Analysis of Weak-Interaction Effects in High Energy Hadron-Hadron Collisions

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Abstract

Parity-violating (pv) effects in inclusive hadron and jet productions in high energy hadron-hadron collisions are analyzed. Such effects arise from the interference between strong and weak amplitudes. This interference gives rise to a nonzero value of the pv parameters $A_L$ and $P_L$, where $A_L$ measures the difference in the inclusive cross sections of, for example, $p + p \rightarrow \text{jet} + X$ ($X=$anything), with one of incident proton beams in a state of $\pm$ helicity, and $P_L$ denotes the longitudinal polarization of a high-energy baryon (e.g., $\Lambda$) produced in $p + p \rightarrow \Lambda + X$ with the initial proton beams unpolarized. In the present paper, the single helicity asymmetry $A_L$ in one-jet, two-jet and two-jet plus photon productions as well as in the Drell-Yan process $p + p \rightarrow \ell^+\ell^- + \text{jet} + X$ is probed, and the longitudinal polarization $P_L$ of the $\Lambda$ produced in unpolarized $pp$ collisions is studied. We conclude that the pv effects in high energy proton-proton collisions are in general only sensitive to the spin dependent valence quark distributions.
I. Introduction

Hadron colliders with high-energy polarized proton beams are conceivably available in the future at RHIC, SSC and LHC. Depending on whether the proton beams are polarized longitudinally or transversely, parton spin densities of the proton can be probed via the studies of helicity or transverse spin asymmetries. With longitudinal polarization, the double helicity asymmetry defined by

\[ A_{LL} = \frac{d\sigma^{++} - d\sigma^{+-}}{d\sigma^{++} + d\sigma^{+-}} \]  

is the observable most commonly discussed in the literature, where \( d\sigma^{++} \) (\( d\sigma^{+-} \)) denotes the inclusive cross section for the configuration where the incoming hadron’s longitudinal polarizations are parallel (antiparallel). Double asymmetries at high energies have been investigated for various processes, such as single-jet \([1]\), two-jet \([1,2]\), two-jet plus photon \([3]\) and three-jet \([3,4]\) productions, double-photon production \([5]\), direct photon production at large transverse momentum \([6,7]\), and the Drell-Yan process \([7]\). Most recent works were motivated by the European Muon Collaboration (EMC) measurement of the polarized proton structure function \( g_1^p(x) \) \([8]\). The central issue of much theoretical controversy is whether or not gluons contribute to the first moment of \( g_1^p(x) \). Two extreme possibilities for the explanation of the EMC experiment have been explored in the past: large (negative) sea polarization \([9]\) or large (positive) gluon polarization \([10]\). Measurements of aforementioned processes will help determine the spin dependent parton distributions and shed light on the interpretation of the EMC results.

Contrary to the previous works, the purpose of the present paper is to analyze the single helicity asymmetry \( A_L \) defined in Eq.(2.1) in high energy proton-proton collisions. Experimentally, it should be easier to measure \( A_L \) than the double helicity asymmetry. However, theoretically a nonzero \( A_L \) can occur only if some of the parton-parton scatterings involve parity-violating weak interactions. Therefore, single helicity asymmetry can be used to probe parity violation in parton-parton subprocesses. Another party-violating (pv) effect of interest is the longitudinal polarization \( P_L \) of a high-energy baryon [see Eq.(2.2)] produced from unpolarized incident proton beams. Owing to the small size of weak effects, pv parameters \( A_L \) and \( P_L \) arise from the coherent interference between the strong-QCD and weak amplitudes.
Such pv effects were first analyzed in Ref.[11] and subsequently in Ref.[13]. Specifically, the asymmetry parameter $A_L$ for the processes $p + p \to \pi^+ + X$ and $p + p \to \text{jet} + X$, and the longitudinal polarization of Λ’s in $p + p \to \Lambda + X$ were studied in Ref.[11]. The content of the present work is in some sense the extension of the previous analysis of Ref.[11].

This paper is organized as follows. In Section II we discuss the general formulism for calculating the pv parameters $A_L$ and $P_L$. It is stressed that inspired by the EMC experiment and armed with the phenomenologically determined valence-quark spin densities, two of us (H.Y.C. and C.F.W.) have extracted the polarized sea and gluon distributions from the EMC data for several different possibilities [14]. The pv effects in one-jet, two-jet, two-jet plus photon productions are investigated in Sections III and IV. The Drell-Yan process $p + p \to \ell^+\ell^- + \text{jet} + X$ is studied in Section V. The results are discussed in Section VI, where a series of figures will be presented.

II. General formalism

In high energy proton-proton collisions, there are two single-helicity asymmetry observables which we are interested in:

$$A_L = \frac{d\sigma^+ - d\sigma^-}{d\sigma^+ + d\sigma^-}, \quad (2.1)$$

and

$$P_L = \frac{d\sigma_+ - d\sigma_-}{d\sigma_+ + d\sigma_-}. \quad (2.2)$$

In Eq.(2.1) $d\sigma^\pm$ denote the inclusive cross sections for $pp$ scattering where one of the initial proton beams is longitudinally polarized and has $\pm$ helicity. In Eq.(2.2) $d\sigma_\pm$ denote the cross sections for producing a high-energy baryon (e.g. Λ) in a state of $\pm$ helicity from unpolarized proton beams. Both parameters are expected to vanish to all orders in strong interactions. This can be easily seen in the quark-parton model where the unpolarized inclusive differential cross section for $pp$ collisions is given by

$$d\sigma = \sum_{i,j} \int dx_1 dx_2 f_i(x_1, Q^2) f_j(x_2, Q^2) \left( \frac{d\hat{s}_{ij}}{d\hat{t}} \right) d\hat{t}, \quad (2.3)$$

1 This interference effect was also briefly discussed in Ref.[12].

2 See Section VI for comments on the pv parameter $A_L$ for $W^\pm$ and $Z^0$ productions in $pp$ collisions.
where \( f_i(x, Q^2) \) is the unpolarized distribution function of the parton \( i \) in a proton with momentum fraction \( x \), and \( d\hat{\sigma}_{ij} \) is the cross section for the interaction of two partons \( i \) and \( j \). When one of the initial proton beams is longitudinally polarized, we have

\[
\begin{align*}
d\sigma^+ - d\sigma^- &= \sum_{i,j} \int dx_1 dx_2 \Delta f_i(x_1, Q^2) f_j(x_2, Q^2) \left[ \frac{d\hat{\sigma}}{dt}(i^+ j \rightarrow kl) - \frac{d\hat{\sigma}}{dt}(i^- j \rightarrow kl) \right] dt, (2.4)
\end{align*}
\]

with

\[
\Delta f(x, Q^2) = f_+(x, Q^2) - f_-(x, Q^2), \quad f(x, Q^2) = f_+(x, Q^2) + f_-(x, Q^2), \quad (2.5)
\]

where \( f_\pm(x, Q^2) \) is the parton distribution function in a polarized proton with helicity parallel (antiparallel) to the proton spin, and \( d\hat{\sigma}(i^\pm j \rightarrow kl) \) is the cross section for the scattering \( ij \rightarrow kl \) when parton \( i \) has \( \pm \) helicity. Since parity is conserved by strong interactions, it is evident that \( d\sigma^+ = d\sigma^- \) and hence \( A_L = 0 \).

