Non-perturbative structure of the polarized nucleon sea

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Abstract

We investigate the flavour and quark-antiquark structure of the polarized nucleon by calculating the parton distribution functions of the nucleon sea using the meson cloud model. We find that the SU(2) flavor symmetry in the light antiquark sea and quark-antiquark symmetry in the strange quark sea are broken, i.e. $\Delta \bar{u} < \Delta \bar{d}$ and $\Delta s < \Delta \bar{s}$. The polarization of the strange sea is found to be positive, which is in contradiction to previous analyses. We predict a much larger quark-antiquark asymmetry in the polarized strange quark sea than that in the unpolarized strange quark sea. Our results for both polarized light quark sea and polarized strange quark sea are consistent with the recent HERMES data.

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I. INTRODUCTION

Since the famous EMC experiment [1] at CERN in 1988 there has been great interest in the polarized quark distributions of the proton and neutron. The reason for this interest in the spin dependent structure of the nucleon is that the EMC experiment (and other polarized DIS experiments [2]) could be interpreted as showing that the quarks carry only a small proportion of the total angular momentum of the proton. A further conclusion, made using SU(3) flavour arguments, is that the strange sea quarks in the proton are strongly polarized opposite to the polarization of the proton [3]. Both these results are much at odds with expectations based on constituent quark models of the nucleon. Further experiments have generally confirmed the EMC results for the proton - photon asymmetry \( A_1 \) and proton spin structure function \( g_1(x) \), as well as extending these to deuteron and neutron targets [4] and to the second nucleon - photon asymmetries \( A_2^N \) and the related structure functions \( g_2^N(x) \) [5].

All these experiments have been inclusive scattering using the virtual photon as the probe of nucleon structure. Unfortunately no experiments have been performed using neutrino scattering, where one of the vector bosons is the probe. Thus there is no information on the spin dependent analogue of the unpolarized structure function \( F_3(x) \). This makes the decomposition of the measured structure functions into ‘valence’ and ‘sea’ parts very difficult, and flavour decomposition much harder to do than in the unpolarized case. However the HERMES collaboration [6] has recently reported the first results from its semi-inclusive scattering measurements, where final state pions and kaons are measured. This provides a method to extract the spin dependent quark distribution of a given flavour, supposing that we have enough information about the relevant fragmentation functions \( D_{qh}^h(Q^2, z) \). The HERMES collaboration’s first semi-inclusive results have been able to extract the ratios of spin dependent to spin independent quark distributions (\( \Delta u, \Delta \bar{u}, \Delta d, \Delta \bar{d}, \Delta s, \Delta \bar{s} \), \( 0 \pm \frac{1}{3} \)). In contrast to earlier flavour decompositions [7, 8], which have needed SU(3) flavour symmetry assumptions for the sea distributions \( \Delta \bar{u} = \Delta \bar{d} = \Delta s = \Delta \bar{s} \), the only symmetry assumption made in the HERMES data analysis is that \( \frac{\Delta \bar{s}(x)}{\Delta \bar{s}(x)} = 0 \). In the context of the earlier results from inclusive DIS, the results for the spin dependent sea quark distributions \( \Delta \bar{u}(x), \Delta \bar{d}(x), \Delta s(x) \) are rather surprising:
the polarization of each flavour is very small, and compatible with zero. This may imply a breaking of SU(3) flavour symmetry.

With this new data in mind, it is important that other sources of information on the spin dependent sea quark distributions be examined. In particular model calculations of the parton distributions can provide some insight into whether the antiquark polarizations can be expected to be large or small, and whether the sign of the polarization is positive or negative. In this paper we will investigate the spin dependent sea quark distribution functions ($\Delta \bar{u}, \Delta \bar{d}, \Delta s, \Delta \bar{s}$) in the context of the meson cloud model (MCM). In this model the physical nucleon wavefunction contains virtual meson - baryon components which ‘dress’ the bare nucleon. The MCM provides a natural explanation for symmetry breaking among the parton distributions, in particular the flavour asymmetry in the unpolarized nucleon sea ($\bar{d}(x) > \bar{u}(x)$) [9] seen in the NMC [10] and E866 [11] experiments. Thus it is reasonable to ask whether the violation of the Ellis-Jaffe sum rule [12] can also be explained in this model. In previous calculations [13] it has been shown that including effects of the meson cloud significantly lowers the value of $\Delta \Sigma = \sum_f (\Delta Q_f + \Delta \bar{Q}_f)$, which is the total spin carried by the quarks and anti-quarks. This arises firstly because of ‘dilution’ of the bare proton quark distributions by those of the cloud and secondly from the inclusion of non s-wave components in the cloud wavefunction, which effectively increases the proportion of the proton spin due to orbital angular momentum. Boros and Thomas [13] studied the effects of strange mesons and baryons in the cloud by considering $\Lambda K$ and $\Sigma K$ components in the Fock wavefunction of the proton and found that the polarization of the strange sea was small ($\Delta S + \Delta \bar{S} < 0.01$). However higher mass components in the Fock wavefunction, in particular $\Lambda K^*$ and $\Sigma K^*$, can have important effects on the strange see. Our recent investigation of the unpolarized strange sea [14] has shown that although these higher mass components may be kinematically suppressed, they have large couplings to the nucleon, and numerically give amplitudes of similar size to the lowest mass states. Also the $K^*$ is a vector meson, so any polarization of the anti-strange quark in the $K^*$ will give a contribution to the $\Delta \bar{s}(x)$ distribution of the proton, whereas there can be no such contribution from the $\bar{s}$ of the Kaon.

