

PUBLISHED VERSION

Melnitchouk, Wolodymyr; Speth, J.; Thomas, Anthony William
[Dynamics of light antiquarks in the proton](#) Physical Review D, 1998; 59(1):014033

©1998 American Physical Society

<http://link.aps.org/doi/10.1103/PhysRevD.59.014033>

PERMISSIONS

<http://publish.aps.org/authors/transfer-of-copyright-agreement>

“The author(s), and in the case of a Work Made For Hire, as defined in the U.S. Copyright Act, 17 U.S.C.

§101, the employer named [below], shall have the following rights (the “Author Rights”):

[...]

3. The right to use all or part of the Article, including the APS-prepared version without revision or modification, on the author(s)' web home page or employer's website and to make copies of all or part of the Article, including the APS-prepared version without revision or modification, for the author(s)' and/or the employer's use for educational or research purposes.”

17th April 2013

<http://hdl.handle.net/2440/10907>

Dynamics of light antiquarks in the proton

W. Melnitchouk*

Institut für Kernphysik, Forschungszentrum Jülich, D-52425 Jülich, Germany

J. Speth†

Institut für Kernphysik, Forschungszentrum Jülich, D-52425 Jülich, Germany

A. W. Thomas‡

Department of Physics and Mathematical Physics, and Special Research Centre for the Subatomic Structure of Matter, University of Adelaide, Adelaide 5005, Australia

(Received 4 June 1998; published 8 December 1998)

We present a comprehensive analysis of the recent data from the E866 experiment at Fermilab on Drell-Yan production in pD and pp collisions, which indicates a non-trivial x dependence for the asymmetry between \bar{u} and \bar{d} quark distributions in the proton. The relatively fast decrease of the asymmetry at large x suggests the important role played by the chiral structure of the nucleon, in particular the πN and $\pi\Delta$ components of the nucleon wave function. At small x the data require an additional non-chiral component, which may be attributed to the Pauli exclusion principle, as first suggested by Field and Feynman. [S0556-2821(99)05701-X]

PACS number(s): 12.39.Ba, 13.85.Qk, 14.20.Dh, 24.85.+p

I. INTRODUCTION

The recent Drell-Yan experiment by the E866/NuSea Collaboration at Fermilab [1] provides the best information yet on the detailed structure of the light antiquark sea of the proton. Although previous experiments by the New Muon Collaboration [2] on the difference between the proton and neutron structure functions established that an asymmetry in the sea exists, they yielded direct information only on the first moment of the antiquark asymmetry. The earlier Drell-Yan experiment by the NA51 Collaboration at CERN [3] measured the up and down antiquark ratio, though at zero rapidity [4], but in order to improve the statistical accuracy the data were binned to a single value of Bjorken- x . The E866 experiment, on the other hand, has for the first time mapped out the shape of the \bar{d}/\bar{u} ratio over a large range of x , $0.02 < x < 0.345$.

The relatively large asymmetry found in these experiments implies the presence of non-trivial dynamics in the proton sea which does not have a perturbative QCD origin. From the symmetry properties of QCD, we know that one source of non-perturbative quark-antiquark pairs is the pion cloud associated with dynamical chiral symmetry breaking. The novel and unexpected feature of the E866 data is that the \bar{d}/\bar{u} asymmetry peaks at rather small values of x , and drops quite rapidly with increasing x , approximately like $(1-x)^n$ with $n \sim 10$. This behavior had not been anticipated in global parametrizations of data, and is softer than pion cloud models of the nucleon would generally predict. While the former is just an artifact of an extrapolation of a parametric function into an unmeasured region, with no physics implications, the

reason for the latter is the rather hard valence antiquark distribution in the pion, $\bar{q}_v^\pi \sim (1-x)^n$ with $n \sim 1$. Although this x dependence is softened somewhat by the probability of finding the pion in the nucleon in the first place, the resulting convoluted distribution still does not die off as rapidly, at large x , as the new E866 data would suggest.

Our analysis suggests that a quantitative description of the entire region of x covered in the experiment requires a delicate balance between several competing mechanisms, which leads us to speculate about the possibility of a two-phase picture of the non-perturbative sea of the nucleon. At larger x , the dynamics of the pion cloud of the nucleon come to the fore, with deep-inelastic scattering from the πN component of the nucleon wave function providing the bulk of the $\bar{d} - \bar{u}$ asymmetry. On the other hand, the important role of the Δ isobar in nuclear physics has been known for a long time, and it proves to be of some importance in this case too. If a part of the $\pi\Delta$ distribution happens to be harder than the πN , there would be cancellation of some of the \bar{d} excess at large x . Such a piece does indeed arise in the light-cone formulation of the meson-cloud model, provided the $\pi N\Delta$ form factor, parametrized for example by a dipole form, is harder than the πNN form factor [5]—something which is consistent with the measured difference between the nucleon and $N\Delta$ transition axial form factors.

To be consistent with the measured sum, $\bar{u} + \bar{d}$, it is known that both the πNN and $\pi N\Delta$ form factors need to be relatively soft [6]. The best fit, within the pion cloud framework, to the data on both the sum and difference of \bar{u} and \bar{d} at large x accounts for around half of the integrated asymmetry, leaving room for possible other, non-pionic, mechanisms to provide the missing strength at smaller x . (While the usual discussions of the pionic contribution focus on the valence quarks in the pion, there is also some theoretical argument for a $\bar{d} - \bar{u}$ asymmetry in the sea of the pion. We

*Email address: w.melnitchouk@fz-juelich.de

†Email address: j.speth@fz-juelich.de

‡Email address: athomas@physics.adelaide.edu.au

also estimate how this might affect the analysis.)

It should be noted that, aside from the flavor-asymmetric πN and $\pi\Delta$ components of the nucleon, there is no *a priori* reason why the “bare” (non-pion-dressed) nucleon state itself cannot have an intrinsic asymmetric sea associated with it. In fact, this is actually what is expected from the Pauli exclusion principle, as anticipated long ago by Field and Feynman [7] on the simple basis that the u and d valence quark sectors are unequally populated in the proton ground state. Although more difficult to estimate model-independently, the contribution to the $\bar{d}-\bar{u}$ difference from antisymmetrization has been calculated within a non-perturbative model of the nucleon [8]. Along the lines of the model estimates, we find that the effects of antisymmetrization are most relevant at small x , with normalization such that they can account for most of the remaining half of the integrated asymmetry. Indeed, our analysis suggests that the best fit to the E866 data is obtained when both mechanisms play a role, consistent with the conclusions of earlier analyses [9,10] of the NMC data on the proton-neutron structure function difference.

In the following, we first outline in Sec. II the experimental status of the \bar{d}/\bar{u} asymmetry, including a comparison of the Drell-Yan data with deep-inelastic scattering data. In Sec. III the question of the possible origin of the asymmetry is addressed, in the form of the chiral structure of the nucleon and the associated pion cloud, while Sec. IV deals with the role of the Pauli exclusion principle as a source of flavor asymmetry. In Sec. V we summarize our conclusions.

II. THE LIGHT ANTIQUARK ASYMMETRY

The E866/NuSea Collaboration measured $\mu^+\mu^-$ Drell-Yan pairs produced in pp and pD collisions. In the parton model the Drell-Yan cross section is proportional to

$$\sigma^{ph} \propto \sum_q e_q^2 (q^p(x_1)\bar{q}^h(x_2) + \bar{q}^p(x_1)q^h(x_2)), \quad (1)$$

where $h=p$ or D , and x_1 and x_2 are the light-cone momentum fractions carried by partons in the projectile and target hadron, respectively.

