upon as being so puzzling. This is certainly not a crisis.

(2) Trace anomaly equation and the proton mass

Given the new perspectives of current divergence equations and decoupling theorem, it is natural to make a comparison of another static property of the proton: its mass. Here the relevant current is the dilation current with a divergence given by the trace of the energy momentum tensor.

$$\partial^\mu D_\mu = \theta^\mu{}_\mu = \Sigma m_i \bar{q}_i q_i - \left[11 - \frac{2}{3}n_f\right]\frac{\alpha_s}{8\pi}TrG^{\mu\nu}G_{\mu\nu} \tag{32a}$$

Again, the naive divergence is due to the quark masses, and there is also an anomaly term, the *trace anomaly* [38].* When taken between the proton states, this divergence equation yields

$$M = <m_u \bar{u}u + m_d \bar{d}d + m_s \bar{s}s> + \sum_{heavies} <m_i \bar{q}_i q_i> - \left[9 - \frac{2}{3}n_h\right]\frac{\alpha_s}{8\pi}<TrGG> \tag{32b}$$

where n_h is the number of heavy quarks. Again the heavy quarks decouple through the cancellation of the quark mass terms and their contributions to the anomaly. As for the surviving term, it has been tempting to say that the $<m\bar{q}q>$ terms are negligible because m_u, m_d, & m_s are small. Thus one ends up with the proton mass being almost all given by the gluon contribution. This view has been advocated by Shifman, Vainshtein, and Zakharov [39].

Actually there is experimental information on the quark contributions. The u, d term is measured by the πN sigma term [ignoring a small correction of the form $(m_u - m_d)<N|\bar{u}u - \bar{d}d|N>$]

$$<N|m_u \bar{u}u + m_d \bar{d}d|N> \cong 60\ MeV. \tag{33}$$

And just as we can deduce the strange quark contribution from the EMC result, here one can use SU(3), the octet baryon mass differences, and the canonical quark mass ratio $\frac{2m_s}{m_u + m_d} \cong 25$ to get the fraction

$$\frac{2<N|\bar{s}s|N>}{<N|\bar{u}u + \bar{d}d|N>} \cong 0.47. \tag{34}$$

* Trace anomaly equation has also been discussed in Ref. [8] albeit for the purpose of illustrating the 'ambiguity' of idendifying a gluonic contribution (see our comments in Sec. V.(2)). For another discussion, see also Ref. [40].

This is the so called '$\sigma_{\pi N}$ problem' [41]: nucleon seems to have surprisingly significant amount of the strange quark sea. But from the perspective of the decoupling theorem it really is not so odd. Incidently with such interpretation of the $\sigma_{\pi N}$ value, one finds that the proton mass is about equally shared between the quarks and the gluons [42].

IV. Estimating ΔG & ΔQ

In the context of divergence equation, we can use the standard current algebra approach of Goldstone pole saturation to estimate ΔG and ΔQ [35] as defined in Eq. (28). The basic idea is a simple one: In the divergence equations,

$$\Delta u' = \frac{m_u}{M}\nu_u - \frac{\alpha_s}{2\pi}\Delta G \tag{35}$$

$$\Delta d' = \frac{m_d}{M}\nu_d - \frac{\alpha_s}{2\pi}\Delta G \tag{36}$$

(where we have written out ΔQ in terms of the quark masses and the pseudoscalar density matrix elements $\nu_q \bar{u} i\gamma_5 u = \langle p|\bar{q}i\gamma_5 q|p\rangle$) one can solve for ΔG and $\nu_u + \nu_d$ (the "unknowns") if we regard as known quantities: $\Delta u'$, $\Delta d'$ (as given *via* the EMC result) and $\nu_u - \nu_d$ (*via* saturation of the nonsinglet channel by the Goldstone poles). Let us first recall that the Goldberger-Treiman relation [43] can be derived in the charged channel by the $\pi^\pm$ pole-dominance of the pseudoscalar density. After taking the nucleon matrix element of the divergence equation,

