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Chiral Substructure of the Nucleon*

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The representation mixing induced by the spontaneous breaking of chiral symmetry is a spatially varying phenomenon. A model for this effect predicts the form of the spin-dependent structure functions of deep inelastic lepton scattering and leads to the relation $1-\alpha_{A_1}(0)=2[1-\alpha_\rho(0)]$.

The Nambu-Goldstone realization of chiral symmetry leads naturally[1] to representation mixing of quarks and their bound states. In this Letter we describe this representation mixing in the context of the parton model[2] and point out how this phenomenon is probed in polarized electroproduction experiments. We suggest that the mixing exhibits spatial variation[3] and that the nucleon consequently possesses an interesting chiral substructure. Our description of this structure is summarized in predictions of the spin-dependent structure functions $G(x)$ and a relation between the A_1 and ρ Regge trajectories, $1-\alpha_{A_1}(0)=2[1-\alpha_\rho(0)]$.

Chiral symmetry is spontaneously broken[4] when some chirally noninvariant triggering field acquires a vacuum expectation value v. This phenomenon leads in quark models to the spontaneous generation of a quark mass m. In bound-state models[5] one anticipates that v, and hence m, will be spatially dependent quantities, approaching some asymptotic values in regions far from the localized bound state. This means that the chiral representation mixing which results from the spontaneously generated quark mass will also vary across the region of the bound state. The domain in which the mixing is a significant phenomenon is simply that region in which v is appreciably large.

In the quark model, the composite field $\bar{\psi}\psi$ is the triggering field for spontaneous chiral symmetry breaking, and the vacuum expectation value v is a measure of the density of quark-antiquark pairs. Phenomenologically,[6] this density is large only for small values of the longitudinal momentum fraction x. We expect, therefore, that the effects of chiral representation mixing should be significant only in the small-x region.

A concrete model for this mixing can be constructed if we consider a single valence quark embedded in a sea of gluons and quark-antiquark pairs. In the absence of interactions the valence quark's spin would be a constant of motion and, in the context of some broken SU(6) symmetry scheme, the distribution of quark spins would be completely determined. Interactions with the sea will, however, transfer the valence quark's spin to particles in the sea. We assume that these interactions are of short range in rapidity and discuss them, for simplicity, as if they were local in x. Our results are thus most reliable for small x, which, as we have noted, is the only region where we expect significant mixing to occur.

Let $\sin^2\theta$ denote the probability that the valence quark's spin is altered[7] by its interactions with the sea. If the density of the sea relative to valence quarks is $N(x)$, and if $\mathcal{H}(x)$ denotes the probability of a spin-flip interaction between valence

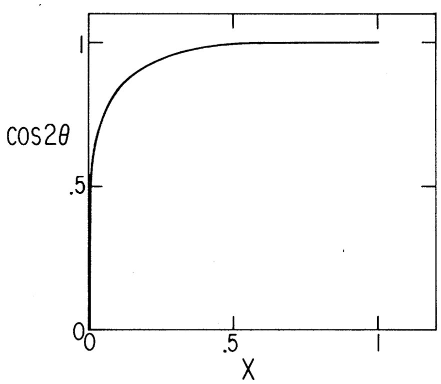

FIG. 1. Spin dilution factor $\cos 2\theta(x)$.

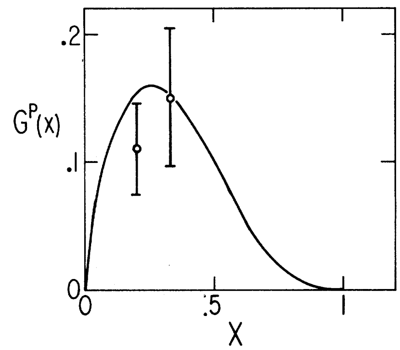

FIG. 2. Predicted spin-dependent structure function for protons. Data points are from Ref. 15.

and sea, then $\sin^2\theta$ has the form $\frac{1}{2}\mathcal{H}N/(\mathcal{H}N+1)$.[8] A measure of the spin dilution induced by these interactions is given by

$$\cos 2\theta(x) = [\mathcal{H}(x)N(x)+1]^{-1}. \qquad (1)$$

Regge arguments suggest that as $x \to 0$, $N(x) \sim x^{\alpha_f - \alpha_P} \sim x^{-1/2}$, where α_f and α_P denote the zero intercepts of the f (or A_2) and Pomeron Regge trajectories, respectively. Near $x=1$, $N(x)$ is quite small as we have already noted. Dimension counting rules[9] suggest a falloff[10] as $(1-x)^2$ for the gluon component of the sea and give us the simple parametrization

$$\mathcal{H}(x)N(x) = \mathcal{H}_0(1-x)^2 x^{-1/2}. \qquad (2)$$

Quark-antiquark pairs are presumably less numerous than gluons and thus provide only a small correction to this expression.

