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# Chiral Symmetry and the Charge Asymmetry of the $s\bar{s}$ Distribution in the Proton[+]

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## ABSTRACT

Recently it has been suggested that a sizeable fraction of the strange and charm quarks in a nucleon—the so-called "intrinsic strangeness or charm"— have momentum distributions which extend to large $x_{bj}$. This effect is enhanced if these virtual heavy quarks live long enough such that many interactions with the rest of the nucleon can occur. It is shown that the same mechanism responsible for the intrinsic component also leads to a sizable charge asymmetry of the corresponding spin and momentum distributions.

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## INTRODUCTION

Most virtual $s\bar{s}$ pairs[#1] in a proton have a very short lifetime (of the order $\tau \sim \frac{1}{\sqrt{-q^2}}$, where $q$ is the momentum transfer in the deep inelastic scattering process). They are concentrated at small $x_{bj}$ and arise primarily from logarithmic $QCD$ evolution.[2] The underlying microscopic process is the incoherent fragmentation of a gluon into a $Q\bar{Q} = (s\bar{s}, c\bar{c})$ pair where interactions with other partons (spectators) are neglected. The resulting $Q\bar{Q}$ sea is then characterized by the following properties:

inclusive:

- The spin and momentum distribution of the $Q$ and $\bar{Q}$ are the same, by charge conjugation and using that the $Q\bar{Q}$ pair is too short-lived to interact with the rest of the proton (and thus cannot find out whether it has been created in a proton or antiproton).
- The spin and momentum carried by the pair are proportional to the gluon spin and momentum and thus the $Q\bar{Q}$ pairs are typically concentrated at low $x_{bj}$.

exclusive:

- The sum of the magnetic moment contributions of $s$ and $\bar{s}$ is zero by charge conjugation (see above).

Besides these perturbative or extrinsic $s\bar{s}$ pairs the proton is expected to contain also a more long-lived[#2] component of virtual pairs.[1,4] Of course the initial

#1 Most of the conclusions in this work remain qualitatively correct if we replace $s\bar{s}$ by $c\bar{c}$ though there will be a quantitative difference.

#2 Long-lived means here a lifetime of the order $M_{Q\bar{Q}}^{-1}$ [3]

process for creation of $s\bar{s}$ pairs is always the same: fragmentation of a gluon. However, a few of these sea quarks—the "intrinsic" component—do not immediately recombine, and interact for some time with other quarks and gluons in the hadron. One major difference between extrinsic and intrinsic $s\bar{s}$ pairs is that intrinsic ones can be found at larger values of $x_{bj}$. This is because they have time to reach an energetically more favorable (i.e. less off-shell) state, where the light-cone kinetic energy

$$P^{-}_{\text{kin}} = \sum_i \frac{m_i^2 + k_{i\perp}^2}{x_i} \tag{1}$$

is close to the minimum value[1]. Thus small values of $x_{bj}$ — in particular for heavy quarks — are supressed in these long-lived components. In this work we will concentrate on this component and see what general features of the corresponding distribution functions we can predict. Unless otherwise stated, all following remarks concerning sea quarks will refer to this intrinsic component.

In order to reach large values of $x_{bj}$ (i.e. $x_{bj} \gtrsim 0.2$) a sea quark has to undergo several interactions while accumulating more and more momentum fraction.[#3] During that process the $qqqQ\bar{Q}$ fluctuation tends to arrange itself into energetically more favorable states. In the case of $Q = s$ the lowest contributing state with the right quantum numbers is a $\Lambda K^+$ state[#4], which is thus expected to play an important role for $s\bar{s}$ production at large $x_{bj}$. In order to understand the qualitative implications of this picture let us assume for the moment that the $p \to \Lambda K^+$ fluctuation is the only source of virtual $s\bar{s}$ pairs in a proton. The conse-

#3 This indicates already that perturbative QCD is not appropriate to describe this component of the proton wavefunction and we have to use other approximation schemes for these large $x_{bj}$ sea quarks.

#4 It is assumed here that the lifetime of the fluctuation is large enough to allow formation of these hadrons.

quences for some spin observables are then clear. Angular momentum and parity conservation require the $K^+$ to be emitted in an $\ell = 1$ state and the total angular momentum wavefunction reads

$$\left|J = \frac{1}{2}, J_z = \frac{1}{2}\right\rangle = \left[\sqrt{2}\,|\ell = 1, \ell_z = 1\rangle \left|s = \frac{1}{2}, s_z = -\frac{1}{2}\right\rangle \right. \\ \left. - |\ell = 1, \ell_z = 0\rangle \left|s = \frac{1}{2}, s_z = \frac{1}{2}\right\rangle\right] \frac{1}{\sqrt{3}} . \tag{2}$$

In a constituent quark model the $\Lambda$ spin is carried by its $s$ quark. It is thus most likely to find the $s$-quark with polarization antiparallel to the initial proton spin.[#5] In this oversimplified picture the $\bar{s}$ is unpolarized because the $K^+$, where it is contained, is spinless. Later we will see that the chiral symmetry of the interaction demands an additional scalar meson which, through interference with the pseudoscalar $K^+$, yields $\bar{s}$ quarks polarized parallel to the initial proton spin. Also vector mesons, like the $K^*$, which have been neglected here, can yield polarized $\bar{s}$.

