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# Flavor Asymmetry in the Light-Quark Sea of the Nucleon

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Parton distributions with an excess of down quarks over up quarks in the sea can reproduce data on the structure functions $F_2^{\mu p}$, $F_2^{\mu n}$, and $x\mathcal{F}_3^{\nu N}$. A model calculation in chiral field theory shows how an up-down asymmetry can arise from the dissociation of a quark into a quark plus a pion within the nucleon. This effect is large enough to account for the Gottfried sum rule defect reported by the New Muon Collaboration. Similar calculations may advance the understanding of other quasistatic properties of hadrons.

Recent measurements of the structure function ratio $F_2^{\mu n}/F_2^{\mu p}$ by the New Muon Collaboration at CERN[1] have renewed interest in the structure of the light-quark sea of the nucleon. In the quark-parton model, the integral

$$I_G(x_0, x_1; Q^2) = \int_{x_0}^{x_1} dx \, \frac{[F_2^{\mu p}(x, Q^2) - F_2^{\mu n}(x, Q^2)]}{x} \tag{1}$$

can be expressed in terms of the parton distribution functions $q_i^{(p,n)}(x, Q^2)$ of the proton and neutron as

$$I_G(x_0, x_1; Q^2) = \int_{x_0}^{x_1} dx \sum_i e_i^2 \left[q_i^{(p)}(x, Q^2) + \bar{q}_i^{(p)}(x, Q^2) - q_i^{(n)}(x, Q^2) - \bar{q}_i^{(n)}(x, Q^2)\right] \quad , \tag{2}$$

where $e_i$ is the charge (in units of the proton charge) of a quark of flavor $i$. Isospin invariance relates the parton distribution functions of the proton and neutron through the interchange $u \leftrightarrow d$. Separating the quark distributions of the proton into valence and sea contributions by writing $q_i = q_{iv} + \bar{q}_i$, we obtain

$$I_G(x_0, x_1; Q^2) = \frac{1}{3}\int_{x_0}^{x_1} dx \left[u_v(x, Q^2) - d_v(x, Q^2)\right] + \frac{2}{3}\int_{x_0}^{x_1} dx \left[\overline{u}(x, Q^2) - \overline{d}(x, Q^2)\right] \quad . \tag{3}$$

The Gottfried sum rule (GSR),[2]

$$I_G(0, 1; Q^2) = \frac{1}{3} \quad , \tag{4}$$

is obtained by noting that the first term in (3) becomes simply the difference between the number of up and down valence quarks in a proton, times one-third, while the second term vanishes for a $u$-$d$ flavor-symmetric sea.[3]

Using their own measurements of $F_2^{\mu n}/F_2^{\mu p}$ and the absolute deuteron structure function $F_2^{(e,\mu)d}$ from a parametrization of earlier measurements, the New Muon Collaboration[1] has determined $I_G(0.004, 0.8; Q^2 = 4\ \mathrm{GeV}^2) = 0.227 \pm 0.007$ and estimated

$$I_G(0, 1; Q^2 = 4\ \mathrm{GeV}^2) = 0.240 \pm 0.016 \quad . \tag{5}$$

This result implies a substantial violation of the Gottfried sum rule (4) derived under the assumption that the light-quark sea is flavor-symmetric.[4,5]

The GSR defect invites two interpretations. The contribution from the unmeasured region of small $x$, $\int_0^{x_0} dx \left[F_2^{\mu p} - F_2^{\mu n}\right]/x$, could make up the defect. Martin, Stirling, and Roberts[6] have displayed structure functions with this property that provide satisfactory fits to most observables. However, preliminary data of Fermilab experiment E-665,[7] which cover the range $0.001 \lesssim x \lesssim 0.1$, suggest that the NMC extrapolation is not grossly in error. The other interpretation, which we examine in this Letter, is that the light-quark sea is flavor-asymmetric, i.e., that $\overline{u}(x) \neq \overline{d}(x)$.

All modern fits to nucleon structure functions incorporate the simplifying assumption that $\overline{u}(x) = \overline{d}(x)$.[8] In light of the NMC data it is reasonable to ask[9] whether an up-down asymmetry can be accommodated without disrupting the character of the fits. It is interesting to recall that early data on the GSR led Field and Feynman[10] to argue that the Pauli principle would inhibit the development of up quarks and antiquarks in the proton sea; they accordingly built an asymmetry into their schematic parton distributions. We follow their example and write a variant of the EHLQ[11] Set 1 parton distributions, with valence distributions adjusted to reproduce the large-$x$ behavior of $F_2^{\mu p}$ and $F_2^{\mu n}$. The modified distributions are characterized at $Q_0^2 = 5\ \mathrm{GeV}^2$ by

