Toward Understanding the Quark Parton Model of Feynman

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Abstract

The quark parton model of Feynman, which has been used for analyses of high energy physics experiments, invokes a set of parton distributions in the description of the nucleon structure (the probability concept), contrary to the traditional use of wave functions in nuclear and medium energy physics for structural studies (the amplitude concept). In this paper, I first review briefly how the various parton distributions of a nucleon may be extracted from high energy physics experiments. I then proceed to consider how the sea distributions of a free nucleon at low and moderate Q^2 (e.g., up to $20\,GeV^2$), may be obtained in the meson-baryon picture, a proposal made by Hwang, Speth, and Brown. Using the form factors associated with the couplings of mesons to baryons such as πNN, $\pi N\Delta$, and $KN\Lambda$ couplings which are constrained by the CCFR neutrino data, we find that the model yields predictions consistent with the CDHS and Fermilab E615 data on the sea-to-valence ratio. We also find that the recent finding by the New Muon Collaboration (NMC) on the violation of the Gottfried sum rule can be understood quantitatively. Finally, we consider, using the pion as the example, how valence quark distributions of a hadron may be linked to the hadron wave function written in the light-cone language. Specifically, we use the leading pion wave function that is constrained by the QCD sum rules, and find that, at $Q^2 \approx (0.5\,GeV)^2$, the leading Fock component accounts for about 40 % of the observed valence quark distributions in the pion. The question of how to generate the entire valence quark distributions from the valence quark distribution calculated from the leading Fock component is briefly considered again using the specific proposal of Hwang, Speth, and Brown.

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I. Introduction

The physics Nobel prize in 1990 was awarded to Jerome Friedman, Henry Kendall, and Richard Taylor for the celebrated SLAC-MIT experiments[1] which were carried out in late 1960's. The experiments demonstrate that, at large Q^2 with Q^2 the four-momentum transfer squared or, equivalently, at very high resolution (typically, in the sub-fermi range), a nucleon (proton or neutron) looks like a collection of (almost non-interacting) pointlike partons. Nowadays, we identify these partons as quarks, antiquarks, and gluons. Consequently, a nucleon is described by a set of structure functions or distributions:

$$\begin{aligned}&\{u(x,Q^2),\ \bar{u}(x,Q^2),\ d(x,Q^2),\ \bar{d}(x,Q^2),\ s(x,Q^2),\ \bar{s}(x,Q^2),\ \cdots\\&g(x,Q^2)\}.\end{aligned} \tag{1}$$

Here x is the fraction of the hadron longitudinal momentum carried by the parton as visualized[2] in the "infinite momentum frame" ($P_z \to \infty$).

High energy physics experiments with hadrons, including deep inelastic scattering (DIS) by charge leptons (e or μ), Drell-Yan (DY) production in hadronic collisions ($A+B \to \ell^+ + \ell^- + X$), experiments with high energy neutrino beams, charm production with high energy neutrinos, and others, all have customarily been analyzed in the framework of the quark parton model, as proposed by R. P. Feynman.[2]

On a different front, the idea of quarks or antiquarks as the buliding blocks of hadrons was proposed[3] in 1964 by Gell-Mann and, independently, Zweig. Since then, quark models in a variety of detailed forms have been introduced and shown to describe quite successfully the structure of hadrons. For instance, a proton (or neutron) at low Q^2 may be treated approximately as a system of three quarks (u,u,d) (or (d,d,u)) confined to within a region defined by the hadron size (perhaps with each of the quarks dressed by quark-antiquark pairs or gluons). Unlike what is in the parton model where one adopts distributions to describe a nucleon (at large Q^2), one uses wave functions in the quark model to characterize the structure of a nucleon (at low Q^2). The dual picture for describing the nucleon structure has generated the impression that the information acquired from high energy physics experiments seems, to some extent, irrelevant for low energy strong interaction physics which are described very well by the meson-baryon picture via wave functions (rather than via distributions).

Nowadays, it is believed that quantum chromodynamics (QCD)[4] describes strong interactions among quarks and gluons. The "existence" of quarks or antiquarks is beyond doubt even though a quark or antiquark *in isolation* is not found experimentally. The experimental information points to the confinement of color - namely, quarks, antiquarks, and gluons all carry colors and only color-singlet, or colorless, objects may be found in the true ground state (vacuum) of QCD. Accordingly, quarks, antiquarks, and gluons must organize themselves into colorless clusters (hadrons) in the true vacuum, leading to the observation that low energy strong interaction physics (or hadron physics) can *successfully* be described *effectively* in terms of mesons and baryons (the "meson-baryon picture").

Although the gap between high energy (particle) physics and nuclear physics is quite understandable owing to the dual picture for describing the nucleon structure at large Q^2 and small Q^2, we nevertheless believe that such "gap" is caused primarily by our ignorance toward the physics associated with the quark parton model of Feynman.[2] Should it become possible to obtain quantitatively the various parton distributions from a quark model (perhaps augmented with the meson-baryon picture), the information extracted from particle physics experiments is then complementary, rather than orthogonal (as of today), to that deduced from low energy strong interaction physics. It is the purpose of the present paper to provide an overview of recent efforts, mostly of our own as much more space and time would be needed for a truly comprehensive review, in trying to understand the quark parton model of Feynman.[2]

II. Parton Distributions as Inferred from Experiments

To unravel the structure functions of a nucleon, it is customary to employ deep inelastic scatterings (DIS) of the various kind, including (a) DIS by charge leptons, (b) experiments with high energy neutrino beams, (c) heavy flavor production by high energy neutrinos, and others. In addition, Drell-Yan (DY) lepton pair production in hadronic collisions has also yielded indispensible information concerning parton distributions associated with a hadron. In what follows, we illustrate briefly these reactions in connection with parton distributions of a proton.

(a) Deep Inelastic Scatterings by Charge Leptons:

Consider the DIS by electrons or muons:

$$\ell(\ell) + p(P) \rightarrow \ell'(\ell') + X. \tag{2}$$

We may adopt the following kinematic variables:

$$q \equiv \ell - \ell',$$

$$Q^2 \equiv q^2 \approx 4E_\ell E_{\ell'} \sin^2\frac{\theta_L}{2}, \tag{3a}$$

$$\nu \equiv -\frac{q \cdot P}{m_N} = E_\ell - E_{\ell'}, \tag{3b}$$

$$x \equiv \frac{Q^2}{2m_N\nu}, \tag{3c}$$

$$y \equiv \frac{q \cdot P}{\ell \cdot P} = \frac{\nu}{E_\ell}, \tag{3d}$$

$$s \equiv -(\ell + P)^2 = \frac{Q^2}{xy} + m_N^2. \tag{3e}$$

Here the metric is chosen to be pseudo-Euclidean such that $q^2 \equiv \vec{q}^{\,2} - q_0^2$, but we shall use wherever possible the variables Q^2, ν, x, y and others which are metric-independent and are adopted fairly universally in the literature. Note that, in Eqs. (3a), (3b), and (3d), the last equality holds only in the laboratory frame. The differential cross section for the DIS process, Eq. (2), may be cast in any of the following forms:

$$\begin{aligned}\frac{d^2\sigma}{dx\,dy} &= \nu(s-m_N^2)\frac{d^2\sigma}{dQ^2\,d\nu}\\ &= x(s-m_N^2)\frac{d^2\sigma}{dQ^2\,dx}\\ &= \frac{2\pi m_N \nu}{E_{\ell'}}\frac{d^2\sigma}{d\Omega_{\ell'}\,dE_{\ell'}}.\end{aligned} \tag{4}$$

In particular, we have

$$\frac{d^2\sigma}{dx\,dy} = \frac{2\pi\alpha^2}{sx^2y^2}\{1+(1-y)^2\}F_2(x,Q^2), \tag{5a}$$

with

$$F_2(x,Q^2) = \sum_i e_i^2 x f_i(x,Q^2). \tag{5b}$$

Here the summation over the flavor index i extends over all quarks and antiquarks. (See Eq. (1).) Thus, DIS with a proton target allows for the measurement of the structure function $F_2^{ep}(x,Q^2)$:

$$\begin{aligned}F_2^{ep}(x,Q^2) = {}&\frac{4}{9}x(u^p(x,Q^2)+\bar{u}^p(x,Q^2)) + \frac{1}{9}x(d^p(x,Q^2)+\bar{d}^p(x,Q^2))\\ &+\frac{1}{9}x(s^p(x,Q^2)+\bar{s}^p(x,Q^2))+\cdots.\end{aligned} \tag{6}$$

On the other hand, DIS with a neutron target (often with a deuteron target with the contributions from the proton suitably subtracted) yields

$$\begin{aligned}F_2^{en}(x,Q^2) = {}&\frac{1}{9}x(d^n(x,Q^2)+\bar{d}^n(x,Q^2)) + \frac{4}{9}x(u^n(x,Q^2)+\bar{u}^n(x,Q^2))\\ &+\frac{1}{9}x(s^n(x,Q^2)+\bar{s}^n(x,Q^2))+\cdots.\end{aligned} \tag{7}$$

The proton and neutron are two members of an isospin doublet. Assuming isospin symmetry, we have $d^n(x) = u^p(x)$, $\bar{d}^n(x) = \bar{u}^p(x)$, $u^n(x) = d^p(x)$, ...

$s^n(x) = s^p(x)$, etc. Eq. (7) becomes

$$F_2^{en}(x,Q^2) = \frac{1}{9}x(u^p(x,Q^2)+\bar{u}^p(x,Q^2)) + \frac{4}{9}x(d^p(x,Q^2)+\bar{d}^p(x,Q^2)) + \frac{1}{9}x(s^p(x,Q^2)+\bar{s}^p(x,Q^2)) + \cdots. \tag{8}$$

Taking the sum of Eqs. (6) and (8), we find

$$\begin{aligned}&\int_0^1 \{F_2^{ep}(x,Q^2)+F_2^{en}(x,Q^2)\}dx \\ =&\frac{5}{9}\{<x>_u + <x>_{\bar{u}} + <x>_d + <x>_{\bar{d}} + <x>_s + <x>_{\bar{s}} + \cdots\} \\ &-\frac{1}{3}\{<x>_s + <x>_{\bar{s}} + \cdots\},\end{aligned} \tag{9}$$

where $<x>_i \equiv \int_0^1 f_i(x,Q^2)dx$ is the total momentum fraction carried by the quarks (antiquarks) of flavor i. Using the deuteron as the target and neglecting the very small correction from heavy quarks, we conclude from the SLAC-MIT experiments[1] that in the proton the total momenta carried by all quarks and antiquarks add up to only about (40 − 50)% of the proton momentum, leaving the rest of the momentum for electrically neutral partons (such as gluons).

