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Pion cloud of the nucleon and its effect on deep-inelastic structure

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In this paper we present the calculation of deep-inelastic scattering structure functions of nucleons dressed by pions. The calculation is performed within the convolution model. We present analytic results for quark and antiquark distributions of a given flavor and spin and numerical results on those structure functions that are likely to be most affected by the dressing of the nucleon. We allow for the probe scattering off both nucleons and Δ's. Diagrams where it scatters off the pions themselves are not taken into account. For the numerical results we look only at combinations of structure functions where these do not contribute.

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I. INTRODUCTION

There are two different reasons why the measurement of deep-inelastic scattering (DIS) of leptons on single nucleons is of fundamental importance. On the one hand it is possible, by comparing structure functions at various Q^2, to test perturbative quantum chromodynamics (PQCD) in a model independent manner. The success of these tests so far is a major reason why QCD has generally come to be accepted as the correct theory of the strong interactions. On the other hand, DIS also provides us, once the PQCD corrections have been incorporated, with important information on the substructure of the nucleons themselves. Indeed, the first DIS experiments [1–3] predate QCD and established, through the observation of scaling, that the nucleon consists of pointlike entities, the partons. It was shown that the quantity $x = Q^2/2p\cdot q$, restricted to lie between 0 and 1, is a measure of the fraction of the target's light-cone momentum carried by the struck quark [$q(p)$ is the four-momentum of the lepton (target); $Q^2 = -q^2$]. The measured cross sections are directly proportional to the probabilities $q_a(x)$ that a quark of type a carries momentum fraction x.

Even though the DIS process is one of high momentum transfer, these probabilities are essentially given by low energy properties of the nucleon. The reason for this is that the mathematical description of DIS factorizes into perturbative (short distance, high energy) and non-perturbative (long distance, low energy) ingredients. If one has information on the relevant low energy matrix elements this may be used to make predictions for DIS as the perturbative part is fully calculable. A solution to QCD in the non-perturbative region, however, still eludes us, so in order to calculate the low energy matrix elements we are forced to use models that incorporate some of QCD's key features, such as confinement. The matrix elements are model dependent and comparison with DIS data can therefore act as a constraint on some of the parameters involved.

One of the simplest models of the nucleon in terms of relativistic massless quarks is the MIT bag model [4, 5]. In this the nucleon consists of three massless relativistic quarks, carrying its quantum numbers such as its spin and flavor, confined within a finite volume by a "bag" surface. Deep-inelastic scattering within this model has been calculated previously (see [6] and references contained therein). It is, however, clear from numerous experimental results that this simple model of nucleon structure is not really adequate. Let us briefly mention two results from DIS experiments in particular. The first is the cause of the so-called proton-spin crisis [7, 8]. The European Muon Collaboration (EMC) results indicate that the major part of the proton spin does not seem to be carried by (current) quarks and that there is a large strange-quark component in the proton. This is in direct contradiction to the expectations from the MIT bag model. The second is the measurement of the isospin distribution within the nucleon by the New Muon Collab-

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oration (NMC) group [9]. The (sea) antiquarks seem to carry net isospin—also at odds with simple expectations, as the perturbatively generated QCD sea is an isospin singlet.

Experimental results such as those mentioned above provide an indication that the nucleon wave function is not as simple as one might have thought. That important physics is missing in the MIT bag model is of course not a surprise—it is well known that it violates PCAC (partial conservation of axial-vector current), and refinements that incorporate this were introduced some time ago [10–15]. It is the purpose of this paper to examine the effect of these refinements on DIS. In particular, we shall examine the consequences of a pionic component in the nucleon wave function on DIS. The effect of this on both the spin and the Gottfried sum rules has been discussed previously [16, 17]. It was found that the former was only weakly influenced by this extension of the wave function. Around 10% of the strength of the proton structure function is shifted to that of the neutron. This is about a factor of 4 too small if the experimental value of the EMC group [7, 8] is accepted at face value. On the other hand, the experimentally observed decrease in the Gottfried sum rule [9] is of about the same order as that expected from the modification of the nucleon wave function by a pion cloud.

The results of [16] and [17] are for the sum rules only. In this paper we shall extend this work to the actual distributions themselves. We shall provide the formalism to do this in Sec. II, followed by some numerical results in Sec. III.

The calculation of Sec. II will be done in the convolution model approximation. Before we proceed we want to give a justification for its use. It is difficult to estimate the validity of the convolution model as it is not possible to calculate corrections in any sensible manner. If the objects interacting with the probe are pions there are even some heuristic arguments [18] why the convolution model should not be a good approximation (final state interactions being important in this case). Given a typical final state interaction time $\tau_{\rm int}$ of the order of the inverse of a few hundred MeV (a typical nuclear binding energy), then the convolution model might be reasonable if the typical time scale for the emittance and reabsorption of the target τ is much less than this. The latter is of the order of $\tau \approx 1/p^{+}_{\rm target} \approx 1/M_{\rm target}$. Hence the heavier the struck object, the more likely it is that the convolution model will be a good approximation. For the pion $\tau_{\rm int} \approx \tau$, so in this case there would seem to be little justification to neglect final state interactions from a physical point of view.

Within the convolution model there are several structure functions, or combinations of structure functions, where the contribution from the pions [19] cancel—for example spin structure functions (a pion, being a spinless object, contains equal numbers of quarks of a given flavor spinning up or down), $F_3(x)$ (a π^+ contains equal numbers of u and $\overline{d}$ quarks, etc.), and $F_2^p(x) - F_2^n(x)$ (for the same reason). More explicitly, let us define the distribution of quarks in the dressed proton in terms of the functions $f_{q/h}^{hH}(x)$, where the proton consists of hadrons h and H and the probe interacts with the hadron h. We have

$$\begin{aligned} u(x) &= 2f_{q/N}^{N}(x) + \tfrac{4}{3}f_{q/N}^{N\pi}(x) + \tfrac{7}{3}f_{q/\Delta}^{\Delta\pi}(x) \\ &\quad + \tfrac{5}{6}f_{q/\pi}^{N\pi}(x) + \tfrac{1}{3}f_{q/\pi}^{\Delta\pi}(x), \\ d(x) &= f_{q/N}^{N}(x) + \tfrac{5}{3}f_{q/N}^{N\pi}(x) + \tfrac{2}{3}f_{q/\Delta}^{\Delta\pi}(x) \\ &\quad + \tfrac{1}{6}f_{q/\pi}^{N\pi}(x) + \tfrac{2}{3}f_{q/\pi}^{\Delta\pi}(x), \\ \overline{u}(x) &= \tfrac{1}{6}f_{q/\pi}^{N\pi}(x) + \tfrac{2}{3}f_{q/\pi}^{\Delta\pi}(x), \\ \overline{d}(x) &= \tfrac{5}{6}f_{q/\pi}^{N\pi}(x) + \tfrac{1}{3}f_{q/\pi}^{\Delta\pi}(x). \end{aligned} \tag{1}$$

We have explicitly written the terms where the probe interacts with the pionic and baryonic component of the dressed proton wave function. We have at this stage assumed, for the sake of simplicity, that the distribution of u quarks in the bare proton is the same as that of the d quarks. We shall not impose this condition in the rest of the paper—it will not change the argument. Furthermore, Eq. (1) is a leading order expression; higher order contributions are not taken into consideration in this paper, and so the bare distributions do not themselves have a pionic component. Finally, Eq. (1) does of course not include perturbative QCD corrections—at this stage we are discussing sum rules intrinsic to the model. In order to compare with data we will need to take these into account—we shall do so in Sec. III. The integrals of the distributions (which we denote by capital F's) are normalized—for example, $F_{q/\Delta}^{\Delta\pi}$ is the probability to find a Δ in the dressed nucleon, etc. In the convolution model this probability is independent of which one of the two component objects is struck by the probe, i.e.,

$$F_{q/\Delta}^{\Delta\pi} = F_{q/\pi}^{\Delta\pi} \equiv \mathrm{Prob}(\Delta\pi), \quad \text{etc.} \tag{2}$$

The distributions $f(x)$ themselves of course do not in general satisfy this equality.

