Spin asymmetry in proton-proton collisions as a probe of sea and gluon polarization in a proton

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Quark and gluon spin densities in a proton are phenomenologically parametrized based on the European Muon Collaboration (EMC) data and on some plausible theoretical arguments. Four different characteristic values of gluon and sea polarizations suggested by various theoretical conjectures are considered. The sea polarization in a proton is probed by measuring the spin-spin asymmetry $A_{LL}^{\rm DY}$ in the Drell-Yan process, while the helicity asymmetry A_{LL}^{γ} in direct photon production at high p_T is employed to test the gluon spin content. Helicity asymmetries in both processes are quite sizable. $A_{LL}^{\rm DY}$ is *positive* and of order 10^{-1} if the sea is polarized opposite to the proton spin, as suggested by the EMC data. However, even in the absence of the sea polarization at the EMC energies, we find $A_{LL}^{\rm DY}$ to be large and *negative*. Experimental measurements of $A_{LL}^{\rm DY}$ and A_{LL}^{γ} together will not only provide a clean probe of sea and gluon polarizations, but also test whether the combination $\Delta s-(\alpha_s/4\pi)\Delta G$ inferred from the EMC data is valid, i.e., whether gluons contribute to the spin-dependent structure function $g_1^p(x,Q^2)$ via the triangular anomaly.

I. INTRODUCTION

The recent measurements of the proton polarized structure function by the European Muon Collaboration[1] (EMC), when combined with the data of neutron and hyperon β decays, imply the existence of a substantial negative polarization of sea quarks and/or a sizable net spin carried by the gluons in a polarized proton.[2] In order to know the spin content of the proton it is very important to measure the relative weight of spin densities of sea quarks and gluons. We propose in the present paper various schemes for the parton spin densities. We then proceed to study the longitudinal spin asymmetry in the Drell-Yan process and in direct photon production at large transverse momentum in pp collisions with polarized beams and targets to test the polarization of gluons and sea quarks.

In deep-inelastic polarized lepton scattering off polarized protons, the first moment of the spin-dependent structure function is related to the proton matrix elements of the axial-vector currents via the operator-product expansion (OPE):

$$\begin{aligned}M_1^p(Q^2)&\equiv\int_0^1 dx\, g_1^p(x,Q^2)\\&=\frac{1}{12}\left[1-\frac{\alpha_s(Q^2)}{\pi}\right]\\&\quad\times(\langle p^\uparrow|A_\mu^3|p^\uparrow\rangle\\&\quad+\tfrac{1}{3}\langle p^\uparrow|A_\mu^8|p^\uparrow\rangle+\tfrac{4}{3}\langle p^\uparrow|A_\mu^0|p^\uparrow\rangle)_{\mu=3}\end{aligned}\tag{1.1}$$

with

$$\begin{aligned}A_\mu^3&=\bar u\gamma_\mu\gamma_5u-\bar d\gamma_\mu\gamma_5 d\ ,\\A_\mu^8&=\bar u\gamma_\mu\gamma_5u+\bar d\gamma_\mu\gamma_5 d-2\bar s\gamma_\mu\gamma_5 s\ ,\\A_\mu^0&=\bar u\gamma_\mu\gamma_5u+\bar d\gamma_\mu\gamma_5 d+\bar s\gamma_\mu\gamma_5 s\ ,\end{aligned}\tag{1.2}$$

where the proton is polarized in the z direction. It should be stressed that the matrix elements $\langle p|A_\mu^a|p\rangle$ in Eq. (1.1) are evaluated at $\mu^2=Q^2$ (Ref. 3) and have the following parton-model interpretation:[4]

$$\Delta q_i+\Delta\bar q_i-\frac{\alpha_s}{2\pi}\Delta G=\langle p^\uparrow|\bar q_i\gamma_3\gamma_5 q_i|p^\uparrow\rangle\ ,\tag{1.3}$$

where we have dropped $\bar\psi\gamma_3\gamma_5\psi$ for convenience,

$$\Delta q=\int_0^1 dx[q^\uparrow(x)-q^\downarrow(x)]\equiv\int_0^1 dx\,\Delta q(x)$$

is the helicity of the quark q in the infinite-momentum frame of the proton with positive helicity,[5] and

$$\Delta G=\int_0^1 dx[G^\uparrow(x)-G^\downarrow(x)]\equiv\int_0^1 dx\,\Delta G(x)$$

is the net spin carried by the gluons. A nice feature of the OPE is that if the anomalous dimensions of A_μ^a vanish, then the first moment of g_1^p at high Q^2 attributed from A_μ^a are related to the low-energy axial charges

$$g_A^a=\left\langle p\left|\int d^3x\ A_0^a(\mathbf{x},t)\right|p\right\rangle\ .\tag{1.4}$$

There are two arguments indicating that the matrix element $\langle p^\uparrow|\bar q\gamma_3\gamma_5 q|p^\uparrow\rangle$ and hence M_1^p should include anomalous gluon contributions. In the OPE approach, we notice that A_μ^0 has an anomalous dimension which first occurs at two-loop order.[6] On the one hand, since Δq is not affected by gluon emissions[7] $\langle p|A_\mu^0|p\rangle$ should not be just identified with the quark's helicity since the former evolves with Q^2. On the other hand, although there are no twist-two gluonic operators contributing to M_1^p, the matrix element $\langle G|A_\mu^0|G\rangle$ does not vanish at order α_s due to the triangular anomaly. To be more specific, we decompose the axial-vector current in the form (n_f being the number of quark flavors)

$$A_\mu^0=\tilde{A}_\mu^0+n_f K_\mu\ ,$$
$$K_\mu=\frac{\alpha_s}{2\pi}\epsilon_{\mu\nu\alpha\beta}A_a^\nu\left[\partial^\alpha A_a^\beta-\frac{g}{3}f_{abc}A_b^\alpha A_c^\beta\right]\ , \tag{1.5}$$

so that the gauge-noninvariant current $\tilde{A}_\mu^0$ is conserved in the chiral limit. The proton matrix elements are thus identified with[8]

$$\langle p^\uparrow|\tilde{A}_3^0|p^\uparrow\rangle=\Delta\Sigma\equiv\sum^{n_f}(\Delta q_i+\Delta\bar{q}_i)\ ,$$
$$\langle p^\uparrow|K_3|p^\uparrow\rangle=-\Delta\Gamma\equiv-\frac{\dot{\alpha}_s}{2\pi}\Delta G\ . \tag{1.6}$$

As pointed out by Altarelli and Ross,[8] though the current K_μ is gauge variant, its diagonal matrix elements are gauge invariant as the gauge-dependent part of K_μ can be recast as a four-divergence. It follows from Eq. (1.1) that

$$M_1^p(Q^2)=\frac{1}{2}\sum e_i^2\left[1-\frac{\alpha_s(Q^2)}{\pi}\right]\times\left[\Delta q_i+\Delta\bar{q}_i-\frac{\alpha_s(Q^2)}{2\pi}\Delta G(Q^2)\right]\ . \tag{1.7}$$

The next crucial observation, as first pointed out by Ratcliffe[9] and stressed by Altarelli and Ross,[8] is that $\Delta G(Q^2)$ increases logarithmically with Q^2 and hence $\alpha_s\Delta G$ is conserved to the leading QCD evolution; that is, $\alpha_s\Delta G$ is really of order $(\alpha_s)^0$ rather than α_s.

