A Chiral Theory of Strange Sea Distributions in the Nucleon

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Theoretical predictions are given for the strange sea distributions in the nucleon on the basis of the flavor SU(3) chiral quark soliton model, with emphasis upon the asymmetry of quark and antiquark distributions. We find that the quark-antiquark asymmetry of the strange sea is much larger for longitudinally polarized distribution functions than for unpolarized ones. A preliminary comparison with the CCFR data for the unpolarized $s$-quark distribution and with the LSS fits of the longitudinally polarized distribution functions is encouraging.

§1. Introduction

A unique feature of the chiral quark soliton model (CQSM) as compared with many other effective models of QCD, like the MIT bag model, is that it can give reasonable predictions not only for the quark distributions but also for the antiquark distributions.1), 2) This is due to the field theoretical nature of the model that enables us to carry out a nonperturbative evaluation of the parton distribution functions with full inclusion of the vacuum polarization effects in the rotating mean field of hedgehog form.1), 3) It has already been shown that, without introducing any adjustable parameter, except for the initial-energy scale of the $Q^2$-evolution, the CQSM can describe almost all the qualitatively prominent features of the recent high-energy deep-inelastic scattering observables. It naturally explains the excess of the $d$-sea over the $u$-sea in the proton.4), 5) It also reproduces the qualitative behavior of the observed longitudinally polarized structure functions for the proton, the neutron and the deuteron.1), 2) The most puzzling observation, i.e. the unexpectedly small quark spin fraction of the nucleon, can also be explained with no need of a large gluon polarization at the low renormalization scale.6), 7) Finally, the model predicts a quite large isospin-asymmetry also for the spin-dependent sea-quark distributions,1)−3), 8) which we expect will be confirmed by future experiments. (The asymmetry between the $u$-sea and $d$-sea, which we discuss in this paper, should not be confused with the breakdown of the $SU(2)$ isospin symmetry, which relates distributions in the proton and the neutron.)

Our theoretical analyses to this time have been based on the flavor $SU(2)$ CQSM, so that no account has been taken of the possibility of strange quark excitations in the nucleon. However, there have been several experimental indications that $s\bar{s}$ pairs in the nucleon are responsible for a number of non-trivial effects.9) An interesting question here is how large the magnitude of this admixture is and/or how large the quark-antiquark asymmetry of the nucleon strange sea distributions

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is. Also interesting is whether we do expect asymmetry of $s$- and $\bar{s}$-quarks also for the spin-dependent distributions. The purpose of the present paper is to answer these questions, based on the CQSM generalized to the case of flavor $SU(3)$.\textsuperscript{10,11} We emphasized that the field theoretical nature of the model, which is taken from the $SU(2)$ model, plays a decisively important role, since the strange quark distributions, which are our main concern here, have a completely non-valence character.

\section*{§2. Brief description of the model}

Since the flavor $SU(3)$ CQSM is constructed on the basis of the flavor $SU(2)$ model, we first recall some basics of the $SU(2)$ CQSM. It is specified by the effective Lagrangian,

$$L_0 = \bar{\psi}(i\not\partial - Me^{i\gamma_5 \cdot \pi(x)/f_\pi})\psi,$$

which describes the effective quark fields with a dynamically generated mass $M$, interacting with massless pions. The nucleon (or $\Delta$) in this model appears as a rotational state of a symmetry-breaking hedgehog object, which itself is obtained as a solution of the self-consistent Hartree problem with infinitely many Dirac-sea quarks.\textsuperscript{12,6} The theory is not renormalizable, and it is defined with an ultraviolet cutoff. In the Pauli-Villars regularization scheme, which is used throughout the present analysis, that which plays the role of the ultraviolet cutoff is the Pauli-Villars mass $M_{PV}$ obeying the relation $(N_cM^2/4\pi^2) \ln (M_{PV}/M)^2 = f_\pi^2$ with $f_\pi$ the pion weak decay constant.\textsuperscript{3} Using the value $M \simeq 375$ MeV, which is obtained from the phenomenology of nucleon low energy observables, this relation fixes the Pauli-Villars mass as $M_{PV} \simeq 562$ MeV. (Here, the regularization is introduced only to the real part of the Euclidean action, not to the imaginary part.) Since we are to use these values of $M$ and $M_{PV}$, there is no free parameter additionally introduced into the calculation of distribution functions.\textsuperscript{2}