From Eq. (1) we may read off the result for the difference $F_2^p(x) - F_2^n(x)$:

$$\begin{aligned} \frac{F_2^p(x) - F_2^n(x)}{x} &= \tfrac{1}{3}[u(x) - d(x) + \overline{u}(x) - \overline{d}(x)] \\ &= \tfrac{1}{3}[f_{q/N}^{N}(x) - \tfrac{1}{3}f_{q/N}^{N\pi}(x) + \tfrac{5}{3}f_{q/\Delta}^{\Delta\pi}(x)]. \end{aligned} \tag{3}$$

As discussed, the contribution from the pions cancels for all x. The Gottfried sum is given by the integral of Eq. (3):

$$\int_0^1 dx \frac{F_2^p(x) - F_2^n(x)}{x} = \tfrac{1}{3}(F_{q/N}^{N} - \tfrac{1}{3}F_{q/N}^{N\pi} + \tfrac{5}{3}F_{q/\Delta}^{\Delta\pi}). \tag{4}$$

Using equalities such as those in Eq. (2) this is equal to

$$\int_0^1 dx \frac{F_2^p(x) - F_2^n(x)}{x} = \tfrac{1}{3}(F_{q/N}^{N} - \tfrac{1}{3}F_{q/\pi}^{N\pi} + \tfrac{5}{3}F_{q/\pi}^{\Delta\pi}). \tag{5}$$

Note that in Ref. [17] the Gottfried sum has been calculated using Eq. (4), while in [20–22] it has been calculated according to Eq. (5). As demonstrated, this makes no difference to the integral. It would, however, be wrong to use the unintegrated version of Eq. (5) to describe the x distribution as seems to have been done in Ref. [23]—the decrease in the Gottfried sum from $\frac{1}{3}$ corresponds to a decrease in the "valence" contribution of the dressed proton, not a pionic "sea" with unequal $\overline{u}$ and $\overline{d}$ components. In order to calculate $F_2^p(x)-F_2^n(x)$ [and $g_1^n(x)$] we shall therefore restrict ourselves to calculating the contribution of the virtual probe interacting with the baryon, with possibly a spectator pion being present.

It should be pointed out that we would naturally not expect to be able to saturate momentum and various number sum rules by neglecting the contribution from the pions—after all, they will carry a part of the nucleon's momentum and a measurement of, for example, their net u-quark content would indeed be nonzero. This will not concern us here, but should be kept in mind when applying the results presented in this paper to a calculation of, for example, $F_2^p(x)$. Furthermore, as has already been pointed out in Ref. [24], the convolution model does not automatically satisfy sum rules. This is investigated in some detail in Ref. [25]. In our case this effect is numerically rather small.

II. STRUCTURE FUNCTIONS IN THE CLOUDY BAG MODEL

We now turn towards the calculation of quark distributions within a model which includes a pionic component in the wave function. In order to be definite we concentrate on the quark distributions inside a dressed proton with four-momentum p (we take it to be at rest) and with spin pointing in the positive z direction (see Fig. 1). The quark distribution of flavor f and helicity positive ($\uparrow$) or negative ($\downarrow$) is given by

$$\tilde{q}_f^{\uparrow\downarrow}(x) = \frac{p^+}{2\pi}\int_{-\infty}^{\infty} d\xi^- e^{-ixp^+\xi^-}\langle\tilde{p}|\Psi_{+,f}^{\uparrow\downarrow\,\dagger}(\xi^-)\Psi_{+,f}^{\uparrow\downarrow}(0)|\tilde{p}\rangle_c$$
$$\equiv \langle\tilde{p}|\mathcal{O}_f^{\uparrow\downarrow}(x)|\tilde{p}\rangle. \tag{6}$$

The projections on the field operator are defined by

$$\Psi_{+,f}^{\uparrow\downarrow}(\xi) = \frac{1}{2}\gamma^-\gamma^+\frac{1\pm\gamma^5}{2}\Psi_f(\xi) \tag{7}$$

and the tilde on the q and in the bra and ket distinguishes

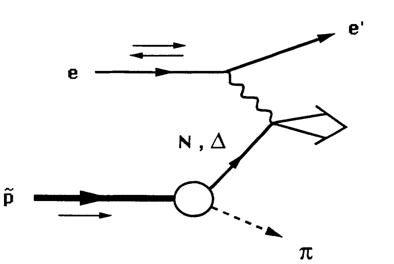

FIG. 1. The convolution diagram for the dressed proton in the one-pion approximation. The arrows indicate the relative directions of the spin of the proton and electron.

quantities dressed by pions from bare quantities $|p\rangle$. The subscript c signifies a connected matrix element and it is understood that it is to be evaluated at the spatial coordinates $\xi^+ = \xi^\perp = 0$.

A. The dressed proton wave function

Let us review the Hamiltonian formalism of the cloudy bag model which is necessary to derive expressions for the dressed proton wave functions. Further details may be found in [13–15, 26]. We shall only work in the space of nonstrange baryons, in which case the full Hamiltonian may be written as

$$H = H_0 + H_I \quad . \tag{8}$$

The kinetic part is

$$H_0 = M_{0N}N^\dagger N + M_{0\Delta}\Delta^\dagger\Delta + \int d\mathbf{k}\,\omega_{0k}\,\mathbf{a_k}^\dagger\mathbf{a_k} \quad . \tag{9}$$

Here N, Δ ($N^\dagger$, $\Delta^\dagger$) are bare baryon annihilation (creation) operators, M_{0N} and $M_{0\Delta}$ are the bare baryon masses, and $\mathbf{a_k}$ ($\mathbf{a_k}^\dagger$) annihilates (creates) a pion with momentum $\mathbf{k}$ and energy ω_{0k} (we shall neglect the kinetic energy of the baryons with respect to their masses).

