Pion Content of the Nucleon in Polarized Semi-Inclusive DIS *

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Abstract

An explicit pionic component of the nucleon may be identified by measuring polarized $\Delta^{++}$ fragments produced in deep-inelastic scattering (DIS) off polarized protons. The pion-exchange model predicts highly correlated polarizations of the $\Delta^{++}$ and target proton, in marked contrast with the competing diquark fragmentation process.

I. INTRODUCTION

Much has been learned about the internal structure of the nucleon from inclusive deep-inelastic lepton scattering since the first such experiments were carried out at SLAC in the late 1960s. These and many subsequent DIS experiments [1] were instrumental to the development of the quark-parton model of hadrons.

Successful as it has been in describing much of the DIS data in the perturbative region, a quark-parton based picture of the nucleon is not entirely complete — it is unable to explain the non-perturbative structure of the nucleon. A very good example of this is the deviation from the parton model prediction for the Gottfried sum rule seen in the recent NMC experiment [2]. The most natural explanation of this result is that there exists an excess of $\bar{d}$ quarks over $\bar{u}$ in the proton, something which is clearly impossible to obtain from perturbative QCD alone. A non-perturbative pionic cloud

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in the nucleon on the other hand offers a simple explanation of SU(2) flavor symmetry breaking in the proton sea [3].

The pion content of the nucleon has been investigated in previous studies of inclusive DIS [4], although only upper bounds on the average pion number were able to be extracted. In addition, since the pion contribution to the nucleon structure function appears at relatively small Bjorken $x$ ($x \sim 0.1$), its signal is lost amongst the perturbative sea background. It is therefore a challenge to seek direct experimental confirmation of pionic effects in inclusive nucleon DIS.

The pertinent question is whether pions leave any unique traces in other processes which cannot be understood in terms of quarks and gluons alone. A possible candidate for this is semi-inclusive DIS, where a hadron is detected in the final state in coincidence with the scattered electron. We shall demonstrate that pions give rise to characteristic fragmentation distributions in comparison with the predictions of parton model hadronization. These differences are significantly enhanced when polarization degrees of freedom are considered. We focus on semi-inclusive polarized $\Delta^{++}$ electroproduction from a polarized proton, $e\bar{p} \rightarrow e'\Delta^{++}X^-$. (The decay products of $\Delta^+$ or $\Delta^0$ include neutrals whose detection would be more difficult.)

Because the $g_1$ structure function of a pion is zero, an unpolarized electron beam will suffice for this purpose. The produced $\Delta$’s will appear predominantly in the target fragmentation region, or backward hemisphere in the $\gamma^*p$ c.m. frame, and hence will appear slow in the laboratory frame. With the high luminosity beam available at CEBAF the rate of $\Delta^{++}$ production will generally be high. Even though the efficiency with which low momentum $\Delta$’s can be accurately identified is lower than for fast baryons produced in the forward $\gamma^*p$ c.m. hemisphere, their detection may still be feasible with the CEBAF Large Acceptance Spectrometer.

II. KINEMATICS OF TARGET FRAGMENTATION

We define our variables in the laboratory frame as follows: $l, l'$ are the momentum four-vectors of the initial and final electrons, $P_\mu = (M; 0, 0, 0)$ and $p_\mu = (p_0; |\mathbf{p}| \sin \alpha \cos \phi, |\mathbf{p}| \sin \alpha \sin \phi, |\mathbf{p}| \cos \alpha)$ are the momentum vectors of the target proton and recoil $\Delta$, respectively, and $q_\mu = (\nu; 0, 0, \sqrt{\nu^2 + Q^2})$ denotes the photon 4-momentum, defined to lie along the positive $z$-axis. Then $\nu = E - E'$ is the
energy transferred to the target, \( y = \nu/E = 1 - E'/E \) is the fractional energy transfer relative to the incident energy, and \( Q^2 = -q^2 = 2MExy \) is minus the four-momentum squared of the virtual photon, with \( x = Q^2/2P\cdot q \). With the CEBAF upgrade to \( E \approx 8 - 10 \) GeV, values of \( x \approx 0.13 - 0.14 \) can be reached in the deep inelastic region for \( \nu \approx 8 \) GeV and \( Q^2 \approx 2 \) GeV².