Now, the principle dynamical assumption of the $SU(3)$ CQSM is as follows. The first is the embedding of the $SU(2)$ self-consistent mean-field (of hedgehog shape) into the $SU(3)$ matrix as

$$U_0^{\gamma_5}(x) = \begin{pmatrix} e^{i\gamma_5 \cdot \hat{r}} F(r) & 0 \\ 0 & 1 \end{pmatrix},$$

which is believed to give the lowest energy classical configuration, as can be deduced from a simple variational argument. In fact, an arbitrary slight variation of the $(3,3)$ component of $U_0^{\gamma_5}(x)$ induces changes of the strange-quark single-particle spectra in such a way that weak bound states appear from the positive energy Dirac continuum as well as from the negative energy one in a charge-conjugation symmetric way. Since the negative energy continuum is originally occupied, this necessarily increases the total energy of the system of baryon number one.

The next assumption is the semiclassical quantization of the rotational motion in the $SU(3)$ collective space represented as

$$U_0^{\gamma_5}(x,t) = A(t) U_0^{\gamma_5}(x) A^\dagger(t),$$

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with
\[ A(t) = e^{-i\Omega t}, \quad \Omega = \frac{1}{2} \lambda^a \Omega_a \in SU(3). \] (2.4)

The semiclassical quantization of this collective rotation leads to a systematic method for calculation of any nucleon observables, including the parton distribution functions, which is given as 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 order \( 1/N_c \) quantity.) In the present study, all the terms up to first order in \( \Omega \) are consistently taken into account, according to the formalism explained in Ref. 1. Since the resultant expressions for the quark (and/or antiquark) distribution functions are quite lengthy, we present them in detail elsewhere and demonstrate only the main results here. Several comments are in order, however. The \( SU(3) \) symmetry breaking effects arising from the effective mass difference between the strange and nonstrange quarks (which takes the form of an additional parameter of the model) can in principle be taken into account by using a perturbation method. (The perturbative treatment of the \( SU(3) \) symmetry breaking term would be justified (though not completely), since the effective mass difference between the strange and nonstrange quarks of order 150 MeV is much smaller than the typical energy scale of the model, which may roughly be specified by the Pauli-Villars cutoff mass around 600 MeV.) To carry it out for parton distribution functions taking full account of the vacuum polarization effects is quite involved, however. We therefore leave this to future studies. (Probabilistically, the magnitude of the strange quark mixture under this approximation should be taken as the upper limit.) Secondly, some inconsistency is known to exist between the basic dynamical assumption of the \( SU(3) \) CQSM and the time-order-preserving collective quantization procedure of the rotational motion, although the latter is believed to resolve the long-standing \( g_A \) problem in the \( SU(2) \) model.\textsuperscript{13,14} Here, we simply follow the symmetry preserving approach advocated in Ref. 15, which amounts to dropping some theoretically contradictory terms by hand.

§3. Numerical results

We display in Fig. 1 the theoretical \( s \)- and \( \bar{s} \)-quark distribution functions evaluated at the model energy scale. Here, (a) represents the unpolarized distributions, and (b) represents the longitudinally polarized distributions. It is confirmed that both \( s(x) \) and \( \bar{s}(x) \) satisfy the positivity constraint, as they should, in sharp contrast to the previous result obtained by the Tübingen group in the so-called “valence-quark-only” approximation.\textsuperscript{16} This proves our assertion that proper account of the vacuum polarization effects is vital to give a reliable prediction for sea-quark distributions. It is also noticed that the distributions \( \bar{s}(x) \) have softer (lower-\( x \)) components than \( s(x) \), qualitatively consistent with the argument of Brodsky and Ma based on the light-cone meson-baryon fluctuation model.\textsuperscript{17} Note, however, that the asymmetry cannot be extremely large, due to the restriction of the conservation of the strange quantum number in the nucleon, i.e. \( \int_0^1 s(x) \, dx = \int_0^1 \bar{s}(x) \, dx \), with the positivity constraints \( s(x) > 0 \) and \( \bar{s}(x) > 0 \).

