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Structure functions in the polarized Drell-Yan processes with spin-1/2 and spin-1 hadrons: II. parton model

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Abstract

We analyze the polarized Drell-Yan processes with spin-1/2 and spin-1 hadrons in a parton model. Quark and antiquark correlation functions are expressed in terms of possible combinations of Lorentz vectors and pseudovectors with the constrains of Hermiticity, parity conservation, and time-reversal invariance. Then, we find tensor polarized distributions for a spin-1 hadron. The naive parton model predicts that there exist 19 structure functions. However, there are only four or five non-vanishing structure functions, depending on whether the cross section is integrated over the virtual-photon transverse momentum $\vec{Q}_T$ or the limit $Q_T \to 0$ is taken. One of the finite structure functions is related to the tensor polarized distribution b_1, and it does not exist in the proton-proton reactions. The vanishing structure functions should be associated with higher-twist physics. The tensor distributions can be measured by the quadrupole polarization measurements. The Drell-Yan process has an advantage over the lepton reaction in the sense that the antiquark tensor polarization could be extracted rather easily.

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I. INTRODUCTION

Spin structure of the proton has been studied extensively. Although it is not still obvious how to interpret the proton spin in terms of partons, it is important to test our knowledge of spin physics by using other observables. The spin structure of spin-1 hadrons is a good example. In particular, tensor structure exists for the spin-1 hadrons as a new ingredient. However, it is not clear at this stage how to describe it in a parton model although there are some initial studies [1–4]. In this sense, the studies of spin-1 hadrons should be a challenging experience.

On the other hand, the Relativistic Heavy Ion Collider (RHIC) will be completed soon and new proton-proton (pp) polarization experiments will be done. As a next-generation project, the spin-1 deuteron acceleration may be possible [5]. With this future project and others [6] in mind, we have completed general formalism for the polarized Drell-Yan processes with spin-1/2 and spin-1 hadrons in Ref. [7]. It was revealed that 108 structure functions exist in the reactions and that the number becomes 22 after integrating the cross section over the virtual-photon transverse momentum $\vec{Q}_T$ or after taking the limit $Q_T \to 0$.

The purpose of this paper is to clarify how these structure functions are related to the parton distributions in the colliding hadrons. In a parton model, it is known that the unpolarized distribution f_1 (or denoted as q), the longitudinally polarized one g_1 (Δq), and the transversity one h_1 ($\Delta_T q$) are studied in the pp Drell-Yan processes. In addition to these, the tensor distribution b_1 (δq) should contribute to our Drell-Yan processes. In the following sections, we obtain the cross section in the parton model and then discuss how b_1 is related to one of the structure functions in Ref. [7]. We also study how spin asymmetries can be expressed by the parton distributions.

First, the hadron tensor is obtained in the parton model by finding possible Lorentz vector and pseudovector combinations in section II. Then, the cross section is calculated by using the hadron tensor. Second, we derive the cross-section expression by integrating over $\vec{Q}_T$ or by taking the limit $Q_T \to 0$ in section III. Third, using the $\vec{Q}_T$-integrated cross

section, we express the spin asymmetries in terms of the parton distributions in IV. Finally, our studies are summarized in section V.

II. PARTON-MODEL DESCRIPTION OF THE DRELL-YAN PROCESS

A. Drell-Yan process

Parton-model analyses of the Drell-Yan process with spin-1/2 hadrons were reported by Ralston-Soper [8], Donohue-Gottlieb [9], and Tangerman-Mulders (TM) [10]. In order to clarify the difference from the spin-1/2 case, we discuss our Drell-Yan process along the TM formalism. We consider the following process

$$A\,(spin\,1/2) + B\,(spin\,1) \to \ell^+\ell^- + X \ , \tag{2.1}$$

where A and B are spin-1/2 and spin-1 hadrons, respectively. They could be any hadrons; however, realistic ones are the proton and the deuteron experimentally. The description of this paper can be applied to any hadrons with spin-1/2 and spin-1.

Although the formalism of this subsection is discussed in Ref. [10], we explain it in order to present the definitions of kinematical variables and functions for understanding the subsequent sections. The cross section is written in terms of the lepton tensor $L_{\mu\nu}$ and the hadron tensor $W^{\mu\nu}$ as [7]

$$\frac{d\sigma}{d^4Q\,d\Omega} = \frac{\alpha^2}{2\,s\,Q^4}\,L_{\mu\nu}\,W^{\mu\nu} \ , \tag{2.2}$$

where $\alpha = e^2/(4\pi)$ is the fine structure constant, s is the center-of-mass energy squared $s = (P_A+P_B)^2$, Q is the total dilepton momentum, and Ω is the solid angle of the momentum $\vec{k}_{\ell^+} - \vec{k}_{\ell^-}$. The hadron and lepton masses are neglected in comparison with s and Q^2: $M_A^2, M_B^2 \ll s$ and $m_\ell^2 \ll Q^2$. In order to compare with the TM results, we use the same notations as many as we can. Three vectors $\hat{x}^\mu$, $\hat{y}^\mu$, and $\hat{z}^\mu$ are defined as

$$\hat{x}^\mu = -\frac{X^\mu}{\sqrt{-X^2}} \ , \quad \hat{y}^\mu = -\frac{Y^\mu}{\sqrt{-Y^2}} \ , \quad \hat{z}^\mu = +\frac{Z^\mu}{\sqrt{-Z^2}} \ , \tag{2.3}$$

where X^μ, Y^μ, and Z^μ are given in Ref. [7]. The Z^μ axis corresponds to the Collins-Soper choice [11] and these axes are spacelike in the dilepton rest frame. In the same way, $\hat{Q}$ is defined by

$$\hat{Q}^\mu = \frac{Q^\mu}{\sqrt{Q^2}} \ . \tag{2.4}$$

It is convenient to introduce the lightcone notation $a = [a^-, a^+, \vec{a}_T]$, where $a^\pm = (a^0 \pm a^3)/\sqrt{2}$. The transverse vector $a_T = [0, 0, \vec{a}_T]$ is projected out by

$$g_T^{\mu\nu} = g^{\mu\nu} - n_+^\mu n_-^\nu - n_+^\nu n_-^\mu \ , \tag{2.5}$$

with $n_+ = [0, \kappa, \vec{0}_T]$ and $n_- = [1/\kappa, 0, \vec{0}_T]$, where $\kappa = \sqrt{x_A/x_B}$ in the center-of-momentum (c.m.) frame. The transverse vector is orthogonal to the hadron momenta: $a_T \cdot P_A = a_T \cdot P_B = 0$. Furthermore, we define $g_\perp^{\mu\nu}$ as

$$g_\perp^{\mu\nu} = g^{\mu\nu} - \hat{Q}^\mu \hat{Q}^\nu + \hat{z}^\mu \hat{z}^\nu \ , \tag{2.6}$$

which projects out the perpendicular vector $a_\perp^\mu \equiv g_\perp^{\mu\nu} a_{T\nu}$. It is orthogonal to the vectors $\hat{Q}$ and $\hat{z}$: $a_\perp \cdot \hat{Q} = a_\perp \cdot \hat{z} = 0$. Because the transverse projection is equal to the perpendicular one if the $1/Q$ term can be neglected [$g_T^{\mu\nu} = g_\perp^{\mu\nu} + O(1/Q)$], we use the approximation $a_\perp \approx a_T$ in the following calculations. Nevertheless, the perpendicular vectors are often used because they are convenient due to the orthogonal relation. The lepton tensor is expressed in terms of these quantities as [10]

$$\begin{aligned} L^{\mu\nu} &= 2\, k_{\ell+}^\mu\, k_{\ell-}^\nu + 2\, k_{\ell+}^\nu\, k_{\ell-}^\mu - Q^2\, g^{\mu\nu} \\ &= -\frac{Q^2}{2} \Bigg[(1 + \cos^2\theta)\, g_\perp^{\mu\nu} - 2\sin^2\theta\, \hat{z}^\mu\, \hat{z}^\nu + 2\sin^2\theta\, \cos 2\phi\, (\hat{x}^\mu \hat{x}^\nu + \frac{1}{2}\, g_\perp^{\mu\nu}) \\ &\qquad + \sin^2\theta\, \sin 2\phi\, \hat{x}^{\{\mu} \hat{y}^{\nu\}} + \sin 2\theta\, \cos\phi\, \hat{z}^{\{\mu} \hat{x}^{\nu\}} + \sin 2\theta\, \sin\phi\, \hat{z}^{\{\mu} \hat{y}^{\nu\}} \Bigg] \ , \end{aligned} \tag{2.7}$$

where θ and ϕ are polar and azimuthal angles of the vector $\vec{k}_{\ell+} - \vec{k}_{\ell-}$, and the notation $A^{\{\mu} B^{\nu\}}$ is defined by

$$A^{\{\mu} B^{\nu\}} \equiv A^\mu B^\nu + A^\nu B^\mu \ . \tag{2.8}$$

The hadron tensor is given by

$$W^{\mu\nu} = \int \frac{d^4\xi}{(2\pi)^4}\, e^{iQ\cdot\xi} \, < P_A S_A; P_B S_B \,|\, J^\mu(0)\, J^\nu(\xi) \,|\, P_A S_A; P_B S_B > \,, \tag{2.9}$$

where J^μ is the electromagnetic current, and the hadron momenta and spins are denoted as P_A, P_B, S_A, and S_B. The analysis of the hadron tensor is more complicated than the one in deep inelastic lepton-hadron scattering because it contains the currents with two-hadron states. The leading lightcone singularity originates from the process that a quark emits a virtual photon, which then splits into $\ell^+\ell^-$. However, it does not contribute to the cross section significantly because the quark should be far off-shell [12]. The dominant contribution comes from quark-antiquark annihilation processes. In the following, we discuss the hadron tensor and the cross section due to the annihilation process: q(in A)+$\bar{q}$(in B)$\rightarrow \ell^+ + \ell^-$ in Fig. 1. Of course, the opposite process $\bar{q}$(in A)+q(in B)$\rightarrow \ell^+ + \ell^-$ should be taken into account in order to compare with the experimental cross section. Its contribution is included in discussing the spin asymmetries in section IV. The first process contribution to the hadron tensor in Eq. (2.9) is