On the other hand, we may take the difference between Eqs. (6) and (8) and obtain

$$\begin{aligned}S_G &\equiv \int_0^1 \frac{dx}{x}\{F_2^{ep}(x) - F_2^{en}(x)\} \\ &= \frac{1}{3}\int_0^1 \{u^p(x) - d^p(x) + \bar{u}^p(x) - \bar{d}^p(x)\} \\ &= \frac{1}{3} + \frac{2}{3}\int_0^1 \{\bar{u}^p(x) - \bar{d}^p(x)\}.\end{aligned} \tag{10}$$

The Gottfried sum rule[5] (GSR) may then be derived by assuming *isospin independence* of the sea distributions in the *proton*:

$$\bar{u}^p(x) = \bar{d}^p(x), \tag{11}$$

so that $S_G = \frac{1}{3}$. Here we emphasize that *isospin invariance*, or *isospin symmetry*, is *not* violated even if Eq. (11) is *not* true, since as a member of an isospin doublet the proton already has different valence u and d quark distributions. Thus, perhaps the preliminary value[6] reported by the New Muon Collaboration (NMC), at $< Q^2 >= 4\,GeV^2$, should not be considered as a major surprise:

$$\int_{0.004}^{1} \frac{dx}{x}\{F_2^{ep}(x) - F_2^{en}(x)\} = 0.230 \pm 0.013(stat.) \pm 0.027(syst.). \qquad (12a)$$

Analogously, using only the NMC F_2^n/F_2^p ratio and the world average fit to F_2 on deuterium the following value has been obtained:[7]

$$\int_{0}^{1} \frac{dx}{x}\{F_2^{ep}(x) - F_2^{en}(x)\} = 0.240 \pm 0.016. \qquad (12b)$$

The significance of the finding by the NMC group on the violation of the Gottfried sum rule has been discussed by many authors,[8−13] especially concerning the possible origin of the observed violation. As emphasized by Preparata, Ratcliffe, and Soffer[9], the observed deviation of S_G from $\frac{1}{3}$ is at variance with the standard hypothesis used by many of us for years:

$$\bar{u}_p(x) = \bar{d}_p(x) \approx 2\bar{s}_p(x). \qquad (13)$$

Such difference may be considered as a *surprise* but, as mentioned earlier, the deviation of S_G from the value of $\frac{1}{3}$ is *not* a signature for violation of isospin invariance or isospin symmetry. It is therefore helpful to caution that the words used by some authors, such as "isospin violation" by Preparata *et al.*[9] or "isospin symmetry violation" by Anselmino and Predazzi[10], are in fact somewhat misleading.

As the standard hypothesis Eq. (11) is nullified by the NMC data, Eqs. (12a) and (12b), it implies that almost all of the existing parametrized parton distributions, as well as the analyses of the high energy experiments in extracting the sea quark distributions for the nucleon or other hadrons, suffer from the commonly accepted *bias*. In this regard, we may echo the criticisms raised by Preparata *et al.*[9]

(b) Experiments with High Energy Neutrino Beams:

DIS with high energy neutrino beams on a proton target may either proceed with charged-current weak interactions:

$$\nu_\mu(\ell) + p(P) \rightarrow \mu^-(\ell') + X, \tag{14a}$$

$$\bar{\nu}_\mu(\ell) + p(P) \rightarrow \mu^+(\ell') + X, \tag{14b}$$

or proceed with neutral-current weak interactions:

$$\nu_\mu(\ell) + p(P) \rightarrow \nu_\mu(\ell') + X, \tag{15a}$$

$$\bar{\nu}_\mu(\ell) + p(P) \rightarrow \bar{\nu}_\mu(\ell') + X. \tag{15b}$$

The subject has been reviewed by many authors; we recommend the published lecture delivered by J. Steinberger[14] at the occasion of the presentation of the 1988 Nobel Prize in Physics.

The cross sections for Eqs. (14a) and (14b) are given by

$$\frac{d^2\sigma^\nu}{dxdy} = \frac{G^2 E m_p}{\pi} x\{q(x) + (1-y)^2\bar{q}(x)\}, \tag{16a}$$

$$\frac{d^2\sigma^{\bar{\nu}}}{dxdy} = \frac{G^2 E m_p}{\pi} x\{\bar{q}(x) + (1-y)^2 q(x)\}, \tag{16b}$$

with $q(x) = u(x) + d(x) + s(x) + \cdots$ and $\bar{q}(x) = \bar{u}(x) + \bar{d}(x) + \bar{s}(x) + \cdots$. It is clear that the measurements with both the neutrino and antineutrino beams offer a means to determine the quantity $R_{\bar{Q}}$:

$$R_{\bar{Q}} \equiv \frac{<x>_{\bar{u}} + <x>_{\bar{d}} + <x>_{\bar{s}}}{<x>_u + <x>_s + <x>_s}, \tag{17}$$

since contributions from heavy c quarks and others are negligibly small. The quantity $R_{\bar{Q}}$ is the ratio of the total momentum fraction carried by antiquarks to that carried by quarks, a ratio that sets the constraint for the amount of the antiquark sea in the nucleon. Experimentally, the CCFR Collaboration obtained[15]

$$R_{\bar{Q}} = 0.153 \pm 0.034, \tag{18}$$

a result similar to what obtained earlier by the CDHS and HPWF Collaborations.

The reactions (15a) and (15b) have been used to determine[16] the couplings of Z^0 boson to the *up* and *down* quarks, thereby providing tests of the standard electroweak theory.[17]

(c) Charm Production by High Energy Neutrinos:

Charm production by high energy neutrinos proceeds with the reaction:

$$\begin{aligned} &\nu_\mu + s(d) \rightarrow c + \mu^- \\ &\qquad\qquad c \rightarrow s + \mu^+ + \nu_\mu, \end{aligned} \tag{19a}$$

or, at the hadronic level,

$$\nu_\mu + p \rightarrow \mu^+\mu^- + K + X. \tag{19b}$$

Analogously, charm production by antineutrinos proceeds with the reaction:

$$\begin{aligned} &\bar{\nu}_\mu + \bar{s}(\bar{d}) \rightarrow \bar{c} + \mu^+ \\ &\qquad\qquad \bar{c} \rightarrow \bar{s} + \mu^- + \bar{\nu}_\mu, \end{aligned} \tag{19c}$$

or, at the hadronic level,

$$\bar{\nu}_\mu + p \rightarrow \mu^+\mu^- + \bar{K} + X. \tag{19d}$$

Thus, production of $\mu^+\mu^-$ pairs together with strange hadrons serves as a signature for charm production in inclusive $\nu_\mu(\bar{\nu}_\mu) + p$ reactions. Neglecting the charm quark mass for the sake of simplicity (which may be introduced in a straightforward manner), we may write the cross section as follows:

$$\begin{aligned}
&\frac{d^2\sigma}{dxdy}(\nu_\mu(\bar{\nu}_\mu) + p \rightarrow \mu^+\mu^- X) \\
&= \frac{G^2 E_\nu m_p}{\pi}\{sin^2\theta_c x[d(x) + \bar{d}(x)] + cos^2\theta_c x[u(x) + \bar{u}(x)]\},
\end{aligned} \tag{20}$$

with θ_c the Cabibbo angle ($sin\theta_c = (0.221 \pm 0.003)$).[16] It follows that such reactions provide an effective means to pin down the strange content of the proton. Indeed, CCFR Collaboration obtains[18]

$$\kappa \equiv \frac{2 < x >_s}{< x >_{\bar{u}} + < x >_{\bar{d}}} = 0.44^{+0.09+0.07}_{-0.07-0.02}; \tag{21a}$$

$$\eta_s \equiv \frac{2 < x >_s}{< x >_u + < x >_d} = 0.057^{+0.010+0.007}_{-0.008-0.002}, \tag{21b}$$

where the first errors are statistical while the second ones are systematic.

(d) Drell-Yan Production in Hadronic Collisions:

The Drell-Yan (DY) lepton-pair production in the hadronic collision:[19]

$$A(P_1) + B(P_2) \rightarrow \ell^+\ell^- + X, \tag{22}$$

proceeds, at the quark level, with the process

$$q + \bar{q} \rightarrow \gamma^* \rightarrow \ell^+ + \ell^-. \tag{23}$$

This is the leading-order process at the quark level, which yields

$$\begin{aligned}
&\frac{d^2\sigma}{dM^2 dx_F}(A + B \rightarrow \ell^+\ell^- + X) \\
&= \frac{4\pi\alpha^2}{9SM^2}\sum_q Q_q^2\{f_{q/A}(x_1)f_{\bar{q}/B}(x_2) + f_{\bar{q}/A}(x_1)f_{q/B}(x_2)\},
\end{aligned} \tag{24}$$

with $S \equiv -(P_1+P_2)^2$ (total center-of-mass energy squared), $M^2 \equiv -(\ell^+ + \ell^-)^2$ (invariant mass squared of the lepton pair), and $x_F \equiv x_1 - x_2$ (Feynman x_F). Energy-momentum conservation yields

$$\tau \equiv \frac{M^2}{S} = x_1 x_2. \tag{25}$$

Note that, in Eq. (23), the reaction may also proceed through other vector-meson resonances, such as ρ, ω, ϕ, J/ψ, ψ', ..., Υ, ..., Z^0, etc. Historically, this has led to the major discoveries of J/ψ (a $c\bar{c}$ system), Υ (a $b\bar{b}$ system), and Z^0 (neutral weak boson).

So far the DY production has been the only direct way to measure the structure functions for those hadrons, such as $\pi^\pm$ and $K^\pm$, which can be extracted from a proton synchrotron as a beam but never as a target (due to the very short lifetime). Specifically, the quark distributions of the pion have been determined from the DY production in pion-nucleon collisions - such as the earlier NA3 experiment[20] or the more recent CERN NA10 and Fermilab E615 experiments.[21,22] The form of the distribution is assumed to be the one dictated by naive counting rules, such as $xv(x) = a_v x^\alpha (1-x)^\beta$ with $v(x)$ the valence quark distribution normalized to unity. (Note that such counting rules are valid presumably as $Q^2 \to \infty$.) The prompt photon production in pion-nucleon collisions such as the WA70 experiment[23], as dominated at the parton level by the gluon-photon Compton process $g + Q(\bar{Q}) \to \gamma + Q(\bar{Q})$, may shed some light on the gluon distribution in the pion. While there is room for improvent on the overall quality of the data, extraction of the parton distributions of a pion is based upon the assumption[24] that the parton distributions in a nucleon are well determined from experiments (and thus can be used as the input). As a numerical example, the NA3 Collaboration obtains, for the valence distribution in a pion, $\alpha = 0.45 \pm 0.12$ and $\beta = 1.17 \pm 0.09$ while an analysis[24] of the NA10

Drell-Yan data[21] yields $\alpha = 0.64 \pm 0.03$ and $\beta = 1.08 \pm 0.02$.

