QUARKS AND GLUONS IN HADRONS AND NUCLEI

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ABSTRACT

These lectures discuss (1) the particle-nuclear interface — a general introduction to the ideas and application of colored quarks in nuclear physics, (2) color, the Pauli principle, and spin flavor correlations — this lecture shows how the magnetic moments of hadrons relate to the underlying color degree of freedom, and (3) the proton's spin — a quark model perspective. This lecture reviews recent excitement which has led some to claim that in deep inelastic polarized lepton scattering very little of the spin of a polarized proton is due to its quarks.

THE PARTICLE-NUCLEAR INTERFACE

Once upon a time nuclear physics was the study of nucleons and pions vibrating and oscillating at the center of the atom; particle physics was the study of nucleons and pions interacting and producing resonances. From the latter, people gradually realized that hadrons are built from quarks; the fundamental rules governing their interactions were deduced ("QCD" — quantum chromodynamics) and the similarities with QED suggested the possibility that all of the natural forces can be described in a grand unified theory. Today, particle physics deals with questions ranging from the origins of matter in the first microseconds of the universe, whose experimental investigation requires the energies of the SSC, down to the quark structure of protons and nuclei that can be studied at relatively low energies.

Nuclear physics theory has taken QCD on board. There are still detailed studies going on in nuclear excitations and there are important overlaps with nuclear astrophysics. The field is very rich. The attempt to understand nuclear phenomena at the quark level causes many nuclear physicists to be concerned with the same sort of problems as their colleagues in particle physics.

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During the last three years, there has been a blossoming in this overlap area which I would like to call "hadron physics". I would like to make the following definitions.

Nuclear Physics deals with the collective properties of the many-body nucleus consisting of nucleons. Hadron physics deals with the structure and interaction dynamics of those hadrons. Particle physics uses these particles as tools with which to elucidate deeper truths, seeking the origins of matter, the nature of mass, and the several other parameters which presently have to be invoked ad hoc (the weak mixing angles, fermion masses, etc.).

With tongue in cheek, one might contrast the two extremes. In extreme high-energy particle physics, there is infinite theorizing but almost no data; in the nuclear structure field, we have copious data but no truly fundamentally useful theory. In hadron physics we have much data and the hope of confronting them with the fundamental QCD theory of interacting quarks and gluons. This is stimulating but also difficult. A major problem is confinement of colored quarks and gluons within the nucleons — this is simulated on computers where space-time is described as a discrete lattice, but its origins analytically are still rather poorly understood.

This confinement phenomenon is also the catch-22 of "deriving" nuclear physics from QCD. As Bob Jaffe once remarked, "Looking for evidence of quarks in nuclei is like looking for the mafia in Sicily: everyone knows they are there, but it's hard to find the evidence." The quarks are confined in nucleons, and so any successful description of the nuclear structure must reduce to that of quarks clustered inside individual nucleons. Thus we have first to understand the proton and neutron — hence the interest in hadron physics. There is the interesting possibility that the interactions and overlaps among closely packed nucleons in large or dense nuclei, or in high-energy collisions of heavy ions, may disturb the distributions of quarks — their spatial or momentum distributions — relative to their behavior when confined in isolated free nucleons. There are indeed hints of such behavior, e.g., the "EMC effect" where quarks in iron have a slightly different momentum distribution relative to that in the deuteron. Their mean momenta are reduced in iron and other heavy nuclei suggesting greater spatial freedom. Is this "liberation" a "cold" precursor to "hot" deconfinement? Are nucleons in nuclei physically "enlarged", or is this a manifestation of quarks exchanged between nucleons, tying the nucleus together and having more spatial mobility than when in a single free nucleon? Future experiments may help to answer these questions — questions raised, in large degree, by the underlying quark theory.

I would like now to draw some analogies between QED (electrical charges, atoms, and molecules) on the one hand, and on the other, QCD (color, hadrons, and nuclei). The similarity is such that one could rewrite Bjorken and Drell's QED text by inserting a traceless 3×3 matrix (λ of SU(3)) at the fermion gauge boson vertices and let QED become QCD (with α replaced by $\alpha_s \approx 1/10$). However, the gluons themselves have color and so mutually interact via the color forces (contrast the photon of QED which transmits but does not directly "feel" the electromagnetic force). These new intergluon interactions give rise to vertices involving three or four gluons at a point, and so a text on QCD requires more than just a coloring of Bjorken and Drell.

Now let's make a matrix to summarize how systems variously react to the forces.

Notice that the gluon and photon are in different slots. This small difference gives rise to the different long-range phenomena in QCD compared to QED (e.g. confinement versus ionization).

QED	QCD
electric charge	3 colors
attraction of opposites	attraction of unlike colors
electrically zero atoms	colorless hadrons
radiation photon	radiation gluons
magnetic effects	chromomagnetic effects
hyperfine splitting $^3S_1 - {}^1S_0$	color hyperfine $m_\rho - m_\pi$, $m_\Delta - m_N$
Fermi-Breit in hydrogen	Fermi-Breit splittings in hadron spectroscopy

		QED	QCD
Feel the force	Carry the charge	e^- Z^+, Na^+ Cl^-	Quarks, Gluons
	Contain the charge	Atoms, Molecules	Hadrons, Nuclei
Do not feel the force; Do not contain the charge		Neutrinos, Photon	Leptons (ν,e)

Within atoms and hadrons one finds analogues. The Coulomb potential of hydrogen has an analogue in quark systems: as the quarks' relative separations $r \to o$, $V(r) \sim 1/r$, but at large r, $V(r) \sim r$, presumably due to the detailed self-interactions among the gluons that are transmitting the force. In the ground state of hydrogen the magnetic interaction ("one photon exchange") splits the $^3S_1$ and $^1S_0$ levels. In the ground state quark conglomerates, the high-J combinations have enhanced masses relative to their low-J counterparts due to "one gluon exchange"; thus the $\Delta(1230)$ resonance, with $J = 3/2$, has greater mass than the $J = 1/2$ nucleon.

Now move up a layer in complexity to the world of molecules (QED) and nuclei (QCD).

At the risk of being accused of oversimplification by the atomic experts, I will divide the interatomic forces into three broad classes, then make analogy in the QCD world with interhadronic forces at the level of the quarks and gluons.