Note that both $s$ and $\bar{s}$ contribute to the proton's magnetic moment with the same sign (both parallel to the proton spin). This is because the $s$ has negative charge, thus compensating for the antiparallel spin, and the positively charged $\bar{s}$ has orbital angular momentum parallel to the proton spin.

Binding of the quarks in pseudoscalar mesons is (due to chiral symmetry) usually stronger than in baryons. This has striking consequences for the (unpolarized) momentum distributions.[14] In order to see this let us assume that the momentum

#5 Strictly speaking this nonrelativistic reasoning used here cannot be applied to the structure functions. However, our explicit calculations in the context of the Gross-Neveu model confirmed these heuristically obtained results.

in the $\Lambda K$ system is shared such that the light-cone kinetic energy is minimized,[4] i.e.[#6]

$$\frac{\langle x_\Lambda \rangle}{m_\Lambda} \approx \frac{\langle x_K \rangle}{m_K} \quad \text{or} \quad \langle x_\Lambda \rangle \approx \frac{m_\Lambda}{m_\Lambda + m_K}\,, \quad \langle x_K \rangle \approx \frac{m_K}{m_\Lambda + m_K} \tag{3}$$

and that a corresponding relation is valid for the quarks inside the $\Lambda$ and $K$, i.e.

$$\langle x_s \rangle \approx \frac{M_s}{M_s + 2M_u} \langle x_\Lambda \rangle \approx 0.3 \qquad \langle x_{\bar{s}} \rangle \approx \frac{M_s}{M_s + M_u} \langle x_K \rangle \approx 0.2 \tag{4}$$

(in these estimates, involving long-lived fluctuations, it is appropriate to use the constituent quark masses $M_u \approx 350\ MeV$ and $M_s \approx 500\ MeV$). Using

$$\frac{\langle x_s \rangle}{\langle x_{\bar{s}} \rangle} = \frac{\frac{m_\Lambda}{M_s + 2M_u}}{\frac{m_K}{M_s + M_u}} > 1$$

it is evident that the stronger binding of the $\bar{s}$ in the $K$ allows smaller values of $x_{bj}$ for the $\bar{s}$ than for the $s$.[#7] Although this (very crude) picture cannot be taken more than qualitative, more realistic models should exhibit a similar trend, and we give an example.

We should emphasize the role of chiral symmetry in this context[8]. The most important point here are the low masses[#8] of the pseudoscalar octet which make those mesons the source of the energetically lowest excitation of nucleons with intrinsic sea quarks. It is the low mass of these mesons which is responsible for the

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#6 The momentum fractions computed here are momentum fractions in the $\Lambda K^+$ system. Of course, in order to estimate the absolute momentum fraction in a proton which is carried by $s$ or $\bar{s}$ one has to multiply these numbers by the probability to find intrinsic sea quarks, i.e. by the probability by which the proton is in a virtual $\Lambda K$ state.

#7 Although perturbative $QCD$ predicts the same scaling power $p$ for $s(x)$ and $\bar{s}(x)$ as $x \rightarrow 1$[15], the coefficient of $(1-x)^p$ can be quite different for quarks and antiquarks and does not follow from simple counting rules.

#8 For zero quark mass those mesons would be Goldstone bosons.

peculiar asymmetry in the momentum splitting between $s$ and $\bar{s}$ quark. Furthermore, for the predictions concerning spin and magnetic moment of the $s$ quark it was important that the kaon is spinless.

For heavy quarks ($Q = c, b$) chiral symmetry is badly broken However, for charm quarks there is some evidence for an intrinsic component[1] and, like chiral symmetry, the color-hyperfine splitting yields comperatively light pseudoscalars. But the effect is relatively small which results in a decreased importance of the $\Lambda_c D$ component. Hence, the charge asymmetry should be smaller for these quarks.

In the following some model calculations will be used to demonstrate what size of effects one can expect. For this purpose one could have in mind to develop some kind of convolution model where the (dressed) nucleon structure function is given by a convolution of the bare nucleon structure function with the bare meson structure function and a relative wavefunction.[5,9]

Here one faces immediately some conceptual difficulties. E.g. the kaons inside a nucleon are off-shell and it is a priori not clear how the off-shell structure function of a kaon relates to its on-shell structure function[#9] and how this depends on the off-shellness. Furthermore, it is not clear how many mesons, besides the pseudoscalar octet, one should take into account. So far such questions have made it very difficult to study the impact of chiral symmetry and chiral symmetry breaking on the structure function of a nucleon.[8]

#9 The latter is also not known but could be determined in a fit procedure.

## Charge Asymmetry of the Strange Sea in the Chiral Gross-Neveu Model

In order to avoid the above mentioned difficulties we start with studying a chirally symmetric generalization of the $1+1$-dimensional Gross-Neveu (GN) model,[10] which can be described in terms of quark degrees of freedom only.

This model is relevant for the above discussion since it is an example with spontaneous chiral symmetry breaking. Furthermore it is renormalizable and asymptotically free (in $1+1$ dimensions), hence deep inelastic structure functions scale in the Bjorken limit, and it makes sense to relate deep inelastic scattering observables to parton distributions.

Since we will define the model in terms of quark degrees of freedom only, the Goldstone bosons will be automatically composite. Most importantly a consistent and physically simple interpretation of the parton distribution arising from the meson cloud becomes possible within this model.