If some of the parton-parton scattering subprocesses involve parity-violating weak interactions, then in general \( d\hat{\sigma}(i^+ j \rightarrow kl) \neq d\hat{\sigma}(i^- j \rightarrow kl) \) and thus a nonzero \( A_L \) is expected. Owing to the small size of the weak effects, the party-violating asymmetry \( A_L \) will arise from the coherent interference between the strong-QCD amplitude, which is parity conserving, and the parity-violating weak amplitude. Likewise, the longitudinal polarization \( P_L \) of a high-energy baryon, say \( \Lambda \), produced in \( p + p \rightarrow \Lambda + X \) with unpolarized initial proton beams would also arise from the interference of strong and weak amplitudes. Explicitly,

\[
\begin{align*}
d\sigma_+ - d\sigma_- &= \sum_{i,j,k,l} \int dx_1 dx_2 dz f_i(x_1, Q^2) f_j(x_2, Q^2) \\
&\quad \times \left[ \frac{d\alpha}{dt}(ij \rightarrow k^+l) - \frac{d\alpha}{dt}(ij \rightarrow k^-l) \right] dt \Delta D^\Lambda_k(z), (2.6)
\end{align*}
\]

with

\[
\Delta D^\Lambda_k(z) = D^\Lambda_{k^+}(z) - D^\Lambda_{k^-}(z), \quad (2.7)
\]

where \( d\sigma_\pm \) is the differential cross section for producing parton \( k \) with \( \pm \) helicity, and \( D^\Lambda_{k^\pm}(z) \) is the probability that parton \( k \) with \( \pm \) helicity decays into \( \Lambda \) with \( + \) helicity and fractional momentum \( z \).

In order to estimate the party-violatin effects \( A_L \) and \( P_L \), we need input of the polarized parton distribution functions \( \Delta f_i(x, Q^2) \) and the polarized fragmentation functions \( \Delta D^\Lambda_q(z) \).
Some useful information on the quark and gluon spin densities can be obtained from the measurement of the polarized proton structure function $g_1^p(x)$ in deep inelastic lepton-nucleon scattering. Denoting the spin-dependent parton distributions by

$$\Delta q(x) = q^\uparrow(x) + \bar{q}^\uparrow(x) - q^\downarrow(x) - \bar{q}^\downarrow(x),$$

$$\Delta G(x) = G^\uparrow(x) - G^\downarrow(x),$$

(2.8)

we see that $\Delta q = \int_0^1 \Delta q(x) dx$ and $\Delta G = \int_0^1 \Delta G(x) dx$ represent the net helicities carried by the quark flavor $q$ and the gluon, respectively, in the infinite-momentum frame of the proton with $+\helicity$. Since sea-quark and gluon polarizations are manifest essentially in the region $x < 0.1$, the valence-quark spin densities at $x > 0.1$ are constrained by the SLAC and EMC measurements of $g_1^p(x)$. Following Ref.[14], the valence quark spin densities parametrized at $Q_0^2 = 10 GeV^2$ are given by

$$\Delta u_v(x) = x^{0.287}u_v(x),$$

$$\Delta d_v(x) = \left(\frac{x-x_0}{1-x_0}\right)x^p d_v(x),$$

(2.9)

where $p = 0.03, 0.26, 0.76$, for $x_0 = 0.35, 0.50$ and $0.75$, respectively. For unpolarized parton distribution functions we use the “average” set of parametrization given in Ref.[15] extracted also at the same reference scale $Q_0^2 = 10 GeV^2$ from several experiments. Eq.(2.9) follows from the perturbative QCD suggestion [16] that valence quarks at $x = 1$ remember the spin of the parent proton but become totally unpolarized as $x$ is close to zero, and from the relations

$$\Delta u_v + \Delta d_v = 3F - D, \quad \Delta u_v - \Delta d_v = F + D,$$

(2.10)

with $F$ and $D$ being SU(3) parameters determined from neutron and hyperon $\beta$ decays. We find that this simple parametrization for $\Delta u_v(x)$ and $\Delta d_v(x)$ fits to the data of $g_1^p(x)$ very well for $x > 0.2$.

In principle, the sea-quark polarization function $\Delta q_s(x)$ and the gluon spin density $\Delta G(x)$ are constrained by the EMC measurement of $g_1^p(x)$ at small $x$. However, the issue of whether or not gluons contribute to the first moment of $g_1^p(x)$ has been under hot debate over the past few years. Then it was realized that the size of hard-gluonic contribution to $\int_0^1 g_1^p(x)dx$ is purely a matter of the factorization convention chosen in defining the quark spin densities $\Delta q(x)$ and the hard cross section for photon-gluon scattering $\Delta \sigma^{\gamma G}(x)$ [17]. A change of the
factorization scheme merely shifts the contribution between $\Delta \sigma^{\gamma G}(x)$ and $\Delta q(x)$ in such a way that the polarized proton-photon cross section remains unchanged. Depending on how the data of the polarized proton structure function are explained, two of us (H.Y.C. and C.F.W.) have determined $\Delta q_s(x)$ and $\Delta G(x)$ for three different possibilities from the EMC measurement of $g_1^p(x)$ in conjunction with the above phenomenologically determined valence-quark spin densities and the positivity constraint of the unpolarized parton distributions [14].

In the case of the sea-quark interpretation of the EMC experiment, we obtain

\[
\text{case 1 : } \begin{cases} 
\Delta s(x) = -11.8 x^{0.94} (1 - x)^5 s(x), \\
\Delta G(x) = 0,
\end{cases} \tag{2.11}
\]

where we have assumed SU(3) invariance for sea polarizations $\Delta u_s(x) = \Delta d_s(x) = \Delta s(x)$. On the contrary, if the data of $g_1^p(x)$ are explained in terms of polarized gluon and valence-quark densities, we find that an acceptable $\Delta G(x)$ does exist for properly chosen prescription; more precisely,

\[
\text{case 2 : } \begin{cases} 
\Delta s(x) = 0, \\
\Delta G(x) = 6.0 x^{0.76} (1 - x)^3 G(x),
\end{cases} \tag{2.12}
\]

Since in a realistic case, it is unlikely that $\Delta G(x, Q^2)$ or $\Delta s(x, Q^2)$ vanishes at some scale $Q_0^2$ for all $x$, it is more pertinent to consider the case with non-vanishing $\Delta G(x)$ and $\Delta s(x)$. In Ref.[14] we have proposed more realistic spin-dependent sea and gluon distribution functions which are parametrized in such a way that the first moment of $g_1^p(x)$ receives almost all contributions from the region $x > 0.01$, as indicated by the EMC experiment:

\[
\text{case 3 : } \begin{cases} 
\Delta s(x) = -3.39 x^{0.62} (1 - x)^{1.4} s(x), \\
\Delta G(x) = 2.69 x^{0.76} (1 - x)^3 G(x).
\end{cases} \tag{2.13}
\]

Case 3 is between the sea-quark (case 1) and gluon (case 2) interpretation of the EMC data. The above three sets of polarized sea and gluon distribution functions are all parametrized at the reference scale $Q_0^2 = 10 \text{ GeV}^2$. Their $Q^2$ evolutions are obtained by solving the Altarelli-Parisi equation numerically [18].

