To obtain an upper limit on the branching ratio we use the Poisson distribution

$$P_u(n) = (u^n/n!)\exp(-u),$$

where $u = \epsilon R + b$, $n = 1$ is the observed number of events and, $b = 1.74$ is the expected number of background events. ϵ is the effective number of trials

$$\epsilon = (\Omega/4\pi)\epsilon_c \epsilon_E \epsilon_T N_\pi = 5.96\times 10^8.$$

At a 90% confidence level[12,13] we find an upper limit for the branching ratio for the $\mu^+ \to e^+\gamma$ decay,

$$R_{\mu e\gamma} < 3.6\times 10^{-9}.$$

We would like to thank the staff at TRIUMF for their help in mounting this experiment. This work was supported in part by the National Research Council of Canada.

[a]On sabbatical leave at Physics Department, University of British Columbia, Vancouver, British Columbia V6T 1W5, Canada.

[1]S. Parker, H. Anderson, and C. Rey, Phys. Rev. 133, B768 (1964).

[2]S. M. Korenchenko *et al.*, Zh. Eksp. Theor. Fiz. 70, 3 (1976) [Sov. Phys. JETP 43, 1 (1976)].

[3]D. A. Bryman, M. Blecher, K. Gotow, and R. J. Powers, Phys. Rev. Lett. 28, 1469 (1972).

[4]See for example T. P. Cheng and L.-F. Li, Phys. Rev. Lett. 38, 381 (1977). Also J. D. Bjorken, K. Lane, and S. Weinberg, SLAC Report No. SLAC-PUB-1925 (to be published); their Ref. 2 contains an extensive list of the recent literature on the subject.

[5]G. Opat, Monte Carlo Program, SIMUL 8, TRIUMF Report No. TRI-77-4 (to be published).

[6]G. Källén, in *Springer Tracts in Modern Physics*, edited by G. Höhler (Springer, Berlin, 1968), Vol. 46, p. 96.

[7]Pileup events consisting of two charged particles entering either NaI detector within 400 nsec (twice our integration time for the linear pulses) were detected by the two large plastic counters 9 and 10. Because of the relatively low stopping rate (2×10^5 sec^{-1}), only 1% of the NaI signals were affected by pileup and therefore this effect contributed only a small broadening of our experimental resolution.

[8]J. Spuller *et al.*, Phys. Lett. 67B, 479 (1977).

[9]P. Depommier *et al.*, Nucl. Phys. B4, 189 (1968).

[10]P. Depommier *et al.*, Phys. Lett. 7, 285 (1963). The value R_γ quoted has been calculated using the value of γ (ratio of axial-vector to vector form factor) of P. Depommier *et al.* corrected for the latest value of the π^0 lifetime [$\tau_{\pi^0} = (0.828\pm 0.057)\times 10^{-16}$ sec].

[11]Using the experimentally measure NaI line shape, we estimate that 68% of all particles with incident energy, E, will deposit an energy in the NaI greater than $E - 0.5\Delta E$ (FWHM).

[12]Particle Data Group, Rev. Mod. Phys. 48, 44 (1976).

[13]*Handbook of Mathematical Functions*, edited by M. Abramovitz and I. Stegun (National Bureau of Standards, Washington, D. C., 1964), Vol. 55.

Whirlpools in the Sea: Polarization of Antiquarks in a Spinning Proton

F. E. Close
Rutherford Laboratory,[a] Chilton Didcot, Oxfordshire OX11 0QX, England, and Argonne National Laboratory, Argonne, Illinois 60439

and

Dennis Sivers
Argonne National Laboratory, High Energy Physics Division, Argonne, Illinois 60439
(Received 3 August 1977)

We argue that the sea of virtual quark-antiquark pairs in a proton is polarized. Quantum chromodynamics allows a valence quark to emit a gluon which then produces pairs. If the parent quark in this process exhibits a $(1-x)^3$ behavior as $x \to 1$ then the antiquark with helicity the same as the original quark will have a leading $(1-x)^5$ distribution. Implications for the counting rules and for polarized Drell-Yan annihilation are discussed.

