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Spin asymmetries in lepton-hadron scattering

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Asymmetries arising in deep-inelastic scattering of longitudinally polarized electrons and positrons from polarized nucleons are considered. The behavior of these asymmetries arising from electroweak effects is estimated for collider energies. Such measurements will test models for spin structure of the nucleon and the standard model.

I. INTRODUCTION

Spin effects in the perturbative regime ($Q^2 > 1$ GeV^2) of quantum chromodynamics (QCD) have not been adequately tested at high energies. Nevertheless it has long been realized that the present ideas about the underlying theory of elementary particles can be probed sensitively in measurements with polarized beams and (polarized) targets.

Deep-inelastic lepton-nucleon scattering studies have lent impressive support to the quark-parton model of hadronic structure. It may be recalled that electron-nucleon scattering is essentially described by four independent structure functions. Of these four, two relate to the spin structure of the nucleon. A knowledge of spin distributions of partons (quarks) within a nucleon would lead to an understanding of (surprising) spin-dependent high-energy phenomena involving hadrons.[1]

This paper concerns itself with polarized-e–polarized-p scattering at e-p collider energies and asymmetries arising from the weak-electromagnetic interference therein. The study of such double asymmetries at high energies is bound to provide important information about the internal spin structure of the nucleon.[2] In Sec. II the basic formulas involved are tabulated and Sec. III gives a rather brief exposition of two particular models for spin-dependent quark distributions. The parameters involved in predicting the asymmetry are discussed in Sec. IV and Sec. V provides the results.

II. THE FORMALISM

The various electron-proton differential cross sections are well known and are listed below as a function of the relative energy loss of the lepton y:

$$\frac{d^2\sigma_{LL,LR}}{dx\,dy} = \frac{S_x}{4\pi}\sum_i [f_i^{\pm}(x)\sigma^i_{LL} + (1-y)^2 f_i^{\mp}(x)\sigma^i_{LR}] \;, \tag{1}$$

$$\frac{d^2\sigma_{RL,RR}}{dx\,dy} = \frac{S_x}{4\pi}\sum_i [f_i^{\mp}(x)\sigma^i_{RR} + (1-y)^2 f_i^{\pm}(x)\sigma^i_{RL}] \;,$$

where $f_i^{\pm}(x)$ denotes the probability of finding in the proton a quark of flavor i carrying a fraction x of the proton momentum and spin aligned parallel (antiparallel) to the parent spin. The subscript R (L) denotes right-handed (left-handed) helicity of the colliding particles, and σ^i give the cross sections for the pointlike scattering of polarized electrons with polarized quarks of flavor i:

$$\sigma^i_{L(R)L(R)} = \left| \frac{q_e^{\rm EM} q_i^{\rm EM}}{Q^2} + \frac{q^z_{L(R)e} q^z_{L(R)i}}{Q^2 + M_Z^2} \right|^2 , \tag{2}$$

where $q^{\rm EM}$ denotes the electromagnetic charges and the weak charges are

$$q^Z_{L(R)} = \frac{e}{\sin\theta_W \cos\theta_W}(T_{3L(R)} - q^{\rm EM}\sin^2\theta_W) \;, \tag{3}$$

T_{3L} being the third component of weak isospin for left-handed fermions. In the standard model $T_{3R}=0$.

From the above considerations one can build the double asymmetry for scattering electrons off protons (both polarized) as

$$A(e^-p) = \frac{d^2\sigma_{LL} - d^2\sigma_{LR} + d^2\sigma_{RR} - d^2\sigma_{RL}}{d^2\sigma_{LL} + d^2\sigma_{LR} + d^2\sigma_{RR} + d^2\sigma_{RL}}$$
$$= \sum_i \frac{\Delta f_i(x)[\sigma^i_{LL} + \sigma^i_{RR} - (1-y)^2(\sigma^i_{LR} + \sigma^i_{RL})]}{f_i(x)[\sigma^i_{LL} + \sigma^i_{RR} + (1-y)^2(\sigma^i_{LR} + \sigma^i_{RL})]} \;, \tag{4}$$