In the parton model, the spin structure function of the proton at order α_s reads[10]

$$g_1^p(x,Q^2)=\frac{1}{2}\int_x^1\frac{dy}{y}\sum_i^{2n_f}e_i^2\left\{\Delta q_i(y,Q^2)\left[\delta\left[1-\frac{x}{y}\right]+\frac{\alpha_s(Q^2)}{2\pi}\Delta f_q\left[\frac{x}{y}\right]\right]+\Delta G(y,Q^2)\frac{\alpha_s(Q^2)}{2\pi}\Delta f_G\left[\frac{x}{y}\right]\right\}\ . \tag{1.8}$$

The quantity Δf_G has been computed by Ratcliffe[11] and subsequently by many authors.[8,12] It depends on the regularization adopted for the infrared divergence and on the renormalization scheme. For example, using dimension regularization Ratcliffe obtained

$$\Delta f_G(x)=\frac{1}{2}\left[(2x-1)\ln\frac{1-x}{x}-a(2x-1)\right]-\frac{1}{\epsilon}\Delta P_{qG}(x)\ , \tag{1.9}$$

where the polarized kernel $\Delta P_{qG}(x)$ is given by $\frac{1}{2}(2x-1)$, and the regularization-dependent term $a(2x-1)$ stems from the definition of $\epsilon^{\mu\nu\alpha\beta}$ in $2-\epsilon$ dimensions. Nevertheless, the first moment of $\Delta f_G(x)$ is well defined and is independent of the renormalization scheme chosen:

$$\int_0^1 dx\ \Delta f_G(x)=-\tfrac{1}{2}\ . \tag{1.10}$$

This together with the result $\int_0^1 dx\ \Delta f_q(x)=-2$ (Ref. 11) leads precisely to Eq. (1.7) (Ref. 13). Note that for the unpolarized structure function $F_1(x,Q^2)$, the first moments of $f_G(x)$ and $f_q(x)$ vanish.[11] As a consequence, there is no gluon contribution to the first moment of $F_1(x,Q^2)$, contrary to the spin-dependent structure function $g_1^p(x,Q^2)$. Since only the first moment of the polarized quark densities has been precisely defined, we will follow Ref. 14 to fix the high moments so that the gluon contribution takes the simplest form

$$g_1^p(x,Q^2)=\frac{1}{2}\sum_i^{n_f}\left[1-\frac{\alpha_s(Q^2)}{\pi}\right]e_i^2\times\left[\Delta q_i(x,Q^2)+\Delta\bar{q}_i(x,Q^2)-\frac{\alpha_s(Q^2)}{2\pi}\Delta G(x,Q^2)\right]\ . \tag{1.11}$$

Now the polarized quark and gluon densities are related to the axial charges g_A^a via Eqs. (1.2) and (1.3) by

$$g_A^3=\Delta u+\Delta\bar{u}-\Delta d-\Delta\bar{d}\ ,$$
$$g_A^8=\Delta u+\Delta\bar{u}+\Delta d+\Delta\bar{d}-2\Delta s-2\Delta\bar{s}\ , \tag{1.12}$$
$$g_A^0=\sum_i^{n_f}(\Delta q_i+\Delta\bar{q}_i)-\frac{3\alpha_s}{2\pi}\Delta G\ .$$

The SU(3)-nonsinglet couplings g_A^3 and g_A^8 are Q^2 independent and thus can be related to the usual F and D parameters determined from low-energy neutron and hyperon beta decays by means of SU(3)-flavor symmetry

$$g_A^3=F+D,\quad g_A^8=3F-D\ . \tag{1.13}$$

The isoscalar coupling g_A^0, however, cannot be determined by SU(3) symmetry alone. Nevertheless, it can be extracted from the recent EMC data. The first moment of the spin structure function g_1^p was measured by EMC (Ref. 11) to be $\int_0^1 dx\ g_1^p(x,Q_{\rm EMC}^2)=0.114\pm0.012\pm0.026$ at $Q_{\rm EMC}^2=10.7\ \mathrm{GeV}^2$. The total error can be reduced by combining the older SLAC data[15] over the region $0.1\le x\le 0.7$ with the EMC points at low x $(0.01<x<0.1)$; this yields[16]

$$\int_0^1 dx\ g_1^p(x,Q_{\rm EMC}^2)=0.116\pm0.009\pm0.019\ . \tag{1.14}$$

When combined with Eqs. (1.12) and (1.13), the EMC measurement (1.14) leads to

$$Du=0.75\pm0.07,\quad Dd=-0.51\pm0.07\ ,$$
$$Ds=-0.22\pm0.007\ , \tag{1.15}$$
$$g_A^0=Du+Dd+Ds=0.02\pm0.21\ ,$$

at $Q_{\rm EMC}^2=10.7\ \mathrm{GeV}^2$, where $Dq=\Delta q+\Delta\bar{q}\ -(\alpha_s/2\pi)\Delta G$, and uses of $F+D=1.259\pm0.004$ (Ref. 17) and $F/D=0.61-0.64$ have been made.

The EMC result (1.15) looks bizarre at first glance since it indicates that the proton spin cannot arise solely from the valence quarks even if sea quarks and gluons are unpolarized. Moreover, the helicity of the quark-antiquark sea and gluons are intimately correlated. However, it should be stressed that the SU(6) nonrelativistic quark model (NQM) does not really predict what is the net helicity carried by the current (QCD) quarks in the proton. First, the naive SU(6) NQM prediction $\Delta U+\Delta D=1$ refers to the constituent quarks Q_i rather than the current quarks q_i. Second, NQM applies only to the low-energy nucleon. Nevertheless, if the Ellis-Jaffe ansatz[18] $\Delta s=0$ is made at the EMC energies, then $\Delta u+\Delta d\sim 0.70$ and $M_1^p\approx 0.18$ are inferred from Eqs. (1.12) and (1.13). The result $\Delta u+\Delta d\sim 0.70$ is usually regarded as the "standard prediction" of NQM for the polarization of the valence quarks in a proton. In fact, if the sea polarization is SU(3) invariant, i.e., $\Delta u_s=\Delta\bar{u}_s=\Delta d_s=\Delta\bar{d}_s=\Delta s=\Delta\bar{s}$ (the subscripts v and s denote "valence" and "sea," respectively), it is easily seen from Eq. (1.15) that, irrespective of the value of ΔG,

$$\Delta u_v+\Delta d_v\approx 0.68 ,$$

in accord with the "NQM's" prediction $\Delta u_v+\Delta d_v\approx 0.70$.

What is the magnitude and sign of the gluon spin component in a polarized proton is a subject of hot controversy at present. The EMC leptoproduction data are not adequate to determine the helicity carried by the gluons, but they do provide a useful correlation between Δs and ΔG: namely, $\Delta s+\Delta\bar{s}-\Delta\Gamma=-0.22\pm 0.07$. Three major possibilities for the gluon contribution have been considered in the literature.

(i) The polarized proton does not contain polarized strange quarks at the EMC energies. As a result, $\Delta\Gamma\approx 0.22$ or $\Delta G\approx 5.5$ for $\alpha_s=0.25$. Since ΔG grows logarithmically with Q^2, a large $\Delta G\sim 5$–6 is conceivable. However, the spin sum rule for the proton

$$\tfrac{1}{2}=\tfrac{1}{2}\Delta\Sigma+\Delta G+(L_z)_q+(L_z)_G \tag{1.16}$$

indicates that a large negative orbital angular momentum of order -5.3 due to quarks and/or gluons is required to balance ΔG. In the absence of the sea polarization, the discrepancy between theory based on the Ellis-Jaffe ansatz and experiment for M_1^p arises entirely from the anomalous gluon spin contributions.[8]

(ii) It has been shown that ΔG is very small in some particular versions of the Skyrme model.[19] Consequently, $\Delta s+\Delta\bar{s}\approx -0.22\pm 0.07$ and most of the proton spin comes from the orbital angular momentum, inferred from Eqs. (1.15) and (1.16). We caution that the result $\Delta G\approx 0$ is derived at $Q^2=0$ and is not necessarily valid at the energies $Q^2=10.7$ GeV^2.

(iii) From the anomalous divergence of the axial-vector currents and from the smallness of the pseudoscalar matrix element $\langle p|\bar{u}i\gamma_5 u+\bar{d}i\gamma_5 d+\bar{s}i\gamma_5 s|p\rangle$ owing to the decoupling of a singlet Goldstone boson from the nucleon, we find[20] $\langle p|(\alpha_s/8\pi)G^a_{\mu\nu}\tilde{G}^{a\mu\nu}|p\rangle\approx 268$ MeV. Since $\partial^\mu K_\mu=(\alpha_s/4\pi)G^a_{\mu\nu}\tilde{G}^{a\mu\nu}$, it follows from Eq. (1.6) that[21]

$$\frac{\alpha_s}{2\pi}\Delta G=-\frac{1}{m_N}\left\langle p\uparrow\left|\frac{\alpha_s}{8\pi}G^a_{\mu\nu}\tilde{G}^{a\mu\nu}\right|p\uparrow\right\rangle\approx -0.29 . \tag{1.17}$$

This striking result derived at zero momentum transfer implies a large but *negative* gluon spin contribution at moderate $Q^2=10.7$ GeV^2. The sea-quark polarization $2\Delta s$ is of order -0.5 and hence deviates even further from the naive-quark-model expectations.[21]

Several processes in which the gluon polarization can be extracted directly from experiment have been suggested in the literature. These include heavy-quark pair production in the polarized photon-gluon fusion process (e.g., the photoproduction[22] $\gamma p\to J/\psi+X$, and the leptoproduction[23] $\mu p\to\mu+J/\psi+X$), high-p_T jet production in polarized deep-inelastic lepton scattering,[24] gluon jet oblateness,[25] and direct photon production at large p_T in polarized proton-proton collisions.[26,27,12] In the present paper we shall study the spin-spin asymmetry in the Drell-Yan process and in prompt photon production at large transverse momentum with polarized proton beams and proton targets. It has been known that the former is sensitive to the sea polarization,[27] whereas the latter depends strongly on the gluon spin densities.[26,27,12]

This paper is organized as follows. We shall present in Sec. II several phenomenological parametrizations of parton spin-weighted distributions based on the EMC data and on some plausible theoretical arguments. We then proceed to compute the polarization asymmetry in the Drell-Yan process and in direct photon production at high p_T in Secs. III and IV, respectively, for four different characteristic values of ΔG and Δs. Conclusions are given in Sec. V.

