Do we expect light flavor sea-quark asymmetry also for
the spin-dependent distribution functions of the nucleon?

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Abstract

After taking account of the scale dependence by means of the standard DGLAP evolution
equation, the theoretical predictions of the chiral quark soliton model for the unpolarized
and longitudinally polarized structure functions of the nucleon are compared with the recent
high energy data. The theory is shown to explain all the qualitative features of the experi-
ments, including the NMC data for $F_2^p(x) - F_2^n(x)$, $F_2^p(x)/F_2^n(x)$, the Hermes and NuSea
data for $\bar{d}(x) - \bar{u}(x)$, the EMC and SMC data for $g_1^p(x)$, $g_1^n(x)$ and $g_1^d(x)$. Among others,
flavor asymmetry of the longitudinally polarized sea-quark distributions is a remarkable pre-
diction of this model, i.e., it predicts that $\Delta \bar{d}(x) - \Delta \bar{u}(x) = C x^\alpha [\bar{d}(x) - \bar{u}(x)]$ with a sizable
negative coefficient $C \simeq -2.0$ (and $\alpha \simeq 0.12$) in qualitative consistency with the recent
semi-phenomenological analysis by Morii and Yamanishi.

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Undoubtedly, the most natural and widely-accepted explanation of the light flavor sea-quark asymmetry in the nucleon, revealed by the NMC measurements \[1\], would be the one by the pion cloud effects \[2\]–[5]. Accepting its essential validity, the NMC observation may be taken as the first clear evidence manifesting the nonperturbative QCD dynamics of spontaneous chiral symmetry breaking in high-energy deep-inelastic observables. This flavor asymmetry of the spin-averaged sea-quark distributions could have been predicted by some theoretical models which properly incorporate the chiral symmetry, but unfortunately, most theoretical explanations were given only after the NMC observation.

Now, a natural question is “Do we expect light flavor sea-quark asymmetry not only for the unpolarized distributions but also for the spin-dependent ones?” In our opinion, to make a trustworthy prediction for spin-dependent sea-quark distributions, the chiral symmetry would be a minimum ingredient to be incorporated into model construction. Otherwise, it would be hard to explain the already-established flavor asymmetry of the unpolarized sea-quark distributions. Also desirable is a reasonable reproduction of the existing high energy data for the unpolarized and the longitudinally polarized structure functions.

We claim that the chiral quark soliton model (CQSM) of the nucleon \[6\]–[12] fulfills the above requirements and that it is in this sense qualified as a good candidate which is able to give some definite prediction for the flavor asymmetry of the spin-dependent sea-quark distributions. In the present report, we shall first compare the theoretical predictions of the CQSM with the recent high energy data for the unpolarized as well as the longitudinally polarized structure functions, including the NMC data for \(F_p^2(x) - F_n^2(x)\), \(F_n^2(x)/F_p^2(x)\) \[1\], the HERMES \[14\] and NuSea (E866) data \[15\] for \(\bar{d}(x) - \bar{u}(x)\), the EMC \[16\],\[17\],\[18\] and SMC data \[19\] for \(g_1^p(x)\), \(g_1^n(x)\) and \(g_1^d(x)\). After confirming that the theory is able to explain all the qualitative feature of these experimental data, we finally give a prediction for the longitudinally polarized anti-quark distributions \(\Delta \bar{d}(x) - \Delta \bar{u}(x)\) in the proton.

The chiral quark soliton model (CQSM) is a very simple model of the nucleon \[6\],\[7\], which maximally takes account of the spontaneous chiral symmetry breaking of the QCD vacuum. It is specified by the effective lagrangian:

\[
\mathcal{L}_0 = \bar{\psi} \left( i \not \! D - M e^{i \gamma_5 \mathbf{T} \cdot \mathbf{\pi}(x)/f_{\pi}} \right) \psi,
\]

