SPIN STRUCTURE FUNCTIONS IN LEPTOPRODUCTION

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We discuss a few examples of structure functions for polarized, semi-inclusive scattering processes to show the richness of structure. Then we indicate how polarization and particle production can be used to study the quark and gluon structure of hadrons going further than the well-known parton densities and fragmentation functions.

1 Structure functions

We start our discussion with the object of interest for 1-particle inclusive leptoproduction, the hadronic tensor, given by

\[ 2MW_{\mu\nu}^{(H)}(q; PS; P_hS_h) = \frac{1}{(2\pi)^4} \int \frac{d^3P_X}{(2\pi)^32P_X^0} (2\pi)^4\delta^4(q + P - P_X - P_h) \times \langle PS|J_{\mu}(0)|P_X; P_hS_h\rangle\langle P_X; P_hS_h|J_{\nu}(0)|PS\rangle, \]

(1)

where \( P, S \) and \( P_h, S_h \) are the momenta and spin vectors of target hadron and produced hadron, \( q \) is the (spacelike) momentum transfer with \(-q^2 = Q^2\) sufficiently large. The kinematics is illustrated in Fig. 1, where also the

\[
\begin{align*}
x_B &= \frac{Q^2}{2P\cdot q} \\
y &= \frac{P\cdot q}{P\cdot k} \\
z_h &= \frac{P\cdot P_h}{P\cdot q}
\end{align*}
\]

Figure 1: Kinematics for 1-particle inclusive leptoproduction.

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scaling variables are introduced. For inclusive scattering (unpolarized lepton and hadron, $\gamma$-exchange) the most general symmetric part of the hadronic tensor is

$$2MW^{\mu\nu}(q,P) = \begin{cases} -g^{\mu\nu} + \hat{q}^{\mu} \hat{q}^{\nu} - \hat{t}^{\mu} \hat{t}^{\nu} \\ -g^{\mu\nu}_{\perp} \end{cases} F_1 + \hat{t}^{\mu} \hat{t}^{\nu} \left( F_2 - \frac{2x_B}{F_1} F_1 \right).$$

(2)

Combined with the leptonic part, one obtains the cross section

$$\frac{d\sigma_O}{dx_B dy} = \frac{4\pi \alpha_s^2 x_B s}{Q^4} \left\{ \left( 1 - y + \frac{1}{2} y^2 \right) F_T + (1 - y) F_L \right\}. \tag{3}$$

In order to calculate the hadronic tensor, a diagrammatic expansion is written down starting with the well-known handbag diagram (see Fig. 2, left), yielding the parton model results for the structure functions,

$$F_T(x_B, Q) = F_1(x_B, Q) = \frac{1}{2} \sum_{a, \bar{a}} e_a^2 f_1^a(x_B), \tag{4}$$

$$F_L(x_B, Q) = 0, \tag{5}$$

expressed in terms of the quark distribution $f_1^a$ ($a$ is the flavor index). The summation runs over quarks and antiquarks. The most general antisymmetric

$$\hat{q}^{\mu} = q^{\mu}/Q, \quad \hat{t}^{\mu} = \hat{p}^{\mu}/\sqrt{\hat{p}^2} = \left( p^{\mu} - \frac{P \cdot q}{q^2} q^{\mu} \right)/\sqrt{\hat{p}^2}. \tag{6}$$
part of the hadronic tensor involves polarized leptons and hadrons and is for $\gamma$-exchange given by

$$2MW^\mu_\nu_A(q, P, S) = -i \lambda \frac{\epsilon^{\mu\nu\rho\sigma} P_\rho q_\sigma}{P \cdot q} g_1 + i \frac{2Mx_B}{Q} \hat{t} |\nu e^\rho| S^\rho_\perp g_T \tag{6}$$

with $\lambda \equiv q \cdot S / q \cdot P$ and $S_\perp$ the transverse spin vector obtained with the help of $g_{\mu\nu}$. The cross section becomes

$$\frac{d\sigma_L}{dx_q dy} = \lambda e^2 \frac{4\pi \alpha^2}{Q^2} \left\{ \lambda \left( 1 - \frac{y}{2} \right) g_1 - |S_\perp| \cos \phi_S \frac{2Mx_B}{Q} \sqrt{1 - y} g_T \right\}, \tag{7}$$

with the parton model results

$$g_1(x_B, Q) = \frac{1}{2} \sum_{a, \bar{a}} e^2_a g^a_1(x_B), \tag{8}$$

$$g_T(x_B, Q) = (g_1 + g_2)(x_B, Q) = \frac{1}{2} \sum_{a, \bar{a}} e^2_a g^a_T(x_B). \tag{9}$$

The function $g^a_1$ is the quark helicity distribution. The function $g^a_T$ is a higher twist distribution.