In a previous paper [15] we have calculated the flavour asymmetry $\Delta \bar{u}(x) - \Delta \bar{d}(x)$ in the MCM. These results are consistent with the HERMES data. Here we shall extend
that calculation to consider each of the light antiquark flavours separately. Thus the fluctuations that contribute to $\Delta \bar{u}(x)$ and $\Delta \bar{d}(x)$ equally and make no contributions to $\Delta \bar{u}(x) - \Delta \bar{d}(x)$ are now included.

In addition to the MCM contribution to the light polarised antiquark distributions there may be other contributions to these distributions. These would be mainly non-perturbative, arising from the distributions in the bare nucleon, and would need to be calculated in some model of the bare nucleon e.g. the MIT bag model. These contributions are not necessarily small, however we shall see that the HERMES data does constrain the size of these bare distributions. Perturbative contributions to the light polarised antiquark distributions, which come from QCD evolution from the model scale ($Q_0^2 < 1 \text{ GeV}^2$) up to experimental scales, are expected to be small as the first moment of these distributions changes rather slowly with $Q^2$.

In the Section 2 of this paper we review the meson cloud model, and derive the necessary expressions for calculating the polarised antiquark distributions. This will include the contributions of $\Lambda K^*$ and $\Sigma K^*$ components of the cloud. In Section 3 we present a non-perturbative calculation for the polarized parton distributions (PDFs) of the hyperons and mesons. The numerical results and discussions are given in Section 4. The last section is reserved for a summary.

II. POLARIZED NUCLEON SEA IN THE MESON CLOUD MODEL

In the meson cloud model (MCM) [16] the nucleon can be viewed as a bare nucleon plus some baryon-meson Fock states which result from the fluctuation of nucleon to baryon plus meson $N \rightarrow BM$. The wavefunction of the nucleon can be written as [17],

$$|N\rangle_{\text{physical}} = Z |N\rangle_{\text{bare}} + \sum_{BM} \sum_{\lambda \lambda'} \int dy \, d^2k_\perp \phi_{BM}^{\lambda \lambda'}(y, k_\perp^2) |B^\lambda(y, k_\perp) M^{\lambda'}(1 - y, -k_\perp)\rangle$$

where $Z$ is the wave function renormalization constant, $\phi_{BM}^{\lambda \lambda'}(y, k_\perp^2)$ is the wave function of the Fock state containing a baryon ($B$) with longitudinal momentum fraction $y$, transverse momentum $k_\perp$, and helicity $\lambda$, and a meson ($M$) with momentum fraction $1 - y$, transverse momentum $-k_\perp$, and helicity $\lambda'$. The model assumes that the lifetime of a virtual baryon-meson Fock state is much longer than the interaction time in the deep inelastic
or Drell-Yan process, thus the quark and anti-quark in the virtual baryon-meson Fock states can contribute to the parton distributions of the nucleon. These non-perturbative contributions can be calculated via a convolution between the fluctuation function, which describes the microscopic process $N \rightarrow BM$, and the quark (anti-quark) distributions of the hadrons in the Fock states $|BM\rangle$ in Eq. (1). We consider only the valence quarks of the baryon-meson pair as the small $x$ region, where sea quarks may be important, is kinematically suppressed (see the discussions of the fluctuation functions below).

The contributions to the helicity-dependent parton distribution of the nucleon are

\[ x\delta q^a = \sum _\lambda \int _x^1 dy f^\lambda _{BM/N}(y) \frac{x}{y} q^a _{B,\lambda}(\frac{x}{y}) + \sum _\lambda \int _x^1 dy f^\lambda _{MB/N}(y) \frac{x}{y} q^a _{M,\lambda}(\frac{x}{y}) \]

\[ + \sum _\lambda \int _x^1 dy f^\lambda _{(B,B_2)M/N}(y) \frac{x}{y} q^a _{(B,B_2),\lambda}(\frac{x}{y}) \]

\[ + \sum _\lambda \int _x^1 dy f^\lambda _{(M,M_2)B/N}(y) \frac{x}{y} q^a _{(M,M_2),\lambda}(\frac{x}{y}), \]

(2)

\[ x\delta \bar{q}^a = \sum _\lambda \int _x^1 dy f^{\bar{\lambda}'} _{MB/N}(y) \frac{x}{y} \bar{q}^a _{M,\lambda'}(\frac{x}{y}) + \sum _\lambda \int _x^1 dy f^{\bar{\lambda}'} _{(M,M_2)B/N}(y) \frac{x}{y} \bar{q}^a _{(M,M_2),\lambda'}(\frac{x}{y}), \]

(3)

where $f^\lambda _{BM/N}(y)$, $f^{\bar{\lambda}'} _{MB/N}(y)$, $f^\lambda _{(B,B_2)M/N}(y)$ and $f^{\bar{\lambda}'} _{(M,M_2)B/N}(y)$ are the helicity dependent fluctuation functions

\[ f^\lambda _{BM/N}(y) = \sum _\lambda \int _0^\infty dk_1^2 |\phi _{BM}^\lambda (y,k_1^2)|^2, \]

\[ f^{\bar{\lambda}'} _{MB/N}(y) = \sum _\lambda \int _0^\infty dk_1^2 |\phi _{BM}^{\bar{\lambda}'} (1-y,k_1^2)|^2, \]

\[ f^\lambda _{(B,B_2)M/N}(y) = \sum _\lambda \int _0^\infty dk_1^2 \phi _{BM}^{\lambda \lambda'} (y,k_1^2) \phi _{B_2M}^{\lambda \lambda'} (y,k_1^2), \]

\[ f^{\bar{\lambda}'} _{(M,M_2)B/N}(y) = \sum _\lambda \int _0^\infty dk_1^2 \phi _{BM}^{\lambda \lambda'} (1-y,k_1^2) \phi _{B_2M}^{\lambda \lambda'} (1-y,k_1^2). \]

(4)