II. POLARIZED DISTRIBUTIONS

In this section we shall elaborate on the spin densities of quarks and gluons in a polarized proton. To begin with, it is convenient to write Δq in terms of valence and sea contributions

$$\Delta q=\Delta q_v+\Delta q_s . \tag{2.1}$$

Let

$$\Delta u_s=\Delta\bar{u}_s=\Delta d_s=\Delta\bar{d}_s=(1+\epsilon)\Delta s=(1+\epsilon)\Delta\bar{s} , \tag{2.2}$$

where the parameter ϵ characterizes the degree of SU(3) breaking. It is easily seen from Eq. (1.15) that when $\epsilon=0$, $\Delta u_v+\Delta d_v=0.68$ irrespective of the value of ΔG. Obviously, the net spin carried by the valence quarks is very close to what is expected from the NQM in conjunction with the Ellis-Jaffe ansatz. Therefore, we will assume that the sea polarization is SU(3) invariant (unpolarized sea distributions for s and u or d quarks are known to be different, however) and Δq_v are those predicted by the NQM: namely,

$$\begin{aligned}\Delta u_v &\equiv \int_0^1 dx\,\Delta u_v(x)=0.98 ,\\ \Delta d_v &\equiv \int_0^1 dx\,\Delta d_v(x)=-0.28 .\end{aligned} \tag{2.3}$$

To proceed further we follow Ref. 28 assuming that the polarized and unpolarized valence-quark distributions are related via

$$\Delta u_v(x)=\alpha(x)u_v(x), \quad \Delta d_v(x)=\beta(x)d_v(x) \ . \tag{2.4}$$

As $x\to 1$ it is desirable to have $\alpha(x)$, $\beta(x)\to 1$ to reflect the argument[29] that the valence quark at $x=1$ remembers the spin of the parent proton as supported by the perturbative QCD suggestion.[30] By contrast, the region near $x=0$ is expected to be dominated by the sea so that the spin of the parent proton would no longer be reflected in the valence quarks. This suggests that $\alpha(x)$, $\beta(x)\to 0$ as $x\to 0$. Since Δd_v is negative, the boundary condition $\beta(x)\to 1$ as $x\to 1$ implies that $\Delta d_v(x)$ changes sign as a function of x. Hence, $\beta(x)$ can be taken as[31]

$$\beta(x)=\left[\frac{x-x_0}{1-x_0}\right]x^p \tag{2.5}$$

so that the sign of $\Delta d_v(x)$ is flipped at $x=x_0$ when p is positive.

In the following analysis we will use the "average" set of the unpolarized parton distribution functions given by Diemoz, Ferroni, Longo, and Martinelli[32] (DFLM): namely,

$$\begin{aligned} xu_v(x)&=2.26x^{0.54}(1-x)^{2.52}\\ &\quad\times[1-1.617(1-x)+3.647(1-x)^2\\ &\quad-1.998(1-x)^3] \ ,\\ xd_v(x)&=0.57(1-x)xu_v(x) \ ,\\ x\bar{s}(x)&=0.4x\bar{u}(x)=0.4x\bar{d}(x)\\ &=0.1(1-x)^{8.5}\\ &\quad\times(1-4.18x+20.3x^2-15.3x^3) \ ,\\ xG(x)&=3.34(1-x)^{5.06}(1-0.177x) \end{aligned} \tag{2.6}$$

extracted at the reference scale $Q^2=10$ GeV^2 from several experiments. Note that the gluon distribution functions of DFLM, supported by the large-p_T direct photon experiments,[33] are substantially larger than those given by Duke and Owens[34] and by Eichten, Hincliffe, Lane, and Quigg[35] at the small-x region, $x<0.1$. We find that the simple parametrization

$$\Delta u_v(x)=x^{0.26}u_v(x) \tag{2.7}$$

satisfies all constraints mentioned above. For $\Delta d_v(x)$, x_0 must be greater than 0.31 in order for p to be positive. Treating x_0 as a free parameter, we get

$$p=\begin{cases} 0.08 & \text{at } x_0=0.35,\\ 0.31 & \text{at } x_0=0.50,\\ 0.84 & \text{at } x_0=0.75. \end{cases} \tag{2.8}$$

According to the Carlitz-Kaur model,[36] interactions between the valence quark and the sea or gluons cause a spin dilution of the valence quark.[37] Therefore, a spin-dilution factor

$$\cos 2\theta=\frac{1}{1+H(x)N(x)} \tag{2.9}$$

is introduced, where $N(x)$ denotes the density of gluons relative to the valence quarks (sea-quark effects can be neglected as gluons carry $\sim 50\%$ of the proton momentum) and $H(x)$ is the probability of interaction between the valence quark and gluons. The spin-dilution factor is taken to be

$$\cos 2\theta=\frac{1}{1+H_0\dfrac{(1-x)^2}{\sqrt{x}}} \tag{2.10}$$

by Carlitz and Kaur, where the free parameter H_0 is fixed by Eq. (2.3). A fit to $\Delta u_v=0.98$ yields $H_0=0.31$. It is easy to check that this spin-dilution factor is actually quite similar to our $\alpha(x)=x^{0.26}$. However, the d-quark spin density

$$\Delta d_v(x)=-\tfrac{1}{3}\cos 2\theta\, d_v(x) \tag{2.11}$$

proposed in the SU(6) Carlitz-Kaur model does not satisfy the perturbative QCD result $\beta(x)\to 1$ as $x\to 1$. Hence, we will confine ourselves to the polarized d-quark distributions given by Eqs. (2.5) and (2.8).

We next turn to the spin densities of the sea and gluons. it is clear from Eq. (1.15) that the net spin carried by the sea quarks and gluons at energies $Q^2_{\rm EMC}=10.7$ GeV^2 is correlated by the relation

$$\Delta s-\frac{\alpha_s}{4\pi}\Delta G=-0.11\pm 0.03 \ . \tag{2.12}$$

For the purpose of illustration, four different characteristic values of ΔG and hence Δs are chosen in the present paper: (a) $\Delta G=5.5$, $\Delta s=0$, (b) $\Delta G=3.0$, $\Delta s=-0.05$, (c) $\Delta G=0$, $\Delta s=-0.11$, and (d) $\Delta G=-4.5$, $\Delta s=-0.20$. Cases (a), (c), and (d) were already discussed in Sec. I. For completeness we also consider case (b) which is between cases (a) and (c). As elaborated on in Sec. I, case (d) that both gluons and sea quarks are polarized in opposite direction to the proton spin deviates mostly from the expectations of the naive nonrelativistic quark model.