As for the kaon, the NA3 Collaboration[25] observes some difference between the valence distribution in K^- and that in π^-:

$$\frac{\bar{u}^{K^-}(x)}{\bar{u}^{\pi^-}(x)} \approx (1-x)^{0.18\pm0.07}, \tag{26}$$

which is perhaps a manifestation of flavor $SU(3)$ symmetry breaking for parton distributions.

Provided that the valence distributions in the pion and proton are known reasonably well, we may compare the DY cross sections for $\pi^- + p$ and $\pi^+ + p$ collisions so that the sea-to-valence ratio in the proton may be determined experimentally.[26,27] The results confirm the standard wisdom that, for the proton, the ratio becomes negligibly small for $x \geq 0.25$ and it is important only at very small x (e.g. ≈ 0.35 at $x \approx 0.05$).

It is clear that DY production in hadronic collisions offers experimental opportunities, alternative to DIS or sometimes completely new, for unraveling the quark parton distributions of a hadron.

(e) Summary:

It is clear that one of the primary objectives of high energy physics experiments is to offer a clear picture concerning parton distributions associated with the proton. The brief presentation given above serves only as an introduction to the subject, rather than as a review of the vast developments over the last two decades.

To summarize the situation, one possibility is to extract parton distributions of the proton by performing a global fit to the existing data up to the date of fitting, as guided by theoretical constraints imposed, e.g., by QCD. This has resulted in from time to time different phenomenological parton distributions such as those obtained by Duke and Owens,[28,29] by Eichten et al. (EHLQ),[30], by Harriman et al. (HMSR)[31], and by Glück et al.[32] As new data always keep coming in while there are occasionally inconsistencies among the existing data, efforts in trying to fit *all* the data do not always pay off. On the other hand, analyses of a mission-oriented experiment often give rise to parton distributions which incorporate certain important aspects. Thus, lack of a global fit does not necessarily reflect that the parton distributions are less reliable (or less useful). For example, parton distributions obtained from the neutrino DIS data, such as in the work of Mattison et al.[33], in general put slightly more stringent constraints on the amount of the heavy quark sea (s, c, etc.) than those obtained from the electron or muon DIS data.

III. Sea Quark Distributions and Generalized Sullivan Processes

In 1972, Sullivan[34] pointed out that, in deep inelastic scattering (DIS) of a nucleon by leptons, the process shown in Fig. 1, in which the virtual photon strikes the pion emitted by the nucleon and smashes the pion into debris, will scale like the original process where the virtual photon strikes and smashes the nucleon itself. In other words, the process will contribute by a finite amount to cross sections in the Bjorken limit, $Q^2 \to \infty$ and $\nu \equiv E_\ell - E'_\ell \to \infty$ with $x \equiv Q^2/(2m_N\nu)$ fixed. Specifically, Sullivan obtained

$$\delta F^{\pi}_{2N}(x, Q^2) = \int_x^1 dy f_\pi(y) F_{2\pi}(\frac{x}{y}, Q^2), \tag{27a}$$

$$f_\pi(y) = \frac{3}{4\pi}\frac{1}{4\pi}(f_{\pi NN}\frac{2m_N}{\mu})^2 y \int_{-\infty}^{t^m} dt \frac{(-t) \mid F_\pi(t) \mid^2}{(-t+\mu^2)^2}, \tag{27b}$$

where $t^m = -m_N^2 y^2/(1-y)$ with m_N the nucleon mass. $F_{2\pi}(x)$ is the pion structure function as would be measured in deep inelastic electron (or muon) scattering with the pion as the target. $\delta F_{2N}(x)$ is the correction to the nucleon structure function due to the Sullivan process. $f_\pi(y)$ is the probability of finding a pion carrying the nucleon momentum fraction y. μ is the pion mass. $f_{\pi NN}$ is the πNN coupling in the form of a pseudovector coupling (as dictated by chiral symmetry) with $F(t)$ characterizing its t-dependence. In what follows, we adopt the dipole form for the sake of illustration:

$$F_\pi(t) = (\frac{\Lambda_\pi^2 - \mu^2}{\Lambda_\pi^2 - t})^2. \tag{27c}$$

The observation of Sullivan was not appreciated until, in 1983, Thomas[35] noted that in the Sullivan process the virtual photon will see most of time the valence distributions in the pion as the probability function $f_\pi(y)$ peaks at $y \approx 0.3$, a region where only valence quarks and antiquarks are relevant. By comparing the excess of the momentum fractions carried by $\bar{u}$ and $\bar{d}$ quarks to that of the $\bar{s}$ quarks, Thomas was then able to use the Sullivan process to set a limit on the total momentum fraction carried by those pions which surround the nucleon.

In a recent paper[12], we (in collaboration with J. Speth at Jülich and G. E. Brown at Stony Brook) observed that Thomas' argument is in fact subject to modifications when contributions due to the kaon cloud, introduced in a way analogous to pions, is suitably incorporated. We went much further when we noticed that the *entire* sea distributions of a nucleon at moderate Q^2, e.g., up to $Q^2 = 20\,GeV^2$ can in fact be attributed to the generalized Sullivan processes, i.e., Figs. 2(a) and 2(b) with the meson-baryon pair (M, B) identified as any of (π, N), (ρ, N), (ω, N), (σ, N), (K, Λ), (K, Σ), (K^*, Λ), (K^*, Σ), (π, Δ), and (ρ, Δ).

To simplify the situations, we take all the coupling constants (which are essentially the $SU(3)$ values) and masses from a previous hyperon-nucleon study[36] and assume dipole forms for all couplings. We find that a universal cutoff mass Λ of $1150\,MeV$ in the Δ/N sector and $1400\,MeV$ in the Λ/Σ sector yields very reasonable results. Here the (time-like) form factors needed to describe Fig. 2(b) are obtained by constraining the amount of strangeness to be the same as that of anti-strangeness, as strange baryons, when produced, decay slowly via weak processes and will survive long enough to be struck by the virtual photon (through electromagnetic interactions). Nevertheless, it was shown[12] that some specific parametrizations of the form factors may improve the fine details but the physics picture remains very much intact.

We compare our model predictions with the neutrino DIS data[18] obtained by the CCFR Collaboration at Fermilab, which provide stringent constraints for the model. Specifically, we obtain, with $< Q^2 >= 16.85\,GeV^2$,

$$\kappa = 0.40 \qquad (\text{exp: } 0.44^{+0.09+0.07}_{-0.07-0.02}), \tag{28a}$$

$$\eta_s = 0.056 \qquad (\text{exp: } 0.057^{+0.010+0.007}_{-0.008-0.002}), \tag{28b}$$

$$R_{\bar{Q}} = 0.161 \qquad (\text{exp: } 0.153 \pm 0.034). \tag{28c}$$

It is clear that the agreement between theory and experiment is excellent.

As reported earlier[12], we find that the shape of the various sea distributions obtained in this way are very similar to that in the corresponding phenomenologically parametrized parton distributions of Eichten et al. (EHLQ)[30], lending strong support toward the conjecture that sea distributions of a hadron at moderate Q^2 come almost entirely from the meson cloud. In addition, we show[12] in Fig. 3 that the ratio $\frac{1}{2}\{\bar{u}(x) + \bar{d}(x)\}/\{u_v(x) + d_v(x)\}$ as a function of x for $Q^2 = 16.85\,GeV^2$, as obtained in our model, is in good agreement with the sea-to-valence ratio extracted from the CDHS data[26] (in triangles) and the Fermilab E615 data (in solid squares).[27] Indeed, consistency among the NA3 data for extracting pion[20] and kaon[25] distributions, the CCFR data,[15,18] and the CDHS data[14,26] emerges nicely within our model calculations.

Using the model to determine possible deviation from the Gottfried sum rule, we have obtained[37]

$$\begin{aligned}\int_{0.002}^{1} \frac{dx}{x}\{F_2^{ep}(x) - F_2^{en}(x)\} &= 0.251, \qquad (\text{EHLQ})^{30}\\ &= 0.177, \qquad (\text{HMSR})^{31}\\ &= 0.235, \qquad (\text{NC: Mattison et al.})^{33}\\ &= 0.235, \qquad (\text{CC: Mattison et al.})^{33}\end{aligned} \tag{29}$$

Here the first two calculations are carried out by using the phenomenologically parametrized valence distributions[30,31] (both at $Q^2 = 4\,GeV^2$) while the third and fourth entries are obtained when we adopt the parton distributions (at $Q^2 = 10\,GeV^2$) which T. S. Mattison *et al.*[33] extracted from weak reactions involving neutral currents (NC) or charge currents (CC). It is useful to stress the point that, as the sea quark distributions are now calculated from generalized Sullivan processes, our results are controlled essentially by the input valence distributions (together with the form factors chosen to reproduce the strength of the observed sea). Such input is considered to be the most reliable piece of the various parton distributions.

To investigate the situation in much greater detail, we plot in Fig. 4 the structure function difference $F_2^p(x) - F_2^n(x)$ as a function of x. The four curves are the predictions using four different input distributions for the nucleon - in dash-dotted curve from the distribution extracted from the neutral-current neutrino data,[33] in dashed curve from the charge-current neutrino data,[33] in dotted curve from the input distribution of Harriman *et al.* (HMSR),[31] and in solid curve from the distributions of Eichten *et al.* (EHLQ).[30] It is clear that the shape of the EHLQ valence distributions performs better than that of the HMSR ones. The QCD evolution softens the valence distributions slightly (from $Q^2 = 4\,GeV^2$ to $10\,GeV^2$) so that the results from the NC and CC neutrino data are more or less consistent with the EHLQ prediction.