The interacting part of the Hamiltonian is

$$H_I = \int d\mathbf{k}\,(\mathbf{V_{0k}}\cdot\mathbf{a_k} + \mathbf{V}_{0\mathbf{k}}^\dagger\cdot\mathbf{a_k}^\dagger) \quad , \tag{10}$$

where

$$\mathbf{V_{0k}} = \sum_{\alpha\beta}\alpha^\dagger\mathbf{v}_{0\mathbf{k}}^{\alpha\beta}\,\beta, \quad \alpha,\ \beta\in\{N,\Delta\} \tag{11}$$

and

$$\mathbf{v}_{0\mathbf{k}}^{\alpha\beta} = \frac{if_0^{\alpha\beta}}{m_\pi}\frac{\sqrt{4\pi}u(kR)}{\sqrt{2\omega_{0k}}\sqrt{(2\pi)^3}}(\mathbf{S}^{\alpha\beta}\cdot\mathbf{k})\mathbf{T}^{\alpha\beta}, \tag{12}$$

$$\mathbf{S}^{\alpha\beta} = \sum_m C(S_\beta 1 \to S_\alpha|s_\beta m \to s_\alpha)\hat{\mathbf{s}}_m^*\cdot\sqrt{3}, \tag{13}$$

$$\mathbf{T}^{\alpha\beta} = \sum_n C(T_\beta 1 \to T_\alpha|t_\beta n \to t_\alpha)\hat{\mathbf{t}}_n^*\cdot\sqrt{3}, \tag{14}$$

$$f_0^{NN} = f_0^{\Delta\Delta} = \frac{5}{4\sqrt{2}}f_0^{N\Delta} = \frac{5}{2\sqrt{2}}f_0^{\Delta N} \quad . \tag{15}$$

The quantities $\hat{\mathbf{s}}_m^*$ and $\hat{\mathbf{t}}_m^*$ are spherical spin and isospin vectors, respectively, the C's are Clebsch-Gordan coefficients, $u(kR)$ is the form factor for the cloudy bag [$u(kR)=j_0(kR)+j_2(kR)$], and Ω/R is the lowest energy eigenvalue of the MIT bag. The $f_0^{\alpha\beta}$ are the unrenormalized baryon pion coupling constants. For the physics behind expressions (9) to (15) we refer the reader to [14].

We define the bare and dressed nucleon states through

$$H_0|N_{H_0}\rangle = M_{0N}|N_{H_0}\rangle \tag{16}$$

and

$$H|\tilde{N}\rangle = M|\tilde{N}\rangle \quad , \tag{17}$$

respectively, M being the dressed nucleon's mass. Similar expressions hold for the Δ.

Our aim is to express all the relevant physics of the cloudy bag model in the wave function $|\tilde{N}\rangle$ of the nucleon in terms of the bare baryon states $|N_{H_0}\rangle$. To this end we shall define

$$|\tilde{N}\rangle = |\tilde{N}_0\rangle + |\tilde{N}_I\rangle \ , \tag{18}$$

with

$$|\tilde{N}_0\rangle = (|N_{H_0}\rangle\langle N_{H_0}|)\ |\tilde{N}\rangle \equiv \sqrt{Z_2^N}|N_{H_0}\rangle \ . \tag{19}$$

Z_2^N is the probability that the nucleon is not accompanied by pions. The part of the wave function containing pions is then given by

$$|\tilde{N}_I\rangle = (1 - |N_{H_0}\rangle\langle N_{H_0}|)|\tilde{N}\rangle \equiv \Lambda|\tilde{N}\rangle \ . \tag{20}$$

Using Eqs. (16) and (17) we obtain

$$H_0(|\tilde{N}_0\rangle + |\tilde{N}_I\rangle) + H_I|\tilde{N}\rangle = M(|\tilde{N}_0\rangle + |\tilde{N}_I\rangle); \tag{21}$$

i.e.,

$$(M - H_0)|\tilde{N}_I\rangle = (H_0 - M)|\tilde{N}_0\rangle + H_I|\tilde{N}\rangle \ . \tag{22}$$

Noting that $|\tilde{N}_I\rangle$ is an eigenstate of Λ while $|\tilde{N}_0\rangle$ is annihilated by it, and that Λ commutes with H_0, one then only needs to operate with Λ on both sides of Eq. (22) to obtain

$$|\tilde{N}\rangle = \sqrt{Z_2^N}|N_{H_0}\rangle + \frac{1}{M - H_0}\Lambda H_I|\tilde{N}\rangle \ ; \tag{23}$$

i.e.,

$$|\tilde{N}\rangle = \sqrt{Z_2^N}\ \left(1 - \frac{1}{M - H_0}\Lambda H_I\right)^{-1}\ |N_{H_0}\rangle \ . \tag{24}$$

Assuming that the π-baryon coupling constants are small we shall keep only the terms up to first order in H_I (i.e., we shall keep only the terms with 1 pion):

$$|\tilde{N}\rangle \sim \sqrt{Z_2^N}\ \left(1 + \frac{1}{M - H_0}\Lambda H_I\right)\ |N_{H_0}\ \rangle. \tag{25}$$

Finally then we get

$$|\tilde{N}\rangle \simeq \sqrt{Z_2^N}\ \left(1 + \frac{1}{M - H_0}H_I\right)\ |N_{H_0}\ \rangle \ , \tag{26}$$

where we have used the fact that the matrix element of H_I between bare states vanishes; i.e., $\langle\ N_{H_0}\ |H_I|\ N_{H_0}\ \rangle = 0$.

B. The quark distributions

We are now able to rewrite the quark distributions within the dressed proton [Eq. (6)] in terms of matrix elements of bare baryons. To do this, we insert complete sets of states:

$$\begin{aligned}\tilde{q}_f^{\uparrow\downarrow}(x) = {} & Z_2^N\langle p|\mathcal{O}_f^{\uparrow\downarrow}(x)|p\rangle \\ & + Z_2^N \sum_{\substack{\alpha,\pi \\ \alpha',\pi'}} \int d\mathbf{k}\, d\mathbf{k}' \left\langle p \right| \ H_I\frac{1}{M - H_0}\ \left|\alpha(\mathbf{k})\right\rangle|\pi(-\mathbf{k})\rangle \\ & \times\langle\alpha(\mathbf{k})|\langle\pi(-\mathbf{k})|\ \ \mathcal{O}_f^{\uparrow\downarrow}(x')\ \ |\alpha'(\mathbf{k}')\rangle|\pi'(-\mathbf{k}')\rangle \\ & \times\langle\alpha'(\mathbf{k}')|\left\langle\pi'(-\mathbf{k}')\right|\ \frac{1}{M - H_0}H_I\ \left|p\right\rangle.\end{aligned} \tag{27}$$

It should be noted that the assumption of the convolution model is contained in the factorization of the intermediate states into pion and baryon states [α (α') and π (π') indicate baryon and pion states, respectively]. The variable x' is defined as

$$x' = \frac{Q^2}{2q\cdot k_{\text{target}}}. \tag{28}$$

As has already been mentioned, we shall restrict ourselves to the probe scattering off the baryon, so $k_{\text{target}} = k_{\text{baryon}}$. Using the orthogonality of the wave functions we may reduce Eq. (27) to

$$\begin{aligned}\tilde{q}_f^{\uparrow\downarrow}(x) = Z_2^N\langle p|\mathcal{O}_f^{\uparrow\downarrow}(x)|p\rangle + Z_2^N \sum_{\alpha,\pi,\alpha'} \int\ & d\mathbf{k}\langle\alpha(\mathbf{k})|\mathcal{O}_f^{\uparrow\downarrow}(x')|\alpha'(\mathbf{k})\rangle \\ & \times\frac{\langle p|\ \ H_I\ \ |\alpha(\mathbf{k})\rangle|\pi(-\mathbf{k})\rangle}{M - M_\alpha - E_\pi}\,\frac{\langle\alpha'(\mathbf{k})|\langle\pi(-\mathbf{k})|\ \ H_I\ \ |p\rangle}{M - M_{\alpha'} - E_\pi}.\end{aligned} \tag{29}$$