In order to ensure the positivity constraints $|\Delta s(x)/s(x)|\le 1$ and $|\Delta G(x)/G(x)|\le 1$ be satisfied for all x between 0 and 1, we write the polarized sea-quark and gluon distributions as

$$\Delta s(x)=-x^{\gamma_s}s(x), \quad \Delta G(x)=x^{\gamma_G}G(x) \tag{2.13}$$

with γ_s, $\gamma_G>0$ [$\Delta G(x)=-x^{\gamma_G}G(x)$ for case (d)]. The parametrization of gluon spin densities given by Eq. (2.13) is justified on the grounds that according to the perturbative QCD calculation $\Delta G_\downarrow(x)$ is suppressed relative to $\Delta G^\uparrow(x)$ by a factor of $(1-x)^2$ at large x (Refs. 30 and 38), while $\Delta G(x)$ at small x does not increase as rapidly as $G(x)$. For the polarized sea densities, the more appropriate parametrization is $\Delta s(x)/s(x)=-Ax^m(1-x)^n$ with $A\sim 1$ and $m,n>0$ so that the nonvalence polarization $\Delta s(x)$ vanishes when $x\to 1$. However, we find that it makes no practical differences as to which parametrization for the sea spin densities is employed. Fitting (2.13) to the above four different cases of proton helicity

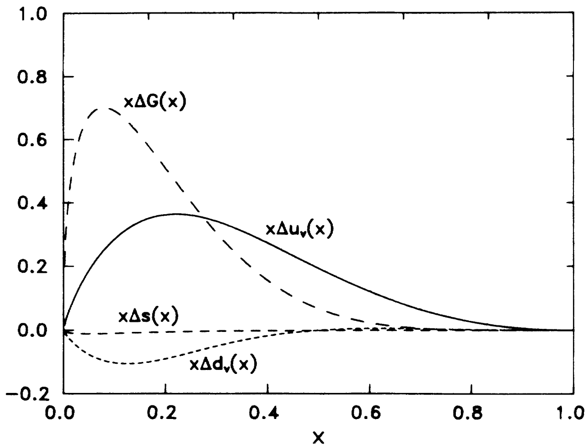

FIG. 1. Parton spin densities with $\Delta G(Q^2_{\rm EMC})=3.0$ and $\Delta s(Q^2_{\rm EMC})=-0.05$. For $\Delta d_v(x)$, the parameter $x_0=0.50$ is used [see Eq. (2.8)].

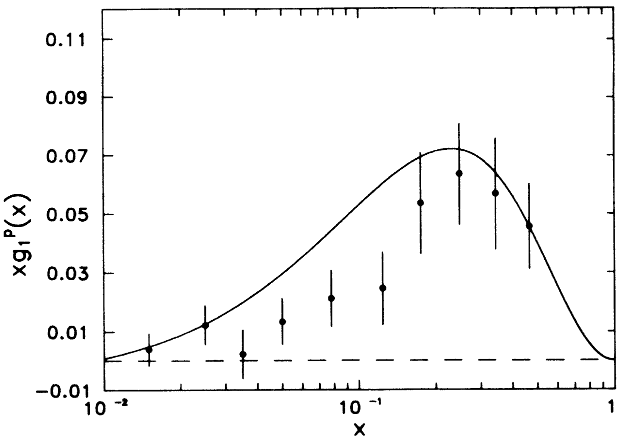

FIG. 3. Predicted spin-dependent structure function $xg_1^p(x,Q^2)$ of the proton for case (d) described in the text.

contents ΔG and Δs gives

$$\begin{aligned}&(\mathrm{a})\ \gamma_G=0.31,\quad (\mathrm{b})\ \gamma_s=0.51,\ \gamma_G=0.45\ ,\\ &(\mathrm{c})\ \gamma_s=0.35,\quad (\mathrm{d})\ \gamma_s=0.25,\ \gamma_G=0.35\ .\end{aligned} \tag{2.14}$$

The spin-dependent quark and gluon distributions for case (b) are plotted in Fig. 1. The proton polarized structure function $xg_1^p(x)$ vs x is displayed in Figs. 2 and 3, respectively, for cases (b) and (d). [The spin structure function $xg_1^p(x)$ for (a) and (c) is very close to that of case (b) and hence is not shown.] It is evident that the EMC data of $xg_1^p(x)$ are well fitted for $\Delta G \gtrsim 0$ (Ref. 39). This implies that if gluons inside the proton are polarized, they prefer to polarize in the same direction as the proton spin.

Once the parton helicity content of the proton is known, it is expected that the neutron polarized structure function $xg_1^n(x)$ can be immediately determined by the relation

$$g_1^n(x)=\tfrac{1}{2}\sum_i^{n_f}\left[1-\frac{\alpha_s}{\pi}\right]\left[\tfrac{4}{9}\Delta d_v(x)+\tfrac{1}{9}\Delta u_v(x)+\tfrac{12}{9}\Delta s(x)-\frac{\alpha_s}{3\pi}\Delta G(x)\right] \tag{2.15}$$

as depicted in Fig. 4. However, it should be stressed that Eq. (2.15) is obtained from Eq. (1.11) via isospin symmetry argument. It has been known for some time that nucleon matrix elements of the anomaly term $\alpha_s \mathrm{Tr} G\tilde{G}$ are far from isoscalar.[40] Consequently, ΔG measured in a polarized neutron is substantially different from the gluon component seen in the proton.[41] The large isospin violation of ΔG, however, does not affect the validity of the Bjorken sum rule for spin-dependent electroproduction: the isospin-violating effects due to the quark mass difference of u and d quarks conspire with the isospin symmetry breaking of the anomaly operator to render all matrix elements of the axial-vector currents free of large isospin violation.[40] Of course, the x dependence (but *not*

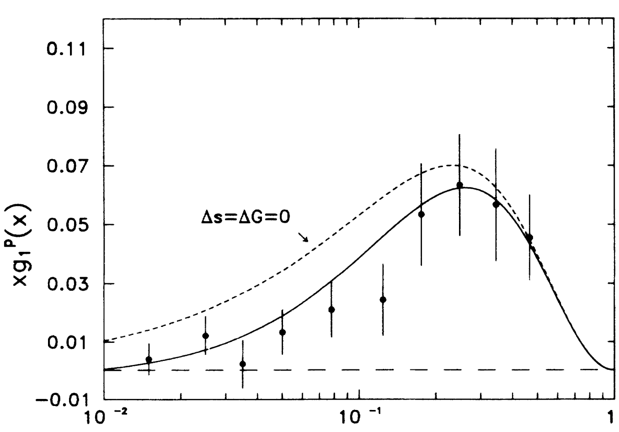

FIG. 2. Predicted spin-dependent structure function $xg_1^p(x,Q^2)$ of the proton for case (b) described in the text. Curves for cases (a) and (c) are quite similar to that of (b) and hence are not shown here. The dashed line is the contribution from valence quarks only.

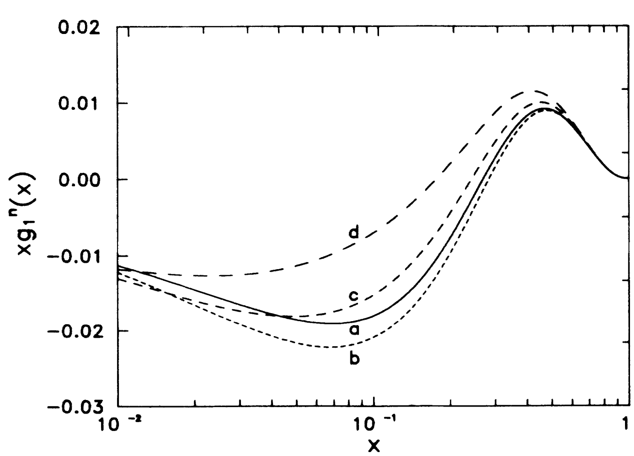

FIG. 4. Predicted spin-dependent structure function $xg_1^n(x,Q^2)$ of the neutron under the assumption of isospin symmetry for four different cases of ΔG and Δs described in the text.

the first moment) of $g_1^n(x)$ may still exhibit considerable amount of isospin nonconservation. Hence, any substantial deviation from the isospin predictions presented in the Fig. 4 will signal the breakdown of isospin symmetry for ΔG and Δq.

