1. INTRODUCTION

To access the transverse distribution in two-particle seminclusive production, we propose to measure the production of a pair of particles in a forward-backward jet. In the forward region, the two-particle production function is measured in the presence of both particles, while in the backward region, the production function is measured in the presence of only one particle. This allows us to determine the differential distribution of the two-particle production function.

We consider the option of extracting the transverse distribution of two-particle production functions.
two unpolarized hadrons at leading twist. We will recover the results originally presented in Ref. [3]. We will devote particular attention to the connection with the helicity basis formalism (see, e.g., Refs. [12, 21]) and for the first time we will deduce positivity bounds on IFF.

In Sec. III, the whole problem is reconsidered by expanding in partial waves the two-hadron system in its center-of-mass frame. If we consider only low invariant masses, the expansion can be truncated to include the first two terms only, as hadron pairs are produced mainly in the s-wave channel or in the p-wave channel (via a spin-1 resonance). We can thus deduce a general unifying formalism that naturally incorporates the specific case of Ref. [2], in the subsector describing the interference between relative s and p waves, as well as the case of spin-1 hadron fragmentation [13], in the subsector of the relative p wave. In particular, we will identify a SSA where the transversity distribution appears in connection with a s-p IFF, and a SSA where the transversity is connected to a pure p-wave IFF. These two asymmetries are completely distinct, they could have different physical origins and different magnitudes.

In Sec. IV we complete our analysis by including the intrinsic partonic transverse momentum and $\vec{k}_T$-unintegrated fragmentation functions. Also in this case, in Sec. V we will present positivity bounds and will carry out the partial wave expansion. The results for the complete cross section for all combinations of beam and target polarizations are listed in the appendices. Finally, some conclusions are drawn in Sec. VI.

![Diagram](image)

**Fig. 1:** The usual quark handbag diagram contributing at leading twist to the semiinclusive DIS into two leading hadrons: a) hadron and parton momenta are shown, in particular the total momentum $P_h = P_1 + P_2$ and relative momentum $R = (P_1 - P_2)/2$ of the two-hadron system; b) target helicity, parton chirality and two-hadron partial wave indices are shown.

### II. TWO-PARTICLE INCLUSIVE DEEP INELASTIC SCATTERING

In the following, we will describe the kinematics and the details of the semi-inclusive production of two unpolarized hadrons in the context of the SIDIS process. However, we point out that the involved fragmentation functions can be used also in the case of reactions with a hadronic probe or in $e^+ e^-$ annihilation [14, 15].

#### A. Kinematics and hadronic tensor

The process is schematically depicted in Fig. 1. An electron with momentum $l$ scatters off a target nucleon with mass $M$, polarization $S$ and momentum $P$, via the exchange of a virtual hard photon with momentum $q = l - l'$ ($q^2 = -Q^2$). Inside the target, the photon hits a quark with momentum $p$, changing its momentum to $k = p + q$. The quark then fragments into a residual jet and two leading unpolarized hadrons with masses $M_1, M_2$, and momenta $P_1$ and $P_2$. We introduce the vectors $P_h = P_1 + P_2$ and $R = (P_1 - P_2)/2$. We describe a 4-vector $a$ as $[a^{-}, a^{+}, \vec{a}_T]$, i.e. in terms of its light-cone components $a^{\pm} = (a^{0} \pm a^{3})/\sqrt{2}$ and the bidimensional vector $\vec{a}_T$. It is convenient to choose the $z$ axis according to the condition $P_{T} = P_h T = 0$. In this case, the virtual photon has a nonvanishing transverse momentum $\vec{q}_T$. However, it is also customary to align the $z$ axis opposite to the direction of the virtual photon, in which case the outgoing hadron has a nonvanishing transverse momentum $P_{hT} = -z\vec{q}_T$. These two directions overlap up to corrections of order $1/Q$, which we will systematically neglect in the following. The $y$ axis is chosen to point in the direction of the vector product $(-\vec{q} \times \vec{P})$ [22].
We define the variables \( x = p^+ / P^+ \), which represents the light-cone fraction of target momentum carried by the initial quark, and \( z = P_h^- / k^- \), the light-cone fraction of fragmenting quark momentum carried by the final hadron pair. Analogously, we define the light-cone fraction \( \zeta = 2 \vec{R}^- / P_h^- \), which describes how the total momentum of the hadron pair is split into the two single hadrons.\(^2\) The relevant momenta can be parametrized as

\[
\begin{align*}
P^\mu &= \begin{bmatrix} M_2^2 2P^+ \\ P^+ \end{bmatrix}, \\
P^\mu &= \begin{bmatrix} \vec{P}_T^2 + \frac{P_T^2}{2} \\ 2x P^+ \end{bmatrix}, \\
k^\mu &= \begin{bmatrix} P_T^- \\ \frac{z(k^2 + \vec{k}_T^2)}{2P_h^-} \end{bmatrix}, \\
P^\mu_h &= \begin{bmatrix} P_T^- \\ M_h^2 2P_h^- \\ 0 \end{bmatrix}, \\
R^\mu &= \begin{bmatrix} \frac{\zeta}{2} P_T^- \\ (M_h^2 - M_1^2) - \frac{\zeta}{2} M_h^2 \\ 2P_T^- \end{bmatrix}. \quad (1)
\end{align*}
\]

Not all components of the 4-vectors are independent. In particular, here we observe that

\[
\begin{align*}
R^2 &= \frac{M_2^2 + M_2^2}{2} - \frac{M_h^2}{4}, \\
R_T^2 &= \frac{1}{2} \left[ \frac{(1 - \zeta)(1 + \zeta)}{2} M_h^2 - (1 - \zeta) M_1^2 - (1 + \zeta) M_2^2 \right], \\
P_h \cdot R &= \frac{M_h^2 - M_1^2}{2}, \\
P_h \cdot k &= \frac{M_h^2}{2z} + \frac{z^2 + [\vec{k}_T]^2}{2}, \\
R \cdot k &= \frac{(M_2^2 - M_1^2) - \frac{\zeta}{2} M_h^2}{2z} + \frac{z \zeta [k^2 + [\vec{k}_T]^2]}{4} - \vec{k}_T \cdot \vec{R}_T. \quad (2)
\end{align*}
\]

The positivity requirement \( R_T^2 \geq 0 \) imposes the further constraint

\[
M_h^2 \geq \frac{2}{1 + \zeta} M_1^2 + \frac{2}{1 - \zeta} M_2^2. \quad (3)
\]

We shall first consider the case when the cross section is integrated over the transverse momentum of the virtual photon, \( \vec{q}_T \), postponing the analysis of the complete case in Sec. IV. Until then, no transverse-momentum dependent distribution and fragmentation functions will appear. The seven-fold differential equation for two-particle-inclusive DIS is

\[
\frac{d^7 \sigma}{d\zeta dM_h^2 d\phi_R dx dy d\phi_S} = \sum_a \frac{\alpha^2 y e^2_a}{32 \zeta Q^2} L_{\mu \nu} 2M W_\alpha^{\mu \nu}, \quad (4)
\]

where \( L_{\mu \nu} \) is the lepton tensor; \( y = (E - E') / E \) is the fraction of beam energy transferred to the hadronic system and it is related to the lepton scattering angle in the center-of-mass (cm) frame; \( \phi_R \) and \( \phi_S \) are the azimuthal angles of \( \vec{R}_T \) and \( \vec{S}_T \) with respect to the lepton scattering plane. At tree level, the leptonic tensor for a flavour \( a \) is given by

\[
2M W_\alpha^{\mu \nu} = 32 \zeta \text{Tr} \left[ \Phi_a(x, S) \gamma^\mu \Delta_a(z, \zeta, M_h^2, \phi_R) \gamma^\nu \right] + \left( \begin{array}{cc} q & q \\ \mu & \nu \end{array} \right), \quad (5)
\]

where

\[
\Phi_a(x, S) = \int d\vec{P}_T dP^- \Phi_a(p, P, S) \bigg|_{p^+ = xP^+},
\]

\[
\Delta_a(z, \zeta, M_h^2, \phi_R) = \left( \frac{z}{32} \right) \int d\vec{k}_T d(k^+) \Delta_a(k; P_h, R) \bigg|_{k^- = P_h^-, k^+ = xP^+}. \quad (7)
\]

\(^2\) Note that \(-1 \leq \zeta \leq 1\), and \( \zeta = 2\xi - 1\), with \( \xi \) defined in Ref. [3].
The quark-quark correlator $\Phi$ describes the nonperturbative processes determining the distribution of parton $a$ inside the spin-1/2 target (represented by the lower shaded blob in Fig. 1) and, similarly, the correlator $\Delta$ symbolizes the fragmentation of quark $a$ producing two tagged leading hadrons in a residual jet (upper shaded blob in Fig. 1).

![Diagram of quark-quark correlator](image)

**FIG. 2:** Kinematics for the SIDIS of the lepton $l$ on a transversely polarized target leading to two hadrons inside the same current jet.

We are going to focus only on the leading twist contributions to the hadronic tensor of Eq. (5). A method to extract these contributions consists in projecting the so-called good light-cone components out of the quark field $\psi$. As it is evident from the kinematics in the infinite momentum frame, the $+$ and the $-$ light-cone components are the dominant ones for the parton entering and exiting the hard vertex, respectively. They can be projected out by means of the operators $\bar{P}_\pm = \frac{1}{2} \gamma^{\pm} \gamma^\mu$. Any other component of $\psi$ is automatically of higher twist. Therefore, the hadronic tensor (5) at leading twist looks like

$$2 M W^{\mu\nu}_a = 32 \left[ \text{Tr} \left[ \bar{P}_+ \Phi_a(x, S) \bar{P}_- \Gamma_m \Delta(z, \zeta, M^2, \phi_R) \right] \right],$$

where $\bar{P}_\pm \equiv \gamma^0 P_\pm \gamma^0$. In the last step the Dirac indices have been explicitly indicated. In the following, we will analyze each contribution to Eq. (8) separately.

### B. The quark-quark correlator $\Phi$

The leading-twist projection of the quark-quark correlator $\Phi$ can be parametrized in terms of the well known distribution functions [23, 24]

$$\begin{align*}
\bar{P}_+ \Phi_a(x, S) \gamma^+ &= \left( f_1(x) + \lambda g_1(x) \gamma_5 + 2 h_1(x) \gamma_5 \gamma_T \right) P_+ \\
&= \begin{pmatrix}
 f_1^+ + \lambda g_1^+ & 0 & 0 & (S_x - i S_y) h_1^+ \\
 0 & 0 & 0 & 0 \\
 (S_x + i S_y) h_1^+ & 0 & f_1^- - \lambda g_1^- 
\end{pmatrix},
\end{align*}$$

where $\lambda = M S^+ / P^+$ and $\gamma_T = (S_x, S_y)$ are the light-cone helicity and transverse components of the target spin, respectively ($\bar{P}_+ \Phi$ corresponds to the $p_T$-integrated parametrization of Eq. (2) in Ref. [6]). It is possible to rewrite

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3 Other common notations are $f_1(x) = a(x)$, $g_1(x) = \Delta a(x)$, $h_1(x) = \delta a(x), \Delta_T a(x)$ [4].
Eq. (9) in a more compact notation, namely in the chiral basis of the good quark fields \( \psi_{ \pm R/L} = \mathcal{P}_{ \pm} \mathcal{P}_{R/L} \psi \), with \( \mathcal{P}_{R/L} = (1 \pm i \gamma_5)/2 \) [24],

\[
[\mathcal{P}_{ \pm} \Phi_\alpha(x, S) \gamma^\mu]_{\chi_1 \chi_1} = \begin{pmatrix}
 f_\alpha^i(x) + \lambda g_\alpha^i(x) & (S x - i S_y) h_\alpha^i(x) \\
 (S x + i S_y) h_\alpha^i(x) & f_\alpha^i(x) - \lambda g_\alpha^i(x)
\end{pmatrix},
\]