$$
\begin{aligned}
xu_v(x,Q_0^2) &= 2.406x^{0.60}(1-x^{1.4})^{3.1} \\
xd_v(x,Q_0^2) &= 2.144x^{0.70}(1-x^{1.2})^{4.8} \\
x\overline{u}(x,Q_0^2) &= 0.20(1-x)^{11.2} \\
x\overline{d}(x,Q_0^2) &= 0.20(1-x)^{5.8} \\
x\overline{s}(x,Q_0^2) &= 0.086(1-x)^{8.5} \\
xG(x,Q_0^2) &= (2.62+9.17x)(1-x)^{5.9} \\
\Lambda_{QCD} &= 200 \text{ MeV} \quad .
\end{aligned}
\tag{6}
$$

This parametrization is not the result of a global fit. It should not be taken as a replacement for the standard distributions.

The difference between the exponents of $(1-x)$ for the up and down antiquarks has been chosen to reproduce within errors the NMC estimate (5) for $I_G(0,1)$. As the integral of a flavor-nonsinglet quantity, $I_G(0,1)$ is independent of $Q^2$. In any event, the evolution of structure functions from $Q^2 = 4$ GeV$^2$, where the NMC results are presented, and $Q^2 = 5$ GeV$^2$, where the parton distributions are defined, is small.

The agreement between the parton distributions (6) and the NMC determinations of $I_G(x_0,1)$ may be judged from Figure 1. The modified EHLQ 1 parametrization (solid curve) yields $I_G(0,1) = 0.251$. In contrast, the HMRS(EB) parametrization (dotted curve), which also describes the data adequately, predicts that $I_G(x_0,1)$ rises rapidly at small $x_0$, so that $I_G(0,1) = 1/3$. We show in Figure 2 data and parametrizations for the difference between proton and neutron structure functions. Our modified distributions reproduce the trend of the data. The unconventional small-$x$ behavior of the HMRS(EB) distributions is apparent. The valence-quark distributions are tested by the large-$x$ behavior of the structure function $x\mathcal{F}_3^{\nu N}$, which is displayed in Figure 3. The parametrizations satisfy the baryon number sum rule, $\int_0^1 dx\mathcal{F}_3^{\nu N}(x) = 3$, whereas preliminary CCFR data[12] yield $2.66 \pm 0.08$. QCD corrections[13] to the structure function $x\mathcal{F}_3^{\nu N}$ would reduce the baryon number sum rule to $3(1-\alpha_s/\pi)$ and bring the overall normalization of the parametrization (6) into agreement with the data. Finally, we show in Figure 4 the ratio of neutron to proton structure functions, which rises smoothly to unity as $x \to 0$ in the modified EHLQ parametrization, but behaves unconventionally at small $x$ in the HMRS(EB) parametrization.

We conclude that parton distributions that differ only in detail from a standard set can reproduce adequately all the important features of the data. We have evolved the parametrization (6) in leading order in QCD,[11] and used the resulting distributions to evaluate a number of $\overline{p}p$ cross sections. We have found no observable differences for the familiar observables. Having satisfied ourselves that these parton distributions lead to no obvious contradictions, we now ask what might be the origin of a flavor asymmetry in the light-quark sea.

A number of authors[14-19] have attributed the apparent excess of $\overline{d}$ over $\overline{u}$ to the pion cloud around the nucleon. The dissociation $p \to n\pi^+$, it is argued, leads to an excess of $\overline{d}$ over $\overline{u}$ because $\pi^+$ contains one valence up quark and one valence down antiquark. This physical picture has merit, but its implementation is problematical. First, a thorough treatment would include not only the dissociation $p \to (n\pi^+, p\pi^0)$, but also $p \to (\Delta^0\pi^+, \Delta^+\pi^0, \Delta^{++}\pi^-)$, and perhaps dissociation into higher isobars as well. Second, the interaction of a virtual photon with the pion does not contribute[16] to $F_2^{\mu p} - F_2^{\mu n}$, because $F_2^{\pi^+} \equiv F_2^{\pi^-}$. Instead, the GSR defect arises entirely from the interaction of a virtual photon with the recoil nucleon or isobar the pion leaves behind.[20]

To compute $F_2^{\mu p}(x) - F_2^{\mu n}(x)$ therefore requires a knowledge of the parton distributions of the nucleon (which we have) or the isobar (which we lack). It is not necessary to know the structure functions of the pions. To estimate the integrated GSR defect in this picture—as opposed to the $x$-dependence of $F_2^{\mu p}(x) - F_2^{\mu n}(x)$—it is only necessary to know the (valence) quark composition of the hadrons. Estimates in the literature account for about half of the reported defect.[14,15,17,19]

The flavor asymmetry in the light-quark sea is an example of quasistatic properties of hadrons, determined by low-energy, nonperturbative phenomena but probed by hard interactions. These properties occupy an intermediate position between static properties—like the mass, magnetic moment, charge radius, and axial charge of a hadron—for which correlations among constituents are all-important, and asymptotic properties—like the partition of momentum among parton species as $Q^2 \to \infty$—that rely on the independence of constituents. Although the measurement of quasistatic properties depends on parton-model incoherence, the properties are determined by correlations among the constituents. The art in understanding quasistatic properties lies in identifying the important correlations and calculating their effect with controlled approximations.