$$\partial^\mu(\bar{u}\gamma_\mu\gamma_5 d) = (m_u + m_d)\bar{u}i\gamma_5 d \tag{37}$$

one obtains

$$2Mg_A = 2f_\pi g_{\pi NN} + \mu_\pm \tag{38}$$

where $\mu_\pm$ denotes the correction to the $\pi^\pm$ pole-dominance, and is the correction to g_A as given by the Goldberger-Treiman relation. Repeating the same for the neutral isovector channel,

$$\partial^\mu\left[\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d\right] = 2m_u\bar{u}i\gamma_5 u - 2m_d\bar{d}i\gamma_5 d \tag{39}$$

we have

$$2Mg_A = 2f_\pi g_{\pi NN} + \mu_0 + (m_u - m_d)(v_u + v_d) \tag{40}$$

Comparing these two expressions for g_A one concludes that the singlet density is small, on the order of *correction* to the GT expression for g_A . If we assume that μ_0 correction to the π^0 PCAC is negligible $\mu_0 \cong 0$, then the isosinglet density is simply given by δg_A, and using the 'new' EMC numbers one gets:

$$\frac{\alpha_s}{2\pi}(\Delta G)_p \cong -0.10 \quad and \quad \Sigma(\Delta Q)_p \cong -0.18 \tag{41}$$

If we assume $\mu_0 = \mu_\pm$, the isosinglet pseudoscalar density vanishes [41], and this leads to

$$\frac{\alpha_s}{2\pi}(\Delta G)_p \cong -0.23 \quad and \quad \Sigma(\Delta Q)_p \cong -0.57 \tag{42}$$

If one repeats the same calculation for the neutron, one finds a significantly different result (for $\mu_0 = 0 \sim \mu_\pm$):

$$\frac{\alpha_s}{2\pi}(\Delta G)_n \cong +0.26 \sim 0.13 \quad and \quad \Sigma(\Delta Q)_n \cong +0.90 \sim +0.51 \tag{43}$$

Thus although the operator $Tr G\tilde{G}$ is manifestly iso-symmetric, its nucleon matrix elements violate isospin symmetry strongly.

If one goes back through the calculation, one finds that the isospin violation comes in the form of

$$\left[\frac{m_u - m_d}{m_u + m_d}\right](\Delta u' - \Delta d'). \tag{44}$$

It is not difficult to see what this represents: the $(\Delta u' - \Delta d')$ factor is the isovector π^0 axial vector coupling to the nucleon, and it gives opposite sign contributions to $(\Delta G)_p$ and $(\Delta G)_n$, respectively. This term comes about as follows: Because we are *not* setting $m_u = m_d$, the isospin violating term in the QCD Hamiltonian causes a $\pi^0\eta$ mixing. This brings about the π^0 pole in the matrix element of the isosinglet operators, with a mixing angle $\sim (m_u - m_d)$. Such a pion pole contribution is enhanced by the smallness of the pion mass, as indicated by the $(m_u + m_d)$ factor in the denominator.

With such a large isospin violation one should worry that this is not compatible with our original assumption of flavor SU(3) symmetry in extracting $\Delta q'$ from the EMC data. But, it turns out that the pion pole contributions (*i.e.* the large isospin violations) in the gluon ΔG and the quark mass terms ΔQ cancel out. The axial vector matrix elements $\Delta u'$ and $\Delta d'$ respect isospin symmetry to a high degree. All this was worked out long ago by Gross, Treiman & Wilczek [45] who pointed out this cancellation as a consequence of the Sutherland theorem [46]. (The Sutherland theorem is usually used in connection with $\pi^0 \to \gamma\gamma$, but here it comes in connection with the strong anomaly.)