The four-momentum transfer squared between the proton and \( \Delta \) is \( t = (P-p)^2 = -p_T^2/\zeta + t_{\text{max}} \), which is bounded from above by \( t_{\text{max}} = -(M_{\Delta}^2 - M^2\zeta)(1 - \zeta)/\zeta \), where \( \zeta = p\cdot q/P\cdot q \) is the light-cone fraction of the target proton’s momentum carried by the \( \Delta \). In terms of \( t \), the three-momentum of the produced \( \Delta \) is given by

\[
|p| = \frac{1}{2M} \sqrt{(M^2 + M_{\Delta}^2 - t)^2 - 4M^2M_{\Delta}^2},
\]  

so that in the lab. frame the slowest \( \Delta \)’s are those for which \( t \to 0 \), which occurs when \( \zeta \to 1 \). As the upper limit on \( \zeta \) is \( 1 - x \), slow \( \Delta \) production also corresponds to the \( x \to 0 \) limit, and the slowest possible particles produced at \( \zeta = 1 \) (at \( x = 0 \)) have momentum \( |p_{\text{min}}| = (M_{\Delta}^2 - M^2)/2M \approx 340 \) MeV. For the pion-exchange process considered here, the peak in the differential cross section occurs at \( |p| \sim 600 \) MeV. This corresponds to a total c.m. energy squared of the photon–proton system of \( W^2 = (P + q)^2 \sim 14 \) GeV², and a missing mass of \( p_X^2 = (k + q)^2 \sim 0.8 \) GeV².

In terms of the laboratory angle \( \alpha \) between the \( \Delta \) and \( \gamma^* \) directions, where

\[
\cos \alpha = \frac{M_{\Delta}^2 + (1 - 2\zeta)M^2 - t}{\sqrt{(M_{\Delta}^2 - M^2 - t)^2 - 4M^2t}},
\]

production of \( \Delta \)'s will occur between \( \alpha = 0 \) and \( \alpha_{\text{max}} = \cos^{-1} \left( \sqrt{1 - (M\zeta/M_{\Delta})^2} \right) \approx 50^\circ \) for \( \zeta \to 1 \).

For studies of the spin dependence of the fragmentation process, we require the target proton polarization to be parallel to the photon direction, with the spin of the produced \( \Delta \) quantized along its direction of motion. Experimentally, the polarization of the produced \( \Delta^{++} \) can be reconstructed from the angular distribution of its decay products (\( p \) and \( \pi^+ \)).
III. DYNAMICS OF PION EXCHANGE

The relevant process is illustrated in Fig. 1, where the dissociation of a physical proton into a $\pi^-$ and $\Delta^{++}$ is explicitly witnessed by the probing photon. In the pion-exchange model the differential cross section is

$$\frac{d^3\sigma}{dxdQ^2d\zeta dp_T^2d\phi} = \left(\frac{\alpha^2}{ME^2x^2Q^4\zeta}\right) \frac{f_{\pi N\Delta}^2}{16\pi^2m_{\pi}^2} \frac{T^S_s(t)T_{\pi\Delta}^2}{(t-m_{\pi}^2)^2} L_{\mu\nu}(l,q) W_{\mu\nu}(k,q)$$

(3)

where $L_{\mu\nu} = 2l'_{\mu}l'_{\nu} + 2l'_{\nu}l'_{\mu} - g_{\mu\nu}Q^2$ is the lepton tensor, and

$$W_{\pi}^{\mu\nu} = -\left(g^{\mu\nu} + \frac{q^\mu q^\nu}{Q^2}\right) W_{1\pi} + \left(k^\mu + \frac{k\cdot q}{Q^2} q^\mu\right) \left(k^\nu + \frac{k\cdot q}{Q^2} q^\nu\right) \frac{W_{2\pi}^{2\pi}}{m_{\pi}^2}$$

(4)

describes the $\gamma^*\pi$ vertex. The pion structure function is extracted from Drell-Yan experiments [5].

For the $\pi N\Delta$ form factor in Eq.(3) we take the form suggested in Ref.[6] in an analysis of the pionic content of the proton in the infinite momentum frame,

$$F_{\pi\Delta}(p_T^2, \zeta) = \left(\frac{\Lambda^2 + M^2}{\Lambda^2 + s_{\pi\Delta}}\right)^2$$

(5)

where $s_{\pi\Delta} \equiv (m_{\pi}^2 + p_T^2)/(1-\zeta) + (M_{\Delta}^2 + p_T^2)/\zeta$. As discussed in Ref.[6], the formulation in the infinite momentum frame allows the use of the experimental structure function of the pion in Eq.(1).