An apt description of the important degrees of freedom at momentum scales relevant to hadron structure—quarks, gluons, and Goldstone bosons—is provided by the effective chiral quark theory formulated by Georgi and Manohar.[21] We will show that the fluctuation of valence quarks into quarks plus Goldstone bosons (specifically pions) is a plausible origin for the observed violation of the Gottfried sum rule. We argue that the important correlations for the flavor asymmetry of the light-quark sea are embodied in the existence of light pseudoscalar mesons.[22]

Before undertaking a full discussion of the effective chiral quark model, we present a toy model that captures the essential effects. A simplified description of a nucleon would contain only valence up and down quark distributions. In the chiral quark model, these valence quarks can emit a pion. Diagrams for the fluctuations $u \to (\pi^+ d, \pi^0 u)$ are shown in Figure 5. Let $a$ denote the probability for an up quark to turn into a down quark with the emission of a $\pi^+$ (containing a valence $u$ quark and a valence $\overline{d}$ antiquark), as depicted in Figure 5(a). We suppose that this fluctuation is small enough to be treated as a perturbation. The up quark can also emit a $\pi^0$, with valence $u\overline{u}$ and $d\overline{d}$ components, as shown in Figures 5(b) and (c). The final state resulting from pion emission by an up quark is

$$\begin{aligned} u &\rightarrow a\pi^+ + ad + \frac{a}{2}\pi^0 + \frac{a}{2}u \\ &= \frac{7a}{4}u + \frac{5a}{4}d + \frac{a}{4}\overline{u} + \frac{5a}{4}\overline{d} \quad . \end{aligned} \tag{7}$$

Isospin symmetry—the interchange $u \leftrightarrow d$—requires that the corresponding down-quark processes produce the final state

$$\begin{aligned} d &\rightarrow a\pi^- + au + \frac{a}{2}\pi^0 + \frac{a}{2}d \\ &= \frac{5a}{4}u + \frac{7a}{4}d + \frac{5a}{4}\overline{u} + \frac{a}{4}\overline{d} \quad . \end{aligned} \tag{8}$$

If the probability for a quark to emit a pion is $3a/2$, then the probability for a quark to do nothing is $1 - 3a/2$. The proton composition after one interaction is therefore $(2+7a/4)u + (1+11a/4)d + (7a/4)\overline{u} + (11a/4)\overline{d}$, which preserves the valence composition $u_v = (u - \overline{u}) = 2$ and $d_v = (d - \overline{d}) = 1$. The parton distributions for the neutron are obtained by isospin symmetry. We then have a GSR defect

$$\Delta I_G = \frac{2}{3}(\overline{u} - \overline{d}) = -\frac{2a}{3} \neq 0 \quad . \tag{9}$$

Chiral field theory ($\chi$FT) enables us to calculate the probability for a valence quark to fluctuate and to compute the $x$-dependence of the resulting parton distributions. We shall formulate the calculation generally and then specialize to $q \rightarrow \pi q'$ transitions.

The effective interaction Lagrangian is

$$\mathcal{L} = \overline{\psi}\,(i\mathcal{D}_\mu + V_\mu)\,\gamma^\mu\psi + ig_A\overline{\psi}A_\mu\gamma^\mu\gamma_5\psi + \ldots \tag{10}$$

where

$$\psi = \begin{pmatrix} u \\ d \\ s \end{pmatrix} \tag{11}$$

is a set of color-triplet Dirac quarks, and $\mathcal{D}_\mu = \partial_\mu + igG_\mu$ is the gauge-covariant derivative of QCD, with $G_\mu$ the gluon field and $g$ the strong coupling constant. The dimensionless axial coupling $g_A = 0.7524$ is determined from the axial charge of the nucleon.[21] The vector and axial-vector currents

$$\begin{pmatrix} V_\mu \\ A_\mu \end{pmatrix} = \frac{1}{2}\left(\xi^\dagger \partial_\mu \xi \pm \xi \partial_\mu \xi^\dagger\right) \tag{12}$$

are expressed in terms of the octet of Goldstone boson fields,

$$\Pi \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & K^0 \\ K^- & \overline{K}^0 & \frac{-2\eta}{\sqrt{6}} \end{pmatrix} \quad , \tag{13}$$