Thus far all parton spin distributions are determined at $\langle Q_0^2 \rangle = 10.7$ GeV^2. Their Q^2 evolutions are governed by the Altarelli-Parisi (AP) equations

$$\frac{d\Delta q_v(x,t)}{dt} = \int_x^1 \frac{dz}{z} \Delta P_{qq}(z) \Delta q_v\left[\frac{x}{z}, t\right],$$

$$\frac{d\Delta\Sigma(x,t)}{dt} = \int_x^1 \frac{dz}{z}\left[\Delta P_{qq}(z)\Delta\Sigma\left[\frac{x}{z},t\right] + 2n_f \Delta P_{qG}(z)\Delta G\left[\frac{x}{z},t\right]\right], \quad (2.16)$$

$$\frac{d\Delta G(x,t)}{dt} = \int_x^1 \frac{dz}{z}\left[\Delta P_{Gq}(z)\Delta\Sigma\left[\frac{x}{z},t\right] + \Delta P_{GG}(z)\Delta G\left[\frac{x}{z},t\right]\right],$$

where

$$t = \frac{1}{b}\ln\left[\frac{\ln(Q^2/\Lambda^2)}{\ln(Q_0^2/\Lambda^2)}\right], \quad b = \frac{33-2n_f}{6}. \quad (2.17)$$

To the first order in α_s, the polarized AP kernels read[10]

$$\begin{aligned}
\Delta P_{qq}(x) &= \frac{4}{3}\left[\frac{1+x^2}{1-x}\right]_+, \\
\Delta P_{qG}(x) &= \tfrac{1}{2}(2x-1), \\
\Delta P_{Gq}(x) &= \tfrac{4}{3}(2-x), \\
\Delta P_{GG}(x) &= 3\left[(1+x^4)\left[\frac{1}{x} + \frac{1}{(1-x)_+}\right] - \frac{(1-x)^3}{x}\right] + b\delta(1-x).
\end{aligned} \quad (2.18)$$

The next-order polarized kernels have not been completely worked out, but their first moments are known.[42] Since

$$\int_0^1 dx\, \Delta P_{qq}(x) = \int_0^1 dx\, \Delta P_{qG}(x) = 0, \quad \int_0^1 dx\, \Delta P_{GG}(x) = b \quad (2.19)$$

it follows from Eq. (2.16) that

$$\begin{aligned}
\Delta q_i(Q^2) &= \Delta q_i(Q_0^2), \\
\Delta G(Q^2) &= \Delta G(Q_0^2)\frac{\alpha_s(Q_0^2)}{\alpha_s(Q^2)} + \frac{9}{2b}\Delta\Sigma(Q_0^2)\left[\frac{\alpha_s(Q_0^2)}{\alpha_s(Q^2)} - 1\right].
\end{aligned} \quad (2.20)$$

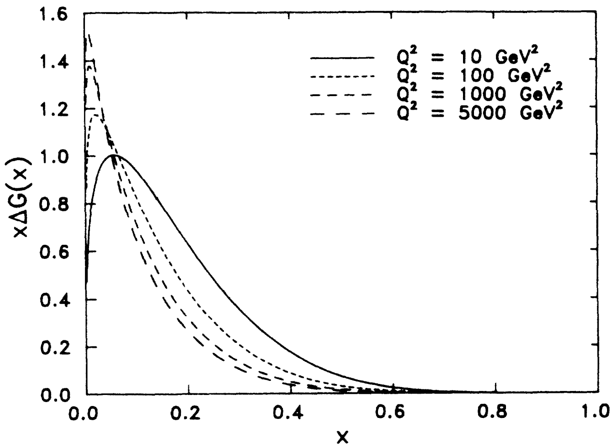

FIG. 5. The evolution of gluon spin densities at various values of Q^2 with $\Delta G(Q^2_{\rm EMC}) = 5.5$.

We see that the net helicity of each quark flavor is conserved as it should be and $\alpha_s(Q^2)\Delta G(Q^2)$ is independent of Q^2 to the leading order in α_s.

Figures 5–8 exhibit the Q^2 evolution of the polarized gluon and sea-quark densities with the initial parton spin content given by cases (a) and (c) (Ref. 43). From Fig. 6 we see that even if the strange sea quarks are initially unpolarized at $Q^2_{\rm EMC} = 10.7$ GeV^2, a *positive* $\Delta s(x,Q^2)$ is nevertheless dynamically generated when $Q^2 > Q^2_{\rm EMC}$ and $x > 0.01$. Moreover, the degree of gluon-induced sea polarization is comparable to that evolved from initially nonzero Δs (see Figs. 6 and 8). Of course, the first moment of $\Delta s(x,Q^2)$ in case (a) must vanish as we did check explicitly. Likewise, shown in Fig. 7 is the x behavior of the nonvanishing spin-dependent gluon densities generated from the bremsstrahlung of the quarks with $\Delta G(x, Q^2_{\rm EMC}) = 0$. Notice that the quark-induced ΔG is small in magnitude (see Figs. 5 and 7). We should also mention that sea quarks polarize positively at large x even if $\Delta s(x)$ is initially negative (though we cannot tell this from Fig. 8 since sea densities are very small at large x). This feature will explain the sign change of spin-spin

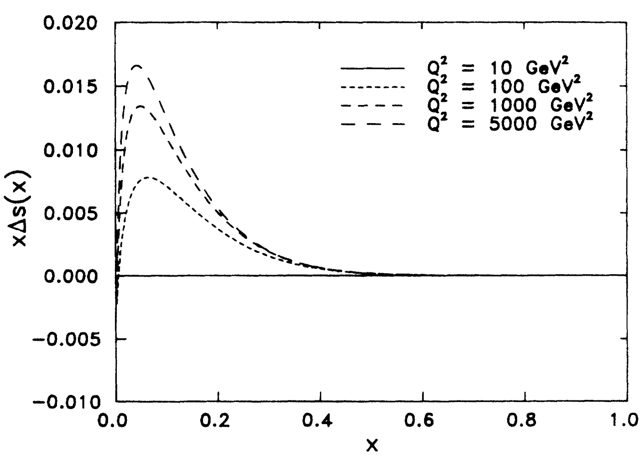

FIG. 6. The evolution of polarized sea densities at various values of Q^2 with $\Delta s(Q^2_{\rm EMC}) = 0$.

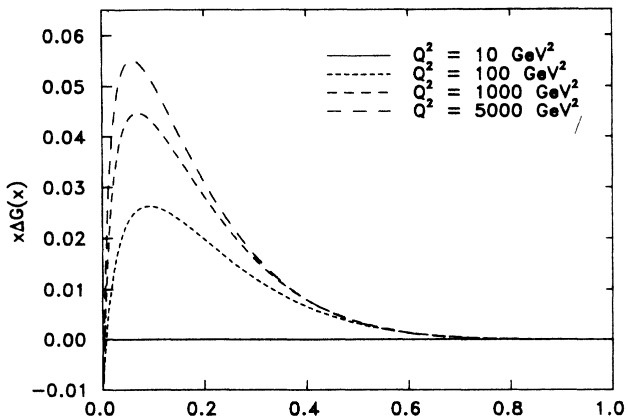

FIG. 7. Same as Fig. 5 but with $\Delta G(Q^2_{\rm EMC})=0$.

asymmetries in the Drell-Yan process as we are going to discuss in the next section.

Following Ref. 14 we have also numerically checked various sum rules derived from the AP evolution equations. The x-integrated distributions ΔG, $\Delta\Sigma$, and $\int g_1^p$ are depicted in Fig. 9 as a function of Q^2 for case (a). Evidently, ΔG increases linearly with $\ln Q^2$ while $\Delta\Sigma$ and the first moment of g_1^p are independent of Q^2, as it should be. The same quantities but integrated over the range $0.01<x<1$ are also presented in the same figure. We see a drastically different Q^2 behavior as first noticed in Ref. 14: the integrals of $\Delta\Sigma(x,Q^2)$ and $g_1^p(x,Q^2)$ over the restricted x region increase with $\ln Q^2$, whereas ΔG remains virtually unchanged. This indicates that when integrated over the range $0<x<0.01$, ΔG increases logarithmically with Q^2, while Δs rises negatively with Q^2, as we have already shown in Figs. 5 and 8, respectively. Quantitatively, Altarelli and Stirling[14] obtained $\int g_1^p=0.20$ and $\Delta\Sigma=1.1$ (integrated over the range $0.01<x<1$) at $Q^2=1000$ GeV^2, whereas we get 0.17 and 0.83, respectively, at the same Q^2. This attributes to the different parametrizations of polarized parton densities employed in the present paper and in Ref. 14. At any rate, the integrals $\int_{0.01}^{1} g_1^p dx$ and $\int_0^{0.01} g_1^p dx$ are sensitive to the way of fitting the polarized EMC data. For example, our fit to $g_1^p(x)$ shown in Fig. 2 gives the respective values 0.136 and -0.018 calculated at the EMC energies. On the other hand, one can adjust the spin densities of gluons and sea quarks at the small-x regime in such a way that a best "EMC fit" to the data (as presented in the original EMC paper[1]) implies a negligible contribution of g_1^p over the region $0<x<0.01$.