Very little is known about the fragmentation function $D_\Lambda^q(z)$. Just as the case of parton distribution functions, we may assume that $\Delta D_\Lambda^q(z)$ is proportional to the unpolarized fragmentation function $D_\Lambda^q(z)$. For these we take the simple parametrization [19]

\[
z D_\Lambda^q(z) \equiv z D_\Lambda^u(z) = z D_\Lambda^d(z) = z D_\Lambda^s(z) = B_v z(c - z)^3 + B_s (1 - z)^4, \tag{2.14}
\]
and
\[
zD^A_q(z) \equiv zD^A_q(z) = zD^A_d(z) = zD^A_s(z) = B_s(1 - z)^4. \tag{2.15}
\]
Assuming \(D^A_q(z) \approx D^B_q(z)\), we find that a fit to the EMC data of \(D^B_q(z)\) \cite{20} yields
\[
c = 1.08, \quad B_v = 0.5, \quad B_s = -0.06. \tag{2.16}
\]

For the polarized fragmentation functions, the construction of \(\Delta D^A_q(z)\) is necessarily somewhat \textit{ad hoc} in the absence of experimental data or a detailed theory. In analog to the quark distributions, we simply assume that
\[
\Delta D^A_q(z) = z^\gamma D^A_q(z). \tag{2.17}
\]
That is, the polarization of the outgoing \(\Lambda\) is equal to that of the parent quark at \(z = 1\) but diminishes as \(z \to 0\). Recently it has been pointed out by Burkardt and Jaffe \cite{21} that a measurement of the helicity asymmetric cross section for semi-inclusive production of \(\Lambda\) in \(e^+e^-\) annihilation near the \(Z^0\) resonance allows a complete determination of \(\Delta D^A_q(z)\). This experiment should be practical at the LEP collider at CERN or at SLC at SLAC.

### III. Hadronic jet and \(\Lambda\) productions

In this section we shall study the asymmetry parameter \(A_L\) for \(pp \to\) jet + \(X\) and \(pp \to 2\) jets + \(X\) and the longitudinal polarization \(P_L\) for \(pp \to \Lambda + X\). Hadronic jet productions are expected to be the main processes in high-transverse-momentum hadron-hadron collisions. At the parton level the parton-parton cross sections are the same for 1-jet, 2-jet and \(\Lambda\) productions, so we put them all together in this section.

As emphasized in passing, we are interested in the differential cross section \(d\sigma/d\hat{t}(ij \to kl)\) arising from the coherent interference between the strong QCD amplitude and the weak amplitude. In order for there to exist any interference, it is necessary that these amplitudes connect the same initial and final states. Since the strong QCD interactions change color but conserve flavor, while weak couplings always conserve color, only a limited number of quark and gluon scattering processes can give rise to a strong-weak interference. Evidently, the external gluons in general cannot make contributions to \(A_L\) and \(P_L\) at tree level. The strong quark-quark scattering processes can be classified into five categories: \(qq \to qq, qq' \to\)
qq' (q' ≠ q), q̅q → q̅q, q̅q → q'̅q', q̅q → q̅q'. The differential cross sections for quark-quark scattering due to QCD and weak interference with one of the initial quarks polarized are given in Ref.[11]. For our purposes, it is more convenient to present these cross sections for both initial quarks or antiquarks having definite helicities. The color-averaged differential cross section due to the interference between the strong and weak amplitudes is given by

\[
\frac{d\hat{\sigma}_{int}}{d\hat{t}} = \frac{1}{32\pi s^2} \left(\frac{4}{9}\right) |M(\lambda_1, \lambda_2, \lambda_3, \lambda_4)|^2,
\]

\[
= \frac{16 G_F}{9 \sqrt{2}} \alpha_s |T(\lambda_1, \lambda_2, \lambda_3, \lambda_4)|^2,
\]

(3.1)

for the quark-quark scattering \(q_1(k_1, \lambda_1)q_2(k_2, \lambda_2) \rightarrow q_3(k_3, \lambda_3)q_4(k_4, \lambda_4)\), where \(\frac{4}{9}\) is a color-averaged factor. The results are summarized in Table I in four different helicity states denoted by

\[
\begin{pmatrix}
q_{1L}q_{2L} & q_{1L}q_{3R} \\
q_{1R}q_{2L} & q_{1R}q_{2R}
\end{pmatrix}.
\]

(3.2)

Note that in Table I,

\[L_q = 2T_3 - 2e_q \sin^2 \theta_W, \quad R_q = -2e_q \sin^2 \theta_W,
\]

(3.3)

are the left-handed and right-handed coupling constants respectively of the quark coupled with the Z boson, where \(T_3\) is an SU(2) isospin quantum number, \(e_q\) is the charge of the quark, and \(\theta_W\) is the Weinberg angle. In Eq.(3.1) \(\hat{s}\), \(\hat{t}\) and \(\hat{u}\) are the usual Mandelstam variables

\[\hat{s} = (k_1 + k_2)^2, \quad \hat{t} = (k_1 + k_3)^2, \quad \hat{u} = (k_2 + k_3)^2,
\]

(3.4)

with \(k_1 + k_2 + k_3 + k_4 = 0\), \(\Gamma_W (\Gamma_Z)\) is the decay width of the W (Z) boson, and \(V_{qq'}\) is a quark mixing matrix element. It should be stressed that our results for \(d\hat{\sigma}_{int}/d\hat{t}\) are in agreement with that first derived in Ref.[11] but differ from Ref.[13] in signs for some of the subprocesses.

The weak amplitudes for the processes \(qq' \rightarrow qq'\), \(q̅q \rightarrow q'̅q'\) and \(qq' \rightarrow q̅q'\) receive both charged- and neutral-current contributions. However, it is clear from Table I that only the weak amplitudes due to W exchange interfer with the strong amplitude. Since only left-handed quarks and right-handed antiquarks participate in weak couplings with the W boson, this explains why there is only one non-vanishing matrix element for above-mentioned quark-quark scattering processes.
Table I. The differential cross section $d\hat{\sigma}_{int}/d\hat{t}$ for various quark-quark scattering processes due to the coherent interference between strong and weak amplitudes. Shown are the matrix elements absolute squared $|T|^2$ [see Eq.(3.1)] in four different helicity states denoted by (3.2).