	Covalent	van der Waals	Ionic
Atoms Molecules	e^- exchange	"two photon"	Na^+Cl^-
Hadrons Nuclei	quark exchange	"two gluon"	no analogue if color confined in neutral clusters

If this was the whole story, then nuclear physics from QCD would be a rerun of molecules from QED. However, confinement of color breaks the simple analogy. The quark exchange at large distances (> 1 fm) is contained within the confined packages, dominantly pions. The confinement of gluons in glueballs also breaks the analogy with van der Waals' forces. The hope that QCD would predict observable color van der Waals' forces in nuclei is most probably flawed, as the gluons will be confined within colorless glueballs. Computer simulations of QCD suggest that the lightest glueballs have masses in excess of 1 GeV and so transmit forces over much less than a nucleon radius. Thus their presence is hidden in nuclear physics.

It is an open question whether analogues of ionic forces occur in dense or hot nuclear systems; whether multiquark clusters occur within nuclei; whether quark-gluon plasma may form in hot-dense systems.

If color attractions among quarks are the source of internucleon forces, then there could exist analogous clusters of mesons — "meson molecules". The instability of most mesons prevents formation of these systems, but π, K, η are stable on the time scales of the strong interactions and may have the chance to bind. Indeed Weinstein and Isgur find that such attractions occur in S-wave. The $\pi\pi$ system has a strong enhancement above $2m_\pi$ which may be manifested in the $\psi \to \omega\pi\pi$ dipion spectrum. The KK system binds forming nearly degenerate I = 0, 1 systems 10 MeV below $2m_K$. The S^*(975 MeV) and δ(980 MeV), scalar "mesons", thus appear to be meson molecules; meson analogues of the I=0 dueteron (whose I=1 partner is above $2m_N$).

The color attractions among quarks and gluons lead to the prediction of glueballs and hybrid hadrons — the latter where gluons play a dynamical role, attracted to quarks to form hybrid mesons and baryons.

The problem in predicting the masses of these states is that we have to simulate the effects of confinement. Perhaps the simplest way of doing this is to suppose that the constituent quarks or gluons are free until they hit an infinitely high wall. This is the essence of cavity or bag models. Confine a massless J=1/2 quark in a radius, R, and it gains an energy that scales as 1/R. This energy becomes of the order of 350 MeV if R is of the order of the proton radius, hence the proton mass may be modeled. For gluons, one solves the eigenvalue equations for J = 1 rather than J = 1/2 confined fields. There are electric or magnetic modes (actually TE and TM in the language of classical electrodynamics) with different eigenvalues. If R is the same as for quark systems, the typical confined-mass-scale is some 500 MeV per TE mode and 750 MeV per TM mode. Thus follows the prediction that the lightest systems consisting of at least two confined gluons weigh in at O(1 GeV) and that the lightest hybrid baryons weigh in at O(1.5 GeV). A problem is that as soon as the hyperfine shifts in energy are taken into account (this involves one first calculating the propagators of confined quarks and gluons), the lowest spin-J systems are pulled down significantly in mass. The lightest hybrid baryon might thus appear to have a mass near that of the proton which suggests either a profound rethink of baryon spectroscopy or that we have unearthed a naievity.

I suspect it is the latter. No one yet has convincingly set up a study of loop effects with renormalization within a cavity. These loop diagrams enter at the same order in perturbation theory to which the hyperfine shifts have been calculated and may alter the naive "effective" energies per confined gluon. In the case of quarks, their effects were subsumed in the MIT bag by an input mass parameter for the quark; this mass fitted to the overall mass scale of the spectroscopy. In the gluonic sector we have no mass scale to set the scale, and until we make sense of the (infinite!) self-energy diagrams, we cannot predict the absolute scale. So the mass separations among the various states may be reliable, but the absolute mass scale is beyond present analysis. To predict the masses of glueballs and hybrid hadrons, we have to resort to computer simulations — lattice QCD. This has proved to be a harder task than was originally thought.

The eventual discovery of the gluonic spectroscopy may give important insights into the nature of confinement of gluons. If lattice calculations, including quarks <u>and</u> gluons (to date, people work in the "quenched" approximation, which roughly translated means "ignore the quarks") merely point out masses of states that correspond to the particle data tables, we will confirm QCD but may still require much study to elucidate the analytic dynamics of confinement. The main outcome of such a success may be the advances that will have come in the art of computation and design of machines. Thus the significant questions posed by hadron physics are having a spinoff in the intellectual stimulation they provide to computational science and, in turn, the subsequent ability to encode problems in field theory, condensed matter, and other areas of science. I am reminded of the title of Tony Hey's talk at a recent meeting of the British Association for the Advancement of Science, and it provides an apt one-line summary of the multidisciplinary efforts flowing from computation at the nuclear-particle interface. It was: "Quarks, Supercomputers, and Oil Prospecting".

COLOR, THE PAULI PRINCIPLE, AND SPIN-FLAVOR CORRELATIONS

Color

If quarks possess a property called color, any quark being able to carry any one of three colors (say red, yellow, blue), then the Ω^- (and any baryon) can be built from distinguishable quarks:

$$\Omega^- \left(s_R^{\uparrow}\ s_Y^{\uparrow}\ s_B^{\uparrow}\right).$$

If quarks carry color but leptons do not, then it is natural to speculate that color may be the property that is the source of the strong interquark forces — absent for leptons.

Electric charges obey the rule "like repel, unlike attract" and cluster to net uncharged systems. Colors obey a similar rule: "like colors repel, unlike (can) attract". If the three colors form the basis of an SU(3) group, then they cluster to form "white" systems — viz. the singlets of SU(3). Given a random soup of colored quarks, the attractions gather them into white clusters, at which point the color forces are saturated.

The residual forces among these clusters are the nuclear forces whose origin will be mentioned later.

If quark (Q) and antiquark ($\bar{Q}$) are the $\underline{3}$ and $\bar{\underline{3}}$ of color SU(3), then combining up to three together gives SU(3) multiplets of dimensions as follows (see Ref. 3):

$$QQ = \underset{\sim}{3} \times \underset{\sim}{3} = \underset{\sim}{6} + \underset{\sim}{\bar{3}}$$

$$Q\bar{Q} = \underset{\sim}{3} \times \underset{\sim}{\bar{3}} = \underset{\sim}{8} + \underset{\sim}{1}$$

The $Q\bar{Q}$ contains a singlet — the physical mesons

$$QQ\bar{Q} = \underset{\sim}{15} + \underset{\sim}{6} + \underset{\sim}{3} + \underset{\sim}{3}$$

$$QQQ = \underset{\sim}{10} + \underset{\sim}{8} + \underset{\sim}{8} + \underset{\sim}{1}.$$

Note the singlet in QQQ — the physical baryons.