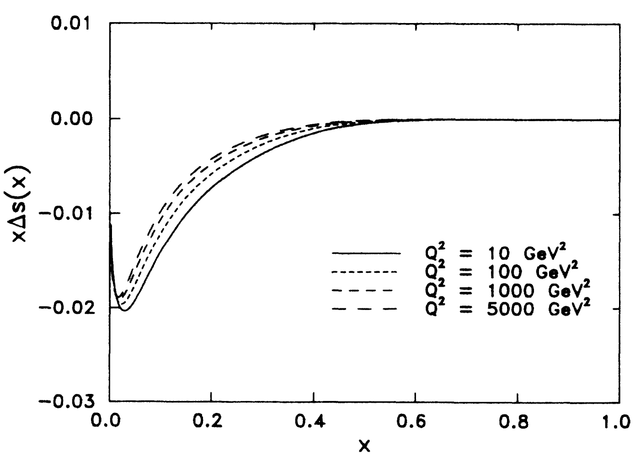

FIG. 8. Same as Fig. 6 but with $\Delta s(Q^2_{\rm EMC})=-0.11$.

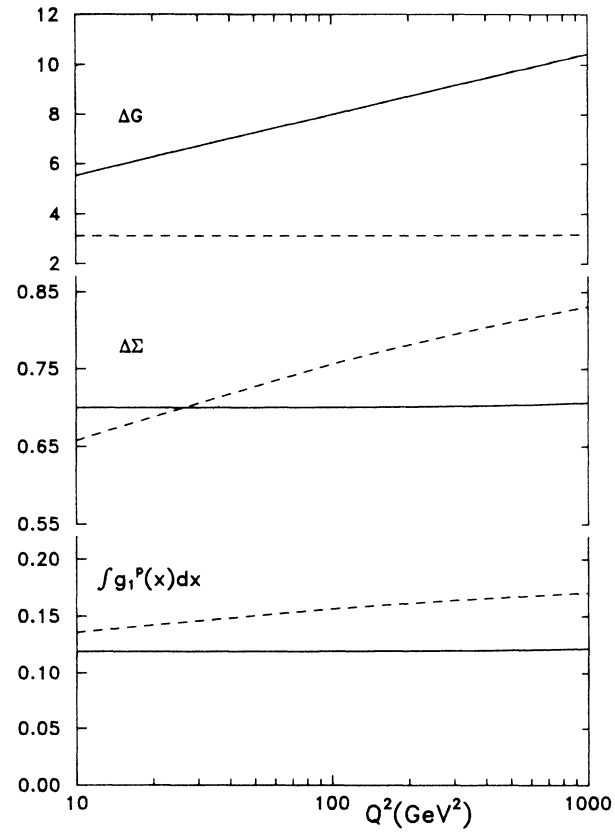

FIG. 9. The x-integrated distributions ΔG, $\Delta\Sigma$, and $\int g_1^p dx$ as a function of Q^2 with $\Delta G(Q^2_{\rm EMC})=5.5$ and $\Delta s(Q^2_{\rm EMC})=0$. Dashed lines are the same quantities integrated over the range $0.01<x<1$.

III. SPIN-SPIN ASYMMETRIES IN THE DRELL-YAN PROCESS

From the EMC data we learn that the net spin carried by gluons and the sea of $\bar{q}q$ pairs is constrained by the relation $\Delta s-(\alpha_s/4\pi)\Delta G=-0.11\pm0.03$. This suggests that gluons and/or sea quarks are strongly polarized in a polarized proton. To test various proposed schemes for Δs, the Drell-Yan experiments with polarized proton beams and targets are suitable for this purpose since the sea polarization dominates over the gluon contribution for the Drell-Yan spin asymmetry.

The longitudinal spin-spin asymmetry in the Drell-Yan process is defined by

$$A_{LL}^{\rm DY}=\frac{d\Delta\sigma/dQ^2}{d\sigma/dQ^2}=\frac{d\sigma^{\uparrow\uparrow}/dQ^2-d\sigma^{\uparrow\downarrow}/dQ^2}{d\sigma^{\uparrow\uparrow}/dQ^2+d\sigma^{\uparrow\downarrow}/dQ^2}\,, \quad (3.1)$$

where Q^2 is the invariant mass squared of lepton pairs, $d\sigma^{\uparrow\uparrow}$ ($d\sigma^{\uparrow\downarrow}$) denotes the Drell-Yan cross section for the configuration where the incoming proton spins are parallel (antiparallel). The cross section $d\sigma/dQ^2$ to next-to-leading order reads[44]

$$\frac{d\sigma}{dQ^2}=\frac{4\pi\alpha^2}{9sQ^2}\int_0^1\frac{dx_1}{x_1}\frac{dx_2}{x_2}\left[\left\{\sum_i^{n_f}e_i^2q_i(x_1,Q^2)\bar{q}_i(x_2,Q^2)\right\}\left\{\delta(1-z)+\frac{\alpha_s(Q^2)}{2\pi}\theta(1-z)w_{q\bar{q}}(z)\right\}\right.$$
$$\left.+\left\{\sum_i^{2n_f}e_i^2q_i(x_1,Q^2)G(x_2,Q^2)\right\}\frac{\alpha_s(Q^2)}{2\pi}\theta(1-z)w_{qG}(z)+(1\leftrightarrow 2)\right], \tag{3.2}$$

where $z=Q^2/(sx_1x_2)\equiv\tau/(x_1x_2)$, and

$$w_{q\bar{q}}(z)=\frac{4}{3}\left[(1+\tfrac{4}{3}\pi^2)\delta(1-z)+\frac{3}{(1-z)_+}+2(1+z^2)\left[\frac{\ln(1-z)}{1-z}\right]_+-6-4z\right],$$
$$w_{qG}(z)=\tfrac{1}{2}[(2z^2-2z+1)\ln(1-z)+\tfrac{9}{2}z^2-5z+\tfrac{3}{2}]\;. \tag{3.3}$$

Because of the large coefficient of the $\delta(1-z)$ term, it has been known for some time that the first-order corrections are quite sizable. The issues of how to sum over the large next-to-leading order corrections and how to tackle with all "plus" contributions in $w_{q\bar{q}}(z)$ were addressed recently in Ref. 45. The spin-weighted cross section $d\Delta\sigma/dQ^2$ is obtained from $-d\sigma/dQ^2$ (notice a sign flip) with the replacement[11]

$$q\to\Delta q,\quad \bar{q}\to\Delta\bar{q},\quad G\to\Delta G\;,$$
$$w_{q\bar{q}}\to\Delta w_{q\bar{q}}=w_{q\bar{q}}+\tfrac{8}{3}(1+z)\;, \tag{3.4}$$
$$w_{qG}\to\Delta w_{qG}=\tfrac{1}{2}[(2z-1)\ln(1-z)-\tfrac{3}{2}z^2+3z-\tfrac{1}{2}]\;.$$

The resulting spin-spin asymmetries[46] calculated at $\sqrt{s}=27,100$ GeV for four different characteristic values of ΔG and Δs [i.e., Eq. (2.14)] are presented in Figs. 10 and 11. The main results are summarized in the following.

(i) *A priori* it is expected that both sea and gluons are equally important for the spin asymmetry since $(\alpha_s/4\pi)\Delta G$ and Δs are of the same order of magnitude. However, owing to the smallness of Δw_{qG} it turns out that the sea polarization plays a dominant role in $A_{LL}^{\rm DY}$. Figure 11 demonstrates clearly that the gluon contribution is small at the experimentally accessible region of τ ($\tau<0.5$). This means that the sign of $A_{LL}^{\rm DY}$ is determined by the behavior of $\Delta s(x)$.

(ii) The main QCD corrections to the (polarized and unpolarized) cross sections arise from the $\delta(1-z)$ term which has a large coefficient. The next-to-leading corrections are, however, less important for helicity asymmetries.

(iii) The polarization asymmetry is negative in the absence of initial sea polarization at $Q^2_{\rm EMC}$ (Ref. 47). This is attributed to the fact that dynamically generated Δs is positive when $x>0.01$, as elaborated on in the previous section. It is also obvious from Fig. 10 that $A_{LL}^{\rm DY}$ is of positive sign in the range $\tau<0.5$ if sea is initially polarized in the direction opposite to the proton spin. Hence, *even a measurement of the sign of $A_{LL}^{\rm DY}$ will be very valuable: it will tell us whether the sea quarks of the proton is unpolarized or polarized negatively at the* EMC *energies.*

(iv) The polarization asymmetry is of order 10^{-1} and changes sign at τ=0.5, 0.7, and 0.77, respectively, for cases (b)–(d). This stems from the fact, as discussed before, that the sign of Δs flips at large x. For example, Δs changes sign at $x=0.65,0.81,0.88$, respectively, for $\Delta s(Q^2_{\rm EMC})=-0.05,-0.11,-0.20$.