Assuming that the deuteron is composed of two bound nucleons, and utilizing isospin symmetry ($u^p=d^n$, etc.), in the limit $x_1 \gg x_2$ [in which $\bar{q}(x_1) \ll q(x_1)$] the ratio of the deuteron to proton cross sections can be written

$$\left. \frac{\sigma^{pD}}{2\sigma^{pp}} \right|_{x_1 \gg x_2} = \frac{1}{2} \left(\frac{\bar{u}(x_2)}{\bar{u}(x_2)} + \frac{\bar{d}(x_2)}{\bar{u}(x_2)} \right) \times \frac{4 + d(x_1)/u(x_1)}{4 + d(x_1)/u(x_1) \cdot \bar{d}(x_2)/\bar{u}(x_2)}, \quad (2)$$

where \bar{q} is the antiquark distribution in the bound proton. Neglecting relativistic and nucleon off-shell effects, this can

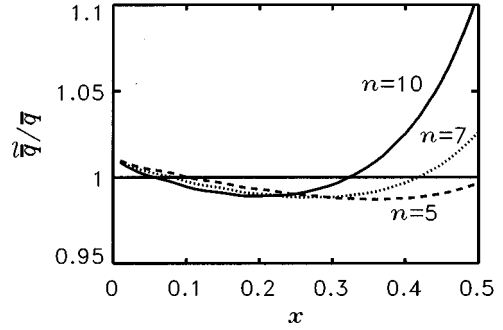


FIG. 1. Ratio of the antiquark distribution in a nucleon bound in the deuteron to that in the free nucleon, for $\bar{q} \sim (1-x)^n$, with $n=5, 7$ and 10 .

be approximated by a convolution of the antiquark distribution in the proton with the proton distribution function in the deuteron [11,12],

$$\bar{q}(x) \approx \int_x^1 \frac{dz}{z} f_{N/D}(z) \bar{q}(x/z), \quad (3)$$

where $f_{N/D}(z)$ is the distribution of nucleons in the deuteron with light-cone momentum fraction z . In practice we use the function $f_{N/D}$ from Ref. [11], where it is given in terms of a realistic deuteron wave function that has been constrained to reproduce the static deuteron properties and nucleon-nucleon phase shifts.

In the absence of nuclear effects, $\bar{q} = \bar{q}_0$, one would have

$$\frac{\sigma^{pD}}{2\sigma^{pp}} = \frac{1}{2} \left(1 + \frac{\bar{d}(x_2)}{\bar{u}(x_2)} \right) \times \frac{4 + d(x_1)/u(x_1)}{4 + d(x_1)/u(x_1) \cdot \bar{d}(x_2)/\bar{u}(x_2)}, \quad (4)$$

so that the ratio would be unity if $\bar{d} = \bar{u}$. On the other hand, we know that nuclear shadowing exists in the deuteron at small x (see [13] and references within), and at large x nuclear binding and Fermi motion effects come into prominence [14]. Since the bulk of the effect is observed outside the very small x region, shadowing will not be important in these data. However, Fermi smearing is potentially more relevant, as what enters in the parentheses in Eq. (2) is the ratio of smeared to unsmeared *sea* quark distributions, for which smearing effects should come into play at much smaller x than for the total structure function [11]. In Fig. 1 we show the ratio of the antiquark distributions in a proton bound in the deuteron to that in a free proton, \bar{q}/\bar{q}_0 . Parametrizing the antiquarks for illustration purposes at large x by a simple $\bar{q} \sim (1-x)^n$ form, with $n=5, 7$ and 10 , the ratio is seen to rise rapidly above $x \sim 0.4$, though in the measured region it only deviates from unity by a few percent. From this one can conclude that Eq. (4) should be a reasonable approximation to Eq. (2).

In the extreme large- x_1 limit, where $d(x_1) \ll u(x_1)$, the cross section ratio would directly give \bar{d}/\bar{u} :

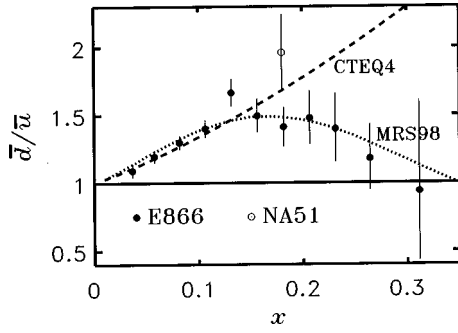


FIG. 2. The x dependence of the \bar{d}/\bar{u} ratio from the E866 [1] (filled circles) and NA51 [3] (open circle) experiments, compared with the CTEQ4 [15] and MRS98 [16] parametrizations. Note that the MRS98 parametrization included the E866 data in their fits, while CTEQ4 predates the experiment.

$$\frac{\sigma^{pD}}{2\sigma^{pp}} \Big|_{d(x_1) \ll u(x_1)} \rightarrow \frac{1}{2} \left(1 + \frac{\bar{d}(x_2)}{\bar{u}(x_2)} \right). \quad (5)$$

For the E866 data the criterion $x_1 \gg x_2$ is not always satisfied, however, so that Eq. (5) gives only an indication of the sensitivity of the Drell-Yan cross section to \bar{d}/\bar{u} , and in practice the full parton model cross section is used together with an iterative procedure in which the valence and total sea distributions are assumed known and the extracted \bar{d}/\bar{u} ratio adjusted to fit the data [1]. The resulting \bar{d}/\bar{u} ratio is shown in Fig. 2, where the asymmetry is found to peak at relatively small x , $x \sim 0.15$, dropping rapidly to unity by $x \sim 0.3$. Also shown for comparison is the NA51 data point [3], extracted from the Drell-Yan pp and pn asymmetry at $x=0.18$, which lies slightly above the E866 data. The two fits are from the CTEQ4 [15] parametrization (dashed line), which has the asymmetry extending out to larger x , and the more recent MRS98 analysis [16] (dotted line), which included the E866 data.

From the \bar{d}/\bar{u} ratio the E866/NuSea Collaboration further extract the difference, $\Delta \equiv \bar{d} - \bar{u}$, assuming the sum $\bar{d} + \bar{u}$ from the CTEQ4 fit,

$$\Delta = \left(\frac{\bar{d}/\bar{u} - 1}{\bar{d}/\bar{u} + 1} \right) (\bar{d} + \bar{u})_{\text{CTEQ4}}. \quad (6)$$

The resulting points are shown in Fig. 3, again in comparison with the CTEQ4 and Martin-Roberts-Stirling 1998 (MRS98) parametrizations. It should be noted, however, that the sea quark distribution in these fits is not very well determined at large x . In particular, a larger total sea at $x \sim 0.2-0.3$ would result in a larger asymmetry Δ in the region just where the E866 data appear to drop rapidly to zero. Further data from the E866 experiment on the total antiquark distribution at large x should help to clarify the issue. The Drell-Yan data

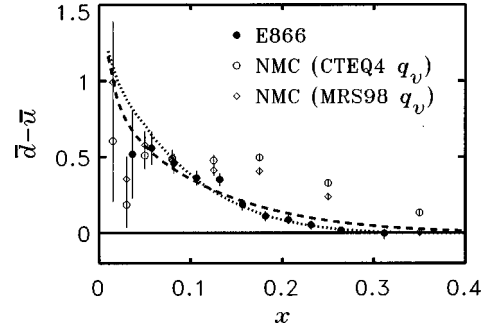


FIG. 3. Comparison of $\bar{d} - \bar{u}$ from the E866 experiment [1] with the values extracted from the NMC measurement of the proton-neutron structure function difference [2], using the CTEQ4 [15] and MRS98 [16] parametrizations for the valence quark distributions. Also shown are the parametrizations of $\bar{d} - \bar{u}$ from CTEQ4 (dashed) and MRS98 (dotted).