In our current picture of the proton it contains not only three valence quarks but also an indefinite number of color-SU(3) gauge bosons and quark-antiquark pairs. We usually refer to these extra objects as "gluons" and "the sea", respectively. It is a widely held belief that the sea is unpolarized.[1]

The simplest nucleon configuration consists of just the valence quarks. The sea can be generated by the valence quarks emitting gluons which then produce pairs. Explicit calculation[2] shows that the pairs produced in this manner carry only a small fraction of the nucleon's momentum, in agreement with data from deep-inelastic scattering.

Using the rules of quantum chromodynamics (QCD) we explicitly study the production of the sea. If valence quarks have probabilities $(1-x)^3$ as $x \to 1$, then the leading component of the sea is found to contain q's and $\bar{q}$'s with helicities predominantly the same as the nucleon and with a $(1-x)^5$ behavior as $x \to 1$. There will, in general, be a net polarization of the $\bar{q}$'s, which can be tested by producing massive lepton pairs from polarized proton beams and targets. Under certain reasonable hypotheses, we shall see, the observable asymmetry in such experiments can be fairly large.

First we show how this spinning sea might arise. For simplicity we will illustrate this for a Δ with all three valence quarks spinning parallel to its direction of motion. We suppress the flavor degrees of freedom and write the wave function in terms of "dressed" or constituent quarks in the form

$$|\Delta^{(+3)}\rangle = |q^{(+)}q^{(+)}q^{(+)}\rangle. \tag{1}$$

Our results follow if we assume the relationship between the dressed constituents and the parton quarks, gluons, and the sea to be approximately given by QCD diagrams such as those shown in Fig. 1. Because of the γ_μ coupling at the quark-gluon vertex the helicity of the original quark is conserved if we neglect the quark mass. The helicity of the gluon is then determined by angular-momentum conservation and hence depends, in particular, on the fraction of the quark momentum carried off. A $q\bar{q}$ pair produced by the gluon will also inherit a memory of its valence-quark progenitor.

In the case of the polarized Δ above, if the distribution of positive-helicity quarks near $x=1$ is given by the constituent-counting rules,[3-5]

$$q_\Delta^{(+)}(x) \sim c^{(+)}(1-x)^3, \tag{2}$$

then we can use the diagrams in Fig. 1 to show that the leading behavior near $x=1$ of the vector gluon distributions depends on helicity,

$$\begin{aligned} V_\Delta^{(+)}(x) &\sim v^{(+)}(1-x)^4, \\ V_\Delta^{(-)}(x) &\sim v^{(-)}(1-x)^6, \end{aligned} \tag{3}$$

while the leading behavior of the antiquark distributions is given by

$$\begin{aligned} \bar{q}_\Delta^{(+)}(x) &\sim s^{(+)}(1-x)^5, \\ \bar{q}_\Delta^{(-)}(x) \approx q_\Delta^{(-)}(x) &\sim s^{(-)}(1-x)^7. \end{aligned} \tag{4}$$

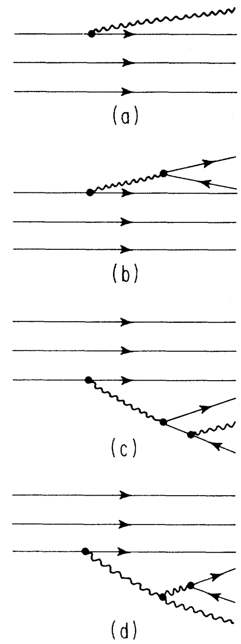

FIG. 1. Schematic diagrams which suggest how the sea is generated in QCD. A valence quark emits a gluon (a) which produces a pair (b) yielding the distributions (8) and (11). More complicated diagrams such as (c) and (d) do not affect the leading component of the $\bar{q}$ distributions and the polarization.