Experimental information on the spin orientation of the nucleon's constituents is provided by the spin-dependent structure functions of deep inelastic scattering. In terms of our parton model, these functions are given simply as the product of a function describing the asymmetries of the valence-quark spins in the absence of interactions with the spin dilution factor $\cos 2\theta$. The relevant valence-quark distributions can be deduced from the unpolarized structure functions in a broken SU(6) model[11] which retains the (spin × isospin) symmetry of the valence-quark wave functions. The structure functions for scattering off longitudinally polarized nucleons thus have the form

$$G^p(x) = \cos 2\theta(x)[\tfrac{4}{9}A_0(x) - \tfrac{2}{27}A_1(x)],$$
$$G^n(x) = \cos 2\theta(x)[\tfrac{1}{9}A_0(x) - \tfrac{1}{9}A_1(x)], \qquad (3)$$

where A_0 and A_1 are the valence-quark distribution functions of Carlitz.[12] Given the functions A_0 and A_1[13] and our parametrization [Eqs. (1) and (2)] of $\cos 2\theta$, the Bjorken sum rule,[14]

$$\int_0^1 (dx/x)[G^p(x) - G^n(x)] = \tfrac{1}{3}|g_A/g_V| \simeq 0.42, \qquad (4)$$

enables us the fix the value of $\mathcal{H}_0$, (namely $\mathcal{H}_0 \simeq 0.052$). The resulting form of $\cos 2\theta$ is shown in Fig. 1. One notes that, as expected, the representation mixing is appreciable only for small x. In the limit $x \to 0$, the angle θ approaches 45°, indicating that valence quarks in the small-x region lose all memory of the spin orientation of the nucleon to which they belong.

In Figs. 2 and 3 we show our predictions for $G^p(x)$ and $G^n(x)$, together with some experimental determinations[15] of these quantities. Our predictions are distinguished from other published work[11,16] by our careful attention to SU(6)-symmetry-breaking effects and, as discussed below, to the proper small-x behavior of the structure functions. The dramatic structure in $G^n(x)$ is due

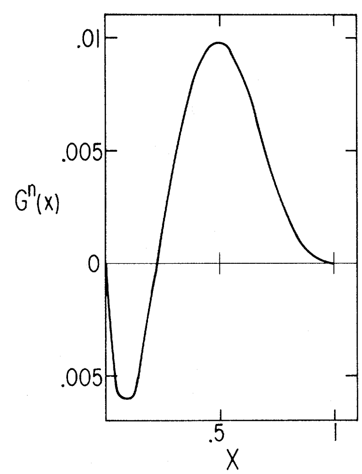

FIG. 3. Predicted spin-dependent structure function for neutrons.

largely to SU(6)-symmetry-breaking effects[12] which give $A_0(x)$ and $A_1(x)$ slightly different shapes. Note, however, that these effects are numerically quite small and that the deviation of $G^n(x)$ from zero will be quite hard to measure.

Near $x=1$, the polarization asymmetries[11] $A^{\gamma p}(x)$ and $A^{\gamma n}(x)$ should both approach unity, while as $x \to 0$ the asymmetries should vanish. The values of $\int G^p dx/x$ and $\int G^n dx/x$ calculated from the curves of Figs. 2 and 3 are 0.41 and -0.005, respectively. This gives[17] an F/D ratio for the axial-vector coupling of 0.65, to be compared with the SU(6) prediction of $\frac{2}{3}$ and an experimental value[18] 0.58 ± 0.03.

The rate at which $G(x)$ vanishes as $x \to 0$ reveals an interesting prediction of our model. The small-x behavior of $A_0(x)$ and $A_1(x)$ is controlled by the exchange-degenerate f and A_2 trajectories and has the form

$$A(x) \underset{x\to 0}{\simeq} x^{1-\alpha_f}. \tag{5}$$

Similar Regge arguments[19] suggest that the A_1 trajectory controls the small-x behavior of $G(x)$ and hence that

$$G(x) \underset{x\to 0}{\simeq} x^{1-\alpha_{A_1}}. \tag{6}$$

Having argued that $\cos 2\theta(x)$ vanishes at small x as $x^{1-\alpha_f}$ we thus have the relation

$$1-\alpha_{A_1} = 2(1-\alpha_f) = 2(1-\alpha_\rho). \tag{7}$$

The last equality follows from the ρ-f exchange degeneracy implicit in our model.