The last two terms in Eq. (2) and the last term in Eq. (3) are the interference contributions [15] which result from the possibility that interactions between the photon and different baryons (mesons) can lead to the same final states. The interference distributions $q^a _{(B,B_2),\lambda}$, $q^a _{(M,M_2),\lambda'}$ and $\bar{q}^a _{(M,M_2),\lambda'}$ which do not have simple interpretations in the quark-parton model, can be related to the PDFs of the vector mesons using SU(6) quark model wavefunctions [13, 15].
From Eqs. (2), (3) and (4) one can obtain the contributions to the polarized parton distributions of the nucleon,

\[ x \delta q = \int_x^1 dy \Delta f_{BM/N}(y) \frac{x}{y} \Delta q_B\left( \frac{x}{y} \right) + \int_x^1 dy \Delta f_{VBM}(y) \frac{x}{y} \Delta q_V\left( \frac{x}{y} \right) + \int_x^1 dy \Delta f_{(B_1B_2)M/N}(y) \frac{x}{y} \Delta q_{(B_1B_2)\left( \frac{x}{y} \right)} + \int_x^1 dy \Delta f_{(V_1V_2)B/N}(y) \frac{x}{y} \Delta q_{(V_1V_2)}\left( \frac{x}{y} \right) + \int_x^1 dy f^0_{(PV)B/N}(y) \frac{x}{y} \Delta q_{(PV)}\left( \frac{x}{y} \right) \]  

(5)

\[ x \delta \bar{q} = \int_x^1 dy \Delta f_{VBM}(y) \frac{x}{y} \Delta \bar{q}_V\left( \frac{x}{y} \right) + \int_x^1 dy \Delta f_{(V_1V_2)B/N}(y) \frac{x}{y} \Delta \bar{q}_{(V_1V_2)}\left( \frac{x}{y} \right) + \int_x^1 dy f^0_{(PV)B/N}(y) \frac{x}{y} \Delta \bar{q}_{(PV)}\left( \frac{x}{y} \right), \]  

(6)

where \( V \) (\( P \)) indicates the meson being a vector (pseudoscalar) meson, and

\[ \Delta f_{BM/N} = f^{1/2}_{BM/N} - f^{-1/2}_{BM/N} \]
\[ \Delta f_{(B_1B_2)M/N} = f^{1/2}_{(B_1B_2)M/N} - f^{-1/2}_{(B_1B_2)M/N} \]
\[ \Delta f_{VBM} = f^{1/2}_{VBM} - f^{-1/2}_{VBM} \]
\[ \Delta f_{(V_1V_2)B/N} = f^{1}_{(V_1V_2)B/N} - f^{-1}_{(V_1V_2)B/N} \]

(7)

are the polarized fluctuation functions. The various polarized parton distribution functions in Eqs. (5) and (6) are defined as follows:

\[ \Delta q_B = q_{B,1/2}^\uparrow - q_{B,1/2}^\downarrow = q_{B,-1/2}^\downarrow - q_{B,-1/2}^\uparrow, \]
\[ \Delta q_V = q_{V,1}^\uparrow - q_{V,1}^\downarrow = q_{V,-1}^\downarrow - q_{V,-1}^\uparrow, \]
\[ \Delta q_{(B_1B_2)} = q_{(B_1B_2),1/2}^\uparrow - q_{(B_1B_2),1/2}^\downarrow = q_{(B_1B_2),-1/2}^\downarrow - q_{(B_1B_2),-1/2}^\uparrow, \]
\[ \Delta q_{(V_1V_2)} = q_{(V_1V_2),1}^\uparrow - q_{(V_1V_2),1}^\downarrow = q_{(V_1V_2),-1}^\downarrow - q_{(V_1V_2),-1}^\uparrow, \]
\[ \Delta \bar{q}_V = \bar{q}_{V,1}^\uparrow - \bar{q}_{V,1}^\downarrow = \bar{q}_{V,-1}^\downarrow - \bar{q}_{V,-1}^\uparrow, \]
\[ \Delta \bar{q}_{(V_1V_2)} = \bar{q}_{(V_1V_2),1}^\uparrow - \bar{q}_{(V_1V_2),1}^\downarrow = \bar{q}_{(V_1V_2),-1}^\downarrow - \bar{q}_{(V_1V_2),-1}^\uparrow, \]
\[ \Delta \bar{q}_{(PV)} = \bar{q}_{(PV),0}^\uparrow - \bar{q}_{(PV),0}^\downarrow. \]  

(8)

For the polarized light quark sea of the proton, we consider fluctuations \( p \rightarrow N\pi, N\rho, N\omega \) and \( p \rightarrow \Delta \pi, \Delta \rho \). The fluctuation \( p \rightarrow \Delta \omega \) is neglected due to isospin symmetry. For the polarized strange sea we consider fluctuations \( p \rightarrow \Lambda K, \Sigma K \) and
\[ p \rightarrow \Lambda K^*, \Sigma K^*. \] Due to \( \pi \) and \( K \) mesons being pseudoscalar, the Fock states involving these mesons do not contribute to the polarization directly. However the interactions between photon and different mesons could lead to the same final states, so there are contributions from the interference effects between \(|B\pi\rangle\) and \(|B\rho(\omega)\rangle\), and \(|BK\rangle\) and \(|BK^*\rangle\) (see Eqs. (9), (10) and (12)). The final expressions for the polarized sea quark distributions are

\[
x \Delta u = \int_x^1 dy \left[ \frac{1}{6} \Delta f_{\rho N/N}(y) + \frac{2}{3} \Delta f_{\rho \Delta N}(y) + \frac{1}{2} \Delta f_{\omega N/N}(y) + \frac{1}{2} \Delta f_{(\omega N/N)(y)} \right] \frac{x}{y} \Delta V_{\rho}(\frac{x}{y}),
\]