$$W^{\mu\nu} = \frac{1}{3}\sum_{a,b} \delta_{b\bar{a}}\, e_a^2 \int d^4k_a\, d^4k_b\, \delta^4(k_a + k_b - Q)\, Tr[\Phi_{a/A}(P_A S_A; k_a)\gamma^\mu \bar{\Phi}_{b/B}(P_B S_B; k_b)\gamma^\nu] \,, \tag{2.10}$$

where k_a and $k_b = k_{\bar{a}}$ are the quark and antiquark momenta, the color average is taken by the factor $1/3 = 3\cdot(1/3)^2$, and e_a is the charge of a quark with the flavor a. The correlation functions $\Phi_{a/A}$ and $\bar{\Phi}_{\bar{a}/B}$ are defined by [8]

$$\begin{aligned}
\Phi_{a/A}(P_A S_A; k_a)_{ij} &\equiv \int \frac{d^4\xi}{(2\pi)^4}\, e^{ik_a\cdot\xi} < P_A S_A \,|\, \bar{\psi}_j^{(a)}(0)\, \psi_i^{(a)}(\xi) \,|\, P_A S_A > \,, \\
\bar{\Phi}_{\bar{a}/B}(P_B S_B; k_{\bar{a}})_{ij} &\equiv \int \frac{d^4\xi}{(2\pi)^4}\, e^{ik_{\bar{a}}\cdot\xi} < P_B S_B \,|\, \psi_i^{(a)}(0)\, \bar{\psi}_j^{(a)}(\xi) \,|\, P_B S_B > \,.
\end{aligned} \tag{2.11}$$

Link operators should be introduced in these matrix elements so as to become gauge invariant [10] although they are not explicitly written in the above equations. It is also known that they become identity in the lightcone gauge. In any case, such link operators do not alter the following discussions of this paper in a naive parton model. Figure 1 suggests that the hadron

tensor could be written by a product of quark-hadron amplitudes. However, the matrix indices in the trace of Eq. (2.10) look like $[\bar{\psi}_j^{(a)}(0)\,\psi_i^{(a)}(\xi)]_A\,[\bar{\psi}_\ell^{(a)}(\xi)\,\psi_k^{(a)}(0)]_B\,(\gamma^\mu)_{jk}\,(\gamma^\nu)_{\ell i}$, which is not in a separable form. A Fierz transformation [13] is used so that the index summations are taken separately in the hadrons A and B. Using the relation

$$4(\gamma^\mu)_{jk}(\gamma^\nu)_{li} = [\mathbf{1}_{ji}\mathbf{1}_{lk} + (i\gamma_5)_{ji}(i\gamma_5)_{lk} - (\gamma^\alpha)_{ji}(\gamma_\alpha)_{lk} - (\gamma^\alpha\gamma_5)_{ji}(\gamma_\alpha\gamma_5)_{lk} + \frac{1}{2}(i\sigma_{\alpha\beta}\gamma_5)_{ji}(i\sigma^{\alpha\beta}\gamma_5)_{lk}]\,g^{\mu\nu}$$
$$+ (\gamma^{\{\mu})_{ji}(\gamma^{\nu\}})_{lk} + (\gamma^{\{\mu}\gamma_5)_{ji}(\gamma^{\nu\}}\gamma_5)_{lk} + (i\sigma^{\alpha\{\mu}\gamma_5)_{ji}(i\sigma^{\nu\}}{}_\alpha\gamma_5)_{lk}\ , \tag{2.12}$$

we factorize the hadron tensor as

$$W^{\mu\nu} = \frac{1}{3}\sum_{a,b}\delta_{b\bar{a}}\,e_a^2\int d^2\vec{k}_{aT}\,d^2\vec{k}_{bT}\,\delta^2(\vec{k}_{aT}+\vec{k}_{bT}-\vec{Q}_T)\left[\left\{ -\Phi_{a/A}[\gamma^\alpha]\,\bar{\Phi}_{b/B}[\gamma_\alpha]\right.\right.$$
$$\left.\left. - \Phi_{a/A}[\gamma^\alpha\gamma_5]\,\bar{\Phi}_{b/B}[\gamma_\alpha\gamma_5] + \frac{1}{2}\Phi_{a/A}[i\sigma_{\alpha\beta}\gamma_5]\,\bar{\Phi}_{b/B}[i\sigma^{\alpha\beta}\gamma_5]\right\} g^{\mu\nu} + \Phi_{a/A}[\gamma^{\{\mu}]\,\bar{\Phi}_{b/B}[\gamma^{\nu\}}]\right.$$
$$\left. + \Phi_{a/A}[\gamma^{\{\mu}\gamma_5]\,\bar{\Phi}_{b/B}[\gamma^{\nu\}}\gamma_5] + \Phi_{a/A}[i\sigma^{\alpha\{\mu}\gamma_5]\,\bar{\Phi}_{b/B}[i\sigma^{\nu\}}{}_\alpha\gamma_5]\right] + O(1/Q)\ . \tag{2.13}$$

The definition of the brace is given in Eq. (2.8). For example, the last term in the bracket is explicitly written as $\Phi_{a/A}[i\sigma^{\alpha\{\mu}\gamma_5]\,\bar{\Phi}_{b/B}[i\sigma^{\nu\}}{}_\alpha\gamma_5] = \Phi_{a/A}[i\sigma^{\alpha\mu}\gamma_5]\,\bar{\Phi}_{b/B}[i\sigma^\nu{}_\alpha\gamma_5] + \Phi_{a/A}[i\sigma^{\alpha\nu}\gamma_5]\,\bar{\Phi}_{b/B}[i\sigma^\mu{}_\alpha\gamma_5]$. The functions $\Phi_{a/A}[\Gamma]$ and $\bar{\Phi}_{b/B}[\Gamma]$ are defined by

$$\Phi_{a/A}[\Gamma](x,\vec{k}_T) \equiv \frac{1}{2}\int dk_a^-\,Tr[\Gamma\,\Phi_{a/A}]\ , \tag{2.14}$$

$$\bar{\Phi}_{b/B}[\Gamma](x,\vec{k}_T) \equiv \frac{1}{2}\int dk_b^+\,Tr[\Gamma\,\bar{\Phi}_{b/B}]\ , \tag{2.15}$$

and we assume $k_b^+ \ll k_a^+$ and $k_a^- \ll k_b^-$ in obtaining the factorized expression. The functions $\Phi_{a/A}[\mathbf{1}]$ and $\Phi_{a/A}[i\gamma_5]$ are obtained as $\Phi_{a/A}[\mathbf{1}] \sim O(1/Q)$ and $\Phi_{a/A}[i\gamma_5] = 0$ according to the calculations in the next subsection, so that they are not explicitly written in Eq. (2.13). As it is obvious from Eq. (2.13), we do not address ourselves to the higher-twist terms. Because the Drell-Yan process of spin-1/2 and spin-1 hadrons is not investigated at all in any parton model, we first discuss the leading contributions in this paper. We found 108 (22) structure functions in general (after integration over $\vec{Q}_T$), and most of them are related to the higher-twist physics as it becomes obvious in the later sections of this paper. Although it is interesting to study the higher-twist terms [14], we leave this topic as our future project.

B. Correlation functions and parton distributions

The hadron tensor is expressed by the correlation functions in Eq. (2.10). We expand them in terms of the possible Lorentz vectors and pseudovectors. The correlation function $\Phi(PS;k)$ is a matrix with sixteen components, so that it can be written in terms of sixteen 4×4 matrices [13]:

$$\mathbf{1},\ \gamma_5,\ \gamma^\mu,\ \gamma^\mu\gamma_5,\ \sigma^{\mu\nu}\gamma_5\ . \tag{2.16}$$

Because the spin-1/2 hadron case was already discussed in Refs. [8–10], we investigate the correlation function for a spin-1 hadron. We discuss it in the frame where the spin-1 hadron is moving in the z direction. Because the z_{cm} direction is taken as the direction of the hadron A momentum in Ref. [7], it could mean that we assume the hadron A as if it were a spin-1 hadron in this subsection. Of course, appropriate spin-1/2 and spin-1 expressions are used for the hadron A and B, respectively, in calculating the hadron tensor and the cross section. The correlation function is expanded in terms of the matrices of Eq. (2.16) together with the possible Lorentz vectors and pseudovector: P^μ, k^μ, and S^μ. However, these combinations have to satisfy the usual conditions of Hermiticity, parity conservation, and time-reversal invariance. In finding the possible combinations, we should be careful that the rank-two spin terms are allowed for a spin-1 hadron. The reader may look at Ref. [7] for the detailed discussions on this point. Considering these conditions, we obtain the possible Lorentz scalar quantities. Then, the coefficient A_i is assigned for each term:

$$\begin{aligned}\Phi(PS;k) = {} & A_1\mathbf{1} + A_2\not{P} + A_3\not{k} + A_4\gamma_5\not{S} + A_5\gamma_5[\not{P},\not{S}] + A_6\gamma_5[\not{k},\not{S}] + A_7 k\cdot S\gamma_5\not{P} + A_8 k\cdot S\gamma_5\not{k} \\ & + A_9 k\cdot S\gamma_5[\not{P},\not{k}] + A_{10}(k\cdot S)^2\mathbf{1} + A_{11}(k\cdot S)^2\not{P} + A_{12}(k\cdot S)^2\not{k} + A_{13}\, k\cdot S\,\not{S}\ .\end{aligned} \tag{2.17}$$

For simplicity, the subscripts a and A in Φ, momenta, momentum fraction, spin, and helicity are not written in this subsection. The spin dependent factors $(k\cdot S)^2$ and $k\cdot S\,\not{S}$ do not exist in a spin-1/2 hadron, so that the terms A_{10}, A_{11}, A_{12}, and A_{13} are the additional ones to the spin-1/2 expression in Ref. [10]. It means that the interesting tensor structure of the spin-1 hadron is contained in these new terms.