(1)

which describes the effective quark fields with a dynamically generated mass \(M\), interacting with massless (or nearly massless) pions. The theory is not a renormalizable one and it is defined with some ultraviolet cutoff \(\Lambda\). The above effective lagrangian was originally introduced by Diakonov and Petrov on the basis of the instanton-liquid picture of the QCD vacuum \[20\]. In this derivation, the dynamical quark mass \(M\) and the physical cutoff \(\Lambda\) can both be estimated as functions of the average instanton size \(\bar{\rho}\) with \(1/\bar{\rho} \simeq 600\) MeV and the average separation between instantons \(\bar{R}\) with \(1/\bar{R} \simeq 200\) MeV. Their estimate gives \(M \simeq 350\) MeV...
and $\Lambda \simeq 1/\bar{\rho} \simeq 600\text{MeV}$. To be more precise, the predicted dynamical quark mass $M$ is momentum dependent, but in its application to the chiral soliton problem, it is customary to approximate it by a momentum independent constant. At the same time, the ultraviolet cutoff is fixed by the condition that the effective meson lagrangian derived from (1) reproduces the correct normalization of the pion kinetic term. For instance, in the Pauli-Villars regularization scheme, which is used throughout the present analysis, what plays the role of a ultraviolet cutoff is the Pauli-Villars mass $M_{PV}$ obeying the relation:

$$
\frac{N_c M^2}{4\pi^2} \ln \left( \frac{M_{PV}}{M} \right)^2 = f_{\pi}^2,
$$

with $f_{\pi}$ the pion weak decay constant. Using the value of $M \simeq 375\text{MeV}$, which is favored from the phenomenology of nucleon low energy observables, this gives $M_{PV} \simeq 562\text{MeV}$ in qualitative consistency with the estimate based on the instanton picture [8]. (We refer to [13] for the justification of the single-subtraction Pauli-Villars scheme in the studies of quark distribution functions.) Since we are to use these values of $M$ and $M_{PV}$, there is no free parameter additionally introduced in the calculation of distribution functions. The model is also compatible with the idea of large $N_c$ QCD advocated by Witten many years ago [22]. According to him, in the limit of $N_c = \infty$, the QCD is equivalent to an effective theory of mesons, and a baryon appears as a topological soliton of this effective meson lagrangian. Though an appealing idea, a practical problem of this scenario is that the relevant effective meson theory may not be so simple as that of the Skyrme model. In fact, who can imagine an effective meson theory, which is able to describe deep-inelastic scattering observables of the nucleon? On the contrary, the chiral quark soliton model is an effective theory, which realizes the idea of large $N_c$ QCD in more economical way. In fact, in the stationary-phase approximation, which is exact in the large $N_c$ limit [3], the theory is known to have a self-consistent soliton-like solution of hedgehog shape. The energy degeneracy of this soliton configuration under the spatial and isospin rotation induces spontaneous rotational motion of hedgehog soliton. The semiclassical quantization of this collective rotation leads to a systematic method of calculation of any nucleon observables, which is given in a perturbative series in the collective angular velocity operator $\Omega$. (This takes the form of a $1/N_c$ expansion, since $\Omega$ itself is an $1/N_c$ quantity.) The numerical method of calculation of the static nucleon observables was established many years ago [4], and it has been successfully applied to many interesting observables [21]. However, the calculation of quark distribution functions, especially with full inclusion of the vacuum polarization contributions, is quite involved. The calculation of all the combinations of the twist-2 distributions has been completed only recently [8], [10]–[12].

We first summarize in Fig.1 our final results for the unpolarized and longitudinally polarized distribution functions of the leading twist 2. (In the present investigation, some refinement has been achieved in the numerical method of calculation of these distribution functions in the
small $x$ region. As a consequence, there are some minor changes from the results reported in our previous paper [12], especially for the distribution $\Delta u(x) + \Delta d(x)$. However, these changes are so small that they hardly leave appreciable influence on the structure functions obtained after $Q^2$-evolution.)

Figure 1: The theoretical predictions of the CQSM for the unpolarized distributions $u(x) + d(x)$ and $u(x) - d(x)$ as well as for the longitudinally polarized distributions $\Delta u(x) + \Delta d(x)$ and $\Delta u(x) - \Delta d(x)$. In all the figures, the long-dashed and dash-dotted curves respectively stand for the contributions of the discrete valence level and that of the Dirac continuum in the self-consistent hedgehog background, whereas their sums are shown by the solid curves. The distribution functions in the negative $x$ region are to be interpreted as the antiquark ones according to the rules (7)–(10).