The function $T^S_s(t)$ is obtained by evaluating the trace over the target nucleon spinor and the Rarita-Schwinger spinor-vector for the recoil $\Delta$. Because it is emitted collinearly with the pion, production of $\Delta$ baryons with spin projection $\pm 3/2$ is forbidden, which leads to the selection rule:
The yield of spin projection $\pm 1/2$ states is given by

$$T^{\pm \frac{1}{2}}(t) = \frac{1}{12M_{\Delta}^2} \left[ (M - M_{\Delta})^2 - t \right] \left[ (M + M_{\Delta})^2 - t \right] (1 \pm \cos \alpha)$$

where $\alpha$ is also the angle between the polarization vectors $S$ and $s$. Because the production of $\Delta$ baryons is limited to forward laboratory angles, the presence of the $(1 \pm \cos \alpha)$ factor associated with the final state polarization will significantly suppress the $s = -1/2$ yield relative to that of $s = +1/2$ final states. This suppression leads to strikingly different predictions for the polarization asymmetry compared with those of the competing parton fragmentation process discussed below.

IV. BACKGROUNDS

At the large $Q^2$ and $\nu$ possible with an 8–10 GeV electron beam, the resonance backgrounds should not pose a major problem in identifying the required signal. Firstly, interference from quasi-elastic $\Delta^{++}$ production will not be present because of charge conservation. Secondly, the large $W(\sim 3 - 4$ GeV) involved means that interference from excited $\Delta^*$ states (with subsequent decay to $\Delta^{++}$ and pions) will be negligible. In addition, any such resonance contributions will be strongly suppressed by electromagnetic form factors at large $Q^2$ ($Q^2 \gtrsim 2$ GeV$^2$).

A potentially more important background will be that due to uncorrelated spectator fragmentation. Within the parton model framework the cross section for this process (assuming factorization of the $x$ and $\zeta$ dependence) is proportional to $[7] F_{p^\uparrow}(x, Q^2) \widetilde{D}_{p^\uparrow-q^\uparrow\downarrow}(z, p_T^2)$, where $z = \zeta/(1 - x)$ is the light-cone momentum fraction of the produced baryon carried by the spectator system. The fragmentation function $\widetilde{D}(z, p_T^2)$ gives the probability for the polarized ($p^\uparrow$ minus $q^\downarrow\downarrow$) spectator system to fragment into a $\Delta^{++}$ with polarization $s$. The usual assumption is that the transverse momentum distribution of the $\Delta$ also factorizes, $\widetilde{D}_{p^\uparrow-q^\uparrow\downarrow}(z, p_T^2) = D_{p^\uparrow-q^\uparrow\downarrow}(z) \varphi(p_T^2)$, with $\int dp_T^2 \varphi(p_T^2) = 1$.

The function $F_{p^\uparrow}(x, Q^2)$ is proportional to the spin-weighted interacting-quark momentum distribution functions, $q^\uparrow\downarrow(x)$, where $\uparrow\downarrow$ denote quark spins parallel or antiparallel to the spin of the proton. For estimation purposes we consider the model for $q^\uparrow\downarrow(x)$ of Carlitz & Kaur [8]. The results change little if one uses the model of Schäfer [9].
At large \( z \) the fragment \( \Delta^{++} \) carries most of the parent system’s momentum, and therefore contains both valence \( u \) quarks from the target proton. In our region of interest (\( z \gtrsim 0.6 \)) by far the most important contributions to \( D(z) \) come from the process whereby the \( \Delta^{++} \) is formed after only one \( u\bar{u} \) pair is created [10]. As a consequence DIS from valence \( u \) quarks will be unimportant.