For clusters of three or less, only $Q\bar{Q}$ and QQQ contain color singlets and, moreover, these are the only states realized physically. Thus are we led to hypothesize that only color singlets can exist free in the laboratory; in particular, the quarks will not exist as free particles.

<u>Symmetries and Correlations in Baryons</u>

To have three quarks in color singlet:

$$\underset{\sim}{1} \equiv \frac{1}{\sqrt{6}} [(RB-BR)Y + (YR-RY)B + (BY-YB)R] \tag{1}$$

any pair is in the $\underset{\sim}{\bar{3}}$ and is antisymmetric. Note that $\underset{\sim}{3} \times \underset{\sim}{3} = \underset{\sim}{6} + \underset{\sim}{\bar{3}}$. These are explicitly

$\underset{\sim}{\bar{3}}_{anti}$	$\underset{\sim}{6}_{sym}$
RB−BR	RB+BR
RY−YR	RY+YR
BY−YB	BY+YB
	RR
	BB
	YY

(2)

Note well: <u>Any Pair is Color Antisymmetric</u>

The Pauli principle requires total antisymmetry and therefore any pair must be:

<u>Symmetric in all else</u>

("else" means "apart from color").

This is an important difference from nuclear clusters where the nucleons have no color (hence are trivially symmetric in color!). Hence for nucleons Pauli says

Nucleons are Antisymmetric in Pairs (3)

and for quarks

Quarks are Symmetric in Pairs (4)

If we forget about color (color has taken care of the antisymmetry and won't affect us again), then

(i) Two quarks can couple their spins as follows

$$\left|\begin{array}{ll} S = 1: & \text{symmetric} \\ S = 0: & \text{antisymmetric} \end{array}\right| \tag{5}$$

(ii) Two u,d quarks similarly form isospin states

$$\left\{\begin{array}{ll} I = 1: & \text{symmetric} \\ I = 0: & \text{antisymmetric} \end{array}\right\} \tag{6}$$

(iii) In the ground state L = 0 for all quarks; hence the orbital state is trivially symmetric. Thus for pairs in L = 0, we have

$$\left\{\begin{array}{l} S = 1 \text{ and } I = 1 \ \text{ correlate} \\ S = 0 \text{ and } I = 0 \ \text{ correlate} \end{array}\right\}. \tag{7}$$

Thus the Σ^0 and Λ^0 which are distinguished by their u,d being I = 1 or 0 respectively also have the u,d pair in spin = 1 or 0 respectively:

$$\left|\begin{array}{l} \Sigma^0 (u,d)_{I=1} s \rightarrow (u,d)_{S=1} S \\ \Lambda^0 (u,d)_{I=0} s \rightarrow (u,d)_{S=0} S \end{array}\right|. \tag{8}$$

Thus, the spin of the Λ^0 is carried entirely by the strange quark.

This is the source of the Σ–Λ mass difference. The $\vec{S}\cdot\vec{S}$ interaction acts between all possible pairs; thus

$$\Sigma^0 \ [(u,d)_1 s]: \ \langle \vec{S}\cdot\vec{S}\rangle_1 + \langle \vec{S}\cdot\vec{S}\rangle_{s,1} \tag{9}$$

$$\Lambda^0 \ [(u,d)_0 s]: \ \langle \vec{S}\cdot\vec{S}\rangle_0 \tag{10}$$

(note $\langle \vec{S}\cdot\vec{S}\rangle$ between a spinless diquark and anything vanishes; hence the absence of $\langle S\cdot S\rangle_{s,o}$).

Now

$$\langle \vec{S}\cdot\vec{S}\rangle_0 = -3 \ \langle \vec{S}\cdot\vec{S}\rangle_1 , \tag{11}$$

(see p. 91 of Ref. 3). Further, if $m_s = m_{u,d}$, the Σ and Λ become mass degenerate, and so in this limit

$$\langle \vec{S}\cdot\vec{S}\rangle_{s,1} = -4 \ \langle \vec{S}\cdot\vec{S}\rangle_1 . \tag{12}$$

For unequal masses of u and s, the magnetic interaction scales as the inverse mass. Hence finally

$$\Sigma^0 \sim \langle \vec{S}\cdot\vec{S}\rangle_1 \left\{1-4\,\frac{m_u}{m_s}\right\} \tag{13}$$

$$\Lambda^0 \sim \langle \vec{S}\cdot\vec{S}\rangle_0 \,\{-3 \quad\}. \tag{14}$$

Then with $m_s > m_u$, we find $m_\Sigma > m_\Lambda$ as observed. Increasing m_s/m_u enhances the effect (e.g., for the charmed analogues $\Sigma_c[(u,d)c]$ and $\Lambda_c[(u,d)c]$ the splitting will be larger — again observed).

Color, the Pauli Principle, and Magnetic Moments

The electrical charge of a baryon is the sum of its constituent quark charges. The magnetic moment is an intimate probe of the correlations between the charges and spins of the constituents. Being wise, today we can say that the neutron magnetic moment was the first clue that the nucleons are not elementary particles. Conversely the fact that quarks appear to have $g \simeq 2$ suggests that they are elementary (or that new dynamics is at work if composite).

A very beautiful demonstration of symmetry at work is the magnetic moment of two similar sets of systems of three, viz.

$$\begin{Bmatrix} N & ; & P \\ ddu; & & uud \end{Bmatrix} \qquad \mu_P/\mu_N = -3/2$$

and the nuclei

$$\begin{Bmatrix} H^3 & ; & He^3 \\ NNP; & & PPN \end{Bmatrix} \qquad \mu_{He}/\mu_H = -2/3.$$

The Pauli principle for nucleons requires He^4 to have no magnetic moment:

$$\mu[He^4;\; P^{\uparrow}P^{\downarrow}N^{\uparrow}N^{\downarrow}] = 0.$$

Then

$$He^3 \equiv He^4 - N$$

$$H^3 \equiv He^4 - P$$

and so

$$\frac{\mu_{He3}}{\mu_{H^3}} = \frac{\mu_N}{\mu_P}$$

To get at this result in a way that will bring best comparison with the nucleon three-quark example, let's study the He^3 directly.

He^3 = ppn: pp are flavor symmetric; hence, spin antisymmetric; i.e., S = 0.

Thus

$$[He^3]^{\uparrow} \equiv (pp)_o\; n^{\uparrow} \tag{15}$$

and so the pp do not contribute to its magnetic moment. The magnetic moment (up to mass scale factors) is

$$\mu_{He^3} = 0 + \mu_N. \tag{16}$$

Similarly,

$$\mu_{H^3} = 0 + \mu_P. \tag{17}$$

Now let's study the nucleons in an analogous manner.