through the field

$$\xi = \exp\left(i\Pi/f\right) \quad , \tag{14}$$

where the pseudoscalar decay constant is $f \approx 93$ MeV. An expansion of the currents in powers of $\Pi/f$ yields $V_\mu = 0 + O\left(\Pi/f\right)^2$ and $A_\mu = i\partial_\mu \Pi/f + O\left(\Pi/f\right)^2$, so the effective interaction between Goldstone bosons and quarks becomes

$$\mathcal{L}_{\Pi q} = -\frac{g_A}{f}\overline{\psi}\partial_\mu \Pi \gamma^\mu \gamma_5 \psi \quad . \tag{15}$$

In chiral field theory, the kinematical dependence of the splitting function for finding a Goldstone boson of mass $M_\Pi$ carrying momentum fraction $z$ of a valence quark $q$, leaving a recoil quark $q'$ with momentum fraction $(1-z)$, is given by

$$P_{\Pi q' \leftarrow q}(z) = \frac{g_A^2}{f^2}\,\frac{z\left(m_q + m_{q'}\right)^2}{32\pi^2}\int_{-\Lambda^2}^{t_{\min}} dt\, \frac{\left[\left(m_q - m_{q'}\right)^2 - t\right]}{\left(t - M_\Pi^2\right)^2} \quad , \tag{16}$$

where $t$ is the square of the Goldstone-boson four-momentum and $t_{\min} = m_q^2 z - m_{q'}^2 z/(1-z)$. The quark masses $m_q$ and $m_{q'}$ are constituent masses appropriate below the scale of chiral symmetry breaking. The integral in (16) requires an ultraviolet cutoff because the effective field theory is not valid at arbitrarily high energies. It is conventional to use

$$\Lambda = \Lambda_{\chi \mathrm{SB}} \equiv 4\pi f \approx 1169 \text{ MeV} \tag{17}$$

as an estimate of the chiral symmetry breaking scale and the ultraviolet cutoff.

Fluctuation of a valence quark in the nucleon into a Goldstone boson and a recoil quark will induce changes in the nucleon's parton distributions. Following Altarelli and Parisi,[23] we can write the $\chi$FT contributions of an $SU(n_f)$ adjoint of Goldstone bosons as

$$\overline{q}_i^{(\chi\mathrm{FT})}(x) = \sum_{j=u,d}\sum_{k,\ell}\left(\delta_j^\ell \delta_i^k - \frac{1}{n_f}\delta_j^k \delta_i^\ell\right)^2 \int_x^1 \frac{dy}{y}\int_{x/y}^1 \frac{dz}{z}\,\overline{q}_i^{(\Pi)}\left(\frac{x}{yz}\right) P_{\Pi k \leftarrow j}(z)\, q_{vj}^{(N)}(y) \tag{18}$$

and

$$q_i^{(\chi\mathrm{FT})}(x) = \sum_{j=u,d} \sum_{k,\ell} \left( \delta_j^i \delta_\ell^k - \frac{1}{n_f} \delta_j^k \delta_i^\ell \right)^2 \int_x^1 \frac{dy}{y} \int_{x/y}^1 \frac{dz}{z} q_i^{(\Pi)}(\frac{x}{yz}) P_{\Pi k \leftarrow j}(z) q_{vj}^{(N)}(y) \qquad (19)$$

$$+ \sum_{j=u,d} \sum_{k,\ell} \left( \delta_j^\ell \delta_k^i - \frac{1}{n_f} \delta_j^i \delta_k^\ell \right)^2 \int_x^1 \frac{dy}{y} \int_{x/y}^1 \frac{dz}{z} \delta(1 - \frac{x}{yz}) P_{\Pi i \leftarrow j}(1-z) q_{vj}^{(N)}(y) \quad ,$$

where $q_i^{(\Pi)}$ and $\overline{q}_i^{(\Pi)}$ are the parton distributions of the Goldstone boson and the summations over $k$ and $\ell$ run over the number of quark flavors, $n_f$. The antiquarks arise only from the structure of the Goldstone boson, whereas quarks arise also from the recoil quark. With $n_f = 2$ we recover the quark counting summarized in (7) and (8).