We should emphasize that, despite the fact that the integrated value as listed in Eq. (29) may come close to the data, it is *nontrivial* to reproduce as well the shape of the experimental data as a function of x. The curves shown in Fig. 4 reflect directly the shape of the proposed valence distribution convoluted according to Sullivan processes. To see this more clearly, we show in Fig. 5 the structure function difference $F_2^p(x) - F_2^n(x)$ as a function of x in the case of EHLQ[30] by decomposing it into two contributions, the dotted curve from the valence contribution $\frac{1}{3}(u_v(x) - d_v(x))$ and the dashed curve from the calculated sea distribution $\frac{2}{3}(\bar{u}(x) - \bar{d}(x))$. In any event, the general agreement may be taken as an *additional* evidence toward the suggestion that the sea distributions of a hadron, at low and moderate Q^2 (at least up to a few GeV^2), may be attributed primarily to generalized Sullivan processes. This then gives the sea distributions which are not biased by the standard hypthesis Eq. (13) and may be used as input for QCD evolutions to higher Q^2.

The significance of the conjecture of attributing the sea distributions of a hadron at low and moderate Q^2 to its associated meson cloud as generated by strong interaction processes at the hadron level is that we are now able to determine the sea distributions of a hadron from the knowledge of the valence distributions of the various hadrons. The QCD evolution equations then take us from low or moderate Q^2 to very high Q^2. The previously very "fuzzy" gap between low Q^2 (nuclear) physics and large Q^2 (particle) physics is now linked nicely together. In other words, while high energy physics experiments with large Q^2 place stringent constraints on the basic input parameters for nuclear physics, the information gained from nuclear physics experiments such as nucleon-nucleon and hyperon-nucleon scatterings allows us to "predict", among others, the sea distributions of a nucleon, including the detailed strangeness, isospin, and spin information.

Nevertheless, it is of importance to note that, in obtaining our results, we have adjusted the cutoffs to values somewhat below those used for fitting the nucleon-nucleon and hyperon-nucleon scattering data. This of course destroys the existing fits. However, with now the various cutoffs constrained by the deep inelastic scattering data and with the coupling constants (previously fixed to the $SU(3)$ values) adjusted slightly to allow for small flavor $SU(3)$ symmetry breaking, it will be of great interest[38] to see if fits of similar quality may still be obtained.

It is known[39] that the meson-baryon picture with a relatively hard πNN form factor (such as the one expected from our calculations) provides a quantitative understanding of *nuclear physics* phenomena such as the electromagnetic form factors of the deuteron, those of the triton or helium-3, and the near-threshold electrodisintegration of the deuteron, all up to a few GeV^2. Accordingly, we have in mind that the conjecture of generating sea quark distributions in a hadron via generalized Sullivan processes is valid at a few GeV^2 and the QCD evolution then takes us to higher Q^2. The issue is when we should start doing QCD evolution via Altarelli-Parisi equations. In our opinion, the Q^2 must be high enough (or the resolution is good enough) in order to see the substructure (or occurrence of subprocesses) at the quark-gluon level. For studying the violation of the Gottfried sum rule, we take it to be $4\,GeV^2$, which we believe is a reasonable guess. For Q^2 below such value, we believe that quarks and gluons exist only by associating themselves with hadrons and thus Sullivan processes provide a natural way for obtaining parton distributions at these Q^2.

As a summary of what we have described in this section, we note that the meson-exchange model for generating the sea distributions of a nucleon at low and moderate Q^2, say up to $20\,GeV^2$, is capable of *not only* accounting for a variety of high energy physics measurements related to free nucleons *but also* providing a simple framework to understand quantitatively the recent finding by the New Muon Collaboration (NMC) on the violation of the Gottfried sum rule.

IV. Valence Quark Distributions and Light Cone Wave Functions

After having considered in the previous section the possible origin of the sea quark distributions associated with a nucleon (or any other hadron), we turn our attention to the physics related to valence quark distributions. To this end, we first take note of the fact that, to shed light on the physical meaning of the parton model, there were attempts to study field theories, with quantum electrodynamics (QED) in particular, in the infinite-momentum frame, leading eventually to adoption of the light-cone language. Using the ϕ^3 theory as an illustrative example, S. Weinberg[40] showed that many undesirable Feynman diagrams disappear in a reference frame with infinite total momemtum while the contribution of the remaining diagrams is characterized by a new set of rules. However, S. J. Chang and S. K. Ma[41] pointed out that in ϕ^3 theory vacuum diagrams (i.e., diagrams with no external lines) which should vanish according to Weinberg's rule acquire nonvanishing contributions from end points of allowed longitudinal momenta carried by internal particles. Nevertheless, Drell, Levy, and Yan[42] noted that if it is possible to restrict our attention to the time and third components of the electromagnetic current (and inferring the contributions from the transverse components using covariance requirement), then Weinberg's argument holds and no particle of negative longitudinal momentum may enter or leave the electromagnetic vertex. For this reason, the time and third components of the electromagnetic current are referred to as "good currents", suggesting the advantage of quantizing the field theory adopting the light-cone language. Subsequently, Kogut and Soper[43] and later Lepage and Brodsky[44] obtained the complete set of Feynman rules for QED and QCD, respectively. These Feynman rules define the so-called "light-cone perturbation theory"[44] which, as suggested by Lepage and Brodsky[44], may be used for obtaining the hard-scattering amplitudes for high-energy exclusive processes. In conjunction with the suggestion, it was proposed that the hadron wave function may be represented as an infinite series of Fock components. For instance, the pion π^+ may be described in the light-cone language as follows:

$$| \pi^+(P) >= C_0 \mid u\bar{d} > +C_g \mid u\bar{d}g > +C_Q \mid u\bar{d}Q\bar{Q} > + \cdots . \tag{30}$$

where the coefficients C_0, C_g, and C_Q are functions of Q^2.

Specifically, it was shown[44] that, for exclusive processes at sufficiently large Q^2, the contribution from the leading Fock component $\mid u\bar{d} >$ dominates over all the others. Considerable progresses were made by Chernyak and Zhitnitsky[45] who were able to improve upon the hadron wave functions making use of the results from QCD sum rule studies[46] of properties of low-lying hadrons. To describe a specific component in the wave function, we adopt, using light-cone variables,

$$x_i \equiv \frac{p_i^+}{P^+} = \frac{p_{i0} + p_{i3}}{P_0 + P_3}, \qquad \vec{p}_{i\perp} \equiv (p_{i1}, p_{i2}). \tag{31}$$

which are invariant under Lorentz boosts in the z direction. In the light-cone language, moreover, Lorentz boosts in the z direction are purely kinematical - that is, no particles are created nor destroyed. Thus, there is an *invariant* description of the complicated hadron wave function such as Eq. (30) in all frames which are related by Lorentz boosts in the z direction. In this way, we may eventually go over to the infinite-momentum limit ($P_3 \to \infty$) to study Bjorken scaling and its violations. For these reasons, we expect that, if the quark distributions in the parton model can ever be described in terms of wave functions of any sort, the hadron wave functions written in the light-cone language appear to be the best candidate for such a description. Indeed, such aspect has been taken up by different authors.[44,47]

The aim of the present section is as follows: First, we wish to investigate in some detail how quark distributions of a hadron may be linked to the hadron wave function written in the light-cone language, using the pion as our explicit example and keeping track of technical details and approximations. We shall make precise identification of what to calculate and then keep track of terms in transverse momenta. Next, we use the leading pion wave function as constrained by QCD sum rules to determine the fraction of the valence distribution that may be attributed to the leading Fock component in the pion wave function. We then apply the specific proposal[12] of using generalized Sullivan processes to generate the entire valence quark distributions from the valence quark distributions calculated from the leading Fock component.

What is implicit in our approach is that, similar to the study of QCD sum rules,[44] there is an assumed optimal region of Q^2, say around about $1\,GeV^2$, in which we believe our procedure of obtaining valence and sea quark distributions is best justified. This may be explained as follows: At very large Q^2 (say, $>> 1\,GeV^2$), the coefficient C_0 for the leading Fock component is tiny so that determination of the valence quark distributions from it is a completely inefficient task - yet, there is not any efficient way to obtain contributions from the very complicated nonleading Fock components. On the other hand, if Q^2 is small (say, comparable with the confinement scale Λ^2_{QCD}), the description of the hadron wave function in terms of different Fock components, such as Eq. (30), no longer makes much sense either. Accordingly, one need to work with an intermediate scale, such as $Q^2 \approx 1\,GeV^2$, such that the contribution from the leading Fock component can be calculated while effects from the rest of the wave function may be organized naturally using the meson-baryon picture.[12]

The approach may be contrasted with the pioneering work of Jaffe and

Ross[48], who considered how the structure functions of a hadron can be linked to the hadron wave function in a bag model. It is clear that there are uncertainties, many of which are difficult to resolve, in trying to understand the distributions in the parton model using a naive quark model often phrased in configuration space. For instance, the center-of-mass (CM) problem is a nasty problem to resolve especially in a relativistic model. By proposing to solve the problem using light-cone wave functions obtained via QCD sum rule studies, we may in fact bypass many ambiguities involved in the Jaffe-Ross procedure.

We begin by considering the derivation of the differential cross section for the deep inelastic scattering (DIS) $e(\ell) + h(P,\lambda) \to e(\ell') + X$,

$$d\sigma = \frac{d^3\ell'}{2\ell_0'(2\pi)^3} \frac{1}{((\ell\cdot P)^2 - m_\ell^2 M^2)^{1/2}} \frac{e^4}{q^4} L^{\mu\nu} 4\pi M W_{\mu\nu}, \tag{32}$$

where $L^{\mu\nu}$ is the tensor for the probing lepton while the hadronic tensor $W_{\mu\nu}$ is specified by

$$\begin{aligned} W_{\mu\nu} &\equiv \frac{1}{4\pi M}\sum_X (2\pi)^4 \delta^4(P+q-P_X) < P,\lambda \mid J_\mu(0) \mid X >< X \mid J_\nu(0) \mid P,\lambda > \\ &= \frac{1}{4\pi M}\int d^4x e^{-iq\cdot x} < P,\lambda \mid [J_\mu(x),\, J_\nu(0)] \mid P,\lambda > \qquad (33a) \\ &\equiv W_1(q^2,\nu)(\delta_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}) + W_2(q^2,\nu)(P_\mu - \frac{(P\cdot q)q_\mu}{q^2})(P_\nu - \frac{(P\cdot q)q_\nu}{q^2}). \qquad (33b) \end{aligned}$$

Note that the definition (33b) is for a spinless target and may easily be generalized to include the spin for hadrons such as the proton.