The proton contains u,u flavor symmetric and <u>color antisymmetric</u>; thus the spin of the "like" pair is symmetric (S = 1) in contrast to the nuclear example where this pair had S = 0. Thus coupling spin 1 and spin 1/2 together, the Clebsches yield

$$p^{\uparrow} = \frac{1}{\sqrt{3}} (u,u)_0 d^{\uparrow} + \sqrt{\frac{2}{3}} (u,u)_1 d^{\downarrow} \tag{18}$$

(contrast Eq. (15)), and (up to mass factors)

$$\mu_P = \frac{1}{3} (0+d) + \frac{2}{3} (2u-d). \tag{19}$$

Suppose that $\mu_{u,d} \propto e_{u,d}$, then

$$\mu_u = -2 \mu_d \tag{20}$$

so

$$\frac{\mu_P}{\mu_N} = \frac{4u-d}{4d-u} = -\frac{3}{2} \tag{21}$$

(the neutron follows from proton by replacing $u \to d$).

I cannot overstress the crucial, hidden role that color played here in getting the flavor-spin correlation right.

We can extend this discussion to the full baryon octet. Six of these states contain two identical quark flavors (which by symmetry necessitates that this pair have total spin S = 1):

$$\left.\begin{array}{l} \left[P(uu)_1 d\right] \\ \left[N(dd)_1 u\right] \\ \left[\Sigma^+(uu)_1 s\right] \\ \left[\Sigma^-(dd)_1 s\right] \\ \left[\Xi^0(ss)_1 u\right] \\ \left[\Xi^-(ss)_1 d\right] \end{array}\right\} \tag{22}$$

The remaining pair are Σ^0 and Λ^0, both u,d's. In the former, the (ud) have I = 1 and hence S = 1. For the Λ^0, on the other hand, the (ud) have I = 0 and hence S = 0.

Thus the Λ^0 $\left[(u,d)_0 s\right]$ is analogous to the He^3 nuclear example. The magnetic moment is carried entirely by the third quark, namely s. The data yield[8]

$$\mu_\Lambda + \mu_s \simeq \frac{3}{5}\,\mu_d. \tag{23}$$

The strange and down quarks have the same charge (-1/3) and so the datum fits with

$$m_d = \frac{3}{5}\,m_s \tag{24}$$

as already noted.

If we approximate $m_u = \frac{1}{3}\,m_p$ (thus the proton would have $g = 3$), then we can do a quick computation of baryon magnetic moments where the individual contributions to the g factors are

$$\left.\begin{aligned} u &= 2 \\ d &= -1 \quad \left(\text{ratio of } e_d/e_u\right) \\ s &= -3/5 \quad \left(\text{ratio of } m_d/m_s\right) \end{aligned}\right\}. \tag{25}$$

From the general spin structure of Eq. (18)

$$B^{\uparrow} = \frac{1}{\sqrt{3}}\,(q_1q_2)_{+}\,q_3^{\uparrow} + \sqrt{\frac{2}{3}}\,(q_1q_2)_{+}\,q_3^{\downarrow}$$

we have

$$\mu = \frac{2}{3}\,(q_1+q_2) - \frac{1}{3}\,q_3 \tag{26}$$

into which the (25) are to be substituted as required. The resulting pattern is as follows

	Prediction	Data	
$P[(uu)_1d]$	$\frac{1}{3}$ (4u-d) = 3	2.79	(27)
$N[(dd)_1u]$	$\frac{1}{3}$ (4d-u) = -2	-1.9	(28)
$\Sigma^+[(uu)_1s]$	$\frac{1}{3}$ (4u-s) = 2.8	2.33 ± 0.13	(29)
$\Sigma^-[(dd)_1s]$	$\frac{1}{3}$ (4d-s) = -1.1	-1.41 ± 0.25	(30)
$\Xi^0[(ss)_1u]$	$\frac{1}{3}$ (4s-u) = -1.5	-1.25 ± 0.02	(31)
$\Xi^-[(ss)_1d]$	$\frac{1}{3}$ (4s-d) = -0.5	-0.75 ± 0.06	(32)

The trend is exceptionally well described. There are undeniably 20% effects not fully accounted for.

Encouraged by this success, we might look further at this problem since, after all, there are exchange effects in nuclei that cause 20% deviations from the naive additive approach analogous to that which we have used for quarks.

We can form contributions of u and d quarks from (27) and (28) for nucleon, (29,30) for Σ, (31,32) for Ξ, and the data yield

$$(u-d)_N = 2.9$$

$$(u-d)_\Sigma = 1.7 \pm 0.15$$

$$(u-d)_\Xi = 0.9 \pm 0.12.$$

As we go to systems with more strange quarks, the u,d quarks act as if their effective mass increases (by a factor of three??). There is a systematic trend but extremely dramatic.

Now let's study the strange quark. We can do this by supposing that the environmental dependence for u and d flavors is the same. Then

$$(s)_\Xi \equiv -\left(\Xi^0 + 2\Xi^-\right) = -0.69 \pm 0.03$$

$$(s)_\Sigma \quad \left(\Sigma^+ + 2\Sigma^-\right) \quad = -0.5 \pm 0.5$$

$$(s)_\Lambda = -0.6$$

So the strange quark gives its "canonical" contribution to baryons containing either one or two strange quarks.

THE PROTON'S SPIN: A QUARK MODEL PERSPECTIVE

Inelastic lepton scattering from nucleons at high momentum transfer measures the number densities of charged constituents, $q(x)$, $\bar{q}(x)$, as a function of the Bjorken variable x (essentially the ratio of the constituent and target longitudinal momenta in an infinite momentum frame). There is a weak dependence of these distributions on the momentum transfer, Q^2, but I shall suppress this in much of what follows.

If the beam and target are polarized, one can extract the helicity-dependent distributions for quarks or antiquarks polarized parallel $\left(q^{\uparrow}(x)\right)$ or antiparallel $\left(q^{\downarrow}(x)\right)$ to the target polarization. I shall define $\Delta q(x) \equiv q^{\uparrow}(x) - q^{\downarrow}(x)$; $q(x) \equiv q^{\uparrow}(x) + q^{\downarrow}(x)$, and similarly for antiquarks, $\bar{q}$.