We now specialize to the emission of pions, the lightest of the pseudoscalar mesons, and take $m_u = m_d = 350$ MeV$/c^2$. By integrating the splitting function (16) over possible values of the momentum fraction $z$, we can evaluate the probability for an up quark to emit a $\pi^+$ as

$$a = \frac{g_A^2 m_u^2}{8\pi^2 f^2} \int_0^1 dz \Theta\left(\Lambda - \tau(z)\right) z \left\{ \ln\left[\frac{\Lambda^2 + M_\pi^2}{\tau(z) + M_\pi^2}\right] + M_\pi^2 \left[\frac{1}{\Lambda^2 + M_\pi^2} - \frac{1}{\tau(z) + M_\pi^2}\right] \right\} \quad , \qquad (20)$$

where $\tau(z) = m_u^2 z^2/(1-z)$. With the standard choice of ultraviolet cutoff $\Lambda = \Lambda_{\chi SB}$, we find $a = 0.083$, which leads to $I_G(0,1) = (1-2a)/3 = 0.278$. Considering the simplicity of our model, we regard this as an encouraging result.

To study the $x$-dependence of the flavor asymmetry, we adjust the ultraviolet cutoff $\Lambda$ to better fit the observed GSR defect. The value $\Lambda = 1800$ MeV leads to $I_G(0,1) = 0.252$. For simplicity, and to verify that the $\chi$FT process can be regarded as a perturbation, we calculate the fluctuation of the proton valence-quark distributions (6). To compute $F_2^{\mu p}$ and $F_2^{\mu n}$, we choose the pion structure function[24]

$$x q_v^{(\pi^+)}(x, Q_0^2) = 0.99 x^{0.58} (1-x)^{1.26} \qquad (21)$$

and add the flavor-symmetric nucleon sea distributions of EHLQ Set 1,[11] $xu_s(x) = xd_s(x) = 0.182(1-x)^{8.54}$.

The results are shown as dashed curves in Figures 1 to 4. The $\chi$FT structure functions are in satisfactory agreement with the NMC data on $I_G(x_0, 1)$ and $F_2^{\mu p} - F_2^{\mu n}$. The nonsinglet structure function $x\mathcal{F}_3^{\nu N}$, shown in Figure 3, is slightly degraded from the input valence distribution shown as the solid curve. The small difference between the two shows that it is indeed reasonable to treat pion emission from valence quarks as a perturbation. Somewhat harder input valence distributions would be required to reproduce the observed dependence of $x\mathcal{F}_3^{\nu N}$ at large $x$. The $\chi$FT prediction for the ratio $F_2^{\mu n}/F_2^{\mu p}$ shows why it is difficult to find evidence of an excess of $\overline{d}$ over $\overline{u}$ without forming cross-section differences: the ratio $\overline{u}(x)/\overline{d}(x)$ is close to unity. We look forward to more extensive results from E-665 and to measurements at HERA that will extend our knowledge of the structure functions to smaller values of $x$.[25]

Our analysis can be extended to include the effects of the full $SU(3)$ chiral effective Lagrangian (15). An up or down valence quark can change into a strange quark by emitting a kaon. Since the initial state has no strangeness, the generated strange quark distributions satisfy $s(x) - \overline{s}(x) = 0$ and will not contribute to the GSR defect.

In the generalization to a $U(n_f)$ chiral field theory, emission of the $SU(n_f)$-singlet Goldstone boson cancels the $-(1/n_f)\times$Kronecker delta terms in (18) and (19). Up and down antiquarks are created in equal numbers, so the Gottfried sum rule is respected. It is instructive to see this explicitly for the case of $U(2)$ flavor symmetry by considering the toy model with degenerate $\pi^0 = (u\overline{u} - d\overline{d})/\sqrt{2}$ and $\eta = (u\overline{u} + d\overline{d})/\sqrt{2}$. The emission of $\pi^0$ and $\eta$ lead to the same final states, so the two paths must be added coherently. After one iteration, in the units introduced in the discussion surrounding (7) and (8), the composition of the proton would be $(2+3a)u + (1+3a)d + 3a(\overline{u}+\overline{d})$, which respects the GSR. For an $SU(3)$ octet of degenerate Goldstone bosons, the proton composition after one iteration is $(2+2a)u + (1+8a/3)d + 2a\overline{u} + (8a/3)\overline{d} + (10a/3)(s+\overline{s})$, which leads to a GSR defect $\Delta I_G = -4a/9$. In a $U(3)$ nonet theory, the proton composition would be $(2+3a)u + (1+3a)d + 3a(\overline{u}+\overline{d}+s+\overline{s})$, so that the GSR defect is cancelled. This cancellation does not occur to any significant degree in nature, because the $\eta$ and $\eta'$ mesons are not degenerate with the $\pi^0$. In the $\chi$FT picture, the GSR defect can thus be traced to the breaking of flavor-$SU(3)$ symmetry by the quark masses and the axial anomaly contribution to the $U(1)$-meson mass.