Using the definition of H_I we find

$$\langle p|\ \ H_I\ \ |\alpha(\mathbf{k})\rangle|\pi(-\mathbf{k})\rangle = \mathbf{v}_{\mathbf{k}}^{N\alpha}\cdot\mathbf{I}_\pi \ \ \text{and}\ \ \langle\alpha'(\mathbf{k})|\langle\pi(-\mathbf{k})|\ \ H_I\ \ |p\rangle = \mathbf{I}_\pi^\dagger\cdot\mathbf{v}_{\mathbf{k}}^{N\alpha'\,\dagger} \tag{30}$$

where $\mathbf{I}_\pi$ is the isospin vector for the pion. The remaining matrix element in Eq. (29) is the quark distribution in the bare baryon. We shall assume it to be the same as that within a free, on-mass-shell, baryon. It is a quantity invariant under boosts and may be evaluated in any frame, depending only on

$$x' = \frac{Q^2}{2q\cdot k} = \frac{Q^2}{2q\cdot p}\frac{2q\cdot p}{2q\cdot k} = \frac{x}{y}, \tag{31}$$

where y is the fraction of the dressed nucleon's light-cone momentum carried by the interacting baryon, i.e., $y = k^+/p^+$. Let us define the quark distributions within the baryons as

$$\langle\alpha(\mathbf{k})|\mathcal{O}_f^{\uparrow\downarrow}(x')|\alpha'(\mathbf{k})\rangle = \frac{1}{y}q_{\alpha\alpha'}\left(\frac{x}{y}\right), \tag{32}$$

and note that the phase space in Eq. (29) is given by

$$\int d\mathbf{k} = 2\pi M\int_x^1 dy\int_{k_{\min}}^\infty k\,dk \tag{33}$$

(we have performed the azimuthal integration because the integrand is independent of ϕ) with

$$k_{\min} = \left|\frac{(1-y)^2M^2 - M_\pi^2}{2M(1-y)}\right| . \tag{34}$$

The quark distributions may be written in the familiar form of the convolution model:

$$\tilde{q}_f^{\uparrow\downarrow}(x) = Z_2^N q_{f,\,pp}^{\uparrow\downarrow}(x) + \sum_{\alpha,\alpha'}\int_x^1 dy\, f_{\alpha\alpha'}(y)\frac{q_{f,\,\alpha\alpha'}^{\uparrow\downarrow}(\frac{x}{y})}{y} , \tag{35}$$

where we define the light-cone momentum distribution of baryons within the nucleon as

$$f_{\alpha\alpha'}(y) = 2\pi M\int_{k_{\min}}^\infty k\,dk\,F_{\alpha\alpha'}(\mathbf{k}) , \tag{36}$$

with

$$\begin{aligned} F_{\alpha\alpha'}(\mathbf{k}) &= Z_2^N\sum_\pi \frac{\mathbf{v}_{\mathbf{k}}^{N\alpha}\cdot\mathbf{I}_\pi}{M - M_\alpha - E_\pi}\frac{\mathbf{I}_\pi^\dagger\cdot\mathbf{v}_{\mathbf{k}}^{N\alpha'\,\dagger}}{M - M_{\alpha'} - E_\pi} \\ &= Z_2^N\frac{9f_0^{N\alpha}f_0^{N\alpha'}4\pi u(kR)^2\mathcal{S}^{\alpha\alpha'}\mathcal{T}^{\alpha\alpha'}}{(2\pi)^3 2E_\pi M_\pi^2(M - M_\alpha - E_\pi)(M - M_{\alpha'} - E_\pi)} . \end{aligned} \tag{37}$$

In Eq. (37) we have collected spin and isospin factors as

$$\mathcal{S}^{\alpha\alpha'} = C(S_\alpha 1 \to S_N|s_\alpha m \to s_N)C(S_{\alpha'}\,1 \to S_N|s_{\alpha'}\,m' \to s_N)\hat{\mathbf{s}}_m^*\cdot\mathbf{k}\;\hat{\mathbf{s}}_{m'}\cdot\mathbf{k} \tag{38}$$

and

$$\mathcal{T}^{\alpha\alpha'} = T(T_\alpha 1 \to T_N|t_\alpha n \to t_N)T(T_{\alpha'}\,1 \to T_N|t_{\alpha'}\,n' \to t_N)\hat{\mathbf{t}}_n^*\cdot\hat{\mathbf{t}}_{n'} ,$$

respectively. For convenience we have listed these in Table I. The general structure of the quark distributions within the dressed proton may therefore be written in terms of the bare distributions as

$$\begin{aligned}\tilde{q}_f^{\uparrow\downarrow}(x) = Z_2^N q_{f,\,pp}^{\uparrow\downarrow}(x) + \sum_{\alpha,\alpha'}\int_x^1 \frac{dy}{y}[c_\perp(t_\alpha,t_{\alpha'},s_\alpha,s_{\alpha'})f_\perp^{T_\alpha,T_{\alpha'}}(y) + c_z(t_\alpha,t_{\alpha'},s_\alpha,s_{\alpha'})f_z^{T_\alpha,T_{\alpha'}}(y)] \\ \times q_{f,t_\alpha,t_{\alpha'},s_\alpha,s_{\alpha'}}^{\uparrow\downarrow}\left(\frac{x}{y}\right),\end{aligned} \tag{39}$$

where we have defined $f_\perp^{T_\alpha,T_{\alpha'}}(y)$ and $f_z^{T_\alpha,T_{\alpha'}}(y)$ to be given by Eq. (36) with the integrand multiplied by $\mathbf{k}_\perp^2/(2\mathcal{S}^{\alpha\alpha'}\mathcal{T}^{\alpha\alpha'})$ and $\mathbf{k}_z^2/(\mathcal{S}^{\alpha\alpha'}\mathcal{T}^{\alpha\alpha'})$, respectively. The coefficients $c_\perp(t_\alpha,t_{\alpha'},s_\alpha,s_{\alpha'})$ and $c_z(t_\alpha,t_{\alpha'},s_\alpha,s_{\alpha'})$ are tabulated in Tables II and III. The quark distributions $q_{f,t_\alpha,t_{\alpha'},s_\alpha,s_{\alpha'}}^{\uparrow\downarrow}(\frac{x}{y})$ are evaluated between baryon states with third components of spin (isospin) equal to s_α and $s_{\alpha'}$ (t_α and $t_{\alpha'}$)—we suppress the indicies indicating the dependence on the total spin and isospin.