Proceeding to the 1-particle inclusive case for unpolarized lepton and hadron, we obtain generally for the symmetric part of the hadronic tensor

$$2MW^\mu_\nu^S(q, P, P_h) = -g^\mu_\perp \mathcal{H}_T + \hat{\mu} \hat{\nu} \mathcal{H}_L$$

$$+ \hat{t} (\hat{\mu} \hat{\nu}) \mathcal{H}_{LT} + \left\{ 2 \hat{h}^\mu \hat{h}^\nu + g^\mu_\perp \right\} \mathcal{H}_{TT}, \tag{10}$$

leading to the unpolarized cross section

$$\frac{d\sigma_O}{dx_q dy dx_h dz_h d^2q_T} = \frac{4\pi \alpha^2 s}{Q^4} x_B z_h \left\{ \left( 1 - y + \frac{1}{2} y^2 \right) \mathcal{H}_T + (1 - y) \mathcal{H}_L \right. $$

$$\left. - (2 - y) \sqrt{1 - y} \cos \phi^\ell \mathcal{H}_{LT} + (1 - y) \cos 2\phi^\ell \mathcal{H}_{TT} \right\}. \tag{11}$$

where

$$\hat{q}^\mu = q^\mu / Q, \quad \hat{p}^\mu = (q^\mu + 2x_B P^\mu) / Q,$$

$$\hat{q}_T^\mu = q^\mu + x_B P^\mu - \frac{P_h^\mu}{z_h} = -P_{h_\perp}^\mu / z_h \equiv -Q_T \hat{h}^\mu.$$
We will come back to the parton expressions for these structure functions later with emphasis on the azimuthal dependence, the $\cos \phi_h^\ell$ and $\cos 2\phi_h^\ell$ parts depending on the azimuthal angle between the lepton scattering plane and the production plane (see Fig. 1). Limiting ourselves to unpolarized leptons, the antisymmetric part of the hadronic tensor is

$$2M W_A^{\mu\nu}(q, P, P_h) = -i \hat{t}^{[\mu} \hat{H}^{\nu]} \mathcal{H}_{LT}' ,$$  \hspace{1cm} (12) $$

leading to the cross section

$$\frac{d\sigma_L}{dx_d dy d z_h d^2 q_T} = \lambda e \frac{4\pi \alpha^2}{Q^2} z_h \sqrt{1 - y} \sin \phi_h^\ell \mathcal{H}_{LT}' .$$  \hspace{1cm} (13) $$

Our aim in studying leptoproduction is the study of the quark and gluon structure of the hadronic target using the known framework of Quantum chromodynamics (QCD). Thus, as a theorist the aim is to calculate the hadronic tensor $W_{\mu\nu}$ by making a diagrammatic expansion. Already at the simplest level (Fig. 2) a problem is encountered, namely there are hadrons involved for which QCD does not provide rules. Thus, soft parts are identified that allow inclusion of hadrons in the field theoretical framework. Furthermore it will turn out that for $Q^2 \to \infty$ only a limited number of diagrams is needed.

## 2 Soft parts

### 2.1 Definition as quark operators

Next, we look in more detail to the soft parts, such as appear for instance in the parton diagram. They can be written down in terms of quark and gluon fields as illustrated below. They are characterized by the fact that the momenta are soft with respect to each other. We have for the distribution part

$$\Phi_{ij}(p, P, S) = \frac{1}{(2\pi)^4} \int d^4x \ e^{ip\cdot x} \langle P, S|\bar{\psi}_j(0)\psi_i(x)|P, S\rangle ,$$  \hspace{1cm} (14) $$

and the fragmentation part
represented by

\[ \Delta_{ij}(k, P_h, S_h) = \sum_X \frac{1}{(2\pi)^4} \int d^4x \ e^{ikx} \langle \psi_i(x) | P_h, S_h; X | \psi_j(0) \rangle \langle P_h, S_h; X | \psi_j(0) \rangle. \]  

(15)

In order to find out which information in the soft parts \( \Phi \) and \( \Delta \) is important in a hard process one needs to realize that the hard scale \( Q \) leads in a natural way to the use of lightlike vectors \( n_+ \) and \( n_- \) satisfying \( n_+^2 = n_-^2 = 0 \) and \( n_+ \cdot n_- = 1 \). For 1-particle inclusive scattering one parametrizes the momenta

\[
\begin{align*}
q^2 &= -Q^2 \\
P^2 &= M^2 \\
P_h^2 &= M_h^2 \\
2P \cdot q &= \frac{Q^2}{2} \\
2P_h \cdot q &= -z_h Q^2
\end{align*}
\]  

Comparing the power of \( Q \) with which the momenta in the soft and hard part appear one immediately is led to \( \int dp^- \Phi(p, P, S) \) and \( \int dk^+ \Delta(k, P_h, S_h) \) as the relevant quantities to investigate.