(v) The relation for the nonvalence quarks $0.4\bar{u}(x)=0.4\bar{d}(x)=\bar{s}(x)$ [see Eq. (2.6)] is no longer held at large Q^2. As a consequence, $|A_{LL}^{\rm DY}|$ is not bound to be less than 0.4.

To conclude, the spin-spin asymmetry in the Drell-Yan

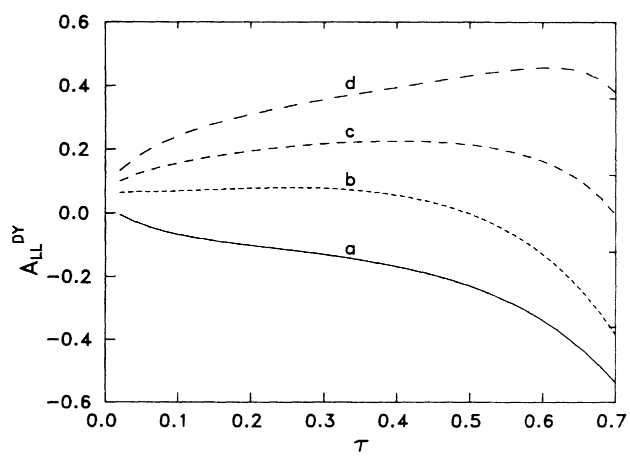

FIG. 10. Predicted Drell-Yan spin asymmetry $A_{LL}^{\rm DY}$ as a function of $\tau(=Q^2/s)$ at $\sqrt{s}=27$ GeV for four different values of ΔG and Δs described in the text.

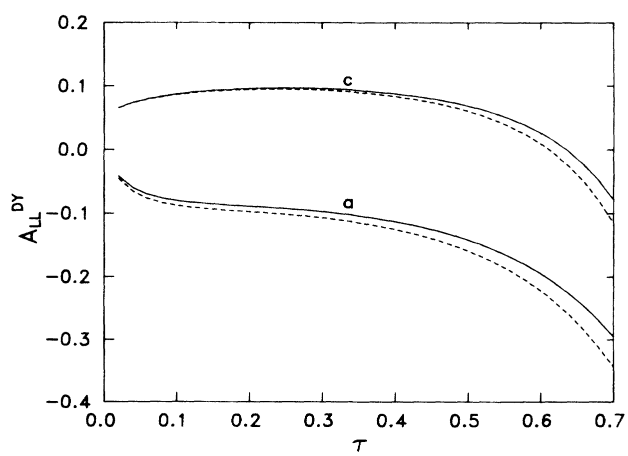

FIG. 11. Same as Fig. 10 but for $\sqrt{s}=100$ GeV. Dashed curves are the helicity asymmetries without gluon contributions.

process is sizable (of order 10^{-1}) and it provides the unique advantage: a measurement of A_{LL}^{DY} will supply valuable information on the importance of the sea polarization in a polarized proton.

IV. SPIN-SPIN ASYMMETRY IN HIGH-P_T PHOTON PRODUCTION

In the previous section we saw that the sea polarization plays an essential role in determining the spin asymmetry in the Drell-Yan process. We shall use in this section direct photon production in proton-proton collisions at high p_T as a probe of gluon polarization. At leading order, prompt photons are produced at the parton level via Compton scattering $qG\to\gamma q$ and annihilation $q\bar{q}\to\gamma G$. Since the Compton subprocess dominates over annihilation (see below), the spin asymmetry in this process depends strongly on the spin densities of the gluon.

The unpolarized invariant cross section for the reaction $p+p\to\gamma+X$ has the form[26]

$$E_\gamma \frac{d^3\sigma}{dp_\gamma^3}=\frac{\alpha\alpha_s}{s}\int_{x_{\min}}^1 dx_1\frac{1}{x_1x_2(x_1s+u)}\left[\sum_i^{n_f} e_i^2[q_i(x_1,Q^2)\bar{q}_i(x_2,Q^2)+(1\leftrightarrow 2)]\frac{d\hat\sigma}{d\hat t}(q\bar q\to\gamma G)\right.$$
$$\left.+2F_1(x_1,Q^2)G(x_2,Q^2)\frac{d\hat\sigma}{d\hat t}(qG\to\gamma q)+2G(x_1,Q^2)F_1(x_2,Q^2)\frac{d\hat\sigma}{d\hat t}(Gq\to\gamma q)\right], \tag{4.1}$$

where

$$x_2=-\frac{x_1t}{x_1s+u},\quad x_{\min}=-\frac{u}{s+t},$$
$$2F_1(x,Q^2)=\sum e_i^2q_i(x,Q^2), \tag{4.2}$$

and $d\hat\sigma/d\hat t$ is the parton-parton cross section given by

$$\frac{d\hat\sigma}{d\hat t}(q\bar q\to\gamma G)=\frac{8}{9}\left[\frac{\hat t}{\hat u}+\frac{\hat u}{\hat t}\right],$$
$$\frac{d\hat\sigma}{d\hat t}(qG\to\gamma q)=-\frac{1}{3}\left[\frac{\hat t}{\hat s}+\frac{\hat s}{\hat t}\right], \tag{4.3}$$
$$\frac{d\hat\sigma}{d\hat t}(Gq\to\gamma q)=-\frac{1}{3}\left[\frac{\hat u}{\hat s}+\frac{\hat s}{\hat u}\right],$$

with $\hat s,\hat u,\hat t$ being the Mandelstam variables for the parton subprocess

$$\hat s=x_1x_2s,\quad \hat t=x_1t,\quad \hat u=x_2u. \tag{4.4}$$

The polarized cross section $E_\gamma d^3\Delta\sigma/dp_\gamma^3$ is obtained from Eq. (4.1) with the replacement

$$q\to\Delta q,\quad \bar q\to\Delta\bar q,\quad G\to\Delta G,\quad F_1\to g_1^p \tag{4.5}$$

and $d\hat\sigma\to d\Delta\hat\sigma$, where

$$\frac{d\Delta\hat\sigma}{d\hat t}(q\bar q\to\gamma G)=-\frac{8}{9}\left[\frac{\hat t}{\hat u}+\frac{\hat u}{\hat t}\right],$$
$$\frac{d\Delta\hat\sigma}{d\hat t}(qG\to\gamma q)=-\frac{1}{3}\left[-\frac{\hat t}{\hat s}+\frac{\hat s}{\hat t}\right], \tag{4.6}$$
$$\frac{d\Delta\hat\sigma}{d\hat t}(Gq\to\gamma q)=-\frac{1}{3}\left[-\frac{\hat u}{\hat s}+\frac{\hat s}{\hat u}\right].$$

In terms of the longitudinal and transverse momenta of the photon

$$x_F=\frac{2p_\parallel}{\sqrt{s}},\quad x_T=\frac{2p_T}{\sqrt{s}} \tag{4.7}$$

one has

$$t=-\tfrac{1}{2}s[(x_F^2+x_T^2)^{1/2}-x_F],$$
$$u=-\tfrac{1}{2}s[(x_F^2+x_T^2)^{1/2}+x_F]. \tag{4.8}$$

We choose $Q^2=p_T^2/2$ in practical calculations.[12] The helicity asymmetry A_{LL}^γ is defined as in Eq. (3.1).

Since $(\alpha_s/4\pi)\Delta G$ is of the same order of magnitude as Δs, one should in principle include higher-order corrections to $E_\gamma d^3\Delta\sigma/dp_\gamma^3$ to take into account contributions from $\alpha_s\Delta G$ terms. Unfortunately, next-to-leading-order terms have been worked out only for unpolarized pp collisions[48] but not for polarized ones. Nevertheless, we will follow Ref. 12 assumming that contributions to $E_\gamma d^3\Delta\sigma/dp_\gamma^3$ up to second order in α_s amount to replacing Δq by $\Delta q'=\Delta q-(\alpha_s/4\pi)\Delta G$, as in the polarized deep-inelastic scattering [cf. Eqs. (1.8)–(1.11)]. Since the final expression of higher-order corrections to the unpolarized cross section $E_\gamma d^3\sigma/dp_\gamma^3$ is rather lengthy and complicated,[48] we will not include them here as they will not affect the main features of the spin asymmetry we are interested in.