1. The bare distributions

In order to proceed further we need to calculate the quark distributions $q(\frac{x}{y})$ within the bare targets. As already stated, we shall assume that these are the same as if the bare targets were on shell. For the nucleon we shall use those of Ref. [6] [Eqs. (30), (31), and (43)]. They are of the form

TABLE I. The nonzero spin and isospin factors $\mathcal{S}^{\alpha\alpha'}$ and $\mathcal{T}^{\alpha\alpha'}$.

s or t	$S(T)_\alpha$	$S(T)_{\alpha'}$	$\mathcal{S}^{\alpha\alpha'}$	$\mathcal{T}^{\alpha\alpha'}$
$s_{\alpha'} = s_\alpha = \frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{3}\mathbf{k}_z^2$	$\frac{1}{3}$
	$\frac{1}{2}$	$\frac{3}{2}$	$\frac{1}{3}\mathbf{k}_z^2$	$\frac{1}{3}$
	$\frac{3}{2}$	$\frac{1}{2}$	$\frac{1}{3}\mathbf{k}_z^2$	$\frac{1}{3}$
	$\frac{3}{2}$	$\frac{3}{2}$	$\frac{1}{3}\mathbf{k}_z^2$	$\frac{1}{3}$
$s_{\alpha'} = s_\alpha = \frac{3}{2}$	$\frac{3}{2}$	$\frac{3}{2}$	$\frac{1}{2}\frac{\mathbf{k}_\perp^2}{2}$	$\frac{1}{2}$
$s_{\alpha'} = s_\alpha = -\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{2}{3}\frac{\mathbf{k}_\perp^2}{2}$	$\frac{2}{3}$
	$\frac{1}{2}$	$\frac{3}{2}$	$\frac{1}{3}\frac{\mathbf{k}_\perp^2}{2}$	$\frac{1}{3}$
	$\frac{3}{2}$	$\frac{1}{2}$	$\frac{1}{3}\frac{\mathbf{k}_\perp^2}{2}$	$\frac{1}{3}$
	$\frac{3}{2}$	$\frac{3}{2}$	$\frac{1}{6}\frac{\mathbf{k}_\perp^2}{2}$	$\frac{1}{6}$

TABLE II. The coefficients $c_\perp(t_\alpha, t_{\alpha'}, s_\alpha, s_{\alpha'})$.

t_α	$t_{\alpha'}$	$s_\alpha = s_{\alpha'}$	$c_\perp(t_\alpha, t_{\alpha'}, s_\alpha, s_{\alpha'})$
Δ^{++}	Δ^{++}	$\frac{3}{2}$	$\frac{1}{4}$
Δ^{++}	Δ^{++}	$-\frac{1}{2}$	$\frac{1}{12}$
Δ^{+}	Δ^{+}	$\frac{3}{2}$	$\frac{1}{6}$
Δ^{+}	Δ^{+}	$-\frac{1}{2}$	$\frac{1}{18}$
Δ^{0}	Δ^{0}	$\frac{3}{2}$	$\frac{1}{12}$
Δ^{0}	Δ^{0}	$-\frac{1}{2}$	$\frac{1}{36}$
p	p	$-\frac{1}{2}$	$\frac{2}{9}$
n	n	$-\frac{1}{2}$	$\frac{4}{9}$
n	Δ^{0}	$-\frac{1}{2}$	$\frac{2}{9}$
p	Δ^{+}	$-\frac{1}{2}$	$\frac{2}{9}$

TABLE III. The coefficients $c_z(t_\alpha, t_{\alpha'}, s_\alpha, s_{\alpha'})$.

t_α	$t_{\alpha'}$	$s_\alpha = s_{\alpha'}$	$c_z(t_\alpha, t_{\alpha'}, s_\alpha, s_{\alpha'})$
Δ^{++}	Δ^{++}	$\frac{1}{2}$	$\frac{1}{6}$
Δ^{+}	Δ^{+}	$\frac{1}{2}$	$\frac{1}{9}$
Δ^{0}	Δ^{0}	$\frac{1}{2}$	$\frac{1}{18}$
p	p	$\frac{1}{2}$	$\frac{1}{9}$
n	n	$\frac{1}{2}$	$\frac{2}{9}$
n	Δ^{0}	$\frac{1}{2}$	$\frac{2}{9}$
p	Δ^{+}	$\frac{1}{2}$	$\frac{2}{9}$

$$
\begin{aligned}
u^{\uparrow\downarrow}_{p,p,\frac{1}{2},\frac{1}{2}}(x) &= w_2^u(p,p,\tfrac{1}{2})F_{(2)}(x) \\
&\quad \pm w_2^{\Delta u}(p,p,\tfrac{1}{2})G_{(2)}(x) + 3F_{(4)}(x), \\
d^{\uparrow\downarrow}_{p,p,\frac{1}{2},\frac{1}{2}}(x) &= w_2^d(p,p,\tfrac{1}{2})F_{(2)}(x) \\
&\quad \pm w_2^{\Delta d}(p,p,\tfrac{1}{2})G_{(2)}(x) + 3F_{(4)}(x), \\
\overline{u}^{\uparrow\downarrow}_{p,p,\frac{1}{2},\frac{1}{2}}(x) &= w_4^{\overline{u}}(p,p,\tfrac{1}{2})F_{(4)}(x) \\
&\quad \pm w_4^{\Delta\overline{u}}(p,p,\tfrac{1}{2})G_{(4)}(x), \\
\overline{d}^{\uparrow\downarrow}_{p,p,\frac{1}{2},\frac{1}{2}}(x) &= w_4^{\overline{d}}(p,p,\tfrac{1}{2})F_{(4)}(x) \\
&\quad \pm w_4^{\Delta\overline{d}}(p,p,\tfrac{1}{2})G_{(4)}(x) \quad .
\end{aligned}
\tag{40}
$$

$F(x)$ and $G(x)$ allow for the different shapes of the spin independent and spin dependent distributions, respectively. (The shapes are necessarily different in any model where the quarks are confined and thus have nonzero perpendicular momentum. The difference is indicative of the fact that the former is related to the number of quarks in the target, which is independent of whether or not the quarks are relativistic, while the latter is proportional to $|g_a/g_v|$, which is reduced in the bag as compared to a nonrelativistic model of the nucleon.) The subscripts on these distributions refer to the mass of the spectator system of the hard scattering process. For further details we refer the reader to Ref. [6]. The expressions for the other baryons may be written down in analogy. The coefficients w are tabulated in Table IV. The coefficients for the antiquarks are given by

$$
\begin{aligned}
&w_4^{\overline{q}}(T_\alpha = T_{\alpha'}, s_\alpha) = 3 - w_2^q(T_\alpha = T_{\alpha'}, s_\alpha), \\
&w_4^{\overline{q}}(T_\alpha \neq T_{\alpha'}, s_\alpha) = 0, \\
&w_4^{\Delta\overline{q}}(T_\alpha, T_{\alpha'}, s_\alpha) = w_2^{\Delta q}(T_\alpha, T_{\alpha'}, s_\alpha) \quad .
\end{aligned}
\tag{41}
$$