With the general expression of Eq. (2.17), we can calculate $\Phi[\Gamma]$ in Eq. (2.13). Because the new terms contribute only to $\Phi[\gamma^\alpha]$, its calculation procedure is discussed in the following by using the lightcone representation. The γ matrices are given by [12]

$$\gamma^0 = \begin{pmatrix} 0 & \sigma^3 \\ \sigma^3 & 0 \end{pmatrix} \ , \quad \vec{\gamma}_{_T} = \begin{pmatrix} i\,\vec{\sigma}_{_T} & 0 \\ 0 & i\,\vec{\sigma}_{_T} \end{pmatrix} \ , \quad \gamma^3 = \begin{pmatrix} 0 & -\sigma^3 \\ +\sigma^3 & 0 \end{pmatrix} \ , \quad \gamma^\pm = \frac{1}{\sqrt{2}}\,(\gamma^0 \pm \gamma^3) \ . \tag{2.18}$$

It is necessary to calculate $\Phi[\gamma^+]$, $\Phi[\gamma^-]$, and $\Phi[\vec{\gamma}_{_T}]$ for obtaining the $\Phi[\gamma^\alpha]$ term in Eq. (2.13). Substituting Eq. (2.17) into Eq. (2.14), we have the expression for $\Phi[\gamma^+]$ as

$$\begin{aligned} \Phi[\gamma^+] &= \frac{1}{2}\int dk^- Tr[\gamma^+\Phi] \\ &= \frac{1}{2}\int dk^- Tr\left[\,\gamma^+\{A_2\not{P} + A_3\not{k} + A_{11}(k\cdot S)^2\not{P} + A_{12}(k\cdot S)^2\not{k} + A_{13}\,k\cdot S\,\not{S}\}\,\right] \ . \end{aligned} \tag{2.19}$$

The traces of the lightcone γ matrices are

$$\begin{aligned} &Tr(\gamma^+\gamma^+) = Tr(\gamma^-\gamma^-) = 0, \quad Tr(\gamma^+\gamma^-) = Tr(\gamma^-\gamma^+) = 4, \\ &Tr(\gamma^+\vec{\gamma}_{_T}) = Tr(\gamma^-\vec{\gamma}_{_T}) = 0, \quad Tr(\gamma^i_{_T}\gamma^j_{_T}) = -4\,\delta^{ij}_T, \ . \end{aligned} \tag{2.20}$$

Therefore, the trace $Tr(\gamma^+\not{P})$ becomes $Tr(\gamma^+\not{P}) = 4P^+$. Then, introducing the momentum fraction x and the helicity λ by $k^+ = xP^+$ and $S^+ = \lambda P^+/M$, we obtain

$$\Phi[\gamma^+] = \int d(2k\cdot P)\left[A_2 + xA_3 + (\vec{k}_T\cdot\vec{S}_T)^2\,(A_{11} + xA_{12}) - \vec{k}_T\cdot\vec{S}_T\,\frac{\lambda}{M}\,A_{13}\right] , \tag{2.21}$$

where M is the hadron mass. This equation is written as

$$\Phi[\gamma^+] = f_1(x,\vec{k}_T^{\,2}) + b_1(x,\vec{k}_T^{\,2})\left[\frac{4(\vec{k}_T\cdot\vec{S}_T)^2}{\vec{k}_T^{\,2}} - \frac{2}{3}\right] + c_1(x,\vec{k}_T^{\,2})\,\lambda\,\frac{\vec{k}_T\cdot\vec{S}_T}{M} \ , \tag{2.22}$$

with the parton distributions

$$\begin{aligned} f_1(x,\vec{k}_T^{\,2}) &= \int d(2k\cdot P)\,\left\{A_2 + \vec{k}_T^{\,2}A_{11}/6 + x\left(A_3 + \vec{k}_T^{\,2}A_{12}/6\right)\right\} \ , \\ b_1(x,\vec{k}_T^{\,2}) &= \int d(2k\cdot P)\,\vec{k}_T^{\,2}\,(A_{11} + xA_{12})/4 \ , \\ c_1(x,\vec{k}_T^{\,2}) &= -\int d(2k\cdot P)\,A_{13} \ . \end{aligned} \tag{2.23}$$

Because the integration variable $2k\cdot P$ is constrained by the relation $\vec{k}_T^{\,2}+k^2-2xk\cdot P+x^2M^2=0$, it is of the order of 1 despite $P^+\sim O(Q)$. In addition to the usual unpolarized parton distribution f_1, there appear new distributions b_1 and c_1 which do not exist for a spin-1/2 hadron. The correlation function $\Phi[\gamma^+]$ can be also expressed by the quark probability density $\mathcal{P}(x,\vec{k}_T)$ according to Ref. [10]:

$$\Phi[\gamma^+]=\mathcal{P}(x,\vec{k}_T)\ . \tag{2.24}$$

Integrating over the transverse momentum $\vec{k}_T$, we obtain

$$\mathcal{P}(x)=f_1(x)+b_1(x)\,\frac{2}{3}\,(2\,|\vec{S}_T|^2-\lambda^2)\ , \tag{2.25}$$

where the relation $\lambda^2+|\vec{S}_T|^2=1$ is used and the functions $\mathcal{P}(x)$, $f_1(x)$, and $b_1(x)$ are defined by

$$\begin{aligned}\mathcal{P}(x)&=\int d^2\vec{k}_T\,\mathcal{P}(x,\vec{k}_T)\ ,\\ f_1(x)&=\int d^2\vec{k}_T\,f_1(x,\vec{k}_T^{\,2})\ ,\\ b_1(x)&=\int d^2\vec{k}_T\,b_1(x,\vec{k}_T^{\,2})\ .\end{aligned} \tag{2.26}$$

It is obvious from the spin combination in Eq. (2.25) that b_1 is related to the tensor structure. Equation (2.25) means that the b_1 distribution is given by

$$b_1(x)=\frac{1}{2}\left[\mathcal{P}(x)_{\lambda=0}-\frac{\mathcal{P}(x)_{\lambda=+1}+\mathcal{P}(x)_{\lambda=-1}}{2}\right]\ . \tag{2.27}$$

Indeed, this definition of the tensor distribution agrees with the one in the lepton-deuteron studies [2]. On the other hand, c_1 should be related to the intermediate polarization according to Ref. [7]. This distribution has never been discussed as far as we are aware, so that it is simply named c_1. This is a brand-new one in this paper. The interesting point of this distribution is that it cannot be measured by the longitudinally and transversely polarized reactions. The optimum way of observing it is to polarize the spin-1 hadron with the angles 45° and 135° with respect to the hadron momentum direction. However, as we mention in section IV, it is difficult to attain the intermediate polarization in the collider experiments

if the deuteron is used as a spin-1 hadron because of its small magnetic moment. It is not unique to express c_1 in terms of $\mathcal{P}(x, \vec{k}_T)$. For example, it is written by the distributions in the intermediate polarizations I_1 and I_3 of Ref. [7] as

$$\mathcal{P}(x, \vec{k}_T)_{I_1} - \mathcal{P}(x, \vec{k}_T)_{I_3} = \frac{|\vec{k}_T|}{M} \sin\phi_k \, c_1(x, \vec{k}_T^{\,2}) \ , \tag{2.28}$$

where ϕ_k is the azimuthal angle of the vector $\vec{k}_T$. Because its contribution vanishes by integrating the correlation function over $\vec{Q}_T$ or by taking the limit $Q_T \to 0$, it could be related to higher-twist distributions. This point should be clarified by our future project. At this stage, we should content ourselves with the studies of leading contributions because there exists no parton-model analysis of the Drell-Yan process with a spin-1 hadron. The same calculations are done for the functions $\Phi[\gamma^-]$ and $\Phi[\gamma_T]$. However, these are proportional to $O(1/Q)$, so that they are ignored in the following discussions.

The other correlation functions are calculated in the same way; however, the results are the same as those in Ref. [10]:

$$\begin{aligned}
\Phi[\gamma^+\gamma_5] &= \mathcal{P}(x, \vec{k}_T) \, \lambda(x, \vec{k}_T) = g_{1L}(x, \vec{k}_T^{\,2}) \, \lambda + g_{1T}(x, \vec{k}_T^{\,2}) \, \frac{\vec{k}_T \cdot \vec{S}_T}{M} \ , \\
\Phi[i\sigma^{i+}\gamma_5] &= \mathcal{P}(x, \vec{k}_T) \, \vec{s}_T^{\,i}(x, \vec{k}_T) = h_{1T}(x, \vec{k}_T^{\,2}) \, \vec{S}_T^{\,i} + \left[h_{1L}^\perp(x, \vec{k}_T^{\,2}) \, \lambda + h_{1T}^\perp(x, \vec{k}_T^{\,2}) \, \frac{\vec{k}_T \cdot \vec{S}_T}{M} \right] \frac{\vec{k}_T^{\,i}}{M} \ .
\end{aligned} \tag{2.29}$$

Here, $\lambda(x, \vec{k}_T)$ and $\vec{s}_T^{\,i}(x, \vec{k}_T)$ are the quark helicity and transverse polarization densities. The longitudinally and transversely polarized distributions are given as

$$\begin{aligned}
g_{1L}(x, \vec{k}_T^{\,2}) &= \int d(2k\cdot P) \, (A_4/M) \ , & g_{1T}(x, \vec{k}_T^{\,2}) &= -\int d(2k\cdot P) \, M(A_7 + xA_8) \ , \\
h_{1T}(x, \vec{k}_T^{\,2}) &= -\int d(2k\cdot P) \, 2 \, (A_5 + xA_6) \ , & h_{1L}^\perp(x, \vec{k}_T^{\,2}) &= \int d(2k\cdot P) \, 2 \, A_6 \ , \\
h_{1T}^\perp(x, \vec{k}_T^{\,2}) &= \int d(2k\cdot P) \, 2 \, M^2 A_9 \ .
\end{aligned} \tag{2.30}$$

Because the explanations were given for the g_1 and h_1 distributions in Ref. [10], we do not repeat them in this paper. The A_{10} term contributes to $\Phi[\mathbf{1}]$ as an additional one to the spin-1/2 case; however, it is proportional to $O(1/Q)$. The terms with $\Phi[\mathbf{1}]$ and $\Phi[i\gamma_5]$

are excluded from our formalism because the relations $\Phi[\mathbf{1}] \sim O(1/Q)$ and $\Phi[i\gamma_5] = 0$ are obtained by the similar calculations.