Data are presented in two ways.[1-3] One is in terms of the polarization asymmetry

$$A(x) = \sum_i e_i^2 \left(\Delta q_i(x) + \Delta\bar{q}_i(x)\right)/\sum_i e_i^2 \left(q_i(x) + \bar{q}_i(x)\right), \tag{33}$$

(note that $-1 < A < +1$). The other involves the polarized structure function

$$g_1(x) = \frac{1}{2} \sum_i e_i^2 \left(\Delta q_i(x) + \Delta\bar{q}_i(x)\right), \tag{34}$$

thus

$$g_1(x) \equiv A(x)\ F_1(x). \tag{35}$$

In advance of the data, the expectations were that

(i) At $x > 0.2$ where valence quarks dominate, $A(x)$ should be large and positive.[3,4] This follows from intuition developed for constituent valence quarks in baryon spectroscopy where the Pauli principle requires $\Delta u > 0$, $\Delta d < 0$. As the charge-squared weighting of Δu is four times that of Δd in protons, so $A^p(x > 0.2) > 0$. Data confirm this brilliantly. For a neutron target, it is Δd that is weighted 4:1 relative to Δu, hence these tend to cancel and one predicts[4] a small (zero?) asymmetry on the neutron.

(ii) Form $g_1(x)$, which directly shows the charge weighted helicity-dependent distributions and integrate over all x.[5,6] If it were not for the charge weightings, this would measure the net $\Delta q + \Delta\bar{q}$ $\left(\Delta q \equiv \int_0^1 dx \Delta q(x) = \text{net quark polarization}\right)$.

Explicitly, in the quark parton model

$$I^p \equiv \int dx\ g_1^p(x) = \frac{1}{2}\left\{\frac{3}{9}\ \Delta u + \frac{1}{9}\ (\Delta u+\Delta d+\Delta s)\right\} + \left(\Delta q_i - \Delta\bar{q}_i\right). \tag{36}$$

The surprise[2] is that $I^p(\text{EMC}) \approx 0.12$ and is almost saturated by[7] $\Delta u(\approx 0.75)$ leaving

$$\sum_i (\Delta q + \Delta\bar{q})_i \approx 0, \tag{37}$$

hence, the much-advertised claim that maybe "none of the proton's spin polarization is carried by quarks". This is a misinterpretation of Eq. (37). The valence quarks are highly polarized (point (i) above); thus, the interpretation of Eq. (37) is that something cancels or hides it. Candidates include a highly polarized sea spinning opposite to the valence quarks, orbital angular momentum, or gluon polarization.[8-10]

One can cancel out some charge weighting effects by looking at the difference of proton and neutron for which

$$I^p - I^n = \frac{1}{6}\ (\Delta u-\Delta d) \equiv \frac{1}{6}\left|\frac{g_A}{g_V}\right|, \tag{38}$$

which is Bjorken's sum rule.[5] The various g_A in the baryon octet give information on the <u>differences</u> of Δu, Δd, and Δs which are summarized by a measured parameter known as F/D. To extract the <u>sum</u>, Δq, we need the proton integral (Eq. (36)) or information on neutral current form factors

$$\tilde{g}_A(\nu p \to \nu p) = \Delta u - \Delta d - \Delta s. \tag{39}$$

I shall discuss this at the end of the talk. Preceding that, I shall discuss the question of Δs, since the measured F/D and the measured I^p can be combined to extract a value for Δs. This appears to be substantial; EMC claiming that

$$\Delta s = -0.23 \pm 0.08. \tag{40}$$

Implications and criticisms of this startling result will occupy the latter half of this talk. First, I will discuss what we know about the (constituent) quark polarization from static properties of the nucleon (magnetic moments, g_A/g_V) and review the extent to which the new insights do or do not require revision of this simple picture.

Spin Polarization of Valence (Constituent) Quarks

In the constituent quark model where $L_Z = 0$ the charges and the magnetic moments of neutron and proton place the following constraints on the probabilities for finding the flavors and spin correlations of "valence" quarks,

$$u_v = 2d_v \quad \frac{\mu_n}{\mu_p} = -\frac{2}{3} \quad \Delta u_v = -4\Delta d_v. \tag{41}$$

The 56, $L_Z = 0$ wave function of the nonrelativistic quark model (NRQM) satisfies (41) but it is by no means unique. A hybrid state, where a gluon ($J_Z = \pm 1$) is partnered by qqq in 70 (required by the Pauli principle for qqq in color 8) satisfies Eq. (41) for the coherent combination[11] $g(^28+^48)$ where the superscripts refer to the 2S+1 of the net spin of the qqq system. The "valence quarks" here are significantly depolarized relative to 56. One can also have significant polarized sea without destroying the magnetic moment relations. This is because

$$\frac{\mu_n}{\mu_p} = \frac{2\Delta d-\Delta u+(-2\Delta\bar{u}+\Delta\bar{d}+R\Delta\bar{s})}{2\Delta u-\Delta d+(-2\Delta\bar{u}+\Delta\bar{d}+R\Delta\bar{s})}, \tag{42}$$

where $R = m_d/m_s = 3/5$. The electrical neutrality of the sea tends to shield its contribution. A detailed fit is made in Ref. 12.

The $(g_A g_V)$ for the octet of baryons also relate to the spin polarized probabilities such as

$$\left(\frac{g_A}{g_V}\right)_{np} = \Delta u_v - \Delta d_v + -5\Delta d_v, \tag{43}$$

where we used Eq. (41). Thus immediately

$$\Delta d_v = -0.25; \quad \Delta u_v = 1. \tag{44}$$

In the 56 NRQM one would have[3]

$$\Delta d_v = -1/3; \quad \Delta u_v = 4/3; \quad \Delta u_v + \Delta d_v = 1, \tag{45}$$

and the entire spin polarization comes from the quarks. However, from Eq. (44), we see that

$$\Delta u_v + \Delta d_v = 3/4, \tag{46}$$

and so, in advance of the EMC data, only naive "quarkists" would have expected 100% for Δq_v. Anyone who worked with four-component spinors, of which the MIT bag is a specific model example, knew that the "orbital dilution" in the lower components played an essential role.[13] In fact, the Δq_v expectation is even less than Eq. (46). When one makes a best fit to all of the baryon octet g_A/g_V, one finds

$$\Delta q_v (\equiv 3F-D) = 0.55 \pm 0.10. \tag{47}$$

Note the appearance of F and D which summarizes the g_A/g_V. This parameter will appear later. Note that many analyses of the polarization data use[2,6,10,15] F/D = 0.63 (Ref. 16). However, this value fitted a value of the neutron lifetime that we now know to have been incorrect.[17,18] The correct current value[19-25] is lower than 0.63 and is dependent upon assumptions about SU(3) flavor breaking.