\[
x \Delta d = \int_x^1 dy \left[ \frac{5}{6} \Delta f_{\rho N/N}(y) + \frac{1}{3} \Delta f_{\rho \Delta N}(y) + \frac{1}{2} \Delta f_{\omega N/N}(y) - \frac{1}{2} \Delta f_{(\omega N/N)(y)} \right] \frac{x}{y} \Delta V_{\rho}(\frac{x}{y}),
\]

\[
x \Delta s = \int_x^1 dy \left[ \Delta f_{\Lambda K/N}(y) + \Delta f_{\Lambda K^*/N}(y) \right] \frac{x}{y} \Delta s_{\Lambda}(\frac{x}{y})
\]

\[
+ \left[ \Delta f_{\Sigma K/N}(y) + \Delta f_{\Sigma K^*/N}(y) \right] \frac{x}{y} \Delta s_{\Sigma}(\frac{x}{y}),
\]

\[
x \Delta \bar{s} = \int_x^1 dy \left[ \Delta f_{K^* \Lambda N}(y) + \Delta f_{K^* \Sigma N}(y) \right] \frac{x}{y} \Delta s_{K^*}(\frac{x}{y})
\]

\[
+ \left( f_{(K^* K^* N)} \right) \Delta s_{K^*}(\frac{x}{y}).
\]

We have employed the relations among the helicity-dependent fluctuation functions \([18, 19]\) resulting from isospin symmetry and \(SU(3)\) flavour symmetry,

\[
\Delta f_{\rho^+ n/p} = 2 \Delta f_{\rho \rho n/p} = \frac{2}{3} \Delta f_{\rho N/N},
\]

\[
\Delta f_{\rho^- \Delta+/p} = \frac{3}{2} \Delta f_{\rho \Delta +/p} = 3 \Delta f_{\rho^+ \Delta^0/p} = \frac{1}{2} \Delta f_{\rho \Delta N},
\]

\[
f_{(\pi^+ \rho^0)n/p} = 2 f_{(\pi^0 \rho^0)p/p} = f_{(\pi \rho)n/N},
\]

\[
f_{(\pi^- \rho^-)\Delta+/p} = \frac{3}{2} f_{(\pi^0 \rho^0)\Delta+/p} = 3 f_{(\pi^+ \rho^0)\Delta^0/p} = \frac{1}{2} f_{(\pi \rho)\Delta N},
\]

\[
\Delta f_{\Sigma^+ K^0n/p} = 2 \Delta f_{\Sigma^0 K^+n/p} = \frac{2}{3} \Delta f_{\Sigma^+ K/N},
\]

\[
\Delta f_{\Sigma^+ K^-n/p} = 2 \Delta f_{\Sigma^0 K^+n/p} = \frac{2}{3} \Delta f_{\Sigma^+ K^*n},
\]

\[
\Delta f_{K^+ \Sigma^0n/p} = 2 \Delta f_{K^+ \Sigma^0n/p} = \frac{2}{3} \Delta f_{K^+ \Sigma^+ N},
\]

\[
f_{(K^0 K^0)\Sigma^+n/p} = 2 f_{(K^+ K^+ \Sigma^0n/p} = f_{(K K^*)\Sigma^0 N}.
\]

\[7\]
The polarized parton distributions and ‘interference’ distributions of different charge states of baryons (mesons) are related using SU(6) wavefunction of the baryons (mesons) [15],

\[
\Delta \bar{d}^\rho = \Delta \bar{u}^\rho = 2\Delta \bar{d}^\rho = 2\Delta \bar{u}^\rho = \Delta V^\rho,
\]

\[
\Delta \bar{d}^\rho_\omega = -\Delta \bar{u}^\rho_\omega = -\frac{1}{2}\Delta V^\rho,
\]

\[
\Delta \bar{d}^{(\pi^+\rho^+)}_o = \Delta \bar{u}^{(\pi^-\rho^-)}_o = 2\Delta \bar{d}^{(\pi^0\rho^+)}_o = 2\Delta \bar{u}^{(\pi^0\rho^-)}_o = \Delta V^\rho,
\]

\[
\Delta \bar{d}^{(\pi^0\omega)}_o = -\Delta \bar{u}^{(\pi^0\omega)}_o = -\frac{1}{2}\Delta V^\rho,
\]

\[
\Delta s_{K^*} = \Delta s_{K^*}^+ = \Delta s_{K^*}^0,
\]

\[
\Delta s_{K^*K^*} = \Delta s_{K^*K^*}^+ = \Delta s_{K^*K^*}^0 = \Delta s_{K^*},
\]

\[
\Delta s_{\Sigma} = \Delta s_{\Sigma^+} = \Delta s_{\Sigma^0},
\]

\[
\Delta s_{\Lambda\Sigma^0} = 0.
\]  

The interference distribution \(\Delta s_{\Lambda\Sigma^0}\) vanishes so there is no interference contribution to the polarized strange sea of the nucleon (see Eq. (11)).

The wave functions \(\phi^{\lambda\lambda'}_{BM}(y, k^2_\perp)\) which determine the fluctuation functions (see Eq. (4)) are calculated using time-order perturbation theory in the infinite momentum frame (the meson being treated as on the energy-shell i.e. \(E_M = E_N - E_B\)) [17],

\[
\phi^{\lambda\lambda'}_{BM}(y, k^2_\perp) = \frac{1}{2\pi\sqrt{y(1-y)}} \sqrt{m_N m_B} \frac{V^{\lambda\lambda'}_{IMF}(y, k^2_\perp) G_{BM}(y, k^2_\perp)}{m_N^2 - m_B^2(y, k^2_\perp)},
\]  

where \(m_{BM}^2\) is the invariant mass squared of the BM Fock state,

\[
m_{BM}^2(y, k^2_\perp) = \frac{m_B^2 + k^2_\perp}{y} + \frac{m_M^2 + k^2_\perp}{1 - y}.
\]  

