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Orbital angular momentum, spin fractions, and scenarios for the proton's spin-weighted parton distributions

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New data from the European Muon Collaboration (EMC) on the spin-spin asymmetry in deep-inelastic lepton-proton scattering suggests that the total spin carried by valence quarks in a polarized proton may be approximately canceled by a strong negative polarization of the sea of $\bar{q}q$ pairs. The evolution in Q^2 of the fraction of proton spin carried by gluons depends on the initial spin fractions and it is possible to avoid the introduction of large orbital angular momentum by adopting a "hybrid" quark-Skyrme picture of the proton suggested by the EMC results.

Recent data from the European Muon Collaboration[1] (EMC) on the small-x behavior of the deep-inelastic scattering asymmetry for polarized leptons and polarized protons suggest an aspect of proton structure not easily understood in terms of traditional quark-model ideas. These data can also have significant implications for many other processes involving polarized protons and their impact on our ideas of hadron structure can be demonstrated by a simple exercise involving Q^2 evolution of parton spin fractions and orbital angular momentum.

For a proton with large momentum and positive helicity, it is possible to define the spin-weighted constituent distribution functions

$$\Delta q^i(x,Q^2) = q^i_{+/+}(x,Q^2) - q^i_{-/+}(x,Q^2)\,,\quad i = u^v, d^v, u^s, \bar{u}, d^s, \bar{d}, s, \bar{s}, \ldots, \qquad (1)$$
$$\Delta G(x,Q^2) = G_{+/+}(x,Q^2) - G_{-/+}(x,Q^2)\,.$$

The spin-weighted distributions enter into the spin-spin asymmetries of processes involving protons of definite helicity.[2,3] For example if we adopt a "process-independent" definition of the Δq_i, the structure function measured in leptoproduction with polarized leptons and a polarized proton target is given by

$$g_1^p(x,Q^2) = \sum_i e_i^2 \Delta q_i(x,Q^2)\left[1 + \frac{\alpha_s}{\pi}\tau(x,Q^2) + \cdots\right], \qquad (2)$$

where the $\tau(x,Q^2)$ is a calculable hard-scattering factor.[4] The Q^2 evolution of the $\Delta q_i(x,Q^2)$ is calculable using the Altarelli-Parisi formalism[5] and has been discussed by many people.[2,6] One result, which appears in the study of the Q^2 evolution of these functions and tells us a good deal about hadron structure, involves the orbital angular momentum of the partons.

To understand constituent models, it is frequently convenient to write sum rules involving conserved quantum numbers. For a proton with momentum and spin in the z direction, O(2) invariance gives the conservation of J_z. Writing $J_z = \frac{1}{2}$ in terms of the parton distributions defined above gives the sum rule[2]

$$\tfrac{1}{2} = \tfrac{1}{2}\sum_i \left(\langle \Delta q_i^{\rm val}\rangle + \langle \Delta q_i^{\rm sea} + \Delta\bar{q}_i\rangle\right) + \langle \Delta G\rangle + \langle L_z\rangle\,, \qquad (3)$$

where $\langle \Delta q_i\rangle = \int_0^1 dx\, \Delta q_i(x,Q^2)$, etc., can be interpreted as the net z component of spin carried by the various constituents and $\langle L_z\rangle$ is the total z component of orbital angular momentum. The term involving orbital angular momentum arises because there is no guarantee *a priori* that the z-component parton spins should sum to $\frac{1}{2}$. In fact, a quick study of the evolution equations suggests that orbital angular momentum is necessary. To see this we can write the evolution of the individual terms involving the constituents using the moment equations

$$\frac{\partial}{\partial t}\langle \Delta q_i\rangle = \Delta P^{(1)}_{q/q}\langle \Delta q_i\rangle + \Delta P^{(1)}_{q/G}\langle \Delta G\rangle\,,$$
$$\frac{\partial}{\partial t}\langle \Delta G\rangle = \sum_j \Delta P^{(1)}_{G/q}\langle \Delta q_j\rangle + \Delta P^{(1)}_{G/G}\langle \Delta G\rangle\,, \qquad (4)$$