We recall that, within the theoretical framework of the flavor SU(2) CQSM, any nucleon observables are evaluated in either of the isoscalar combination or the isovector one [8,11]. In fact, theoretical formulas for them have totally dissimilar form reflecting the fact that they
have different dependence on the collective angular velocity $\Omega$ given as $[8],[11],[12]$

$$u(x) + d(x) \sim N_c O(\Omega^0) \sim O(N_c^1),$$

$$u(x) - d(x) \sim N_c O(\Omega^1) \sim O(N_c^0),$$

$$\Delta u(x) + \Delta d(x) \sim N_c O(\Omega^1) \sim O(N_c^0),$$

$$\Delta u(x) - \Delta d(x) \sim N_c [O(\Omega^0) + O(\Omega^1)] \sim O(N_c^1) + O(N_c^0),$$

where use has been made of the fact that $\Omega$ scales as $1/N_c.$ In Fig.1, quark distributions with negative $x$ are to be interpreted as antiquark distributions according to the rule $[8]:$

$$u(-x) + d(-x) = -[\bar{u}(x) + \bar{d}(x)], \quad (0 < x < 1),$$

$$u(-x) - d(-x) = -[\bar{u}(x) - \bar{d}(x)], \quad (0 < x < 1),$$

$$\Delta u(-x) + \Delta d(-x) = \Delta \bar{u}(x) + \Delta \bar{d}(x), \quad (0 < x < 1),$$

$$\Delta u(-x) - \Delta d(-x) = \Delta \bar{u}(x) - \Delta \bar{d}(x), \quad (0 < x < 1).$$

Here, the long-dashed and dash-dotted curves respectively stand for the contribution of the discrete valence level (it is a deep bound state emerging from the positive energy continuum under the influence of hedgehog mean field) and that of the negative energy Dirac sea (the latter is sometimes called the vacuum polarization contribution), while their sums are shown by the solid curves. We emphasize that the separation into the valence and sea-quark contributions is highly model-dependent concept, and too much significance should not be given to it. Let us first look into $u(x) + d(x)$ in Fig.1(a). This distribution, which emerges at the leading order of $N_c,$ was first evaluated by Diakonov et al. $[8].$ As emphasized by them, the proper account of the vacuum polarization contribution is essential here. In fact, the “valence-quark-only” approximation would lead to $\bar{u}(x) + \bar{d}(x) < 0$ $[9],)$ thereby violating the positivity of the parton distribution. After including the vacuum polarization contribution, on the other hand, this fundamental requirement is satisfied finely. Next, Fig.1(b) shows the isovector combination of the unpolarized distribution function $u(x) - d(x).$ The vacuum polarization contribution is sizable also for this quantity. Among others, the significant positivity of $u(x) - d(x)$ in the negative $x$ region denotes that $\bar{u}(x) - \bar{d}(x) < 0$ for physical value of $x$ with $0 < x < 1,$ thereby indicating the excess of $\bar{d}$ sea over the $\bar{u}$ sea in the proton. This result follows from the fact that the effect of pion cloud is automatically incorporated into the framework of the CQSM $[23].$ We shall later compare the above prediction of the CQSM directly with the existing high energy data.