where

$$\Delta f_i(x) = f_i^+(x) - f_i^-(x) \;,$$
$$f_i(x) = f_i^+(x) + f_i^-(x) \;, \tag{5}$$

and one can define asymmetries at the quark level

$$A_i(x) = \frac{\Delta f_i(x)}{f_i(x)} \tag{6}$$

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III. MODELS FOR SPIN-DEPENDENT DISTRIBUTIONS

A measurement of the double asymmetry $A(e^-p)$ in the range $0.18<x<0.70$ and an average $Q^2 \approx 6.3$ GeV^2 exists[3] (Fig. 1). Further constraints on model building are provided by the Bjorken sum rule for the difference in polarization asymmetry in deep-inelastic scattering of electrons from protons and neutrons.[4] Confining ourselves to the parton picture two models are singled out by the data.

(i) The Carlitz-Kaur model.[5] Carlitz and Kaur take into account the interactions between the valence quarks, $q\bar{q}$ sea, and gluons by incorporating a phenomenological factor. This factor is a measure of the transfer of spin from the valence quarks to gluons and $q\bar{q}$ pairs, and is significant at small x.

In the scaling limit the nucleon's spin-dependent structure function is given by

$$G_1(x)=\tfrac{1}{2}\sum_i q_i^{\mathrm{EM}}\Delta f_i(x) \tag{7}$$

which in the Carlitz-Kaur picture is given as

$$2xG_1^p(x)=\cos 2\theta[\tfrac{4}{9}A_0(x)-\tfrac{2}{27}A_1(x)] \tag{8}$$

for the proton and

$$2xG_1^n(x)=\cos 2\theta[\tfrac{1}{9}(A_0-A_1)] \tag{9}$$

with

$$A_0=xu(x)-\tfrac{1}{2}x\,d(x)\ , \quad A_1=\tfrac{3}{2}x\,d(x) \tag{10}$$

being the contributions from valence quark (u,d) terms in which the noninteracting valence quarks are in a state with isospin 0 and 1, respectively. The spin dilution factor $(\cos 2\theta)$ is argued from Regge theory and large-x behavior $(x\to 1)$ to be

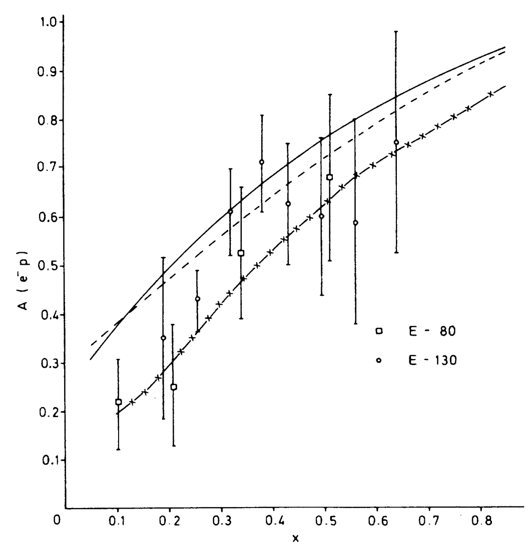

FIG. 1. Available data compared with model predictions. The dash-cross curve is the Carlitz-Kaur model. The dashed curve is for the Callaway-Ellis prediction with $x_0=0.75$ and the full line is with $x_0=0.5$.

$$\cos 2\theta=[H_0x^{-1/2}(1-x)^2+1]^{-1}\ . \tag{11}$$

The constant H_0 is fixed by demanding that Bjorken sum rule

$$\int_0^1 dx[G_1^p(x)-G_1^n(x)]=\frac{g_A}{g_V}$$

be satisfied. The ratio of axial-vector to vector couplings in neutron β-decay (g_A/g_V) is determined by experiment to be ~ 1.254. This analysis yields

$$\begin{aligned}\Delta u(x)&=\cos 2\theta[u(x)-\tfrac{2}{3}d(x)]\ ,\\ \Delta d(x)&=-\tfrac{1}{3}\cos 2\theta\, d(x)\ .\end{aligned} \tag{12}$$