can also be compared with the proton-neutron¹ structure function difference measured previously by the New Muon Collaboration [2], if one assumes that the valence quark distributions in the proton are known. In this case the antiquark asymmetry can be written

$$\bar{d}(x) - \bar{u}(x) = \frac{1}{2} (u_v(x) - d_v(x)) - \frac{3}{2x} (F_2^p(x) - F_2^n(x))_{\text{NMC}}. \quad (7)$$

The $\bar{d} - \bar{u}$ difference extracted from the New Muon Collaboration (NMC) data is shown in Fig. 3 using both the CTEQ4 (open circles) and MRS98 (diamonds) fits to the valence quark distributions. The NMC values appear to lie consistently above the E866 data for x above ~ 0.1 , although there is some sensitivity to the choice of valence quark parametrization. Not surprisingly, the E866 integrated value (after extrapolating down to $x=0$ and up to $x=1$) is found to be: $\Delta^{\text{E866}} = \int_0^1 dx (\bar{d} - \bar{u}) = 0.100 \pm 0.018$ [17], somewhat smaller than the NMC value, which is $\Delta^{\text{NMC}} = 0.148 \pm 0.039$ [2], although still consistent within errors. This difference will be further enhanced if one corrects the NMC data for shadowing in the deuteron, omission of which underestimates the violation of the Gottfried sum rule [13].

The observation of a large asymmetry between \bar{u} and \bar{d} , now both at CERN and Fermilab, provides theorists with a challenge to better understand the internal, non-perturbative structure of the nucleon, as the asymmetry due to perturbative effects is known to be very small [18]. In the next section we consider the possible origin of this asymmetry, in the form of the non-perturbative chiral structure of the nucleon.

III. CHIRAL SYMMETRY AND THE MESON CLOUD

The simplest and most obvious source of a non-perturbative asymmetry in the light quark sea is the chiral

¹Note that the NMC extracted the neutron structure function from proton and deuteron data assuming $F_2^n = F_2^D - F_2^p$.

structure of QCD. From numerous studies in low energy physics, including chiral perturbation theory, pions are known to play a crucial role in the structure and dynamics of the nucleon. However, there is no reason why the long-range tail of the nucleon should not also play a role at higher energies. This was first alluded to by Sullivan [19], who argued that deep-inelastic scattering from the pion cloud of the nucleon is a scaling contribution to the nucleon structure function. Indeed, expectations for the ratio of nuclear to nucleon structure functions, based on arguments that nuclear scales are much smaller than typical deep-inelastic scales, and therefore irrelevant, were proved to be dramatically wrong by the observation of the nuclear European Muon Collaboration (EMC) effect [20].

As pointed out by Thomas [6], if the proton's wave function contains an explicit π^+n Fock state component, a deep-inelastic probe scattering from the virtual π^+ , which contains a valence \bar{d} quark, will automatically lead to a \bar{d} excess in the proton. This is the essential physical idea behind these expectations, and has been used to address not only the \bar{d}/\bar{u} asymmetry [5,9,10,21–27], but also SU(3) flavor symmetry breaking in the proton sea [6], as well as asymmetries in the strange [28] and heavier flavor sectors [29,30]. In recent years this picture has been refined and elaborated with inclusion of additional meson and baryon states [5,24], and constraints put on many of the model parameters by comparisons of the predictions of the model with other processes [31,32].

The basic hypothesis of the meson cloud model is that the physical nucleon state can be expanded [in an infinite momentum frame (IMF) and in the one-meson approximation] in a series involving bare nucleon and two-particle, meson-baryon states. The essential ingredients are the meson-baryon distribution functions, $f_{MB}(y)$, which give the probability to find a meson, M , in the nucleon carrying a fraction y of the nucleon's light-cone momentum. As discussed at length in the literature, for these functions to have the correct probabilistic interpretation in the IMF, they must be related to the distributions of baryons in the nucleon, $f_{BM}(y)$, via:

$$f_{MB}(y) = f_{BM}(1-y). \quad (8)$$

This constraint can be verified easily in the IMF, but is not entirely clear in covariant formulations (see Refs. [5,24,27] for further discussion on this point). The IMF treatment has the additional advantage that the meson and baryon are on-mass-shell and so one has no ambiguities associated with the possible off-mass-shell behavior of their structure functions that are encountered in the covariant treatments.

The contribution to the antiquark distribution in the proton, $\delta^{(MB)}\bar{q}$, can then be written in the IMF as a convolution of the meson distribution function and the antiquark distribution in the (on-mass-shell) pion:

$$\delta^{(MB)}\bar{q}(x) = \int_x^1 \frac{dy}{y} f_{MB}(y) \bar{q}^M(x/y). \quad (9)$$

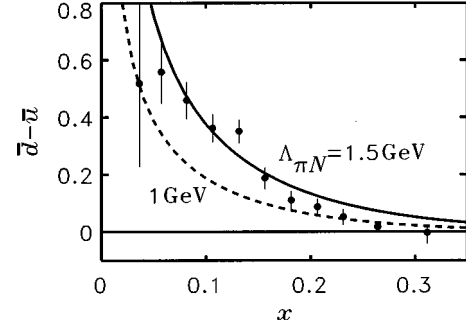


FIG. 4. Calculated $\bar{d}-\bar{u}$ difference arising from the πN component of the proton's wave function, for cut-off masses $\Lambda_{\pi N} = 1$ GeV (dashed) and 1.5 GeV (solid).

Note that this is the leading contribution to the antiquark distribution, and is independent of the model of the bare nucleon states.

In earlier studies it was found, not surprisingly, that pions are the most important mesons, and that the dominant contributions are those associated with the πN component of the proton's wave function. The distribution of pions with a recoiling nucleon is given by [5,27,31]

$$f_{\pi N}(y) = \frac{3g_{\pi NN}^2}{16\pi^2} \int_0^\infty \frac{dk_T^2}{(1-y)} \times \frac{\mathcal{F}_{\pi N}^2(s_{\pi N})}{y(M^2 - s_{\pi N})^2} \left(\frac{k_T^2 + y^2 M^2}{1-y} \right), \quad (10)$$

so that $f_{\pi^+n} = 2f_{\pi^0 p} = (2/3)f_{\pi N}$ for the respective charge states. The invariant mass squared of the πN system is given by $s_{\pi N} = (k_T^2 + m_\pi^2)/y + (k_T^2 + M^2)/(1-y)$, and for the functional form of the πNN vertex form factor $\mathcal{F}_{\pi N}(s_{\pi N})$ we take a simple dipole parametrization:

$$\mathcal{F}_{\pi N}(s_{\pi N}) = \left(\frac{\Lambda_{\pi N}^2 + M^2}{\Lambda_{\pi N}^2 + s_{\pi N}} \right)^2, \quad (11)$$

normalized so that the coupling constant $g_{\pi NN}$ has its standard value ($= 13.07$) at the pole [$\mathcal{F}(M^2) = 1$]. The symmetry relations (8) are automatically guaranteed with this type of form factor, whereas in the earlier covariant formulations [9,10,21–23], with t -dependent form factors, this could not be achieved. Note that the E866 group also utilized the formulation in terms of t -dependent form factors in their recent theoretical analysis [17] of the Drell-Yan data.