The vertex function \(V^{\lambda\lambda'}_{IMF}(y, k^2_\perp)\) depends on the effective interaction Lagrangian that describes the fluctuation process \(N \rightarrow BM\). From the meson exchange model for hadron production [17, 21] we have

\[
\mathcal{L}_1 = ig \bar{N} \gamma_5 \pi B,
\]

\[
\mathcal{L}_2 = f \bar{N} \partial_\mu \pi \Delta^\mu + h.c.,
\]

\[
\mathcal{L}_3 = g \bar{N} \gamma_\mu \theta^\mu B + f \bar{N} \sigma_{\mu\nu} B (\partial^\mu \theta^\nu - \partial^\nu \theta^\mu),
\]

\[
\mathcal{L}_4 = if \bar{N} \gamma_5 \gamma_\mu \Delta (\partial^\mu \theta^\nu - \partial^\nu \theta^\mu) + h.c.,
\]  

where \(m_B^2, m_M^2\) are the masses of the baryon and meson, respectively.
where $N$ and $B = \Lambda, \Sigma$ are spin-1/2 fields, $\Delta$ a spin-3/2 field of Rarita-Schwinger form ($\Delta$ baryon), $\pi$ a pseudoscalar field ($\pi$ and $K$ mesons), and $\theta$ a vector field ($\rho$, $\omega$ and $K^*$). The coupling constants for various considered fluctuations are taken to be [20, 21],

\[
\begin{align*}
    g_{N\pi}^2/4\pi &= 13.6, \\
    g_{NN\rho}^2/4\pi &= 0.84, \\
    g_{NN\omega}^2/4\pi &= 8.1, \\
    f_{N\pi}/g_{NN\rho} &= 6.1/4m_N, \\
    f_{N\omega}/g_{NN\omega} &= 0, \\
    f_{N\Delta\pi}/4\pi &= 12.3 \text{ GeV}^{-2}, \\
    f_{N\Delta\rho}/4\pi &= 34.5 \text{ GeV}^{-2}, \\
    g_{N\Lambda K} &= -13.98, \\
    g_{N\Sigma K} &= 2.69, \\
    g_{N\Lambda K^*} &= -5.63, \\
    g_{N\Sigma K^*} &= -3.25, \\
    f_{N\Sigma K^*} &= 2.09 \text{ GeV}^{-1}.
\end{align*}
\]

The phenomenological vertex form factor $G_{BM}(y, k_{\perp}^2)$ in Eq. (16) is introduced to describe the unknown dynamics of the fluctuation $N \rightarrow BM$, for which we adopt the exponential form

\[
G_{BM}(y, k_{\perp}^2) = \exp \left[ \frac{m_N^2 - m_{BM}(y, k_{\perp}^2)}{2\Lambda^2} \right],
\]

with $\Lambda_C$ being a cut-off parameter. This form factor satisfies the relation $G_{BM}(y, k_{\perp}^2) = G_{MB}(1 - y, k_{\perp}^2)$.

III. POLARIZED PARTON DISTRIBUTION FUNCTIONS OF HADRONS

The polarized parton distribution functions of the hyperons $\Lambda$ and $\Sigma$ and mesons $\rho$, $\omega$ and $K^*$ are largely unknown. In order to estimate these distributions we extend the method of the Adelaide group [22] which uses the bag model to evaluate the parton distributions of baryons. The bag model calculations give results consistent with the experimental data for the parton distributions of the nucleon, and the calculations have been extended to other baryons [13]. As the bag model gives a good description of many non-perturbative properties (e.g. the mass spectrum) of the mesons except for the pion [23], we argue that bag model calculations of the parton distributions of the mesons should give a reasonably good approximation to these distributions.
In the bag model the dominant contributions to the parton distribution functions of a hadron in the medium-$x$ range come from intermediate states with the lowest number of quarks, so the intermediate states we consider contain one quark (or anti-quark) for the mesons and two quarks for the hyperons. Following [22] we can write these contributions as

$$q_{h,f}^{\uparrow\downarrow}(x) = \frac{M_h}{(2\pi)^3} \sum_m \langle \mu | P_{f,m} | \mu \rangle \int d\mathbf{p}_n \frac{\left| \phi_i(\mathbf{p}_n) \right|^2}{\left| \phi_j(0) \right|^2} \delta(M_h(1-x) - p_n^+) |\tilde{\Psi}_{+}^{\uparrow\downarrow}(\mathbf{p}_n)|^2,$$  \hspace{1cm} (21)

where $M_h$ is the hadron mass, ‘$+$’ components of momenta are defined by $p^+ = p^0 + p^3$, and $\mathbf{p}_n$ is the 3-momentum of the intermediate state. $\tilde{\Psi}$ is the Fourier transform of the MIT bag ground state wavefunction $\Psi(\mathbf{r})$, and $\phi_m(\mathbf{p})$ is the Fourier transform of the Hill-Wheeler overlap function between $m$-quark bag states:

$$|\phi_m(\mathbf{p})|^2 = \int d\mathbf{r} e^{-i\mathbf{p} \cdot \mathbf{r}} \left[ \int d\mathbf{r} \Psi^\dagger(\mathbf{r} - \mathbf{R}) \Psi(\mathbf{r}) \right]^m.$$ \hspace{1cm} (22)

In Eq. (21) one takes $i = 1$, $j = 2$ for the mesons ($\rho, \omega, K$, and $K^*$) and $i = 2$, $j = 3$ for the hyperons ($\Lambda$ and $\Sigma$). The matrix element $\langle \mu | P_{f,m} | \mu \rangle$ appearing in Eq. (21) is the matrix element of the projection operator $P_{f,m}$ onto the required flavour $f$ and helicity $m$ for the $SU(6)$ spin-flavour wavefunction $|\mu\rangle$ of the hadron under consideration.