Hatsuda [32] using an effective Lagrangian approach showed that to the leading order in the $\frac{1}{N}$ expansion, ΔQ & ΔG are given in terms of the η', η, & π^0 pole contributions. Again they mostly cancel in the combination of Eq. (1a) and thus the sum of the axial vector current matrix elements $\Delta\Sigma'$ is given by the η' nucleon coupling. Thus the EMC result of Eq. (9) implies that η' essentially decouples from the nucleon. Veneziano [33] has argued that this relation between $\Delta\Sigma'$ and $g_{\eta' NN}$ is independent of the $\frac{1}{N}$ expansion, and is a generalized Goldgerger-Treiman relation in the singlet channel. That $\Delta\Sigma' \cong 0$ and the decoupling of η' from the nucleon is first derived in the Skyrmion model by Brodsky, Ellis, and Karliner [47].

V. Summary & discussion

(1) Proton matrix element of the axial vector current

This is perhaps the least controversial part of the discussion on the problem of the proton spin contents:

(i) That the EMC experiment measures the proton matrix element of the axial vector current is not in serious doubt: the data on $g_1(x)$ are in the scaling region, and there is no good reason to believe that the extrapolation to $x = 0$ point is unreliable. Operator product expansion then relates the first moment of $g_1(x)$ to the matrix element of the axial vector current operator.

(ii) Given that in the measured axial charges the flavor SU(3) appears to be a good symmetry, the extraction of $\Delta q'$ (for the flavor q) from the EMC data should be dependable. In particular the result of a significant $\Delta s'$ should be taken seriously. As for the sum $\Delta u' + \Delta d' + \Delta s'$, it is most likely suppressed, perhaps not as strongly as indicated in the original analysis, see Eq. (9).

(iii) Since the connection of axial vector current matrix element to the measurement of spin asymmetry in the deep inelastic scattering is not in question. We have every reason to expect that the Bjorken sum rule, Eq. (5), to hold when the corresponding measurements on the neutron target are carried out in future experiments. (Namely, one should keep in mind that the discussion in Sec. IV of large isopin violation concerns the separate quark and gluon contributions to $\Delta q'$ rather than $\Delta q'$ itself, which enters the sum rule directly.)

(iv) By the same reasoning, at sufficient high Q^2 (above the heavy quark thresholds) the first moment of $g_1(x)$ will in principle receive contributions from $\Delta c'$, $\Delta b'$, & $\Delta t'$, respectively. However by our decoupling argument, independent of the specific quark-gluon decompositions, matrix elements of such heavy quark local operators should be suppressed by inverse powers of the heavy quark masses. The key question that determines whether a quark is heavy or not is the following one: what's the relevant scale for the proton matrix element of a local operator? Is it the QCD confinement scale ≅ 1 GeV? Or, is it the Q^2 of the virtual photon? One would think that the issue of the relative sizes of the quark mass vs. Q^2 has to do with the Wilson coefficients (with its flavor thresholds) multiplying the local operator. Once the heavy flavor threshold is passed, one then gets a clear-cut connection between the first moment and the matrix element of the heavy quark bilinear. If the relevant scale is the confinement radius, such matrix elements should all be suppressed,

$$\Delta c' \cong \Delta b' \cong \Delta t' \cong 0, \tag{45}$$

independent of the value Q^2 (besides it's sufficiently above the corresponding quark threshold). Several authors offer a different prediction on this point [16, 48]. See paragraph (3) below for further discussions.

(2) The quark and gluon contributions: variant perspectives

In this report we have focused our discussions on the quark and gluon contributions to $\Delta q'$. It has been pointed out that there are two closely related, yet different, decompositions of $\Delta q'$:

$$\Delta q' = \Delta q - \frac{\alpha_S}{2\pi}\Delta g \tag{46a}$$

$$\Delta q' = \Delta Q - \frac{\alpha_S}{2\pi}\Delta G \tag{46b}$$

where Δq and Δg are spin-dependent parton distributions, as given in Eq. (11), and ΔQ and ΔG are matrix elements of local operators bilinear in the quark and gluon fields, as given in Eq. (28). Thus we can regard the two equations (46 a,b) as being related by shifting a piece from the quark to the gluon term, as in the manner of Eq. (30):