| process            | $|T|^2$                                                                 |
|--------------------|------------------------------------------------------------------------|
| $qq \rightarrow qq$| $\left[ \frac{M_Z^2}{t(u-M_Z^2)} + \frac{M_Z^2}{u(t-M_Z^2)} \right] \begin{pmatrix} L_q^2 & 0 \\ 0 & R_q^2 \end{pmatrix}$ |
| $qq' \rightarrow qq'$| $-\frac{8M_W^2}{t(u-M_W^2)} |V_{qq'}|^2 \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix}$ |
| $q\bar{q} \rightarrow q\bar{q}$| $\frac{\hat{s}^2}{s^2} \left[ \frac{(\hat{s}-M_Z^2)M_Z^2}{t[(\hat{s}-M_Z^2)^2+\Gamma_Z^2M_Z^2]} + \frac{M_Z^2}{s(t-M_Z^2)} \right] \begin{pmatrix} 0 & L_q^2 \\ R_q^2 & 0 \end{pmatrix}$ |
| $q\bar{q} \rightarrow q'\bar{q}'$| $-\frac{8\hat{s}^2}{s^2} \frac{M_W^2}{s(t-M_W^2)} |V_{qq'}|^2 \begin{pmatrix} 0 & 1 \\ 0 & 0 \end{pmatrix}$ |
| $q'\bar{q} \rightarrow q'\bar{q}'$| $-\frac{8\hat{s}^2}{s^2} \frac{(\hat{s}-M_W^2)M_W^2}{t[(\hat{s}-M_W^2)^2+\Gamma_W^2M_W^2]} |V_{qq'}|^2 \begin{pmatrix} 0 & 1 \\ 0 & 0 \end{pmatrix}$ |

It follows from Eq.(2.4) that the differential cross section for a single jet production in $pp$ collisions at transverse momentum $p_T$ and rapidity $y$ is given by

$$E \frac{d\sigma^+}{d^3p} - E \frac{d\sigma^-}{d^3p} = \frac{1}{\pi} \sum_{i,j} \int_{x_0}^1 dx_a \frac{x_a x_b}{2x_a - x_T e^y} \Delta f_i(x_a, Q^2) f_j(x_b, Q^2) \times \left[ \frac{d\hat{\sigma}}{dt}(i^+ j \rightarrow kl) - \frac{d\hat{\sigma}}{dt}(i^- j \rightarrow kl) \right],$$

(3.5)

where

$$x_T = \frac{2p_T}{\sqrt{s}}, \quad x_0 = \frac{x_T e^y}{2 - x_T e^{-y}}, \quad x_b = \frac{x_a x_T e^{-y}}{2x_a - x_T e^y},$$

(3.6)

and it is to be understood that

$$\frac{d\hat{\sigma}}{dt}(i^+ j \rightarrow kl) - \frac{d\hat{\sigma}}{dt}(i^- j \rightarrow kl) = \frac{1}{2} \left[ \frac{d\hat{\sigma}}{dt}(i^+ j^+ \rightarrow kl) + \frac{d\hat{\sigma}}{dt}(i^+ j^- \rightarrow kl) - \frac{d\hat{\sigma}}{dt}(i^- j^- \rightarrow kl) - \frac{d\hat{\sigma}}{dt}(i^- j^+ \rightarrow kl) \right].$$

(3.7)

For the production of two jets with rapidities $y_1$ and $y_2$, and with equal and opposite trans-
verse momentum $p_T$, the differential cross section reads [22]

$$\frac{d\sigma^+}{dy_1dy_2dp_T^2} - \frac{d\sigma^-}{dy_1dy_2dp_T^2} = x_ax_b \sum_{i,j} \Delta f_i(x_a, Q^2)f_j(x_b, Q^2)$$

$$\times \left[ \frac{d\hat{\sigma}}{dt}(i^+ j \to kl) - \frac{d\hat{\sigma}}{dt}(i^- j \to kl) \right], \quad (3.8)$$

where

$$x_a = \frac{1}{2} x_T \left( \tan^{-1} \frac{\theta_1}{2} + \tan^{-1} \frac{\theta_2}{2} \right), \quad x_b = \frac{1}{2} x_T \left( \tan \frac{\theta_1}{2} + \tan \frac{\theta_2}{2} \right), \quad (3.9)$$

with $\theta_1$ ($\theta_2$) being the angle between the first (second) jet and one of the incident proton beams (see Ref.[23] for detail), and $\theta_1 + \theta_2 = 180^\circ$. As for the longitudinal polarization $P_L$ of $\Lambda$ produced in $pp \to \Lambda + X$, the numerator of Eq.(2.2) is given by [23]

$$E \frac{d\sigma^+}{d^3p} - E \frac{d\sigma^-}{d^3p} = \frac{1}{\pi} \sum_{i,j,k,l} \int_{x_{a_{\min}}}^1 dx_a \int_{x_{b_{\min}}}^1 dx_b f_i(x_a, Q^2)f_j(x_b, Q^2)$$

$$\times \left[ \frac{d\hat{\sigma}}{dt}(ij \to k^+l) - \frac{d\hat{\sigma}}{dt}(ij \to k^-l) \right] \frac{1}{z} \Delta D^\Lambda_k(z), \quad (3.10)$$

where

$$x_{a_{\min}}^{\min} = -\frac{u}{s + t}, \quad x_{b_{\min}}^{\min} = -\frac{tx_a}{u + sx_a}, \quad z = -\frac{t}{sx_b} - \frac{u}{sx_a}. \quad (3.11)$$

The denominator of (2.1) and (2.2) is twice the unpolarized cross section, whose explicit expression is not written here. The general features of unpolarized cross sections are known.

Since gluon densities are large at small $x$, gluon-gluon scattering dominates the underlying parton-parton interaction subprocesses at small $x_T$. As the jet momentum increases, quark-gluon scattering becomes more and more important due to the relatively fast decrease of the gluon distribution with increasing $x$. It is finally governed by quark-quark scattering at large $x_T$.

IV. Two-jet plus photon production

At the parton level, the production of 2-jet plus photon final state at large transverse momentum can proceed via processes (denoted by $H_i$) where a photon is emitted from a quark or antiquark, and processes (denoted by $W_i$) where a pair $V - \gamma$ is produced with $V(= W^\pm$ or $Z^0)$ decaying hadronically. The amplitudes for $H_i$ and all other amplitudes for 2-2 parton scattering involving at least one quark or anti-quark off which $V$ is radiated, give
rise to the dominant QCD background. The amplitudes for $W_i$ with $V$ being far off shell yield a much smaller cross section of order $\alpha^3$, compared with the amplitudes $H_i$, whose corresponding cross sections are of order $\alpha\alpha_s^2$. However, the former cross section can be possibly greatly enhanced if $V$ is on its mass shell. This can be seen by comparing the boson propagator $i/(k^2 - M_V^2 + iM_V\Gamma_V)$ with that of a massless quark or gluon $i/(k^2 + i\epsilon)$ with roughly the same momentum $k$. Since $M_V^2/\Gamma_V^2 \approx 1.4 \times 10^3$ for $V = W^\pm$ or $Z^0$, it should help this production cross section rise above the QCD background.

The tree-level processes which give rise to two-jet plus one real photon production events, have five possible classes, $qq \rightarrow qq\gamma$, $q'q' \rightarrow qq'\gamma$, $q\bar{q} \rightarrow q\bar{q'}\gamma$, and $q\bar{q'} \rightarrow q\bar{q'}\gamma$. We shall only write down those interference terms involving QCD and electroweak interactions.