(ii) The Callaway-Ellis model.[6] Assuming an unpolarized sea the calculation for the expectation value of the contribution of quark (valence) spins (S_3) to the total proton spin gives

$$\begin{aligned}\int_0^1 \Delta u(x)dx&=S_3+\frac{1}{2}\frac{g_A}{g_V}\ ,\\ \int_0^1 \Delta d(x)dx&=S_3-\frac{1}{2}\frac{g_A}{g_V}<0\ .\end{aligned} \tag{13}$$

In this case one argues that the spin of the proton arises entirely from spin of the quarks then one expects $S_3=\frac{1}{2}$. However there exists a sum rule derived by Ellis and Jaffe[7] for deep-inelastic scattering of polarized electrons off polarized protons which constrains $S_3<\frac{1}{2}$. Further a QCD-based analysis[8] also yields this constraint. Data[9] suggest $S_3=0.3$, which we shall use.

Parton-model arguments yield the interesting limits

$$\lim_{x\to 0}\begin{bmatrix}A_u(x)\\ A_d(x)\end{bmatrix}=0$$

and (14)

$$\lim_{x\to 1}\begin{bmatrix}A_u(x)\\ A_d(x)\end{bmatrix}=1\ .$$

Comparing Eqs. (13) and (14), Callaway-Ellis parametrize

$$\begin{aligned}A_u(x)&=\frac{\Delta u(x)}{u(x)}=x^p\ ,\\ A_d(x)&=\frac{\Delta d(x)}{d(x)}=\left[\frac{x-x_0}{1-x_0}\right]x^{p'}\ .\end{aligned} \tag{15}$$

One treats x_0 (d quark spin flips at $x=x_0$) as an adjustable parameter which p and p' are determined from the sum rules [Eqs. (13)] above.

Both the models outlined above need spin-averaged valence-quark distributions as input to fix the parameters involved. In what follows the parametrization of Duke and Owens[10] is used for valence-quark distribution in a proton.

IV. THE PARAMETERS

(i) Carlitz-Kaur model. The sole parameter to be fixed in this case is H_0 [Eq. (11)]. As the predictions we wish

TABLE I. The Calitz-Kaur parameter H_0.

Q^2 (GeV^2)	H_0
6.3	0.069
10^2	0.050
10^3	0.0427
10^4	0.040

to arrive at involve variation over an enormous energy scale ($Q^2 \sim 10-10^4$ GeV^2), we find H_0 from the Bjorken sum rule including order-α_s corrections:

$$2\int_0^1 dx[G_1^p(x,Q^2)-G_1^n(x,Q^2)]$$
$$=\frac{1}{3}\frac{g_A}{g_V}\left[1-\frac{\alpha_s(Q^2)}{\pi}+\cdots\right]. \quad (16)$$

It has been argued, specifically in the context of the Carlitz-Kaur picture, that one has available the option of first QCD evolving spin-averaged distributions before extracting the explicit spin dependence.[11] Use is made of this assertion and values of $H_0=H_0(Q^2)$ are tabulated in Table I.

(ii) Callaway-Ellis model. In this case we have three parameters as noted earlier. Two of these, p and p', are fixed using the sum rules [Eqs. (11)] above with appropriate Q^2-evolved spin-averaged quark distributions at each Q^2 and are given in Table II. The fraction of proton momentum at which the quark spin flips x_0 is fixed in our analysis $x_0=0.75$ (see Fig. 1).

V. RESULTS

In Fig. 1 the predictions of the two models are compared with the available data. Keeping in view the low statistics and consequently wide error bars the agreement can at best be termed satisfactory. Clearly more data is needed.

If one examines the behavior of the double asymmetry $A(e^-p)$ in the limit of an infinitely heavy Z^0 ($M_Z \to \infty$), which corresponds to scattering mediated by photon exchange only, a very weak dependence on Q^2 results. This is apparent as scaling is broken only logarithmically. Interesting behavior results on restoring the coupling to the Z^0 boson. As shown in Fig. 2, the electroweak effects enhance the asymmetry as one goes from one momentum scale to a higher-momentum scale, at a fixed relative ener-

TABLE II. Parameters p and p for Callaway-Ellis model ($x_0=0.75$).