In contrast, a rather large quark-antiquark asymmetry is observed for the lon-
gitudinally polarized distributions. It is seen that the $s$- and $\bar{s}$-quarks are both negatively polarized, but $|\Delta s(x)|$ is much smaller than $|\Delta \bar{s}(x)|$. This feature too is consistent with the conjecture of Brodsky and Ma, at least qualitatively. In fact, they argue that, if the intrinsic strange fluctuations in the proton are mainly due to the intermediate $K^+\Lambda$ configuration, the $s$-quark is negatively polarized, but the polarization of $\bar{s}$ is zero. Their argument goes as follows. Since the $K^+$ meson is a pseudoscalar particle with negative parity, and the parity of $\Lambda$ is positive, parity conservation dictates that the relative orbital angular momentum of the intermediate $K^+\Lambda$ state must be odd, and most probably is a $p$-wave state. This gives the total angular momentum wave function in the following form:

$$
|K^+\Lambda \left( J = \frac{1}{2}, J_z = \frac{1}{2} \right) \rangle = \sqrt{\frac{2}{3}} |L = 1, L_z = 1 \rangle |\Lambda \left( S = \frac{1}{2}, S_z = -\frac{1}{2} \right) \rangle
- \sqrt{\frac{1}{3}} |L = 1, L_z = 0 \rangle |\Lambda \left( S = \frac{1}{2}, S_z = \frac{1}{2} \right) \rangle.
$$

(3.1)

The point here is that the probability of the $\Lambda$-spin being opposite to the proton spin is twice as large as that of it being parallel. Combining this with the observation that the spin of $\Lambda$ almost completely comes from its constituent $s$-quark, we immediately conclude that the virtually mixed $s$-quark in the proton is negatively polarized with respect to the proton spin direction. The situation is quite different for the $\bar{s}$-quark generated through the same intrinsic fluctuation $p \rightarrow K^+\Lambda$. Since $\bar{s}$ is contained in the pseudoscalar meson $K^+$ without spin, the net spin of $\bar{s}$ in $K^+$, and consequently in the proton, would be zero. Although it is qualitatively consistent with this argument of Brodsky and Ma, the CQSM predicts a sizable amount of negative polarization also for the $\bar{s}$-quark. Such a nonzero polarization of the $\bar{s}$-quark may be obtained by introducing more complicated virtual process like $p \rightarrow K^{*+}\Lambda$. However,
the precise estimation of the size of polarization in such meson cloud models would
be quite difficult, since there are many competing processes.

Just by considering intermediate $p\pi^0$ and $n\pi^+$ configurations instead of the
$K^+\Lambda$ fluctuation, the meson-baryon fluctuation model (or the meson cloud convolu-
tion model) can naturally explain the excess of the $\bar{d}$-sea over the $\bar{u}$-sea in the
proton.\textsuperscript{18) Note, however, that for the same reason that the net polarization of $\bar{s}$ is
zero, one must conclude that the net polarizations of $\bar{d}$- and $\bar{u}$-seas are zero (or at
least very small). This clearly contradicts the previously-mentioned predictions of
the $SU(2)$ CQSM, in which $\Delta\bar{u}(x)$ is large and positive, while $\Delta\bar{d}(x)$ is large
and negative.\textsuperscript{2) In our opinion, that which is responsible for this remarkable difference
is the nontrivial correlation between spin and isospin quantum numbers embedded
in the CQSM. At least, one should recognize that the physical content of the pion
cloud model and the CQSM is not necessarily the same, as naively believed.

There are two common ways to extract unpolarized strange sea distributions
from the deep-inelastic-scattering data. The first method uses neutrino-induced
charm production, while the second relies upon a global fit (as in the case of the
other flavor densities). The first direct determination of the strange quark distri-
bution based on the neutrino-induced charm productions was carried out by the
CCFR collaboration some years ago.\textsuperscript{19) Here, we carry out a very preliminary com-
parison of the theoretical predictions of the $SU(3)$ CQSM with the strange quark
distribution obtained with the CCFR next-to-leading-order (NLO) analysis. This
comparison should be regarded as preliminary for several reasons. First, the hidden
strangeness excitation in the nucleon is thought to be very sensitive to the inclusion
of the $SU(3)$ breaking effects due to the mass difference between the strange and non-
strange quarks which we have not yet included.\textsuperscript{16) Second, to be more strict, the dis-
tribution functions evaluated within the model should be taken as constituent quark
distributions rather than current quark distributions. For lack of precise knowledge
relating these two distributions, a common approach is to compare the predictions
of the effective models directly with the experimental observables, like the structure
functions. Still, we believe that the comparison at the distribution function level is
not completely meaningless, since our main concern here is to understand qualita-
tively prominent features of the quark and antiquark distributions, especially those
of the strange seas. To carry out the comparison, we have taken account of the scale
dependence of the distribution functions, using the Fortran code for NLO evolution
provided by the Saga group.\textsuperscript{20) The initial energy scale of this evolution is taken to
be $Q_{\text{ini}}^2 = 0.25\text{GeV}^2$ and the gluon distribution at this scale is simply set to zero,
although this may not be completely justified. In Fig. 2, we plot the theoretical
distributions $s(x)$ and $\bar{s}(x)$ together with the result of the CCFR NLO analysis at
$Q^2 = 4\text{GeV}^2$ with the constraint $s(x) = \bar{s}(x)$. Considering that yet-to-be-included
$SU(3)$ breaking effects is expected to suppress the magnitude of strange quark ex-
citations, it can be said that the theory reproduces the order of magnitude of the
observed strange sea distribution.