For scattering from sea quarks we assume the same fragmentation probabilities for \( uuq\bar{q} \) spectator states as for \( uu \) (although at \( Q^2 \approx 2 \) GeV\(^2 \) the sea constitutes at most \( \sim 15\% \) of the cross section at \( x \sim 0.1 \)). We thus parametrize the (very limited) EMC data [11] on unpolarized \( \Delta^{++} \) muon production for \( z \rightarrow 1 \) as:

\[
D_{uu}(z \rightarrow 1) = a(1 - z)^b,
\]

where \( a \approx 0.68 \) and \( b \approx 0.3 \).

For the spin dependence of the fragmentation process we follow the approach taken by Bartl et al. [12] in their study of polarized quark \( \rightarrow \) baryon fragmentation. Namely, the diquark is assumed to retain its helicity during its decay, and the \( q\bar{q} \) pair creation probability is independent of the helicity state of the quark \( q \). At leading order this means that the produced baryon contains the helicity of the diquark, so that, for example, a \( \Delta^{\uparrow} \) or \( \Delta^{\downarrow} \) can emerge from a \( q^\uparrow q^\uparrow \) diquark, whereas a \( \Delta^{\downarrow} \) cannot. (Notation here is that \( \uparrow, \uparrow, \downarrow, \downarrow \) represent \( s = +3/2, +1/2, -1/2, -3/2 \) states, respectively.)

The overall normalization of the spin-dependent fragmentation functions is fixed by the condition

\[
q(x) \, D_{p-q}(z) + \bar{q}(x) \, D_{p-q}(z) = q^\uparrow(x) \, D_{p^\uparrow-q^\uparrow}(z) + q^\downarrow(x) \, D_{p^\uparrow-q^\downarrow}(z) \\
+ \bar{q}^\uparrow(x) \, D_{p^\uparrow-q^\uparrow}(z) + \bar{q}^\downarrow(x) \, D_{p^\uparrow-q^\downarrow}(z)
\] (8)

where \( D(z) = \sum_{s=-3/2}^{+3/2} D^s(z) \). In relating the production rates for various polarized \( \Delta^{++} \) we employ SU(6) spin-flavor wavefunctions, from which simple relations among the valence diquark \( \rightarrow \Delta^{++} \) fragmentation functions, \( D^{s}_{qq_{j_{z}}}(z) \), can be deduced (the diquark state \( qq_{j_{z}} \) is labeled by its spin \( j \) and spin projection \( j_{z} \)). The leading functions are related by:

\[
D^{\uparrow}_{uu_{(1)}}(z) = 3 \, D^{\uparrow}_{uu_{(1)}}(z) = 3 \, D^{\uparrow}_{uu_{(0)}}(z) = 3 \, D^{\downarrow}_{uu_{(0)}}(z) = \frac{3}{2} D^{\uparrow}_{uu_{(0)}}(z),
\] (9)

with normalization determined from:

\[
D^{\uparrow}_{uu_{(1)}}(z) = \frac{3}{4} D_{uu}(z).
\] (10)

(Note that this is true only when the spin projections of the diquark and \( \Delta \) are in the same direction.)

The non-leading fragmentation functions are those which require at least two \( q\bar{q} \) pairs to be created.
from the vacuum, namely \( D_{uu(0)}^{\uparrow/\downarrow} \), \( D_{ud(0)}^{\uparrow/\downarrow} \), \( D_{uu(0)}^{\uparrow/\downarrow} \), and \( D_{ud(1)}^{\uparrow/\downarrow} \), and those which require 3 such pairs, \( D_{uu(1)}^{\uparrow/\downarrow} \) and \( D_{ud(1)}^{\uparrow/\downarrow} \). Except at very small \( z \) (\( z \lesssim 0.2 \)) the latter functions are consistent with zero [10]. For the 2-\( q\bar{q} \) pair fragmentation functions, we also expect that \( D_{uu(0)}^{\uparrow/\downarrow}(z) = D_{uu(1)}^{\downarrow}(z) \).

For \( z > \sim 0.2 \) the unpolarized model fragmentation functions of Ref.[10] requiring two \( q\bar{q} \) pairs (e.g. \( D_{ud}(z) \)) are quite small compared with the leading fragmentation functions, \( D_{ud}(z) \simeq 0.1 \ D_{uu}(z) \).

For spin-dependent fragmentation we therefore expect a similar behavior for those decay probabilities requiring two \( q\bar{q} \) pairs created in order to form the final state with the correct spin and flavor quantum numbers. This then allows for a complete description of polarized fragmentation at large \( z \) in terms of only the 4 fragmentation functions in Eq.(9).