Finally, it is useful to project out also the target helicity density matrix \( \rho_{ \Lambda \Lambda'} \), by

\[
[\mathcal{P}_{ \pm} \Phi_{ \alpha} \gamma^\mu]_{\chi_1 \chi_1} = \rho_{ \Lambda \Lambda'} [\mathcal{P}_{ \pm} \Phi_{ \alpha} \gamma^\mu]_{\chi_1 \chi_1}^\Lambda \Lambda',
\]

with

\[
\rho_{ \Lambda \Lambda'} = \frac{1}{2} \begin{pmatrix}
 1 + \lambda & S x - i S_y \\
 S x + i S_y & 1 - \lambda
\end{pmatrix},
\]

\[
[\mathcal{P}_{ \pm} \Phi_{ \alpha} \gamma^\mu]_{\chi_1 \chi_1}^\Lambda \Lambda' = \begin{pmatrix}
 f_\alpha^i + g_\alpha^i & 0 & 0 & 0 \\
 0 & f_\alpha^i - g_\alpha^i & 0 & 0 \\
 0 & 0 & 2 h_\alpha^i & 0 \\
 0 & 0 & 0 & f_\alpha^i + g_\alpha^i
\end{pmatrix}.
\]

In Eq. (13) the pair of indices \((\Lambda, \Lambda')\) identifies each component of the \(2 \times 2\) submatrices and indicates the spin state of the target; they are attached to each corresponding nucleon leg in the diagram of Fig. 1b. The pair \((\chi_1, \chi_1')\) identifies each submatrix and indicates the parton chirality; they are attached to the emerging quark legs in Fig. 1b. Equation (13) satisfies general requirements, such as the angular momentum conservation \((\chi_1 + \Lambda = \chi_1' + \Lambda')\), hermiticity and parity invariance. The chiral transposed matrix is also positive semidefinite, from which the well known Soffer bound [25], among others, is obtained:

\[
\begin{align*}
 f_\alpha^i(x) & \geq 0 \\
 |h_\alpha^i(x)| & \leq \frac{1}{2} \sqrt{f_\alpha^i(x) + g_\alpha^i(x)}.
\end{align*}
\]

C. The quark-quark correlator \( \Delta \) and positivity bounds

The most general parametrization of the quark-quark correlator \( \Delta(k, P_h, R) \) entering Eq. (7), compatible with hermiticity and parity invariance, is [3]

\[
\Delta(k, P_h, R) = M_B C_1 I + C_2 T_h + C_3 R + C_4 k
\]

\[
+ \frac{C_5}{M_B} \sigma_{\mu \nu} P_h^\mu k^\nu + \frac{C_6}{M_B} \sigma_{\mu \nu} R^\nu k^\mu + \frac{C_7}{M_B} \sigma_{\mu \nu} P_h^\mu R^\nu
\]

\[
+ \frac{C_8}{M_B} \gamma_\mu \gamma_\nu \sigma_{\mu \nu} \gamma_5 P_h \cdot R k \sigma,
\]

where the amplitudes \( C_i(k^3, k \cdot P_h, k \cdot R, R^2) \) are dimensionless real scalar functions. By using Eqs. (15, 7) the leading-twist projection becomes

\[
\mathcal{P}_{-} \Delta_\alpha(z, \zeta, M_B^2, \phi R) \gamma^\mu = \frac{1}{8 \pi} \left( D_1^\mu(z, \zeta, M_B^2) + i H_1^\mu a(z, \zeta, M_B^2) \frac{B_P}{M_B \phi R} \right) \mathcal{P}_{-} \mathcal{P}_{-}
\]

\[
= \frac{1}{8 \pi} \begin{pmatrix}
 0 & 0 & 0 & 0 \\
 0 & D_1^\mu & 0 & 0 \\
 0 & 0 & 0 & 0 \\
 0 & 0 & 0 & 0
\end{pmatrix},
\]

where

\[
D_1(z, \zeta, M_B^2) = \frac{z \pi}{4} \int d^2 \vec{k}_T d^2 k d^2 (2 k \cdot P_h) \delta \left( \frac{\vec{k}_T^2 + k^2}{z^2} + \frac{M_B^2}{z^2} - \frac{2 k \cdot P_h}{z} \right) \left[ C_2 + \frac{\zeta}{2} C_3 + \frac{1}{z} C_4 \right]
\]

\[
H_1^\mu(z, \zeta, M_B^2) = \frac{z \pi}{4} \int d^2 \vec{k}_T d^2 k d^2 (2 k \cdot P_h) \delta \left( \frac{\vec{k}_T^2 + k^2}{z^2} + \frac{M_B^2}{z^2} - \frac{2 k \cdot P_h}{z} \right) \left[ \frac{1}{z} C_6 - C_7 \right]
\]
The prefactors have been chosen to have a better connection with the one-hadron results, i.e. after integrating over $\zeta, M_b^2$ and $\phi_R$. In Eq. (16), $P_\perp \Delta$ corresponds to the parametrization of Eq. (3) in Ref. [6].

The fragmentation function $H_1^\perp$ is chiral odd and represents a possible partner to isolate the transverse distribution inside the cross section at leading twist [6]. Moreover, it is also odd with respect to naive time-reversal transformations (for brevity, T-odd) [3]. Noteworthy, it is the only example of leading-twist T-odd function surviving the integration upon the quark transverse momentum $k_T$. It would be interesting to investigate it in order to understand what is the relevance of the transverse-momentum dependence in generating T-odd effects [26, 27]. As a consequence of the $k_T$-integrated $H_1^\perp$ could have simpler evolution equations than the ones of the Collins function. Since $H_1^\perp$ has the same operator structure as the transversity, it has been suggested that it could have the same evolution equations [28, 29, 30]. However, the situation is complicated by the presence of two hadrons [31, 32, 33] except for the component of $H_1^\perp$ describing the production of a spin-1 resonance (see Sec. III B).

Again, in the chiral basis for the good light-cone components Eq. (16) is simplified to

$$\left[ P_\perp \Delta \bar{e}(z; \zeta, M_b^2, \phi_R) \gamma_{-}^{-} |_{\chi_2 \chi_3} = \frac{1}{8\pi} \begin{pmatrix} \frac{D^2_1(z; \zeta, M_b^2)}{M_b^2} + i e^{i \phi_R} \frac{R_1^2}{M_b} H_1^\perp(z; \zeta, M_b^2) \\ -i e^{-i \phi_R} \frac{R_1^2}{M_b} H_1^\perp(z; \zeta, M_b^2) \\ \frac{D^2_1(z; \zeta, M_b^2)}{M_b^2} \end{pmatrix} \right],$$

(19)

where $(\chi_2, \chi_3)$ are the quark chiralities to be attached to the parton legs entering the $\Delta$ blob in Fig. 1b.

The matrix in Eq. (19) is positive-semidefinite, from which the following bounds can be derived:

$$D_1^2(z; \zeta, M_b^2) \geq 0 \quad D_1^2(z; \zeta, M_b^2) \geq \frac{|R_1^2|}{M_b} |H_1^\perp(z; \zeta, M_b^2)|.$$  

(20)

### D. Cross section and transverse spin asymmetry

Using the previous results, we can now rewrite the leading-twist cross section for unpolarized two-hadron SIDIS in the helicity basis. In fact, after inserting Eqs. (11) and (19) inside Eq. (8), the cross section in Eq. (4) becomes

$$\frac{d^7\sigma}{d\zeta \, dM_b^2 \, d\phi_R \, dz \, dx \, dy \, d\phi_S} = \sum_a |\rho_{AA}(S)| \left[ P_+ \Phi_d(x) \gamma^+ \lambda^A_{\chi_1} \right] \left( \frac{d\sigma^{\perp q}}{dy} \right)^{\chi_1 \chi_2 \chi_3} \left[ P_\perp \Delta \bar{e}(z; \zeta, M_b^2, \phi_R) \gamma_{-}^{-} \right]^{\chi_2 \chi_3}$$

(21)

where

$$\frac{d\sigma^{\perp q}}{dy}^{\chi_1 \chi_2 \chi_3} = \frac{2e^2_+ \alpha^2 \gamma^2 L_{\mu \nu} \left( \frac{\gamma^- \gamma^\mu}{2} P_+ - \frac{\gamma^+ \gamma^\nu}{2} P_- \right)^\chi_1 \chi_2 \chi_3}$$

(22)

represents the elementary electron-quark scattering. Strictly speaking, this is not a scattering matrix, but a scattering amplitude times the conjugate of a different scattering amplitude [12]. However, for conciseness we follow the notation of Ref. [2]. The polarization of the incident beam is indicated with $\lambda_x$ and

$$A(y) = 1 - y + \frac{y^2}{2}, \quad B(y) = 1 - y, \quad C(y) = y(2 - y).$$

(23)

In Eq. (22), the indices $(\chi_1, \chi_2)$ refer to the chiralities of the entering quarks and identify each submatrix, while $(\chi_2, \chi_3)$ refer to the exiting quarks and point to the elements inside each submatrix. By expanding the sum over repeated indices in Eq. (21), we get the expression

$$\frac{d^7\sigma}{d\zeta \, dM_b^2 \, d\phi_R \, dz \, dx \, dy \, d\phi_S} = \sum_a e^2_+ \frac{2e^2 \alpha^2}{4 \pi Q^2 y} \left\{ A_+(y) f^e(x) \frac{D^2_1(z; \zeta, M_b^2)}{M_b^2} + \lambda_x C(y) \frac{C(y)}{2} y f^e(x) \frac{D^2_1(z; \zeta, M_b^2)}{M_b^2} \right\} + B(y) \frac{\sqrt{\mathcal{S}_T}}{M_b} \sin(\phi_R + \phi_S) h_1^\perp(z; \zeta, M_b^2).$$

(24)

We thank D. Boer for pointing out this detail.
For an unpolarized beam ($\lambda_e = 0$, indicated with $O$) and a transversely polarized target ($\lambda = 0$, indicated with $T$), Eq. (24) corresponds to Eq. (10) of Ref. [6] after integrating over all transverse momenta. The following SSA can be built:

$$A_{OT}^{\sin(\phi_R + \phi_S)}(y, x, z, M_h^2) = \frac{\int d\phi_S \, d\phi_R \, d\zeta \, \sin(\phi_R + \phi_S) \, d^7\sigma_{OT}}{\int d\phi_S \, d\phi_R \, d\zeta \, d^7\sigma_{OT}}$$

$$= \left[ \frac{S_T}{A(y)} \right] \frac{\sum_s e_s^2 \, h_1^s(x) \int d\zeta \, \frac{1}{M_h^2} \, H_1^{s\ast}(z, \zeta, M_h^2)}{\sum_s e_s^2 \, f_s^2(x) \int d\zeta \, D_1^s(z, \zeta, M_h^2)} ,$$

which allows to isolate the transversity $h_1$ at leading twist. Apart from the usual variables $x, y, z$, the only other variable to be measured is the angle $\phi_R + \phi_S$. Instead of using the scattering plane as a reference to measure azimuthal angles, it is sometimes convenient to use the directions of the beam and of the transverse component of the target spin. The new plane is rotated by the angle $\phi_S = -\phi_t$ with respect to the scattering plane; therefore, we have

$$\phi_R \equiv \phi_R - \phi_t$$

and

$$\phi_S \equiv \phi_t + 2\phi_t [6].$$

The asymmetry described in Eq. (25) is the most general one at leading twist for the case of two-hadron production when an unpolarized lepton beam scatters off a transversely polarized target. No assumptions are made on the behavior of the fragmentation functions. However, as we shall see in the next Section, it is useful and desirable to understand how different partial waves contribute to the above fragmentation functions.