Chiral field theory may also yield insight into the distribution of spin in polarized nucleons. Fluctuation of a valence quark into a recoil quark and Goldstone boson will reduce the spin fraction carried by the valence quark. In the language of our toy model, this time for the full octet of Goldstone bosons, a positive-helicity valence up quark will produce a negative-helicity strange-quark distribution when it emits a kaon. Because the strange antiquark is a constituent of a spinless Goldstone boson, it will be unpolarized. The difference between strange quark and antiquark invalidates a key assumption needed to derive the polarized-leptoproduction sum rules:[26] the spins of partons in the sea are paired.

In the elementary quark model with $SU(6)$ wavefunctions, the composition of a polarized proton with positive helicity is $p_\uparrow = \frac{5}{3}u_\uparrow + \frac{1}{3}u_\downarrow + \frac{1}{3}d_\uparrow + \frac{2}{3}d_\downarrow$. Denoting as usual the probability for an up quark to fluctuate into a down quark and a $\pi^+$ as $a$, we can compute the total probability for a quark to fluctuate into a quark and a

Goldstone boson as $8a/3$, so the probability that nothing happens is $(1-8a/3)$. This model reproduces the observed GSR defect with $a = 0.21$, a rather large probability. The spin asymmetries of the three quark flavors after one iteration are

$$\begin{aligned}\Delta_u &= u_\uparrow - u_\downarrow = \frac{4}{3} - \frac{37a}{9} \approx 0.47 \\ \Delta_d &= d_\uparrow - d_\downarrow = -\frac{1}{3} - \frac{2a}{9} \approx -0.38 \\ \Delta_s &= s_\uparrow - s_\downarrow = -a \approx -0.21 \quad ,\end{aligned} \tag{22}$$

while there is no spin asymmetry for the antiquarks. The integrals of the spin-dependent structure functions are

$$\begin{aligned}\int_0^1 dx g_1^{(p)}(x) &= \frac{1}{2}\left(\frac{4}{9}\Delta_u + \frac{1}{9}\Delta_d + \frac{1}{9}\Delta_s\right) \\ &= \frac{5}{18} - \frac{53a}{54} \approx 0.07 \ ,\end{aligned} \tag{23}$$

$$\begin{aligned}\int_0^1 dx g_1^{(n)}(x) &= \frac{1}{2}\left(\frac{1}{9}\Delta_u + \frac{4}{9}\Delta_d + \frac{1}{9}\Delta_s\right) \\ &= -\frac{a}{3} \approx -0.07 \quad ,\end{aligned} \tag{24}$$

and the total spin carried by quarks in a polarized nucleon is

$$I_0 = \Delta_u + \Delta_d + \Delta_s = 1 - \frac{16a}{3} \approx -0.12 \quad . \tag{25}$$

In contrast to the GSR case, the order-$a$ corrections do not vanish in the $U(3)$ limit.

In the toy $SU(2)$ chiral model, for which $a = 0.14$ reproduces the GSR defect, we find $\Delta_u = 4/3 - 7a/3 \approx 1.007$, $\Delta_d = 1/3 - 2a/3 \approx -0.427$, $\Delta_s = 0$, which leads to numerical results for the sum rules that are close to the standard values.[26] We would expect a complete calculation including physical masses for the pseudoscalars to give results in between these two cases.

These simple considerations show that the same dynamical correlations that can explain the GSR defect may also help us to understand the fraction of the proton's spin carried by quarks. In this picture, it is natural for the sea to reflect the nucleon polarization.

In conclusion, flavor asymmetry for the sea quark distributions arises naturally through isopsin-conserving interactions if the quark distributions are generated by processes that are correlated with the total isospin of the nucleon. We have accounted for these correlations using the chiral quark model. The resulting quark distributions give a value for the Gottfried sum rule in agreement with experiment without significantly affecting other well-measured processes. What this model shows to be unnatural—or at least unnecessary—is the idealization that the sea carries no knowledge of the flavor of the valence quarks.

The Gottfried Sum Rule is representative of a class of observables that measure the quasistatic properties of hadrons. Quasistatic properties are determined from amplitudes involving the perturbative degrees of freedom, quarks and gluons, for which strong (nonperturbative) correlations between these degrees of freedom enter in an essential way. A related example of an observable depending on quasistatic properties is the spin sum rule for a polarized nucleon.

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FIG. 1. Gottfried sum rule integral $I_G(x_0, 1)$ vs. $x_0$, measured by the New Muon Collaboration, Ref. 1. Solid curve: modified EHLQ structure functions, Eq. (6); dotted curve: HMRS(EB) structure functions, Ref. 6; dashed curve: chiral field theory.

FIG. 2. NMC measurements of the difference $F_2^{\mu p}(x) - F_2^{\mu n}(x)$ vs. $x$, Ref. 1. Solid curve: modified EHLQ structure functions, Eq. (6); dotted curve: HMRS(EB) structure functions, Ref. 6; dashed curve: chiral field theory.