For the case of the subprocess $q(k_1)q(k_2) \rightarrow q(k_3)q(k_4)\gamma(k_5)$, the absolute value squared of spin dependent matrix-elements is given by

$$|M(++)|^2 = \frac{\hat{s}_{12}^2 + \hat{s}_{34}^2}{\hat{s}_{45}} e_q^2 R_q \left\{ (f_3 - g_2) \left[ \frac{1}{\hat{s}_{23}(\hat{s}_{24} - M_W^2)} + \frac{1}{\hat{s}_{24}(\hat{s}_{23} - M_W^2)} \right] 
- g_3 \left[ \frac{1}{\hat{s}_{23}(\hat{s}_{13} - M_Z^2)} + \frac{1}{\hat{s}_{13}(\hat{s}_{23} - M_Z^2)} \right] 
+ (-g_1 + g_2 + g_3 - f_1) \left[ \frac{1}{\hat{s}_{14}(\hat{s}_{24} - M_Z^2)} + \frac{1}{\hat{s}_{24}(\hat{s}_{14} - M_Z^2)} \right] 
- (f_1 + g_1) \left[ \frac{1}{\hat{s}_{14}(\hat{s}_{13} - M_Z^2)} + \frac{1}{\hat{s}_{13}(\hat{s}_{14} - M_Z^2)} \right] \right\},$$

where $\hat{s}_{ij} = (k_i + k_j)^2$, and

$$f_1 = 2\frac{\hat{s}_{12}}{\hat{s}_{35}}, \quad f_2 = 2\frac{\hat{s}_{13}}{\hat{s}_{25}}, \quad f_3 = 2\frac{\hat{s}_{23}}{\hat{s}_{15}},$$

$$g_1 = \frac{\hat{s}_{25}\hat{s}_{13} + \hat{s}_{35}\hat{s}_{12} - \hat{s}_{23}\hat{s}_{15}}{\hat{s}_{25}\hat{s}_{35}}, \quad g_2 = \frac{\hat{s}_{25}\hat{s}_{13} - \hat{s}_{15}\hat{s}_{23} - \hat{s}_{25}\hat{s}_{12}}{\hat{s}_{35}\hat{s}_{15}},$$

$$g_3 = \frac{\hat{s}_{15}\hat{s}_{23} + \hat{s}_{25}\hat{s}_{13} - \hat{s}_{12}\hat{s}_{35}}{\hat{s}_{15}\hat{s}_{25}},$$

and $|M(--)|^2$ is obtained from $|M(++)|^2$ with the replacement $R_q \rightarrow L_q$. For the case of $qq' \rightarrow qq'\gamma$ scattering, the non-vanishing matrix element absolute squared is

$$|M(--)|^2 = \frac{\hat{s}_{12}^2}{\hat{s}_{45}} \left\{ -\frac{e_q}{\hat{s}_{23}} (Ag_2 + Bg_3 + Cf_3) 
+ \frac{e_{q'}}{\hat{s}_{14}} \left[ A(f_1 + g_1) + B(f_2 + g_1) + C(g_2 + g_3) \right] \right\} |V_{qq'}|^2$$

$$+ \frac{\hat{s}_{34}^2}{\hat{s}_{45}} \left[ \frac{e_{q'}}{\hat{s}_{13} - M_W^2} g_2 - \frac{e_{q'}}{\hat{s}_{14}} (f_1 + g_1) \right]$$

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+ \frac{e_q g_3 - e_{q'} (g_1 + f_2)}{s_{14}} - C \left[ \frac{e_q}{s_{14}} f_3 - \frac{e_{q'}}{s_{14}} (g_2 + g_3) \right] \right) |V_{qq'}|^2, \quad (4.3)

where

\begin{align*}
A &= \frac{\hat{s}_{35} (\hat{s}_{25} + \hat{s}_{45})}{\hat{s}_{12} (\hat{s}_{13} - M_w^2) (\hat{s}_{24} - M_w^2)} + \frac{\hat{s}_{35}}{\hat{s}_{12} (\hat{s}_{13} - M_w^2)} \\
&\quad - \left( \frac{\hat{s}_{34}}{\hat{s}_{12}} e_q + e_{q'} \right) \frac{1}{\hat{s}_{24} - M_w^2}, \\
B &= \frac{e_{q'}}{s_{13} - M_w^2} + \frac{\hat{s}_{25}}{(\hat{s}_{13} - M_w^2) (\hat{s}_{24} - M_w^2)}, \\
C &= \frac{e_q}{\hat{s}_{12} (\hat{s}_{13} - M_w^2)}. \quad (4.4)
\end{align*}

For $\bar{q}q' \to \bar{q}'q' \gamma$ scattering, the nonvanishing matrix element squared is $|M(++)|^2$, which has the same expression as Eq.(4.3). Note that in Eq.(4.3) we have applied the relation $e_q + e_{q'} = 1$ whenever it holds.

For the case of $q\bar{q} \to q\bar{q}'\gamma$ subprocess, we have

\begin{align*}
|M(+-)|^2 &= e_q^2 R^2 \frac{s_{13}}{s_{45}} \left\{ g_2 \left[ \frac{r_{12}^z}{s_{23}} + \frac{1}{s_{12} (\hat{s}_{23} - M_Z^2)} \right] \\
&\quad + (f_1 + g_1) \left[ \frac{r_{12}^z}{s_{14}} + \frac{1}{s_{12} (\hat{s}_{14} - M_Z^2)} \right] \\
&\quad + (f_2 + f_3 + g_3) \left[ \frac{r_{34}^z}{s_{23}} + \frac{1}{s_{34} (\hat{s}_{23} - M_Z^2)} \right] \\
&\quad + (g_1 + g_2 + g_3) \left[ \frac{r_{34}^z}{s_{14}} + \frac{1}{s_{34} (\hat{s}_{14} - M_Z^2)} \right] \right\} \\
&\quad + e_{q'}^2 R^2 \frac{s_{13} - s_{24}}{s_{45}} \left\{ \frac{1}{D_{12}^z} \left( \frac{g_1'}{s_{14}} + \frac{g_2'}{s_{23}} \right) + \frac{1}{D_{34}^z} \left( \frac{g_3'}{s_{23}} - \frac{g_1' + g_2' + g_3'}{s_{14}} \right) \right\}, \quad (4.5)
\end{align*}

where $D_{ij}^z = (\hat{s}_{ij} - M_Z^2)^2 + M_Z^2 \Gamma_Z^2$, $r_{ij}^z = (\hat{s}_{ij} - M_Z^2) / D_{ij}^z$, and

\begin{equation}
g_1' = \frac{4}{\hat{s}_{25} \hat{s}_{35}} \ell, \quad g_2' = \frac{4}{\hat{s}_{35} \hat{s}_{15}} \ell, \quad g_3' = \frac{4}{\hat{s}_{15} \hat{s}_{25}} \ell, \quad (4.6)
\end{equation}

with $\ell \equiv \epsilon_{\mu \nu \alpha \beta} k_1^\mu k_2^\nu k_3^\alpha k_4^\beta$, and $|M(+-)|^2$ is obtained from (4.5) by the replacement of $R_q$ by $L_q$.