### III. PARTIAL-WAVE EXPANSION FOR THE TWO-HADRON SYSTEM

If the invariant mass $M_h$ of the two hadrons is not very large, the pair can be assumed to be produced mainly in the relative $s$-wave channel, with a typical smooth distribution, or in the $p$-wave channel with a Breit-Wigner profile [32]. Therefore, it is useful to expand Eq. (18) — or equivalently Eq. (19) — in relative partial waves keeping only the first two harmonics. To this purpose, in the following we reformulate the kinematics in the cm frame of the two-hadron system. Then, the leading-twist projection for the quark-quark correlator $\Delta$ is conveniently expanded deducing a more detailed structure than Eq. (19). A set of new bounds is derived and the corresponding expression for the cross section is discussed.

![Fig. 3](image_url)

**Fig. 3:** The hadron pair in the cm frame; $\theta$ is the cm polar angle of the pair with respect to the direction of $P_h$ in the target rest frame.

In the cm frame the emission of the two hadrons occurs back-to-back. The direction identified by this emission forms an angle $\theta$ with the direction of $P_h$ in the target rest frame (see Fig. 3). In this frame, the relevant variables become

$$P_h^\mu = \left[ \frac{M_h}{\sqrt{2}}, \frac{M_h}{\sqrt{2}}, 0, 0 \right]$$

$$R^\mu = \left[ \sqrt{M_p^2 + |\vec{R}|^2} - \sqrt{M_p^2 + |\vec{R}|^2 - 2|\vec{R}| \cos \theta}, \sqrt{M_p^2 + |\vec{R}|^2} - \sqrt{M_p^2 + |\vec{R}|^2 + 2|\vec{R}| \cos \theta}, \sqrt{M_p^2 + |\vec{R}|^2} - \sqrt{M_p^2 + |\vec{R}|^2 - 2|\vec{R}| \cos \theta} \right],$$

$$\zeta = \frac{2R^\perp}{P_h^\perp} = \frac{1}{M_h} \left( \sqrt{M_p^2 + |\vec{R}|^2} - \sqrt{M_p^2 + |\vec{R}|^2 - 2|\vec{R}| \cos \theta} \right) ,$$

(26)
where
\[ |\vec{R}| = \frac{1}{2M_h} \sqrt{M_h^2 - 2(M_1^2 + M_2^2) + (M_1^2 - M_2^2)^2}. \]  

(27)

The crucial remark is that in this frame \( \zeta \) depends linearly on \( \cos \theta \), i.e. \( \zeta = a + b \cos \theta \), with \( a, b \), functions only of the invariant mass. This suggests that any function of \( \zeta \) can be conveniently expanded in the basis of Legendre polynomials in \( \cos \theta \), as discussed in the following.

A. Partial-wave expansion of the quark-quark correlator\( \Delta \) and positivity bounds

We first express the leading-twist quark-quark correlator \((16)\) in terms of the cm variables. The connection between the two representations is defined as

\[ \Delta(z, \cos \theta, M_h^2, \phi_R) \equiv \frac{2|M_h|}{M_h} \Delta(z, \zeta, M_h^2, \phi_R). \]  

(28)

to take into account the Jacobian of the transformation, \( \left. d\zeta = \frac{2|M_h|}{M_h} d\cos \theta. \right) \) Therefore

\[ \mathcal{P}_- \Delta(z, \cos \theta, M_h^2, \phi_R) = \frac{2|M_h|}{8\pi M_h} \left( D_1(z, \zeta(\cos \theta), M_h^2) + iH_1^{\delta}(z, \zeta(\cos \theta), M_h^2) \frac{|\vec{R}|}{M_h} \sin \theta \phi_{\phi_R} \right) \mathcal{P}_-, \]  

(29)

where \( \phi_{\phi_R} = \left[ 0, 0, \cos \phi_R, \sin \phi_R \right]. \)

The fragmentation functions can be expanded in Legendre polynomials as

\[ \frac{2|M_h|}{M_h} D_1(z, \zeta(\cos \theta), M_h^2) = \sum_n D_{1,n}(z, M_h^2) P_n(\cos \theta) \]  

\[ \frac{2|M_h|}{M_h} H_1^{\delta}(z, \zeta(\cos \theta), M_h^2) = \sum_n H_{1,n}^{\delta}(z, M_h^2) P_n(\cos \theta) \]  

(30)

with

\[ D_{1,n}(z, M_h^2) = \int_{-1}^{1} d\cos \theta \ P_n(\cos \theta) \frac{2|M_h|}{M_h} D_1(z, \zeta(\cos \theta), M_h^2) \]  

\[ H_{1,n}^{\delta}(z, M_h^2) = \int_{-1}^{1} d\cos \theta \ P_n(\cos \theta) \frac{2|M_h|}{M_h} H_1^{\delta}(z, \zeta(\cos \theta), M_h^2). \]  

(31)

We can truncate the expansion to the first three terms only \((n \leq 2)\), which are the minimal set required to describe all the “polarization” states of the system in the cm frame for relative partial waves with \( L = 0, 1 \). In fact, for \( n = 0 \) \((P_0 = 1)\) the corresponding term in the correlator does not depend on \( \theta \), it is “unpolarized”. For \( n = 1 \), a term linear in \( \cos \theta \) \((P_1 = \cos \theta)\) describes the interference between an “unpolarized” hadron pair in \( s\)-wave, for example on the left hand side of Fig. 1b, and a “longitudinally polarized” pair in \( p\)-wave on the right hand side. Whenever in the correlator we encounter a term linear in \( \sin \theta \), we will interpret it as the interference between a “unpolarized” pair in \( s\)-wave and a “transversely polarized” pair in \( p\)-wave. Similarly, a term proportional to \( \sin \theta \cos \theta \) indicates the interference between “longitudinally” and “transversely polarized” pairs always in a relative \( p\)-wave. The last case corresponds to \( n = 2 \), that is interpreted as a “tensor polarization” still related to the interference between pairs in a relative \( p\)-wave. With notations that are consistent with previous arguments, the correlator \((29)\) is expanded as

\[ \mathcal{P}_- \Delta(z, \zeta(\cos \theta), M_h^2, \phi_R) \sim \frac{1}{8\pi} \left[ D_{1,0}(z, M_h^2) + D_{1,1}(z, M_h^2) \cos \theta + D_{1,2}(z, M_h^2) \frac{1}{2} (3\cos^2 \theta - 1) \right. \]  

\[ + i \left( H_{1,0}^{\gamma}(z, M_h^2) + H_{1,1}^{\gamma}(z, M_h^2) \cos \theta \right) \sin \theta \frac{1}{M_h} \phi_{\phi_R} \]  

\[ \equiv \frac{1}{8\pi} \left[ D_{1,0}(z, M_h^2) + D_{1,1}(z, M_h^2) \cos \theta + D_{1,2}(z, M_h^2) \frac{1}{4} (3\cos^2 \theta - 1) \right. \]  

\[ + i \left( H_{1,0}^{\gamma}(z, M_h^2) + H_{1,1}^{\gamma}(z, M_h^2) \cos \theta \right) \sin \theta \frac{1}{M_h} \phi_{\phi_R} \]  

\[ \mathcal{P}_- . \]  

(32)
Consequently, the same correlator in the chiral basis becomes

\[
\begin{pmatrix}
  D_1^{\text{OO}}(z, M_0^2) + D_1^{\text{OL}}(z, M_0^2) \cos \theta
  + D_1^{\text{LL}}(z, M_0^2) \frac{1}{2}(3 \cos^2 \theta - 1)
  - i e^{-i \phi n} \frac{\omega_0}{M_0} \sin \theta \\
  \times (H_{1,\text{OT}}^0(z, M_0^2) + H_{1,\text{LT}}^0(z, M_0^2) \cos \theta)
  \\
  \times \left( H_{1,\text{OT}}^0(z, M_0^2) + H_{1,\text{LT}}^0(z, M_0^2) \right) \frac{1}{2}(3 \cos^2 \theta - 1)
\end{pmatrix}
\]  

It is useful to project out of Eq. (33) the information about the orbital angular momentum of the system, which is encoded in the angular distribution of the hadron pair. In fact, for \( L \leq 1 \) the decay matrix for the hadron pair is given by the following bilinear combination of spherical harmonics:

\[
D^{LL}_{MM}(\theta, \phi_R) = Y_{LM} Y_{LM}^\ast \equiv 
\begin{pmatrix}
  1 & -\sqrt{3} \sin \theta \cos \phi n & \sqrt{3} \cos \theta \\
  -\sqrt{3} \sin \theta \cos \phi n & -\frac{3}{2} \sin^2 \theta & -\frac{3}{2} \cos \theta \sin \theta \cos \phi n \\
  \sqrt{3} \cos \theta & -\frac{3}{2} \cos \theta \sin \theta \cos \phi n & \frac{3}{2} \cos \theta \cos \phi n
\end{pmatrix}
\]  

with \( L, L' \leq 1 \) and \([M^{(i)}] \leq L^{(i)}\). The upper left block corresponds to \( L = L' = 0 \), i.e. to the system being in relative \( s \) wave. The lower right block instead corresponds to \( L = L' = 1 \), i.e. to the system being in relative \( p \) wave, including all the contributions corresponding to different \( M, M' \) projections and their interferences. The off-diagonal blocks indicate, obviously, the interference between the \( s \) and \( p \) waves. Using the decay matrix, it is possible to represent the fragmentation in the basis of the quark chirality and of the pair orbital angular momentum. In fact, Eq. (35) can be written as

\[
[P_{\pm} \Delta(z, \zeta, M_0^2, \phi_R) \gamma^-]_{\chi_2/\chi_2} = [P_{\pm} \Delta(z, M_0^2) \gamma^-]_{M_0^2}^{LL} \chi_2 \chi_2 \mathcal{D}^{LL}_{MM}(\theta, \phi_R),
\]  

where

\[
[P_{\pm} \Delta(z, M_0^2) \gamma^-]_{M_0^2}^{LL} \chi_2 \chi_2 = \frac{1}{8} \begin{pmatrix}
  A_{MM}^{LL} & B_{MM}^{LL} \\
  B_{MM}^{LL} & A_{MM}^{LL}
\end{pmatrix}
\]

and

\[
A_{MM}^{LL} = \begin{pmatrix}
  D_1^{\text{OO}} & 0 & \frac{2}{\sqrt{3}} D_1^{\text{OL}} \\
  0 & D_1^{\text{OO}} - \frac{2}{3} D_1^{\text{LL}} & 0 \\
  \frac{2}{\sqrt{3}} D_1^{\text{OL}} & 0 & D_1^{\text{OO}} + \frac{2}{3} D_1^{\text{LL}}
\end{pmatrix},
\]

and \( B_{MM}^{LL} \) is given by

\[
B_{MM}^{LL} = \begin{pmatrix}
  0 & 0 & H_{1,\text{OT}}^0 \\
  -\frac{2\sqrt{3}}{\sqrt{3} M_0} H_{1,\text{OT}}^0 & 0 & 0 \\
  0 & 0 & H_{1,\text{LT}}^0
\end{pmatrix}.
\]

The fragmentation matrix \([P_{\pm} \Delta(z, M_0^2) \gamma^-]_{M_0^2}^{LL} \chi_2 \chi_2\) fulfills all the fundamental properties, namely Hermiticity, parity invariance \( [33] \) and angular momentum conservation \( (\chi_2 + M = \chi_2 + M') \). The imaginary entries in its off-diagonal submatrix are T-odd fragmentation functions. It is worth noticing that with the projection (35) we gained a further
information on the “unpolarized” state of the hadron pair. In fact, we see from the diagonal of Eq. (37) that the spherically symmetric state in the cm frame receives contributions from both the relative s and p waves, such that when performing the matrix multiplication of Eq. (35) we get

\[
D_{1,00}(z, M_b^2) = \frac{1}{4} D_{1,00}^s(z, M_b^2) + \frac{3}{4} D_{1,00}^p(z, M_b^2) .
\]  

(39)

However, in an actual cross section the two contributions are merged together and are kinematically indistinguishable, unless a specific hypothesis on the dependence upon the invariant mass \( M_b \) is assumed for the two different partial waves, e.g. a resonant contribution for the p wave and a continuum background for the s wave.