FIG. 3. Preliminary data of the CCFR Collaboration on the structure function $x\mathcal{F}_3^{\nu N}(x)$ at $Q^2 = 3$ GeV$^2$, Ref. 12. Solid curve: modified EHLQ structure functions, Eq. (6); dotted curve: HMRS(EB) structure functions, Ref. 6; dashed curve: chiral field theory.

FIG. 4. NMC measurements (•) of the ratio $F_2^{\mu n}(x)/F_2^{\mu p}(x)$ vs. $x$, Ref. 1. Solid curve: modified EHLQ structure functions, Eq. (6); dotted curve: HMRS(EB) structure functions, Ref. 6; dashed curve: chiral field theory. Also shown (■ and $\times$) are preliminary data of Fermilab Experiment E-665, from Ref. 7.

FIG. 5. Fluctuations of a valence up quark in chiral field theory.

Figure 1

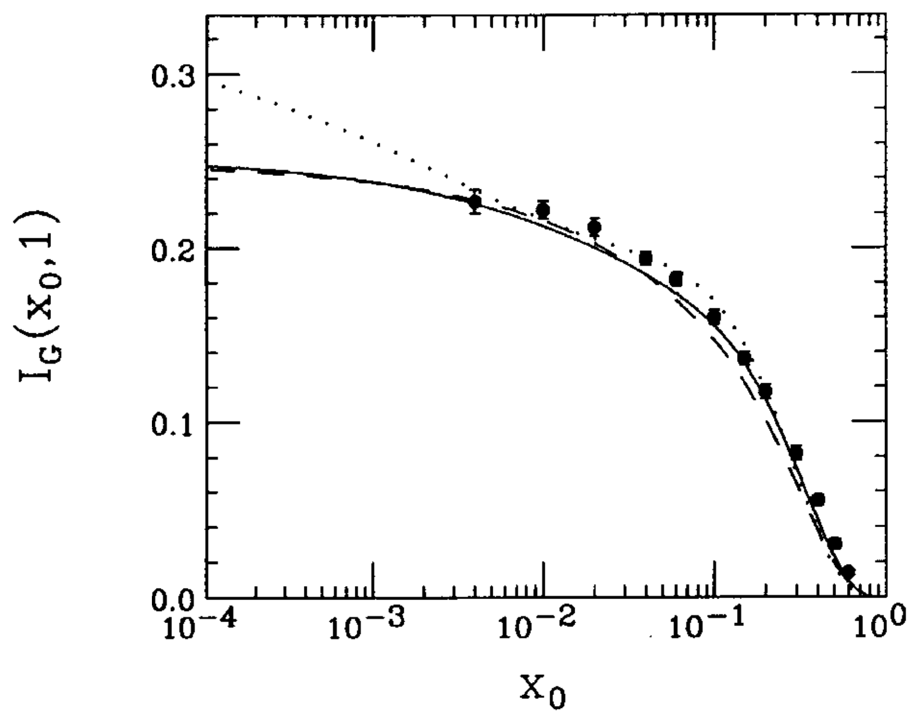
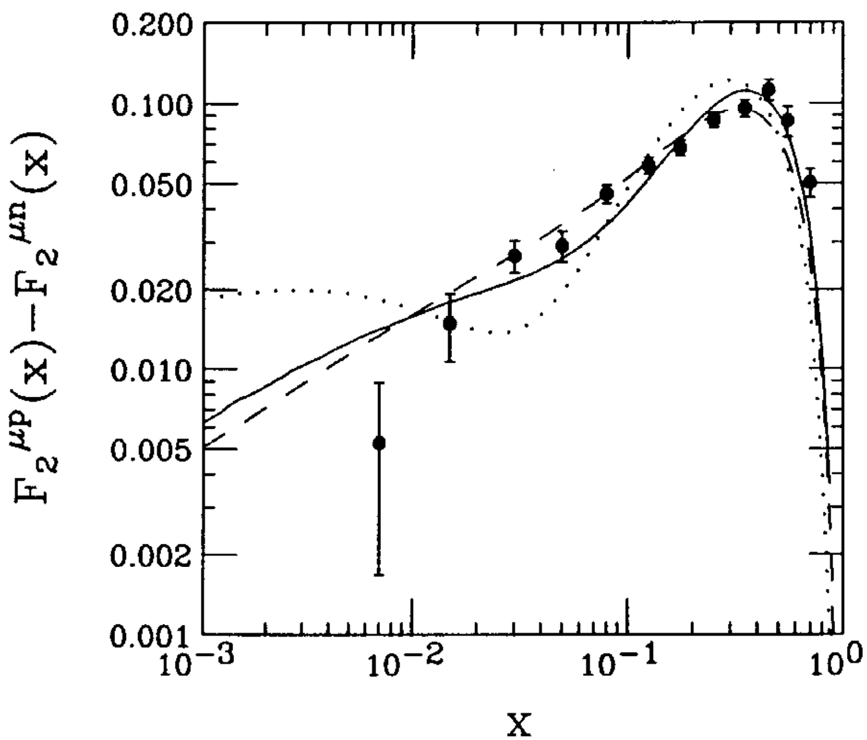