For the case of $q\bar{q} \to q'\bar{q}'\gamma$ subprocess,

\begin{align*}
|M(+-)|^2 &= - \frac{s_{13}^2}{s_{45}} \left\{ \frac{e_{q'}}{s_{12}} (\tilde{A} g_1 + \tilde{B} g_2 + \tilde{C} f_1) \\
&\quad + \frac{e_q}{s_{34}} [\tilde{A} (f_2 + g_3) + \tilde{B} (f_3 + g_2) + \tilde{C} (g_1 + g_2)] \right\} |V_{qq'}|^2
\end{align*}
\[ -\frac{s_{24}^2}{s_{45}} \left\{ \frac{e_{q'}}{s_{14} - M_W^2} \left[ \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} \right] + \frac{e_q}{s_{14} - M_W^2} \left[ \frac{e_{q'}/s_{34}}{s_{12}} g_1 + \frac{e_q}{s_{34}} (f_2 + g_3) \right] \right\} \{V_{qq'}\}^2, \quad (4.7) \]

where

\[ \tilde{A} = \frac{s_{25}}{s_{13} (s_{23} - M_W^2)} + \frac{(1 + \frac{2}{s_{14}} e_{q'})}{s_{14} - M_W^2} + \frac{s_{15} + s_{45}}{(s_{14} - M_W^2) (s_{23} - M_W^2)} \frac{s_{25}}{s_{13}}, \]

\[ \tilde{B} = \frac{s_{25}}{s_{13} (s_{23} - M_W^2)} - \frac{s_{15} + s_{45}}{(s_{14} - M_W^2) (s_{23} - M_W^2)} \frac{s_{25}}{s_{13}}, \]

\[ \tilde{C} = \frac{e_{q'}}{s_{14} - M_W^2}. \quad (4.8) \]

For the case of \( qq' \rightarrow qq'\gamma \) subprocess,

\[ |M(-+)|^2 = \frac{s_{13}^2}{s_{45}} \left\{ \frac{s_{25}}{s_{13}} \left( \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} + \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} + \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} \right) \right\} \{V_{qq'}\}^2 \]

\[ + \frac{M_W \Gamma_W}{D_{12}^w D_{34}^w} \left( \frac{s_{35} + s_{45}}{s_{13}} \frac{s_{25}}{s_{13}} \frac{f_2}{s_{35} f_1} \right) \frac{e_{q'}}{s_{14}} \{V_{qq'}\}^2 \]

\[ - \frac{s_{24}^2}{s_{45}} \left( \frac{2 g_2 e_q}{s_{23}} + \frac{g_1 e_q}{s_{14}} \right) e_{q'} + \frac{s_{24}^2}{s_{45}} \left( \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} + \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} + \frac{e_{q'} f_1 + e_q (g_1 + g_2)}{s_{12}} \right) \]

\[ \quad + \frac{s_{12}^2}{s_{34}^2} \left( \frac{f_1 e_q}{s_{14}} + g_1 e_q \frac{s_{35}}{s_{14}} - g_2 e_q \frac{s_{35}}{s_{23}} \right) \]

\[ - \frac{M_W \Gamma_W}{D_{12}^w D_{34}^w} \left( \frac{f_1 e_q^2}{s_{14}} + \frac{f_2 e_q^2}{s_{14}} + \frac{f_3 e_q^2}{s_{14}} \right) \]

\[ + \frac{M_W^2 \Gamma_W^2}{D_{12}^w D_{34}^w} \left( g_1 e_q + g_2 e_q - f_1 e_q' \frac{s_{35}}{s_{14}} \right) \{V_{qq'}\}^2 \]

\[ + N(-+), \quad (4.9) \]

where \( N(-+) \) involves the totally antisymmetric tensor \( \epsilon_{\mu \nu \rho \sigma} \) and reads

\[ \frac{s_{13}^2}{s_{45}} \left\{ \left( \frac{g_2 - s_{25}}{s_{13}} \frac{e_{q'} s_{35}}{s_{23}} + \frac{1}{s_{14}} \left( \frac{g_1 e_q}{s_{14}} + \frac{g_1 e_q}{s_{14}} \right) \right) \frac{1}{D_{12}^w} + \left( \frac{s_{25}}{s_{14}} + e_{q'} \right) \frac{g_2 e_q}{s_{23}} - \frac{g_1 e_q}{s_{14}} \right\} \]

\[ + \frac{s_{24}^2}{s_{45}} \left\{ \left( \frac{g_1 e_q}{s_{14}} + \frac{g_1 e_q}{s_{14}} \right) \frac{1}{D_{12}^w} + \left( \frac{s_{25}}{s_{14}} + e_{q'} \right) \frac{g_2 e_q}{s_{23}} - \frac{g_1 e_q}{s_{14}} \right\} \]

\[ 13 \]
\[-(g'_2 + g'_3)\frac{e_q e'_q}{s_{14}} \frac{1}{D_{34}} \] \[M_W \Gamma_W + \frac{M_W^2}{D_{12}^2 D_{34}^2} \left( g'_3 (\hat{s}_{35} + \hat{s}_{45}) \frac{\hat{s}_{25}}{\hat{s}_{13}} - g'_2 \hat{s}_{35} \right) \frac{e_q}{\hat{s}_{23}} \]

\[+ \left( (\hat{s}_{35} + \hat{s}_{45}) \frac{\hat{s}_{25}}{\hat{s}_{13}} - \hat{s}_{35} \right) \frac{g'_1 e'_q}{\hat{s}_{14}} \right] \right\} |V_{qq}|^2 \]

\[\frac{\hat{s}_{24}^2}{\hat{s}_{45}} \left\{ M_W \Gamma_W \left[ \frac{e_q}{D_{12}} \left( g'_2 e_q \frac{\hat{s}_{23}}{\hat{s}_{14}} + g'_1 e'_q \frac{\hat{s}_{14}}{\hat{s}_{14}} \right) + \frac{e'_q}{D_{34}^2} \left( g'_3 e_q \frac{\hat{s}_{23}}{\hat{s}_{14}} + g'_2 e'_q \frac{\hat{s}_{14}}{\hat{s}_{14}} + g'_3 e'_q \frac{\hat{s}_{23}}{\hat{s}_{14}} \right) \right] + \left( \frac{r_{12}^w}{D_{34}^w} + \frac{r_{34}^w}{D_{12}^w} \left( g'_2 e_q \frac{\hat{s}_{23}}{\hat{s}_{14}} + g'_1 e'_q \frac{\hat{s}_{14}}{\hat{s}_{14}} \right) \right) \right\} |V_{qq}|^2, \]

(4.10)

with \( D_{ij}^w = (\hat{s}_{ij} - M_W^2)^2 + M_W^2 \Gamma_W^2 \) and \( r_{ij}^w = (\hat{s}_{ij} - M_W^2)/D_{ij}^w \).