Thus an antiquark detected at large x in a polarized Δ is more likely to have the same helicity as the parent hadron than the opposite helicity. To work out more precisely what we expect, it is convenient to take a simple concrete model. Suppose we parametrize the distributions of the original constituent quarks in the form

$$q_\Delta^{0(+)}(x, k_T) = 60 f(k_T) x(1-x)^3, \tag{5}$$

with $\int d^2k_T f(k_T) = 1$. We adopt the assumption of factorization between the x and k_T dependence of the distribution for computational convenience. We are aware that this factorization might not be realized in nature[6] but our results do not depend sensitively on it and the assumption enables us to suppress k_T integrations. The behavior near $x=1$ of the distribution (5) is chosen to agree with the usual constituent-counting rules. The behavior near $x=0$ can also be understood in the framework of these rules since when $x \to 0$ the other two quarks considered as a unit must carry all the momentum. The parametrization (5) is normalized so that these dressed constituent quarks carry all of the Δ's momentum. We now want to investigate how the gluons and the sea are generated dynamically from the constituent quarks. We assume that the fluons are generated by diagrams like that in Fig. 1(a). We can calculate the probability that one of the valence quarks will undergo "bremsstrahlung" to give a gluon of momentum x:

$$V_\Delta^{(j)}(x)=\int_x^1 (dy/y)P_{q^{(+)}\to V^{(j)}}(x/y)q_\Delta^{(+)0}(y), \qquad (6)$$

where, in the interest of compact notation, we do not explicitly display the integration over transverse momentum. By analysis of the leading perturbation-theory diagrams, we can calculate[7]

$$P_{q^{(+)}\to V^{(+)}}(z)=\beta/z,$$
$$P_{q^{(+)}\to V^{(-)}}(z)=\beta(1-z)^2/z, \qquad (7)$$

where β contains factors related to color averaging, the effective strong-interaction coupling constant, and the k_T integrations. Using (7) and (5) we can integrate (6) to get

$$V_\Delta^{(+)}(x)=3\beta(1/x)(1-x)^4[1+4x],$$
$$V_\Delta^{(-)}(x)=60\beta\frac{1}{x}\left[x^2\left(-\ln x-\sum_{k=1}^{3}\frac{(1-x)^k}{k}\right)+\frac{(1-x)^4}{4}(1-2x)-\frac{(1-x)^5}{5}\right]. \qquad (8)$$

Expanding the logarithm in (8) displays the leading $(1-x)^6$ behavior as $x\to 1$.

These distributions are displayed graphically in Fig. 2. In order to obtain the antiquarks, we need to know the probability for a gluon to give a fraction z of its momentum to an antiquark. This is found to be[7]

$$P_{V^{(+)}\to \bar q^{(+)}}(z)=P_{V^{(-)}\to \bar q^{(-)}}(z)=\beta\tfrac{1}{2}z^2, \quad P_{V^{(+)}\to \bar q^{(-)}}(z)=P_{V^{(-)}\to q^{(+)}}(z)=\beta\tfrac{1}{2}(1-z)^2, \qquad (9)$$

and we can write

$$\bar q_\Delta^{(+)}(x)=\sum_{i=+,-}\int_x^1\frac{dy}{y}P_{V^{(i)}\to \bar q^{(+)}}(x/y)\,V_\Delta^{(i)}(y). \qquad (10)$$

Using (8) and (9) we integrate and find

$$\bar q_\Delta^{(+)}(x)=(3\beta^2/6x)[1-30x^2-10x^3+45x^4-6x^5-60x^3\ln x]$$
$$+(3\beta^2/6x)[1+107\tfrac{1}{2}x-283\tfrac{1}{3}x^2+170x^3+5x^4-\tfrac{1}{6}x^5+(30x-40x^2-60x^3)\ln x-60x^2\ln^2 x],$$
$$\bar q_\Delta^{(-)}(x)=2\cdot(3\beta^2/6x)[1-15x-80x^2+80x^3+15x^4-x^5-60x^2(x+1)\ln x]. \qquad (11)$$

Expanding the logarithms we can verify that the leading behaviors near $x=1$ of the terms in square brackets are $(1-x)^5$ ($\bar q^+$ from V^+), $(1-x)^9$ ($\bar q^+$ from $V^{(-)}$), and $(1-x)^7$ ($\bar q^{(-)}$ from $V^{(-)}$ and $\bar q^{(-)}$ from $V^{(+)}$), respectively, so that Eqs. (4) are satisfied. These distributions are also plotted in Fig. 2. Note that when we add in diagrams such as that in Fig. 1(c) the helicity of the antiquark is not changed and the leading behavior of the distributions (13) remain the same. The diagram 1(d) contributes a component $\sim(1-x)^9$ to $\bar q_\Delta^{(-)}(x)$.