The earliest predictions for the deep inelastic polarization asymmetry in the valence-dominated region assumed that all Δq and q (valence) have the same x dependence. Thus (see Refs. 3 and 4 for origins of these formulae)

$$A^n(x) \simeq 4\Delta d + \Delta u \to 0,$$

(the zero following immediately from Eq. (41)) and

$$A^p(x) = \frac{5}{3}(-\Delta d) + \frac{1}{3}(g_A/g_V).$$

The prediction that $A^p > 0$ is non-trivial as a priori it could be anywhere in the range $-1 < A < +1$. The presence of a $q\bar{q}$ sea as $x \to 0$ was expected to cause $A(x \to 0) \to 0$. The other qualitative expectation[26,27] was that $A(x \to 1) \to 1$ as follows.

The valence picture above implicitly assumed that $u_v(x) = 2d_v(x)$ for all x. However, unpolarized data show this to be untrue in that it would require that

$$\frac{F_1^n(x)}{F_1^p(x)} = 2/3.$$

In practice, this ratio drops as $x \to 1$, suggesting that the $u(x \to 1) \gg d(x \to 1)$, a phenomenon which follows from spin dependence via single gluon exchange. Chromomagnetic hyperfine energy shifts split the Δ-N masses and elevate $u(x \to 1)$ over $d(x \to 1)$. They also cause $u^\uparrow(x \to 1)$ to dominate over $u^\downarrow(x \to 1)$, which the consequence that $A^{p,n}(x \to 1) \to 1$. Thus, a qualitative expectation for A^p emerged:

$$A^p(x \to 0) \to 0;\; A^p(x \simeq 1/3) \simeq 1/3\left|\frac{g_A}{g_V}\right|;\; A^p(x \to 1) \to 1.$$

These predictions turned out to be remarkably well verified and even agree with the latest EMC data.

Recently Close and Thomas[28] showed that, within the framework of the MIT bag model, one could relate the x-dependent distortion of the valence distributions to the measured chromomagnetic energy shift in the Δ-N masses. All of this suggests that the valence quark polarizations measured in polarized deep inelastic scattering are similar to the polarizations of the constituent quarks manifested in low-energy spectroscopy. This is an important constraint on model builders. The memory of the constituent quark spins is not lost as one proceeds to the deep inelastic: the valence quarks are highly polarized.

If, as is being claimed, the quarks and antiquarks contribute (within errors) nothing to the net spin polarization of the proton, then we must conclude that something is canceling the contribution of the valence quarks. Candidates include orbital angular momentum polarized gluons or a negatively polarized sea.

We already noted that in the constituent limit it is over naive to ignore orbital angular momentum. The presence of polarized gluons may be probed by studying the polarization dependence of direct photon production or spin dependence of heavy flavor production; a polarized sea may affect the inclusive production of hadrons[33,34] and fast $K^-(s\bar{u})$ production may be a tag for scattering from the sea.[34]

Dziembowski et al.[35] have studied the relation between constituent quarks and partons. They view the constituent quarks as being a conglomerate of partons—quarks, antiquarks, and gluons, thus

$$q_i^{\lambda_i}(x,Q^2) = \sum_{v,\lambda_v} \int_x^1 \frac{dy}{y} G_{v/N}^{\lambda_v}(y)\, q_{i/v}^{\lambda_i \lambda_v}\left(\frac{x}{y}, Q^2\right),$$

where the λ are helicity labels. The constituent quark distributions $G_{v/N}(y)$ reflect the dynamics that binds the quarks to form hadrons, and are determined by a light cone nucleon wavefunction. The constituent quark structure functions $q(x/y, Q^2)$ are adapted from Altarelli et al.[36] together with Carlitz and Kaur's ansatz[37] for the spin of soft valence partons within a polarized constituent quark.

This picture of partons convoluted within constituents generates some effective $L_z \neq 0$ but not enough to account for the spin deficit claimed by EMC. The data seem to fall below the model systematically for $x < 0.1$. If these small x data survive further experiments, then it seems that polarization of the sea (not included in Ref. 35) must be allowed for. This naturally leads to the question of whether there are polarized strange quarks in the proton.

Polarized Strange Quarks?

One exciting possibility is that the EMC data imply a large polarization of strange quarks and/or antiquarks within the proton. If true, this could have significant consequences. In particular, it could modify earlier analyses of electroweak parity violation in deuterium where Campbell et al. argue,[15] the polarized strange quarks could give contributions that dominate over electroweak radiative corrections. An extreme claim has appeared in the literature that the large value for Δs is in conflict with perturbative QCD. If true, this would be devastating. This claim comes about, in part, because an incorrect value of F/D has been used in the analyses. It is this parameter, and its implications for Δs, that I will now discuss.

Given the integral, I_p, of the polarized structure function $g_1^p(x,Q^2)$, one extracts Δs (including new QCD corrections)

$$I_p \equiv \int dx g_1^p(x,Q^2) = \frac{1}{18}\left(\frac{g_A}{g_V}\right)\left[\frac{9f-1}{f+1} - \frac{\alpha_s(Q^2)}{\pi}\frac{3f+1}{f+1}\right] + \frac{\Delta s}{3}\,, \qquad (48)$$

where $f \equiv F/D$ with $\alpha_s(Q^2) = 0.27$, $g_A/g_V = 1.254 \pm 0.006$ and $I_p = 0.126 \pm 0.022$. A feeling for the sensitivity of Δs to f can be gauged from the approximate relation

$$\Delta s \simeq (f-0.40) \pm 0.07. \qquad (49)$$

The widely used value, following the much-quoted fit of Ref. 16 has been

$$F/D = 0.63 \pm 0.02 \rightarrow \Delta s = -0.23 \pm 0.09. \qquad (50)$$

If the sea is flavor-independent, then Eq. (50) summarizes the widely accepted interpretation of the EMC polarized structure function data where a significant negative polarization of the sea cancels out the positive polarization of the valence quarks.

This value was based on the original value for I_p quoted by EMC,[2] namely $I_p = 0.116 \pm 0.022$. However, the revised value,[29] $I_p = 0.126 \pm 0.022$, reduces the magnitude of Δs by 0.03, and so $\Delta s = -0.20 \pm 0.09$ should replace Eq. (50).