$$\Delta Q = \Delta q - \sigma_q \qquad \Delta G = \Delta g - \sigma_g \tag{47}$$

with the σ's cancel in $\Delta q'$. In fact a plausible expectation is that $(\Delta q,\ \Delta g)$ corresponds to those parts of $(\Delta Q,\ \Delta G)$ that can be reached in perturbative QCD, while $(\sigma_q,\ \sigma_g)$ are purely non-perturbative (instantcn interactions, ghost poles, *etc.*). Efremov, Soffer, and Törnqvist [31] have pointed out that the result obtained by Veneziano [30] should really be interpreted as the ghost-pole residue in the chiral limit; we can then interprete this to mean:

$$\sum_q \sigma_q = n_f \frac{\alpha_s}{2\pi}\sigma_g = \frac{F_{\eta'}}{2M} g_{\eta' NN} \tag{48}$$

where the decay constant $F_{\eta'} \cong 287\ MeV$. Because ΔQ is explicitly proportional to the quark mass, it vanishes in the chiral limit, leading through Eq. (47) to:

$$\sum_q \Delta Q = 0, \quad or \quad \Delta u + \Delta d + \Delta s = \frac{F_{\eta'}}{2M} g_{\eta' NN}. \tag{49}$$

However one must be very careful in making such an estimate. As the calculation in Sec. VI showed that $\Sigma\Delta Q$ is not particularly small and it is probably dominated by the non-perturbative σ_q, therefore using such an approximation to determine $\Sigma\Delta q$ is not likely to be reliable (besides the difficulty of an uncertain $g_{\eta' NN}$).

We also comment that that if one does not set the quark masses to zero, chiral and isospin symmetry breaking will bring about mixings between η' with η and π^0 [32, 33]. The pion pole contribution is the source of the large isospin violation discussed in the last section. Thus this large isospin violation is associated with σ_q & σ_g, which cancel out in axial vector matrix elements $\Delta q'$, Eq. (30). This is again in agreement with the result of ref. [45], and with our expectation that there should not be such large isospin violation in the parton distributions Δq & Δg. However it does not mean that this isospin violation is unphysical because there are situations where it is ΔQ *or* ΔG that is the physical observable. For example ΔQ's enter directly in the pseudoscalar Higgs boson nucleon couplings [44, 49], ΔG can in principle contribute to the neutron electric dipole moment, etc.

We should emphasize that the existence of these two variant approaches Eqs. (46a,b) does not imply there are unavoidable ambiguities so that it is meaningless to talk about any separate quark and gluon contributions. (For such a viewpoint, see ref. [8].) Both decompositions are meaningful because in each case the quark and gluon terms can be physically defined. (It does imply that one has to be careful in distinguishing between these two types of spin contributions.*) Each decomposition emphasizes a particular aspects of the spin content problem; each definition has its own advantage: Clearly Δq & Δg emphasizes the parton aspect of the problem, while ΔQ & ΔG are directly given by matrix elements of gauge invariant local operators, which are related to a variety of physically measurable quantities.

(3) Divergence equations and the decoupling theorem

It has been suggested that one works directly with the divergences rather with the current, to obtain Eq. (46b), in which the quark and gluon terms ΔQ & ΔG can be expressed in terms of local gauge invariant operators. In certain situation this is of considerable advantage as is demonstrated in our discussion of the heavy quark contributions and the decoupling theorem.