2. *The dressed distributions*

With these coefficients we may now write down the final expressions for the quark and antiquark distributions. Defining

$$u(x) = u_{\text{bare}}(x) + u_{\text{dressed}}(x), \quad \text{etc.,} \tag{42}$$

we have

$$
\begin{aligned}
&u_{\text{bare}}(x) = [2F_N^{(2)}(x) + 6F_N^{(4)}(x)]Z_2^N, \quad d_{\text{bare}}(x) = [F_N^{(2)}(x) + 6F_N^{(4)}(x)]Z_2^N, \\
&\Delta u_{\text{bare}}(x) = \frac{4}{3}G_N^{(2)}(x)Z_2^N, \quad \Delta d_{\text{bare}}(x) = -\tfrac{1}{3}G_N^{(2)}(x)Z_2^N, \\
&\overline{u}_{\text{bare}}(x) = 4F_N^{(4)}(x)Z_2^N, \quad \overline{d}_{\text{bare}}(x) = 5F_N^{(4)}(x)Z_2^N, \\
&\Delta\overline{u}_{\text{bare}}(x) = \tfrac{4}{3}G_N^{(4)}(x)Z_2^N, \quad \Delta\overline{d}_{\text{bare}}(x) = -\tfrac{1}{3}G_N^{(4)}(x)Z_2^N,
\end{aligned}
\tag{43}
$$

and

$$
\begin{aligned}
u_{\text{dressed}}(x) = \int_x^1 \frac{dy}{y}\Bigg\{ &[f_z^{NN}(y) + 2f_\perp^{NN}(y)]\left[\frac{4}{9}F_N^{(2)}\left(\frac{x}{y}\right) + 2F_N^{(4)}\left(\frac{x}{y}\right)\right] \\
&+[f_z^{\Delta\Delta}(y) + 2f_\perp^{\Delta\Delta}(y)]\left[\frac{7}{9}F_\Delta^{(2)}\left(\frac{x}{y}\right) + 2F_\Delta^{(4)}\left(\frac{x}{y}\right)\right]\Bigg\},
\end{aligned}
$$

$$
\begin{aligned}
d_{\text{dressed}}(x) = \int_x^1 \frac{dy}{y}\Bigg\{ &[f_z^{NN}(y) + 2f_\perp^{NN}(y)]\left[\frac{5}{9}F_N^{(2)}\left(\frac{x}{y}\right) + 2F_N^{(4)}\left(\frac{x}{y}\right)\right] \\
&+[f_z^{\Delta\Delta}(y) + 2f_\perp^{\Delta\Delta}(y)]\left[\frac{2}{9}F_\Delta^{(2)}\left(\frac{x}{y}\right) + 2F_\Delta^{(4)}\left(\frac{x}{y}\right)\right]\Bigg\},
\end{aligned}
$$

$$\Delta u_{\rm dressed}(x) = \int_x^1 \frac{dy}{y}\left\{ \frac{2}{27}\left[f_z^{NN}(y) - 2f_\perp^{NN}(y)\right] G_N^{(2)}\left(\frac{x}{y}\right) + \frac{7}{27}\left[f_z^{\Delta\Delta}(y) + 4f_\perp^{\Delta\Delta}(y)\right] G_\Delta^{(2)}\left(\frac{x}{y}\right) + \frac{8\sqrt{2}}{27}\left[f_z^{N\Delta}(y) + f_\perp^{N\Delta}(y)\right] G_{N\Delta}^{(2)}\left(\frac{x}{y}\right)\right\},$$

$$\Delta d_{\rm dressed}(x) = \int_x^1 \frac{dy}{y}\left\{ \frac{7}{27}\left[f_z^{NN}(y) - 2f_\perp^{NN}(y)\right] G_N^{(2)}\left(\frac{x}{y}\right) + \frac{2}{27}\left[f_z^{\Delta\Delta}(y) + 4f_\perp^{\Delta\Delta}(y)\right] G_\Delta^{(2)}\left(\frac{x}{y}\right) - \frac{8\sqrt{2}}{27}\left[f_z^{N\Delta}(y) + f_\perp^{N\Delta}(y)\right] G_{N\Delta}^{(2)}\left(\frac{x}{y}\right)\right\},$$

$$\overline{u}_{\rm dressed}(x) = \int_x^1 \frac{dy}{y}\left\{ \frac{14}{9}\left[f_z^{NN}(y) + 2f_\perp^{NN}(y)\right] F_N^{(4)}\left(\frac{x}{y}\right) + \frac{11}{9}\left[f_z^{\Delta\Delta}(y) + 2f_\perp^{\Delta\Delta}(y)\right] F_\Delta^{(4)}\left(\frac{x}{y}\right)\right\}, \quad (44)$$

$$\overline{d}_{\rm dressed}(x) = \int_x^1 \frac{dy}{y}\left\{ \frac{13}{9}\left[f_z^{NN}(y) + 2f_\perp^{NN}(y)\right] F_N^{(4)}\left(\frac{x}{y}\right) + \frac{16}{9}\left[f_z^{\Delta\Delta}(y) + 2f_\perp^{\Delta\Delta}(y)\right] F_\Delta^{(4)}\left(\frac{x}{y}\right)\right\},$$

$$\Delta \overline{u}_{\rm dressed}(x) = \int_x^1 \frac{dy}{y}\left\{ \frac{2}{27}\left[f_z^{NN}(y) - 2f_\perp^{NN}(y)\right] G_N^{(4)}\left(\frac{x}{y}\right) + \frac{7}{27}\left[f_z^{\Delta\Delta}(y) + 4f_\perp^{\Delta\Delta}(y)\right] G_\Delta^{(4)}\left(\frac{x}{y}\right) + \frac{8\sqrt{2}}{27}\left[f_z^{N\Delta}(y) + f_\perp^{N\Delta}(y)\right] G_{N\Delta}^{(4)}\left(\frac{x}{y}\right)\right\},$$

$$\Delta \overline{d}_{\rm dressed}(x) = \int_x^1 \frac{dy}{y}\left\{ \frac{7}{27}\left[f_z^{NN}(y) - 2f_\perp^{NN}(y)\right] G_N^{(4)}\left(\frac{x}{y}\right) + \frac{2}{27}\left[f_z^{\Delta\Delta}(y) + 4f_\perp^{\Delta\Delta}(y)\right] G_\Delta^{(4)}\left(\frac{x}{y}\right) - \frac{8\sqrt{2}}{27}\left[f_z^{N\Delta}(y) + f_\perp^{N\Delta}(y)\right] G_{N\Delta}^{(4)}\left(\frac{x}{y}\right)\right\} .$$

TABLE IV. The nonzero coefficients $w(t_\alpha, t_{\alpha'}, s_\alpha)$.