Q^2 (GeV^2)	p	p'
6.3	0.284	0.957
10^2	0.253	0.864
10^3	0.239	0.816
10^4	0.230	0.777

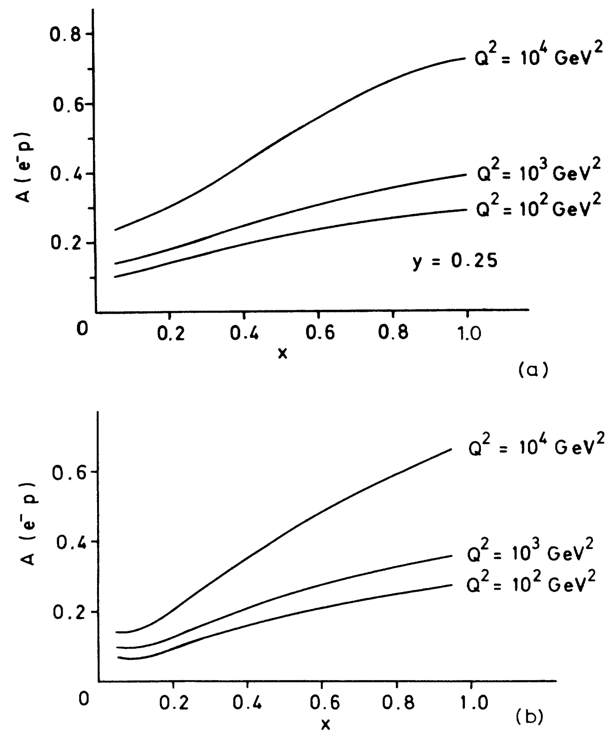

FIG. 2. $A(e^-p)$ at $y=0.25$, vs momentum fraction at three Q^2 values. (a) The Callaway-Ellis model, (b) the Carlitz-Kaur model.

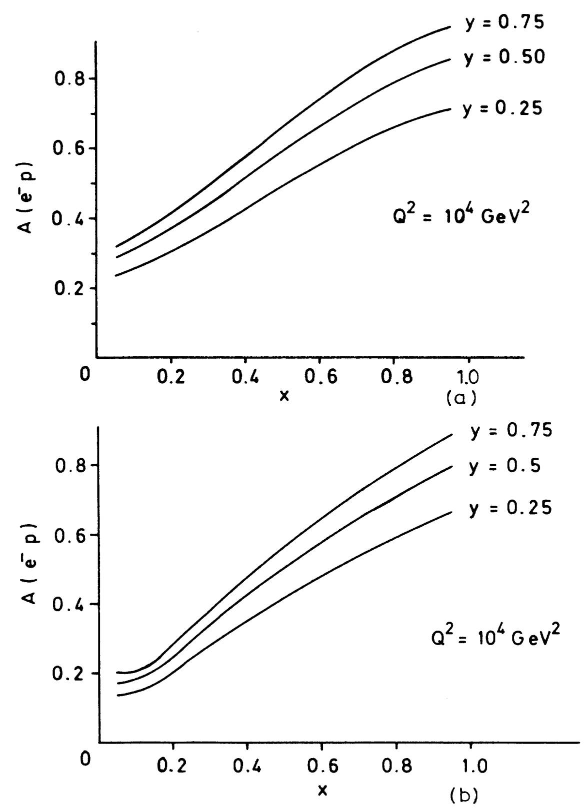

FIG. 3. $A(e^-p)$ at $Q^2=10^4$ GeV^2 at different y values. (a) The Callaway-Ellis and (b) Carlitz-Kaur.