Turning to the spin-dependent distribution functions, the quality of the presently
existing semi-inclusive data is rather poor, so that the analyses is mainly lim-
ited to the inclusive DIS data alone. This forces us to introduce several sim-
the results of LSS analyses. and $\Delta \bar{s}$

\[
\Delta s(x) = \Delta \bar{d}(x) = \Delta \bar{s}(x),
\]

In the other analyses, it has been assumed that
\[
\Delta q_3(x, Q^2) = c \Delta q_8(x, Q^2)
\]
with $c$ a constant. Probably, the most ambitious analyses free from these ad hoc assumptions on the distribution functions are those of Leader, Sidrov and Stamenov (LSS). They also investigated the sensitivity of their analysis to the size of the $SU(3)$ symmetry breaking effect. (Although they did not take account of the possibility that
\[
\Delta s(x) \neq \Delta \bar{s}(x),
\]

this simplification is harmless, because only the combination $\Delta s(x) + \Delta \bar{s}(x)$ appears in their analysis of inclusive DIS data.)

To compare the theoretical distributions of the $SU(3)$ CQSM with the LSS fits given at $Q^2 = 1\text{ GeV}^2$, we must consider the fact that their analyses are carried out in the so-called JET scheme (or the chirally invariant scheme). To take account of this, we start with the theoretical distribution functions $\Delta u(x)$, $\Delta \bar{u}(x)$, $\Delta d(x)$, $\Delta \bar{d}(x)$, $\Delta s(x)$, $\Delta \bar{s}(x)$, which are taken as the initial distribution functions given at $Q^2_{\text{ini}} = 0.25\text{ GeV}^2$. Under the assumption that $\Delta g(x) = 0$ at this initial energy scale, we solve the DGLAP equation in the standard $\overline{MS}$ scheme to obtain the distributions at $Q^2 = 1\text{ GeV}^2$. The corresponding distribution functions in the JET scheme are then obtained by using the transformation

\[
\Delta \Sigma(x, Q^2)_{\text{JET}} = \Delta \Sigma(x, Q^2)_{\overline{MS}} + \frac{\alpha_s(Q^2)}{\pi} N_f (1 - x) \otimes \Delta g(x, Q^2)_{\overline{MS}},
\]
\[
\Delta g(x, Q^2)_{\text{JET}} = \Delta g(x, Q^2)_{\overline{MS}},
\]

with
\[
\Delta \Sigma(x, Q^2) = \sum_{i=1}^{N_f} (\Delta q_i(x, Q^2) + \Delta \bar{q}_i(x, Q^2)).
\]

The solid curves in Fig. 3 represent the theoretical distributions $x(\Delta u(x) + \Delta \bar{u}(x))$, $x(\Delta d(x) + \Delta \bar{d}(x))$, $x(\Delta s(x) + \Delta \bar{s}(x))$ and $x\Delta g(x)$ at $Q^2 = 1\text{ GeV}^2$ in comparison with the corresponding LSS fits. The long-dashed, dotted and dash-dotted curves in (c) and (d) are their fits, respectively, obtained by imposing the constraint on the value of the axial charge $a_8$ to be 0.58 [the $SU(3)$ limit], 0.86 and 0.40. [Only the case of $a_8 = 0.58$ is shown in (a) and (b), since these distributions are insensitive to the variation of $a_8$.] We see that the distributions $\Delta s(x) + \Delta \bar{s}(x)$ as well as $\Delta g(x)$ are fairly sensitive to the effects of the $SU(3)$ symmetry breaking, and their magnitude cannot be determined with good precision from inclusive DIS data alone. Taking into account the large uncertainties in the magnitudes of $x(\Delta s(x) + \Delta \bar{s}(x))$ and $x\Delta g(x)$, the predictions of the $SU(3)$ CQSM are qualitatively consistent with the results of LSS analyses.
A Chiral Theory of Strange Sea Distributions in the Nucleon