where $t=(1/b_0)\ln[\alpha_s(Q_0)/\alpha_s(Q)]$, $b_0=\frac{11}{6}C_2(G)-\frac{2}{3}T(R)$. The $\Delta P_{i/j}^{(1)}$ are the first moments of the Altarelli-Parisi kernels. To lowest order these are[5]

$$\Delta P_{q/q}^{(1)}=0,\ \Delta P_{G/q}^{(1)}=\tfrac{3}{2}C_2(R)\,,\qquad \Delta P_{q/G}^{(1)}=0,\ \Delta P_{G/G}^{(1)}=b_0\,. \tag{5}$$

The vanishing of $(\partial/\partial t)\langle\Delta q_i\rangle$ can be understood in terms of the chiral properties of the theory with massless quarks. The equations for the amount of spin carried by the gluons is more complicated. Its solution depends on the values of the spin fractions on the right-hand side (RHS) of (4). Consistency of Eq. (3) requires $(\partial/\partial t)\langle L_z\rangle=-(\partial/\partial t)\langle\Delta G\rangle$, and we see that orbital angular momentum cannot be completely absent.[2,6,7] The idea of increasing $\langle L_z\rangle$ has been considered controversial. As presented above, there is no definition of the operator which defines $\langle L_z\rangle$ and the increase of $\langle L_z\rangle$ with resolution is inferred from the growth of $\langle\Delta G\rangle$. To avoid the growth of $\langle L_z\rangle$, a modification of the Altarelli-Parisi kernels has been suggested.[8] This suggestion, in its original form, can be rejected for a number of reasons.[7] However, the necessity for large orbital angular momentum can influence our picture of hadronic structure. Ratcliffe[9] has recently examined this problem and defined a measure of $\langle L_z\rangle$. Within this formalism, he was able to derive the evolution of $\langle L_z\rangle$ without reference to the $J_z=\frac{1}{2}$ sum rule.

There is one way to avoid $\langle L_z\rangle\to\infty$. If we write

$$\frac{\partial}{\partial t}\langle L_z\rangle=-\left[\tfrac{3}{2}C_2(R)\sum_i\langle\Delta q_i\rangle+b_0\langle\Delta G\rangle\right], \tag{6}$$

we see that the growth of orbital angular momentum is strongly connected to the net spin carried by the constituents and it is possible that the RHS of (6) vanishes.

We have, as yet, no direct or indirect experimental evidence concerning the polarized gluon distribution. We do, however, have some information concerning the spin fractions of the quarks in (6) from weak baryonic decays and from asymmetry measurements in electroproduction. Let us summarize this information in a simple parametrization. If we define

$$\langle\Delta u\rangle=\langle\Delta u^v\rangle+\langle\Delta u^s\rangle\,,\qquad \langle\Delta d\rangle=\langle\Delta d^v\rangle+\langle\Delta d^s\rangle\,, \tag{7}$$

then for the average polarization of the sea we can write

$$\langle\Delta d^s+\Delta\bar{d}\rangle=\langle\Delta u^s+\Delta\bar{u}\rangle=(1+\epsilon)\langle\Delta s+\Delta\bar{s}\rangle\,, \tag{8}$$

where ϵ is a parameter which measures the SU(3) violation. In the nomenclature of (7) and (8) we get two important constraints from the weak decays of baryons:[10,11]

$$A_3=\langle\Delta u^v-\Delta d^v\rangle=g_A=1.258\pm0.004\,,\qquad A_8=\langle\Delta u^v+\Delta d^v\rangle+2\epsilon\langle\Delta s+\Delta\bar{s}\rangle=0.54\pm0.10\,, \tag{9}$$

so that

$$\langle\Delta u_v\rangle=\tfrac{1}{2}(A_8+A_3)-\epsilon\langle\Delta s+\Delta\bar{s}\rangle\,,\qquad \langle\Delta d_v\rangle=\tfrac{1}{2}(A_8-A_3)-\epsilon\langle\Delta s+\Delta\bar{s}\rangle\,. \tag{10}$$