Finally, from the matrix (36) being positive semidefinite the following bounds are derived:\(^5\)

\[
D_{1,00}^s \geq 0, \quad D_{1,00}^p \geq 0,
\]

\[
-\frac{3}{2} D_{1,00}^s \leq D_{1,LL} \leq 3 D_{1,00}^s,
\]

\[
D_{1,0L} \leq \sqrt{\frac{3}{4} D_{1,00}^s \left( D_{1,00}^p + \frac{2}{3} D_{1,LL} \right) - \frac{3}{2} D_{1,00}},
\]

\[
\frac{|\vec{R}|}{M_h} H_{1,OT}^s \leq \sqrt{\frac{3}{4} D_{1,00}^s \left( D_{1,00}^p - \frac{1}{3} D_{1,LL} \right) - \frac{3}{2} D_{1,00}},
\]

\[
\frac{|\vec{R}|}{M_h} H_{1,LT}^s \leq \frac{3}{2 \sqrt{2}} \sqrt{\left( D_{1,00}^s + \frac{2}{3} D_{1,LL} \right) \left( D_{1,00}^p - \frac{1}{3} D_{1,LL} \right) - \frac{9}{8} D_{1,00}} .
\]  

(40)

B. Cross section and transverse spin asymmetries

Using Eq. (35) inside Eq. (21), we can take advantage of the full power of the analysis in the helicity formalism. In fact, the cross section can be expressed in the density matrices for the target helicity, for the chirality of the initial and fragmenting quark, and for the relative orbital angular momentum of the leading hadron pair [2]. Inserting the corresponding expressions (12,13,22,36,34), we get

\[
\frac{d^7 \sigma}{d \zeta d M_h^2 d \phi_R d z d x d y d \phi_S} = \sum_a \rho_{\Lambda,n}(S) \left[ \mathcal{P}_+ \Phi_a(s) \gamma^{+ \lambda \lambda_1} \chi_1 \chi_\lambda_2 \chi_\lambda \mathcal{P}_- \Delta(z, M_h^2) \gamma^- L M \chi_\lambda \chi_2 \chi_\lambda \right] D_{LL}^L M \left( \theta, \phi_R \right)
\]

\[
= \sum_a \epsilon_a^2 \frac{\alpha^2}{2 \pi Q^2} \left[ \mathcal{A}(y) f_1^R(x) + \lambda_{\Lambda} \lambda C(y) g(y) \right] \left[ D_{1,00}^s + 3 D_{1,00}^p \right] \left( \phi_R + \phi_S \right) h_1^s(x) \sin \theta \left[ H_{1,OT}^s + H_{1,LT}^s \cos \theta \right] ,
\]  

(41)

where all the fragmentation functions depend just on \((z, M_h^2)\).

Replacing \( \lambda = \lambda_{\Lambda} = |S_T| = 0 \) in the previous equation, we get the unpolarized cross section \(d^7 \sigma_{OO}\). However, it is particularly interesting to consider the case for an unpolarized beam and a transversely polarized target, i.e.

\[
\frac{d^7 \sigma_{OT}}{d \zeta d M_h^2 d \phi_R d z d x d y d \phi_S} = \sum_a \epsilon_a^2 \frac{\alpha^2}{2 \pi Q^2} B(y) \left| \vec{R} \right| \left( \phi_R + \phi_S \right) h_1^s(x) \sin \theta \left[ H_{1,OT}^s + H_{1,LT}^s \cos \theta \right] ,
\]  

(42)

because we can see that the transversity \(h_1\) can be matched by two different chiral-odd, T-odd IFF: one \((H_{1,OT}^s)\) pertaining to the interference between s and p-wave states of the hadron pair, the other \((H_{1,LT}^s)\) pertaining to the p wave only. The partial-wave analysis allows us for the first time to comprehend different theoretical analyses in a

\[\boxed{}\]

\(^5\) Note that the bounds involving the pure p-wave functions correspond to those obtained in Ref. [34]
unifying framework. In fact, $H_{1,OT}^q$ corresponds to the hypothesis first formulated in Ref. [2] and later reconsidered in Ref. [6], where the necessary spin asymmetry is generated by the interference between two channels describing two leading pions in the relative $s$ and $p$ waves, respectively. As a simple cross-check, taking Eq. (42) and integrating the $\theta$ dependence away yields

$$
\int_{-1}^{1} d \cos \theta \frac{d^7 \sigma_{OT}}{d \cos \theta \, d M^2_R \, d \phi_R \, d \phi_S \, dz \, dx \, dy} = \int_{-1}^{1} d \cos \theta \frac{d^7 \sigma_{OT}}{d \cos \theta \, d M^2_R \, d \phi_R \, d \phi_S \, dz \, dx \, dy}
$$

$$
= \sum_a \alpha^2 \frac{a^2}{4Q^2 y} B(y) \left[ \frac{S_T \| \mathbf{\vec{r}} \|}{M_h} \sin(\phi_R + \phi_S) h_1(x) H_{1,OT}^q(z, M_h^2) \right]. \quad (43)
$$

This asymmetry corresponds to the one studied in Ref. [2], although in that paper several assumptions were made. Firstly, the IFF was factorized in a part dependent only on the variable $z$, designated as $Q_T(z)$, and in a part containing the $M_h$-dependent $\pi-\pi$ phase shifts. Due to neglecting the scattering angle (see Fig. 2). The azimuthal angle of the hadron pair defined in Ref. [2] is $\phi = \pi/2 - \phi_R$. It is worth to note that the peculiar behavior in the invariant mass discussed in Ref. [2] relies on the assumption that only the $\pi-\pi$ rescattering can generate the T-odd character of the IFF. It has already been shown, however, that a different model with more general assumptions leads to a factorized dependence on the fragmentation function and to a completely different behaviour of the SSA [6]. Therefore, it is of great interest to experimentally explore the production of two unpolarized hadrons, e.g. two pions, in the relevant kinematic range, namely with an invariant mass around the $p$ resonance.

As for the function $H_{1,LT}^q$, it naturally links with the analysis developed in the case of a spin-1 hadron fragmentation [13], because the two spinless hadrons, e.g., two pions, can be considered as the decay product of a spin-1 resonance, e.g., a $\rho$ particle. The T-odd IFF arises from the interference between different channels on the relative $p$ wave. To the purpose of isolating an asymmetry containing the function $H_{1,LT}^<$, we show that integrating Eq. (42) upon $\theta$ in a different range, namely in the interval $[-\pi/2, \pi/2]$, yields

$$
\int_{-\frac{\pi}{2}}^{\frac{\pi}{2}} \frac{d \theta \sin \theta}{4Q^2 y} \frac{d^7 \sigma_{OT}}{d \cos \theta \, d M^2_R \, d \phi_R \, d \phi_S \, dz \, dx \, dy} = \int_{-\frac{\pi}{2}}^{\frac{\pi}{2}} d \theta \sin \theta \frac{d^7 \sigma_{OT}}{d \cos \theta \, d M^2_R \, d \phi_R \, d \phi_S \, dz \, dx \, dy \, d \phi_S}
$$

$$
= \sum_a \alpha^2 \frac{a^2}{4Q^2 y} B(y) \left[ \frac{S_T \| \mathbf{\vec{r}} \|}{M_h} \sin(\phi_R + \phi_S) h_1(x) \left[ H_{1,OT}^<(z, M_h^2) + \frac{A}{2\pi} H_{1,LT}^<(z, M_h^2) \right] \right], \quad (44)
$$

where both kinds of IFF appear at leading twist and can contribute to a SSA isolating the transversity $h_1$. Although spin-1 fragmentation functions have already been proposed in the past as possible chiral-odd partners for the transversity [10, 11, 12, 13], to our knowledge this is the first time that the asymmetry where they occur is explicitly identified and a clear distinction from the s-p interference is made.

There are not yet quantitative model predictions for $H_{1,LT}^<$. On the other hand, since the $p$-wave production of two hadrons becomes significant only when it proceeds via a spin-1 resonance, we can expect that the shape of this function in the invariant mass corresponds to a parton-Wigner curve peaked at the resonance mass. Moreover, it has the same features as a single-particle fragmentation function, unlike $H_{1,OT}^<$: its evolution equations can be expected to be analogous to that of the transversity [28, 29, 30]; it does not require a rescattering of the hadrons after they are produced and its physical origin could have something in common with the one of the Collins function. However, it should be noticed that in the case of the Collins function an essential role is played by the partonic transverse momentum, which in the case of $H_{1,LT}^<$ is replaced by the relative transverse momentum of the hadron pair.

It would be interesting to elaborate these topics since data for the electromagnetic $p$ production and decay are already available in the diffractive regime [35, 36, 37], and they could be available in the DIS regime as well in the near future.

IV. EXPLICIT DEPENDENCE ON THE TRANSVERSE MOMENTA

For sake of completeness, in this Section we extend the previous results to the case where the transverse momenta are not integrated away. In this case, the cross section is nine-fold and reads

$$
\frac{d^9 \sigma}{d \zeta \, d M^2_R \, d \phi_R \, d z \, d P_{h\perp} \, d x \, d y \, d \phi_S} = \sum_a \alpha^2 \frac{a^2}{32 \, Q^2 \, z \, \epsilon_a^2 \, \epsilon_z^2 \, L_{\mu \nu}} \, 2M \, W^{\mu \nu}. \quad (45)
$$
The hadronic tensor takes the form

\[ 2 M W_{\alpha}^{\mu \nu} = 32 \pi \mathcal{I} \left[ \text{Tr} \left[ \Phi_\alpha(x, p_T, S) \gamma^\mu \Delta_\alpha(z, k_T^z, \zeta, M^2 \omega \phi_R) \gamma^\nu \right] \right], \tag{46} \]

where we introduced the shorthand notation

\[ \mathcal{I}[f] \equiv \int d\vec{p}_T d\vec{k}_T \delta(\vec{p}_T - \vec{p}_T) [f]. \tag{47} \]

and where the transverse momentum dependent correlation functions are

\[ \Phi_\alpha(x, \vec{p}_T, S) = \int dp_T^+ \Phi_\alpha(p, P, S) \bigg|_{p^+ = xP^+}, \tag{48} \]

\[ \Delta_\alpha(z, \vec{k}_T^z, \zeta, M^2 \omega \phi_R) = \frac{1}{2z} \int dk^+ \Delta_\alpha(k; P, P) \bigg|_{k^+ = -P^+_\omega}. \tag{49} \]

The leading-twist projection of \( W^{\mu \nu} \) proceeds in an analogous way to Eq. (8); we usually have \cite{22}

\[ \mathcal{P}_+ \Phi_\alpha(x, \vec{p}_T, S) \gamma^+ = \left\{ \mathcal{F}_1^\alpha \left( x, \vec{p}_T^2 \right) + \frac{iG_F^\alpha \bar{N}_L}{M} f_1^\alpha(x, \vec{p}_T^2) \bigg| \gamma\gamma \bigg| \right. \]

\[ + \left. \frac{\lambda g_{1L}^2 \bar{N}_R}{M} \gamma^\mu \gamma^\nu \left[ \frac{\bar{p}_T \cdot \vec{k}_T}{M} \right] \gamma_5 \right\}, \tag{50} \]

where \( \eta^\mu = \epsilon^+ \bar{\eta}^\nu. \) Equation (50) corresponds to Eq. (2) of Ref. [6]. Again, similarly to Eq. (11) and following ones, we project out the density matrix of the target helicity so that Eq. (50) in the basis of quark chirality and target helicity becomes

\[ \left[ \mathcal{P}_+ \Phi_\alpha \gamma^+ \right]^\Lambda_{\Lambda} \gamma_{\chi_1 \chi_1} = \left( \begin{array}{cc}
\frac{\mathcal{F}_1^\alpha \gamma^+}{M} \left( \frac{\bar{p}_T + \gamma^5}{\gamma^\nu} \right) & \frac{1}{\bar{N}_L \gamma^\nu} e^{-2i\phi_p} \left( h_{1L}^\alpha + ih_{1L}^\alpha \right) \\
\frac{1}{\bar{N}_L \gamma^\nu} e^{-2i\phi_p} \left( h_{1L}^\alpha + ih_{1L}^\alpha \right) & -\frac{\mathcal{F}_1^\alpha \gamma^+}{M} \left( \frac{\bar{p}_T + \gamma^5}{\gamma^\nu} \right)
\end{array} \right) \tag{51} \]

where \( \phi_p \) is the azimuthal angle of \( \vec{p}_T. \) The matrix is Hermitian, respects parity invariance and conservation of total angular momentum. Introducing the dependence upon the quark transverse momentum \( \vec{p}_T \) modifies the conditions for angular momentum and parity conservation, which now read, respectively,

\[ \lambda \left\{ \lambda \right\} + l_{\lambda \pi} = \Lambda_{\lambda} + \chi_{\lambda}, \]

\[ \left[ \mathcal{P}_+ \Phi_\alpha \gamma^+ \right]^\Lambda_{\Lambda} \gamma_{\chi_1 \chi_1} = \left( -1 \right)^{l_{\lambda \pi}} \left[ \mathcal{P}_+ \Phi_\alpha \gamma^+ \right]^\Lambda_{\Lambda} \gamma_{\chi_1 \chi_1}, \tag{52} \]

where \( l_{\lambda \pi} \) denotes the units of angular momentum introduced by \( \vec{p}_T. \) The chiral transposed matrix is still positive definite, so that the bounds on the various distribution functions can be obtained \cite{24}.