Next, we turn to the longitudinally polarized distributions. The isoscalar combination $\Delta u(x) + \Delta d(x)$ and the isovector one $\Delta u(x) - \Delta d(x)$ are respectively shown in Fig.1(c) and Fig.1(d). The isoscalar distribution survives only at the first order of $\Omega$ or equivalently at $O(N_c^0).$ This is because the leading-order term vanishes identically because of the hedgehog
symmetry. On the other hand, the leading contribution to the isovector distribution arises from the $O(\Omega^0)$ or $O(N_1^1)$ term. Although the main feature of the isovector longitudinally polarized distributions are controlled by this leading order contribution, the next-to-leading-order contribution is also important. In fact, without this $1/N_c$ correction, the first moment of $\Delta u(x) - \Delta d(x)$, i.e. the isovector axial charge of the nucleon $g_A^{(3)}$ would be largely underestimated. (Incidentally, the presence of this first order rotational correction to $g_A^{(3)}$ is a distinguishable feature of our effective quark theory as compared with a deeply connected soliton model of the nucleon, i.e. the Skyrme model [24],[25]. This breakdown of the so-called Cheshire-cat principle can be understood as originating from the noncommutativity of the bosonization procedure and the collective quantization procedure of the rotational motion [26].) The numerical results also shows that the isoscalar and isovector distributions have rather different behaviors as functions of $x$. The first thing one may notice for the isoscalar distribution is the smallness of the sea quark contribution. (At least the smallness of the sea quark contribution to the first moment of $\Delta u(x) + \Delta d(x)$ has long been known to us [7].) The smallness of the sea quark contribution here does not necessarily mean negligible role of the vacuum polarization effect itself. Since our valence quark level is a deep bound state generated by the strong mean-field of hedgehog shape, it is legitimate to think that the valence level contribution above also contains some sort of vacuum polarization effect due to this strong mean field. In fact, the first moment of $\Delta u(x) + \Delta d(x)$, which can be interpreted as the quark spin content of the nucleon $\Delta \Sigma_3$ or the flavor-singlet axial charge at the energy scale of the present model, turns out to be fairly small as

$$\Delta \Sigma_3 \equiv \int_{-1}^{1} dx \, [\Delta u(x) + \Delta d(x)]$$

$$= \int_{0}^{1} dx \, [\Delta u(x) + \Delta \bar{u}(x) + \Delta d(x) + \Delta \bar{d}(x)] \simeq 0.35. \quad (11)$$

As advocated in [7], the symmetry breaking strong mean-field of hedgehog shape must be responsible for this smallness of the quark spin content. A quite interesting feature of the theoretical distribution $\Delta u(x) + \Delta d(x)$ is that it changes sign from positive to negative as $x$ becomes smaller than about 0.1. As already pointed out in our previous paper [12], what causes this sign change is the first order rotational correction, which arises from the proper account of the nonlocality (in time) of the quark bilinear operator entering into the theoretical definition of the quark distribution function. We shall later show that this sign change of the isoscalar distribution function in the small $x$ region is just what is required by the recent experimental data for the deuteron structure function $g_1^d(x)$.

The isovector combination $\Delta u(x) - \Delta d(x)$ illustrated in Fig.1(d) shows totally different $x$ dependence from the isoscalar one. The vacuum polarization contribution to this distribution function has a sizable magnitude of peak with positive sign around $x \simeq 0$. A sizable nonzero
value of $\Delta u(x) - \Delta d(x)$ in the negative $x$ domain has a far-reaching physical consequence. That is, it means the violation of SU(2) symmetry for the longitudinally polarized sea-quark (anti-quark) distributions. In consideration of the relation (11), we can read from Fig.1(d) that $\Delta\bar{u}(x) - \Delta\bar{d}(x) > 0$ for physical range of $x$ between 0 and 1.

To sum up, within the theoretical framework of the CQSM, the isoscalar and isovector distribution functions turn out to have fairly different $x$ dependence reflecting the fact that they have totally dissimilar $N_c$ dependence. Now we shall compare these predictions of the CQSM with the existing high energy data for the unpolarized as well as the longitudinally polarized structure functions. Before carrying out this comparison, we must first take account of the scale dependence of the distribution functions. We have done it by using the Fortran programs given by Saga group [27], which provide us with solutions of Dokshitzer-Gribov-Lipatov-Altarrelli-Parisi (DGLAP) evolution equation at the next-to-leading order (NLO). The question here is what value we should use for the initial energy scale of this $Q^2$ evolution. Here we have tried to see the effect of variation of $Q^2_{\text{init}}$ in a range centered around the model energy scale of $M^2_{PV} \simeq (0.56 \text{ GeV})^2 \simeq 0.31 \text{ GeV}^2$. We find that the final results are not significantly different in the range $0.25 \text{ GeV}^2 \leq Q^2_{\text{init}} \leq 0.35 \text{ GeV}^2$. We shall use the value $Q^2_{\text{init}} = 0.30 \text{ GeV}^2$ throughout the following analyses. In solving the DGLAP equation, we set $N_f = 3$ and assume $s(x) = \bar{s}(x) = 0$, $g(x) = 0$ and $\Delta s(x) = \Delta\bar{s}(x) = 0$, $\Delta g(x) = 0$ at the initial energy scale.