Figure 2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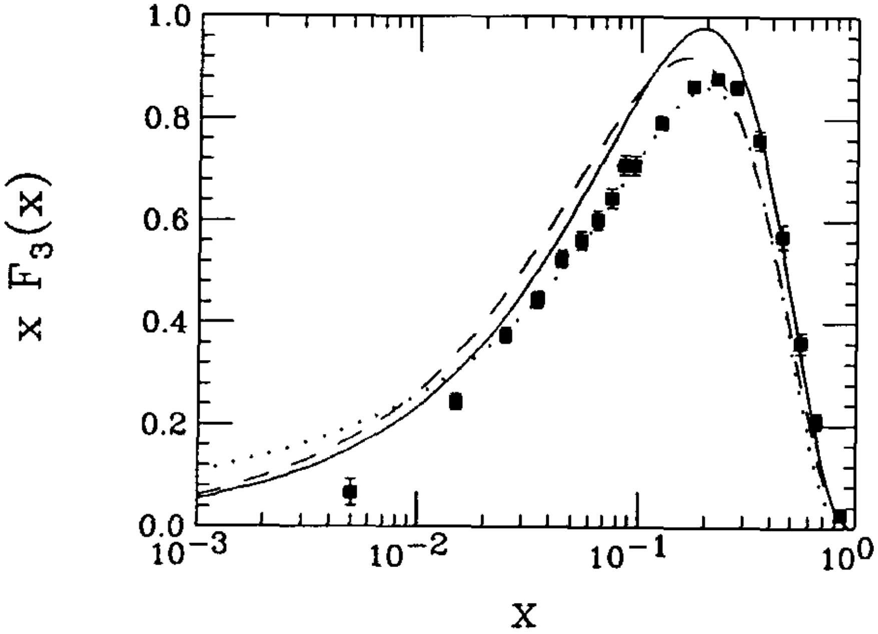
0.100
0.050
0.020
0.010
0.005
0.002
0.001
$10^{-3}$
$10^{-2}$
$10^{-1}$
$10^{0}$
X
$F_2^{\mu p}(x)-F_2^{\mu n}(x)$

Figure 3



FERMILAB-PUB-91/272-T

Figure 4

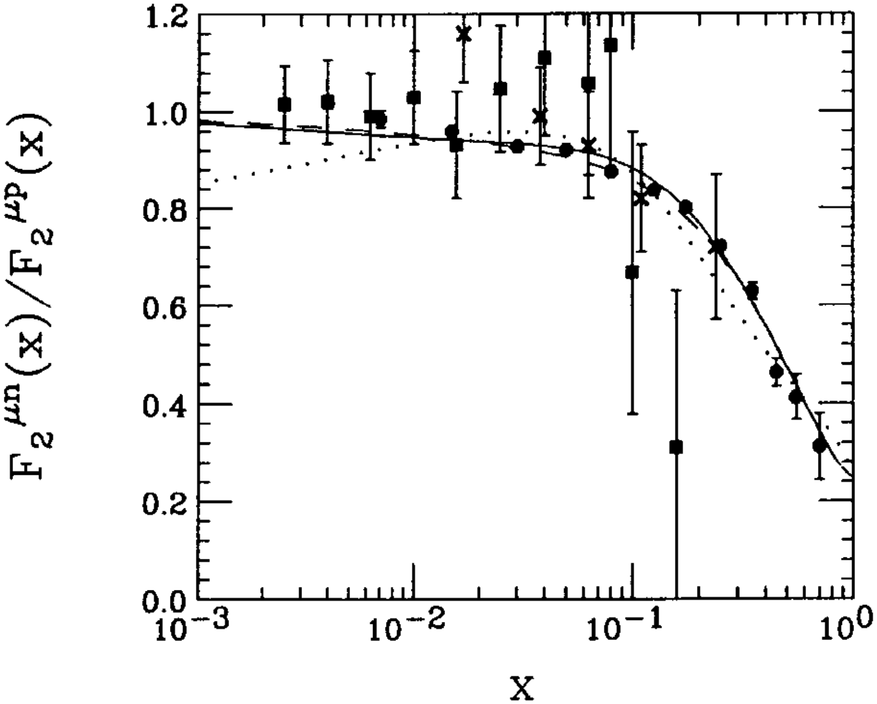


Figure 5