The basic idea consists primarily of calculation of the matrix element $< P,\lambda \mid [J_\mu(x),\, J_\nu(0)] \mid P,\lambda >$ using as the input the operator obtained from the light-cone perturbation theory (augmented with effects due to quark and gluon condensates, if so desirable) and the wave function of the leading Fock component (which is constrained by QCD sum rules). From the results, we may then identify the structure function $W_2(q^2,\nu)$ (or $F_2(x,Q^2)$) and sort out the exact relation between a specific valence quark distribution and the given light-cone wave function.

For a spinless target such as a pion, we may consider the frame in which the hadron is co-moving with the virtual photon:

$$q_\mu = (0,0,q_3,iq_0), \qquad P_\mu = (0,0,P_3,iE), \tag{34}$$

so that

$$\begin{aligned} F_1(x_B,Q^2) &\equiv MW_1(Q^2,\nu) = MW_{11}, \\ F_2(x_B,Q^2) &\equiv M^2\nu W_2(Q^2,\nu), \\ W_2(Q^2,\nu)(M^2+\frac{Q^2}{4x_B^2}) &= 2W_{+-} + W_{11}, \end{aligned} \tag{35}$$

with $x_B \equiv Q^2/(2M\nu)$ (the Bjorken x) and the $(+-)$ component is specified in the same sense as we define $(\eta,\, H)$ or $(p^+,\, p^-)$:

$$\begin{aligned} \eta &\equiv \frac{1}{\sqrt{2}}(\hat{p}_0+\hat{p}_3) \equiv \frac{1}{\sqrt{2}}p^+, \\ H &\equiv \frac{1}{\sqrt{2}}(\hat{p}_0-\hat{p}_3) \equiv \frac{1}{\sqrt{2}}p^-, \end{aligned} \tag{36}$$

with the old variables (in the usual instant-form dynamics) denoted by careted symbols. In what follows, we use the notation of Kogut and Soper[43] whenever the light-cone language is adopted. In particular, we write

$$\begin{aligned} \psi(x) = \frac{1}{(2\pi)^{3/2}}\frac{1}{\sqrt{2}}\int d^2p_\perp \int \frac{d\eta}{\eta}\sum_\lambda \{&u(p,\lambda)e^{-ip\cdot x}b(p,\lambda) \\ &+ v(p,\lambda)e^{ip\cdot x}d^\dagger(p,\lambda)\}, \end{aligned} \tag{37}$$

$$\{b(p,\lambda),\, b^\dagger(p',\,\lambda')\} = \delta_{\lambda\lambda'}(2\pi)^3 2\eta\delta(\eta-\eta')\delta^2(p_\perp - p'_\perp), \tag{38a}$$

$$\{d(p,\lambda),\, d^\dagger(p',\,\lambda')\} = \delta_{\lambda\lambda'}(2\pi)^3 2\eta\delta(\eta-\eta')\delta^2(p_\perp - p'_\perp), \tag{38b}$$

$$\sum_\lambda u(p,\lambda)\bar{u}(p,\lambda) = \gamma^\mu p_\mu + m, \quad \sum_\lambda v(p,\lambda)\bar{v}(p,\lambda) = \gamma^\mu p_\mu - m. \tag{38c}$$

The electromagnetic current $J_\mu(x)$ is specified by

$$J_\mu(x) =: \bar{\psi}(x)\gamma_\mu\psi(x) :, \tag{39}$$

where $: AB :$ denotes the normal-ordered product of A and B. (Note that the electric charge Q_i may be inserted at the end of manipulations.) For the commutator $[: AB :, : CD :]$, we may apply Wick's theorem separately to the products $: AB :: CD :$ and $: CD :: AB :$ and then take the difference. In this way, we obtain, with A, B, C, and D fermion operators,

$$\begin{aligned}&[: AB :, : CD :]\\ =&(< AD >_0 + : AD :)\{B, C\} + (< CA >_0 + : CA :)\{B, D\}\\ &-\{A, D\}(< CB >_0 + : CB :) - \{A, C\}(< BD >_0 + : BD :),\end{aligned} \tag{40}$$

where we have adopted the standard definitions, $[x, y] \equiv xy - yx$, $\{x, y\} \equiv xy + yx$, and $< xy >_0 \equiv < 0 \mid xy \mid 0 >$ (with $\mid 0 >$ the vacuum or the ground state).

Eq. (40) may be used to obtain the commutator $[J_\mu(x), J_\nu(y)]$, yielding

$$\begin{aligned}&[J_\mu(x), J_\nu(y)]\\ =&\gamma_{\mu ab}\gamma_{\nu cd}(< \bar{\psi}_a(x)\psi_d(y) >_0 + : \bar{\psi}_a(x)\psi_d(y) :)\{\psi_b(x), \bar{\psi}_c(y)\}\\ &-\gamma_{\mu ab}\gamma_{\nu cd}\{\bar{\psi}_a(x), \psi_d(y)\}(< \bar{\psi}_c(y)\psi_b(x) >_0 + : \bar{\psi}_c(y)\psi_b(x) :).\end{aligned} \tag{41}$$

Now, consider π^+ as example. The leading Fock component of the pion may be described as follows, with the $\vec{p}_\perp$ dependence explicitly taken into account:

$$\begin{aligned}\mid P >\equiv \int \frac{d\eta_1 d^2p_{1\perp}}{(2\pi)^3 2\eta_1} \int \frac{d\eta_2 d^2p_{2\perp}}{(2\pi)^3 2\eta_2}(2\pi)^3 2\eta\delta(\eta_1 + \eta_2 - \eta)\delta^2(\vec{p}_{1\perp} + \vec{p}_{2\perp} - \vec{P}_\perp)\\ \cdot \phi_\pi(p_1, \lambda_1; p_2, \lambda_2; P)d^\dagger(p_2, \lambda_2)b^\dagger(p_1, \lambda_1) \mid 0 >,\end{aligned} \tag{42a}$$

$$\phi_\pi(p_1, \lambda_1; p_2, \lambda_2; \vec{P}_\perp = 0) = \Psi(x, k_\perp^2)\bar{u}(p_1, \lambda_1)(\gamma \cdot P + M)\gamma_5 v(p_2, \lambda_2), \tag{42b}$$

$$\Psi(x, k_\perp^2) = A_\pi(\frac{1-x}{x})^{1/2} exp(-\frac{k_\perp^2 + m^2}{8\beta^2 x(1-x)}), \tag{42c}$$

where $\vec{k}_\perp \equiv \vec{k}_{1\perp} \equiv \vec{p}_{1\perp} - x_1\vec{P}_\perp$ and $x \equiv x_1 \equiv \eta_1/\eta$, etc. Instead of the wave function adopted by Chernyak and Zhitnitsky[45] or by Dziembowski[49] which reproduces the QCD sum rules approximately, we may choose an explicit form for the wave function $\psi(x, k_\perp^2)$ that reproduces better the sum rules at $\mu_0^2 = (500\,MeV)^2$, with $\xi \equiv x_1 - x_2 = 2x - 1$:

$$< \xi^2 >= 0.46, \qquad < \xi^4 >= 0.30, \qquad < \xi^6 >= 0.21. \tag{43}$$

For a choice of $m = 330\,MeV$ (constituent quark mass), $M = 600\,MeV$ ("mock" pion mass, i.e., the mass before the $\rho - \pi$ mass splitting as may be caused by some spin-spin interaction), and $\beta = 500\,MeV$ (which characterizes the size of the pion), we obtain exactly the values listed in Eq. (43). Note that these input parameters are very similar to what was used by Dziembowski[49] although the detailed form for $\Psi(x, k_\perp^2)$ is different. Note that[46] the normalization is fixed by the condition:

$$< 0 \mid \bar{d}(0)\gamma_\mu\gamma_5 u(0) \mid \pi^+(P) >= \sqrt{2} f_\pi P_\mu, \tag{44}$$

with f_π the pion decay constant ($94\,MeV$). This may be justified in perturbative QCD as the current cannot connect the nonleading Fock components (as given by Eq. (30)) to the vacuum.

It is straightforward, albeit a little tedious, to evaluate the matrix element $< P \mid : [J_\mu(x), J_\nu(0)] : \mid P >$ using Eqs. (41) and (42) with the aid of Eqs. (37) and (38). First, we substitute Eqs. (42a) and (41) into Eq. (33a) and then use Eqs. (37), (38a), and (38b) to eliminate all creation and annihilation operators. Subsequently, we make the substitution indicated by Eq. (42b) and use Eq. (38c) to sum up all the spin indices, leading to the various traces of products of γ-matrices. These traces can easily be evaluated and a relatively compact form for $4\pi M W_{\mu\nu}$ is then obtained. According to Eqs. (35), it is sufficient to pick out only the components W_{11} and W_{+-}. The final results for $W_{\mu\nu}$ can be split into a contribution from the u-quark and another one from the $\bar{d}$-quark, with the two contributions equal to each other apart from the charge factors Q_i^2. The contribution from the u-quark, with the charge factor $(2/3)^2$ neglected from the expressions, is recorded below:

$$
\begin{aligned}
W_{11}^u = & \frac{8\pi}{x_0 M} \int \frac{dx}{x(1-x)} \int dk_\perp^2 \mid \Psi(x, k_\perp^2) \mid^2 \cdot \\
& \left(\frac{k_\perp^2 + m^2}{x} + xM^2 + 2Mm\right)\left(\frac{k_\perp^2 + m^2}{1-x} + (1-x)M^2 + 2Mm\right)\left(k_\perp^2 - \frac{1}{2}Q^2\right) \\
& \{\delta(k_\perp^2 + m^2 - \frac{x(x+x_0)}{x_0^2} Q^2) - \delta(k_\perp^2 + m^2 - \frac{x(x-x_0)}{x_0^2} Q^2)\}, \qquad (45a) \\
W_{+-}^u = & \frac{8\pi}{x_0 M} \int \frac{dx}{x(1-x)} \int dk_\perp^2 \mid \Psi(x, k_\perp^2) \mid^2 \cdot \\
& \left(\frac{k_\perp^2 + m^2}{x} + xM^2 + 2Mm\right)\left(\frac{k_\perp^2 + m^2}{1-x} + (1-x)M^2 + 2Mm\right)\left(k_\perp^2 + m^2\right) \\
& \{\delta(k_\perp^2 + m^2 - \frac{x(x+x_0)}{x_0^2} Q^2) - \delta(k_\perp^2 + m^2 - \frac{x(x-x_0)}{x_0^2} Q^2)\}, \qquad (45b)
\end{aligned}
$$

with $x_0 \equiv q^+/P^+$. Note that Eqs. (45a) and (45b) contain nontrivial factors other than $\mid \Psi(x, k_\perp^2) \mid^2$ - a fact making the linkage between the wave function and the corresponding valence quark distributions somewhat nontrivial. Nevertheless, derivation of Eqs. (45a) and (45b) from Eqs. (41) and (42), with the aid of Eqs. (37) and (38), is indeed a straightforward task - that is, it does not involve any approximation nor assumption. It is essential to keep in mind this specific aspect when one tries to distinguish our results from those obtained or adopted previously.[47,44]