Figure 2: The predictions for $F^p_2(x) - F^n_2(x)$ and $F^p_2(x)/F^n_2(x)$ at $Q^2 = 4 \text{ GeV}^2$ are compared with the NMC data given at the corresponding energy scale [1].

Fig.2 shows the theoretical predictions for $F^p_2(x) - F^n_2(x)$ and $F^p_2(x)/F^n_2(x)$ at $Q^2 = 4 \text{ GeV}^2$ in comparison with the NMC data. One sees that the NMC data for $F^p_2(x) - F^n_2(x)$ are well
reproduced by the theory, except for a small discrepancy between the peak positions of the theoretical curve and the experimental data. Next, we compare our theoretical prediction for the ratio $F^n_2(x)/F^p_2(x)$ with the corresponding NMC data. One can say that the agreement between the theory and the experiment is fairly good at least for $x \leq 0.25$. The theoretical prediction deviates from the data as $x$ becomes larger than 0.3. This might be connected with the fact that our treatment of the soliton center-of-mass motion is essentially nonrelativistic. Note, however, that a reliable theoretical prediction of the ratio $F^n_2(x)/F^p_2(x)$ near $x \simeq 1$ is extremely difficult, because both of $F^n_2(x)$ and $F^p_2(x)$ damps rapidly as $x$ becomes larger.

Figure 3: The theoretical predictions for the unpolarized antiquark distribution $\bar{d}(x) - \bar{u}(x)$ at $Q^2 = 2.3$ GeV and $Q^2 = 54$ GeV in comparison with the HERMES [14] and the E866 data [15]. Here the open squares represent the HERMES data given at $Q^2 = 2.3$ GeV, while filled squares denote the E866 data given at $Q^2 = 54$ GeV.

Next, in Fig.3, we compare the theoretical predictions for the anti-quark distribution $\bar{d}(x) - \bar{u}(x)$ evaluated at two different values of $Q^2$ with the HERMES [14] and NuSea data [15] given at the corresponding energy scales. (One certainly confirms that the scale dependence of the quantity $\bar{d}(x) - \bar{u}(x)$ is rather small.) The general trend of the experimental data is well reproduced by the theory, although the theory has a tendency to slightly overestimate the magnitude of the flavor asymmetry of the sea-quark distributions.

Now we turn to the discussion of the longitudinally polarized distributions. Shown in Fig.4 are the longitudinally polarized structure functions for the proton, the neutron and the deuteron, $g^n_1(x), g^n_1(x)$ and $g^n_1(x)$. The results for $g^n_1(x)$ and $g^n_1(x)$ were already reported in
but we show them again for the purpose of comparison, since \( g_1^d(x) \) is essentially given as their average. To be more precise, we evaluate \( g_1^d(x, Q^2) \) in Fig. 4(c) by making use of the standardly-used formula

\[
g_1^d(x, Q^2) = \frac{1}{2} \left( g_1^p(x, Q^2) + g_1^n(x, Q^2) \right) (1 - 1.5 \omega_D),
\]

with \( \omega_D \) the probability that the deuteron is in a D-state.

Figure 4: The theoretical predictions for the longitudinally polarized structure functions for the proton, the neutron and the deuteron at \( Q^2 = 5 \text{ GeV}^2 \) in comparison with the corresponding experimental data. The filled circles in (a) and (b) respectively correspond to the E143 \[16\] and the E154 data \[17\], whereas the filled circles, the open circles and the open squares in (c) and (d) represent the E143, the E155 \[18\] and the SMC data \[19\].