However, it does not seem to be widely appreciated that the F/D of Ref. 16 was much constrained by an outdated value of the neutron lifetime, and that Ref. 16 chose "to omit from (their) fit the neutron decay correlation (which yields) $g_A = 1.258 \pm 0.009$, which differs significantly from the result 1.239 ± 0.009 required by the neutron lifetime measurements". The value accepted as correct today[18] differs by some 3σ from the old value, and this, together with other data on hyperon beta decays,[16,18,19] shows that F/D is much smaller than the old value. Flavor symmetry breaking causes a spread in values of F/D, depending on which partial set of data one uses; indeed, the symmetry breaking even calls into question the utility of the F/D parameter,[20] and so Refs. 17 and 21 set up their analyses without direct reference to F/D. Translating their work into F/D, one finds that the value subsumed in Ref. 17 is F/D = 0.56 consistent with that implicit in Ref. 21 and, within errors, with the fitted value in Ref. 22. Reference 23 obtained an even smaller value of F/D = 0.545 ± 0.02. Recent improvements in the Σn beta decay data, in particular, may raise F/D to 0.58 (Ref. 24), but nowhere as high as the 0.63 used previously.

The magnitudes for Δs implied by these values for F/D are

$$F/D = 0.548 \pm 0.01 \rightarrow \Delta s = -0.12 \pm 0.06 \text{ Ref. 23} \quad (51)$$

$$F/D = 0.58 \pm 0.01 \rightarrow \Delta s = -0.15 \pm 0.08 \text{ Ref. 24} \quad (52)$$

Thus we see that the magnitude of the (negative) strange polarization may be only half as big as that previously assumed. The QCD-corrected value of the Ellis-Jaffe sum rule falls from 0.19 (the cited value when F/D = 0.63) to 0.17 if F/D = 0.56, thereby reducing the statistical significance of the much-advertised failure of this sum rule.

What independent information exists on Δs? Elastic neutrino-proton scattering can, in principle, probe this quantity,[32] and a fit to these data give

$$\Delta s = -0.15 \pm 0.09 .$$

Note that this agrees with the revised value in the present paper arising from the smaller F/D and the revised EMC integral (Ref. 29).

One should also be aware that the neutrino experiment is also consistent with $\Delta s = 0$ which, in advance of the controversial EMC experiment, was the expectation.

Flavor-changing weak interactions, such as neutron beta decay, can yield

$$\frac{g_A}{g_V} \simeq 1.25 = \Delta u - \Delta d ,$$

while the zero momentum limit of $\nu p \to \nu p$ can probe

$$\tilde{g}_A(0) = \Delta u - \Delta d - \Delta s \equiv \left(\frac{g_A}{g_V}\right)\left(1 - \frac{\Delta s}{1.25}\right),$$

and so a difference between $\tilde{g}_A(0)$ and g_A/g_V can, after radiative corrections, reveal nonzero Δs. (Our $\Delta s \equiv 1.25\eta$ of Ref. 32.)

A practical problem is that $\nu p \to \nu p$ is detected by proton recoil and so an extrapolation to $q^2 = 0$ is needed. One fits the $q^2 \neq 0$ data with a form factor, in essence

$$\frac{1-\Delta s/1.25}{\left(1 + Q^2/M_A^2\right)^2},$$

where M_A is a mass scale to be fitted. Other experiments have determined this to have the value $M_A = 1.032 \pm 0.036$ GeV. If one fixes M_A to equal the world average, then $\Delta s = -0.15 \pm 0.09$; hence the claim to support the nonzero strange polarization. However, Ref. 32 also makes another, less well-advertised, fit. They constrain $\Delta s = 0$ and find that in this case $\Delta s = 0$; $M_A = 1.06 \pm 0.05$ GeV. Thus, one sees that $\Delta s = 0$ yields M_A consistent with the world average and hence is equally acceptable as a solution. The crucial statement in Ref. 32 is that "M_A and $\eta(\Delta s)$ are strongly correlated". Thus, Ref. 32 does not require $\Delta s < 0$ and thereby does not necessarily lend support to those who desire $\Delta s \neq 0$. Thus the question of the magnitude of the (strange) sea polarization is open. It is likely to be significantly nearer to zero than is being assumed in much of the current literature. Some of the inferences claimed from the EMC polarization data may need re-evaluation therefore. In particular, there need be no conflict with perturbative quantum chromodynamics.[30]

Polarized Gluons?

It has recently been realized[8] that the perturbative QCD correction to the singlet part of $g_1^P(x)$ effectively scales (to $O(\alpha_s^2)$) and may be important. This may be incorporated by replacing the Δq in Section 1 by $\Delta\tilde{q} \equiv \Delta q - \alpha_s/2\pi\ \Delta G$, where $\Delta G \equiv \int_0^1 dx\Delta g(x)$ and $\Delta G(x) = g_\uparrow(x) - G_\downarrow(x)$ is the polarized gluon distribution. This modifies the polarized lepton analysis, but cancels out in the expressions for (g_A/g_V) and does not enter the magnetic moment (Section 2) analysis.

One consequence is that there may be a continuity between the low-energy polarization revealed in constituent quarks (magnetic moments and spin dependence of resonance excitation) and the deep inelastic polarization.

First of all, we summarize the data on the Δq (or equivalently $\Delta\tilde{q}$) from the various (g_A/g_V).

If we assume $SU(3)_F$ symmetry in the sense that $s(\Sigma^+) \equiv d(P)$, then we may write the various g_A in terms of F, D, or Δq as follows:

g_A	F,D	$\Delta q^{(P)}$	Data
np	F + D	$\Delta u - \Delta d$	1.26 ± 0.005
Λp	$F + \frac{1}{3} D$	$\frac{1}{3}(2\Delta u - \Delta d - \Delta s)$	0.72 ± 0.02
ΞΛ	$F - \frac{1}{3} D$	$\frac{1}{3}(\Delta u + \Delta d - 2\Delta s)$	0.25 ± 0.05
Σn	F − D	$\Delta d - \Delta s$	-0.33 ± 0.02

Thus $F/D \equiv (\Delta u - \Delta s)/(\Delta u + \Delta s - 2\Delta d)$. Extracting the individual contributions involves a correlated fit. The EMC values, corrected for F/D, become $\tilde{\Delta} u = 0.80 \pm 0.06$, $\tilde{\Delta} d = -0.45 \pm 0.06$, and $\tilde{\Delta} s = -0.15 \pm 0.06$. One possibility is that $\Delta s = 0$, so that $\tilde{\Delta} s = -\alpha/2\pi \, \Delta G$. In this case, we obtain for

$$\Delta u \equiv \tilde{\Delta} u - \tilde{\Delta} s = 0.95 \pm 0.06 \qquad (53)$$

$$\Delta d \equiv \tilde{\Delta} d - \tilde{\Delta} s = -0.30 \pm 0.06 \qquad (54)$$

It is interesting to note that these values are consistent with those extracted from the magnetic moments (Eq. (4)) viz

$$\Delta u_v = 1, \ \Delta d_v = -0.25$$

The proton helicity is given by

$$\frac{1}{2} = \frac{1}{2} \Delta q + (\Delta G + L_z) \qquad (55)$$

Hence $\Delta G = 3.5$ and $L_z = -3.35$ at $Q^2 \simeq 10$ GeV^2. As one devalues to lower Q^2, $d/dQ^2(\Delta G + L_z) = 0$, and the individual contributions fall. It is an open question whether the "passive" L_z in the constituent model (i.e., the dilution of S_z due to relativistic spinors) at low Q^2 provides a consistent picture between constituent spin polarization and "parton" polarization.

