Chiral symmetry and the charge asymmetry of the $s\bar{s}$ distribution in the proton

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Recently it has been suggested that a sizeable fraction of the strange and charm quarks in a nucleon—the so-called "intrinsic strangeness of charm"—have momentum distributions which extend to large x_{bj}. This effect is enhanced if these virtual heavy quarks live long enough so 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 sizeable charge asymmetry of the corresponding spin and momentum distributions.

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I. INTRODUCTION

Most virtual $s\bar{s}$ pairs[1] in a proton have a very short lifetime (of the order $\tau \sim 1/\sqrt{-q^2}$, 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 [1]. 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 the fact 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 [3,4]. Of course, the initial 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 [3]. Thus, small values of x_{bj}—in particular for heavy quarks—are suppressed 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 ΛK^+ 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 consequences for some spin observables are then clear. Angular momentum and parity conservation require the K^+ to be emitted in an $l=1$ state, and the total angular momentum wave function reads

$$|J=\tfrac{1}{2},J_z=\tfrac{1}{2}\rangle = (\sqrt{2}|l=1,l_z=1\rangle|s=\tfrac{1}{2},s_z=-\tfrac{1}{2}\rangle - |l=1,l_z=0\rangle|s=\tfrac{1}{2},s_z=\tfrac{1}{2}\rangle)\frac{1}{\sqrt{3}} . \tag{2}$$

In a constituent-quark model the Λ spin is carried by its s

[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 $M_{Q\bar{Q}}^{-1}$ [2].

[3]This indicates already that perturbative QCD is not appropriate to describe this component of the proton wave function, 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.

quark. Thus, finding the s quark with polarization antiparallel to the initial proton spin is most likely.[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 positively charged $\bar{s}$ has orbital angular momentum parallel to the proton spin.

The binding of quarks in pseudoscalar mesons is (due to chiral symmetry) usually stronger than in baryons. This has striking consequences for the (unpolarized) momentum distributions [5]. In order to see this let us assume that the momentum in the ΛK system is shared such that the light-cone kinetic energy is minimized [4], i.e.,[6]

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

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

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

(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

$$\langle x_s \rangle / \langle x_{\bar{s}} \rangle = \left[\frac{m_\Lambda}{M_s + 2M_u} \right] \Big/ \left[\frac{m_K}{M_s + M_u} \right] > 1 ,$$

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 as 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 [7]. The most important point here is 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 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 [3] and, like chiral symmetry, the color-hyperfine splitting yields comparatively light pseudoscalars. But the effect is relatively small, which results in 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 developing 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 wave function [8,9].

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 [7].

II. CHARGE ASYMMETRY OF THE STRANGE SEA IN THE CHIRAL GROSS-NEVEU MODEL

In order to avoid the above-mentioned difficulties we start by studying a chirally symmetric generalization of the $(1+1)$-dimensional Gross-Neveu (GN) model [11] 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 de-

[5]Strictly speaking, the 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.

[6]The momentum fractions computed here are momentum fractions in the ΛK^+ 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 of finding intrinsic sea quarks, i.e., by the probability that the proton is in a virtual ΛK state.

[7]Although perturbative QCD predicts the same scaling power p for $s(x)$ and $\bar{s}(x)$ as $x \to 1$ [16], 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.

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

grees 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],

$$\mathcal{L}_0=i\bar{\psi}\partial\!\!\!/\psi+\frac{g^2}{2N_c}[(\bar{\psi}\tau^i\psi)^2-(\bar{\psi}\tau^i\gamma_5\psi)^2] , \quad (5)$$

where the quark fields carry both color and flavor [the τ^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 nonzero vacuum expectation value, and hence, an effective mass for the fermions is generated. Now, quarks (in the real world) have nonzero current masses, i.e., chiral symmetry is explicitly broken. This phenomenology is incorporated into the model by adding fermion mass terms to the Lagrangian

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

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 are reproduced [11].

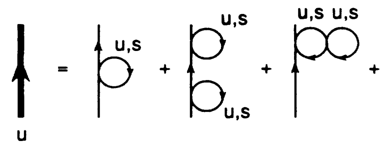

FIG. 1. Typical $O(N_c^0)$ contributions to the u-quark propagator.

(Here we use $M_u=M_d=340$ MeV, $M_s=540$ MeV, and $M_c=1800$ MeV.)