2.2 Analysis of soft parts: distribution and fragmentation functions

Hermiticity, parity and time reversal invariance (T) constrain the quantity \( \Phi(p, P, S) \) and therefore also the Dirac projections \( \Phi^{[\Gamma]} \) defined as

\[ \Phi^{[\Gamma]}(x, p_T) = \int dp^- \frac{Tr[\Phi^\Gamma]}{2} \]
\[
\int \frac{d\xi d^2\xi_T}{2(2\pi)^3} e^{ip\cdot\xi} \langle P, S | \overline{\psi}(0) \Gamma \psi(\xi) | P, S \rangle \bigg|_{\xi^+ = 0},
\]
which is a lightfront \((\xi^+ = 0)\) correlation function. The relevant projections in \(\Phi\) that are important in leading order in \(1/Q\) in hard processes are

\[
\begin{align*}
\Phi^{[\gamma^+]}(x, p_T) &= f_1(x, p_T^2) \\
\Phi^{[\gamma^+\gamma_5]}(x, p_T) &= \lambda g_{1L}(x, p_T^2) + \frac{(p_T \cdot S_T)}{M} g_{1T}(x, p_T^2) \\
\Phi^{[\gamma^+\gamma_5]}(x, p_T) &= h_1(x, p_T^2) + \frac{\lambda p_T^2}{M} h_{1L}(x, p_T^2) \\
&- \left( \frac{p_T^2 p_T^2 + 2 p_T^2 g_T^2}{M^2} \right) h_{1T}(x, p_T^2)
\end{align*}
\]

Here \(x = p^+/P^+, \lambda = MS^+/P^+\) and \(S_T\) is the spin-component projected out by \(g_T^{\mu\nu} = g^{\mu\nu} - n_+^{(\mu)} n_+^{(\nu)}\). They satisfy \(\lambda^2 + S_T^2 = 0\).

All functions appearing above can be interpreted as momentum space densities, as illustrated in Fig. 3. The ones denoted \(f_{\ldots}\) involve the operator structure \(\overline{\psi} \gamma^+ \psi = \psi^\dagger \gamma^+ \psi\), where \(\psi_+ = P_+ \psi\) with \(P_+ = \gamma^- \gamma^+/2\). This operator projects on the so-called good component of the Dirac field, which can be considered as a free dynamical degree of freedom in front form quantization. It is precisely in this sense that partons measured in hard processes are free.

The functions \(g_{\ldots}\) and \(h_{\ldots}\) appearing above are differences of densities involving good fields, but in addition projection operators \(P_{R/L} = (1 \pm \gamma_5)/2\) and \(P_{1/1} = (1 \pm \gamma^+ \gamma_5)/2\), all of which commute with \(P_+\). To be precise for the functions \(g_{\ldots}\) one has \(\psi_\gamma^+ \gamma_5 \psi = \psi^\dagger_{+,R} \psi_{+,R} - \psi^\dagger_{+,L} \psi_{+,L}\) while in the case of \(h_{\ldots}\) one has \(\psi \sigma^+ \gamma_5 \psi = \psi_{+,1}^\dagger \psi_{+,1} - \psi_{+,1}^\dagger \psi_{+,1}\).

Figure 3: Interpretation of the functions in the leading Dirac projections of \(\Phi\).
It is useful to remark here that flavor indices have been omitted, i.e. one has \( f_1^c, f_1^d, \) etc. At this point it may also be good to mention other notations used frequently such as \( f_1^c(x) = u(x), g_1^c(x) = \Delta u(x), h_1^c(x) = \Delta_T u(x), \) etc. These \( x \)-dependent functions are the ones obtained after integration over \( p_T \).