We start from a "chirally"-symmetric generalization of the Gross-Neveu model[10]

$$\mathcal{L}_0 = i\bar{\psi}\, \not{\partial}\psi + \frac{g^2}{2N_c}\, \left[(\bar{\psi}\tau^i\psi)^2 - (\bar{\psi}\tau^i\gamma_5\psi)^2\right] \tag{5}$$

where the quark fields carry both color and flavor (the $\tau^i$ generate the $U(N)$ flavor symmetry subgroup). In leading order in $1/N_C$ the ground state of $\mathcal{L}_0$ breaks chiral symmetry, i.e. $\bar{\psi}\psi$ develops a non-zero vacuum expectation value and hence an effective mass for the fermions is generated. Now, quarks (in the real world) have non-zero current masses, i.e. chiral symmetry is explicitly broken. This phenomenology is incorporated into the model by adding fermion mass terms to the Lagrangian

$$\mathcal{L}_M = m_u^0 \bar{\psi}_u \psi_u + m_d^0 \bar{\psi}_d \psi_d + m_s^0 \bar{\psi}_s \psi_s + m_c^0 \bar{\psi}_c \psi_c \, . \tag{6}$$

As is the case for the coupling constant, these bare masses are tuned such that the pseudoscalar meson spectrum ($m_\pi = 139\ MeV$, $m_K = 494\ MeV$, $m_D = 1.87\ GeV$) as well as the effective quark masses[11] (here $M_u = M_d = 340\ MeV$, $M_s = 540\ MeV$ and $M_c = 1800\ MeV$ were used) are reproduced.

Since the GN model is 1+1 dimensional there are no rotations, hence no notion of spin, in 1+1 dimensions we will restrict ourselves in this section to the unpolarized structure functions. Note that the GN model does not confine the constituent fermions (which we will call quarks in the following). This allows us to simplify the discussion by considering the meson cloud around a single constituent quark instead of the meson cloud around a nucleon.[16] Furthermore we will perform a $1/N_c$-expansion and evaluate the structure functions only to first order in $1/N_c$ to which the quark propagator is modified by tadpole type graphs (Fig. 1) as well as virtual emission of bubble chains (Fig. 2). Only the latter contribute sea quarks to the structure functions yielding for the wavefunctions,[#10]

$$\psi_{s\bar{s}u} = \frac{1}{N_c} \times \frac{D_K(q^2)\left(M_u - \frac{M_s}{x}\right)\left(\frac{M_s}{y} + \frac{M_u}{1-x-y}\right) - D_\kappa(q^2)\left(M_u + \frac{M_s}{x}\right)\left(\frac{M_s}{y} - \frac{M_u}{1-x-y}\right)}{M_u^2 - \frac{M_s^2}{x} - \frac{M_s^2}{y} - \frac{M_u^2}{1-x-y}} \, . \tag{7}$$

#10 It is convenient to replace the chain of $\bar{Q}u$ pairs by an effective meson propagator. We should emphasize that this is a mere rewriting of the sum of $\mathcal{O}(1/N_c)$ diagrams and not an approximation.

There

$$q^2 = (1-x)\left(M_u^2 - \frac{M_s^2}{x}\right) \tag{8}$$

and $D_K$, $D_\kappa$ are effective meson propagators

$$\begin{aligned} D_K^{-1}(q^2) &= \left[(M_u - M_s)^2 - q^2\right]\ B(M_u^2, M_s^2, q^2) \\ &- \left[(M_u - M_s)^2 - \mu_K^2\right]\ B(M_u^2, M_s^2, \mu_K^2) \end{aligned} \tag{9}$$

$$\begin{aligned} D_\kappa^{-1}(q^2) &= \left[(M_u + M_s)^2 - q^2\right]\ B(M_u^2, M_s^2, q^2) \\ &- \left[(M_u - M_s)^2 - \mu_K^2\right]\ B(M_u^2, M_s^2, \mu_K^2) \end{aligned} \tag{10}$$

with

$$B(M_u^2, M_s^2, q^2) = \int_0^1 dx\, \frac{1}{M_u^2 x + M_s^2(1-x) - x(1-x)q^2}. \tag{11}$$

One now evaluates the structure functions, from the defining equation

$$s(x) = N_c \int_0^{1-x} dy\, |\psi_{s\bar{s}u}(x,y)|^2 \qquad \bar{s}(y) = N_c \int_0^{1-y} dx\ |\psi_{s\bar{s}u}(x,y)|^2\ . \tag{12}$$

Typical numerical results are displayed in Fig. 3, where also contributions from light ($d\bar{d}$) and heavy ($c\bar{c}$) quarks are shown. Note that, although the charge asymmetry decreases as we go from light to heavy quarks,[#11] there is still a sizeable effect — even for $c$-quarks.

At this point one might be tempted to extract the contribution from the pion pole from the full calculation (7) - (12) , but this would go beyond the scope of this work.