Since the GN model is $1+1$ dimensional there are no rotations; hence, there is no notion of spin, so 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 [12]. Furthermore, we will perform a $1/N_c$ expansion and evaluate the structure functions only to first order in $1/N_c$. To this order the quark propagator is modified by tadpole-type graphs (Fig. 1) as well as virtual emission of bubble chains (Fig. 2).[10] Only the latter contribute sea quarks to the structure functions, yielding for the wave functions

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

There

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

and D_K,D_κ are effective meson propagators:

$$D_K^{-1}(q^2)=[(M_u-M_s)^2-q^2]B(M_u^2,M_s^2,q^2)-[(M_u-M_s)^2-\mu_K^2]B(M_u^2,M_s^2,\mu_K^2) , \quad (9)$$

$$D_\kappa^{-1}(q^2)=[(M_u+M_s)^2-q^2]B(M_u^2,M_s^2,q^2)-[(M_u-M_s)^2-\mu_K^2]B(M_u^2,M_s^2,\mu_K^2) , \quad (10)$$

with

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

One now evaluates the structure functions from the defining equation

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

Typical numerical results are displayed in Fig. 3, where contributions from light ($d\bar{d}$) and heavy ($c\bar{c}$) quarks are also 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.

III. 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

[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.

[11] This is as we expect since the splitting between the pseudoscalar meson and its first excitation decreases.

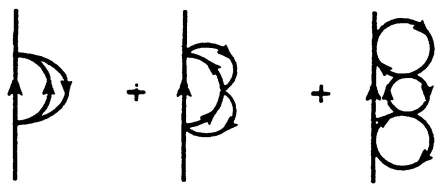

FIG. 2. Typical $O(N_c^{-1})$ contributions to the u-quark propagator.

spin distribution associated 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, and 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 (NJL) [14]. 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 $1/N_c$ expansion, those models are not identical, though they are similar in a random-phase approximation. Since the NJL model is nonrenormalizable,

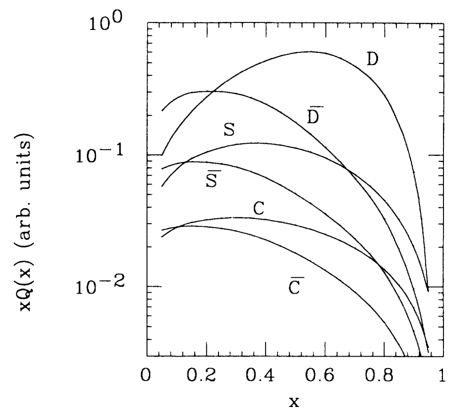

FIG. 3. Numerical results for the sea-quark structure functions in the Gross-Neveu model.

one has to work with a fixed cutoff that is typically taken at the order of 1 GeV [15]. To leading order in $1/N_c$, 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

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

where $\hat{k}_\perp=k_\perp+[y/(1-x)]q_\perp$ and

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

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

$$s^\uparrow(x)=\sum_{\lambda_2,\lambda_3}\int\frac{d^2q_\perp}{(2\pi)^2}\int\frac{d^2k_\perp}{(2\pi)^2}\int_0^{1-x}dy\,|\psi_u^{\uparrow\lambda_2\lambda_3}(x,q_\perp,y,k_\perp)|^2 \tag{15}$$

has to be kept finite.

Equation (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.,

$$\Theta(\Lambda_K^2-|q^2|)\ , \tag{16}$$

and the Brodsky-Lepage cutoff [16] for the internal $u\bar{s}$ loop,[13] i.e.,

$$\Theta\left[\frac{P_{K\perp}^2+\Lambda_{lc}^2}{P_K^+}-\frac{P_{u\perp}^2+M_u^2}{P_u^+}-\frac{P_{s\perp}^2+M_s^2}{P_s^+}\right]\ , \tag{17}$$

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 Λ_{lc}^2 to choose. The light-cone cutoff essentially gives the restriction $\Lambda_{lc}^2>k_\perp^2/[x(1-x)]$, where x and $1-x$ are the light-cone momentum fraction carried, respectively, by the u and the $\bar{s}$. But $1/[x(1-x)]\geq 4$, 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$^2\leq\Lambda_{lc}^2\leq 6$ GeV^2. As far as the numerical value for Λ_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) from below. This leaves us typically the freedom to choose any value 1 GeV$^2\leq\Lambda_K^2\leq 2$ GeV^2, as explained

[12]For technical reasons we preferred 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.

in the Appendix.

The following qualitative results turned out to be cutoff independent: s quarks carry more momentum than $\bar{s}$. The polarization of the s quarks is mostly antiparallel to the initial u spin, whereas the $\bar{s}$ are polarized parallel to the u spin. However, it was not possible to make an unambiguous 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 interactions in the scalar and pseudoscalar channels are of equal strength.[14] With these assumptions, and the helicity amplitudes listed in the Appendix, 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 γ^0 and the projection properties of $1\pm\gamma_5$, one obtains a $s\bar{s}u$ state where the s has negative and 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 infinite-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* polarization depends on the cutoff—mainly due to uncertainties associated 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 to the magnetic moment of the proton.