In calculating the Drell-Yan cross section, the antiquark correlation function should be also expressed in terms of the antiquark distributions. It is obtained by using the charge-conjugation property [10,15]. The antiquark correlation function in the hadron A is related to the quark one in the antihadron $\bar{A}$ by the charge-conjugation matrix $C = i\gamma^2\gamma^0$:

$$\bar{\Phi}_{\bar{a}/A} = -\, C^{-1} \left(\Phi_{a/\bar{A}}\right)^T C \ , \tag{2.31}$$

where the superscript T indicates the transposed matrix. Because the quark distribution in the antihadron is equal to the antiquark distribution in the hadron, the antiquark correlation functions is expressed as

$$\begin{aligned}
\bar{\Phi}[\gamma^+] =& \bar{f}_1(x, \vec{k}_T^{\,2}) + \bar{b}_1(x, \vec{k}_T^{\,2}) \left[\frac{4(\vec{k}_T \cdot \vec{S}_T)^2}{\vec{k}_T^{\,2}} - \frac{2}{3}\right] + \bar{c}_1(x, \vec{k}_T^{\,2})\, \lambda\, \frac{\vec{k}_T \cdot \vec{S}_T}{M} \ , \\
\bar{\Phi}[\gamma^+\gamma_5] =& -\, \bar{g}_{1L}(x, \vec{k}_T^{\,2})\, \lambda - \bar{g}_{1T}(x, \vec{k}_T^{\,2})\, \frac{\vec{k}_T \cdot \vec{S}_T}{M} \ , \\
\bar{\Phi}[i\sigma^{i+}\gamma_5] =& \bar{h}_{1T}(x, \vec{k}_T^{\,2})\, \vec{S}_T^i + \left[\bar{h}_{1L}^{\perp}(x, \vec{k}_T^{\,2})\, \lambda + \bar{h}_{1T}^{\perp}(x, \vec{k}_T^{\,2})\, \frac{\vec{k}_T \cdot \vec{S}_T}{M}\right] \frac{\vec{k}_T^{\,i}}{M} \ .
\end{aligned} \tag{2.32}$$

Furthermore, the anticommutation relations for fermions indicate $\bar{\Phi}_{ij}(PS; k) = -\Phi_{ij}(PS; -k)$, so that the distributions satisfy the relation

$$\bar{f}(x, \vec{k}_T^{\,2}) = \begin{cases} -f(-x, \vec{k}_T^{\,2}) & f = f_1,\, b_1,\, g_{1T},\, h_{1T}, \text{ and } h_{1T}^{\perp} \\ +f(-x, \vec{k}_T^{\,2}) & f = c_1,\, g_{1L}, \text{ and } h_{1L}^{\perp} \ . \end{cases} \tag{2.33}$$

In this way, we have derived the expressions of the correlation functions in terms of the quark and antiquark distributions. As new distributions, the b_1, $\bar{b}_1$, c_1, and $\bar{c}_1$ distributions appear in the correlation functions. In particular, the b_1 distribution is a leading-twist one and it is associated with the tensor structure of the spin-1 hadron. On the other hand, the c_1 distribution is related to the intermediate polarization of Ref. [7] and it could be associated with higher-twist physics. The other longitudinally and transversely polarized distributions exist in the same way as those of a spin-1/2 hadron.

C. Structure functions and the cross section

Because the correlation functions in Eq. (2.13) are calculated, the hadron tensor can be expressed in terms of the parton distributions. Neglecting the higher-twist contributions, we write the hadron tensor as

$$W^{\mu\nu} = -\frac{1}{3}\sum_{a,b}\delta_{b\bar{a}}\,e_a^2\int d^2\vec{k}_{aT}\,d^2\vec{k}_{bT}\,\delta^2(\vec{k}_{aT}+\vec{k}_{bT}-\vec{Q}_T)\Big\{\big(\,\Phi_{a/A}[\gamma^+]\,\bar{\Phi}_{b/B}[\gamma^-] \\ +\Phi_{a/A}[\gamma^+\gamma_5]\,\bar{\Phi}_{b/B}[\gamma^-\gamma_5]\,\big)\,g_T^{\mu\nu}+\Phi_{a/A}[i\sigma^{i+}\gamma_5]\,\bar{\Phi}_{b/B}[i\sigma^{j-}\gamma_5]\,\Big(g_{Ti}{}^{\{\mu}g_T{}^{\nu\}}{}_j-g_{Tij}\,g_T^{\mu\nu}\Big)\Big\}\,. \tag{2.34}$$

The correlation functions in the previous subsection are substituted into the above equation. Then, the integrals over $\vec{k}_{aT}$ and $\vec{k}_{aT}$ are manipulated by using the equations in Appendix of Ref. [10]. The calculations are rather lengthy particularly in the transverse h_1 part. However, because the g_1 and h_1 portions are the same as the ones in a spin-1/2 hadron [10] and the calculations of f_1, b_1, and c_1 terms are rather simple, we do not explain the calculation procedure. Noting that the b_1 and c_1 terms do not exist for the spin-1/2 hadron A, we obtain

$$W^{\mu\nu} = -g_\perp^{\mu\nu}\Big[W_T+\frac{1}{4}\,\lambda_A\,\lambda_B\,V_T^{LL}-\frac{2}{3}\,V_T^{UQ_0(1)}-2\,S_{B\perp}^2\,V_T^{UQ_0(2)}+2\,(\hat{x}\cdot S_{B\perp})^2\,V_T^{UQ_0(3)} \\ -\lambda_B\,\hat{x}\cdot S_{B\perp}\,V_T^{UQ_1}-\lambda_A\,\hat{x}\cdot S_{B\perp}\,V_T^{LT}-\hat{x}\cdot S_{A\perp}\,\lambda_B\,V_T^{TL}-S_{A\perp}\cdot S_{B\perp}\,V_T^{TT(1)}+\hat{x}\cdot S_{A\perp}\,\hat{x}\cdot S_{B\perp}\,V_T^{TT(2)}\Big] \\ -\Big(\hat{x}^\mu\hat{x}^\nu+\frac{1}{2}\,g_\perp^{\mu\nu}\Big)\Big[\frac{1}{4}\,\lambda_A\,\lambda_B\,V_{2,2}^{LL}-\lambda_A\,\hat{x}\cdot S_{B\perp}\,V_{2,2}^{LT}-\hat{x}\cdot S_{A\perp}\,\lambda_B\,V_{2,2}^{TL}-S_{A\perp}\cdot S_{B\perp}\,V_{2,2}^{TT(1)} \\ +\hat{x}\cdot S_{A\perp}\,\hat{x}\cdot S_{B\perp}\,V_{2,2}^{TT(2)}\Big]+\Big(\hat{x}^{\{\mu}S_{A\perp}^{\nu\}}-\hat{x}\cdot S_{A\perp}\,g_\perp^{\mu\nu}\Big)\Big(\hat{x}\cdot S_{B\perp}\,U_{2,2}^{TT(A)}-\lambda_B\,U_{2,2}^{TL}\Big) \\ +\Big(\hat{x}^{\{\mu}S_{B\perp}^{\nu\}}-\hat{x}\cdot S_{B\perp}\,g_\perp^{\mu\nu}\Big)\Big(\hat{x}\cdot S_{A\perp}\,U_{2,2}^{TT(B)}-\lambda_A\,U_{2,2}^{LT}\Big)-\Big(S_{A\perp}^{\{\mu}S_{B\perp}^{\nu\}}-S_{A\perp}\cdot S_{B\perp}\,g_\perp^{\mu\nu}\Big)\,U_{2,2}^{TT}\,. \tag{2.35}$$

The structure functions are expressed by the integral

$$I[d_1\bar{d}_2]\equiv\frac{1}{3}\sum_{a,b}\delta_{b\bar{a}}\,e_a^2\int d^2\vec{k}_{aT}\,d^2\vec{k}_{bT}\,\delta^2(\vec{k}_{aT}+\vec{k}_{bT}-\vec{Q}_T)\,d_1(x_A,\vec{k}_{aT}^{\,2})\,\bar{d}_2(x_B,\vec{k}_{bT}^{\,2})\,. \tag{2.36}$$

First, the unpolarized structure function is given by

$$W_T = I[f_1 \bar{f}_1] \ , \tag{2.37}$$

where the subscript T of W_T corresponds to the index combination of $(0,0) - (2,0)/3$ in the expressions of Ref. [7]. The structure functions associated with the factors $-g_\perp^{\mu\nu}$ and $\hat{z}^\mu \hat{z}^\nu$ in $W^{\mu\nu}$ are denoted as W_T and W_L. Obviously, W_L vanishes in the parton model. The longitudinal structure functions V_T^{LL} and $V_{2,2}^{LL}$ are

$$\begin{aligned} V_T^{LL} &= -\,4\, I[g_{1L}\, \bar{g}_{1L}] \ , \\ V_{2,2}^{LL} &= I\left[\left(\alpha + \beta - \frac{(\alpha-\beta)^2}{Q_T^2}\right) \frac{4\, h_{1L}^\perp\, \bar{h}_{1L}^\perp}{M_A M_B}\right] \ , \end{aligned} \tag{2.38}$$

where Q_T^2 is given by $\vec{Q}_T^2$ (note $Q_T^2 \neq -\vec{Q}_T^2$) and the variables α and β are defined by $\alpha = \vec{k}_{aT}^{\,2}$ and $\beta = \vec{k}_{bT}^{\,2}$. The tensor structure functions become

$$\begin{aligned} V_T^{UQ_0(1)} &= I[f_1 \bar{b}_1] \ , \\ V_T^{UQ_0(2)} &= I\left[\left(-Q_T^2 + 2(\alpha+\beta) - \frac{(\alpha-\beta)^2}{Q_T^2}\right) \frac{f_1\, \bar{b}_1}{2\,\beta}\right] \ , \\ V_T^{UQ_0(3)} &= I\left[\left(Q_T^2 - 2\alpha + \frac{(\alpha-\beta)^2}{Q_T^2}\right) \frac{f_1\, \bar{b}_1}{\beta}\right] \ , \\ V_T^{UQ_1} &= I\left[(Q_T^2 - \alpha + \beta) \frac{f_1\, \bar{c}_1}{2\, M_B\, Q_T}\right] \ . \end{aligned} \tag{2.39}$$