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#11 Aas we expect since the splitting between the pseudoscalar meson and its first excitation decreases

## Charge Asymmetry of the Strange Sea in the 3+1 dimensional Gross-Neveu Model

The simple kaon cloud picture, presented in the introduction, suggested already some charge asymmetry of the spin distribution assotiated with strange quarks in a proton. In order to go beyond heuristic arguments we have to study a microscopic model. Since there are no rotations, hence no spin, in $1+1$ dimensions we have to proceed to a $3+1$ dimensional model. A simple case, which has been quite helpful in understanding the implications of chiral symmetry and chiral symmetry breaking is due to Nambu and Jona-Lasinio[13] . For $N_c = 1$ it can be considered a $3+1$ dimensional generalization of the chiral Gross-Neveu model (5) . However, since we will perform a $\frac{1}{N_c}$ expansion, those models are not identical, though similar in a random phase approximation. Since it is non-renormalizable one has to work with a fixed cut-off that is typically taken at the order of $1.0GeV$.[7] To leading order in $\frac{1}{N_c}$, to which we will restrict ourselves, the general features are rather similar to the $1+1$-dimensional GN model and we omit the details. One finds

$$\psi_{s\bar{s}u}^{\lambda_1\lambda_2\lambda_3}(x,q_\perp,y,k_\perp) = \frac{1}{N_c} \times \frac{D_K(q^2)\times T^{ps}_{\lambda_1\lambda_2\lambda_3} + D_\kappa(q^2)\times T^{s}_{\lambda_1\lambda_2\lambda_3}}{M_u^2 - \frac{M_s^2}{x} - \frac{q_\perp^2}{x(1-x)} - \frac{M_s^2+\hat{k}_\perp^2}{y} - \frac{M_u^2+\hat{k}_\perp^2}{1-x-y}}, \tag{13}$$

where $\hat{k}_\perp = k_\perp + \frac{y}{1-x}q_\perp$ and

$$q^2 = M_u^2(1-x) + M_s^2(1-\frac{1}{x}) - \frac{q_\perp^2}{x}. \tag{14}$$

The effective meson propagators, $D_K(q^2)$ and $D_\kappa(q^2)$, as well the helicity amplitudes $T^s$ and $T^{ps}$, are given in the appendix. Again, the structure functions are given by integrating the amplitudes squared, e.g.

$$s\uparrow(x) = \sum_{\lambda_2,\lambda_3} \int \frac{d^2 q_\perp}{(2\pi)^2} \int \frac{d^2 k_\perp}{(2\pi)^2} \int_0^{1-x} dy \left| \psi_{s\bar{s}u}^{\uparrow\lambda_2\lambda_3}(x, q_\perp, y, k_\perp) \right|^2 . \qquad (15)$$

has to be kept finite.

Eq.(15) contains two logarithmic divergences: one from integrating over the internal $u\bar{s}$-loop momentum and the other one from integrating over the kaon momentum. The cutoffs which we used[#12] were a euclidean momentum cutoff on the kaon line, i.e.

$$\Theta(\Lambda_K^2 - |q^2|) \qquad (16)$$

and the Brodsky-Lepage cutoff[17] for the internal $u\bar{s}$-loop,[#13] i.e.

$$\Theta\left(\frac{P_{K\perp}^2 + \Lambda_{lc}^2}{P_K^+} - \frac{P_{u\perp}^2 + M_u^2}{P_u^+} - \frac{P_{\bar{s}\perp}^2 + M_{\bar{s}}^2}{P_{\bar{s}}^+}\right), \qquad (17)$$

which is invariant under all kinematic transformations in the light-cone formalism. Before applying the latter cutoff we should be careful about which value of $\Lambda_{cov}^2$ to choose. The light-cone cutoff essentially gives the restriction: $\lambda_{lc}^2 > \frac{k_\perp^2}{x(1-x)}$, where $x$ and $(1-x)$ are the light-cone momentum fraction carried respectively by the $u$ and the $\bar{s}$. But $\frac{1}{x(1-x)} \geq 4$, so if the typical transverse momentum is $k_\perp^2 \approx 0.5-1(GeV)^2$ then we must choose an appropriately larger value $2(GeV)^2 \leq \Lambda_{cov}^2 \leq 6(GeV)^2$. As far as the numerical value for $\Lambda_K^2$ is concerned we are bounded by the Landau pole from above and by a "typical mass scale" (e.g. the pseudoscalar meson masses)

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#12 For technical reasons we prefered to work with a cutoff procedure that is easy to implement, once one has performed the light-cone quantization.

#13 Alternatively we also used a simple cutoff only on the transverse component of the internal loop, which provided qualitatively similar results.

from below This leaves us typically the freedom to choose any value $1(GeV)^2 \leq \Lambda_K^2 \leq 2(GeV)^2$.[#14]

The following qualitative results turned out to be cutoff independent: $s$-quarks carry more momentum that $\bar{s}$. The polarization of the $s$-quarks is mostly anti-parallel to the initial $u$-spin, whereas the $\bar{s}$ are polarized parallel to the $u$-spin. However, it was not possible to make an unambigous statement about the *net* polarization of the strange sea.

Typical structure functions for $s$ and $\bar{s}$ quarks around a u quark are shown in Fig.4.

The interpretation of the unpolarized distributions is the same as in the $1+1$-dimensional model. The strong negative polarization of the $s$ quarks at large $x_{bj}$ can be understood if one assumes that the kaon dominates the meson cloud. However, this approximation is too crude to understand the (positive) polarization of the $\bar{s}$. Here one has to take the interference between scalar and pseudoscalar degrees of freedom into account.