The analysis of the soft part \( \Phi \) can be extended to other Dirac projections. Limiting ourselves to \( p_T \)-averaged functions one finds

\[
\Phi^{[1]}(x) = \frac{M}{P^+} e(x),
\]

\[
\Phi^{[\gamma \gamma]}(x) = \frac{M S^i}{P^+} g_T(x),
\]

\[
\Phi^{[i\sigma^- \gamma_5]}(x) = \frac{M}{P^+} \lambda h_L(x).
\]

Lorentz covariance requires for these projections on the right hand side a factor \( M/P^+ \), which as can be seen from the earlier given parametrization of momenta produces a suppression factor \( M/Q \) and thus these functions appear at subleading order in cross sections. The constraints on \( \Phi \) lead to relations between the above higher twist functions and \( p_T/M^2 \)-weighted functions, e.g.

\[
g_2 = g_T - g_1 = \frac{d}{dx} g^{(1)}_{1T},
\]

where

\[
g^{(1)}_{1T}(x) = \int d^2 p_T \frac{p_T^2}{2M^2} g_{1T}(x, p_T).
\]

We will use the index \((1)\) to indicate a \( p_T^2 \)-moment of the above type.

Just as for the distribution functions one can perform an analysis of the soft part describing the quark fragmentation. The Dirac projections are

\[
\Delta^{[\Gamma]}(z,k_T) = \int dk^+ \frac{Tr[\Delta^\Gamma]}{4z} = \sum_X \int \frac{d\xi^+ d^2 \xi_T}{4z (2\pi)^3} e^{ik_\xi} Tr\langle 0|\psi(x)|P_h, X\rangle \langle P_h, X|\bar{\psi}(0)\Gamma|0\rangle
\]

\[
\xi^- = 0
\]

The relevant projections in \( \Delta \) that appear in leading order in \( 1/Q \) in hard processes are for the case of no final state polarization,

\[
\Delta^{[\gamma^-]}(z,k_T) = D_1(z,-z k_T),
\]

\[
\Delta^{[i\sigma^- \gamma_5]}(z,k_T) = \frac{\gamma^j_{ij} k_T}{M_h} H^+_1(z,-z k_T). \quad [T\text{-odd}]
\]
The arguments of the fragmentation functions $D_1$ and $H_1^\perp$ are chosen to be $z = P_h^- / k^- \equiv P_{h_\perp} = - z k_T$. The first is the (lightcone) momentum fraction of the produced hadron, the second is the transverse momentum of the produced hadron with respect to the quark. The fragmentation function $D_1$ is the equivalent of the distribution function $f_1$. It can be interpreted as the probability of finding a hadron $h$ in a quark. Noteworthy is the appearance of the function $H_1^\perp$, interpretable as the different production probability of unpolarized hadrons from a transversely polarized quark (see Fig. 4). This function has no equivalent in the distribution functions and is allowed because of the non-applicability of time reversal invariance because of the appearance of out-states $|P_{h_\perp},X\rangle$ in $\Delta$, rather than the plane wave states in $\Phi$.

After $k_T$-averaging one is left with the functions $D_1(z)$ and the $k_T/M$-weighted result $H_1^{1(1)}(z)$. We summarize the full analysis of the soft part with a table of distribution and fragmentation functions for unpolarized (U), longitudinally polarized (L) and transversely polarized (T) targets, distinguishing leading (twist two) and subleading (twist three, appearing at order $1/Q$) functions and furthermore distinguishing the chirality $\Gamma$. The functions printed in boldface survive after integration over transverse momenta. We have for the distributions included a separate table with distribution functions that can exist without the T constraint, suggested to explain single spin asymmetries $^{7,8,9}$. We have included them in our complete classification scheme.

### Classification of distribution and fragmentation functions:

<table>
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<th>DISTRIBUTIONS (T-even)</th>
<th>DISTRIBUTIONS (T-odd)</th>
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</tr>
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<tr>
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<tr>
<td>$h_T$</td>
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</tbody>
</table>
3 Cross sections for lepton-hadron scattering

Having completed the analysis of the soft parts, the next step is to find out how one obtains the information on the various correlation functions from experiments, in this paper in particular lepton-hadron scattering via one-photon exchange as discussed in section 1. To get the leading order result for semi-inclusive scattering it is sufficient to compute the diagram in Fig. 2 (right) by using QCD and QED Feynman rules in the hard part and the matrix elements $\Phi$ and $\Delta$ for the soft parts, parametrized in terms of distribution and fragmentation functions. The results are:

\[
\begin{align*}
\frac{d\sigma_{OO}}{dx_B dy dz_h} &= \frac{2\pi\alpha^2 s}{Q^4} \sum_{a,a'} c_a^2 \left( 1 + (1 - y)^2 \right) x_B f^0_T(x_B) D^{a}_T(z_h) \tag{28} \\
\frac{d\sigma_{LL}}{dx_B dy dz_h} &= \frac{2\pi\alpha^2 s}{Q^4} \lambda c \lambda \sum_{a,a'} c_a^2 y(2 - y) x_B g_T^0(x_B) D^{a}_T(z_h) \tag{29}
\end{align*}
\]

Comparing with the expressions in section 1, one can identify the structure function $H_T$ and deduce that in leading order $\alpha_s^0$ the function $H_L = 0$.