As already argued in \[12\], a prominent feature of the CQSM is the good reproduction of the neutron data, which are known to have large magnitudes with negative sign. We have interpreted this as a manifestation of the chiral symmetry, maximally incorporated into this effective quark model. Now a direct comparison with the deuteron data reveals another
An interesting aspect of physics. An noticeable characteristic of the recent data for $g_1^d(x,Q^2)$ is that it appears to show sign change as $x$ approaches zero, although care must
be paid to the fact that the precision of the experimental data in the small $x$ region is still poor. Quite interestingly, the theoretical prediction of the CQSM closely follows the variation of this $g_1^d(x,Q^2)$ at least qualitatively. What is the origin of this behavior of the theoretical structure function? Undoubtedly, it can be traced back to the behavior of the isoscalar distribution function $\Delta u(x) + \Delta d(x)$ shown in Fig.1(c). In fact, as emphasized by Windmolders [28], the singlet distribution is strongly constrained by the measured value of $g_1^d$. (This seems only natural, since the deuteron is an isoscalar target.) To see it more closely, we recall the general expression for $g_1^d$ given as

$$g_1^d \sim \frac{1}{9} \left( \frac{1}{4} C_{NS} \otimes \Delta q_8 + C_S \otimes \Delta \Sigma + 2 N_f C_g \otimes \Delta g \right),$$  \hspace{1cm} (13)

where

$$\Delta q_8 = \Delta u + \Delta d - 2 \Delta s,$$  \hspace{1cm} (14)
$$\Delta \Sigma = \Delta u + \Delta d + \Delta s,$$  \hspace{1cm} (15)

and $C_{NS}, C_S, C_G$ are the non-singlet, singlet and gluon Wilson coefficients and the symbol $\otimes$ represents convolution with respect to $x$.

Assuming minor role of the $s$-quark degrees of freedom at the energy scale in question, the difference between $\Delta q_8$ and $\Delta \Sigma$ would be small. The $g_1^d$ is then determined by $\Delta u + \Delta d$ and $\Delta g$ at this energy scale. Now, somewhat peculiar small-$x$ behavior of the isoscalar longitudinally polarized distribution function predicted by the CQSM seems to get a phenomenological support of the latest deuteron data. In any case, the $g_1^d$, especially in a smaller $x$ region, appears to be extremely sensitive to the detailed dynamical content of the theories. To be able to reproduce it or not would then be a good touchstone of model selection.

Now that we have confirmed that the CQSM can explain the principle experimental data for the unpolarized and the longitudinally polarized structure functions of the nucleon as well as the deuteron, we come back to an interesting prediction of the CQSM given by Fig.1(d), i.e. the possible violation of the light flavor sea-quark asymmetry for the longitudinally polarized distributions. Since the sea-quark distributions are not well constrained by the inclusive data alone, very little is known as to them. In a recent report [29], Morii and Yamanishi proposed new formulas for extracting a difference, $\Delta \bar{d} - \Delta \bar{u}$, from data of the polarized semi-inclusive processes $\vec{l} + \vec{N} \rightarrow l' + H + X$ and estimated the value of it from the available data of the SMC and Hermes groups. The results of their analyses are shown in Fig.5 together with the prediction of the CQSM at the corresponding energy scale, $Q^2 = 4 \text{ GeV}^2$. Unfortunately, the uncertainties of the fit are so large that nothing definite can be said at the present status.
Figure 5: The theoretical prediction for the longitudinally polarized antiquark distribution \( \Delta \bar{d}(x) - \Delta \bar{u}(x) \) at \( Q^2 = 4 \text{ GeV}^2 \) is compared with the recent semi-phenomenological fit by Morii and Yamanishi. See figure caption of [29] for the detailed meaning of three different marks in their results.

Note, however, that Morii and Yamanishi parameterized \( \Delta \bar{d}(x) - \Delta \bar{u}(x) \) as

\[
\Delta \bar{d}(x) - \Delta \bar{u}(x) = C \cdot x^\alpha \left( \bar{d}(x) - \bar{u}(x) \right)
\]

and determined the value of \( C \) and \( \alpha \) from the \( \chi^2 \)-fit of the results presented in the above figure. The results were \( C = -1.00 \) and \( \alpha = 0.18 \), thereby indicating an asymmetry of \( \bar{d} \) and \( \bar{u} \). The negative value of \( C \) denotes that

\[
\Delta \bar{d}(x) - \Delta \bar{u}(x) < 0,
\]

since we already know that \( \bar{d}(x) - \bar{u}(x) > 0 \). This is qualitatively consistent with the prediction of the CQSM. In fact, the values of \( C \) and \( \alpha \) determined from our theoretical distributions at \( Q^2 = 4 \text{ GeV}^2 \) turn out to be

\[
C \simeq -2.0, \quad \alpha \simeq 0.12,
\]

in qualitative consistency with the result of their analysis.