In the SU(6) limit where we neglect spin-spin forces, a polarized nucleon differs from our Δ in that it has two quarks with spin aligned ($q^{(+)}$) and one opposite ($q^{(-)}$). Using our simple distributions (11) we would estimate the antiquark asymmetry,

$$\bar a(x)=[\bar q^{(+)}(x)-\bar q^{(-)}(x)]/[\bar q^{(+)}(x)+\bar q^{(-)}(x)]\equiv\tilde{\bar q}(x)/\bar q(x), \qquad (12)$$

displayed in Fig. 2(b) where $q^{(+)}$ in $N^{(+)}$ is equal to $q^{(-)}$ in $N^{(-)}$ and the tilde signals subtraction. This asymmetry is, in principle, measured in the Drell-Yan[8] production of massive lepton pairs. As in e^+e^- annihilation where the massive photon is transversely polarized, because of the γ_μ electromagnetic coupling the probability for $\bar q q$ annihilation is greatest when the q and $\bar q$ have the opposite helicity. Hence

$$A^{\rm DY}(x_a,x_b)=\frac{d\sigma(p^{(+)}p^{(+)}\to l\bar l x)-d\sigma(p^{(+)}p^{(-)}\to l\bar l x)}{d\sigma(p^{(+)}p^{(+)}\to l\bar l x)+d\sigma(p^{(+)}p^{(-)}\to l\bar l x)}\cong\frac{\sum e_i^2[\tilde q_i(x_a)\tilde{\bar q}_i(x_b)+\tilde{\bar q}_i(x_a)\tilde q_i(x_b)]}{\sum e_i^2[q_i(x_a)\bar q_i(x_b)+\bar q_i(x_a)q_i(x_b)]}, \qquad (13)$$

where $x_a x_b=Q^2/s$ and $x_a x_b=x_L=2Q_L/\sqrt{s}$. For $x_a=x_b$ (at 90°) or if we take $x_a\gg x_b$ we can write this in

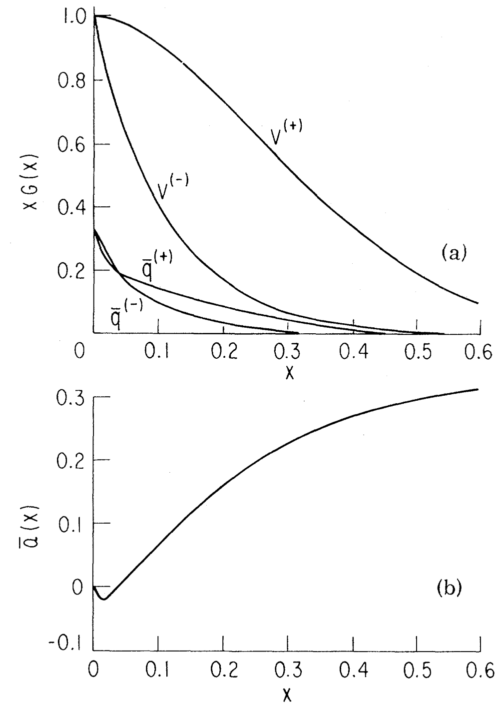

FIG. 2. (a) Gluon and antiquark distributions. (b) Antiquark asymmetry for an SU(6) proton.