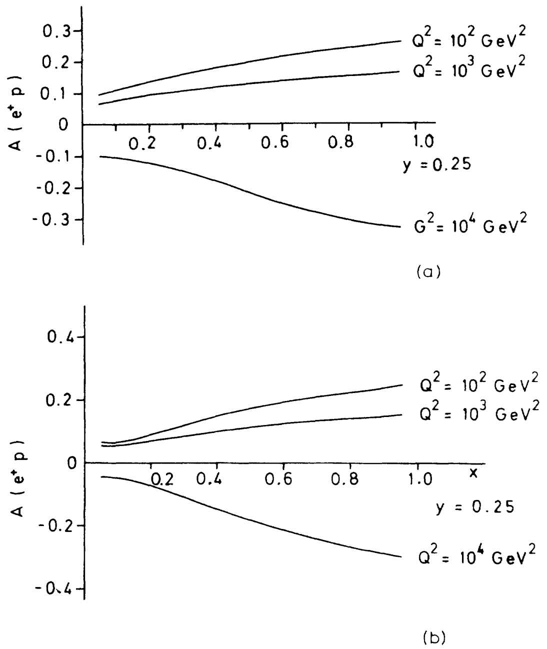

FIG. 4. $A(e^+p)$ at $y = 0.25$ at different Q^2. (a) Callaway and (b) Carlitz-Kaur models.

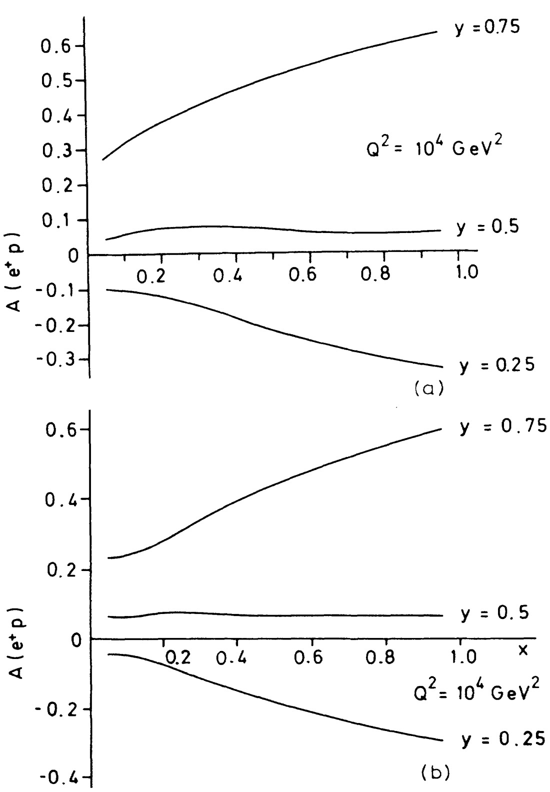

FIG. 5. $A(e^+p)$ at fixed Q^2 and varying y. (a) Callaway-Ellis and (b) Carlitz-Kaur models.

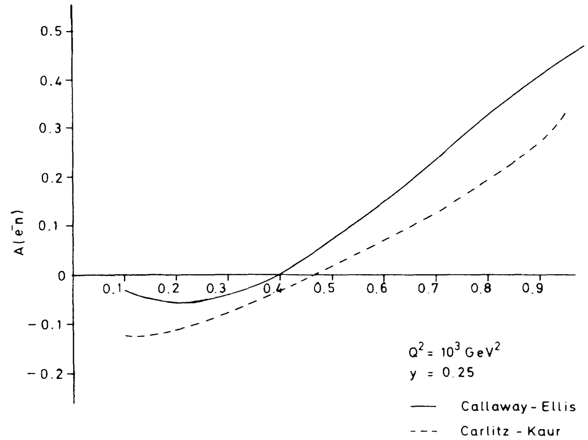

FIG. 6. $A(e^-n)$ for the two models compared.