The superscripts U, Q_0, and Q_1 indicate the unpolarized state U, quadrupole polarization Q_0, and quadrupole polarization Q_1, respectively [7]. For example, $V_T^{UQ_1}$ indicates that the hadron A is unpolarized and B is polarized with the quadrupole polarization Q_1. The

longitudinal-transverse structure functions are

$$
\begin{aligned}
V_T^{LT} &= I\left[(-Q_T^2+\alpha-\beta)\frac{g_{1L}\,\bar{g}_{1T}}{2\,M_B\,Q_T}\right]\ ,\\
V_T^{TL} &= I\left[(-Q_T^2-\alpha+\beta)\frac{g_{1T}\,\bar{g}_{1L}}{2\,M_A\,Q_T}\right]\ ,\\
V_{2,2}^{LT} &= I\left[\left(\alpha\,Q_T^2+\beta^2+\alpha\beta-2\alpha^2+\frac{(\alpha-\beta)^3}{Q_T^2}\right)\frac{h_{1L}^{\perp}\,\bar{h}_{1T}^{\perp}}{M_A\,M_B^2\,Q_T}\right]\ ,\\
V_{2,2}^{TL} &= I\left[\left(\beta\,Q_T^2+\alpha^2+\alpha\beta-2\beta^2-\frac{(\alpha-\beta)^3}{Q_T^2}\right)\frac{h_{1T}^{\perp}\,\bar{h}_{1L}^{\perp}}{M_A^2\,M_B\,Q_T}\right]\ ,\\
U_{2,2}^{LT} &= I\left[(Q_T^2+\alpha-\beta)\frac{h_{1L}^{\perp}\,\bar{h}_{1T}}{2M_AQ_T}-\left(Q_T^2(\alpha-\beta)-2(\alpha^2-\beta^2)+\frac{(\alpha-\beta)^3}{Q_T^2}\right)\frac{h_{1L}^{\perp}\,\bar{h}_{1T}^{\perp}}{4\,M_A\,M_B^2\,Q_T}\right]\ ,\\
U_{2,2}^{TL} &= I\left[(Q_T^2-\alpha+\beta)\frac{h_{1T}\,\bar{h}_{1L}^{\perp}}{2\,M_B\,Q_T}+\left(Q_T^2(\alpha-\beta)-2(\alpha^2-\beta^2)+\frac{(\alpha-\beta)^3}{Q_T^2}\right)\frac{h_{1T}^{\perp}\,\bar{h}_{1L}^{\perp}}{4\,M_A^2\,M_B\,Q_T}\right]\ .
\end{aligned}
\tag{2.40}
$$

In the Drell-Yan process of identical hadrons, the structure functions V^{TL} and U^{TL} are equal to V^{LT} and U^{LT}. However, they are different in our reactions, so that both types are listed. The transverse structure functions are

$$
\begin{aligned}
V_T^{TT(1)} &= I\left[\left(-Q_T^2+2\alpha+2\beta-\frac{(\alpha-\beta)^2}{Q_T^2}\right)\frac{g_{1T}\,\bar{g}_{1T}}{4\,M_A\,M_B}\right]\ ,\\
V_T^{TT(2)} &= I\left[\left(-\alpha-\beta+\frac{(\alpha-\beta)^2}{Q_T^2}\right)\frac{g_{1T}\,\bar{g}_{1T}}{2\,M_A\,M_B}\right]\ ,\\
V_{2,2}^{TT(1)} &= I\left[\left(Q_T^2(\alpha+\beta)-(\alpha-\beta)^2-2(\alpha+\beta)^2+\frac{3(\alpha+\beta)(\alpha-\beta)^2}{Q_T^2}-\frac{(\alpha-\beta)^4}{Q_T^4}\right)\frac{h_{1T}^{\perp}\,\bar{h}_{1T}^{\perp}}{4\,M_A^2\,M_B^2}\right]\ ,\\
V_{2,2}^{TT(2)} &= I\left[\left(\alpha^2+\beta^2-\frac{2(\alpha+\beta)(\alpha-\beta)^2}{Q_T^2}+\frac{(\alpha-\beta)^4}{Q_T^4}\right)\frac{h_{1T}^{\perp}\,\bar{h}_{1T}^{\perp}}{M_A^2\,M_B^2}\right]\ ,\\
U_{2,2}^{TT(A)} &= I\Bigg[\left(Q_T^2-2\alpha+\frac{(\alpha-\beta)^2}{Q_T^2}\right)\frac{h_{1T}\bar{h}_{1T}^{\perp}}{2\,M_B^2}+\Bigg((\alpha-\beta)Q_T^2+(\alpha-\beta)^2-4\alpha(\alpha-\beta)\\
&\qquad+\frac{2\alpha(\alpha-\beta)^2+(\alpha-\beta)^2(\alpha+\beta)}{Q_T^2}-\frac{(\alpha-\beta)^4}{Q_T^4}\Bigg)\frac{h_{1T}^{\perp}\,\bar{h}_{1T}^{\perp}}{8\,M_A^2\,M_B^2}\Bigg]\ ,\\
U_{2,2}^{TT(B)} &= I\Bigg[\left(Q_T^2-2\beta+\frac{(\alpha-\beta)^2}{Q_T^2}\right)\frac{h_{1T}^{\perp}\bar{h}_{1T}}{2\,M_A^2}+\Bigg(-(\alpha-\beta)Q_T^2+(\alpha-\beta)^2+4\beta(\alpha-\beta)\\
&\qquad+\frac{2\beta(\alpha-\beta)^2+(\alpha-\beta)^2(\alpha+\beta)}{Q_T^2}-\frac{(\alpha-\beta)^4}{Q_T^4}\Bigg)\frac{h_{1T}^{\perp}\,\bar{h}_{1T}^{\perp}}{8\,M_A^2\,M_B^2}\Bigg]\ ,\\
U_{2,2}^{TT} &= I\left[h_{1T}\bar{h}_{1T}-\left(Q_T^2-2\alpha-2\beta+\frac{(\alpha-\beta)^2}{Q_T^2}\right)\left(\frac{h_{1T}^{\perp}\bar{h}_{1T}}{4\,M_A^2}+\frac{h_{1T}\bar{h}_{1T}^{\perp}}{4\,M_B^2}\right)\right]\ .
\end{aligned}
\tag{2.41}
$$

Substituting the hadron tensor of Eq. (2.35) and the lepton tensor of Eq. (2.7) into Eq.

(2.2), we obtain the cross section

$$\frac{d\sigma}{d^4Q\,d\Omega} = \frac{\alpha^2}{2\,s\,Q^2}\Bigg[(1+\cos^2\theta)\Bigg\{ W_T + \frac{1}{4}\lambda_A\lambda_B\,V_T^{LL} - \frac{2}{3}\,V_T^{UQ_0(1)} + 2\,|\vec{S}_{BT}|^2\,V_T^{UQ_0(2)}$$
$$+ 2\,|\vec{S}_{BT}|^2\cos^2\phi_B\,V_T^{UQ_0(3)} + \lambda_B\,|\vec{S}_{BT}|\cos\phi_B\,V_T^{UQ_1} + \lambda_A\,|\vec{S}_{BT}|\cos\phi_B\,V_T^{LT}$$
$$+ \lambda_B\,|\vec{S}_{AT}|\cos\phi_A\,V_T^{TL} + |\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos(\phi_A-\phi_B)\,V_T^{TT(1)} + |\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos\phi_A\cos\phi_B\,V_T^{TT(2)}\Bigg\}$$
$$+ \sin^2\theta\Bigg\{ \frac{1}{2}\cos 2\phi\left(\frac{1}{4}\lambda_A\lambda_B\,V_{2,2}^{LL} + \lambda_A\,|\vec{S}_{BT}|\cos\phi_B\,V_{2,2}^{LT} + \lambda_B\,|\vec{S}_{AT}|\cos\phi_A\,V_{2,2}^{TL}\right.$$
$$\left. + |\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos(\phi_A-\phi_B)(V_{2,2}^{TT(1)} + U_{2,2}^{TT(A)} + U_{2,2}^{TT(B)}) + |\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos\phi_A\cos\phi_B\,V_{2,2}^{TT(2)}\right)$$
$$+ |\vec{S}_{AT}|\cos(2\phi-\phi_A)\,\lambda_B\,U_{2,2}^{TL} + |\vec{S}_{BT}|\cos(2\phi-\phi_B)\,\lambda_A\,U_{2,2}^{LT}$$
$$+ \frac{1}{2}\sin 2\phi\,|\vec{S}_{AT}|\,|\vec{S}_{BT}|\,\sin(\phi_A-\phi_B)\,(U_{2,2}^{TT(A)} - U_{2,2}^{TT(B)})$$
$$+ |\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos(2\phi-\phi_A-\phi_B)(U_{2,2}^{TT} + U_{2,2}^{TT(A)}/2 + U_{2,2}^{TT(B)}/2)\Bigg\}\Bigg]\,. \qquad (2.42)$$

This equation indicates that $V_{2,2}^{TT(1)}$, $U_{2,2}^{TT(A)}$, $U_{2,2}^{TT(B)}$, and $U_{2,2}^{TT}$ cannot be measured independently. Only the combinations $V_{2,2}^{TT(1)} + U_{2,2}^{TT(A)} + U_{2,2}^{TT(B)}$, $U_{2,2}^{TT(A)} - U_{2,2}^{TT(B)}$, and $U_{2,2}^{TT} + U_{2,2}^{TT(A)}/2 + U_{2,2}^{TT(B)}/2$ can be studied experimentally. In our previous paper [7], we predicted that 108 structure functions exist in the Drell-Yan processes of spin-1/2 and spin-1 hadrons. According to Eq. (2.42), there are 19 independent ones in the naive parton model. It means that the rest of them are related to the neglected $O(1/Q)$ terms, namely the higher-twist structure functions. Although the $V_T^{UQ_0(1)}$ term may seem to contribute to the unpolarized cross section, it is canceled out by the other UQ_0-type terms in taking the spin average.