We can then incorporate the information concerning the weak decays into the integration of $g_1^p(x,\mu^2)$, as proposed by Ellis and Jaffe.[12] With the changes introduced into the formalism above, we define

$$G_1^p\equiv\left[1+\frac{\alpha_s(\mu^2)}{\pi}\right]\int_0^1 dx\,g_1^p(x,\mu^2)\,,\qquad G_1^p=\tfrac{1}{36}(5A_8+3A_3)+\tfrac{1}{3}\langle\Delta s+\Delta\bar{s}\rangle\,. \tag{11}$$

One can use data on the asymmetry $A_1=g_1/F_1$ to estimate the above integral. Using the extrapolation quoted by the experimenters,[1] and a value of $\alpha_s=0.24$, we obtain

$$G_1^p=0.123\pm0.013(\text{stat})\pm0.028(\text{syst})\,.$$

Inserting the values for A_8 and A_3 from weak decays into Eq. (11) gives

$$G_1^p=0.180\pm0.014+\delta G_1^p\,,\qquad \delta G_1^p\equiv\tfrac{1}{3}\langle\Delta s+\Delta\bar{s}\rangle\,. \tag{12}$$

Comparing Eq. (12) with the experimental determination of G_1^p we find

$$\delta G_1^p=-0.057\pm0.020(\text{stat})\pm0.029(\text{syst})\,.$$

The result is sensitive to the extrapolation of $g_1^p(x,Q^2)$ to small x. This point has been addressed in the analysis of Close and Roberts.[11] To illustrate the sensitivity, we assume that for $x<x_0$ we can approximate

$$g_1^p(x)\cong g_1^p(x_0)\left[\frac{x}{x_0}\right]^{-\alpha}, \tag{13}$$

for which the contribution to G_1^p from the small-x regime is given by

$$I(x_0)=\int_0^{x_0}g_1^p(x_0)dx\cong\frac{x_0g_1^p(x_0)}{1-\alpha}\,. \tag{14}$$

The extrapolation uncertainty is then controlled by the measured value of $g_1^p(x_0)$ and the unknown power behavior as $x\to0$. Regge-pole arguments suggest $\alpha\cong\alpha_{A_1}(0)\cong-0.1$ which is close to the value $\alpha=0$ assumed in the experimental analysis.[1] Even if one is wary about applying t-channel, coherent arguments on parton distributions, the assumption that $\Delta q_v(x,Q^2)/q_v(x,Q^2)$ vanish and $\Delta q_s(x,Q^2)/\Delta q_v(x,Q^2)$ be bounded at $x=0$ gives the bound $\alpha<\frac{1}{2}$ in (13) because of the observed small-x behavior of unpolarized nonsinglet distribution. This can be combined with EMC small-x measurements in (14) to give an upper limit of $+0.02$ to the extrapolation uncertainty.[13] The behavior assumed in (13) is subject to modification from the evolution of the quark densities with Q^2. However, a study of the small-x behavior of the nonsinglet evolution equation shows the modification to be negligible in the energy range of the EMC data.[14] The two spin fractions

$$f_v=A_8-6\epsilon\delta G_1^p=\langle\Delta u^v+\Delta d^v\rangle\,,\qquad f_s=(9+6\epsilon)\delta G_1^p=\langle\Delta u^s+\Delta\bar{u}+\Delta d^s+\Delta\bar{d}+\Delta s+\Delta\bar{s}\rangle\,, \tag{15}$$

are shown plotted vs δG_1^p for $\epsilon=0$, $\epsilon=\frac{1}{2}$, and $\epsilon=1$ in Fig.

1. Also shown is the sum

$$f_v+f_s=\sum_i\langle\Delta q_i\rangle=A_8+9\delta G_1^p\,, \tag{16}$$

which enters into the evolution equation for $\langle L_z\rangle$.