Whether a particular flavor can be treated as a heavy quark depends whether the quark mass m is large compared to the relevant scale of the problem so that the 'heavy quark expansion' Eq. (31) is valid. For a general low energy hadronic phenomenon this is set by the QCD confinement scale for which the u, d, & s quarks are light, and the c, b, & t are heavy. However for local operators probed in the deep inelastic processes there is another scale besides the confinement scale of the target hadron, it is the virtual photon mass Q^2. The question arises whether the decoupling takes place only if $m^2 >> Q^2$. If this is the case, then there is the possibility that Eq. (45) may not hold. In fact it has been suggested by Altarelli and Stirling [16] that $\Delta c \cong \Delta s \cong 0$ and

* The failure to distinguish ΔG from Δg in our previous work [35] has caused some confusion.

$$\Delta c' \cong \Delta s' \cong -\frac{\alpha_s}{2\pi}\Delta g \tag{50}$$

Since there is still some uncertainty as to the relevant scale, it will be very interesting to examine the electroproduction data to see which result, Eqs. (45) or (50), is realized in nature for the charm quark contribution.

It is clear that there is no circumstance when the u & d quarks can be regarded as heavy. We have argued above that the approximate validity of chiral SU(3) x SU(3) symmetry implies that the strange quark should also be treated as a light quark.

(4) OZI rule is not generally applicable for the strange quark

The decoupling theorem suggests that there is no *a priori* reason to expect that $\Delta s'$ should be particularly small. In fact it suggests that OZI rule should not generally be applicable to the strange quark,* because strange quark is not a heavy quark on the QCD confinement scale. The nucleon matrix elements of strange operators generally should not be suppressed. (For recent discussions of such matrix elements, see refs. [7, 8].) Of course this does not mean that there is no physical instance in which the nonstrange hadron matrix element of some strange quark bilinear is suppressed. However, only for the truly heavy quarks c, b, t, OZI rule is generally applicable, (*i.e.* it is really a *rule)* because of asymptotic freedom and the decoupling theorem. To put it in another way, what we are suggesting is that for the case of strange quark bilinears there should be a "new baseline of expectations": generally we should expects their matrix elements to be comparable to those of the u, d operators, *except* for cases where there are specific symmetry or dynamic causes for their suppression.

For example, the OZI rule for the vector mesons (ϕ, ω) system may wll be explainable by vector meson dominance of the strange quark vector current, which has zero forward matrix element because it measures the net strangeness quantum number of the hadron state.

$$\langle nonstrange\ hadron|\bar{s}\gamma_\mu s|nonstrange\ hadron\rangle = 0 \tag{51}$$

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* For an alternative view, suggesting that there is no strange quark OZI rule for baryons, as *vs.* mesons, see [50].

Such a picture is also compatible with the observed feature of 'channel-dependence' of the OZI rule for the strange quark. Recall that while there is OZI suppression for the vectors (ϕ, ω) the pseudoscalar system of (η, η') deviates significantly from ideal mixing. A challenge for future investigations will be to discover the specific mechanism for the suppression of the strange quark contribution to the DIS momentum sum rule, especially in view of its large contribution to the nucleon mass in the energy momemtum trace sum rule as discussed in this report.

(5) Strange quark content of the nucleon

It is also of interest to study the phenomenological consequences of a significant strange quark content of the nucleon in specific hadronic (πN, KN, *etc.*) and semi-leptonic (μN, νN, *etc.*) reactions. Further implications in the area of heavy ion physics, Higgs- and Z^0-nucleon coupling, glueball and multiquark spectroscopy should be examined. And, this problem of the flavor content of the nucleon could be of interest to nuclear physicists and astrophysicists studying pulsar and supernovae, as well as for detection of cosmic axions and other dark matter candidates.

The study of the flavor content of the nucleon can be an important probe of non-perturbative QCD. The problem is ideally suited for lattice QCD studies. Various lattice groups are in fact now beginning to investigate the effect of sea quarks on the hadron masses [51]. The result of a significant strange quark content discussed in this report implies such sea quarks should be very important in most lattice calculations. This probe should of course improve our understanding of the various effective low energy theories of QCD: $\frac{1}{N}$ expansion, chiral bag [52], skyrmion, etc. The proton spin and flavor content problem should be a useful handle for a better understanding of the mysterious connection between the constituent quarks and the QCD partons of quarks and gluons [7, 53].

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