The leading-twist projection of the fragmenting quark correlator is

\[ \mathcal{P}_- \Delta_\alpha(z, \vec{k}_T^z, \zeta, M^2 \omega \phi_R) \gamma^+ \gamma^+ = \frac{1}{8 \pi} \left( D_1^\alpha(z, \zeta, M^2 \omega \phi_R) \gamma^+ \gamma^+ \right), \tag{53} \]

where the actual dependence of the fragmentation functions is the most general one possible \cite{3}. In Eq. (53) \( \mathcal{P}_- \Delta_\alpha \) corresponds to Eq (3) of Ref. [6]. New functions appear: \( G_T^1 \) is chiral even but T-odd, \( H_T^1 \) is chiral odd and T-odd and represents the analogue of the Collins effect for a two-hadron emission \cite{3}. Upon integration over \( d\vec{k}_T, G_T^1 \) vanishes.
and the surviving parts of \( \tilde{H}_1^\pm \) and \( H_1^\pm \) merge into the function \( H_1^\pm \) of Eq. (16) keeping the same \( \vec{R}_T/M_h \) structure. In the chiral basis of the fragmenting quark, Eq. (53) becomes

\[
[P_\Delta (z, \vec{k}_T, \zeta, M_h^2, \phi_R) r] H_1^\pm \chi_{\pm 2} = \frac{1}{8\pi} \left( \begin{array}{c}
D_1^+ + \frac{\alpha^e R_T\mu k_T}{M_h^2} G_1^\pm \, \chi_{\pm 2} \\
- i \left( \frac{\alpha^e}{\vec{k}_T} \frac{\partial}{\partial \vec{k}_T} \tilde{R}_1^\pm \, \chi_{\pm 2} + e^{i\phi_0} \frac{\partial}{\partial \phi_0} \tilde{R}_1^\pm \, \chi_{\pm 2} \right) \end{array} \right) \times \left( \begin{array}{c}
\frac{\alpha^e R_T\mu k_T}{M_h^2} \tilde{R}_1^\pm \, \chi_{\pm 2} + e^{i\phi_0} \frac{\partial}{\partial \phi_0} \tilde{R}_1^\pm \, \chi_{\pm 2} \\
- i \left( \frac{\alpha^e}{\vec{k}_T} \frac{\partial}{\partial \vec{k}_T} \tilde{R}_1^\pm \, \chi_{\pm 2} + e^{i\phi_0} \frac{\partial}{\partial \phi_0} \tilde{R}_1^\pm \, \chi_{\pm 2} \right) \end{array} \right).
\]

The following bounds are derived:

\[
\left| \frac{\alpha^e R_T}{M_h^2} \tilde{R}_1^\pm \right|^2 \leq D_1^+ \\
\left| \frac{\alpha^e R_T}{M_h^2} \tilde{R}_1^\pm \right|^2 + \left| \frac{\alpha^e}{\vec{k}_T} \frac{\partial}{\partial \vec{k}_T} \tilde{R}_1^\pm \right|^2 \leq (D_1^+)^2 + \left( \frac{\alpha^e}{\vec{k}_T} \frac{\partial}{\partial \vec{k}_T} \tilde{R}_1^\pm \right)^2 \leq (D_1^+)^2.
\]

Expanding the cross section of Eq. (45) along the same lines leading to Eq. (21), we have

\[
\frac{d^3\sigma}{d\zeta\, dM_h^2 \, d\phi_R \, dP_{\perp h} \, dz \, dx \, dy \, d\phi_S} = \sum \rho_{\Lambda \Lambda} (S) \mathcal{I} \left[ \frac{\partial}{\partial y} \frac{d^3\sigma}{d\zeta \, dM_h^2 \, d\phi_R} \left( \begin{array}{c}
\frac{\partial}{\partial y} \frac{d^3\sigma}{d\zeta \, dM_h^2 \, d\phi_R} \chi_{\pm 2} \\
\chi_{\pm 2} \end{array} \right) \right] \times \left[ \frac{\partial}{\partial y} \frac{d^3\sigma}{d\zeta \, dM_h^2 \, d\phi_R} \chi_{\pm 2} \right],
\]

where \([P_\Delta \Phi_D \gamma]^{-}\) and \([P_\Delta \Delta \gamma]^{-}\) are given by Eqs. (51) and (54), respectively. The complete formula for the cross section is given in App. A.

V. PARTIAL-WAVE EXPANSION WITH TRANSVERSE MOMENTA

It is again useful to expand all the fragmentation functions of Eq. (53) in the relative partial waves of the hadron pair. The dependence on \( \vec{k}_T \cdot \tilde{R}_T \) makes the expansion more involved:

\[
D_1 = D_{1,OO} + D_{1,OL} \cos \theta + D_{1,LL} \left( 3 \cos^2 \theta - 1 \right) \cos (\phi_k - \phi_R) \sin \theta (D_{1,OT} + D_{1,LT} \cos \theta) + \cos (2\phi_k - 2\phi_R) \sin^2 \theta D_{1,TT} \\
G_1^+ = G_{1,OT} + G_{1,LT} \cos \theta + \cos (\phi_k - \phi_R) \sin \theta G_{1,TT} \\
\tilde{R}_1^\pm = \tilde{R}_{1,OT}^\pm + \tilde{R}_{1,LT}^\pm \cos \theta + 2 \cos (2\phi_k - 2\phi_R) \sin \theta \tilde{R}_{1,TT}^\pm \\
H_1^\pm = H_{1,OO}^\pm + H_{1,OL} \cos \theta + H_{1,LL} \left( 3 \cos^2 \theta - 1 \right) + 2 \cos (\phi_k - \phi_R) \sin \theta (H_{1,OT}^\pm + H_{1,LT} \cos \theta) + 2 \cos (2\phi_k - 2\phi_R) \sin^2 \theta H_{1,TT}^\pm - \sin^2 \theta \left( \frac{\vec{k}_T}{\vec{k}_T} \right) H_{1,TT}^\pm,
\]

where all the functions depend on \((z, \vec{k}_T, M_h^2)\). Then, similarly to Eq. (35), Eq. (54) can be further expanded in the basis of the pair orbital angular momentum as

\[
[P_\Delta \Phi_D (z, \vec{k}_T, M_h^2, \phi_R) r] \chi_{\pm 2} = [P_\Delta \Phi_D (z, \vec{k}_T, M_h^2, \phi_R) r] \chi_{\pm 2} \mathcal{D}_{LL}^{LM} (\theta, \phi_R).
\]

The full expression of \([P_\Delta \Delta \gamma]^{-}\) is shown in App. B. The fully expanded differential cross section in the helicity basis of target, initial and final quark, as well as in the basis of orbital angular momentum of the hadron pair is then

\[
\frac{d^3\sigma}{d\zeta\, dM_h^2 \, d\phi_R \, dP_{\perp h} \, dz \, dx \, dy \, d\phi_S} = \sum \rho_{\Lambda \Lambda} (S) \mathcal{I} \left[ \frac{\partial}{\partial y} \frac{d^3\sigma}{d\zeta \, dM_h^2 \, d\phi_R} \left( \begin{array}{c}
\frac{\partial}{\partial y} \frac{d^3\sigma}{d\zeta \, dM_h^2 \, d\phi_R} \chi_{\pm 2} \\
\chi_{\pm 2} \end{array} \right) \right] \times \left[ \frac{\partial}{\partial y} \frac{d^3\sigma}{d\zeta \, dM_h^2 \, d\phi_R} \chi_{\pm 2} \right] \mathcal{D}_{LL}^{LM} (\theta, \phi_R).
\]

Its explicit expression is presented in App. C. The pure \( p \)-wave sector corresponds to the cross section for the production of a polarized spin-1 hadron and has already been fully studied in Refs. [13, 38]. For sake of completeness, we show it in App. C together with the formulae for the pure \( s \) and \( s-p \) interference sectors.
VI. CONCLUSIONS

In this paper, we have reconsidered the option of extracting the transversity distribution \( h_1 \) at leading twist by using the analyzing power of the interference fragmentation functions (IFF) into two leading unpolarized hadrons inside the same current jet. As already shown in Ref. [6], in the process \( e p \rightarrow e' h_1 h_2 X \) the transversity distribution enters a single-spin asymmetry in the azimuthal angle \( \phi_R \) of the hadron pair plane. The effect survives after the integration upon the transverse component of the 
finite cross section \( \mathcal{A}(\mathbf{p}_T) \) and \( \mathcal{A}(\mathbf{p}_T') \). Therefore, no transverse-momentum dependent function is required and the advantage with respect to the Collins effect is evident. A similar situation was known to occur in the case of fragmentation into spin-1 hadrons [10, 11, 12, 13], but it was never fully examined to the extent of defining a specific asymmetry.

Here, we have reanalyzed the whole problem in the helicity formalism by further expanding the IFF in the basis of the relative orbital angular momentum in the cm frame of the hadron pair. New positivity bounds have been derived. If the invariant mass of the pair is not large, the expansion can be limited to the first two modes, namely the relative \( s \) and \( p \) waves.

Off-diagonal elements in the chirality and in the orbital angular momentum \( \frac{\hat{t}}{f} \) represent the IFF of Ref. [2] and [6], where the interference arises from the hadron pair being in a state with either \( s \) or \( p \) relative wave. Elements in the \( L = L' = 1 \) sector correspond to the analysis of spin-1 hadron fragmentation [13]. Therefore, the present formalism represents a unifying framework for the problem of fragmentation into two unpolarized hadrons and can be used to correctly and exhaustively discuss the extraction of transversity from two-hadron lepton production.

In fact, after calculating the complete leading-twist cross section, we have identified a single spin asymmetry containing two distinct chiral-odd partners of the transversity. By integrating the asymmetry over different ranges of the cm polar angle of the hadron pair, the transversity \( h_1 \) can be extracted through the chiral-odd, T-odd fragmentation \( H^2_{h_1} \) (corresponding to the \( s-p \) interference of Ref. [2]) or through the chiral-odd, T-odd fragmentation \( H^0_{h_1} \) (corresponding to the \( p-p \) interference). This second option has been often neglected in the literature, despite the fact that the two functions have, in principle, a different dependence on the invariant mass and a different physical origin.

In conclusion, we believe that the fragmentation into two leading unpolarized hadrons can be a promising tool to measure the transversity distribution, as well as to achieve further comprehension of the hadronization mechanism.