the form

$$A^{\mathrm{DY}}(x_a, x_b) \to A^{\gamma P}(x_a)\bar{a}(x_b). \tag{14}$$

where $A^{\gamma P}(x_a) = \sum_i e_i^2 \bar{q}_i(x_a) / \sum_j e_j^2 q_j(x_a)$ is essentially the asymmetry measured in polarized electroproduction.[9,10] In the limit of exact SU(6) this would be given by

$$A^{\mathrm{DY}}(x_a, x_b) \simeq \tfrac{1}{3}(g_A/g_V)\bar{a}(x_b). \tag{15}$$

There are, however, reasons to believe that spin-spin forces imply that the large-x behavior of $A^{\gamma P}(x)$ is near 1.[10,11] Taking this into account, we may be underestimating the antiquark asymmetry in Fig. 2(b). For example, if $q^{(-)}(x)$ is suppressed as $x \to 1$ like $(1-x)q^{(+)}(x)$ as it is in some models, then the arguments in (6)-(11) would yield a similar suppression for $\bar{q}^{(-)}(x)$ so that $\bar{a}(x) \to 1$ as well. A more quantitative discussion of the x dependence would be model dependent so we merely wish to draw attention to the possibility of a large, observable asymmetry in lepton pair production. With emerging data on the electroproduction asymmetry, $A^{\gamma P}(x)$, it may be possible to make detailed predictions for $A^{\mathrm{DY}}(x)$.

We should also consider the possibility that the relationship between the sea and valence quarks is more complicated than that implied by perturbation theory. For example, gluons and virtual quark-antiquark pairs could be associated with excitations of a bag or some other entity not included in our calculation.[12] The mechanisms we have considered here do generate a component of the sea and the whirlpools we find there represent the way in which colored gluons transmit spin information to this component. We do not know how to quantify corrections to our mechanism at small x where they may be substantial but at large x the effect we have calculated should dominate. It seems possible to consider experimental measurement of the antiquark asymmetry a test of the adequacy of the usual perturbative approach to QCD.

We would like to comment on the $(1-x)^5$ component of the $\bar{q}$ distribution and its relation to the counting rules which assert that for a five-body *irreducible* bound system a $(1-x)^7$ behavior is required.[3] Farrar[5] takes this fact to imply a $(1-x)^7$ distribution for $\bar{q}(x)$ by invoking the usual Drell-Yan-West[13] connection with the t dependence of form factors:

$$\lim_{t\to\infty} F_{(\text{five body})}(t) \sim t^{-4} \longleftrightarrow \lim_{x\to 1} \nu W_2^{(\text{five body})}(x) \sim (1-x)^7. \tag{16}$$

While this is true when the five-body system is irreducible, the $qqq\bar{q}q$ can reduce to qqq if the $\bar{q}$ annihilates. Therefore, the leading form factor is not the "elastic" one but the one for the process $\gamma + (qqq\bar{q}q) \to (qqq)$ for which $F(t) \sim t^{-3}$ and $\nu W_2(x) \sim (1-x)^5$, in agreement with our result. In the approach of Blankenbecler and Brodsky[3] an $e\bar{q} \to e\bar{q}$ subprocess with exclusive limit $eB \to eB^*$ results in a $(1-x)^5$ behavior using the inclusive-exclusive connection while the exclusive limit $eB \to eB^*M^*$ implies $(1-x)^7$. The work of Gunion[4] is closest to our approach. He creates a $\bar{q}$ by first making a physical meson and pays an extra price of $(1-x)^2$ over what we find from an elementary gluon.

This work was performed under the auspices of the U. S. Energy Research and Development Administration.

[(a)]Permanent address.

[1]See, for example, G. W. Look and E. Fischbach, Phys. Rev. D 16, 211 (1977).