t_α	$t_{\alpha'}$	$s_\alpha = s_{\alpha'}$	$w_2^u(t_\alpha, t_{\alpha'}, s_\alpha)$	$w_2^d(t_\alpha, t_{\alpha'}, s_\alpha)$	$w_2^{\Delta u}(t_\alpha, t_{\alpha'}, s_\alpha)$	$w_2^{\Delta d}(t_\alpha, t_{\alpha'}, s_\alpha)$
Δ^{++}	Δ^{++}	$\frac{3}{2}$	$\frac{3}{2}$	0	$\frac{3}{2}$	0
Δ^{++}	Δ^{++}	$\frac{1}{2}$	$\frac{3}{2}$	0	$\frac{1}{2}$	0
Δ^{++}	Δ^{++}	$-\frac{1}{2}$	$\frac{3}{2}$	0	$-\frac{1}{2}$	0
Δ^{+}	Δ^{+}	$\frac{3}{2}$	1	$\frac{1}{2}$	1	$\frac{1}{2}$
Δ^{+}	Δ^{+}	$\frac{1}{2}$	1	$\frac{1}{2}$	$\frac{1}{3}$	$\frac{1}{6}$
Δ^{+}	Δ^{+}	$-\frac{1}{2}$	1	$\frac{1}{2}$	$-\frac{1}{3}$	$-\frac{1}{6}$
Δ^{0}	Δ^{0}	$\frac{3}{2}$	$\frac{1}{2}$	1	$\frac{1}{2}$	1
Δ^{0}	Δ^{0}	$\frac{1}{2}$	$\frac{1}{2}$	1	$\frac{1}{6}$	$\frac{1}{3}$
Δ^{0}	Δ^{0}	$-\frac{1}{2}$	$\frac{1}{2}$	1	$-\frac{1}{6}$	$-\frac{1}{3}$
p	p	$\frac{1}{2}$	1	$\frac{1}{2}$	$\frac{2}{3}$	$-\frac{1}{6}$
p	p	$-\frac{1}{2}$	1	$\frac{1}{2}$	$-\frac{2}{3}$	$\frac{1}{6}$
n	n	$+\frac{1}{2}$	$\frac{1}{2}$	1	$-\frac{1}{6}$	$\frac{2}{3}$
n	n	$-\frac{1}{2}$	$\frac{1}{2}$	1	$\frac{1}{6}$	$-\frac{2}{3}$
n	Δ^{0}	$+\frac{1}{2}$	0	0	$\frac{\sqrt{2}}{3}$	$-\frac{\sqrt{2}}{3}$
n	Δ^{0}	$-\frac{1}{2}$	0	0	$\frac{\sqrt{2}}{3}$	$-\frac{\sqrt{2}}{3}$
p	Δ^{+}	$+\frac{1}{2}$	0	0	$\frac{\sqrt{2}}{3}$	$-\frac{\sqrt{2}}{3}$
p	Δ^{+}	$-\frac{1}{2}$	0	0	$\frac{\sqrt{2}}{3}$	$-\frac{\sqrt{2}}{3}$

Several interesting observations may be made from Eq. (44).

(1) The spin-dependent and spin-independent quark distributions have different shapes. The reason for this is that both the bare quark distributions as well as the baryon distributions themselves are different for the spin-dependent and spin-independent cases. The origin of both of these effects is that the perpendicular momenta (with respect to the incoming photon momentum) are not zero.

(2) The integrals over y of $f_z(y)$ and $f_\perp(y)$ are the same. If the "Pauli defect," which is related to the magnitude of $F_4(x)$ and $G_4(x)$, were absent [i.e., if $F_4(x)=G_4(x)=0$ and hence the distributions $F_2(x)$ and $G_2(x)$ were normalized] then the integrals of the distributions could be expressed directly in terms of probabilities, e.g., $u_{\rm dressed} = \frac{4}{3}{\rm Prob}(N\pi) + \frac{7}{3}{\rm Prob}(\Delta\pi)$, etc.

(3) Interference terms between Δ's and N's only occur for the spin-dependent distributions (see Ref. [16]).

(4) The total number of valence quarks is three—this is of course unaffected by neglecting the pionic contribution and relies on the fact that $F_2(x) + F_4(x)$ is by definition normalized to one.

III. NUMERICAL RESULTS

For most structure functions the effect of including a pionic component to the nucleon wave function amounts to typically a 10–30 % effect. Unfortunately the uncertainties in the "bare" distributions themselves are probably of the same order (Ref. [6]; for example, the mass of the intermediate diquark state is unknown; we shall keep it as a variable parameter). We shall therefore restrict ourselves to calculating those distributions (or combinations thereof) where we might expect to see the largest effect: $[F_2^p(x) - F_2^n(x)]/x$, whose integral is given by the Gottfried sum rule (Ref. [27]), and $g_1^n(x)$. The former has recently been measured by the NMC [9] and was found to be 0.240 ± 0.016, significantly below the expected value in an SU(6) symmetric model of $\frac{1}{3}$, while the latter has not yet been measured and would be 0 for all x if SU(6) symmetry were exact.

In order to be definite we shall assume that $F_N^{(2,4)}(x) = F_\Delta^{(2,4)}(x) = F_{N\Delta}^{(2,4)}(x)$ and $G_N^{(2,4)}(x) = G_\Delta^{(2,4)}(x) = G_{N\Delta}^{(2,4)}(x)$. For further details the reader is directed to Ref. [6]. A source of uncertainty in any model calculation of structure functions is our ignorance about the Q^2 scale at which the model applies. We shall present the results for a variety of scales and use the well known leading order QCD evolution formalism, with three flavors and $\Lambda_{\rm QCD} = 0.2$ GeV, in order to compare our results to the data from Ref. [9]. (Indeed, the experimental distributions in x are a function of Q^2 themselves, with a typical value of $Q^2 = 4$ GeV^2. Our results are evolved to this value of Q^2.)

The results for $F_2^p(x) - F_2^n(x)$ and $g_1^n(x)$, for two values of the $Q^2 = \mu^2$ scale at which the model applies, are shown in Figs. 2 and 3 for bag radii of 0.7 and 1.0 fm and masses of the intermediate diquark states of 550 MeV and $\frac{3}{4}$ M. Increasing the mass of the diquark state decreases the average fraction of momentum carried by the struck quark and therefore shifts the peak of the distributions to smaller x. Decreasing the bag radius, apart from slightly broadening the bare distributions (an effect to be expected from the Heisenberg uncertainty relations), decreases the magnitude of the distribution. This is because it increases both the "Pauli defect" and the pionic dressing.

The predicted value of the Gottfried sum rule, ranging from 0.17 ($M_2 = \frac{3}{4}$ M) to 0.22 (M_2 =550 MeV) for R = 1.0 fm, tends to be somewhat below the experimental value (it is even smaller for R = 0.7 fm). Indeed, just the "Pauli defect" alone (for R = 1.0 fm and $M_2 = \frac{3}{4}$ M) yields a result of 0.21 while the dressing alone would give 0.27. The latter value is essentially given by the "bare" contribution, which has changed from $\frac{1}{3}$ to $Z_2^N/3$ (Z_2^N is 0.75 for R = 1 fm). The additional contributions from the photon scattering off the nucleon and the Δ (with a pion spectator) are of approximately equal magnitude (≈ 0.02) but are of opposite sign and therefore tend to cancel each other.