The input parameters in the bag model calculations are the bag radius $R$, the mass of the quark (anti-quark) $m_q$ for which the parton distribution is calculated, the mass of the intermediate state $m_n$, and the bag scale $\mu^2$ - at this scale the model is taken as a good approximation to their valence structure of the hadron. In Table I we list the values for these parameters adopted in this work. In figure 1 we show the polarized parton distribution functions after NLO evolution [24] to the scale of $Q^2 = 2.5$ GeV$^2$ where the HERMES results are available. As expected, the PDFs of the mesons are harder than those of the hyperons since the dominant intermediate states of the mesons containing one quark (anti-quark) are lighter than those of the hyperons containing two quarks. Also we can see that $x\Delta s_{K^*}$ is harder than $x\Delta V_\rho$ due to the $s$-quark being heavier than the $u$- and $d$-quarks. The polarization of $s$-quark in the $\Lambda$ hyperon is positive while it is negative for the $\Sigma$ hyperon since the $SU(6)$ wavefunction of the $\Sigma$ is dominated by the term $|u^\uparrow d^\uparrow s^\downarrow\rangle$. The polarized PDF of strange quarks in the $\Lambda$ ($x\Delta s_\Lambda$) is harder than that of the $\Sigma$ ($x\Delta s_\Sigma$)
because the two-quark intermediate state for the Λ is a light scalar while it is a vector for the Σ which is heavier by 200 MeV because of the hyperfine splitting between \(qq\) states.

IV. NUMERICAL RESULTS AND DISCUSSIONS

The fluctuation functions depend on the cut-off parameter \(\Lambda_C\) introduced in the phenomenological vertex form factor \(G_{BM}\) (see Eq. (20)). We adopt \(\Lambda_C^{\text{oct}} = 1.08\) GeV and \(\Lambda_C^{\text{dec}} = 0.98\) GeV for the fluctuations involving the octet and decuplet baryons respectively [17]. These values are determined from analysis of high energy \(p-p\) and \(p-\Lambda\) scattering, and give a reasonable fit to \(\bar{d}(x) - \bar{u}(x)\) in the unpolarized nucleon sea [17]. The polarized fluctuation functions needed for calculating the polarized strange sea of the nucleon (see Eqs. (11) and (12)) are shown in figure 2. For the fluctuation functions involved in calculating the polarized light quark sea of the nucleon (\(\Delta \bar{u}, \Delta \bar{d}\)) we refer to [15]. We note that the fluctuation functions \(\Delta f_{\Lambda K^*/N}\) and \(\Delta f_{\Sigma K^*/N}\) are larger in magnitude than \(\Delta f_{\Lambda K/N}\) and \(\Delta f_{\Sigma K/N}\), although one might expect that the \(K^*\) fluctuation functions would be smaller than the \(K\) fluctuation functions due to the higher mass of the \(K^*\). Also the \(K^*\) fluctuation functions are negative while the \(K\) fluctuation functions are positive, so the calculation of \(\Delta s\) is sensitive to whether contributions from fluctuations involving \(K^*\) are included or not. In the case of \(\Delta \bar{s}\), the contributions from \(K^*\) states are the leading contributions in the MCM. The sum of the fluctuation functions \(f^0_{(KK^*)\Lambda/N}\) and \(f^0_{(KK^*)\Sigma/N}\) changes sign and is much smaller in magnitude than the sum of \(\Delta f_{K^*/\Lambda/N}\) and \(\Delta f_{K^*/\Sigma/N}\), which indicates that the contribution from the \(K-K^*\) interference is not significant. The same conclusion is also true in the calculation for the light polarized quark sea [15].

The results for the light quark sea \(x \Delta \bar{u}(x)\) and \(x \Delta \bar{d}(x)\), along with the data from the HERMES collaboration, are presented in figure 3. The calculations show that the polarizations in the light quark sea are positive and the polarization of the anti-up quark is about 10% of that of the anti-down quark. Thus the \(SU(2)\) flavor symmetry (\(\Delta \bar{u}(x) = \Delta \bar{d}(x)\)) in polarized light quark sea is broken and \(\Delta \bar{u}(x) < \Delta \bar{d}(x)\) over the range of \(0.01 < x < 0.6\). The calculations for \(x \Delta \bar{u}(x)\) and \(x \Delta \bar{d}(x)\) are consistent with the data. To highlight the flavour symmetry breaking, we calculate the difference \(x(\Delta \bar{u}(x) - \Delta \bar{d}(x))\) and compare it with the HERMES result in figure 4. Our theoretical calculations are
consistent with the data, although large uncertainties exist in the data. Also we can find the SU(2) flavour symmetry breaking in the polarized nucleon sea is much smaller than in the unpolarized sea, which is in contrast to calculations using chiral quark soliton model [25] which predict the differences \((\Delta \bar{d} - \Delta \bar{u})\) and \((\bar{d} - \bar{u})\) are similar in magnitude.