Equation (7) is the Regge-trajectory equivalent of the Weinberg mass relation[20] $m_{A_1}{}^2 = 2m_\rho{}^2$, derived from low-mass saturation of chiral sum rules. Our relation does not require the existence of an A_1 particle, however, since the trajectory which we label as the A_1 could easily correspond to a Regge cut. This interpretation is supported by an application of our arguments to the structure functions for scattering off transversely polarized nucleons. The relevant amplitude is predicted to have the same Regge behavior as $G(x)$. This implies[19] the existence of a (pole or cut) trajectory π_C with the quantum numbers of a "conspiring" pion and an intercept $\alpha_{\pi_C} = \alpha_{A_1}$. If this π_C is interpreted as a Regge cut, then our A_1 must be a cut as well.

The fact that there is no representation mixing for x near 1 implies, via duality arguments,[21] a particularly simple picture of resonance electroproduction at large momentum transfer. No Melosh transformation should be required for the description of these processes—in sharp contrast with the description[22] of resonance photoproduction. Details of this analysis will be given in a separate paper.

In discussing the chiral substructure of the nucleon we have emphasized the x dependence of chiral representation mixing. The quark mass is, in this picture, also an x-dependent quantity —large at small x, and small at large x. If the mass at small x exceeds 1 GeV/c^2 or so, then we would expect that the approach to scaling of the deep inelastic scattering structure functions in that region would be slower than in the large-x region. Some evidence for this type of behavior has been reported[23] in recent experiments at Fermilab.

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Direct Evidence for a New Giant Resonance at $80A^{-1/3}$ MeV in the Lead Region

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Inelastic α scattering experiments at 120 MeV on ^{206,208}Pb, ^{197}Au, and ^{209}Bi reveal the existence of a new isoscalar resonance located at $\sim 80A^{-1/3}$ MeV in addition to the well-known resonance at $\sim 63A^{-1/3}$ MeV. The angular distribution of this new resonance, though not inconsistent with an $E4$ assignment, is better described by $E0$ or $E2$. It would exhaust approximately 100%, 50%, or 17% of the corresponding isoscalar $E0$, $E2$, or $E4$ energy-weighted sum rules, respectively.

The existence of an isoscalar giant resonance at an excitation energy $E_x = 63A^{-1/3}$ MeV has now been well established. The bulk of the experimental data indicate that this resonance is predominantly quadrupole ($E2$) and that it exhausts a substantial fraction of the isoscalar energy-weighted sum rule (EWSR).[1] An important open question is whether nonquadrupole collective strength is present at high excitation energies. Experimental evidence for nonquadrupole strength in the giant resonance region has been found in ^{16}O [2] and ^{28}Si [3] where the detection of such strength is facilitated because the giant quadrupole resonance (GQR) is fragmented. Some indirect evidence for such strength in heavier nuclei comes from an analysis of inelastic electron scattering experiments. Specifically, the existence of a giant monopole in ^{90}Zr and ^{208}Pb was suggested,[4,5] to explain additional strength at $\sim 80A^{-1/3}$ MeV remaining after subtraction of the strongly excited giant dipole resonance (GDR). Furthermore, by comparing spectra obtained from inelastic deuteron and α scattering, Marty *et al.*[6] found evidence for additional isoscalar strength in ^{40}Ca, ^{90}Zr, and ^{208}Pb which they could explain by assuming the existence of an isoscalar monopole (breathing mode) resonance which would be located at the same energy. Recently the presence of $L=3$ ($\sim 15\%$ of the EWSR) and $L=0$ ($\sim 2\%$ of the EWSR) isoscalar strength in the continuum region of ^{208}Pb has been deduced from a high-resolution inelastic-proton-scattering experiment.[7]

In this Letter we present data on inelastic α scattering on ^{208}Pb at $E_\alpha = 120$ MeV. The spectra show that in addition to the well-known giant resonance located at $E_x = 10.9 \pm 0.3$ MeV with a width of 3.0 ± 0.3 MeV, there is another smaller peak located at $E_x = 13.9 \pm 0.3$ MeV with a width of 2.5 ± 0.6 MeV. This comprises the most direct evidence for the existence of a new giant resonance (GR) at $80A^{-1/3}$ MeV for which a full angular distribution has been obtained. The angular distribution of the 13.9-MeV structure is well described by an $L=0$ or 2 transfer but is also not inconsistent with an $L=4$ transfer. The measured cross section indicates that it exhausts approximately