[2]V. A. Novikov, M. A. Shifman, A. I. Fainshtein, and V. I. Zakharov, Institute for Theoretical and Experimental Physics, Moscow, Report No. ITEP-112 (un-

published).
[3]R. Blankenbecler and S. J. Brodsky, Phys. Rev. D 10, 2973 (1974).
[4]J. Gunion, Phys. Rev. D 10, 242 (1974).
[5]G. Farrar, Nucl. Phys. B77, 429 (1974).
[6]See, for example, F. E. Close, F. Halzen, and D. M. Scott, Phys. Lett. 68B, 447 (1977).
[7]G. Altarelli and G. Parisi, to be published.
[8]S. D. Drell and T.-M. Yan, Phys. Rev. Lett. 25, 316 (1970).
[9]J. Kuti and V. Weisskopf, Phys. Rev. D 4, 3418 (1971); F. E. Close, Nucl. Phys. B80, 269 (1974).
[10]R. Carlitz and J. Kaur, Phys. Rev. Lett. 38, 673, 1102(E) (1977).
[11]F. E. Close, Phys. Lett. 43B, 422 (1973); R. Carlitz, Phys. Lett. 58B, 345 (1975); G. Farrar and D. R. Jackson, Phys. Rev. Lett. 35, 1416 (1975).
[12]However, for a calculation in the M.I.T. bag in which the sea is generated by mechanism similar to ours, see J. F. Donoghue and E. Golowich, Phys. Rev. D 15, 3421 (1977). (i) In the bag, the $q\bar{q}$ sea is always at least partially polarized, regardless of bag excitation (in agreement with the parton analysis). (ii) However, the bag calculation holds for a *rest-frame hadron*, so direct comparison with parton calculations is not possible.
[13]S. D. Drell and T.-M. Yan, Phys. Rev. Lett. 24, 181 (1970); G. B. West, Phys. Rev. Lett. 24, 1206 (1970).

Hadron Production in Nuclear Collisions—a New Parton-Model Approach

Stanley J. Brodsky
Stanford Linear Accelerator Center, Stanford University, Stanford, California 94305

and

John F. Gunion
Department of Physics, University of California, Davis, California 95616

and

J. H. Kühn
Max-Planck-Institut für Physik und Astrophysik, München 40, Germany
(Received 5 May 1977)

We consider a quark-parton model in which wee partons of the projectile interact with wee partons of essentially independent nucleons in the nuclear target. The ratio of multiplicities in the central region is $\langle n\rangle_{HA}/\langle n\rangle_{HN} = \frac{1}{2}\bar{\nu} + \bar{\nu}/(\bar{\nu}+1)$, where $\bar{\nu}$ is the mean number of inelastic collisions of the projectile H, in agreement with experiment. We also predict nucleus-nucleus multiplicities, the multiplicity distribution, and the absence of shadowing in large-q^2 and large-p_T reactions.

Although the quark-parton model has been very successful in predicting the short-distance behavior of hadronic interactions, the underlying mechanisms involved in the production of hadrons in ordinary high-energy collisions have never been specified. In the case of particle production on nuclear targets, this fundamental uncertainty of the parton approach becomes amplified, and this has led to an extraordinary range of divergent predictions for even the most basic experimental parameters.[1] In this Letter we present a new approach to this problem based on a straightforward application of parton-model concepts. The resulting picture for nuclear collisions is very simple and in good agreement with experiment. It is based upon (1) the assumption that each inelastically excited nucleon in the nuclear target produces hadrons independently of the others, and (2) a specific hadronic collision model based on wee-parton interactions[2] analogous to the Drell-Yan[3] pair-production process.

We begin with a simple parton-model description of hadron-hadron interactions. Each hadron has a Fock-space decomposition in terms of multiparton states. An interaction occurs via a collision of a parton in the beam (B) with a parton in the target (A). The cross section takes the typical Drell-Yan form[3,4]

$$\sigma_{BA} = \sum_{\substack{a\in A\\ b\in B}} \int_0^1 dx_a \int_0^1 dx_b\, G_{a/A}(x_a) G_{b/B}(x_b) \hat{\sigma}_{ab}(\hat{s}_{ab}), \tag{1}$$

where $x_b = (k_b^0 + k_b^z)/(p_B^0 + p_B^z)$ and $x_a = (k_a^0 - k_a^z)/(p_A^0 - p_A^z)$ are the light-cone fractions ($p_B^z > 0$; $p_A^z < 0$) of the beam and target momenta, respectively, and $\hat{s}_{ab} = x_a x_b s + m_a^2 m_b^2/x_a x_b s$ is the square of collision energy for the subprocess. (For