Finally, the numerator of the single asymmetry \( A_L \) for 2-jet plus photon production is given by

\[\sigma^+ - \sigma^- = \frac{32}{9} \frac{\alpha^2}{s_{12} x_W (1 - x_W)} \sum_{ij} dx_a dx_b \Delta f_i(x_a, Q^2) f_j(x_b, Q^2) \]

\[\times |\Delta M|^2 \frac{d^3k_1}{E_3} \frac{d^3k_4}{E_4} \frac{d^3k_5}{E_5} \delta(k_1 + k_2 - k_3 - k_4 - k_5), \]

(4.11)

where \( x_W \equiv \sin^2 \theta_W \), \( |\Delta M|^2 = \frac{1}{2} (|M(++)|^2 + |M(+-)|^2 - |M(+\pm)|^2 - |M(-\pm)|^2) \). The unpolarized cross sections for 2-jet plus photon production can be found in Ref.[3].

V. \( \ell^+ \ell^- \) pair plus 1-jet production

Since QCD interactions change color whereas weak couplings always conserve color, pv effects at tree level in general depend only on the polarized valence and sea quark distributions. It is thus desirable to have some processes in which gluons also contribute to the pv asymmetry \( A_L \). In this section we shall study one of such reactions, namely \( pp \to \ell^+ \ell^- + \text{jet} + X \).

At the parton level, there are two subprocesses contributing to the Drell-Yan reaction \( pp \to \ell^+ \ell^- + \text{jet} + X \): \( G + q(\bar{q}) \to \ell^+ \ell^- + q(\bar{q}) \) and \( q\bar{q} \to \ell^+ \ell^- + G \). In this reaction, \( A_L \) arises from the interference between the amplitudes with \( \gamma \) and \( Z^0 \) exchanges. The transition matrix elements absolute squared for the subprocess \( G_{\lambda_1}(k_1) + q_{\lambda_2}(k_2)[\bar{q}_{\lambda_2}(k_2)] \to \ell^+(k_3)\ell^-(k_4) + q(k_5)[\bar{q}(k_5)] \) are (only the interference terms being written down)

\[ |M(++)|^2 = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{\hat{s}_{34}}{\hat{s}_{12} \hat{s}_{15}} \left\{ \frac{R_q^2 (L_1^2 \hat{s}_{24}^2 + R_{12}^2 \hat{s}_{23}^2)}{4x_W (1 - x_W)^2 D_{34}^2} \right\} \]

\[^3\text{The pv effect in the reaction } pp \to \ell^+ \ell^- + X \text{ has been discussed in Ref.[22].}\]
\[ M(+-) = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{s_{34}}{s_{12}s_{15}} \left\{ \frac{L^2_q(L^2_q s_{45}^2 + R^2 q s_{35}^2)}{4x_W^2 (1-x_W)^2 D_{34}^2} \right\}, \] (5.1)

\[ |M(+-)|^2 = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{s_{34}}{s_{12}s_{15}} \left\{ \frac{L^2_q(L^2_q s_{45}^2 + R^2 q s_{35}^2)}{4x_W^2 (1-x_W)^2 D_{34}^2} \right\} + 4 \frac{e_q^2}{s_{34}^2} + 2L_q r_{34}^z e_q \frac{(L_q s_{45}^2 + R_q s_{35}^2)}{x_W (1-x_W)s_{34}}, \] (5.2)

where

\[ R_q = 2 \sin^2 \theta_W, \quad L_q = -1 + 2 \sin^2 \theta_W. \] (5.3)

Using the hermitian conjugate of \( M(++) \) and \( M(+-) \), we obtain

\[ |M(--)|^2 = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{s_{34}}{s_{12}s_{15}} \left\{ \frac{R^2_q(L^2_q s_{45}^2 + R^2 q s_{35}^2)}{4x_W^2 (1-x_W)^2 D_{34}^2} \right\} + 4 \frac{e_q^2}{s_{34}^2} + 2L_q r_{34}^z e_q \frac{(L_q s_{45}^2 + R_q s_{35}^2)}{x_W (1-x_W)s_{34}}. \] (5.4)

and

\[ |M(+-)|^2 = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{s_{34}}{s_{12}s_{15}} \left\{ \frac{L^2_q[R_q^2(s_{14}^2 + s_{24}^2) + L_q^2(s_{13}^2 + s_{23}^2)]}{4x_W^2 (1-x_W)^2 D_{34}^2} \right\} + 4 \frac{e_q^2}{s_{34}^2} + 2L_q r_{34}^z e_q \frac{(L_q s_{14}^2 + s_{24}^2) + R_q s_{13}^2 + s_{23}^2]}{x_W (1-x_W)s_{34}}, \] (5.5)

The transition matrix elements absolute squared for the subprocess of

\[ q_{\lambda_1}(k_1) + \bar{q}_{\lambda_2}(k_2) \rightarrow \ell^+(k_3)\ell^-(k_4) + G(k_5) \]

read

\[ |M(+-)|^2 = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{s_{34}}{s_{15}s_{25}} \left\{ \frac{L^2_q[R_q^2(s_{14}^2 + s_{24}^2) + L_q^2(s_{13}^2 + s_{23}^2)]}{4x_W^2 (1-x_W)^2 D_{34}^2} \right\} + 4 \frac{e_q^2}{s_{34}^2} + 2L_q r_{34}^z e_q \frac{(L_q s_{14}^2 + s_{24}^2) + R_q s_{13}^2 + s_{23}^2]}{x_W (1-x_W)s_{34}}, \] (5.6)

and

\[ |M(+-)|^2 = 2(4\pi \alpha_s)(4\pi \alpha)^2 \frac{s_{34}}{s_{15}s_{25}} \left\{ \frac{R^2_q[R_q^2(s_{14}^2 + s_{24}^2) + L_q^2(s_{13}^2 + s_{23}^2)]}{4x_W^2 (1-x_W)^2 D_{34}^2} \right\} + 4 \frac{e_q^2}{s_{34}^2} + 2L_q r_{34}^z e_q \frac{(L_q s_{14}^2 + s_{24}^2) + R_q s_{13}^2 + s_{23}^2]}{x_W (1-x_W)s_{34}}. \] (5.7)
VI. Results and discussion

The results of our calculations for parity-violating asymmetries $A_L$ and $P_L$ are presented in a series of figures. Shown in Figs.1 and 2 are the $\sqrt{s}$ dependence of the helicity asymmetric cross section $\Delta\sigma_L = \sigma^+ - \sigma^-$ and $A_L$, respectively, at RHIC energies for a single jet production at $90^\circ$ (i.e. $y = 0$) in the c.m. with a jet momentum cutoff at 5 GeV for three different cases of polarized parton distributions (see Sec.II). Note that all parton spin densities constrained by the EMC data are first parametrized at $Q^2_0 = 10$ GeV$^2$ [Eqs.(2.11-2.13)] and then their $Q^2$ evolutions are governed by the Altarelli-Parisi equations. Since at tree level gluons in general do not contribute to $A_L$ and $P_L$, pv asymmetries are only sensitive to polarized valence and sea quark distribution functions. From Figs. 1 and 2 we see that $A_L$ is of order $10^{-5}$ at RHIC energies and is dominated by case (i) with large sea polarization. This is what expected since as far as $A_L$ is concerned, the three different parametrizations of parton spin densities are different only in their sea polarization. The behavior of $A_L$ in the 2-jet production at rapidities $y_1 = y_2 = 0$ with the jet momentum being cut off at 5 GeV has the same pattern as the previous 1-jet case (see Figs.1-4). This is attributed to that the underlying parton-parton scatterings for 1-jet and 2-jet productions are the same.