In the Bjorken limit ($Q^2 \to \infty$, $\nu \to \infty$, $x_B \to x_0$ with x_B held fixed), we obtain from Eqs. (45a), (45b), and (35) the important results:

$$F_2^u(x_0, Q^2) \equiv M^2 \nu W_2^u(q^2, \nu) \rightarrow x_0 u(x_0), \tag{46a}$$

$$F_1^u(x_0, Q^2) \equiv M W_1^u(q^2, \nu) \rightarrow \frac{1}{2} u(x_0), \tag{46b}$$

$$u(x_0) = \frac{8\pi}{x_0(1-x_0)} \int dk_\perp^2 \mid \Psi(x_0, k_\perp^2) \mid^2 \cdot$$
$$\cdot \left(\frac{k_\perp^2 + m^2}{x_0} + x_0 M^2 + 2Mm\right)\left(\frac{k_\perp^2 + m^2}{1-x_0} + (1-x_0)M^2 + 2Mm\right). \tag{46c}$$

Thus, the contribution from the leading Fock component of the hadron wave function to the valence u-quark distribution can be *unambiguously* identified.

It is essential to note that the well-known relation $2x_0 F_1(x_0) = F_2(x_0)$ comes about automatically. The fact that the wave function is subject to the QCD sum rule constraints, Eqs. (43) and (44), adds some credence to the QCD light-cone perturbation theory.[44] We believe it is of great importance to take note that there is in fact a clear linkage between the parton distributions extracted from the DIS experiments and the light-cone wave functions constrained by QCD sum rule studies. Studies along this line will undoubtedly help to unravel the long-standing mystery concerning the physics of the parton model and may in fact lead to unification of the previously loosely related theoretical ideas - the QCD light-cone perturbation theory,[44] the QCD sum rule method,[46] and the quark parton model of Feynman.[2]

Indeed, using the wave function Eq. (42) that is constrained by the QCD sum rules (Eq. (43)), we find that the resultant $u(x_0)$ gives rise to $\int dx_0 u(x_0) = 0.40$ (the number of the valence u-quark). Of course, this is true at $Q^2 = (0.5\, GeV)^2$ where the QCD sum rule result, Eq. (43), has been obtained. There is no *a priori* reason why the leading Fock component in the wave function (i.e. the first term in Eq. (30)) already gives a substantial portion (40%) of the valence distributions in the pion. Nevertheless, this result is quite comforting as one expects[50] that generalized Sullivan processes[34,12], which provide an efficient way to take into account the remaining Fock components in the infinite series (Eq. (30)), may generate about another half of the valence distributions. In this way, we might be able to understand both the valence and sea quark distributions reasonably well.[50] We shall mention here only some typical results related to valence quark distributions.

As discussed elsewhere[12,50], the idea of using generalized Sullivan processes to generate the entire valence and sea quark distributions is based upon the belief that the various Fock components as appearing on the RHS of Eq. (30) should organize themselves naturally into the various hadrons - as known to be true at low and moderate Q^2. For instance, the first term represents the "core" or "bare" pion, the third term a combination of two-meson states, etc. An important aspect is that the non-leading Fock components also contribute to valence distributions. It is clear that the valence distrbutions must obey the valence number sum rules - e.g. adding up to one *up* quark and one *down* antiquark in π^+. The valence number sum rules thus serve as an important guideline when one takes into account contributions from generalized Sullivan processes. However, there is little reason why the naive counting rules, presumably valid at $Q^2 \to \infty$, would be observed at low and moderate Q^2, although one may exploit the uncertainty related to the wave function $\Psi(x, k_\perp^2)$ in order to obtain a valence distribution in reasonable agreement with naive counting rules.

Numerically, we adopt the following ansatz and compare the prediction on valence distributions with the NA3 result: The valence distribution is taken to be the one that is calculated from the ligh-cone pion wave function (which satisfies the QCD sum rules), with the various couplings (including $\rho\pi\pi$, $K^*\bar{K}\pi$, and $\bar{K}^*K\pi$ vertices, which enter the relevant Sullivan processes) adjusted to reproduce the valence number sum rules. The $\rho\pi\pi$, $K^*\bar{K}\pi$, and $\bar{K}^*K\pi$ couplings are taken from meson-meson scattering studies[50]. The Q^2 is taken to be $3\,GeV^2$. Note that the resultant $\rho\pi\pi$ form factor (with ρ in the t channel) is $2450\,MeV$, which is not very far from that[51] obtained from fitting the extracted phase shifts in $\pi\pi$ and πK scatterings ($\Lambda_\rho = 1600\,MeV$). As there is little clue on the amount of the strangeness in a pion, we choose $\Lambda_{K^*} = 4000\,MeV$ which respresents a similar increase over that[51] used in the study of meson-meson scatterings. Note that Λ_K and Λ_π are adjusted to ensure that quarks and antiquarks are produced in pairs. This yields $\Lambda_\pi = 1882\,MeV$ and $\Lambda_K = 3480\,MeV$.

As a result, the integrated numbers of mesons in the "cloud" may be determined as follows: $\int f_\pi(y)dy = \int f_\rho(y)dy = 0.477$ and $\int f_{K^*}(y)dy = \int f_{\bar{K}}(y)dy$ $= 0.271$. The momentum fractions carried by the various partons (in π^+) are $<x>_u = <x>_{\bar{d}} = 21.5\%$, $<x>_{\bar{u}} = <x>_d = 2.5\%$, $<x>_s = <x>_{\bar{s}} = 3.5\%$, and $<x>_g = 45\%$. All these results appear to be rather reasonable.[24]

In Fig. 6, we show the valence momentum distributions obtained from the above calculation. Using our routine to evolve the valence distributions *via* Altarelli-Parisi equations from $Q^2 = 3\,GeV^2$ to $Q^2 = 25\,GeV^2$ (which has the primary effect of softing slightly the distributions (i.e. shifting the weight to the smaller x region), the result is displayed as a solid curve. For comparison, the valence momentum distribution (at $Q^2 \approx 25\,GeV^2$) obtained by the NA3 Collaboration[20] in fitting to their Drell-Yan data using $\pi^\pm$ beams is shown as a dash-dotted curve. Although there is some uncertainty related to the wave function $\Psi(x, k_\perp^2)$ (as QCD sum rule results, Eq. (43), do not fix the wave function unambiguously), our result corresponds to a distribution $xv(x) = a_v x^\alpha (1-x)^\beta$ with a value of α closer to the NA10 data than the NA3 data but with a value of β considerably larger than both data.

To sum up this section, we have considered, using the pion as the example, the question of how valence quark distributions of a hadron may be linked to the hadron wave function written in the light-cone language. Specifically, we use the leading pion wave function that is constrained by the QCD sum rules, and find that, at $Q^2 \approx (0.5\,GeV)^2$, the leading Fock component accounts for about 40 % of the observed valence quark distributions in the pion. The question of how to generate the entire valence quark distributions from the valence quark distribution calculated from the leading Fock component is briefly discussed using the specific ansatz proposed recently by Hwang, Speth, and Brown.[12,50]

V. Summary

The quark parton model of Feynman, which has been used for analyses of high energy physics experiments, invokes a set of parton distributions in the description of the nucleon structure (the probability concept), contrary to the traditional use of wave functions in nuclear and medium energy physics for structural studies (the amplitude concept).

In this paper, I have reviewed in Section 2 briefly how the various parton distributions of a nucleon may be extracted from high energy physics experiments. I then proceed to consider in Section 3 how the sea distributions of a free nucleon at low and moderate Q^2 (e.g., up to $20\,GeV^2$), may be obtained in the meson-baryon picture, a proposal made by Hwang, Speth, and Brown.[12] Using the form factors associated with the couplings of mesons to baryons such as πNN, $\pi N\Delta$, and $KN\Lambda$ couplings which are constrained by the CCFR neutrino data, we find that the model yields predictions consistent with the CDHS and Fermilab E615 data. We also find that the recent finding by the New Muon Collaboration (NMC) on the violation of the Gottfried sum rule can be understood quantitatively.

Finally, we have considered in Section 4, using the pion as the example, how valence quark distributions of a hadron may be linked to the hadron wave function written in the light-cone language. Specifically, we use the leading pion wave function that is constrained by the QCD sum rules, and find that, at $Q^2 \approx (0.5\,GeV)^2$, the leading Fock component accounts for about 40 % of the observed valence quark distributions in the pion. The question of how to generate the entire valence quark distributions from the valence quark distribution calculated from the leading Fock component is briefly discussed again using the specific proposal of Hwang, Speth, and Brown.[12,50]

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Figure Captions

Fig. 1. The processes considered originally by Sullivan.[34]

Fig. 2. The generalized Sullivan processes: (a) the virtual photon strikes the cloud meson, and (b) the virtual photon strikes the recoiling baryons. Both scale in the Bjorken limit. The meson and baryon pair (M, B) includes (π, N), (ρ, N), (ω, N), (σ, N), (K, Λ), (K, Σ), (K^*, Λ), (K^*, Σ), (π, Δ), and (ρ, Δ).

Fig. 3. The ratio $\frac{1}{2}\{\bar{u}(x) + \bar{d}(x)\}/\{u_v(x) + d_v(x)\}$ as a function of x for $Q^2 = 16.85\,GeV^2$, shown as a function of x. The CDHS data[26] are shown in triangles and the Fermilab E615 data in solid squares.[27]

Fig. 4. The structure function difference $F_2^p(x) - F_2^n(x)$ shown as a function of x. The four curves are our predictions using four different input valence distributions for the nucleon - in dash-dotted curve from the distribution extracted from the neutral-current neutrino data,[33] in dashed curve from the charge-current neutrino data,[33] in dotted curve from the input distribution of Harriman *et al.*,[31] and in solid curve from the distributions of Eichten *et al.*.[30]

Fig. 5. The structure function difference $F_2^p(x) - F_2^n(x)$ as a function of x in the case of EHLQ[30] is decomposd into two contributions, the dotted curve from the valence contribution $\frac{1}{3}(u_v(x) - d_v(x))$ and the dashed curve from the calculated sea distribution $\frac{2}{3}(\bar{u}(x) - \bar{d}(x))$.