The contributions to the light polarized antiquark distributions calculated in this work come mainly from the antiquark in the meson cloud. There may be other non-perturbative contributions to flavour symmetry breaking of the parton distribution of the bare nucleon. Some studies [15, 26] estimated that these contributions could be significantly larger than the contributions from the meson cloud by considering Pauli blocking effects. However the HERMES data indicate that these non-perturbative contributions from the bare nucleon cannot be very large. As Pauli blocking effects are expected to be of similar size in both polarized and unpolarized case [15, 26], this conclusion may be of important in discussions of \(\bar{d} - \bar{u}\) difference [9]

In figure 5 we show the polarization of the strange sea calculated both with and without the contributions from Fock states involving \(K^*\) mesons. We can see that the predictions depend strongly on contributions from the \(K^*\) Fock states. We have arrived at a similar conclusion on the importance of considering the \(K^*\) mesons in a recent investigation of the unpolarized strange sea [14]. To study the quark-antiquark symmetry breaking in the polarized strange sea we show the difference \(x(\Delta s(x) - \Delta \bar{s}(x))\) in figure 6. It can be seen that \(x(\Delta s(x) - \Delta \bar{s}(x)) < 0\) when both contributions from \(K\) and \(K^*\) mesons are included, while \(x(\Delta s(x) - \Delta \bar{s}(x)) > 0\) when only \(K\) mesons are considered. In figure 7 we compare theoretical calculations for \(x(\Delta s(x) + \Delta \bar{s}(x))/2\) with the HERMES measurement for \(x\Delta s(x)\). It is interesting to compare the strange-antistrange asymmetry in the polarized sea with that in the unpolarized sea. We present such a comparison in figure 8. We find that the strange-antistrange asymmetry is much more significant in the polarized sea than in the unpolarized sea.

The integrals of polarized parton distribution functions \((\Delta Q = \int_0^1 \Delta q(x) dx)\) give the contribution to the spin of the nucleon carried by each flavor of parton. We found that \(\Delta \bar{U} = 0.001\), \(\Delta \bar{D} = 0.03\) and \(\Delta S + \Delta \bar{S} = 0.01 (0.004)\) with (without) the \(K^*\) Fock states. The total spin carried by charged partons \((\Delta \Sigma)\) is determined by DIS experiments to be about 0.3, so the light antiquark sea and strange quark sea contribute about 10% of
this total spin. The polarization of the strange quark sea is found to be positive which is in contradiction to the previous conclusion that the strange quark and anti-quarks are polarized negatively with respect to the direction of the nucleon spin ($\Delta S + \Delta S^\prime < 0$) based on analyses of inclusive deep inelastic scattering (DIS) [27] and lattice calculations [28]. Also our result for $\Delta S + \Delta S^\prime$ is about 10% of the magnitude found in previous analyses. However our prediction of a positively polarised strange sea ($\Delta S + \Delta S^\prime$) agrees with the HERMES result, $\int_{0.03}^{0.3} \Delta s(x)dx = 0.03 \pm 0.03(\text{stat.}) \pm 0.01(\text{sysrt.})$.

V. SUMMARY

The polarized parton distribution functions of the nucleon sea provide vital information on the non-perturbative structure of the nucleon. In this paper we have calculated the polarized parton distribution functions of the nucleon sea using the meson cloud model and thereby investigated the flavour and quark-antiquark symmetries of the nucleon. Our calculations show that the SU(2) flavour symmetry and quark-antiquark symmetry in the polarized nucleon sea are broken and $\Delta \bar{u}(x) < \Delta \bar{d}(x)$ and $\Delta s(x) < \Delta \bar{s}(x)$. SU(2) flavour symmetry breaking in the polarized nucleon sea is found to be much smaller than in the unpolarized sea. This is in contrast to calculations in the chiral quark soliton model, or calculations based on Pauli blocking, which have found $(\Delta \bar{d} - \Delta \bar{u})$ to be similar in magnitude to $(\bar{d} - \bar{u})$.