The antiquark distribution in the pion has been measured in πN Drell-Yan experiments by the E615 Collaboration at Fermilab [33] and by the NA10 [34] and NA3 [35] Collaborations at CERN. These have been parametrized in next-to-leading order analyses in Refs. [36,37]. Unless stated otherwise, we use the valence part of the pion's antiquark distribution from Ref. [36] throughout this analysis.

Because the meson cloud model is a model of part of the (non-perturbative) sea, it can only be reliably applied to describing the non-singlet $\bar{d}-\bar{u}$ distribution. In Fig. 4 we show the calculated difference arising from the πN component of

the proton's wave function for two different cut-off masses, $\Lambda_{\pi N}=1$ GeV (dashed) and 1.5 GeV (solid), giving average multiplicities $\langle n \rangle_{\pi N} \equiv \int_0^1 dy f_{\pi N}(y) = 13\%$ and 26%, respectively. With the latter one has excellent agreement at intermediate x , $x \lesssim 0.2$, while somewhat overestimating the data at the larger x values. The excess at large x is less severe for the smaller cut-off. However, the strength of the contribution in that case is too small at lower x .

To reconstruct the ratio from the calculated difference, we assume, following E866, that the total $\bar{d} + \bar{u}$ is given by the CTEQ4 parametrization [15], and invert Eq. (6). The resulting ratio is plotted in Fig. 5. At small x the agreement with the data for the larger cut-off is clearly excellent, but at larger x the calculation does not follow the downward trend suggested by the data, similar to the CTEQ4 parametrization in Fig. 2.

Both Figs. 4 and 5 suggest that the excess of \bar{d} over \bar{u} is too strong for $x \gtrsim 0.2$, and that a mechanism which suppresses or cancels this excess could be responsible for the behavior seen in the data. One way to obtain such a suppression would be if the pion structure function were softer. Actually, perturbative QCD suggests that the leading twist part of F_2^π should behave like $(1-x)^2$ [38] at large x (see also Ref. [39]), although the Drell-Yan data [33–35] show that it is closer to $(1-x)$, even with the inclusion of higher twist contributions.

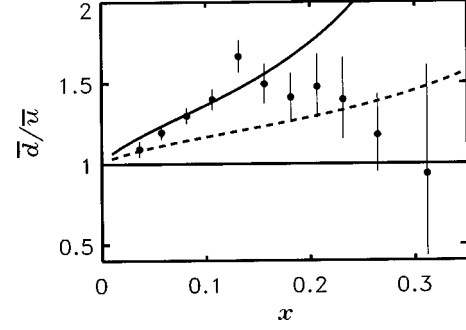


FIG. 5. Extracted \bar{d}/\bar{u} ratio for the πN component, with the curves as in Fig. 4.

Within the pion cloud model another way that some of the cancellation can be understood is through the Δ isobar. Although the $\pi\Delta$ component should be smaller in magnitude than the πN , if some part of the $\pi\Delta$ distribution were to be harder than the πN it would allow for some cancellation of the large- x excess, while preserving more of the asymmetry at smaller x . In fact, qualitatively such behavior is exactly what is seen in the model.

In the IMF, the pion distribution function with a recoil Δ is given by [5,27,31]

$$f_{\pi\Delta}(y) = \frac{2g_{\pi N\Delta}^2}{16\pi^2} \int_0^\infty \frac{dk_T^2}{(1-y)} \frac{\mathcal{F}_{\pi\Delta}^2(s_{\pi\Delta})}{y(s_{\pi\Delta} - M^2)^2} \frac{[k_T^2 + (M_\Delta - (1-y)M)^2][k_T^2 + (M_\Delta + (1-y)M)^2]^2}{6M_\Delta^2(1-y)^3}, \quad (12)$$

where $s_{\pi\Delta}$ is the $\pi\Delta$ invariant mass squared and we take the same functional form [cf. Eq. (11)] for the $\pi N\Delta$ form factor as for πNN . The different $\pi\Delta$ charge states are obtained from Eq. (12) via $f_{\pi^+\Delta^0} = (1/2)f_{\pi^0\Delta^+} = (1/3)f_{\pi^-\Delta^{++}} = (1/6)f_{\pi\Delta}$. The $\pi N\Delta$ coupling constant is defined by:

$$\langle N\pi | H_{int} | \Delta \rangle = g_{\pi N\Delta} C_{(1/2)1(3/2)}^{t_N t_\pi t_\Delta} \bar{u}(p_N, s_N) \times (p_N^\alpha - p_\Delta^\alpha) u_\alpha(p_\Delta, s_\Delta), \quad (13)$$

with u_α the Rarita-Schwinger spinor-vector. The value of $g_{\pi N\Delta}$ can be related to the πNN coupling constant via SU(6) symmetry, $g_{\pi N\Delta} = (6\sqrt{2}/5)f_{\pi NN}/m_\pi \approx 11.8 \text{ GeV}^{-1}$ [with $f_{\pi NN} = (m_\pi/2M)g_{\pi NN}$]. Note that in Ref. [40] (see also [22]) the $\pi N\Delta$ coupling was extracted from the width of the decay $\Delta \rightarrow N\pi$, giving a somewhat larger value $g_{\pi N\Delta} \approx 15.9 \text{ GeV}^{-1}$ compared with the SU(6) value. However, from numerous studies [41] of πN scattering in the Δ resonance region, it is known that up to 50% of the Δ width comes from πN rescattering (for example, through diagrams of the Chew-Low type), so that it would be inappropriate to ascribe the entire width to the tree level process in determining $g_{\pi N\Delta}$. We expect the SU(6) value for the $\pi N\Delta$ coupling to be accurate to within $\sim 10\text{--}20\%$.

In Fig. 6 we show the $\pi\Delta$ distribution function as a function of y , compared with the πN distribution (10). The latter is calculated with a form factor cut-off, $\Lambda_{\pi N}$, of 1 GeV, while the former has $\Lambda_{\pi\Delta} = 1.3$ GeV, which gives $\langle n \rangle_{\pi\Delta} = 11\%$. The $\pi\Delta$ distribution is somewhat broader than the πN , peaking at slightly larger y (0.3 cf. 0.25) and having more strength at larger y , so that relatively more of the πN contribution may be canceled at larger x than at smaller x . (Note that in covariant approaches with a t -dependent dipole form factor the $\pi\Delta$ distribution is softer than the πN , so that there the Δ plays a negligible role at large x .) The cancellation of the \bar{d} excess with the inclusion of $\pi\Delta$ states can be seen explicitly in Fig. 7, where the (positive) πN and (negative) $\pi\Delta$ contributions are shown (dashed lines) for $\Lambda_{\pi N} = 1.5$ GeV and $\Lambda_{\pi\Delta} = 1.3$ GeV, together with the sum (solid). In Fig. 7(a) the $\pi\Delta$ brings the difference $\bar{d} - \bar{u}$ closer to the large- x data points, while allowing for a reasonable fit at smaller x . However, the cancellation is still not sufficient to produce a downturn in the ratio at large x , as the E866 data appear to prefer, Fig. 7(b).