Fig. 6. The valence momentum distributions in the pion. Using our routine to evolve the valence distributions *via* Altarelli-Parisi equations from $Q^2 = 3\,GeV^2$ to $Q^2 = 25\,GeV^2$, the result is displayed as a solid curve. For comparison, the valence momentum distribution (at $Q^2 \approx 25\,GeV^2$) obtained by the NA3 Collaboration[20] in fitting to their Drell-Yan data using $\pi^{\pm}$ beams is shown as a dash-dotted curve.

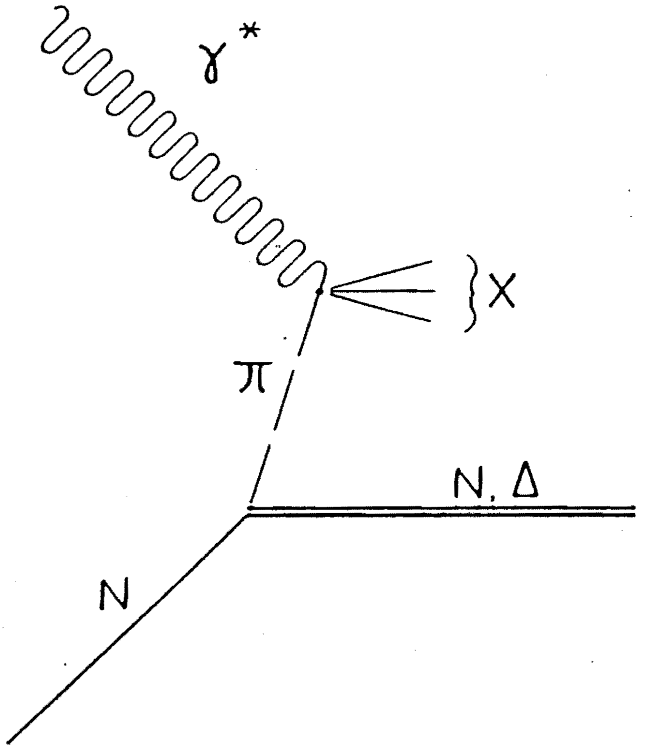

Figure 1

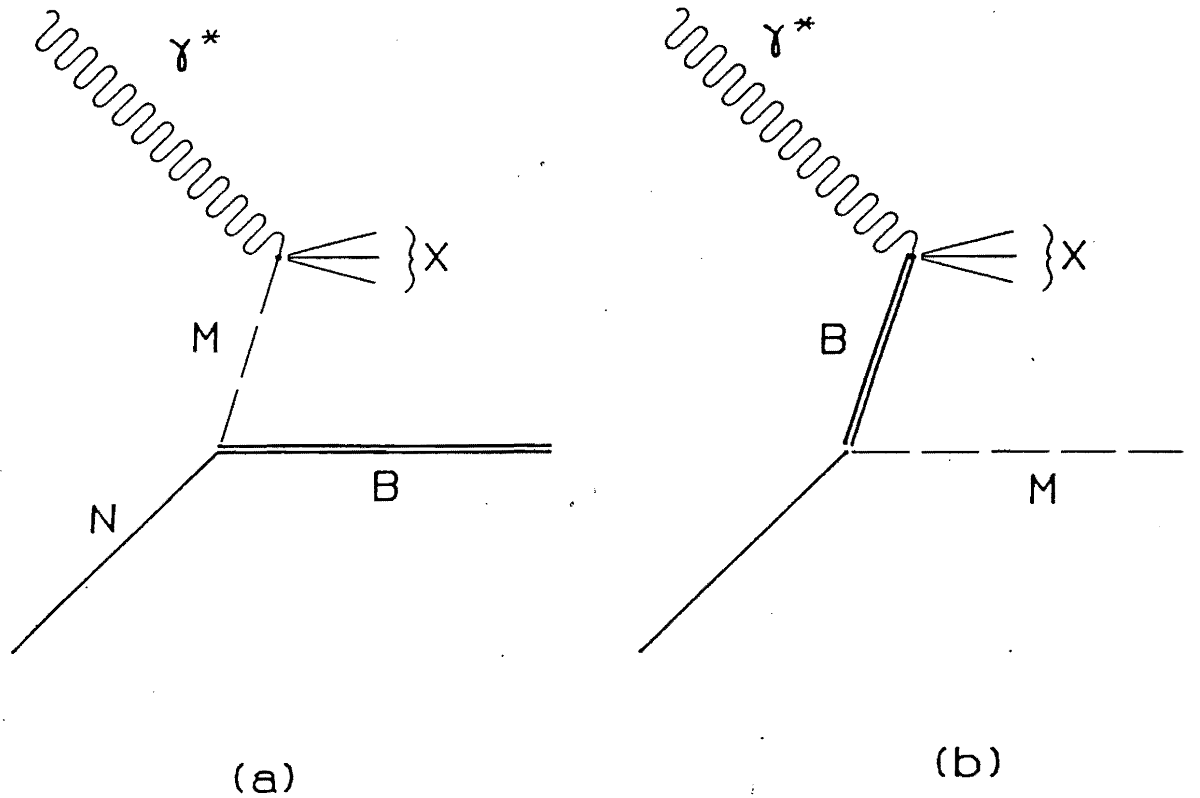

Figure 2

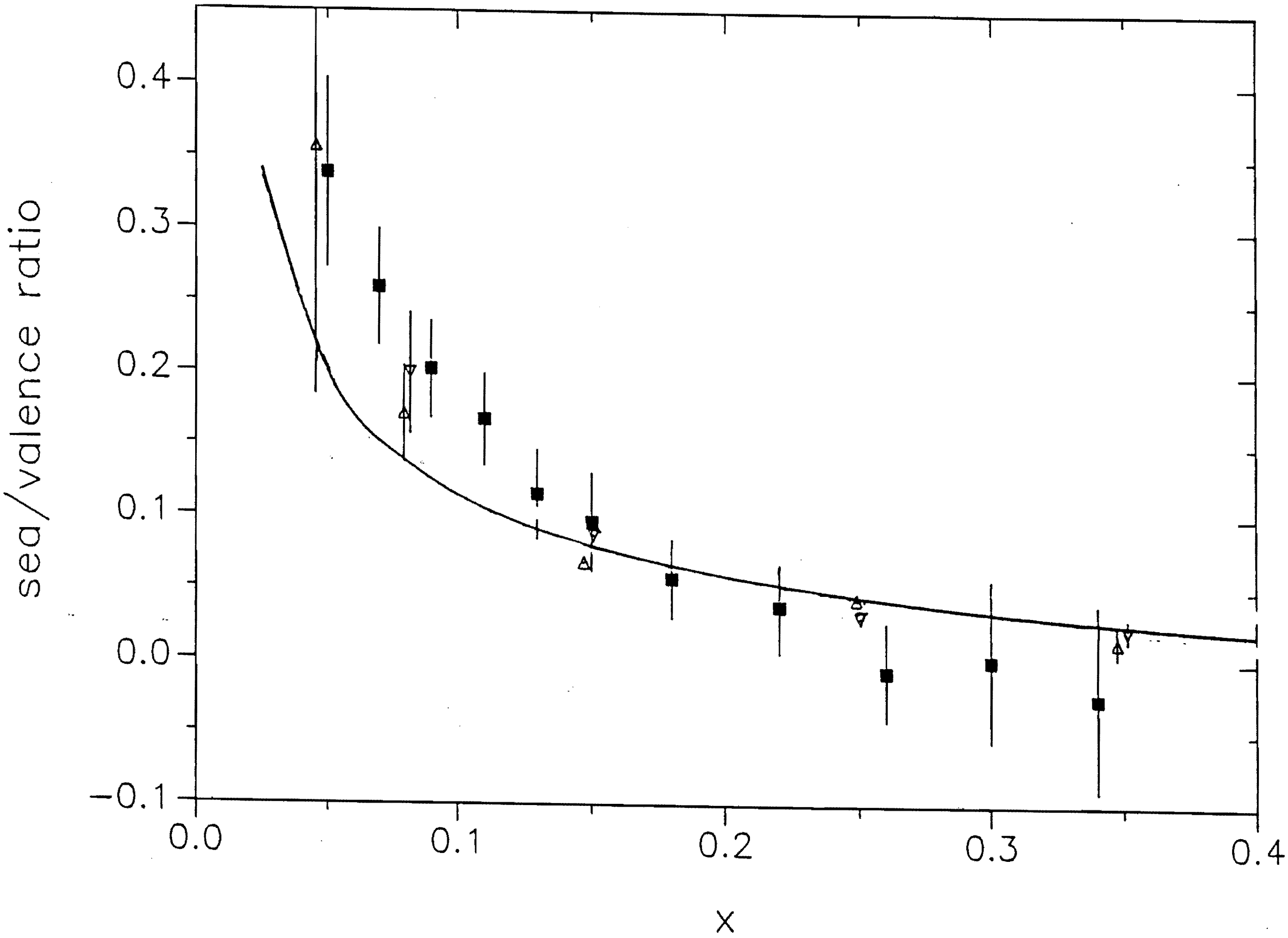

Figure 3

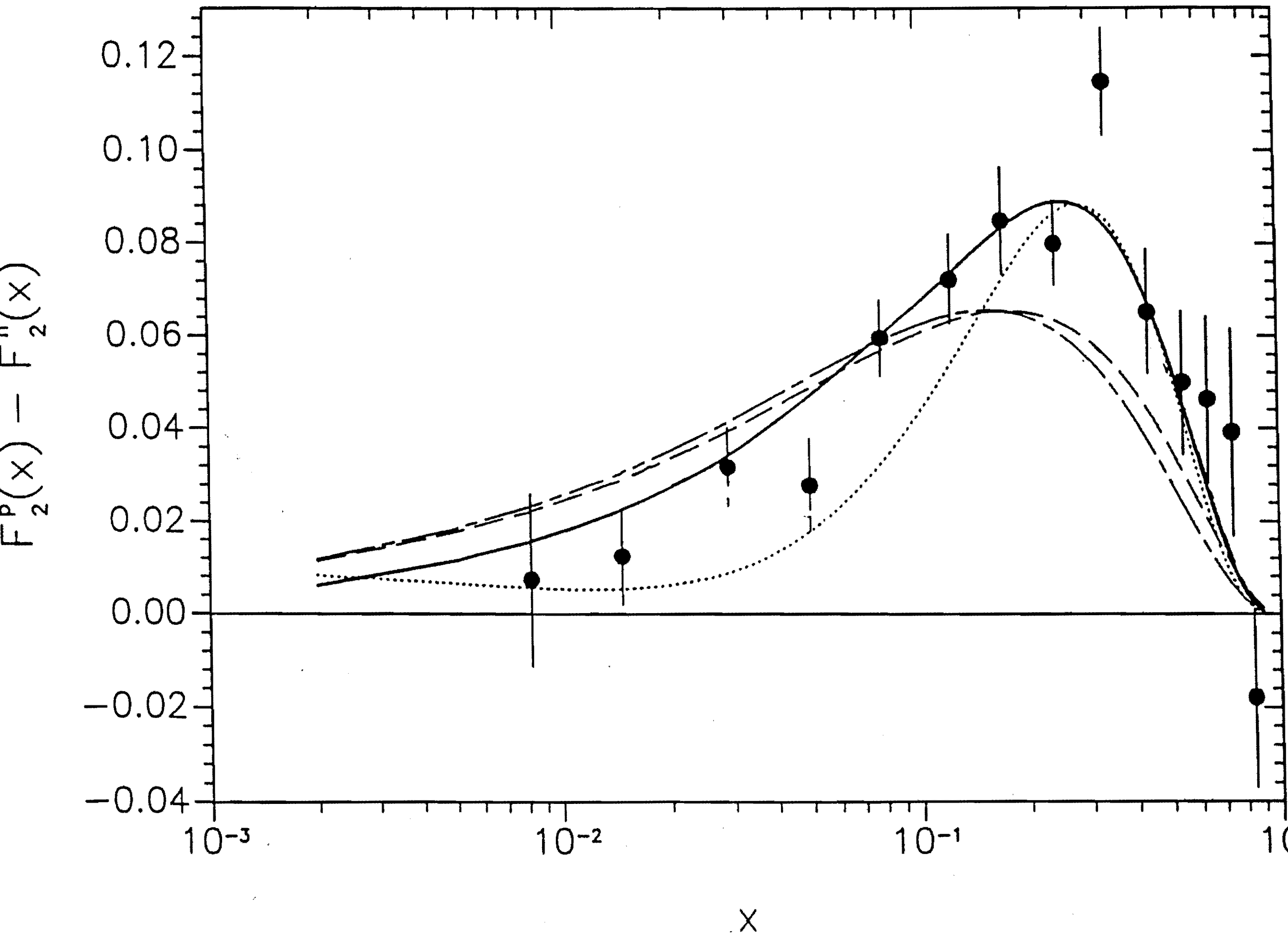

Figure 4

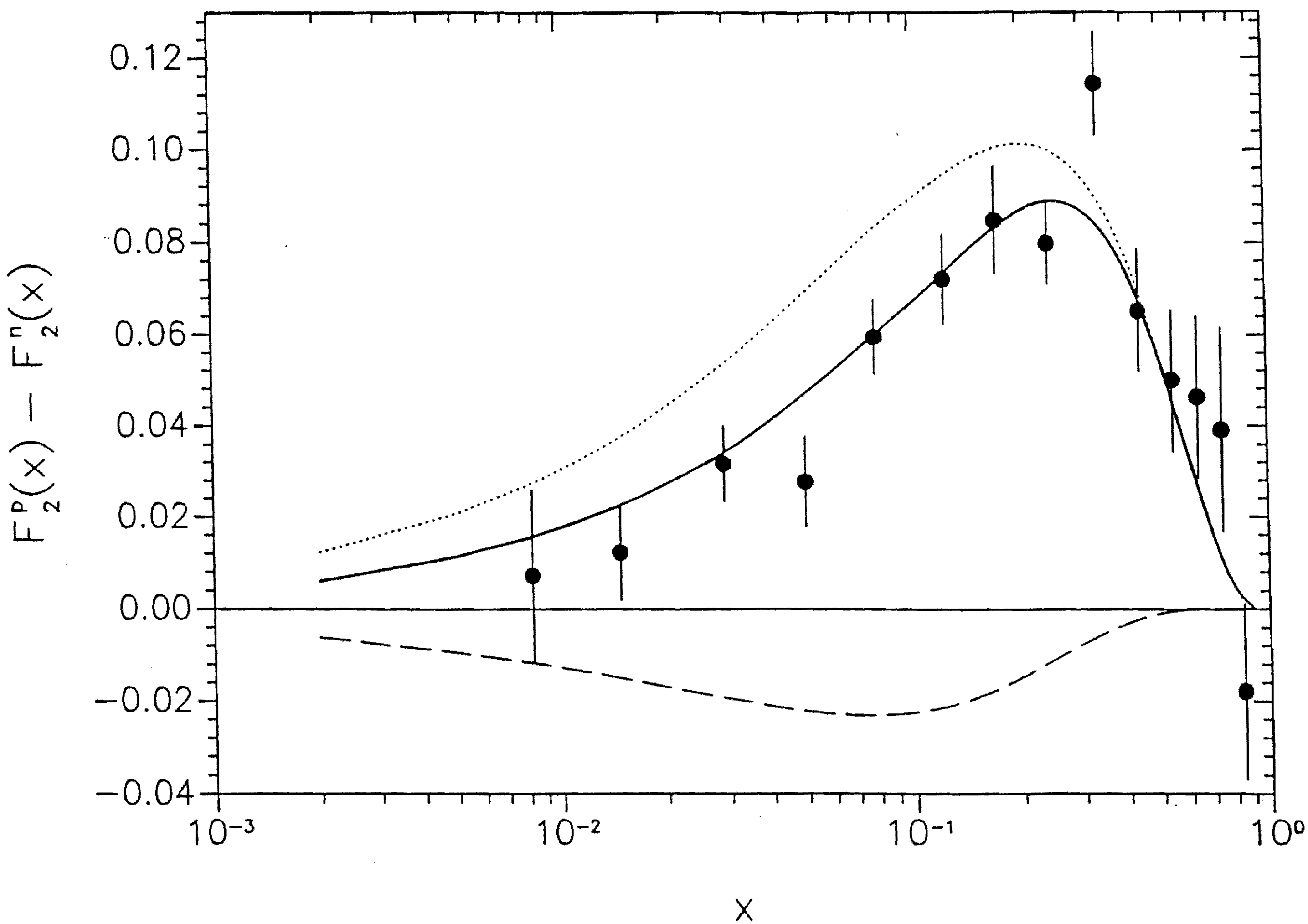

Figure 5

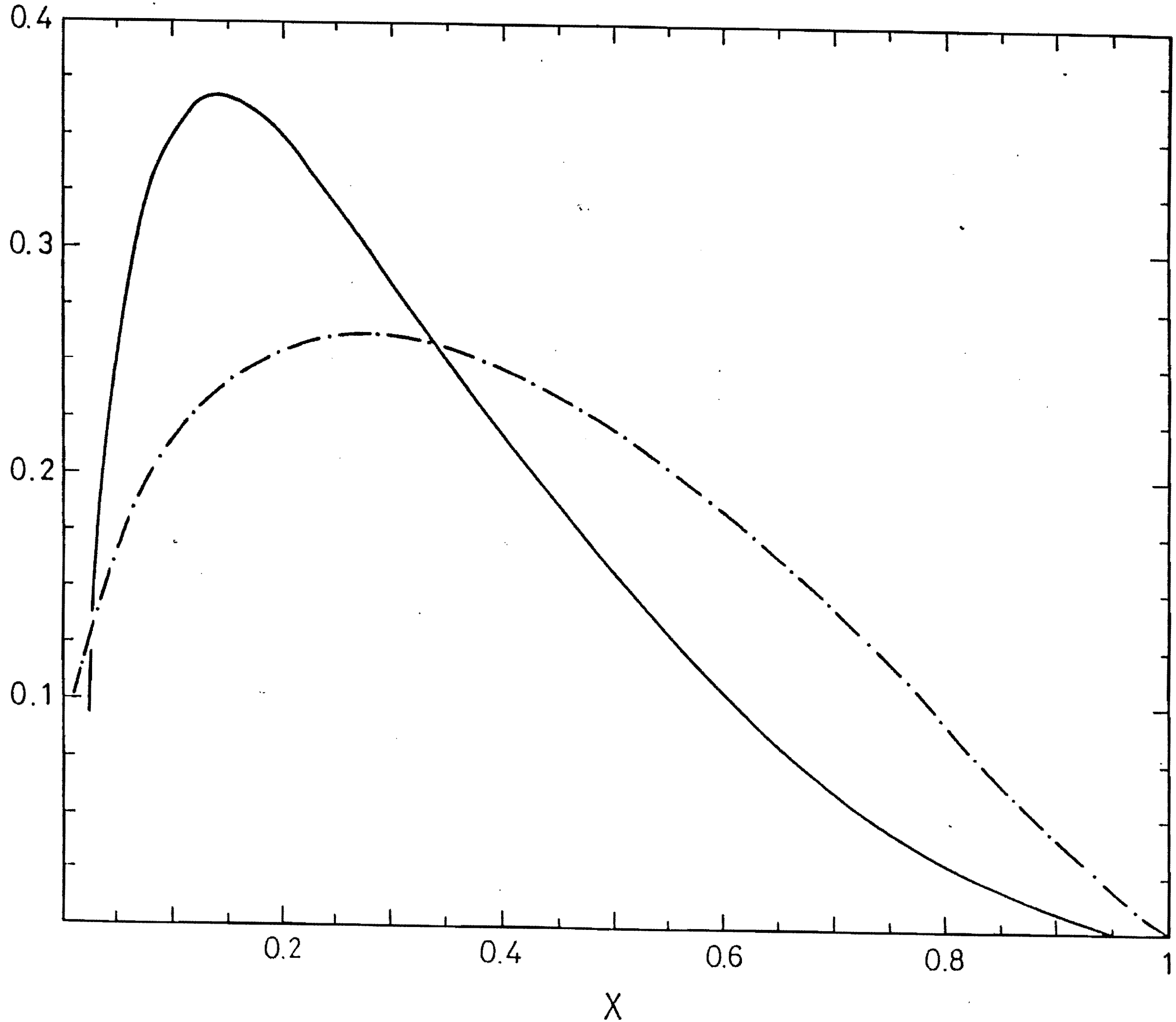

Figure 6