Numerically it turns out that most of the $\bar{s}$ polarization arises from a region with relatively large (compared to the effective quark masses) perpendicular momenta — a kinematic region where the chirally broken and unbroken phases look quite similar. Thus, in order to simplify the argument, let us assume for the moment that we are in the chirally unbroken phase, i.e. that the quarks are massless and that the effective interaction in the scalar and pseudoscalar channels are of equal strength.[#15] With these assumptions, and the helicity amplitudes listed in

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#14 See Appendix

#15 Actually, for the large $k_\perp$ ($k_\perp \approx 0.5 - 1.0 GeV$) component of the wavefunction these are good approximations.

the Appendix A, it is evident that the polarization pattern in this kinematic region is: ($s =\downarrow$, $\bar{s} =\uparrow$, $u =\downarrow$).

A more intuitive way to understand this result is the following. Since helicity and chirality are the same for massless quarks, we can combine the scalar and pseudoscalar amplitudes (Fig.5). Using again the equivalence of helicity and chirality in this limit, as well as the chirality flip property of $\gamma^0$ and the projection properties of $1 \pm \gamma_5$, one obtains a $s\bar{s}u$-state where the $s$ has negative whereas the $\bar{s}$ and the $u$ have positive helicity. Furthermore, the $\bar{s}$ must have the same helicity as the $u$, since in the rest frame of the $K$ they are flying apart. An infintite momentum boost (to the Breit frame) then tilts the spins to be parallel, as shown in Fig.5.

Though the polarization of the strange quark sea at large $x_{bj}$ is dominated by the negative polarization of the $s$, the situation changes for smaller values of $x_{bj}$, where the negatively polarized $\bar{s}$ dominate and a cancellation in the net polarization is conceivable. Numerically it turns out that the sign of the *net* polarizaion depends on the cutoff — mainly due to uncertainties assotiated with the $\bar{s}$ contribution at small $x_{bj}$. However, one should not take the results at small $x_{bj}$ too seriously, since, in this high virtuality region, one does not expect the GN-model to be a good approximation for $QCD$. In fact, in that region one does not have to rely on toy models, because there perturbative $QCD$ is applicable and yields a good description for the structure functions.

Finally we should emphasize that all angular momentum effects discussed so far (spin of $s$ and $\bar{s}$ as well as the orbital angular momentum of the $\bar{s}$) contribute coherently to the magnetic moment of the dressed quark, thus suggesting a relatively large (positive) contribution of strange quarks on the magnetic moment of the proton.

## SUMMARY

Using only heuristic arguments, we argued that a sizeable charge asymmetry in $s$ and $c$ quark structure functions is conceivable. The idea was mainly based on the existence of light pseudoscalars which arise from spontaneous chiral symmetry breaking in QCD. The main predictions are:

- $s$ quarks carry more momentum than $\bar{s}$ quarks, i.e. $\int dx\, x\, s(x) > \int dx\, x\, \bar{s}(x)$.
- $s$ quarks are polarized antiparallel to the initial proton
- $\bar{s}$ quarks carry parallel polarization
- $s$ and $\bar{s}$ quarks contribute both with a magnetic moment parallel to the proton magnetic moment. In the case of the $s$ the magnetic moment arises from the spin whereas the $\bar{s}$ contributes through spin as well as orbital angular momentum. All effects add up coherently to the magnetic moment of the proton.

The first prediction is mainly a consequence of the strong binding in pseudoscalars — making them much lighter than the sum of the effective masses of the valence quarks they are made of. It could be confirmed in the context of the chiral Gross-Neveu as well as the Nambu–Jona-Lasinio model. We should emphasize that, since the integrated sea quark structure functions are dominated by the extrinsic component, the total probability to find intrinsic $s\bar{s}$ pairs might be small. Therefore, even if there is a significant charge asymmetry at large $x_{bj}$, the effect on the total momentum fraction carried by $s$ and $\bar{s}$ can be small.

The second prediction follows also from the pseudoscalar dominance in the meson cloud around a quarks plus angular momentum conservation.[#16] To understand the third prediction is more difficult since it arises as an interference effect between the scalar and the pseudoscalar component in the meson cloud around a quark. Finally the last prediction emerged trivially from the second and third one.

The spin of the $\Lambda$, which is the lightest excited state of the nucleon with strangeness, is carried by its (valence-) $s$ quark. Thus, although above results deal with the $s\bar{s}$ cloud around quarks they should be qualitatively generalizable to nucleons.

An important consequence would be that the usually assumed charge symmetry of the $s\bar{s}$ sea around a nucleon could no longer be used to extract the $s$ distribution from the $\bar{s}$ distribution in di-muon deep inelastic scattering events.[18] We would thus suggest to test this assumption in the large $x_{bj}$ region ($x_{bj} \gtrsim 0.2$) by measuring the $s$ and $\bar{s}$ distributions independently — for example by combining $F_1$ and $F_3$ measurements from charged current $\nu$ and $\bar{\nu}$ scattering experiments on protons and neutrons.[19]

## APPENDIX

Summing the corresponding bubble chains in $3+1$ dimensions gives inverse propagators for the mesons, $D^{-1}(q^2)$, which are logarithmically divergent. These divergences are then to be removed by mass and kinetic energy counterterms, whilst $\frac{d^2}{(dq^2)^2} D^{-1}(q^2)$ remains finite and umambigous. We have for the kaon

$$D_K^{-1}(q^2) = c_1 - c_2 q^2 + \left[(M_u - M_s)^2 - q^2\right] \; B(M_u^2, M_s^2, q^2) \tag{A.1}$$

#16 Note that the explicit GN calculation reproduces the arguments from the kaon cloud picture. Hence it allows indeed enough time for the pseudoscalars to be formed (see footnote 4).