More cancellation can be achieved by either increasing the $\pi\Delta$ contribution, or decreasing the πN contribution. Either is acceptable in the model, as long as the resulting dis-

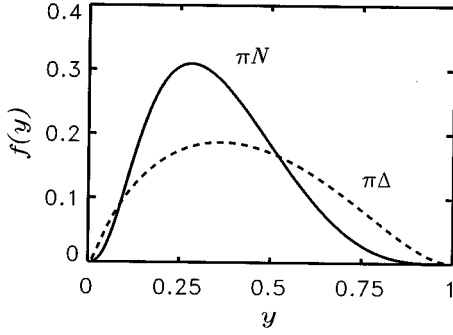


FIG. 6. πN and $\pi\Delta$ momentum distribution functions, with dipole form factor cut-offs $\Lambda_{\pi N}=1$ GeV and $\Lambda_{\pi\Delta}=1.3$ GeV.

tributions do not contradict other observables, such as the total $\bar{d} + \bar{u}$ distribution, which should serve as an absolute upper limit on the strength of the form factor [6]. In Fig. 8(a) we show the contributions to the sum $x(\bar{d} + \bar{u})$ from the πN and $\pi\Delta$ components with $\Lambda_{\pi N}=1.5$ GeV and $\Lambda_{\pi\Delta}=1.3$ GeV, compared with the CTEQ4 [15] and MRS98 [16] parametrizations. While at small x the calculated distributions lie safely below the parametrization (the difference is made up by the perturbatively generated $g \rightarrow q\bar{q}$ antiquark distributions), at large x the pion cloud already saturates the total sea with these cut-offs—although one should add a cautionary note that the antiquark distribution at large x is not determined very precisely. For softer combinations of form factors, namely $\Lambda_{\pi N}=1$ GeV, $\Lambda_{\pi\Delta}=1.3$ GeV and $\Lambda_{\pi N}$

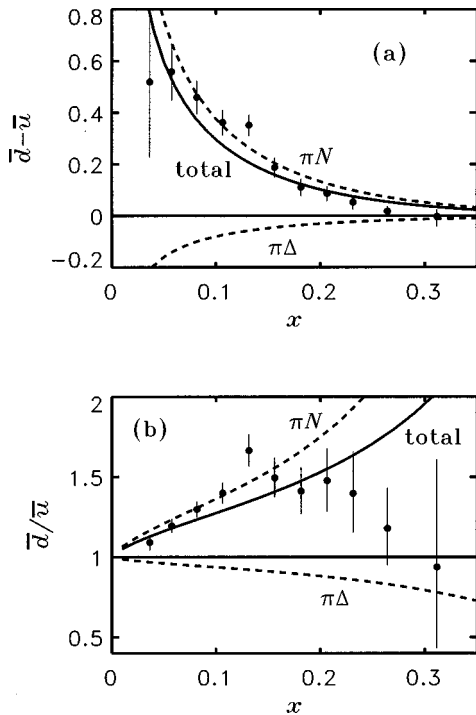


FIG. 7. Contributions from the πN and $\pi\Delta$ components (dashed) and the combined effect (solid) to the (a) $\bar{d} - \bar{u}$ difference and (b) \bar{d}/\bar{u} ratio. The cut-off masses are $\Lambda_{\pi N}=1.5$ GeV and $\Lambda_{\pi\Delta}=1.3$ GeV.

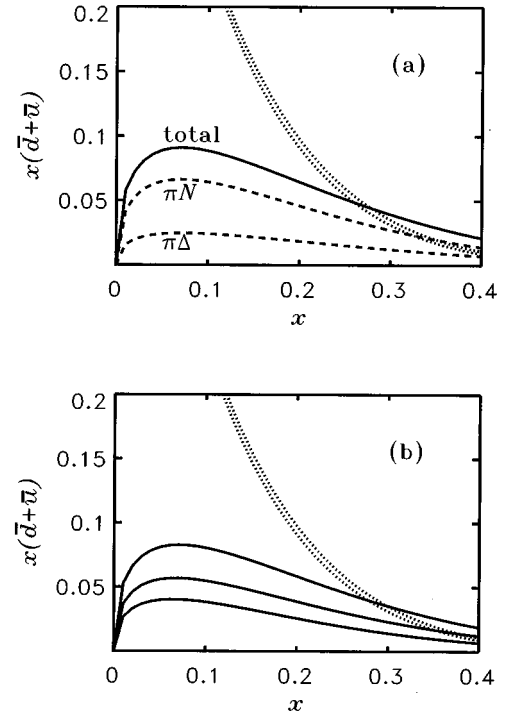


FIG. 8. Total $x(\bar{d} + \bar{u})$ distribution (a) from the πN and $\pi\Delta$ components (dashed), with $\Lambda_{\pi N}=1.5$ GeV, $\Lambda_{\pi\Delta}=1.3$ GeV, and the total (solid), (b) the total contribution for $\Lambda_{\pi N}=1.5$ GeV, $\Lambda_{\pi\Delta}=1.3$ GeV (largest curve), $\Lambda_{\pi N}=1$ GeV, $\Lambda_{\pi\Delta}=1.3$ GeV (middle), and $\Lambda_{\pi N}=\Lambda_{\pi\Delta}=1$ GeV (smallest). The theoretical curves are compared with the CTEQ4 [15] and MRS98 [16] global parametrizations (dotted).

$=\Lambda_{\pi\Delta}=1$ GeV, the total non-perturbative antiquark sea in Fig. 8(b) is below the empirical parametrizations in both cases.

Therefore the only way to obtain a smaller \bar{d} excess at large x and still be consistent with the total antiquark distribution is to reduce the πN component, having a cut-off smaller than for the $\pi N\Delta$ vertex. It was argued in Ref. [17] that the $\pi N\Delta$ form factor should be softer than the πNN , based on the observation that the $M1$ transition form factor was softer for $\gamma N\Delta$ than for γNN . However, there is no clear connection between these form factors, and hence no compelling reason why the $\pi N\Delta$ form factor cannot be harder than that for πNN . Indeed, a comparison of the axial form factors for the nucleon and for the $N-\Delta$ transition strongly favor an $N-\Delta$ axial form factor that is significantly harder than that of the nucleon. In fact, the former is best fit by a 1.3 GeV dipole, while the latter by a 1.02 GeV dipole parametrization [42]. Within the framework of PCAC these form factors are directly related to the corresponding form factors for pion emission or absorption [43].

In Fig. 9 we show the difference and ratio of the \bar{d} and \bar{u} distributions calculated with the softer πNN form factor, $\Lambda_{\pi N}=1$ GeV, and $\Lambda_{\pi\Delta}=1.3$ GeV. The excess at large x now is largely canceled by the $\pi\Delta$. However, the smaller πN contribution means that the asymmetry is underestimated in the intermediate x range, $x \lesssim 0.2$.

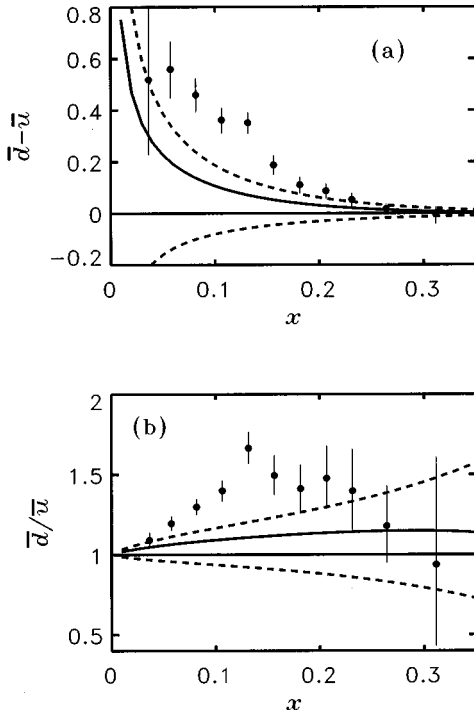


FIG. 9. As in Fig. 7, but for $\Lambda_{\pi N}=1$ GeV, $\Lambda_{\pi\Delta}=1.3$ GeV.