As emphasized above, a noteworthy feature of the $SU(2)$ CQSM is that it predicts a quite large $\bar{u} - \bar{d}$ asymmetry not only for the spin-independent sea-quark distributions but also for the longitudinally polarized one. Why is this observation so important? This is because the NMC observation $\bar{d}(x) - \bar{u}(x) > 0$ in the proton can be described equally well by the CQSM and the naive meson cloud convolution model, whereas the latter essentially predicts $\Delta \bar{u}(x) \simeq \Delta \bar{d}(x) \simeq 0$, in contrast to the prediction of the CQSM, $\Delta \bar{u}(x) > 0 > \Delta \bar{d}(x), |\Delta \bar{u}(x) - \Delta \bar{d}(x)| \simeq |\bar{u}(x) - \bar{d}(x)|$. Now the question is whether this feature of the $SU(2)$ CQSM is shared by the $SU(3)$ CQSM.

In Fig. 4, we compare the predictions of the $SU(2)$ and $SU(3)$ CQSM for $x(\bar{d}(x) - \bar{u}(x))$ and $x(\Delta \bar{u}(x) - \Delta \bar{d}(x))$ with the corresponding predictions obtained by Bhalerao based on what he calls the “statistical model”. Note that his model is semi-phenomenological and uses experimental information as inputs. It is seen that both the predictions of the $SU(2)$ and $SU(3)$ CQSM for $x(\bar{d}(x) - \bar{u}(x))$ are fairly close to Bhalerao’s prediction, which is consistent with the NMC data. However, the magnitude of $x(\Delta \bar{u}(x) - \Delta \bar{d}(x))$ is much smaller in the $SU(3)$ model than in
Fig. 4. The theoretical predictions of the SU(2) and SU(3) CQSM for (a) $x(\bar{d}(x) - \bar{u}(x))$ and (b) $x(\Delta \bar{u}(x) - \Delta \bar{d}(x))$ at $Q^2 = 0.88\text{GeV}^2$, in comparison with Bhalerao’s semi-theoretical predictions.

the SU(2) model. This reduction is due to a delicate cancellation of several terms in this more complicated theory. Unfortunately, we do not yet have any clear physical explanation of this result. However, a possible explanation may be given as follows. As emphasized above, the large isospin asymmetry of both spin-averaged and spin-dependent sea quark distributions has something to do with the strong spin-isospin correlation embedded in the SU(2) CQSM. It may happen that this strong spin-isospin correlation is partially destroyed by the SU(3) symmetric collective quantization introduced in the SU(3) CQSM, which has the effect of generating clouds of kaons in addition to those of pions. Since the SU(3) symmetric collective quantization has a tendency to overestimate the magnitudes for the kaon clouds, we conjecture that the introduction of the SU(3) symmetry breaking effect in the SU(3) model would partially decrease its prediction for $x(\Delta \bar{u}(x) - \Delta \bar{d}(x))$ toward that of the SU(2) model, thereby leading to a result, that is not extremely far from Bhalerao’s prediction.

§4. Summary and conclusion

In summary, we have given theoretical predictions with no free parameters for the strange sea distributions in the nucleon on the basis of the flavor SU(3) CQSM. It has been shown that the $s$- and $\bar{s}$-quarks are both negatively polarized, but the magnitude of $\Delta s(x)$ is much larger than that of $\Delta \bar{s}(x)$, while the quark-antiquark asymmetry of the unpolarized strange sea is not extremely large, because of the conservation of the strange quantum number in the nucleon. A preliminary comparison with the CCFR data for the unpolarized $s$-quark distribution is encouraging. The theory also reproduces the characteristic features of the recent LSS fits of the longitudinally
polarized distribution functions, including the negatively polarized strange sea. We also emphasize that the $SU(2)$ symmetry of the polarized non-strange seas is likely to be significantly violated, such that $\Delta\bar{u}(x) > 0 > \Delta\bar{d}(x)$. At any rate, our analysis clearly shows that the spin and flavor dependences of the antiquark distributions in the nucleon are very sensitive to the nonperturbative dynamics of QCD at low energy. To reveal this interesting aspect of QCD, it is necessary to carry out various types of high-energy DIS experiments, which should enable us to carry out flavor and valence plus sea quark decompositions of the parton distribution functions.

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