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APPENDIX A

In this appendix, we write explicitly the cross section for two-hadron lepton production at leading order in \( 1/Q \) and with the inclusion of partonic transverse momenta. Moreover, we include also T-odd distribution functions, since recently there have been some indications that they are not forbidden by time-invariance [28, 27, 30]. To simplify the notation, we introduce the projection \( \mathbf{\tilde{a}_T} \propto \alpha_T \tilde{q}_T \). Inserting in Eq. (45) the formulae for the target helicity density matrix, Eq. (12), for the distribution correlation matrix, Eq. (51), for the elementary scattering matrix, Eq. (22), and the two-hadron fragmentation matrix, Eq. (54), we obtain the following result:

\[
\begin{align*}
\hat{d}^2\sigma_{GO} &= \sum_a \frac{\alpha_s^2 \hat{s}_a^2}{2\pi Q^2 y} \left\{ A(y) \mathcal{I} [f_1 D_1] - B(y) \frac{\hat{R}_T}{M} \cos(\phi_h + \phi_R) \mathcal{I} \left[ \frac{\hat{f}_T \cdot \hat{P}_{h\perp}}{M} k_1^T \tilde{R}_1^T \right] \right. \\
&\quad + B(y) \left( \frac{\hat{R}_T}{M} \sin(\phi_h + \phi_R) \mathcal{I} \left[ \frac{\hat{P}_{h\perp} \land \hat{f}_T}{M} k_1^T \tilde{R}_1^T \right] \right) \\
&\quad - B(y) \cos(2\phi_h) \mathcal{I} \left[ \frac{2(\hat{f}_T \cdot \hat{P}_{h\perp} \land \hat{k}_T) - (\hat{k}_T \cdot \hat{P}_{h\perp})}{M M_h} k_1^T H_1^T \right] \\
&\quad + B(y) \sin(2\phi_h) \mathcal{I} \left[ \frac{(\hat{f}_T \cdot \hat{P}_{h\perp} \land \hat{k}_T) + (\hat{k}_T \cdot \hat{P}_{h\perp})}{M M_h} k_1^T H_1^T \right] \right\},
\end{align*}
\]

\[
\begin{align*}
\hat{d}^2\sigma_{LO} &= -\sum_a \frac{\alpha_s^2 \hat{s}_a^2}{2\pi Q^2 y} \left| C_a \right| \frac{\hat{R}_T}{M} \left\{ \sin(\phi_h - \phi_R) \mathcal{I} \left[ \frac{\hat{k}_T \cdot \hat{P}_{h\perp} f_1 G_1^T}{M} \right] \right. \\
&\quad + B(y) \cos(2\phi_h) \mathcal{I} \left[ \frac{(\hat{f}_T \cdot \hat{P}_{h\perp} \land \hat{k}_T) - (\hat{k}_T \cdot \hat{P}_{h\perp})}{M M_h} k_1^T H_1^T \right] \\
&\quad + B(y) \sin(2\phi_h) \mathcal{I} \left[ \frac{(\hat{f}_T \cdot \hat{P}_{h\perp} \land \hat{k}_T) + (\hat{k}_T \cdot \hat{P}_{h\perp})}{M M_h} k_1^T H_1^T \right] \right\}.
\end{align*}
\]
\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| \left\{ -A(y) \frac{\vec{R}_T}{M_h} \sin(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
- A(y) \frac{\vec{R}_T}{M_h} \sin(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ B(y) \left[ \frac{\vec{R}_T}{M_h} \sin(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \right] \\
+ B(y) \cos(2\phi_h) \left[ \frac{2(\vec{P}_T \cdot \vec{P}_T) \vec{P}_T \cdot \vec{P}_T - \vec{P}_T \cdot \vec{P}_T}{2 M M_h} \right] \\
+ B(y) \cos(2\phi_h) \left[ \vec{P}_T \cdot \vec{P}_T \g_1 \vec{D}_1 \right] \} \). \tag{A3} \\
\]

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| C(y) \left[ g_1 D_1 \right] \). \tag{A4} \\
\]

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| \left\{ -A(y) \frac{\vec{R}_T}{M_h} \sin(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
- \sin(\phi_h - \phi_R) \left[ \vec{P}_T \cdot \vec{P}_T \g_1 \vec{D}_1 \right] + \cos(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ A(y) \frac{\vec{R}_T}{M_h} \sin(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \\
+ \sin(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \} \]. \tag{A5} 

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| A(y) \left\{ -\frac{\vec{R}_T}{M_h} \sin(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
- \frac{\vec{R}_T}{M_h} \sin(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \cos(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \cos(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
\} \]. \tag{A6} 

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| B(y) \left\{ \sin(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \\
+ \cos(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \} \]. \tag{A7} 

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| \left\{ -\frac{\vec{R}_T}{M_h} \sin(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \sin(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \cos(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \cos(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \} \]. \tag{A8} 

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| C(y) \left[ g_1 D_1 \right] \). \tag{A9} 

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| A(y) \left\{ -\frac{\vec{R}_T}{M_h} \sin(\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
- \frac{\vec{R}_T}{M_h} \sin(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \cos(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \\
+ \frac{\vec{R}_T}{M_h} \cos(2\phi_h - \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \g_1 G_1^1 \right] \} \]. \tag{A10} 

\[ d^3 \sigma_{GL} = \sum_a \frac{\alpha^2 s_a^2}{2 \pi Q^2} \left| S_L \right| B(y) \left\{ \sin(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \\
+ \cos(\phi_h + \phi_R) \left[ \vec{P}_L \cdot \vec{P}_T \bar{\vec{R}}_1 \hat{h}_1 \right] \} \]. \tag{A11}
\[ d^3 \sigma_{LT} = \sum_a \frac{\alpha^2 e_a^2}{2\pi Q^2 y} |S_T|^2 C(y) \left\{ \cos(\phi_b - \phi_s) \mathcal{I} \left[ \frac{p_T^T \cdot p_{hL}}{M} g_{1T} \right] D_1 \right. \\
- \sin(\phi_b - \phi_s) \mathcal{I} \left[ \frac{p_{hL} \wedge p_T^T}{M} g_{1T} \right] D_1 \right. \\
- \left. \frac{|R_T|}{M_h} \cos(\phi_R - \phi_s) \mathcal{I} \left[ \frac{p_T^T \cdot \hat{k}_R}{2M_M} f_{1T}^R G_1^T \right] \right. \\
+ \left. \frac{|R_T|}{M_h} \cos(2\phi_b - \phi_R - \phi_s) \mathcal{I} \left[ \frac{2(p_T^T \cdot p_{hL})(\hat{k}_R \cdot p_{hL})}{2M_M} - \frac{p_T^T \cdot \hat{k}_R}{M_M} f_{1T}^R G_1^T \right] \right. \\
- \left. \frac{|R_T|}{M_h} \sin(\phi_R - \phi_s) \mathcal{I} \left[ \frac{(p_T^T \cdot p_{hL})(p_{hL} \wedge \hat{k}_R) - (\hat{k}_R \cdot p_{hL})(p_{hL} \wedge p_T^T)}{2M_M} f_{1T}^R G_1^T \right] \right. \\
+ \left. \frac{|R_T|}{M_h} \sin(2\phi_b - \phi_R - \phi_s) \mathcal{I} \left[ \frac{(p_T^T \cdot p_{hL})(p_{hL} \wedge \hat{k}_R) + (\hat{k}_R \cdot p_{hL})(p_{hL} \wedge p_T^T)}{2M_M} f_{1T}^R G_1^T \right] \right\}. \]  

(A6)

In the case of \( d^3 \sigma_{OT} \), i.e. for an unpolarized beam and a transversely polarized target, the full expression of the cross section corresponds to the one in Eq. (10) of Ref. [6], apart for a different overall factor, due to slightly different definitions of the hadron tensor and of the fragmentation functions, and the use of \( M_h \) instead of \( M_1(M_2) \) in the denominators, due to a different definition of the expansion (15).

**APPENDIX B**

The full expression of \( \mathcal{P}_- \Delta(z, \vec{k}_T^2, M_h^2)^{-1} L^{LL}_{M_1 M_2} \chi_1^2 \chi_2 \) in Eq. (58) is

\[ \mathcal{P}_- \Delta(z, \vec{k}_T^2, M_h^2)^{-1} L^{LL}_{M_1 M_2} \chi_1^2 \chi_2 = \frac{1}{8} \left( A^{L,L}_{M,M} \right)^T \left( B^{L,L}_{M,M} \right), \]  

(B1)

where

\[
A^{L,L}_{M,M} = 
\begin{pmatrix}
D_{1,0,0}^{L,L} & -\sqrt{2} e^{id} (D_{1,0,T} + i \frac{H_{1,0}^L}{M_1} G_{1,0}^T) & \sqrt{2} e^{id} (D_{1,0,T} - i \frac{H_{1,0}^L}{M_1} G_{1,0}^T) \\
-\sqrt{2} e^{-id} (D_{1,0,L} - i \frac{H_{1,0}^L}{M_1} G_{1,0}^T) & D_{1,0,0}^{L,L} + \frac{i}{2} D_{1,0,L} & -\sqrt{2} e^{-id} (D_{1,0,L} + i \frac{H_{1,0}^L}{M_1} G_{1,0}^T) \\
\sqrt{2} e^{id} (D_{1,0,T} + i \frac{H_{1,0}^L}{M_1} G_{1,0}^T) & -\sqrt{2} e^{id} (D_{1,0,L} + i \frac{H_{1,0}^L}{M_1} G_{1,0}^T) & D_{1,0,0}^{L,L} + \frac{i}{2} D_{1,0,L}
\end{pmatrix}
\]

(B2)

\[ B_{M,M}^{L,L} = 
\begin{pmatrix}
e^{id} H_{1,0}^{L,0} & -2\sqrt{2} e^{2id} H_{1,0}^{L,T} & 2\sqrt{2} e^{id} H_{1,0}^{L,L} \\
-2\sqrt{2} e^{-2id} H_{1,0}^{T,T} & e^{id} (H_{1,0,0}^{L,0} - i H_{1,0}^{L,L}) & -2\sqrt{2} e^{id} H_{1,0}^{L,T} \\
2\sqrt{2} e^{-2id} H_{1,0}^{T,L} & -2\sqrt{2} e^{-id} H_{1,0}^{L,0} & e^{id} (H_{1,0,0}^{L,0} + i H_{1,0}^{L,L})
\end{pmatrix},
\]

(B3)

and \( \phi \equiv \phi_b - \phi_R \). The matrix (B1) respects Hermiticity, angular momentum conservation, and parity invariance.

Due to the explicit dependence upon the transverse momentum \( \vec{k}_T \), the conditions for angular momentum and parity conservation read

\[ M + \chi_2 = M' + \chi_3 + l_{k_T}, \]

\[ \mathcal{P}_- \Delta \gamma^{L,L}_{M,M} \chi_1^2 \chi_2 = (-1)^{L_T} \mathcal{P}_- \Delta \gamma^{L,L}_{M',-M-M} \chi_3 - \chi_3, \]  

(B4)
where $l_k\phi$ denotes the units of angular momentum introduced by $\bar{k}_T$. From the last constraint it is possible to derive the lower right block, i.e. $C_{LM}^{L'} = (-1)^{L''} A_{LM}^{L'} M_{-M}$.

Again, as in the case of Eq. (39), we have

$$H_{1,00}(z, \bar{k}_T^2, M_b^2) = \frac{1}{2} H_{1,00}^{L}(z, \bar{k}_T^2, M_b^2) + \frac{3}{4} H_{1,00}^{L'}(z, \bar{k}_T^2, M_b^2)$$ (B5)

and the functions $H_{1,00}^{L}$, $H_{1,00}^{L'}$ are kinematically indistinguishable unless some hypothesis is made on their $M_b^2$ dependence. The $L = L' = 1$ sector of Eqs. (B2,B3) has been studied in the case of spin-1 fragmentation [13]. The interference ($L = 0, L' = 1$) sector has never been analyzed in this form, namely including the explicit dependence on $\bar{k}_T$. Finally, from $|\mathcal{P}_- A^{L,L'}_{M',M,\chi,\chi}|$ being positive semidefinite, it is possible to derive bounds on each of the displayed fragmentation functions.