Based on these results it would appear difficult to obtain a quantitative description, within the pion cloud model, of both the ratio *and* the difference of the \bar{d} and \bar{u} distributions, together with the total sea. One needs to consider, therefore, the possibility that other mechanisms may at least be partly responsible for the discrepancy. Additional meson-baryon components, such as the ρN , could be included in extended versions of the meson cloud model [5,24]. The ρN , however, has a harder y distribution than the πN , which would lead to an enhancement of asymmetry at large x , in contradiction with the data, even though the ρ structure function may be softer than the π . On the other hand, the magnitude of the ρN contribution is known from previous analyses [5,24] to be significant only for very hard ρNN form factors. For a ρNN vertex of similar shape to that used for the πNN vertex, $\Lambda_{\rho N}\approx\Lambda_{\pi N}$, the contribution from the ρ is unimportant. Furthermore, the $\rho\Delta$ contribution is far too small to be featured in any subsequent cancellation [5,24,27], if the $\rho\Delta$ form factor is comparable to that for $\pi\Delta$.

Another interesting possibility is that the pion sea itself could be asymmetric. In Eq. (9) only the valence structure of the pion was included, though in principle there could even be asymmetric contributions to $\bar{d}-\bar{u}$ in the proton from asymmetric \bar{d}^{π^+} and \bar{u}^{π^+} distributions in a pion. One obvious source of a pion sea asymmetry involves the same physics that is responsible for the charge radius of the π^+ , namely the dissociation into virtual π and ρ mesons, $\pi^+\rightarrow\pi^+\rho^0$ or $\pi^0\rho^+$. The effects of a meson cloud of a pion ($\pi\rho$ as well as $\bar{K}K^*$ and $K\bar{K}^*$) on deep-inelastic structure functions were previously investigated in Ref. [44]. (Of course the Δ -isobar again cancels some of this asymmetry through the π^- , but as with the valence pion contributions,

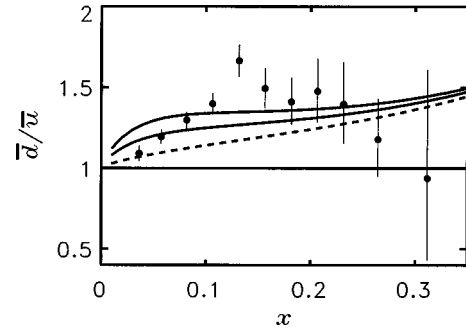


FIG. 10. Effect of an asymmetric pion sea on the \bar{d}/\bar{u} ratio. The dashed curve represents the ratio for a symmetric pion sea with $\Lambda_{\pi N}=\Lambda_{\pi\Delta}=1$ GeV, while the solid curves have asymmetric seas in the ratio $\bar{d}_{\text{sea}}^{\pi^+}:\bar{u}_{\text{sea}}^{\pi^+}=2:1$ (lower curve) and $4:1$ (upper curve).

the sign of the effect remains.)

Unfortunately, nothing is known empirically about the pion sea, so that the shape and normalization of such an asymmetry can at present only be speculative. Rather than construct a detailed model of the pion sea involving additional free parameters (meson-meson vertex functions), at this stage it is more practical to ask how sensitive could the overall \bar{d}/\bar{u} asymmetry be to a possible asymmetric sea, one that is consistent with all the known phenomenological constraints. To address this question we parametrize the non-singlet part of the sea distribution in the pion by a simple form,

$$\bar{d}_{\text{sea}}^{\pi^+}-\bar{u}_{\text{sea}}^{\pi^+}=\mathcal{N}x^\alpha(1-x)^\beta. \quad (14)$$

From low-energy quark models and Regge theory the exponent α is expected to be around 0 and $-1/2$, respectively, while β should be between 5–7 from perturbative QCD arguments and from our knowledge of the nucleon sea quark distribution. The normalization of this component is unknown and given by the parameter \mathcal{N} . Allowing for up to a factor 2 uncertainty in the pion sea (which is related to the uncertainty in the knowledge of the gluon distribution in the pion), we set $\mathcal{N}\sim 4$, and to be definite take $\alpha=0$ and $\beta=5$.

The resulting \bar{d}/\bar{u} ratio is shown in Fig. 10 for a ratio of pion sea distributions $\bar{d}_{\text{sea}}^{\pi^+}:\bar{u}_{\text{sea}}^{\pi^+}=2:1$ (lower solid) and $4:1$ (upper solid), respectively, as well as for a symmetric sea (dashed) for $\Lambda_{\pi N}=\Lambda_{\pi\Delta}=1$ GeV. Clearly one gets appreciable enhancement in the low- and intermediate- x range, bringing the curves to better agreement with the data, even though the ratio is somewhat overestimated at low x for the more asymmetric sea scenario. How reasonable this choice of parameters is can only be ascertained by acquiring data on the pion structure function at values of x smaller than currently available. The phenomenological consequences of an asymmetric pion sea for Drell-Yan πN and other processes will be discussed in more detail elsewhere [45]—see also Ref. [46] for a discussion of measurements which would be sensitive to such an asymmetry.

Going beyond explanations involving meson clouds, one can also investigate the possibility that the bare nucleon itself could be asymmetric with respect to \bar{u} and \bar{d} . As suggested

long ago by Field and Feynman [7], the Pauli exclusion principle can contribute to the asymmetry on the basis of the u and d valence quarks being unequally represented in the proton, thereby affecting the likelihood with which $q\bar{q}$ pairs can be created in different flavor channels. In fact, earlier analyses of the NMC data [9,10] suggested that the best agreement between theory and experiment could be obtained with the combined effects of pions and antisymmetrization, and in the next section we explore this possibility further.

IV. ANTISYMMETRIZATION

Although not directly attributable to the exclusion principle, the perturbative effects of higher-order quark exchange diagrams on the $\bar{d}-\bar{u}$ difference were calculated long ago by Ross and Sachrajda [18]. They found that while they had the correct (positive) sign, their magnitude was insignificantly small, since they only arose at order α_s^2 . The flavor asymmetry of the sea associated with the Pauli principle can therefore only be addressed within non-perturbative approaches to parton distributions, as concluded in [18]. Attempts to calculate the valence distribution of the proton in various quark models began in the mid-70s [47], when it was realized that the relationship to the QCD improved parton model is quite natural at a low scale (below 1 GeV²), where most of the momentum of the nucleon resides on its valence quarks [48]. This observation has been successfully exploited by the Dortmund group [49], for example, in constructing phenomenological, valence dominated, parametrizations in just this region.

Bag model calculations of nucleon structure functions have provided some interesting insights into the non-perturbative parton distributions [8,50]. For any model in which valence quarks are confined by a strong scalar field, the vacuum inside and outside the hadron will be different. From the point of view of an external probe, such as the virtual photon in deep inelastic scattering, the change in the vacuum structure inside the hadron will appear as an intrinsic, non-perturbative sea of $q\bar{q}$ pairs [8]. Because of the Pauli exclusion principle, the presence of two valence u quarks, as opposed to a single valence d quark, in the proton implies an asymmetry in this non-perturbative sea, so that there is a small excess of $d\bar{d}$ pairs over $u\bar{u}$ pairs.