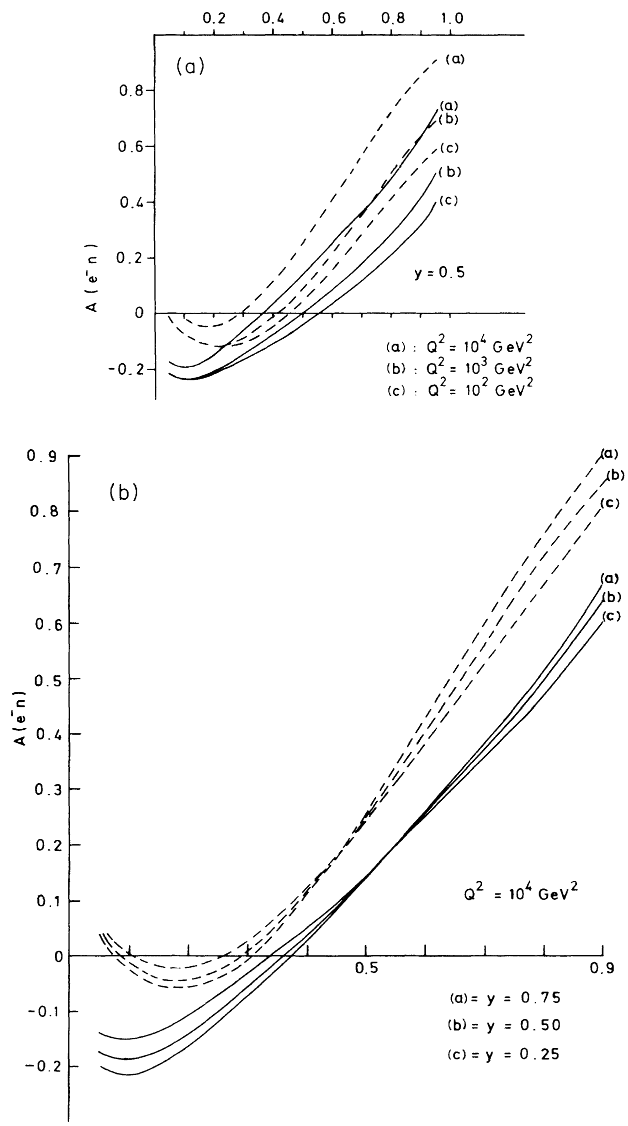

FIG. 7. $A(e^-n)$ for the two models compared; (a) variation with Q^2, (b) variation with y.

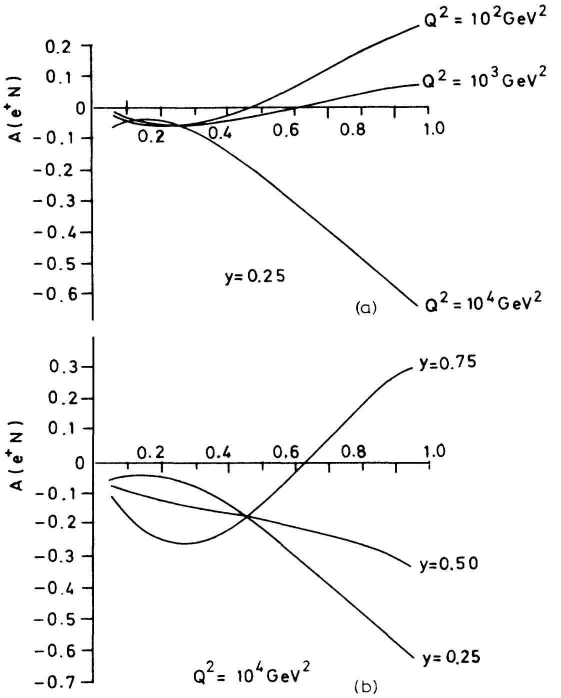

FIG. 8. The double asymmetry $A(e^+n)$ for the Callaway-Ellis model (a) varying Q^2 and (b) varying y.