Finally, the \( p_T \)-integrated differential cross section for the electroproduction of a \( \Delta^{++} \) with spin \( s \) can be written:

\[
\frac{d^3\sigma}{dx dQ^2 d\zeta} = \left( \frac{2\pi\alpha^2}{M^2 E^2 x(1-x)} \right) \left( \frac{1}{2x^2} + \frac{4M^2 E^2}{Q^4} \left( 1 - \frac{Q^2}{2MEx} - \frac{Q^2}{4E^2} \right) \right) \times \left[ \frac{4x}{9} \left( u^+_V D_{ud(1)}^s + 2\tilde{u}^+_V \left( \frac{2}{3} D_{uu(0)}^s + \frac{1}{3} D_{uu(1)}^s \right) + u^+_V D_{ud(1)}^s + 2\tilde{u}^+_V \left( \frac{2}{3} D_{uu(1)}^s + \frac{1}{3} D_{uu(1)}^s \right) \right) \right. \\
+ \left. \frac{x}{9} \left( d^+_V D_{uu(0)}^s + 2\tilde{d}^+_V \left( \frac{2}{3} D_{uu(0)}^s + \frac{1}{3} D_{uu(1)}^s \right) + d^+_V D_{uu(1)}^s + 2\tilde{d}^+_V \left( \frac{2}{3} D_{uu(1)}^s + \frac{1}{3} D_{uu(1)}^s \right) \right) \right].
\]

V. NUMERICAL RESULTS AND DISCUSSION

The differential cross section, \( Q^2 d^3\sigma/dxdQ^2 d\zeta \), for the individual polarization states of the produced \( \Delta^{++} \) (for DIS from a proton with \( S = +1/2 \)) is shown in Fig.2, for \( x = 0.14, Q^2 = 2 \) GeV\(^2 \) and \( E = 8 \) GeV. The pion-exchange model predictions (solid curves) are for a form factor cut-off of \( \Lambda = 800 \) MeV, which gives \( <n>_{p\Delta} \approx 0.02 \) [6]. (For the same \( <n>_{p\Delta} \) the cut-off in a \( t \)-dependent dipole form factor would be \( \sim 700 \) MeV.) The spectrum shows strong correlations between the polarizations of the target proton and \( \Delta^{++} \). In the quark-parton model (dashed curves) the correlations are significantly weaker, and the ratio of polarized \( \Delta \)'s in this case is \( s = +3/2 : +1/2 : -1/2 : -3/2 \approx 3 : 2 : 1 : 0 \).
FIG. 2. Differential electroproduction cross section for various polarization states of the $\Delta^{++}$. The $\pi$-exchange model, with cut-off mass $\Lambda = 800$ MeV, is shown by solid curves. The quark-parton model background (dashed curves) is estimated using the fragmentation functions extracted from the unpolarized EMC data [11] and Eqs. (9) and (10).

The differences between the pion-exchange model and fragmentation backgrounds can be further enhanced by examining polarization asymmetries. In Fig.3 we show the difference $\sigma^+ - \sigma^-$, where $\sigma^\pm \equiv Q^2 d^3/xdx dQ^2 d\zeta (s = \pm 1/2)$, as a fraction of the total unpolarized cross section. The resulting $\zeta$ distributions are almost flat, but significantly different for the two models.

FIG. 3. Polarization asymmetry for the $\pi$-exchange (solid) and parton fragmentation (dashed) models, with $\sigma^\pm$ as defined in the text, and $\sigma_{\text{tot}}$ is the sum over all polarization states.

Of course the two curves in Fig.3 represent extreme cases, in which $\Delta$'s are produced entirely via pion emission or diquark fragmentation. In reality we can expect a ratio of polarization cross
sections which is some average of the curves in Fig.3. The amount of deviation from the parton model curve will indicate the extent to which the pion-exchange process contributes. From this, one can in turn deduce the strength of the $\pi N\Delta$ form factor. Unlike inclusive DIS, which can only be used to place upper bounds on the pion number, the semi-inclusive measurements could pin down the absolute value of $< n >_{\pi\Delta}$. A measurement of this ratio would thus be particularly useful in testing the relevance of non-perturbative degrees of freedom in high energy DIS processes.

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