For details of the quantitative calculation of this effect, which is model dependent, we refer to the original papers [8]. It is enough for us that the shape of $\bar{d}-\bar{u}$ was found to be similar to that of the usual sea quark distributions and the normalization, $\int dx(\bar{d}-\bar{u})$, less than 0.25. With this in mind, we parametrize the Pauli contribution by:

$$(\bar{d}-\bar{u})^{\text{Pauli}} = \Delta^{\text{Pauli}}(n+1)(1-x)^n. \quad (15)$$

Because the E866 data implies a softer asymmetry than typical global fits of total sea quark distributions would give, as Figs. 2 and 3 illustrate, phenomenologically the power n should be ≥ 10 rather than the 5–7 that has been common for

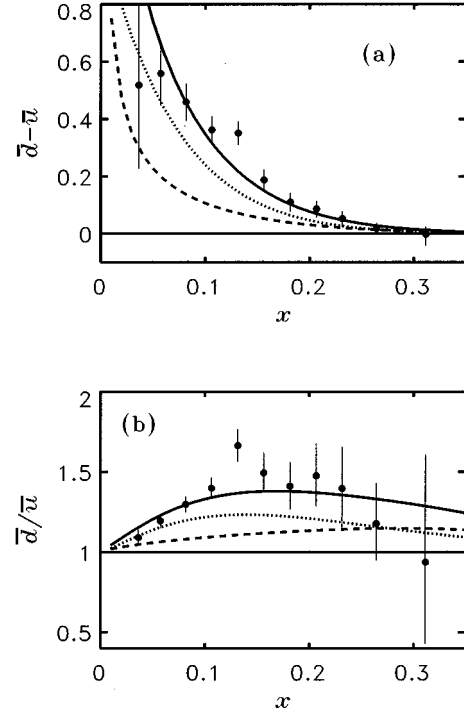


FIG. 11. Contributions from pions with $\Lambda_{\pi N}=1$ GeV and $\Lambda_{\pi\Delta}=1.3$ GeV (dashed) and from antisymmetrization (dotted) to the (a) $\bar{d}-\bar{u}$ difference and (b) \bar{d}/\bar{u} ratio, and the combined effect (solid).

the total \bar{q} fits [15,16]. (Compare also with the original Feynman-Field parametrization [7] which had $n=10$ and 7 for \bar{u} and \bar{d} , respectively.)

The Pauli effect will produce an excess of \bar{d} over \bar{u} over the whole range of x , so that it cannot lead to any cancellation of the large- x asymmetry. To be consistent with the trend of the large- x data, especially for the \bar{d}/\bar{u} ratio, one needs therefore to keep the πNN contribution softer than that from $\pi N\Delta$. Taking the πN and $\pi\Delta$ contributions calculated with $\Lambda_{\pi N}=1$ GeV and $\Lambda_{\pi\Delta}=1.3$ GeV as in Fig. 9 above, we show in Fig. 11 the combined effects of pions and the Pauli effect. For the latter the exponent $n=14$, and the normalization is $\Delta^{\text{Pauli}} \approx 7\%$, which is at the lower end of the expected scale but consistent with the bag model calculations [8]. Together with the integrated asymmetry from pions, $\Delta^{\pi} \approx 0.05$, the combined value $\Delta = \Delta^{\pi} + \Delta^{\text{Pauli}} \approx 0.12$ is in quite reasonable agreement with the experimental result, 0.100 from E866 and 0.148 from NMC. While the quality of the fit in Fig. 11 is quite good, it would be further improved (see Fig. 12) if one were to use the softer pion structure function, $\bar{q}^{\pi} \sim (1-x)^2$, as suggested by perturbative QCD [38].

Before leaving this discussion of antisymmetrization we should also mention the calculations of Donoghue and Golowich [51], and more recently Steffens and Thomas [52], of the one-gluon-exchange corrections to the 3-quark proton wave function. Considering all possible permutations of the 5-quark wave function allowed by Fermi statistics, the antisymmetrization of the $q\bar{q}$ pair split off from the emitted gluon with the quarks in the nucleon ground state, it turns

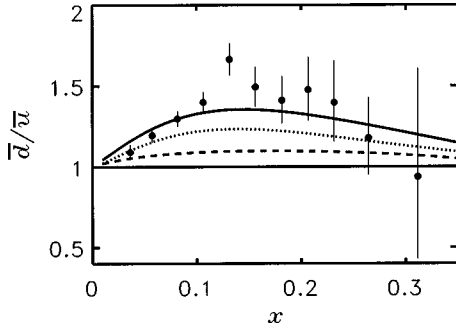


FIG. 12. As in Fig. 11(b), but with an extra power of $(1-x)$ in the pion structure function, according to Ref. [38].

out that there are in fact more diagrams for u quarks than d quarks. This peculiarity results in there actually being an excess of \bar{u} quarks over \bar{d} , albeit a very small one. On the other hand, one should note that this calculation considered just the perturbative contribution, while the Signal-Thomas effect [8] is a totally non-perturbative phenomenon, including all possible non-perturbative interactions between the produced quark (or antiquark) and the confining mean field of the proton. Steffens and Thomas also investigated the effects of antisymmetrization between $q\bar{q}$ pairs arising from one-pion loops with the three quarks in the nucleon ground state [52], although here again the effects were found to be quite small compared with the antisymmetrization for the bare nucleon state, and from the pion cloud contribution discussed in Sec. III.

V. CONCLUSION

We have, for the first time, at our disposal important new data which map out the x -dependence of the asymmetry of the light antiquark sea. Most importantly, the E866 Drell-Yan results confirm the earlier observations that the \bar{d} and \bar{u} content of the proton is not symmetric. One of the more interesting new features of the data is the relatively fast downturn in the \bar{d}/\bar{u} ratio beyond $x \sim 0.15$, which drops rapidly back to unity by $x \sim 0.3$. Taken at face value, this would

appear to provide a challenge to models in which the asymmetry is assumed to arise solely from the pion cloud of the nucleon, and in turn leads us to consider a richer and more complex structure of the non-perturbative sea in which several mechanisms may give competing contributions.

The evidence from the large- x data indicates that a $\pi\Delta$ component in the nucleon wave function may be necessary, one which is harder in momentum space than the πN component. Such a distribution arises naturally in the infinite momentum frame formulation of the pion cloud, unlike in earlier covariant approaches using t -dependent form factors where it was softer than the πN component and hence played no role at large x . Consistency with data for the sum of \bar{d} and \bar{u} at $x \geq 0.2$ requires that both the πNN and $\pi N\Delta$ form factors be relatively soft, making it difficult to avoid underestimating the E866 asymmetry at intermediate x , and leaving room for other effects, such as the Pauli exclusion principle, to make up the difference. Along the lines of previous estimates of the Pauli effect, we find the contribution to the $\bar{d} - \bar{u}$ difference from antisymmetrization to be significant in magnitude, and particularly important at small x . Our final results suggest that the best description of the E866 data is indeed that in which pions and antisymmetrization play roughly equal roles—consistent with the findings of the earlier analysis [10] of the NMC data for $F_2^p - F_2^n$.

In conclusion, we note that it would be helpful to have more data at large x , where the error bars are largest, to verify the downward trend of $\bar{d} - \bar{u}$, and to further explore the possible discrepancy between the Fermilab and CERN data.

ACKNOWLEDGMENTS

W.M. and J.S. would like to thank the Special Research Center for the Subatomic Structure of Matter at the University of Adelaide for hospitality and support. We would also like to thank S. J. Brodsky, G. T. Garvey, E. M. Henley, N. N. Nikolaev, J. Pumplin and F. M. Steffens for helpful discussions. This work was supported by the Australian Research Council.