Figure 12 presents the resulting helicity asymmetry for various gluon and sea spin densities at the energy $\sqrt{s}=100$ GeV and $x_F=0$ (i.e., $\theta_{\mathrm{c.m.}}=90°$, $\theta_{\mathrm{c.m.}}$ being the c.m. production angle). We find that even at $\theta_{\mathrm{c.m.}}=90°$, the annihilation contribution to A_{LL}^γ is at most 10% of the Compton scattering. Since A_{LL}^γ depends predominately on ΔG, the behavior of spin asymmetry shown in Fig. 12 is rather obvious: it is positive for $\Delta G>0$ [curves a and b correspond to $\Delta G(Q_{\mathrm{EMC}}^2)=5.5$, 3.0, respectively], but very small for $\Delta G(Q_{\mathrm{EMC}}^2)=0$, and negative if gluons are initially polarized opposite to the proton spin. We notice a main difference between the case

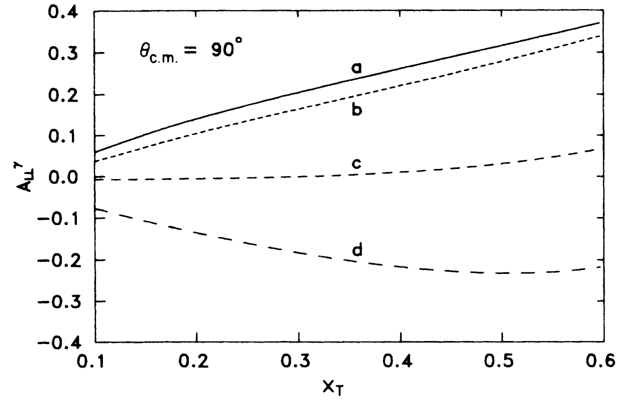

FIG. 12. Predicted direct photon spin asymmetry A_{LL}^{γ} at $\theta_{c.m.}=90°$ as a function of x_T at $\sqrt{s}=100$ GeV for four different values of ΔG and Δs described in the text.

$\Delta G(Q^2_{\rm EMC})=0$ in Fig. 12 and $\Delta s(Q^2_{\rm EMC})=0$ in Fig. 10: bremsstrahlung-induced $(\alpha_s/4\pi)\Delta G$ is very small compared to the gluon-induced Δs. Consequently, A_{LL}^{γ} in case (c) is relatively small compared to $A_{LL}^{\rm DY}$ in case (b).

Shown in Fig. 13 is the spin asymmetry as a function of x_F at fixed $p_T=5$ GeV (or $x_T=0.1$). Evidently, A_{LL}^{γ} grows with x_F at fixed x_T. This is understandable as follows: polarization asymmetry due to the annihilation subprocess alone is -1 as a result of helicity conservation owing to the vector coupling of gluons. As x_F increases (i.e., away from 90°), annihilation contribution becomes weaker and weaker and hence asymmetry is enhanced.

A correlation between $A_{LL}^{\rm DY}$ and A_{LL}^{γ} should also be noticed. For example, the curve d in A_{LL}^{γ} (Fig. 12) with negative gluon polarization must have a large and positive $A_{LL}^{\rm DY}$, and vice versa for curve a. Curve c with gluons unpolarized initially must have a negative A_{LL}^{γ} but positive $A_{LL}^{\rm DY}$. These correlations exist since Δs and ΔG are bound by the constraint Eq. (2.12). Hence, *measurements of A_{LL}^{γ} and $A_{LL}^{\rm DY}$ will not only provide a valuable probe of ΔG and Δs, but also test if gluons contribute to the spin-dependent structure function $g_1^p(x,Q^2)$ via the triangular anomaly.*

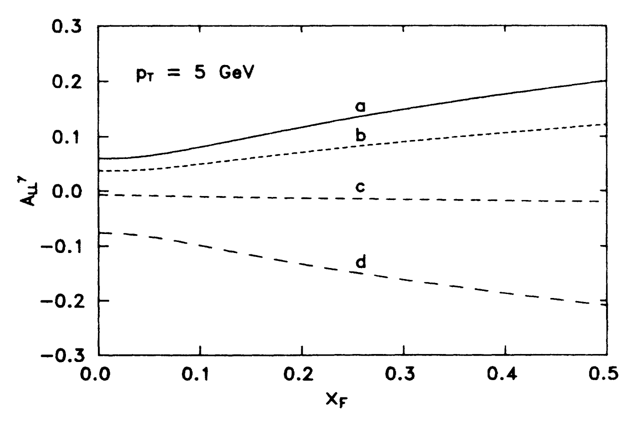

FIG. 13. Predicted direct photon spin asymmetry A_{LL}^{γ} at $p_T=5$ GeV as a function of x_F at $\sqrt{s}=100$ GeV for four different values of ΔG and Δs described in the text.

V. CONCLUSIONS

In order to determine the spin content of the proton due to gluons and sea quarks, we study in the present paper the helicity asymmetry in the Drell-Yan process and in prompt direct photon production at high p_T. The former is sensitive to the sea polarization, whereas the latter depends strongly on the gluon spin densities.

At the next-to-leading-order level of the parton model, there is one crucial difference between the first moments of unpolarized structure function $F_1(x,Q^2)$ and of spin-dependent structure function $g_1^p(x,Q^2)$: gluons contribute to the first moment of the latter but not to the former. Moreover, $[\alpha_s(Q^2)/2\pi]\Delta G(Q^2)$ is of order $(\alpha_s)^0$ rather than α_s. The recent EMC measurement determines the combination $\Delta s-(\alpha_s/4\pi)\Delta G$. Theoretical arguments for ΔG (and hence Δs), however, vary vastly, ranging from 5.5 to -4.5. We consider in this paper four different characteristic values of gluon and sea polarizations suggested by various theoretical conjectures.

To prove the polarizations of sea quarks and gluons, we need to know the evolution of parton spin densities. We have parametrized the quark polarized densities in such a way that they are consistent with the perturbative QCD suggestion and the constraint of positivity. Comparing with the EMC measurement of $xg_1^p(x)$, we find that $\Delta G \gtrsim 0$ is preferred, though the case $\Delta G<0$ is not excluded. Based on isospin symmetry, predictions of spin-dependent polarized structure function for the neutron are presented in Fig. 4. Any substantial deviation from the theoretical curves of Fig. 4 will signal the breakdown of isospin symmetry for ΔG and Δq.

The Drell-Yan process is an ideal place for testing the sea polarization: first, the helicity asymmetry $A_{LL}^{\rm DY}$ is sensitive to the sea spin densities; second, contributions from the gluon component are negligible for most values of τ. For $\Delta s(Q^2_{\rm EMC})<0$ as suggested by the EMC data, $A_{LL}^{\rm DY}$ is positive and of order 10^{-1}. However, even if the sea is initially unpolarized, the spin asymmetry turns out to be *negative* and large due to an appreciable amount of Δs evolved from large ΔG. Hence, a measurement of $A_{LL}^{\rm DY}$ will test if the sea is unpolarized or polarized opposite to the proton spin.

Contrary to the dilepton process, the spin asymmetry A_{LL}^{γ} in the direct photon production at high p_T depends strongly on the polarization of gluons as the Compton subprocess dominates over $q\bar{q}$ annihilation, reflecting by the fact that A_{LL}^{γ} grows with x_F at fixed p_T. The sign of ΔG is directly and unambiguously determined by the helicity asymmetry.

To conclude, experimental measurements of both $A_{LL}^{\rm DY}$ and A_{LL}^{γ} will provide a clean probe of sea and gluon polarizations. Moreover, they will test whether the combination $\Delta s-(\alpha_s/4\pi)\Delta G$ inferred from the EMC data is valid, i.e., whether gluons contribute to the first moment of spin-dependent structure function of the proton g_1^p via the triangular anomaly. Finally, isospin-symmetry breaking in ΔG can be tested by measuring the gluon spin component in the neutron.

ACKNOWLEDGMENTS

We are indebted to Professor W. K. Tung for kindly providing us his programs for parton evolutions. This work was supported in part by the National Science Council of the Republic of China.

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[2]For a review, see, e.g., H. Y. Cheng, talk presented at the Chinese-German Symposium on Medium-Energy Physics, Taipei, Taiwan, 1988 (unpublished).

[3]The perturbative QCD correction for the SU(3)-singlet current A_μ^0 was sometimes written as $1-\alpha_s(Q^2)/3\pi$ if the matrix elements $\langle p^\uparrow | A_\mu^a | p^\uparrow \rangle$ are measured at $\mu^2=\infty$.

[4]This interpretation was criticized recently by R. L. Jaffe and A. Manohar, revised version of Report No. CTP-1706, 1989 (unpublished).

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