where

$$B(M_u^2, M_s^2, q^2) = \frac{1}{16\pi^2}\int_0^1 dx \ \log\left(x + \frac{M_s^2}{M_u^2}(1-x) - x(1-x)\frac{q^2}{M_u^2}\right) \qquad \text{(A.2)}$$

(in a normalization where the baryon-meson coupling is 1). The constants $c_1$ and $c_2$ can then be fit from the physical values for $M_K$ and $f_K$. At this point it appears that the cutoff has disappeared! However, the cutoff $\Lambda_K$ (16) remains implicit in that the meson propagators have Landau poles, and one must choose $\Lambda_K \leq 2GeV$, (where $\Lambda_K$ is euclidean invariant). The propagator for the "scalar kaon partner" is now fixed by chiral symmetry,

$$D_\kappa^{-1}(q^2) = c_1 - c_2 q^2 + \left[(M_u + M_s)^2 - q^2\right] \ B(M_u^2, M_s^2, q^2). \qquad \text{(A.3)}$$

The computation of the spinor matrix elements is straightforward, using the various helicity amplitudes and conventions in Ref. 17 . One finds

$$\begin{aligned}
T^s_{\uparrow\uparrow\uparrow} &= -\left(\frac{M_s}{x} + M_u\right)\hat{k}^*\left(\frac{1}{y} + \frac{1}{1-x-y}\right) \\
T^s_{\uparrow\uparrow\downarrow} &= \left(\frac{M_s}{x} + M_u\right)\left(\frac{M_u}{1-x-y} - \frac{M_s}{y}\right) \\
T^s_{\uparrow\downarrow\uparrow} &= T^s_{\uparrow\uparrow\downarrow}, \quad T^s_{\uparrow\downarrow\downarrow} = -\left(T^s_{\uparrow\uparrow\uparrow}\right)^* \\
T^s_{\downarrow\uparrow\uparrow} &= \frac{q}{x}\hat{k}^*\left(\frac{1}{y} + \frac{1}{1-x-y}\right) \\
T^s_{\downarrow\uparrow\downarrow} &= \frac{q}{x}\left(\frac{M_s}{y} - \frac{M_u}{1-x-y}\right) \\
T^s_{\downarrow\downarrow\uparrow} &= T^s_{\downarrow\uparrow\downarrow}, \quad T^s_{\downarrow\downarrow\downarrow} = -\left(T^s_{\downarrow\uparrow\uparrow}\right)^* \ ,
\end{aligned} \qquad \text{(A.4)}$$

where $\hat{k} = \hat{k}_x + i\hat{k}_y$, $q = q_x + iq_y$. Similarly

$$
\begin{aligned}
T^{ps}_{\uparrow\uparrow\uparrow} &= -\left(\frac{M_s}{x} - M_u\right)\hat{k}^*\left(\frac{1}{y} + \frac{1}{1-x-y}\right) \\
T^{ps}_{\uparrow\uparrow\downarrow} &= \left(\frac{M_s}{x} - M_u\right)\left(\frac{M_u}{1-x-y} + \frac{M_s}{y}\right) \\
T^{ps}_{\uparrow\downarrow\uparrow} &= -T^{ps}_{\uparrow\uparrow\downarrow}\,, \quad T^{ps}_{\uparrow\downarrow\downarrow} = \left(T^{ps}_{\uparrow\uparrow\uparrow}\right)^* \\
T^{ps}_{\downarrow\uparrow\uparrow} &= \frac{q}{x}\hat{k}^*\left(\frac{1}{y} + \frac{1}{1-x-y}\right) \\
T^{ps}_{\downarrow\uparrow\downarrow} &= \frac{q}{x}\left(\frac{M_s}{y} + \frac{M_u}{1-x-y}\right) \\
T^{ps}_{\downarrow\downarrow\uparrow} &= -T^{ps}_{\downarrow\uparrow\downarrow}\,, \quad T^{ps}_{\downarrow\downarrow\downarrow} = \left(T^{ps}_{\downarrow\uparrow\uparrow}\right)^*\ .
\end{aligned}
\tag{A.5}
$$

Note that $T^{ps}_{\downarrow\uparrow\uparrow} = T^{s}_{\downarrow\uparrow\uparrow}$, whereas $T^{ps}_{\downarrow\downarrow\downarrow} = -T^{s}_{\downarrow\downarrow\downarrow}$. Thus, there is constructive interference between $T^s$ and $T^{ps}$ for $T_{\downarrow\uparrow\uparrow}$ but destructive interference for $T_{\downarrow\downarrow\downarrow}$ in the region where $|q|, |k| \gg M_u, M_s$.

## FIGURE CAPTIONS

1) Typical $\mathcal{O}(N_c^0)$ contributions to the propagator of a $u$-quark.

2) Typical $\mathcal{O}(N_c^{-1})$ contribution to the $u$-quark propagator.

3) Numerical results for the sea quark structure functions in the Gross-Neveu model.

4) Leading order $\frac{1}{N_c}$ numerical results for the polarized $s$ and $\bar{s}$ distributions in the meson cloud of a $u$-quark in the $3+1$ dimensional Gross-Neveu model. The parameters $M_u = 360MeV$, $M_s = 500MeV$ and a Brodsky-Lepage cutoff $\Lambda_{lc}^2$ of $4(GeV)^2$ as well as a euclidean cutoff $\Lambda_K^2$ of $1(GeV)^2$ have been used.