It is clear that for reasonable values of the SU(3)-breaking parameter, the data on δG_1^p suggest that f_v and f_s are both large and of opposite sign. They approximately cancel to produce an estimate for $\langle\Delta q_i\rangle$ which is small and independent of ϵ. Such a cancellation is suggestive of that which occurs in a hybrid bag+Skyrme model approach to proton structure in which

$$\tfrac{1}{2}=J_z^{\rm val}+(J_z^{\rm sea}+J_z^{\rm Skyrme})\,. \tag{17}$$

This type of model has been proposed[15] as a way of understanding the behavior of proton structure in which the *same* low-energy property of the proton, such as μ_P, can be explained either as a property of the valence quarks or as a property associated with the Skyrme solution. The dual requirements $J_z^{\rm val}+J_z^{\rm sea}\cong 0$ and $J_z^{\rm sea}+J_z^{\rm Skyrme}\cong 0$ allow each type of approximation to be valid. The fact that $\sum_i\langle\Delta q_i\rangle\cong 0$ in (16) is consistent with the Skyrme model has been suggested by Brodsky, Ellis, and Karliner.[16] The hybrid scenario allows us to reconcile this observation with the strong quark polarization measured at large x.

It is premature to completely adopt the hybrid scenario based on the EMC data because of the large statistical and systematic uncertainties. It is interesting to compare the Q^2 evolution of spin fractions under the hybrid scenario described above, with that of an "old-fashioned" or conventional quark-parton picture in which the polarization of the sea is small. We emphasize that we have no specific information about the makeup of the spin deficit, that is, the fraction of the proton's spin *not* carried by quarks. Perturbative QCD allows us to distinguish between the spin information being carried by polarized gluons and the presence of orbital angular momentum. The gluon polarization enters into measurable hard-scattering symmetries such as $A_{LL}(pp\to\gamma X)$.[2,3] As noted above, the starting value of the spin fractions also enters into the evolution equations for $\langle\Delta G\rangle$ and $\langle L_z\rangle$. The solutions to $\langle\Delta G(Q^2)\rangle$ and $\langle L_z(Q^2)\rangle$ can be written

$$\begin{aligned}\langle\Delta G(Q^2)\rangle&=\langle\Delta G(Q_0^2)\rangle[\alpha_s(Q_0^2)/\alpha_s(Q^2)]\\&\quad+\sum_i\frac{3}{2}\frac{C_2}{b_0}\langle\Delta q_i(Q_0^2)\rangle[\alpha_s(Q_0^2)/\alpha_s(Q^2)-1]\,,\\\langle L_z(Q^2)\rangle&=\tfrac{1}{2}-\tfrac{1}{2}\sum_i\langle\Delta q_i\rangle-\langle\Delta G(Q^2)\rangle\,.\end{aligned} \tag{18}$$

We demonstrate these by plotting the solutions for $\langle\Delta q^v+\Delta q^s\rangle>0$ in Fig. 2(a) and $\langle\Delta q^v+\Delta q^s\rangle=0$ in Fig. 2(b). In each plot we include two possible initial conditions, $\langle\Delta G(Q_0^2)\rangle=0$ and $\langle L_z(Q_0^2)\rangle=0$. It is interesting to note that the starting configuration

$$\sum_i\langle\Delta q_i\rangle\cong 0,\ \ \langle\Delta G(Q_0^2)\rangle\cong 0,\ \ \langle L_z(Q_0^2)\rangle\cong\tfrac{1}{2}$$

is stationary with respect to the Q^2 evolution. In view of the controversy involving the introduction of $\langle L_z\rangle\to-\infty$ as $Q^2\to\infty$, the possibility that nature adopts this evasive