III. $\vec{Q}_T$ INTEGRATION AND THE LIMIT $Q_T \to 0$

Because the 108 structure functions are too many to investigate seriously and many of them are not important at this stage, the cross section is integrated over $\vec{Q}_T$ or it is calculated in the limit $Q_T \to 0$ [7]. There exist 22 structure functions in these cases, and they are considered to be physically significant. On the other hand, the cross section of

Eq. (2.42) is obtained at finite Q_T. In order to compare with the results in Ref . [7], we should investigate the parton-model cross section by taking the $\vec{Q}_T$ integration or the limit $Q_T \to 0$.

A. $\vec{\mathbf{Q}}_\mathbf{T}$-integrated cross section

First, we discuss the integration of the cross section over $\vec{Q}_T$. If the hadron tensor Eq. (2.34) is integrated over $\vec{Q}_T$, the delta function $\delta^2(\vec{k}_{aT}+\vec{k}_{bT}-\vec{Q}_T)$ disappears. It means that the integrations over $\vec{k}_{aT}$ and $\vec{k}_{bT}$ can be calculated separately. Then, the integrals with odd functions of $\vec{k}_T$ vanish: e.g. $\int d^2\vec{k}_{aT} F(\vec{k}_{aT}^{\,2}, \vec{k}_{bT}^{\,2})\vec{k}_{aT} = 0$. Furthermore, the vector Q_T^μ can not be used any longer in expanding, for example, the integral $\int d^2\vec{k}_{aT} d^2\vec{k}_{bT} F(\vec{k}_{aT}^{\,2}, \vec{k}_{bT}^{\,2}) k_{1\perp}^\mu k_{2\perp}^\nu$ in terms of the possible Lorentz-vector combinations. We calculate the hadron tensor in the similar way with the one in section II C. However, the calculations are much simpler because we do not have to take into account Q_T^μ. It does not make sense to explain the calculation procedure again, so that only the final results are shown in the following.

The $\vec{Q}_T$-integrated hadron tensor is expressed as

$$\overline{W}^{\mu\nu} = \int d^2\vec{Q}_T \, W^{\mu\nu} \ , \tag{3.1}$$

and in the same way for the structure functions. Then, the hadron tensor becomes

$$\begin{aligned} \overline{W}^{\mu\nu} = & - g_\perp^{\mu\nu} \left\{ \overline{W}_T + \frac{1}{4}\,\lambda_A\lambda_B\,\overline{V}_T^{LL} - 2\left(S_{B\perp}^2 + \frac{1}{3}\right)\overline{V}_T^{UQ_0} \right\} \\ & - \left(S_{A\perp}^{\{\mu} S_{B\perp}^{\nu\}} - S_{A\perp}\cdot S_{B\perp}\, g_\perp^{\mu\nu} \right) \overline{U}_{2,2}^{TT} \ , \end{aligned} \tag{3.2}$$

where the structure functions are written in terms of the parton distributions as

$$\begin{aligned} \overline{W}_T &= \frac{1}{3}\sum_a e_a^2\, f_1(x_A)\, \bar{f}_1(x_B) \ , \\ \overline{V}_T^{LL} &= -\frac{4}{3}\sum_a e_a^2\, g_1(x_A)\, \bar{g}_1(x_B) \ , \\ \overline{U}_{2,2}^{TT} &= \frac{1}{3}\sum_a e_a^2\, h_1(x_A)\, \bar{h}_1(x_B) \ , \\ \overline{V}_T^{UQ_0} &\equiv \overline{V}_T^{UQ_0(1)} = \overline{V}_T^{UQ_0(2)} = \frac{1}{3}\sum_a e_a^2\, f_1(x_A)\, \bar{b}_1(x_B) \ . \end{aligned} \tag{3.3}$$

The quark distributions are defined in the $\vec{k}_T$-integrated form as

$$f(x) = \int d^2\vec{k}_T\, f(x, \vec{k}_T^{\,2}) \qquad \text{for } f{=}f_1,\, g_1(=g_{1L}), \text{ and } b_1 \ , \tag{3.4}$$

$$h_1(x) = \int d^2\vec{k}_T \left[h_{1T}(x, \vec{k}_T^{\,2}) + \frac{\vec{k}_T^{\,2}}{2M^2} h_{1T}^{\perp}(x, \vec{k}_T^{\,2}) \right] \ , \tag{3.5}$$

and in the same way for the antiquark distributions. The other structure functions vanish by the $\vec{Q}_T$ integration:

$$\begin{aligned} &\overline{V}_T^{UQ_0(3)} = \overline{V}_T^{UQ_1} = \overline{V}_T^{LT} = \overline{V}_T^{TL} = \overline{V}_T^{TT(1)} = \overline{V}_T^{TT(2)} = \overline{V}_{2,2}^{LL} = \overline{V}_{2,2}^{LT} \\ &= \overline{V}_{2,2}^{TL} = \overline{V}_{2,2}^{TT(1)} = \overline{V}_{2,2}^{TT(2)} = \overline{U}_{2,2}^{LT} = \overline{U}_{2,2}^{TL} = \overline{U}_{2,2}^{TT(A)} = \overline{U}_{2,2}^{TT(B)} = 0 \ . \end{aligned} \tag{3.6}$$

With the expression of the hadron tensor in Eq. (3.2), the cross section becomes

$$\begin{aligned} \frac{d\sigma}{dx_A\, dx_B\, d\Omega} = \frac{\alpha^2}{4\,Q^2} \Bigg[&(1+\cos^2\theta) \left\{ \overline{W}_T + \frac{1}{4}\lambda_A\lambda_B\, \overline{V}_T^{LL} + \frac{2}{3}\left(2\,|\vec{S}_{BT}|^2 - \lambda_B^2\right) \overline{V}_T^{UQ_0} \right\} \\ &+ \sin^2\theta\, |\vec{S}_{AT}|\, |\vec{S}_{BT}| \cos(2\phi - \phi_A - \phi_B)\, \overline{U}_{2,2}^{TT} \Bigg] \ . \end{aligned} \tag{3.7}$$

The tensor distribution b_1 contributes to the cross section through the structure function $\overline{V}_T^{UQ_0}$. Because it is given by the multiplication of f_1 and $\bar{b}_1$ ($\bar{f}_1$ and b_1 in the opposite process) in Eq. (3.3), the quark and antiquark tensor distributions could be measured if the unpolarized distributions in the hadron A are well known. The $\bar{b}_1$ is paired with f_1; however, it is not with g_{1L} and h_1. This is because of the Fierz transformation in Eq. (2.13): $\Phi_{a/A}[\gamma^+]$ is multiplied by $\bar{\Phi}_{b/B}[\gamma^-]$ and not by the other factors $\bar{\Phi}_{b/B}[\gamma^-\gamma_5]$ and $\bar{\Phi}_{b/B}[i\sigma^{j-}\gamma_5]$, and the tensor distributions appear only in the functions $\Phi_{a/A}[\gamma^+]$ and $\bar{\Phi}_{b/B}[\gamma^-]$. Therefore, $\bar{b}_1$ can couple only with f_1, b_1, and c_1. Since the hadron A is a spin-1/2 particle, the distributions b_1 and c_1 do not exist. In this way, the only possible combination is $f_1(x_A)\bar{b}_1(x_B)$ for the process $q(\text{in A}){+}\bar{q}(\text{in B}){\to}\, \ell^+ + \ell^-$.

B. Cross section in the limit $\mathbf{Q_T \to 0}$

The transverse momentum Q_T is generally small in comparison with the dilepton mass Q. It originates mainly from intrinsic transverse momenta of the partons, so that its magnitude

is roughly restricted by the hadron size r: $Q_T \lesssim 1/r$. In this respect, it makes sense to consider the limit $Q_T \to 0$ for finding the essential part.

The structure functions in Eqs. (2.37)−(2.41) should be evaluated in this limit. As an example, we show how to take the limit for $V_T^{UQ_0(2)}$ in Eq. (2.39). The integration variables $\vec{k}_{aT}$ and $\vec{k}_{bT}$ in Eq. (2.36) are changed for $\vec{k}_T$ and $\vec{K}_T$, which are defined by $\vec{k}_T = (\vec{k}_{aT} - \vec{k}_{bT})/2$ and $\vec{K}_T = \vec{k}_{aT} + \vec{k}_{bT}$. Then, the delta function is integrated out and the structure function is expressed as

$$\begin{aligned} V_T^{UQ_0(2)} &= \frac{1}{3}\sum_a e_a^2 \frac{1}{2\,(\vec{k}_T - \vec{Q}_T/2)^2} \left\{ 4\,\vec{k}_T^{\,2} - \frac{4\,(\vec{k}_T \cdot \vec{Q}_T)^2}{Q_T^2} \right\} \\ &\qquad \times f_1\left(x_A, (\vec{k}_T + \vec{Q}_T/2)^2\right)\, \bar{b}_1\left(x_B, (\vec{k}_T - \vec{Q}_T/2)^2\right) \\ &= I_0[f_1 \bar{b}_1] \quad \text{in } Q_T \to 0 \ , \end{aligned} \tag{3.8}$$

where the function I_0 is defined by

$$I_0[d_1 \bar{d}_2] = \frac{1}{3}\sum_a e_a^2 \int d^2\vec{k}_T\, d_1(x_A, \vec{k}_T^{\,2})\, \bar{d}_2(x_B, \vec{k}_T^{\,2}) \ , \tag{3.9}$$