We have calculated in Figs.5 and 6 the longitudinal polarization $P_L$ of $\Lambda$ produced at $y = 0$ in unpolarized $pp$ collisions at energies $\sqrt{s} = 200$ GeV and 500 GeV as a function of $x_T = 2p_T/\sqrt{s}$. The dependence of $P_L$ on $x_T$, which is quite similar at the two energies shown in Fig.6, is one of the testable predictions for the polarized $\Lambda$ fragmentation function given in Eq.(2.17). The increase of $P_L$ with $x_T$ basically can be understood from the $z$ dependence of $\Delta D_\Lambda^q(z)/D_\Lambda^q(z) = z^\gamma$ with $\gamma > 0$; especially at large $x_T$ where the unpolarized cross section is dominated, as noted in passing, by quark-quark scattering. In order to have a numerical estimate for $P_L$, we have followed Ref.[11] to choose $\gamma = 10$. The resulting $P_L$ is of order $10^{-2}$ at moderate $x_T$. In practice, it will be easier to measure the $\Lambda$ polarization at small $x_T$ where the signal of $\Delta\sigma_\Lambda$ is large. As stressed in Sec.II, very little is known about the fragmentation functions $D_\Lambda^q(z)$ and $\Delta D_\Lambda^q(z)$. Presumably, the experiment of searching for helicity asymmetric cross section in semi-inclusive production of $\Lambda$ in $e^+e^-$ annihilation

\footnote{For the valence $d$ quark distribution function, we use $x_0 = 0.50$ in realistic calculation [see Eq.(2.9)].}
allows one to measure the polarized fragmentation function $\Delta D_q^A(z)$ [21].

We have also extended our discussions to the hadronic production of 2-jet plus photon final states. We plot in Figs.7 and 8 the differential cross section asymmetry and $A_L$, respectively, as a function of the invariant mass $\sqrt{s_{34}}$ of two jets. Since experimentally it is difficult to distinguish between quark and antiquark jets, we have symmetrized these two jets. The predicted $A_L$ increases from $10^{-4}$ to $10^{-3}$ as $\sqrt{s}$ varies from 100 GeV to 500 GeV, and it is insensitive to the polarized structure functions chosen.

Our investigation of the Drell-Yan type reaction $pp \rightarrow \ell^+\ell^- + \text{jet} + X$ is originally motivated by looking for the processes in which gluons also contribute to pv asymmetry so that a measurement of $A_L$ in such reactions would provide useful information on the gluon polarization. The results are presented in Figs.9 and 10. It is evident from Fig.9 that when the invariant mass of dilepton is around the mass of the $Z^0$ resonance, a bump is shown up in both unpolarized and helicity asymmetric cross sections, as it should be. Unfortunately, the resulting $A_L$ is insensitive to the choice of parton spin densities. This may be ascribed to the fact that contributions to the helicity asymmetric cross section come mainly from the region of moderate value of $x$ where spin dependent gluon and sea distributions are negligible. (Also note that for the subprocess $G + q \rightarrow \ell^+\ell^- + q$, $|\Delta M|^2$ arising from polarized gluons has a sign opposite to that due to polarized quarks [see Eqs.(5.1)-(5.5)].) We conclude that, contrary to the double helicity helicity asymmetry $A_{LL}$, it is unlikely that a measurement of the single helicity asymmetry $A_L$ will bring any new insight into the gluon polarization.

Finally, we note that a large pv asymmetry $A_L$ of order 10% is expected to be seen in $W^\pm$ and $Z^0$ productions with a large $p_T$ at RHIC energies (see Ref.[22] and the first paper in Ref.[1] for details). On the other hand, the helicity asymmetric cross section is estimated to be of order $1 pb$ for $W^+$ (and even smaller for $W^-$ and $Z^0$) in $pp$ collisions at energy, say $\sqrt{s} = 500$ GeV, to be compared with $\Delta \sigma \sim 10^{-1} nb$ for 1-jet or 2-jet production (see Figs.1 and 3). Therefore, it is worth pursuing all possible parity-violating effects in high energy hadron-hadron collisions at the planned hadronic colliders RHIC, SSC and LHC.
Acknowledgments

One of us (H.Y.C.) wishes to thank Prof. C. N. Yang and the Institute for Theoretical Physics at Stony Brook for their hospitality during his stay there for sabbatical leave. C.F.W. would like to thank Profs. Gonsalves and C. Y. Cheung for many helpful discussions. This work was supported in part by the National Science Council of the Republic of China under Contract Nos. NSC82-0208-M001-001Y and NSC82-0112-C001-019.

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FIGURE CAPTIONS

1. The signal of helicity asymmetric cross section for 1-jet production at \( y = 0 \) in \( pp \) collisions as a function of \( \sqrt{s} \) for three different polarized parton distributions as described in the text.

2. Asymmetry \( A_L \) for 1-jet production at \( y = 0 \) in \( pp \) collisions as a function of \( \sqrt{s} \) for three different polarized parton distributions as described in the text.

3. Same as Fig.1 except for 2-jet production at rapidities \( y_1 = y_2 = 0 \).

4. Same as Fig.2 except for 2-jet production at rapidities \( y_1 = y_2 = 0 \).

5. The signal of helicity asymmetric cross section for \( \Lambda \) production at \( y = 0 \) in unpolarized \( pp \) collisions as a function of \( x_T \) at energies \( \sqrt{s} = 200 \) and 500 GeV.

6. Longitudinal polarization \( P_L \) of \( \Lambda \) produced at \( y = 0 \) in unpolarized \( pp \) collisions as a function of \( x_T \) at energies \( \sqrt{s} = 200 \) and 500 GeV.

7. The signal of helicity asymmetric cross section for 2-jet plus photon production in \( pp \) collisions as a function of the invariant mass \( \sqrt{s_{34}} \) of dijet for three different polarized parton distributions.

8. Asymmetry \( A_L \) for 2-jet plus photon production in \( pp \) collisions as a function of the invariant mass \( \sqrt{s_{34}} \) of dijet for three different polarized parton distributions.

9. The signal of helicity asymmetric cross section for \( \ell^+\ell^- \) plus 1-jet production in \( pp \) collisions as a function of the invariant mass \( \sqrt{s_{34}} \) of dilepton for three different polarized parton distributions.

10. Asymmetry \( A_L \) for \( \ell^+\ell^- \) plus 1-jet production in \( pp \) collisions as a function of the invariant mass \( \sqrt{s_{34}} \) of dilepton for three different polarized parton distributions.