**APPENDIX C**

In this appendix, we explicitly present the complete cross section for the production of two unpolarized hadrons in relative $s$ and $p$ waves, at leading order in $1/Q$, including transverse momenta and T-odd distribution and fragmentation functions.

The cross section is obtained by replacing Eqs. (12,51,22,B1,34) in Eq. (59). It is convenient to introduce the following combination of fragmentation functions

$$H_{1,OT}^{L} = \bar{R}_{1,OT}^{L} + \frac{\bar{k}_T}{|\bar{k}_T|} H_{1,OT}^{L},$$

$$H_{1,LT}^{L} = \bar{R}_{1,LT}^{L} + \frac{\bar{k}_T}{|\bar{k}_T|} H_{1,LT}^{L},$$

$$H_{1,TT}^{L} = \bar{R}_{1,TT}^{L} + \frac{\bar{k}_T}{|\bar{k}_T|} H_{1,TT}^{L}.$$ (C3)
\[d^8\sigma_{OO} = \sum_a \frac{\alpha^2 e^2_a}{2\pi s x y^a} A(y) \left\{ \mathcal{I} \left[ f_1 \left( \frac{1}{4} D_1^{\perp,OO} + \frac{3}{4} D_1^{\parallel,OO} \right) \right] + \cos \theta \mathcal{I} \left[ f_1 D_{1,OL} \right] \right. \\
+ \frac{1}{3} \left( 3 \cos^2 \theta - 1 \right) \mathcal{I} \left[ f_1 \left( \frac{3}{4} D_{1,LL} \right) \right] + \sin \theta \cos(\phi_b - \phi_R) \mathcal{I} \left[ \frac{\bar{\vec{k}}_T \cdot \bar{P}_h}{M_h} f_1 \left( -\frac{\bar{R}}{M_h} D_{1,LT} \right) \right] \right. \\
- \sin 2\theta \cos(\phi_b - \phi_R) \mathcal{I} \left[ \frac{\bar{\vec{k}}_T \cdot \bar{P}_h}{M_h} f_1 \left( -\frac{\bar{R}}{M_h} D_{1,LT} \right) \right] - \sin^2 \theta \cos(2\phi_b - 2\phi_R) \\
\times \mathcal{I} \left[ \frac{2(\bar{\vec{k}}_T \cdot \bar{P}_h)^2 - \bar{\vec{k}}_T^2}{M_h^2} f_1 \left( -\frac{\bar{R}}{M_h} D_{1,LT} \right) \right] \right. \\
\left. + \sum_a \frac{\alpha^2 e^2_a}{2\pi s x y^a} B(y) \left\{ -\cos 2\phi_b \mathcal{I} \left[ \frac{2(\bar{\vec{p}}_T \cdot \bar{P}_h)(\bar{\vec{k}}_T \cdot \bar{P}_h) - \bar{\vec{p}}_T \cdot \bar{\vec{k}}_T}{M_{h_h}} \right] h_1^+ \left( \frac{1}{4} H_{1,0}^+ + \frac{3}{4} H_{1,0}^\perp \right) \right. \\
- \frac{1}{3} \left( 3 \cos^2 \theta - 1 \right) \cos 2\phi_b \mathcal{I} \left[ \frac{2(\bar{\vec{p}}_T \cdot \bar{P}_h)(\bar{\vec{k}}_T \cdot \bar{P}_h) - \bar{\vec{p}}_T \cdot \bar{\vec{k}}_T}{M_{h_h}} \right] h_1^+ \left( \frac{3}{4} H_{1,LL}^\perp \right) \\
+ \sin \theta \cos(\phi_b + \phi_R) \mathcal{I} \left[ \frac{\bar{\vec{p}}_T \cdot \bar{P}_h}{M} h_1^+ \left( -\frac{\bar{R}}{M_h} H_{1,0}^\perp \right) \right] + \sin 2\theta \cos(\phi_b + \phi_R) \\
\times \mathcal{I} \left[ \frac{\bar{\vec{p}}_T \cdot \bar{P}_h}{M} h_1^+ \left( -\frac{\bar{R}}{M_h} H_{1,0}^\perp \right) \right] + \sin \theta \cos(3\phi_b - \phi_R) \\
\times \mathcal{I} \left[ \frac{4(\bar{\vec{k}}_T \cdot \bar{P}_h)^2 - 2(\bar{\vec{k}}_T \cdot \bar{P}_h)(\bar{\vec{k}}_T \cdot \bar{P}_h) - \bar{\vec{k}}_T^2(\bar{\vec{p}}_T \cdot \bar{P}_h)}{M_{h_h}^2} \right] h_1^+ \left( -\frac{2M_h}{\bar{R}} H_{1,LT}^\perp \right) \\
+ \sin 2\theta \cos(3\phi_b - \phi_R) \\
\times \mathcal{I} \left[ \frac{4(\bar{\vec{k}}_T \cdot \bar{P}_h)^2 - 2(\bar{\vec{k}}_T \cdot \bar{P}_h)(\bar{\vec{k}}_T \cdot \bar{P}_h) - \bar{\vec{k}}_T^2(\bar{\vec{p}}_T \cdot \bar{P}_h)}{M_{h_h}^2} \right] h_1^+ \left( -\frac{M_h}{\bar{R}} H_{1,LT}^\perp \right) \\
+ \sin^2 \theta \cos(4\phi_b - 2\phi_R) \mathcal{I} \left[ \left( \frac{\bar{\vec{k}}_T^2 - 4(\bar{\vec{k}}_T \cdot \bar{P}_h)^2}{2M_{h_h}^2} \right) \left[ \bar{\vec{p}}_T \cdot \bar{\vec{k}}_T - 4(\bar{\vec{k}}_T \cdot \bar{P}_h)(\bar{\vec{p}}_T \cdot \bar{P}_h) \right] \right. \\
- \frac{8(\bar{\vec{k}}_T \cdot \bar{P}_h)^3 (\bar{\vec{p}}_T \cdot \bar{P}_h)}{M_{h_h}^3} \right) h_1^+ \left( -\frac{2M_h^2}{\bar{R}} H_{1,LT}^\perp \right) \right\}. 
\]

### b. Polarized lepton beam and unpolarized target

\[d^8\sigma_{LO} = -\sum_a \frac{\alpha^2 e^2_a}{2\pi s x y^a} \lambda_c C(y) \left\{ \sin \theta \sin(\phi_b - \phi_R) \mathcal{I} \left[ \frac{\bar{\vec{k}}_T \cdot \bar{P}_h}{M_h} f_1 \left( \frac{\bar{R}}{M_h} G_{1,LT}^\perp \right) \right] \right. \\
+ \sin 2\theta \sin(\phi_b - \phi_R) \mathcal{I} \left[ \frac{\bar{\vec{k}}_T \cdot \bar{P}_h}{M_h} f_1 \left( \frac{\bar{R}}{M_h} G_{1,LT}^\perp \right) \right] \right. \\
+ \sin^2 \theta \sin(2\phi_b - 2\phi_R) \mathcal{I} \left[ \left( \frac{1}{2\bar{R}} G_{1,LT}^\perp \right) \right] \right. \\
\left. + \sin 2\theta \sin(2\phi_b - 2\phi_R) \mathcal{I} \left[ \left( \frac{1}{2\bar{R}} G_{1,LT}^\perp \right) \right] \right\}. 
\]
c. Unpolarized lepton beam and longitudinally polarized target

\[ d^6\sigma_{OL} = -\sum_{\alpha} \frac{\alpha^2 e_{\alpha}^2}{2\pi \sigma y^2} \left| S_L \right| A(y) \left\{ \sin \theta \sin(\phi_h - \phi_R) \right. \mathcal{T} \left[ \frac{\vec{p}_T \cdot P_{h\perp}}{M_h} g_{1L} \left( \frac{1}{M_h} G^1_{1,OL} \right) \right] \\
+ \sin 2\theta \sin(\phi_h - \phi_R) \mathcal{T} \left[ \frac{\vec{E}_T \cdot P_{h\perp}}{M_h} g_{1L} \left( \frac{1}{2M_h} G^1_{1,LT} \right) \right] \\
+ \sin^2 \theta \sin(2\phi_h - 2\phi_R) \mathcal{T} \left[ \frac{2(\vec{E}_T \cdot P_{h\perp})^2 - \vec{E}_T^2}{M_M h_{1L} \left( \frac{1}{4} H^1_{1,OO} + \frac{3}{4} H^1_{1,OO} \right) \right] \} \\
- \sum_{\alpha} \frac{\alpha^2 e_{\alpha}^2}{2\pi \sigma y^2} \left| S_L \right| B(y) \left\{ \sin 2\phi_B \mathcal{T} \left[ \frac{2(\vec{p}_T \cdot P_{h\perp})^2 - \vec{p}_T^2}{M_M h_{1L}} \left( \frac{1}{2M_h} G^1_{1,LL} \right) \right] \\
+ \sin \theta \sin(\phi_h + \phi_R) \mathcal{T} \left[ \frac{\vec{p}_T \cdot P_{h\perp}}{M_h} h_{1L} \left( -\frac{1}{2M_h} H^1_{1,LT} \right) \right] + \sin 2\theta \sin(\phi_h + \phi_R) \mathcal{T} \left[ \frac{\vec{p}_T \cdot \vec{E}_T}{M_M h_{1L}} h_{1L} \left( -\frac{1}{2M_h} H^1_{1,LT} \right) \right] \right\} (C6) \\
+ \sin \theta \sin(3\phi_h - \phi_R) \mathcal{T} \left[ 4(\vec{E}_T \cdot P_{h\perp})^2 \left( \frac{\vec{p}_T \cdot P_{h\perp}}{M_M h_{1L}} - \frac{2M_h}{\vec{E}_T^2} H^1_{1,LL} \right) \right] \\
+ \sin 2\theta \sin(3\phi_h - \phi_R) \mathcal{T} \left[ 8(\vec{E}_T \cdot P_{h\perp})^3 \left( \frac{\vec{p}_T \cdot P_{h\perp}}{M_M h_{1L}} - \frac{2M_h}{\vec{E}_T^2} H^1_{1,LL} \right) \right] \\
+ \sin^2 \theta \sin(4\phi_h - 2\phi_R) \mathcal{T} \left[ \left( \frac{2^{\frac{3}{2}}}{2M_M h_{1L}} \right) \left( -\frac{2M_h}{\vec{E}_T^2} H^1_{1,LT} \right) \right] \\
\] \\
d. Polarized lepton beam and longitudinally polarized target