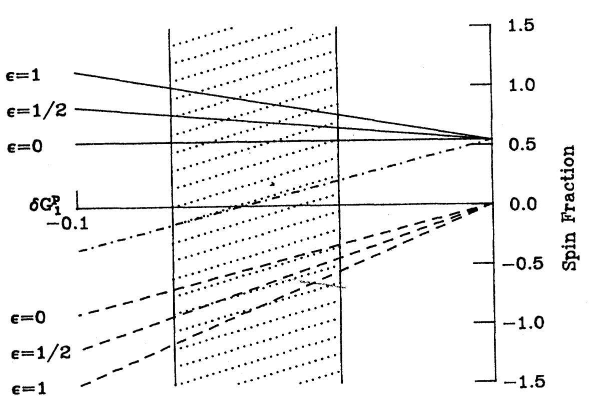

FIG. 1. The spin fractions f_v (solid line) and f_s (dashed line) are shown as a function of δG_1^p for $\epsilon=0$, $\epsilon=\frac{1}{2}$, and $\epsilon=1$. The dot-dashed line shows the total spin fraction f_v+f_s which enters into the evolution equation for $\langle L_z\rangle$. The shaded area denotes the values of δG_1^p covered by the statistical error.

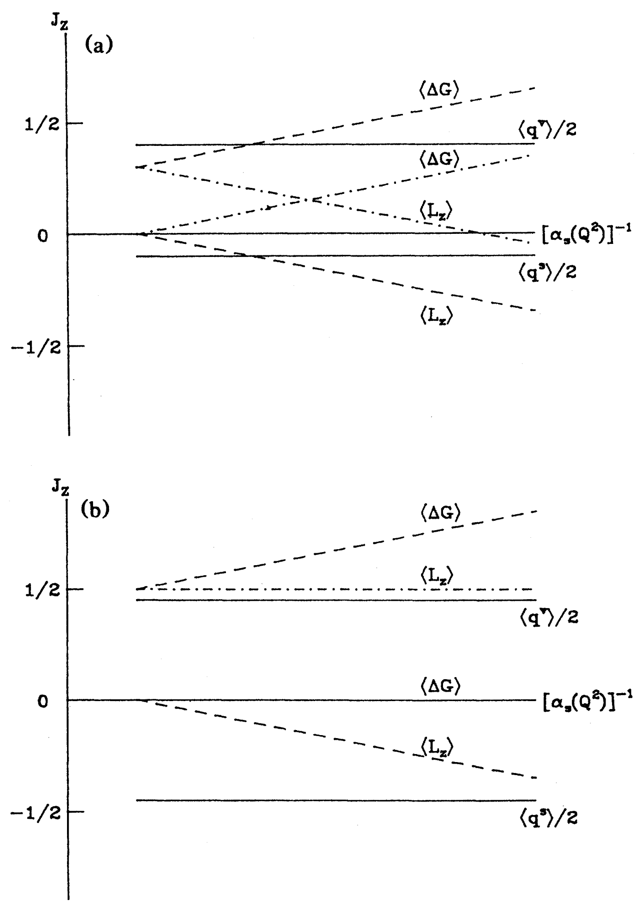

FIG. 2. The spin fractions carried by the valence quarks, sea quarks, gluons, and orbital angular momentum are shown against $[\alpha_s(Q^2)]^{-1}$ for the cases $\langle\Delta q^v+\Delta q^s\rangle>0$ (a), and $\langle\Delta q^v+\Delta q^s\rangle>0$ (b). The dashed and dot-dashed line correspond to the initial conditions $\langle L_z(Q_0^2)\rangle=0$ and $\langle\Delta G(Q_0^2)\rangle=0$, respectively.

approach to orbital angular momentum is worthy of careful consideration.

Although the spin fractions may be Q^2 independent in this final hybrid scenario, the full Altarelli-Parisi equations show that there is substantial structure in the x-dependent spin-weighted structure functions because of the strong correlation between the quark spin and the proton spin at large x. Model distribution functions embodying the large range of conjectured initial conditions will be given in a separate paper.[14] These distribution functions will be used to evaluate the different experimental measurements which may decide between the alternate models of the proton spin structure.

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