5) Graphical representation of the sum of the scalar and pseudoscalar amplitudes for massless quarks.

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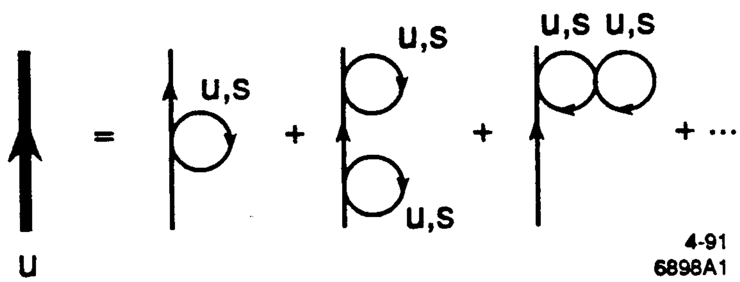



Fig. 1

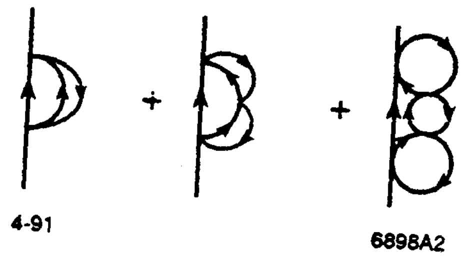



Fig. 2

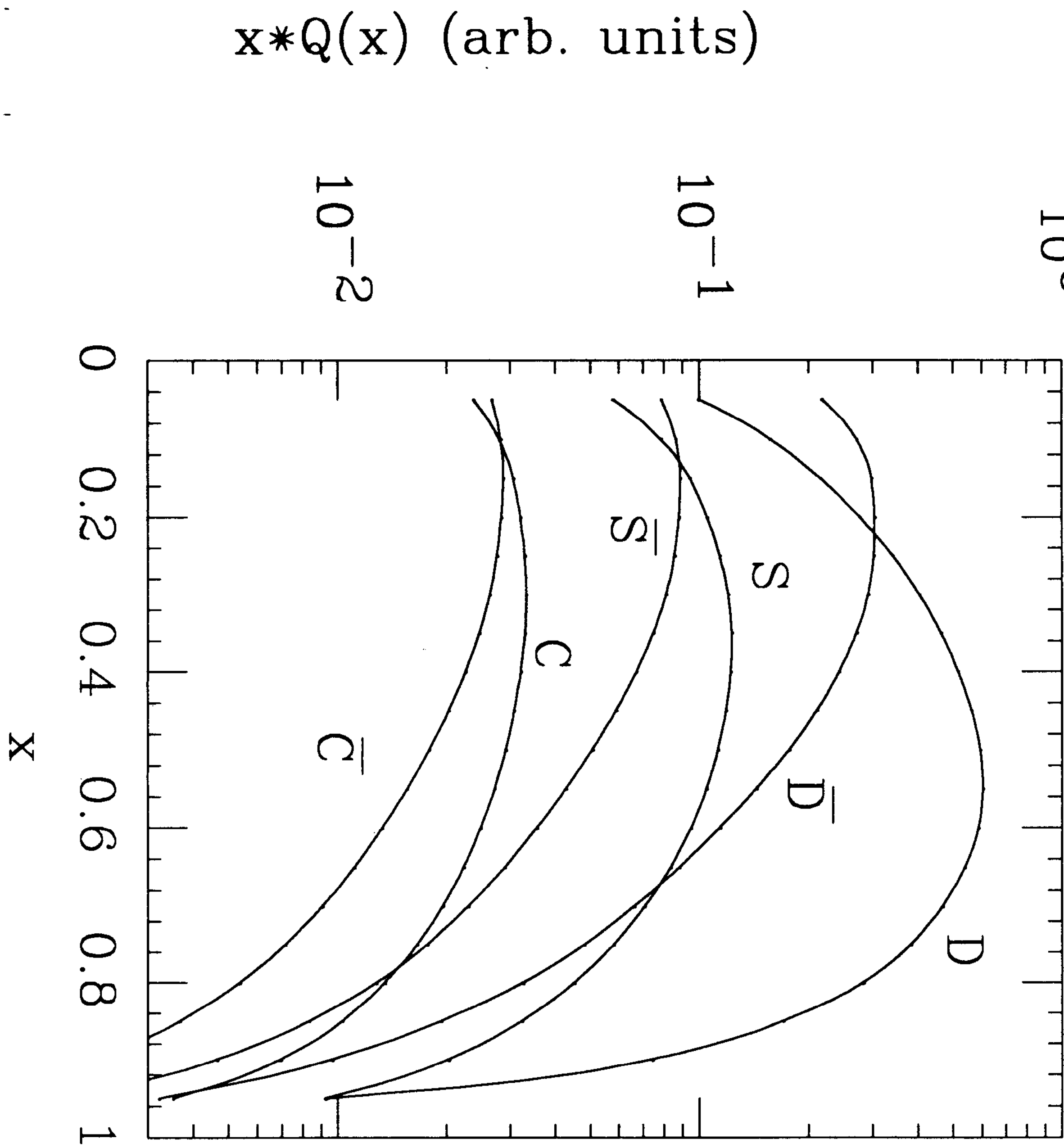
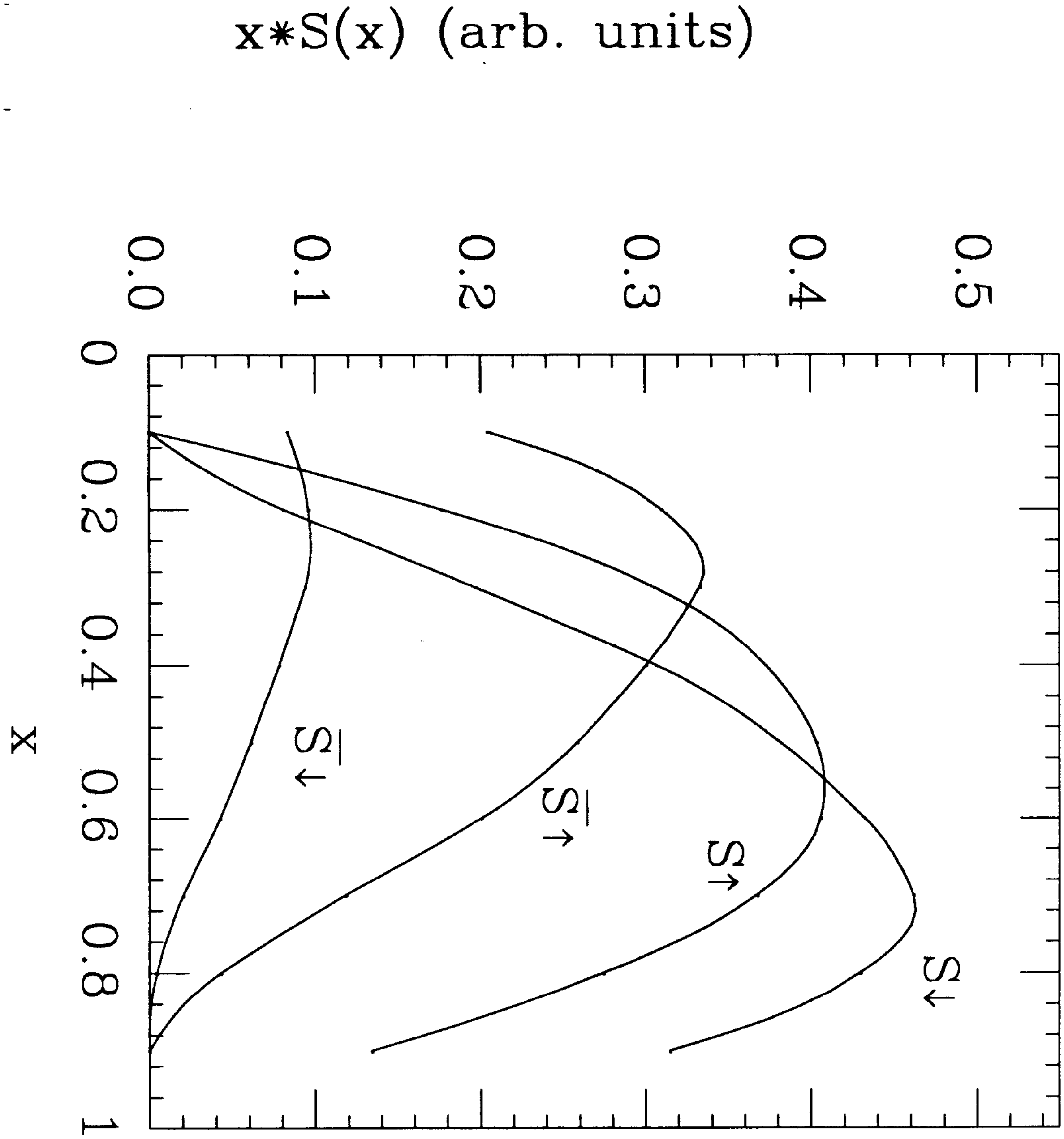