Despite encouraging phenomenological support to the present approach, an important question remains to be answered. The question concerns the role of gluons at the relatively low energy scale of \( Q^2 = 0.3 \sim 0.5 \text{ GeV}^2 \), which may be taken as a starting energy of the DGLAP evolution equation. We have simply assumed zero for the input gluon densities at the initial
energy scale. This assumption may not be necessarily justified, especially for the unpolarized case. In fact, there is some phenomenological indication that the gluon saturates nearly 30% of the nucleon momentum sum rule even at such low energy scale \[30\]. (Note, however, that the isovector-type quantities like \(F_2^p(x) - F_2^n(x)\) and \(\bar{d}(x) - \bar{u}(x)\) investigated in the present paper would be insensitive to these input gluon densities.) On the other hand, very little is known for the polarized gluon densities. A relatively good agreement between the experimental structure functions and our theoretical predictions, which are obtained by assuming zero gluon polarization at \(Q^2_{\text{init}} = 0.3\) GeV\(^2\), indicates that the role of gluons may be less important than the case of the unpolarized distributions. (This never denies the importance of gluon polarization at high energy scale, since the gluons are known to acquire polarization rapidly through the processes of perturbative evolution.) Undoubtedly, it is a very important theoretical question to reliably predict the gluon densities to be used as input distributions of the renormalization-group evolution. The importance of the gluon densities at this matching energy scale of \(Q^2 = 0.3 \sim 0.5\) GeV\(^2\) was also emphasized by by Lampe and Reya in a recent review \[31\]. According to these authors, if low energy models cannot provide necessary input gluon densities at this energy, they would only refer to some nonperturbative input quark scale which cannot be reached by perturbative evolution.

In summary, the CQSM allows us for the calculation of the full \(x\)-dependence of the nucleon structure functions without introducing any new free parameters. In particular, what makes this approach quite promising is the fact that it predicts, besides the valence-like distributions, the sea-quark-like densities concentrated in the small \(x\) domain, if it is treated in a theoretically consistent manner (i.e., not in the “valence-quark-only” approximation). This unique feature of the model is expected to provide us with a good starting point for theoretically understanding unpolarized data as well as yet-to-be-obtained polarized data in the small \(x\) region. Making full use of this advantage of the CQSM, we have made a prediction for the flavor asymmetry of the longitudinally polarized antiquark distributions in the proton. The model has been shown to predict \(\Delta \bar{d}(x) - \Delta \bar{u}(x) < 0\), while at the same time \(\bar{d}(x) - \bar{u}(x) > 0\), i.e., it definitely predicts light flavor sea-quark asymmetry not only for the spin-averaged distributions but also for the longitudinally polarized distributions. We hope that this interesting prediction of the CQSM will be tested through accumulation and analyses of more precise experimental data on the polarized semi-inclusive processes in the near future.

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(a) $u(x) + d(x)$
(b) $u(x) - d(x)$
(c) $\Delta u(x) + \Delta d(x)$
(d) $\Delta u(x) - \Delta d(x)$
\( F_2^p(x) - F_2^n(x) \)
\( Q^2 = 4\text{GeV}^2 \)

\( F_2^n(x) / F_2^p(x) \)
\( Q^2 = 4\text{GeV}^2 \)
\[ Q^2 = 2.3 \text{GeV}^2 \]

\[ Q^2 = 54.0 \text{GeV}^2 \]
\( g_1^p(x) \)
\( Q^2 = 5\text{GeV}^2 \)

\( g_1^n(x) \)
\( Q^2 = 5\text{GeV}^2 \)

\( g_1^d(x) \)
\( Q^2 = 5\text{GeV}^2 \)

\( xg_1^d(x) \)
\( Q^2 = 5\text{GeV}^2 \)
\[ \Delta \bar{d}(x) - \Delta \bar{u}(x) \]

\[ Q^2 = 4 \text{GeV}^2 \]