\[ d^6\sigma_{LL} = \sum_{\alpha} \frac{\alpha^2 e_{\alpha}^2}{2\pi \sigma y^2} \lambda_\alpha \left| S_L \right| C(y) \left\{ \mathcal{T} \left[ g_{1L} \left( \frac{1}{4} D^1_{1,OO} + \frac{3}{4} D^1_{1,OO} \right) \right] + \cos \theta \mathcal{T} \left[ g_{1L} D_{1,OL} \right] \\
+ \frac{1}{3} (3 \cos^2 \theta - 1) \mathcal{T} \left[ g_{1L} \left( \frac{3}{2} D_{1,LL} \right) \right] \\
- \sin \theta \cos(\phi_h - \phi_R) \mathcal{T} \left[ \frac{\vec{E}_T \cdot P_{h\perp}}{M_h} g_{1L} \left( -\frac{M_h}{\vec{E}_T} D_{1,LT} \right) \right] \} (C7) \\
- \sin 2\theta \cos(\phi_h - \phi_R) \mathcal{T} \left[ \frac{\vec{E}_T \cdot P_{h\perp}}{M_h} g_{1L} \left( -\frac{M_h}{2\vec{E}_T} D_{1,LT} \right) \right] \\
- \sin^2 \theta \cos(2\phi_h - 2\phi_R) \mathcal{T} \left[ \frac{2(\vec{E}_T \cdot P_{h\perp})^2 - \vec{E}_T^2}{M_M h_{1L}} g_{1L} \left( -\frac{M_h^2}{\vec{E}_T^2} D_{1,LT} \right) \right] \right\},
\[ d^8 \sigma_{OT} = \sum_{a} \frac{\alpha^2 \alpha_s^2}{2 \pi s x y^2} |S_T| A(y) \left\{ \sin \theta \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{(\vec{p}_T \cdot \vec{k}_T)}{2M_{b\bar{b}}} \frac{g_{1T}}{M_b} \frac{1}{G^1_{1,OT}} \right] \right. \\
- \sin \theta \sin(2\phi_R - \phi_S) \mathcal{I} \left[ \frac{2(\vec{p}_T \cdot P_{b\perp})(\vec{k}_T \cdot P_{b\perp}) - \vec{p}_T \cdot \vec{k}_T}{2M_{b\bar{b}}} \frac{g_{1T}}{M_b} \frac{1}{G^1_{1,LT}} \right] \\
+ \sin 2 \theta \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{(\vec{p}_T \cdot \vec{k}_T)}{2M_{b\bar{b}}} \frac{g_{1T}}{M_b} \frac{1}{G^1_{1,LT}} \right] \\
- \sin 2 \theta \sin(2\phi_R - \phi_S) \mathcal{I} \left[ \frac{2(\vec{p}_T \cdot P_{b\perp})(\vec{k}_T \cdot P_{b\perp}) \cdot P_{b\perp}}{2M_{b\bar{b}}} \frac{g_{1T}}{M_b} \frac{1}{G^1_{1,LT}} \right] \\
- \frac{\vec{k}_T^2 (\vec{p}_T \cdot P_{b\perp})}{2M_{b\bar{b}}} \left( \frac{1}{G^1_{1,TT}} \right) + \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{\vec{p}_T \cdot P_{b\perp}}{M_b} \frac{f_{1T}}{f_{1T}^\perp \left( \frac{\frac{3}{4} D_{1,LO} + \frac{3}{4} D_{1,OO}}{2} \right) \right] \\
+ \cos \theta \mathcal{I} \left[ f_{1T}^\perp D_{1,LL} \right] + \frac{1}{3} (3 \cos^2 \theta - 1) \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{\vec{p}_T \cdot P_{b\perp}}{M_b} \frac{f_{1T}}{f_{1T}^\perp \left( \frac{\frac{3}{4} D_{1,LO} + \frac{3}{4} D_{1,OO}}{2} \right) \right] \\
- \sin \theta \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{(\vec{p}_T \cdot \vec{k}_T)}{2M_{b\bar{b}}} \frac{f_{1T}}{f_{1T}^\perp \left( \frac{M_b}{2} \right) D_{1,OT} \right] \\
- \sin 2 \theta \sin(2\phi_R - \phi_S) \mathcal{I} \left[ \frac{2(\vec{p}_T \cdot P_{b\perp})(\vec{k}_T \cdot P_{b\perp}) \cdot P_{b\perp}}{2M_{b\bar{b}}} \frac{f_{1T}}{f_{1T}^\perp \left( \frac{M_b}{2} \right) D_{1,LT} \right] \\
+ \sin^2 \theta \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{2(\vec{p}_T \cdot P_{b\perp})(\vec{k}_T \cdot P_{b\perp}) \cdot P_{b\perp}}{2M_{b\bar{b}}} \frac{f_{1T}}{f_{1T}^\perp \left( \frac{M_b}{2} \right) D_{1,LT} \right] \\
- \frac{\vec{k}_T^2 (\vec{p}_T \cdot P_{b\perp})}{2M_{b\bar{b}}} \left( \frac{1}{G^1_{1,TT}} \right) \right\} \\
+ \sum_{a} \frac{\alpha^2 \alpha_s^2}{2 \pi s x y^2} B(y) \left\{ \cos 2 \phi_R \mathcal{I} \left[ \frac{2(\vec{p}_T \cdot P_{b\perp})(\vec{k}_T \cdot P_{b\perp}) \cdot P_{b\perp}}{M_{b\bar{b}}} \frac{f_{1T}}{f_{1T}^\perp \left( \frac{M_b}{2} \right) D_{1,LT} \right] \\
+ \sin 2 \theta \sin(\phi_R - \phi_S) \mathcal{I} \left[ \frac{\vec{k}_T \cdot P_{b\perp}}{M_b} \frac{h_1}{h_1} \left( \frac{3}{4} H_{1,LO} + \frac{3}{4} H_{1,OO} \right) \right] \\
- \sin \theta \sin(\phi_R + \phi_S) \mathcal{I} \left[ h_1 \left( \frac{\vec{k}_T}{M_b} H_{1,OT}^{\perp} \right) \right] \right\}
\[-\sin 2\theta \sin(\phi_R + \phi_S) I \left[ b_1 \left( -\frac{\vec{R}_T}{2M_h} H_{1,LT}^\phi \right) \right] \]
\[+ \sin^2 \theta \sin(\phi_h - 2\phi_R - \phi_S) I \left[ \frac{\vec{R}_T \cdot P_{h\perp}}{M_h} h_1 \left( -\frac{\vec{R}_T}{|\vec{k}_T|} H_{1,OT}^\phi \right) \right] \]
\[-\sin \theta \sin(2\phi_h - \phi_R + \phi_S) I \left[ \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{k}_T^2}{2M_h^2} h_1 \left( -\frac{\vec{R}_T}{|\vec{k}_T|} H_{1,OT}^\phi \right) \right] \]
\[-\sin 2\theta \sin(2\phi_h - \phi_R + \phi_S) I \left[ \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{k}_T^2}{2M_h^2} h_1 \left( -\frac{M_h}{|\vec{k}_T|} H_{1,LT}^\phi \right) \right] \]
\[-\sin^2 \theta \sin(3\phi_h - 2\phi_R + \phi_S) I \left[ \frac{A(\vec{R}_T \cdot P_{h\perp})^3 - 3\vec{k}_T^2(\vec{R}_T \cdot P_{h\perp})}{2M_h^3} h_1 \left( -\frac{2M_h^2}{|\vec{k}_T|^2} H_{1,LT}^\phi \right) \right] \]
\[+ \cos 2\phi_h I \left[ \frac{2(\vec{R}_T \cdot P_{h\perp})(\vec{k}_T \cdot P_{h\perp}) - \vec{R}_T \cdot \vec{k}_T}{M_h} h_{1T} \left( \frac{1}{4} H_{1,OO}^{\phi^2} \right) \right] \]
\[+ \frac{1}{3} (3\cos^2 \theta - 1) \sin(3\phi_h - \phi_S) \times I \left[ \frac{A(\vec{R}_T \cdot P_{h\perp})^2(\vec{k}_T \cdot P_{h\perp}) - 2(\vec{R}_T \cdot P_{h\perp})(\vec{R}_T \cdot \vec{k}_T) - \vec{R}_T^2(\vec{k}_T \cdot P_{h\perp})}{2M^2 M_h} h_{1T} \left( \frac{3}{4} H_{1,LT}^\phi \right) \right] \]
\[-\sin \theta \sin(2\phi_h + \phi_R - \phi_S) I \left[ \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{R}_T^2}{2M^2} h_{1T} \left( -\frac{\vec{R}_T}{|\vec{k}_T|} H_{1,OT}^\phi \right) \right] \]
\[-\sin 2\theta \sin(2\phi_h + \phi_R - \phi_S) I \left[ \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{R}_T^2}{2M^2} h_{1T} \left( -\frac{M_h}{|\vec{k}_T|} H_{1,LT}^\phi \right) \right] \]
\[-\sin^2 \theta \sin(\phi_h + 2\phi_R - \phi_S) I \left[ \frac{2(\vec{R}_T \cdot \vec{k}_T)(\vec{R}_T \cdot P_{h\perp}) - (\vec{k}_T \cdot P_{h\perp})\vec{R}_T^2}{2M^2 M_h} h_{1T} \left( -\frac{\vec{R}_T}{|\vec{k}_T|} H_{1,OT}^\phi \right) \right] \]
\[+ \sin \theta \sin(4\phi_h - \phi_R - \phi_S) I \left[ \left( \vec{k}_T \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{R}_T^2}{4M^2 M_h^2} - 2(\vec{k}_T \cdot P_{h\perp}) \right) h_{1T} \left( -\frac{2M_h}{|\vec{k}_T|} H_{1,OT}^\phi \right) \right] \]
\[+ \sin 2\theta \sin(4\phi_h - \phi_R - \phi_S) I \left( \vec{k}_T \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{R}_T^2}{4M^2 M_h^2} - 2(\vec{k}_T \cdot P_{h\perp}) \right) h_{1T} \left( -\frac{M_h}{|\vec{k}_T|} H_{1,LT}^\phi \right) \]
\[+ \sin^2 \theta \sin(5\phi_h - 2\phi_R - \phi_S) I \left[ \left( \vec{k}_T \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{R}_T^2}{4M^2 M_h^2} - 2(\vec{k}_T \cdot P_{h\perp}) \right) h_{1T} \left( -\frac{2M_h}{|\vec{k}_T|^2} H_{1,LT}^\phi \right) \right] \]
\[+ \sin^2 \theta \sin(5\phi_h - 2\phi_R - \phi_S) I \left[ \left( \vec{k}_T \frac{2(\vec{R}_T \cdot P_{h\perp})^2 - \vec{R}_T^2}{4M^2 M_h^2} - 2(\vec{k}_T \cdot P_{h\perp}) \right) h_{1T} \left( -\frac{2M_h}{|\vec{k}_T|^2} H_{1,LT}^\phi \right) \right] \}. \]
\[ d^8 \sigma_{LT} = \sum_a \frac{\alpha_s^2 \alpha_w^2}{2 \pi s y^2} \lambda_c |S_T| C(y) \left\{ \cos(\phi_b - \phi_S) I \left[ \frac{\vec{p}_T \cdot \vec{P}_{b\perp}}{M} g_{1T} \left( \frac{1}{4} D_{1,OO} + \frac{3}{4} D_{1,OO}^* \right) \right] 
+ \frac{1}{3} (3 \cos^2 \theta - 1) \cos(\phi_b - \phi_S) I \left[ \frac{\vec{p}_T \cdot \vec{P}_{b\perp}}{M} g_{1T} \left( \frac{3}{4} D_{1,LL} \right) \right] 
- \sin 2\theta \cos(\phi_R - \phi_S) I \left[ \frac{(\vec{p}_T \cdot \vec{k}_{T})}{2 M M_b} g_{1T} \left( - \frac{M_b}{2 |\vec{k}_T|^2} D_{1,LT} \right) \right] \right\} \]

(C9)

The pure p-wave sector of the previous cross sections corresponds to the results of spin-1 production presented in Refs. [13, 38], once we apply the following identifications

\[ \frac{3}{4} D_{1,OO} = D_1, \quad \frac{3}{4} D_{1,LL} = D_{1,LL}, \quad \frac{M_b}{2 |\vec{k}_T|^2} D_{1,LT} = D_{1,LT}, \quad \frac{M_b}{2 |\vec{k}_T|^2} D_{1,TT} = D_{1,TT}, \quad \frac{3}{4} H_{1,OO}^+ = H_1^+, \quad \frac{3}{4} H_{1,LL}^+ = H_{1,LL}^+, \quad \frac{1}{2 M_b} H_{1,LT}^+ = H_{1,LT}^+, \quad \frac{1}{2 M_b} H_{1,TT}^+ = H_{1,TT}^+. \]

(C10)

Note, however, that while the functions on the left hand side contain a dependence on \( z \) as well as on the invariant mass \( M_b^2 \), the functions on the right hand side depend only on \( z \): it is required to assume that the spin-1 functions