Fig. 3



x*S(x) (arb. units)
0.0
0.1
0.2
0.3
0.4
0.5
x
0
0.2
0.4
0.6
0.8
1
$\bar{S}_\downarrow$
$\bar{S}_\uparrow$
$S_\uparrow$
$S_\downarrow$


Fig. 4

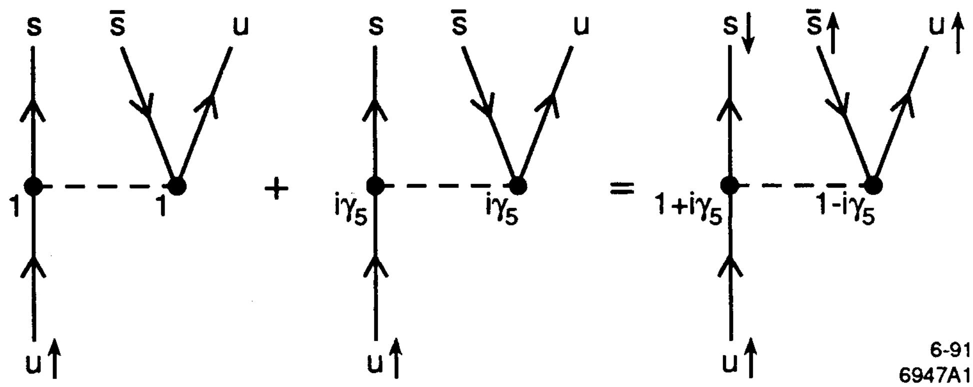



Fig. 5