The strange-antistrange symmetry breaking is much larger in the polarized nucleon than in the unpolarized nucleon. Our finding of a slightly positively polarized strange sea is remarkably different from previous determinations of a significantly negatively polarized strange sea. This may be due to a breaking of SU(3) flavour symmetry, e.g. the F and D values calculated from hyperon decays may not apply to the nucleon. The contributions to the total spin carried by the charged partons from the light antiquark sea and strange sea is about 10%. Our calculations generally agree with recent results from the HERMES Collaboration though large error bars exist in the data. More experimental data with high precision are highly desired and will put more rigorous constraints on models of nucleon structure.
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[1] J. Ashman et al., EMC collaboration, Phys. Lett. B 206, 364 (1988); J. Ashman et al., Nucl. Phys. B328, 1 (1989).
[2] See e.g. K. Abe et al., E-154 collaboration, Phys. Rev. Lett. 79, 26 (1997); B. Adeva, SMC collaboration, Phys. Rev. D 58, 112001 (1998).
[3] J. Ellis and M. Karliner, Phys. Lett. B 341, 397 (1995).
[4] B. Adeva et al., SMC Collaboration, Phys. Rev. D 58, 112001 (1998).
[5] P. L. Anthony, E155 Collaboration, Phys. Lett. B 458, 529 (1999).
[6] A. Airapetian et al., HERMES Collaboration, hep-ex/0307064, submitted to Phys. Rev. Lett.; M. Beckmann, hep-ex/0210049.
[7] K. Ackerstaff et al., HERMES Collaboration, Phys. Lett. B 464, 123 (1999).
[8] B. Adeva et al., SMC Collaboration, Phys. Lett. B 420, 180 (1998).
[9] See e.g., S. Kumano, Phys. Rep. 303, 103 (1998); G. T. Garvey and J. C. Peng, Prog. Part. Nucl. Phys. 47, 203 (2001); W. Melnitchouk, J. Speth, and A. W. Thomas, Phys. Rev. D 59, 014033 (1998).
[10] P. Amaudraz et al., Phys. Rev. Lett. 66, 2712 (1991); M. Arneodo et al., Phys. Rev. D 50, R1 (1994); M. Arneodo et al., Phys. Lett. B 364, 107 (1995).
[11] E. A. Hawker et al., E866/NuSea Collaboration, Phys. Rev. Lett. 80, 3715 (1998); hep-ex/0103030.
[12] J. Ellis and R. Jaffe, Phys. Rev. D 9, 1444 (1974).
[13] C. Boros and A. W. Thomas, Phys. Rev. D 60, 074017 (1999).
[14] F. G. Cao and A. I. Signal, Phys. Lett. B 559, 229 (2003).
[15] F. G. Cao and A. I. Signal, Eur. Phys. J. C 21, 105 (2001).
[16] A. W. Thomas, Phys. Lett. B 126, 97 (1983).
[17] H. Holtmann, A. Szczurek, and J. Speth, Nucl. Phys. A569, 631 (1996).
[18] W. Melnitchouk, J. Speth, and A. W. Thomas, Phys. Rev. D 59, 014033 (1998).
[19] F. G. Cao and A. I. Signal, Phys. Rev. C 62, 015203 (2000).
[20] B. Holzenkamp, K. Holinde and J. Speth, Nucl. Phys. A500, 1989 (485).
[21] R. Machleidt, K. Holinde, and Ch. Elster, Phys. Rep. 149, 1 (1987).
[22] A. I. Signal and A. W. Thomas, Phys. Lett. B 211, 481 (1988); Phys. Rev. D 40, 2832 (1989); A. W. Schreiber, A. W. Thomas, and J. T. Londergan, Phys. Rev. D 42, 2226 (1990); A. W. Schreiber, A. I. Signal and A. W. Thomas, Phys. Rev. D 44, 2653 (1991).
[23] T. DeGrand, R. L. Jaffe, K. Johnson and J. Kiskis, Phys. Rev. D 12, 2060 (1975); A. W. Thomas, Adv. Nucl. Phys. 13, 1 (1984).
[24] M. Hirai, S. Kumano and M. Miyama, Comput. Phys. Commun. 108, 38 (1998).
[25] see e.g. D. I. Diakonov et. al., Nucl. Phys. B480, 341 (1996); Phys. Rev. D 56, 4069 (1997); P. V. Pobylitsa and M. V. Polyakov, Phys. Lett. B 389, 350 (1996); B. Dressler et. al., Eur. Phys. J. C 14, 147 (2000); M. Wakamatsu, hep-ph/0010262.
[26] M. Glück and E. Reya, Mod. Phys. Lett. A 15, 883 (2000).
[27] see e.g., D. Adeva et.al., Phys. Rev. D 58, 112002 (1998); Y. Goto et.al., Phys. Rev. D 62, 034017 (2000); M. Gluck, E. Reya, M. Stratmann and W. Vogelsang, Phys. Rev. D 63, 094005 (2001); J. Blumlein and H. Bottcher, Nucl. Phys. B636, 225 (2002).
[28] M. Fukugita et. al., Phys. Rev. Lett. 75, 2092 (1995); S. J. Dong et. al., Phys. Rev. Lett. 75, 2096 (1995).
Figure Captions

Fig. 1. Bag model calculations for the polarized parton distribution functions of hyperons $\Lambda (x\Delta s_\Lambda$, solid curve) and $\Sigma (x\Delta s_\Sigma$, dashed curve) and mesons $K^*(x\Delta s_{K^*}$, dotted curve) and $\rho (x\Delta V_\rho$, dash-dotted curve). All distributions are evolved to the scale of $Q^2 = 2.5 \text{ GeV}^2$.

Fig. 2. The polarized fluctuation functions. The thin solid and dashed curves are for $\Delta f_{\Lambda K/N}$ and $\Delta f_{\Sigma K/N}$ respectively while the thick solid and dashed curves are for $\Delta f_{\Lambda K^*/N}$ and $\Delta f_{\Sigma K^*/N}$ respectively. The thin and thick dotted curves are for $\Delta f_{K^*/\Lambda/N} + \Delta f_{K^*/\Sigma/N}$ and the interference term $f_{0(KK^*)\Lambda/N}^0 + f_{0(KK^*)\Sigma/N}^0$.

Fig. 3. Comparison of theoretical calculations for the polarized light anti-quark sea and the experimental data from the HERMES collaboration [6].

Fig. 4. Flavour asymmetry in the polarized light anti-quark sea. The solid curve is the theoretical calculation. The HERMES data are taken from [6].

Fig. 5. Polarized strange sea of the nucleon. The thin solid curve is theoretical calculation for $x\Delta s$ when only $K$ Fock states being considered. The thick solid and dashed curves are results for $x\Delta s$ and $x\Delta \bar{s}$ including the contributions from both $K$ and $K^*$ Fock states.

Fig. 6. Strange-antistrange asymmetry $x(\Delta s - \Delta \bar{s})$ in the polarized nucleon sea. The solid curve is the result including the contributions from both $K$ and $K^*$ Fock states, while the dashed curve is the result including only the $K$ Fock states.

Fig. 7. Comparison of theoretical calculations for $x(\Delta s + \Delta \bar{s})/2$ with the HERMES results for $x\Delta s(x)$ [6] at $Q^2 = 2.5 \text{ GeV}^2$.

Fig. 8. Strange-antistrange asymmetry in the unpolarized and polarized sea. $x(s - \bar{s})$ at $Q^2 = 16 \text{ GeV}^2$, dashed curve, and $x(\Delta s - \Delta \bar{s})$ at $Q^2 = 2.5 \text{ GeV}^2$, solid curve.
Table I. Input parameters in the bag model calculation.

|     | $R$(fm) | $m_q$(MeV) | $m_n$(MeV) | $\mu^2$(GeV$^2$) |
|-----|---------|------------|------------|------------------|
| $\Lambda$ | 0.8    | 150        | 650       | 0.23             |
| $\Sigma$   | 0.8    | 150        | 850       | 0.23             |
| $\rho$     | 0.7    | 0          | 425       | 0.23             |
| $K^*$      | 0.7    | 150        | 425       | 0.23             |
Fig. 2
Fig. 3
$x$

$\Phi_{\bar{d}^0\phi_c}$

$(\rho\gamma - \eta\gamma)x$
Fig. 8