Ellis et al.[31] suggest that the modification to $\Delta q(x)$ be driven by evolution

$$\tilde{\Delta} q(x) = \Delta q(x) - \int_x^1 \frac{dy}{y} \Delta G(y) \sigma(x/y)$$

where $\sigma(Z)$ is the cross section for $\gamma^* g \to q\bar{q}$. If so, then $g_1^p(x \to 0) < 0$, the crossover from positive to negative moving to smaller x values as Q^2 increases.

It is tantalizing that such a picture may already be manifested at low Q^2 in the resonance region. It is well known that the prominent $D_{13}(1520)$ and $F_{15}(1690)$ resonances are excited dominantly in $\sigma_{3/2}$ when $Q^2 = 0$, but in $\sigma_{1/2}$ for $Q^2 \neq 0$. The change in helicity structure,[38] or change in sign of $g_1^p(Q^2)$, occurs at $Q^2 \simeq 0.4$ GeV^2 for D_{13} and $Q^2 \simeq 0.7$ GeV^2 for F_{15}. It is amusing that these correspond to $x \simeq 0.2$, and so Bloom-Gilman duality may approximately hold true even for polarized leptoproduction, with a Q^2 dependence to the x_c where $g_1^P(x_c) = 0$. The first resonance $P_{33}(1236)$ sits on top of an S-wave background; the relative Q^2 dependences are not well known. However, perturbative QCD applied to resonance excitation suggests that this excitation may also change its character with Q^2 such that its contribution to $g_1^P(x)$ changes sign at $x \to 0$. It will be interesting at CEBAF to verify if the resonance region indeed matches onto the deep inelastic, and at high energy labs to verify whether $g_1^P(z \to 0) < 0$.

References

1. G. Baum et al., Phys. Rev. Lett. 51:1135 (1983); V. W. Hughes and J. Kuti, Ann. Rev. Nucl. Part. Sci. 33:611 (1983).
2. J. Ashman et al. (EMC), Phys. Lett. B206:364 (1983; V. W. Hughes et al., Phys. Lett. B212:511 (1988).
3. F. E. Close, "Introduction to Quarks and Partons," Academic, New York (1979) Chap. 13.
4. J. Kuti and V. Weisskopf, Phys. Rev. D4:3418 (1971).
5. J. D. Bjorken, Phys. Rev. D1:1976 (1970).
6. J. Ellis and R. L. Jaffe, Phys. Rev. D9:1444 (1984).
7. J. Ellis, R. A. Flores, and S. Ritz, Phys. Lett. B194:493 (1987).
8. G. Altarelli and G. G. Ross, Phys. Lett. B212:391 (1988); R. Carlitz, J. Collins, and A. Mueller, Phys. Lett. B214:229 (1988).
9. L. M. Sehgal, Phys. Rev. D10:1663 (1974).
10. S. J. Brodsky, J. Ellis, and M. Karliner, Phys. Lett. B206:309 (1988).
11. T. Barnes and F. E. Close, Phys. Lett. 128B:277 (1983; F. Wagner, Proc. XVI Rencontre de Moriond (1982), ed. J. Tranthanhvan.
12. C. Carlson and J. Milana, College of William & Mary report, WM-89-101. The role of gluon exchange is discussed by H. Hogassen and F. Myhrer, Phys. Rev. D37:1950 (1988).
13. For example, p. 117 in Ref. 3.
14. This comes from combinations of g_A/g_V for np with that for Λp, Σn, and ΞΛ; see Eq. (11) in Ref. 17.
15. B. A. Campbell, J. Ellis, and R. A. Flores, CERN-TH-5342/89.
16. M. Bourquin et al., Z. Phys. C21:27 (1983).
17. F. E. Close and R. G. Roberts, Phys. Rev. Lett. 60:1471 (1988).
18. M. Aguilar-Benitez et al. (Particle Data Group) Phys. Lett. B204:1 (1988).
19. S. Hsueh et al. (E715 collaboration) Phys. Rev. D38:2056 (1988).
20. H. J. Lipkin, Phys. Lett. 214B:429 (1988).
21. M. Anselmino, B. Ioffe, and E. Leader, Santa Barbara ITP report (1988) unpublished.
22. D. Kaplan and A. Manohar, Nucl. Phys. B310:527 (1988).
23. J. Donoghue, B. Holstein, and S. Klint, Phys. Rev. D35:934 (1987).
24. A. Beretvas, private communication; Z. Dziembowski and J. Franklin, Temple University, Philadelphia report TUHE-89-11 (1989).
25. R. Jaffe and A. Manohar, MIT report, MIT-CTP-1706 (1989) (but note that this inputs outdated neutron lifetime, which artificially increases the errors).
26. F. E. Close, Phys. Lett. 43B:422 (1973).
27. G. Farrar and D. Jackson, Phys. Rev. Lett. 35:1416 (1975).
28. F. E. Close and A. W. Thomas, Phys. Lett. B212:227 (1988).
29. EMC Collaboration, CERN preprint, CERN EP/89-73, June 1989.
30. F. E. Close, Phys. Rev. Letts. (in press).
31. J. Ellis, M. Karliner, and C. Sachrajda, CERN-TH-5471/89.
32. L. Ahrens et al., Phys. Rev. D35:785 (1987).
33. M. Frankfurt et al. (private communication).

34. F. E. Close and R. Milner, Oak Ridge National Laboratory report (1989).
35. Z. Dziembowski et al., Phys. Rev. D 39:3257 (1989).
36. G. Altarelli et al., Nucl. Phys. B69:531 (1974).
37. R. Carlitz and J. Kaur, Phys. Rev. Lett. 38:673 (1977).
38. See, e.g., V. Burkerdt, Research Program at CEBAF (Report of 1985 study group) ed. F. Gross (CEBAF, Newport News, 1986).