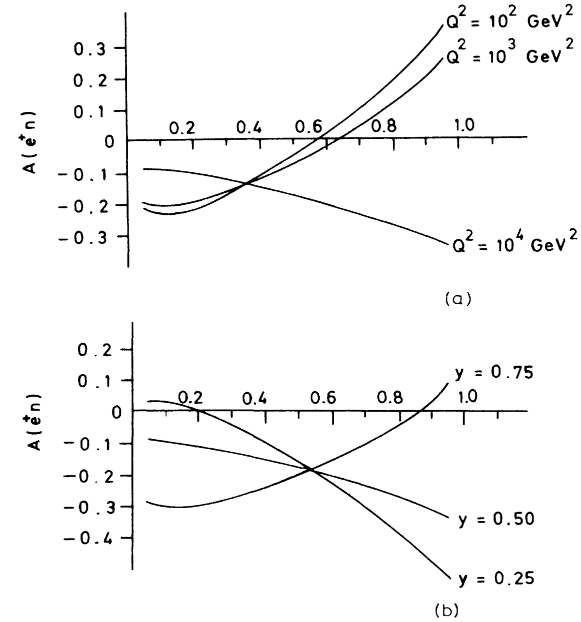

FIG. 9. The double asymmetry $A(e^+n)$ for the Carlitz-Kaur picture (a) varying Q^2 and (b) varying y.

gy loss (y) of the lepton. An enhancement is also observed on varying the lepton energy loss at a fixed Q^2 (Fig. 3).

For $y \neq 0$, additional information results if one considers $e^+ - p$ scattering (both target and beam polarized). The set of cross sections appropriate to $e^+ - p$ are easily obtained from Eqs. (1) by making the replacements

$$\sigma_{LR}(e^-p) \leftrightarrow \sigma_{RL}(e^+p) ,$$

$$1 \leftrightarrow (1-y)^2$$

on the left- and right-hand sides, respectively. Varying Q^2 at fixed y the results are displayed in Fig. 4. There appears a significant difference between the behavior of $A(e^-p)$ and $A(e^+p)$. In fact this reflects the different weights with which weak coupling enters at the lepton-quark vertex. This is also reflected in the variation of $A(e^+p)$, at fixed Q^2, with y (Fig. 5).

The replacement of a polarized proton target by a polarized neutron target is interesting. The double asymmetry $A(e^-n)$ is plotted in Fig. 6. For both the Carlitz-Kaur and Callaway-Ellis parametrizations the asymmetry changes sign. This indicates that the sign change in A [Eqs. (15)] is not responsible for the sign change in $A(e^-n)$. For completeness in Fig. 7 the variation of $A(e^-n)$ with Q^2 and with y are plotted. Finally the double asymmetry $A(e^+n)$ exhibits dramatic behavior as shown in Fig. 8 for the Callaway-Ellis model. The corresponding graph for the Carlitz-Kaur picture is given in Fig. 9.

CONCLUSIONS

We have, in this paper, tried to sample the wealth of information contained in the double asymmetry measurements at collider energies. Estimates of asymmetries one might expect to observe have been presented. Of course, fundamental importance must be attached to a direct QCD calculation of spin-dependent quark distributions inside a nucleon. In the absence of such a rigorous treatment one can only hope to gain further insight by additional testing of the phomenological models.

In the near future such experiments could be carried out at the DESY HERA facility. The major problem to be tackled is of having a high-energy beam of polarized protons. Complications arise due to the depolarizing resonances which occur frequently in the proton acceleration cycle. These were and are quite successfully tackled at the Argonne Zero Gradient Synchrotron and more recently at the Brookhaven Alternating Gradient Synchrotron. The techniques employed were resonance jumping, spin-flip, and orbit corrections. For high-energy machines such as HERA, the depolarizing resonances are probably too many to be handled individually; however, these can be completely avoided by the use of one or more "snakes" (combination of dipoles). These snakes bend the beam and rotate the spin in such a way that the depolarizing effects are eliminated. This method is well suited for high-energy accelerators as these have enough free space available for the magnets of the snakes.

In contrast to the available techniques for accelerating polarized protons the case for polarized neutrons is not as yet clear. Experiments with polarized targets however can be carried out at HERA and also at the Stanford

Linear Collider.

Notwithstanding the difficulties involved in polarized experiments it has to be recognized that in gauge theories the helicity degree of freedom is as fundamental as charge. The physics issues involved certainly warrant a close look at the possibility of doing spin physics at energies and short distances with present-day and upcoming accelerators.

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