As is evident from Fig. 2, the deficit in the sum rule originates largely from the large-x region. This is a rather general feature of these bag model calculations [6]—all structure functions tend to become vanishingly small for

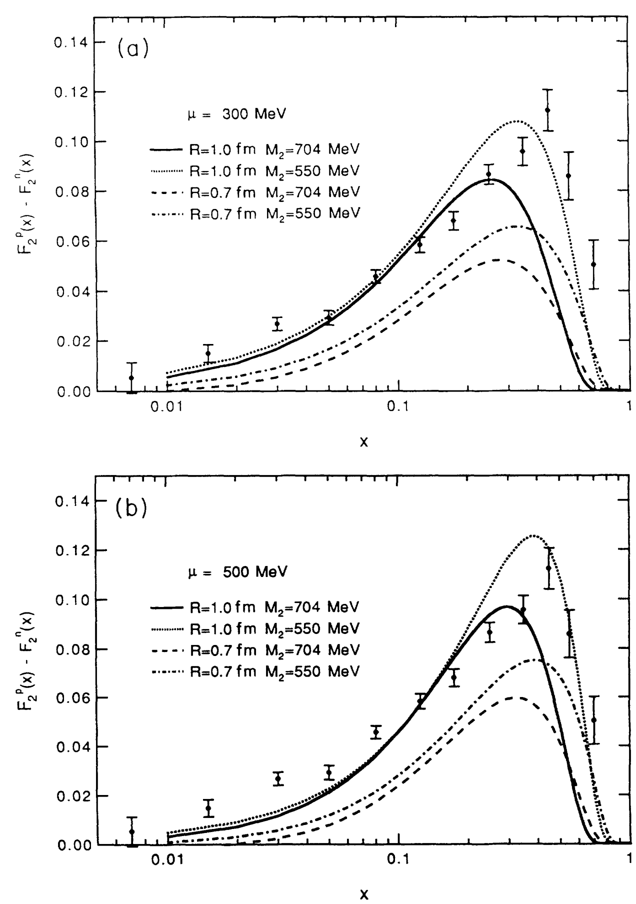

FIG. 2. $F_2^p(x) - F_2^n(x)$ for the model scales (a) μ = 300 MeV and (b) 500 MeV, corresponding to the quarks carrying, at $Q^2 = 4$ GeV^2, respectively 54% and 63% of the momentum that they carried at $Q^2 = \mu^2$. The data are taken from [9] and the curves have been evolved to $Q^2 = 4$ GeV^2.

$x > 0.8$. It is possible that this feature is related to the use of the Peierls-Yoccoz projection which is used in order to obtain the nucleon momentum eigenstates [28]. In any case, it is unrelated to the pion dressing of the nucleon as it is already present in the bare distributions [6].

In contrast with $F_2^p(x) - F_2^n(x)$, $g_1^n(x)$ is not affected by the "Pauli defect" and hence almost the entire effect is due to the pion dressing. There is also a contribution due to QCD evolution: the flavor singlet and flavor nonsinglet evolve differently, so if they cancel precisely at $Q^2 = \mu^2$, they will not do so at other Q^2. As an example of the size of this effect we show the evolved distribution for the bare neutron (i.e., we artificially set $Z_2^N = 1$) for $R = 1.0$ fm and $M_2 = \frac{3}{4}$ M in Fig. 3(a). The integral of $g_1^n(x)$ ranges from -0.015 (R=0.7 fm) to -0.011 (R=1.0 fm). This is much smaller in magnitude than the corresponding number in Ref. [29] as well as the expected value from the Bjorken sum rule (using the EMC's result for the integral of $g_1^p(x)$ [7, 8]).

It is interesting to note that the value of $g_1^n(x)$ is negative for almost all x. This is in contrast with the distributions predicted on the basis of one-gluon exchange within the neutron [6, 29, 30], which tend to be positive in the large-x region and cross over around $x \approx 0.1 - 0.4$. The x distribution of another possible contribution, due to the anomaly [31–33], is essentially unknown because for all but the first moment the division between perturbative and nonperturbative regions is infrared sensitive [34, 35]. Presumably, however, it is concentrated at small x because that is where the gluon distribution is largest.

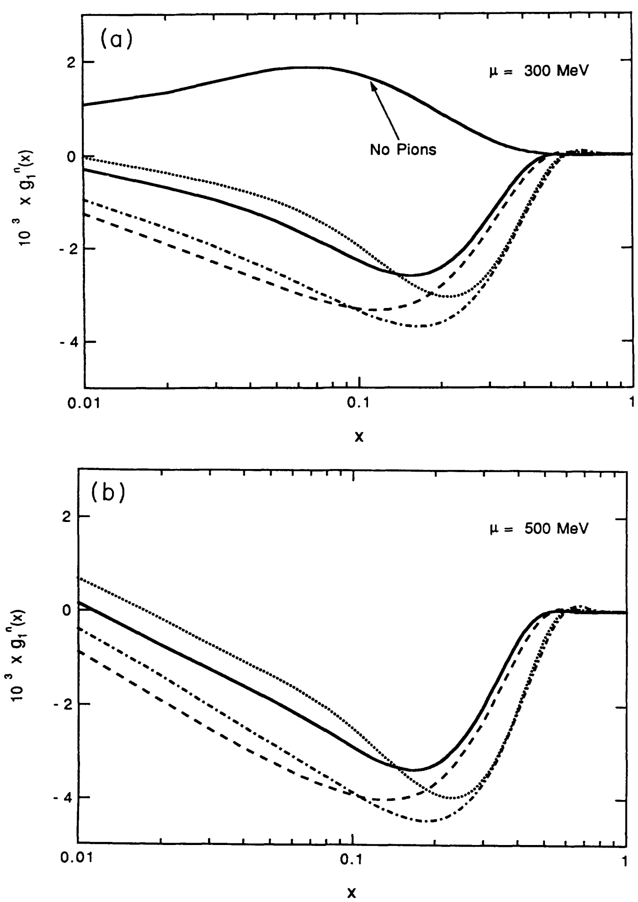

FIG. 3. The neutron structure function $g_1^n(x)$. The curves correspond to those in Fig. 2. Also shown, in (a), is the shape of the distribution expected from QCD evolution alone.

IV. SUMMARY

In this paper we have generalized the calculation of structure functions within the bag model presented in [6] to nucleon targets dressed by a pion cloud. The calculation is done within the convolution model and scattering of the incoming probe off the pion cloud itself is neglected. At small x the latter should be taken into account for general structure functions, for example, in the calculation of $F_2^p(x)$, etc. We have focused on those structure functions that are likely to be most affected by the dressing of the nucleon wave function and are unaffected by neglecting the contribution of the (probe) scattering off the pions. The results for $F_2^p(x) - F_2^n(x)$ are, in particular, for large bag radii, in general agreement with the data except perhaps for large x. The contribution to the neutron spin structure function $g_1^n(x)$ is likely to be rather small and negative over the entire range of x. This may have important experimental consequences. At large x the effect of one-gluon exchange is precisely in the opposite direction, so that one might naively expect the pionic contribution to decrease, or even roughly cancel it, in this region (to see whether this is indeed the case requires the incorporation of both effects within one model, which has not yet been done). Indeed, it might be possible to distinguish the relative importance of the two mechanisms. We recall that in the calculation of the neutron electric form factor, for example, this is not possible as both effects lead to a negative charge square radius [14, 36].

One of the most striking outcomes of the calculation is that the addition of a pionic component to the nucleon wave function leads to significant changes in the valence distribution, i.e., also at large x. As we have not calculated the "sea" ("sea" being defined as those contributions with a "sea"-like shape in x) contribution of the pions, *all* of the changes presented in this paper correspond to changes in the "valence" distribution. In particular, *none* of the decrease of the Gottfried sum rule originates from the "sea." The changes in the valence distribution are due to a reduced probability to observe a bare nucleon and to a change in the spin-isospin structure of the dressed nucleon (as compared to the bare one). At times it has previously been assumed that the chiral structure of the nucleon will only affect the sea distribution at small x [37]. That this is not so has already been pointed out in Refs. [24, 38].

Finally, the inclusion of other mesons in the theory, in particular vector mesons, is likely to lead to smaller corrections, due to their heavier masses.

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