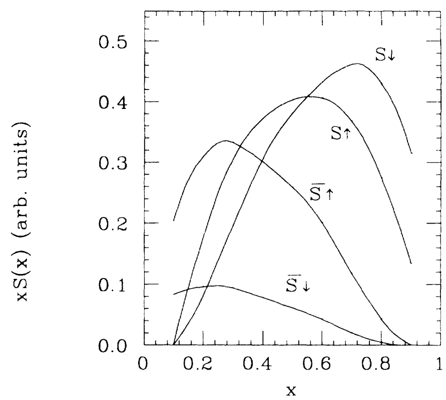

FIG. 4. Leading-order $1/N_c$ numerical results for the polarized s and $\bar{s}$ distributions in the meson cloud of a u quark in the $(3+1)$-dimensional Gross-Neveu model. The parameters $M_u=360$ MeV and $M_s=500$ MeV, and a Brodsky-Lepage cutoff Λ_{lc}^2 of 4 GeV^2 as well as a Euclidean cutoff Λ_K^2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

IV. 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 as follows.

(1) s quarks carry more momentum than $\bar{s}$ quarks, i.e., $\int dx\, xs(x) > \int dx\, x\bar{s}(x)$.

(2) s quarks are polarized antiparallel to the initial proton.

(3) $\bar{s}$ quarks carry parallel polarization.

(4) s and $\bar{s}$ both contribute 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 [17].

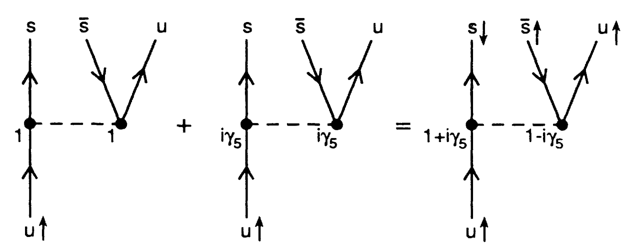

FIG. 5. Graphical representation of the sum of the scalar and pseudoscalar amplitudes for massless quarks.

[14]Actually, for the large-$k_\perp$ ($k_\perp \approx 0.5-1.0$ GeV) component of the wave function these are good approximations.

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 models. We should emphasize that, since the integrated sea-quark structure functions are dominated by the extrinsic component, the total probability of finding 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 quark and from angular momentum conservation.[15] To understand the third prediction is more difficult since it arises as an interference effect between the scalar and the pseudoscalar components in the meson cloud around a quark. Finally, the last prediction emerged trivially from the second and third ones.

The spin of the Λ, which is the lightest excited state of the nucleon with strangeness, is carried by its (valence) s quark. Thus, although the 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 dimuon deep-inelastic scattering events [18]. We would thus suggest testing 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 ν and $\bar{\nu}$ scattering experiments on protons and neutrons [19].

ACKNOWLEDGMENTS

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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, while $d^2D^{-1}(q^2)/(dq^2)^2$ remains finite and unambiguous. We have for the kaon,

$$D_K^{-1}(q^2)=c_1-c_2q^2+[(M_u-M_s)^2-q^2]B(M_u^2,M_s^2,q^2)\ , \tag{A1}$$

where

$$B(M_u^2,M_s^2,q^2)=\frac{1}{16\pi^2}\int_0^1 dx\,\ln\left[x+\frac{M_s^2}{M_u^2}(1-x)-x(1-x)\frac{q^2}{M_u^2}\right] \tag{A2}$$

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

$$D_\kappa^{-1}(q^2)=c_1-c_2q^2+[(M_u+M_s)^2-q^2]B(M_u^2,M_s^2,q^2)\ . \tag{A3}$$

The computation of the spinor matrix elements is straightforward, using the various helicity amplitudes and conventions in Ref. [16]. One finds

$$\begin{aligned}
T^s_{\uparrow\uparrow\uparrow}&=-\left[\frac{M_s}{x}+M_u\right]\hat{k}^*\left[\frac{1}{y}+\frac{1}{1-x-y}\right],\\
T^s_{\uparrow\uparrow\downarrow}&=\left[\frac{M_s}{x}+M_u\right]\left[\frac{M_u}{1-x-y}-\frac{M_s}{y}\right],\\
T^s_{\downarrow\uparrow\uparrow}&=T^s_{\downarrow\uparrow\downarrow},\quad T^s_{\uparrow\downarrow\downarrow}=-(T^s_{\uparrow\uparrow\uparrow})^*\ ,\\
T^s_{\downarrow\uparrow\uparrow}&=\frac{q}{x}\hat{k}^*\left[\frac{1}{y}+\frac{1}{1-x-y}\right],\\
T^s_{\downarrow\uparrow\downarrow}&=\frac{q}{x}\left[\frac{M_s}{y}-\frac{M_u}{1-x-y}\right],\\
T^s_{\downarrow\downarrow\uparrow}&=T^s_{\downarrow\uparrow\downarrow},\quad T^s_{\downarrow\downarrow\downarrow}=-(T^s_{\downarrow\uparrow\uparrow})^*\ ,
\end{aligned} \tag{A4}$$