and $k_T^i k_T^j/\vec{k}_T^{\,2}$ in the second term of Eq. (3.8) is replaced by $\delta_T^{ij}/2$. The other structure functions are calculated in the same way. The finite ones are obtained as

$$\begin{aligned} &W_T = I_0[f_1 \bar{f}_1] \ , \quad V_T^{LL} = -4 I_0[g_{1L}\, \bar{g}_{1L}] \ , \quad V_T^{UQ_0(1)} = I_0[f_1 \bar{b}_1] = V_T^{UQ_0(2)} \equiv V_T^{UQ_0} \ , \\ &V_T^{TT(1)} = I_0\left[\frac{\vec{k}_T^{\,2}}{2\,M_A\,M_B}\, g_{1T}\, \bar{g}_{1T}\right] \ , \quad U_{2,2}^{TT} = I_0\left[h_{1T}\, \bar{h}_{1T} + \frac{\vec{k}_T^{\,2}}{2}\left(\frac{h_{1T}^\perp\, \bar{h}_{1T}}{M_A^2} + \frac{h_{1T}\, \bar{h}_{1T}^\perp}{M_B^2}\right)\right] \ , \\ &U_{2,2}^{TT(A)} = I_0\left[\frac{\vec{k}_T^{\,4}}{4\,M_A^2\,M_B^2}\, h_{1T}^\perp\, \bar{h}_{1T}^\perp\right] = U_{2,2}^{TT(B)} = -\frac{1}{2} V_{2,2}^{TT(1)} \ . \end{aligned} \tag{3.10}$$

The following ones vanish in the $Q_T \to 0$ limit:

$$\begin{aligned} &V_T^{UQ_0(3)} = V_T^{UQ_1} = V_T^{LT} = V_T^{TL} = V_T^{TT(2)} = V_{2,2}^{LL} \\ &= V_{2,2}^{LT} = V_{2,2}^{TL} = V_{2,2}^{TT(2)} = U_{2,2}^{LT} = U_{2,2}^{TL} = 0 \ . \end{aligned} \tag{3.11}$$

The hadron tensor is expressed by the finite structure functions as

$$
\begin{aligned}
W^{\mu\nu} = & -g_\perp^{\mu\nu}\left\{W_T + \frac{1}{4}\,\lambda_A\lambda_B\,V_T^{LL} - 2\left(S_{B\perp}^2 + \frac{1}{3}\right)\,V_T^{UQ_0} - S_{A\perp}\cdot S_{B\perp}\,V_T^{TT(1)}\right\} \\
& + \left(\hat{x}^\mu\hat{x}^\nu + \frac{1}{2}\,g_\perp^{\mu\nu}\right) S_{A\perp}\cdot S_{B\perp}\,V_{2,2}^{TT(1)} + \left(\hat{x}^{\{\mu}S_{A\perp}^{\nu\}} - \hat{x}\cdot S_{A\perp}\,g_\perp^{\mu\nu}\right)\hat{x}\cdot S_{B\perp}\,U_{2,2}^{TT(A)} \\
& + \left(\hat{x}^{\{\mu}S_{B\perp}^{\nu\}} - \hat{x}\cdot S_{B\perp}\,g_\perp^{\mu\nu}\right)\hat{x}\cdot S_{A\perp}\,U_{2,2}^{TT(B)} - \left(S_{A\perp}^{\{\mu}S_{B\perp}^{\nu\}} - S_{A\perp}\cdot S_{B\perp}\,g_\perp^{\mu\nu}\right)U_{2,2}^{TT}\ .
\end{aligned}
\tag{3.12}
$$

Defining $\hat{U}_{2,2}^{TT}$ by

$$
\begin{aligned}
\hat{U}_{2,2}^{TT} &= U_{2,2}^{TT} + U_{2,2}^{TT(A)}/2 + U_{2,2}^{TT(B)}/2 \\
&= I_0[h_1\bar{h}_1]\ ,
\end{aligned}
\tag{3.13}
$$

we obtain the cross section as

$$
\begin{aligned}
\frac{d\sigma}{d^4Q\,d\Omega} = \frac{\alpha^2}{2\,s\,Q^2}\Bigg[& (1+\cos^2\theta)\Bigg\{W_T + \frac{1}{4}\,\lambda_A\lambda_B\,V_T^{LL} + \frac{2}{3}\,(2|\vec{S}_{BT}|^2 - \lambda_B^2)\,V_T^{UQ_0} \\
& + |\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos(\phi_A-\phi_B)\,V_T^{TT(1)}\Bigg\} + \sin^2\theta\,|\vec{S}_{AT}|\,|\vec{S}_{BT}|\cos(2\phi-\phi_A-\phi_B)\,\hat{U}_{2,2}^{TT}\Bigg]\ .
\end{aligned}
\tag{3.14}
$$

Even though $V_{2,2}^{TT(1)}$ is finite, it does not contribute to the cross section. It is canceled out by the term $U_{2,2}^{TT(A)} + U_{2,2}^{TT(B)}$. As it is obvious from the results of this subsection and section III A, the expressions of the hadron tensors and the cross sections are slightly different in the $\vec{Q}_T$ integration and in the $Q_T \to 0$ limit.

Among many structure functions in Eq. (2.42), we have extracted the essential ones by taking the limit $Q_T \to 0$ or by the $\vec{Q}_T$ integration. According to the expressions of the cross section in Eqs. (3.7) and (3.14), merely the four or five structure functions remain finite: the unpolarized structure function W_T, the longitudinal one V_T^{LL}, the transverse one(s) $U_{2,2}^{TT}$ ($V_T^{TT(1)}$), and the unpolarized-quadrupole one $V_T^{UQ_0}$. Most of them are already known in the pp Drell-Yan reactions. The last quadrupole structure function $V_T^{UQ_0}$ is new in the Drell-Yan process of spin-1/2 and spin-1 hadrons.

IV. SPIN ASYMMETRIES AND PARTON DISTRIBUTIONS

We have derived the expressions for the Drell-Yan cross section in the parton model with finite Q_T, $\vec{Q}_T$ integration, and $Q_T \to 0$. Because the spin asymmetries are discussed in Ref. [7] in the latter two cases, they are shown in the $\vec{Q}_T$-integrated case as an example in this section. The polarized parton distributions for a spin-1 hadron are illustrated in Fig. 2. As it is obvious from Eq. (2.29), the longitudinally polarized (transversity) distribution is defined by the probability to find a quark with spin polarized along the longitudinal (transverse) spin of a polarized hadron minus the probability to find it polarized oppositely. On the other hand, the tensor polarized distribution b_1 is very different from these distributions according to Eq. (2.27). It is not associated with the quark polarization as shown by the unpolarized mark $\bullet$ in Fig. 2. It is related to the "unpolarized"-quark distribution in the polarized spin-1 hadron. The tensor distribution is essentially the difference between the unpolarized-quark distributions in the longitudinally and transversely polarized hadron states as it is given in Eq. (2.27).

In Ref. [7], fifteen spin combinations are suggested . However, most of them vanish in the parton model by the $\vec{Q}_T$ integration. There are only four finite structure functions, and three of them exist in the pp Drell-Yan processes. First, the unpolarized cross section is

$$\begin{aligned}\left\langle \frac{d\sigma}{dx_A\, dx_B\, d\Omega} \right\rangle &= \frac{\alpha^2}{4\,Q^2}\,(1+\cos^2\theta)\,\overline{W}_T \\ &= \frac{\alpha^2}{4\,Q^2}\,(1+\cos^2\theta)\,\frac{1}{3}\sum_a e_a^2\,\left[\, f_1(x_A)\,\bar{f}_1(x_B) + \bar{f}_1(x_A)\,f_1(x_B)\,\right]\ . \end{aligned} \tag{4.1}$$

Next, the longitudinal and transverse double spin asymmetries are

$$\begin{aligned} A_{LL} &= \frac{\sigma(\downarrow_L,+1_L)-\sigma(\uparrow_L,+1_L)}{2<\sigma>} = -\frac{\overline{V}_T^{LL}}{4\,\overline{W}_T} \\ &= \frac{\sum_a e_a^2\,\left[\, g_1(x_A)\,\bar{g}_1(x_B) + \bar{g}_1(x_A)\,g_1(x_B)\,\right]}{\sum_a e_a^2\,\left[\, f_1(x_A)\,\bar{f}_1(x_B) + \bar{f}_1(x_A)\,f_1(x_B)\,\right]}\ , \\ A_{TT} &= \frac{\sigma(\phi_A=0,\phi_B=0)-\sigma(\phi_A=\pi,\phi_B=0)}{2<\sigma>} = \frac{\sin^2\theta\,\cos 2\phi}{1+\cos^2\theta}\,\frac{\overline{U}_{2,2}^{TT}}{\overline{W}_T} \\ &= \frac{\sin^2\theta\,\cos 2\phi}{1+\cos^2\theta}\,\frac{\sum_a e_a^2\,\left[\, h_1(x_A)\,\bar{h}_1(x_B) + \bar{h}_1(x_A)\,h_1(x_B)\,\right]}{\sum_a e_a^2\,\left[\, f_1(x_A)\,\bar{f}_1(x_B) + \bar{f}_1(x_A)\,f_1(x_B)\,\right]}\ , \end{aligned} \tag{4.2}$$