It is not difficult to give some general rules on how the distribution and fragmentation functions are encountered in experiments. We will just give a few examples.

In 1-particle inclusive processes, one actually becomes sensitive to quark transverse momentum dependent distribution functions. One finds at order $1/Q$ the following nonvanishing azimuthal asymmetries.
Azimuthal asymmetries for unpolarized targets (higher twist)

\[ \int d^2 q_T \frac{Q_T}{M} \cos(\phi_h^Q) \frac{d\sigma_{OO}}{d x_B \ dy \ dz_h \ d^2 q_T} = - \frac{2\pi\alpha^2 \ s \ Q^4}{Q^4} \ 2(2 - y) \sqrt{1 - y} \]
\[ \times \sum_{a, \bar{a}} e_a^2 \left\{ \frac{2M}{Q} x_B^2 f_{1}^{(1) a}(x_B) D_{1}^{a}(zh) \right. \]
\[ + \frac{2M_h}{Q} x_B f_{1}^{1}(x_B) \left( \frac{\tilde{D}_{1}^{(1) a}(zh)}{z_h} \right) \]  

(30)

Note: \( \tilde{D}_{1}^{a}(z) = D_{1}^{a}(z) - z D_{1}^{1}(z) \),

\[ \int d^2 q_T \frac{Q_T}{M} \sin(\phi_h^Q) \frac{d\Delta\sigma_{LO}}{d x_B \ dy \ dz_h \ d^2 q_T} = \frac{2\pi\alpha^2 \ s \ Q^4}{Q^4} \ \lambda_e \ 2y \sqrt{1 - y} \]
\[ \times \sum_{a, \bar{a}} e_a^2 \frac{2M}{Q} x_B^2 e_a(x_B) H_{1}^{(1) a}(zh) \]  

(31)

Note: \( \tilde{e}^a(x) = e^a(x) - \frac{m_a}{M} f_{1}^{(1) a}(x) \)

The first weighted cross section given here involves the structure function \( H_{LT} \) and contains the twist three distribution function \( f_{1}^{1} \) and the fragmentation function \( D_{1}^{1} \). The second cross section involves the structure function containing the distribution function \( e \) and the time-reversal odd fragmentation function \( H_{1}^{1} \). The tilde functions that appear in the cross sections are in fact precisely the so-called interaction dependent parts of the twist three functions. They would vanish in any naive parton model calculation in which cross sections are obtained by folding electron-parton cross sections with parton densities. Considering the relation for \( \tilde{e} \) one can state it as \( x e(x) = (m/M) f_{1}(x) \) in the absence of quark-quark-gluon correlations. The inclusion of the latter also requires diagrams dressed with gluons.

Azimuthal asymmetries for unpolarized targets (leading twist)

\[ \int d^2 q_T \frac{Q_T^2}{MM_h} \cos(2\phi_h^Q) \frac{d\sigma_{LT}}{d x_B \ dy \ dz_h \ d^2 q_T} = \frac{4\pi\alpha^2 \ s}{Q^4} \ 4(1 - y) \sum_{a, \bar{a}} e_a^2 x_B h_{1}^{(1) a}(x_B) H_{1}^{(1) a} \]  

(32)
4 Concluding remarks

In the previous section some results for 1-particle inclusive lepton-hadron scattering have been presented. Several other effects are important in these cross sections, such as target fragmentation, the inclusion of gluons in the calculation to obtain color-gauge invariant definitions of the correlation functions and an electromagnetically gauge invariant result at order $1/Q$ and finally QCD corrections which can be moved back and forth between hard and soft parts, leading to the scale dependence of the soft parts and the DGLAP equations.

In my talk I have tried to indicate why semi-inclusive, in particular 1-particle inclusive lepton-hadron scattering, can be important. The goal is the study of the quark and gluon structure of hadrons, emphasizing the dependence on transverse momenta of quarks. The reason why this prospect is promising is the existence of a field theoretical framework that allows a clean study involving well-defined hadronic matrix elements.

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References