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

$$\begin{aligned}
T^{ps}_{\uparrow\uparrow\uparrow}&=-\left[\frac{M_s}{x}-M_u\right]\hat{k}^*\left[\frac{1}{y}+\frac{1}{1-x-y}\right],\\
T^{ps}_{\uparrow\uparrow\downarrow}&=\left[\frac{M_s}{x}-M_u\right]\left[\frac{M_u}{1-x-y}+\frac{M_s}{y}\right],\\
T^{ps}_{\uparrow\downarrow\uparrow}&=-T^{ps}_{\uparrow\uparrow\downarrow},\quad T^{ps}_{\uparrow\downarrow\downarrow}=(T^{ps}_{\uparrow\uparrow\uparrow})^*\ ,\\
T^{ps}_{\downarrow\uparrow\uparrow}&=\frac{q}{x}\hat{k}^*\left[\frac{1}{y}+\frac{1}{1-x-y}\right],\\
T^{ps}_{\downarrow\uparrow\downarrow}&=\frac{q}{x}\left[\frac{M_s}{y}+\frac{M_u}{1-x-y}\right],\\
T^{ps}_{\downarrow\downarrow\uparrow}&=-T^{ps}_{\downarrow\uparrow\downarrow},\quad T^{ps}_{\downarrow\downarrow\downarrow}=(T^{ps}_{\downarrow\uparrow\uparrow})^*\ .
\end{aligned} \tag{A5}$$

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 $|q|,|k|\gg M_u,M_s$.

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

[1] V. N. Gribov and L. N. Lipatov, Yad. Fiz. **15**, 781 (1972) [Sov. J. Nucl. Phys. **15**, 438 (1972)]; **15**, 1218 (1972) [**15**, 675 (1970)]; G. Altarelli and G. Parisi, Nucl. Phys. **B126**, 298 (1977).

[2] S. J. Brodsky and I. A. Schmidt, Phys. Rev. D **43**, 179 (1991).

[3] S. J. Brodsky and P. Hoyer, presented at the Second European Workshop on Hadronic Physics with Electrons Beyond 10 GeV, Dourdan, France, 1990 SLAC-PUB-5422 (unpublished).

[4] S. J. Brodsky, P. Hoyer, C. Peterson, and N. Sakai, Phys. Lett. **93B**, 451 (1980); S. J. Brodsky, C. Peterson, and N. Sakai, Phys. Rev. D **23**, 2745 (1981).

[5] F. E. Close and A. W. Thomas, Phys. Lett. B **212**, 27 (1961).

[6] B. L. Ioffe and M. Karliner, Phys. Lett. B **247**, 387 (1990); M. Karliner, presented at QCD '90, Montpellier, France, 1990 (unpublished).

[7] A. W. Thomas, Nucl. Phys. **A518**, 186 (1990).

[8] A. I. Signal and A. W. Thomas, Phys. Lett. B **191**, 205 (1987).

[9] J. D. Sullivan, Phys. Rev. D **5**, 186 (1970); E. L. Berger and F. Coester, *ibid.* **32**, 1071 (1985); R. P. Bickerstaff and A. W. Thomas, J. Phys. G **15**, 1523 (1989); L. L. Frankfurt, L. Mankiewicz, and M. I. Strikman, Z. Phys. A **33**, 343 (1989); H. Jung and G. A. Miller, Phys. Rev. C **41**, 659 (1990).

[10] D. J. Gross and A. Neveu, Phys. Rev. D **10**, 3235 (1974); B. Rosenstein, B. J. Warr, and S. H. Park, Phys. Rep. **205**, 60 (1991).

[11] M. D. Scadron, Rep. Prog. Phys. **44**, 214 (1981).

[12] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).

[13] M. Burkardt, Phys. Lett. B **268**, 419 (1991).

[14] Y. Nambu and G. Jona-Lasinio, Phys. Rev. **122**, 345 (1961); **124**, 246 (1961).

[15] V. Bernard, R. L. Jaffe, and U.-G. Meissner, Phys. Lett. B **198**, 92 (1987); Nucl. Phys. **B308**, 758 (1988).

[16] S. J. Brodsky and G. P. Lepage, Phys. Rev. D **22**, 2157 (1980).

[17] D. Kaplan and A. Manohar, Nucl. Phys. **B310**, 527 (1988); R. L. Jaffe, Phys. Lett. B **229**, 275 (1989).

[18] H. Abramowicz *et al.*, Z. Phys. C **15**, 19 (1982).

[19] E. Leader and E. Predazzi, *An Introduction to Gauge Theories and the New Physics* (Cambridge University Press, Cambridge, England, 1982).