where we have explicitly written the contribution from the process $\bar{q}$(in A)+q(in B)$\to \ell^+ + \ell^-$ in addition to the one from q(in A)+$\bar{q}$(in B)$\to \ell^+ + \ell^-$. The subscripts of $\uparrow_L$, $\downarrow_L$, and $+1_L$ indicate the longitudinal polarization. If ϕ_A or ϕ_B is indicated in the expression of $\sigma(pol_A, pol_B)$, it means that the hadron A or B is transversely polarized with the azimuthal angle ϕ_A or ϕ_B. The above asymmetries are given by the spin-flip cross sections in the hadron A, so that A_{TT} corresponds to $A_{T(T)}$ in Ref. [7]. However, the other asymmetry expressions of Ref. [7] become the same: $A_{TT}^{\parallel} = A_{(T)T} = A_{T(T)}$ in the parton model. Furthermore, the perpendicular transverse-transverse asymmetry is simply given by $A_{TT}^{\perp} = \tan 2\phi\, A_{TT}^{\parallel}$. It should be noted that our definitions of the asymmetries are slightly different from the usual ones in the pp and ep scattering. The asymmetries are defined so as to exclude the b_1 contributions to the denominator. If the usual definition $[\sigma(\uparrow -1) - \sigma(\uparrow +1)]/[\sigma(\uparrow -1) + \sigma(\uparrow +1)]$ is used, the b_1 contributes, for example, to the longitudinal asymmetry A_{LL}. The tensor distribution can be investigated by the unpolarized-quadrupole Q_0 asymmetry

$$A_{UQ_0} = \frac{1}{2 <\sigma>} \left[\sigma(\bullet, 0_L) - \frac{\sigma(\bullet, +1_L) + \sigma(\bullet, -1_L)}{2} \right] = \frac{\overline{V}_T^{UQ_0}}{\overline{W}_T}$$
$$= \frac{\sum_a e_a^2 \left[f_1(x_A)\, \bar{b}_1(x_B) + \bar{f}_1(x_A)\, b_1(x_B) \right]}{\sum_a e_a^2 \left[f_1(x_A)\, \bar{f}_1(x_B) + \bar{f}_1(x_A)\, f_1(x_B) \right]} , \tag{4.3}$$

where the filled circle indicates that the hadron A is unpolarized. The other asymmetries vanish

$$A_{LT} = A_{TL} = A_{UT} = A_{TU} = A_{TQ_0} = A_{UQ_1} = A_{LQ_1} = A_{TQ_1}$$
$$= A_{UQ_2} = A_{LQ_2} = A_{TQ_2} = 0 \quad \text{by the } \vec{Q}_T \text{ integration.} \tag{4.4}$$

The only new finite asymmetry, which does not exist in the pp reactions, is the unpolarized-quadrupole Q_0 asymmetry A_{UQ_0}. In order to measure this quantity, we use a longitudinally and transversely polarized spin-1 hadron with an unpolarized hadron A. Then, the quadrupole Q_0 spin combination [7] should be taken. The b_1 and $\bar{b}_1$ distributions could be extracted from the asymmetry A_{UQ_0} measurements with the information on the unpolarized parton distributions $f_1(x)$ and $\bar{f}_1(x)$ in the hadrons A and B. If it is difficult to attain the longitudinally polarized deuteron in the collider experiment, we may combine

the transversely polarized cross sections with the unpolarized one without resorting to the longitudinal polarization. Alternatively, the fixed deuteron target may be used for obtaining A_{UQ_0}.

There is an important advantage to use the Drell-Yan process for measuring $\bar{b}_1$ over the lepton scattering. In the large Feynman-x ($x_F = x_A - x_B$) region, the antiquark distribution $\bar{f}_1(x_A)$ is very small in comparison with the quark one $f_1(x_A)$. Then, the cross section is dominated by the annihilation process $q(A) + \bar{q}(B) \to \ell^+ + \ell^-$, and the asymmetry becomes

$$A_{UQ_0}(\text{large } x_F) \approx \frac{\sum_a e_a^2 \, f_1(x_A) \, \bar{b}_1(x_B)}{\sum_a e_a^2 \, f_1(x_A) \, \bar{f}_1(x_B)} \ . \tag{4.5}$$

This equation means that the antiquark tensor distributions could be extracted if the unpolarized distributions are well known in the hadrons A and B. We note in the electron-scattering case that the b_1 sum rule is written by the parton model as [3]

$$\int dx \, b_1^e(x) = \lim_{t \to 0} - \frac{5}{3} \frac{t}{4M^2} F_Q(t) + \delta Q_{sea} \ , \tag{4.6}$$

where M is the hadron mass, $F_Q(t)$ is the quadrupole form factor in the unit of $1/M^2$, and δQ_{sea} is the antiquark tensor polarization, for example $\delta Q_{sea} = \int dx [8\delta\bar{u}(x) + 2\delta\bar{d}(x) + \delta s(x) + \delta\bar{s}(x)]/9$ for the deuteron. The first term of Eq. (4.6) vanishes, so that the difference from the sum rule $\int dx \, b_1^e = 0$ could suggest a finite tensor polarization of the antiquarks. Although the sum rule is valid within the quark model, the nuclear shadowing effects should be taken into account properly [4] as far as the deuteron is concerned. On the other hand, the antiquark tensor polarization $\delta\bar{q}$ (or $\bar{b}_1$) could be studied independently by the Drell-Yan process. As the violation of the Gottfried sum rule leads to the light antiquark flavor asymmetry $\bar{u} \neq \bar{d}$ and the difference was confirmed by the Drell-Yan experiments [16], the antiquark tensor polarization could be investigated by both methods: the sum rule of Eq. (4.6) in the lepton scattering and the Drell-Yan process measurement. However, the advantage of the Drell-Yan process is that the antiquark distribution $\bar{b}_1(x)$ is directly measured even though it is difficult in the lepton scattering.

Furthermore, the flavor asymmetry in the polarized antiquark distributions ($\Delta\bar{u} \neq \Delta\bar{d}$, $\Delta_T\bar{u} \neq \Delta_T\bar{d}$) could be investigated by combining pp and pd Drell-Yan data. It is particularly

important for the transversity distributions because the $\Delta_T\bar{u}/\Delta_T\bar{d}$ asymmetry cannot be found in the $W^{\pm}$ production processes due to the chiral-odd nature.

In this way, we find that a variety of interesting topics are waiting to be studied in connection with the new structure functions for spin-1 hadrons. In particular, we have shown in this paper that the tensor distributions b_1 and $\bar{b}_1$ could be measured by the asymmetry A_{UQ_0} in the Drell-Yan process of spin-1/2 and spin-1 hadrons. A realistic possibility is the proton-deuteron Drell-Yan experiment, and it may be realized, for example at RHIC [5]. In order to support the deuteron polarization project in experimental high-energy spin physics, we should investigate theoretically more about the spin structure of spin-1 hadrons.

V. SUMMARY

We have investigated the Drell-Yan processes of spin-1/2 and spin-1 hadrons in a naive parton model by ignoring the $Q(1/Q)$ terms. First, the quark and antiquark correlation functions are expressed by the combinations of possible Lorentz vectors and pseudovectors by taking into account the Hermiticity, parity conservation, and time-reversal invariance. Then, we have shown that the tensor distributions b_1 and c_1 , which are specific for a spin-1 hadron, are involved in the correlation function $\Phi[\gamma^+]$. The expressions of the other functions $\Phi[\gamma^+\gamma_5]$ and $\Phi[i\sigma^{i+}\gamma_5]$ are the same as those of the spin-1/2 hadron in terms of the longitudinal and transverse distributions. Using the obtained correlation functions, we have calculated the hadron tensor and the Drell-Yan cross section. We found that there exist 19 independent structure functions in the parton model. Next, we studied two cases: the $\vec{Q}_T$ integration and the limit $Q_T \to 0$. In these cases, the c_1 contribution vanishes, and there are only four or five finite structure functions: the unpolarized, longitudinally polarized, transversely polarized, and tensor polarized structure functions. The last one is related to the tensor polarized distributions b_1 and $\bar{b}_1$. Although the tensor structure function could be measured in the lepton scattering, the Drell-Yan measurements are valuable for finding particularly the antiquark tensor polarization. In addition to these topics, there are

a number of higher-twist structure functions in the Drell-Yan processes, and they should be also studied in detail.

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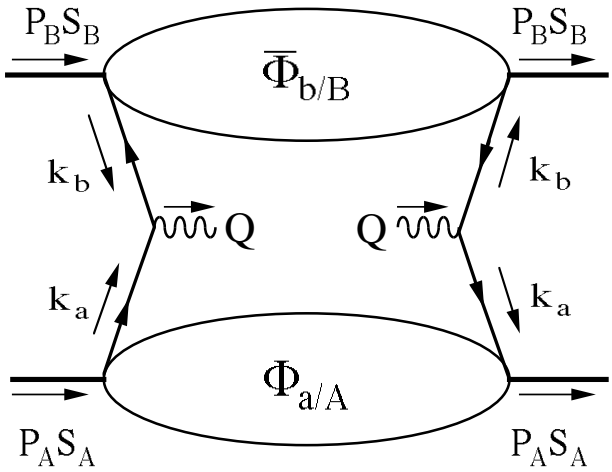

FIG. 1. Drell-Yan process in a parton model.

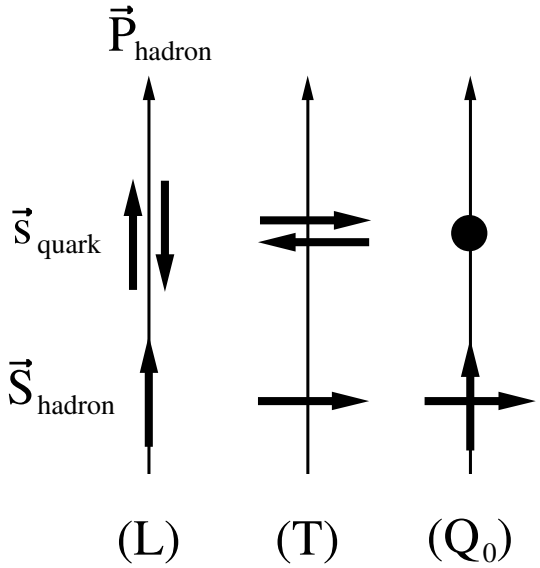

FIG. 2. Polarized parton distributions in a spin-1 hadron are illustrated. The figures (L), (T), and (Q_0) indicate the longitudinally polarized, transversity, and tensor polarized distributions. The mark $\bullet$ indicates the unpolarized state.