- [1] E866/NuSea Collaboration, E. A. Hawker *et al.*, Phys. Rev. Lett. **80**, 3715 (1998).
- [2] P. Amaudruz *et al.*, Phys. Rev. Lett. **66**, 2712 (1991); M. Arneodo *et al.*, Phys. Rev. D **50**, R1 (1994); M. Arneodo *et al.*, Phys. Lett. B **364**, 107 (1995).
- [3] A. Baldit *et al.*, Phys. Lett. B **332**, 244 (1994).
- [4] S. D. Ellis and W. J. Stirling, Phys. Lett. B **256**, 258 (1991).
- [5] A. W. Thomas and W. Melnitchouk, in *Proceedings of the JSPS-INS Spring School*, Shimoda, Japan, edited by S. Homma *et al.* (World Scientific, Singapore, 1993); W. Melnitchouk and A. W. Thomas, Phys. Rev. D **47**, 3794 (1993).
- [6] A. W. Thomas, Phys. Lett. **126B**, 97 (1983).
- [7] R. D. Field and R. P. Feynman, Phys. Rev. D **15**, 2590 (1977).
- [8] A. I. Signal and A. W. Thomas, Phys. Lett. B **211**, 481 (1988);

- A. I. Signal and A. W. Thomas, Phys. Rev. D **40**, 2832 (1989); A. W. Schreiber, A. I. Signal, and A. W. Thomas, *ibid.* **44**, 2653 (1991).
- [9] A. I. Signal, A. W. Schreiber, and A. W. Thomas, Mod. Phys. Lett. A **6**, 271 (1991).
- [10] W. Melnitchouk, A. W. Thomas, and A. I. Signal, Z. Phys. A **340**, 85 (1991).
- [11] W. Melnitchouk, A. W. Schreiber, and A. W. Thomas, Phys. Lett. B **335**, 11 (1994).
- [12] A. Bodek and J. L. Ritchie, Phys. Rev. D **23**, 1070 (1981); G. V. Dunne and A. W. Thomas, Nucl. Phys. **A455**, 701 (1986); L. P. Kaptari and A. Yu. Umnikov, Phys. Lett. B **259**, 155 (1991); M. A. Braun and M. V. Tokarev, *ibid.* **320**, 381 (1994).

- [13] W. Melnitchouk and A. W. Thomas, Phys. Rev. D **47**, 3783 (1993).
- [14] W. Melnitchouk and A. W. Thomas, Phys. Lett. B **377**, 11 (1996).
- [15] H. L. Lai, J. Huston, S. Kuhlmann, F. Olness, J. Owens, D. Soper, W. K. Tung, and H. Weerts, Phys. Rev. D **55**, 1280 (1997).
- [16] A. D. Martin, R. G. Roberts, W. J. Stirling, and R. S. Thorne, Eur. Phys. J. C **4**, 463 (1998).
- [17] E866/NuSea Collaboration, J. C. Peng *et al.*, Phys. Rev. D **58**, 092004 (1998).
- [18] D. A. Ross and C. T. Sachrajda, Nucl. Phys. **B149**, 497 (1979).
- [19] J. D. Sullivan, Phys. Rev. D **5**, 1732 (1972).
- [20] J. J. Aubert *et al.*, Phys. Lett. **123B**, 275 (1983).
- [21] E. M. Henley and G. A. Miller, Phys. Lett. B **251**, 453 (1990).
- [22] S. Kumano, Phys. Rev. D **43**, 3067 (1991); S. Kumano and J. T. Londergan, *ibid.* **44**, 717 (1991).
- [23] W.-Y. P. Hwang, J. Speth, and G. E. Brown, Z. Phys. A **339**, 383 (1991).
- [24] H. Holtmann, A. Szczurek, and J. Speth, Nucl. Phys. **A569**, 631 (1996).
- [25] E. J. Eichten, I. Hinchliffe, and C. Quigg, Phys. Rev. D **47**, R747 (1993).
- [26] W. Koepf, L. L. Frankfurt, and M. I. Strikman, Phys. Rev. D **53**, 2586 (1996).
- [27] J. Speth and A. W. Thomas, Adv. Nucl. Phys. **24**, 83 (1998).
- [28] W. Melnitchouk and M. Malheiro, Phys. Rev. C **55**, 431 (1997); M. Malheiro and W. Melnitchouk, *ibid.* **56**, 2373 (1997).
- [29] W. Melnitchouk and A. W. Thomas, Phys. Lett. B **414**, 134 (1997).
- [30] S. Paiva, M. Nielsen, F. S. Navarra, F. O. Duraes, and L. L. Barz, Sao Paulo U. Report No. IFUSP-P-1240, hep-ph/9610310.
- [31] V. R. Zoller, Z. Phys. C **54**, 425 (1992); **60**, 141 (1993).
- [32] N. N. Nikolaev, J. Speth, and B. G. Zakharov, Jülich Report No. KFA-IKP-TH-1997-17, hep-ph/9708290.
- [33] J. S. Conway *et al.*, Phys. Rev. D **39**, 92 (1989).
- [34] B. Betev *et al.*, Z. Phys. C **28**, 15 (1985).
- [35] J. Badier *et al.*, Z. Phys. C **18**, 281 (1983).
- [36] M. Gluck, E. Reya, and M. Stratmann, Eur. Phys. J. C **2**, 159 (1998).
- [37] P. J. Sutton, A. D. Martin, R. G. Roberts, and W. J. Stirling, Phys. Rev. D **45**, 2349 (1992).
- [38] G. R. Farrar and D. R. Jackson, Phys. Rev. Lett. **35**, 1416 (1975).
- [39] S. J. Brodsky, M. Burkardt, and I. Schmidt, Nucl. Phys. **B441**, 197 (1995).
- [40] J. Pumplin, Phys. Rev. D **8**, 2249 (1973).
- [41] S. Théberge, G. A. Miller, and A. W. Thomas, Can. J. Phys. **60**, 59 (1982); I. R. Afnan and B. C. Pearce, Phys. Rev. C **35**, 737 (1987); C. Schütz, J. W. Durso, K. Holinde, and J. Speth, *ibid.* **49**, 2671 (1994); C. Schütz, J. Haidenbauer, J. Speth, and J. W. Durso, *ibid.* **57**, 1464 (1998).
- [42] G. T. Jones *et al.*, Z. Phys. C **43**, 527 (1989); T. Kitagaki *et al.*, Phys. Rev. D **42**, 1331 (1990).
- [43] K. Holinde and A. W. Thomas, Phys. Rev. C **42**, 1195 (1990); J. Haidenbauer, K. Holinde, and A. W. Thomas, in *Proceedings of the 14th International Conference on Few Body Problems in Physics*, Williamsburg, Virginia, 1994, edited by Franz Gross (AIP, New York, 1995), p. 490.
- [44] W.-Y. P. Hwang and J. Speth, Phys. Rev. D **45**, 3061 (1992).
- [45] W. Melnitchouk, J. Speth, and A. W. Thomas (in preparation).
- [46] J. T. Londergan *et al.*, Phys. Lett. B **361**, 110 (1995).
- [47] A. Le Yaouanc *et al.*, Phys. Rev. D **11**, 680 (1975); G. Parisi and R. Petronzio, Phys. Lett. **62B**, 331 (1976); R. L. Jaffe, Phys. Rev. D **11**, 1953 (1975); J. Bell, Phys. Lett. **74B**, 77 (1978).
- [48] A. W. Thomas, Prog. Part. Nucl. Phys. **20**, 21 (1988).
- [49] M. Gluck, E. Reya, and A. Vogt, Z. Phys. C **48**, 471 (1990).
- [50] A. W. Schreiber, P. J. Mulders, A. I. Signal, and A. W. Thomas, Phys. Rev. D **45**, 3069 (1992); F. M. Steffens, H. Holtmann, and A. W. Thomas, Phys. Lett. B **358**, 139 (1995).
- [51] J. F. Donoghue and E. Golowich, Phys. Rev. D **15**, 3421 (1977).
- [52] F. M. Steffens and A. W. Thomas, Phys. Rev. C **55**, 900 (1997).