Double distributions: Loose ends

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(Dated: March 26, 2022)

We point out that double distributions need not vanish at their boundary. Boundary terms do not change the ambiguity inherent in defining double distributions; instead, boundary conditions must be satisfied in order to switch between different decompositions. We analyze both the spin zero and spin one-half cases.

PACS numbers: 13.60.Fz, 12.38.Lg

QCD factorization provides the way to access information experimentally about the non-perturbative quark and gluon substructure of hadrons. In recent years, much attention has been generated by hard exclusive reactions such as deeply virtual Compton scattering, and the hard electroproduction of mesons. The non-perturbative structure functions entering these reactions are generalized parton distributions (GPDs), see the reviews [2]. Phenomenological modeling of GPDs is almost exclusively done utilizing the formalism of double distributions (DDs) [3]. The polynomiality property required of the Mellin moments of GPDs is elegantly explained by the formalism of double distributions [4, 5].

We point out that double distributions need not vanish at their boundary. Boundary terms do not change the ambiguity inherent in defining double distributions; instead, boundary conditions must be satisfied in order to switch between different decompositions. We analyze both the spin zero and spin one-half cases.

Boundary terms do not change the ambiguity inherent in defining double distributions; instead, boundary conditions must be satisfied in order to switch between different decompositions. We analyze both the spin zero and spin one-half cases.

While the gauge covariant derivative $\vec{D} = (\vec{D} - \vec{D}/2)$, where the gauge covariant derivative $\vec{D} = (\vec{D} - \vec{D})/2$, (…) denotes the symmetrization and trace subtraction performed on Lorentz indices, $\vec{P} = (P^\prime + P)/2$, $\Delta = P^\prime - P$, and $t = \Delta^2$. Often we treat the $t$-dependence as implicit below. Above, $T$-invariance restricts $A_{nk}(t) = 0$ for $k$ odd, and $B_{nk}(t) = 0$ for $k$ even. There is manifest arbitrariness in the twist-two form factors appearing in Eq. (1). The particular decomposition above can be used to define two DDs for the pion. These DDs are generating functions for the twist-two form factors

$$\left(\frac{A_{nk}}{B_{nk}}\right) = \int_{\beta,\alpha} \beta^{n-k} \alpha^k \left(\frac{F(\beta,\alpha)}{G(\beta,\alpha)}\right).$$

Above we have abbreviated the integration as $\int_{\beta,\alpha} = \int_{-1}^{1} d\beta \int_{-1}^{1} d\alpha$. As a consequence of $T$-invariance, the function $F(\beta,\alpha)$ is even in $\alpha$, while $G(\beta,\alpha)$ is odd.

Summing up the moments in Eq. (1), these DD functions then appear in matrix elements of the light-like separated quark bilinear operator $\mathcal{M}(\vec{P} \cdot z, \Delta \cdot z) \equiv \langle P|\overline{\psi}(-z/2)\gamma(z/2)|P\rangle$.

$$\mathcal{M}(\vec{P} \cdot z, \Delta \cdot z) = \int_{\beta,\alpha} e^{-i\alpha_z P z^i \delta \beta_z \Delta z^i} \times \left[2\theta(z) F(\beta,\alpha) - \Delta \cdot z G(\beta,\alpha)\right],$$

where $z^2 = 0$. Now we define the pion GPD

$$H(x, \xi) = \frac{1}{2} \int \frac{dz}{2\pi} e^{iz\vec{P}^\prime z} \mathcal{M}(\vec{P} \cdot z, \Delta \cdot z),$$

with the usual definition $\xi = -\Delta^+ / 2P^\prime^+$. Any other choice of generating functions for the moments must lead to the same GPD. To expose this ambiguity, we follow [6]. Integration of Eq. (2) by parts produces surface terms that in general do not vanish. We are careful about this point because previous models [6, 7] indeed have non-vanishing contributions on the boundary $|\beta| = 1 - |\alpha|$. Mathematically the DDs must vanish only at the corners of support: $\delta(\alpha) \delta(|\beta| - 1)$, $\delta(\beta) \delta(|\alpha| - 1)$, else the form factors in Eq. (2) do not fall off in the space of moments. Physically the DDs’ vanishing at the first set of corners is tied via Eq. (1) to the vanishing of the GPDs at $x = \pm 1$, which is known from perturbative QCD [8]. Vanishing of the DDs at the second set of corners implies the continuity of GPDs at the crossover ($x = \xi$), which in turn ensures factorization. Boundary contributions are thus not ruled out and the DDs appear to be loose at the ends. Phenomenological parameterizations of DDs must generally include such terms. The complete result of integrating by parts
can be expressed in the form
\[ \mathcal{M}(\bar{P} \cdot z, \Delta \cdot z) = -2i \int_{\beta,\alpha} N(\beta, \alpha) e^{-i\beta \bar{P} \cdot z + i\alpha \Delta \cdot z/2} + 8i \int_0^1 d\alpha \cos(\alpha \Delta \cdot z/2) \times \left\{ \cos[(1-\alpha)\bar{P} \cdot z] S^+(\alpha) - i \sin[(1-\alpha)\bar{P} \cdot z] S^-(\alpha) \right\}, \] (5)

where \( N(\beta, \alpha) = \frac{\partial}{\partial \beta} F(\beta, \alpha) + \frac{\partial}{\partial \alpha} G(\beta, \alpha) \) and \( S^\pm(\alpha) = F^\pm(1-\alpha, \alpha) + G^\pm(1-\alpha, \alpha) \). In deriving Eq. (5), we made use of the \( \alpha \)-symmetry of DDs and additionally have used \( \pm \)-distributions \( (F^\pm \text{ and } G^\pm) \), which are even and odd functions of \( \beta \), respectively.

Consider an arbitrary potential function \( \chi(\beta, \alpha) \) which is odd with respect to \( \alpha \). We can decompose \( \chi \) in terms of its even and odd parts with respect to \( \beta \), namely \( \chi^\pm(\beta, \alpha) \) where \( \chi^\pm(-\beta, \alpha) = \pm \chi^\pm(\beta, \alpha) \). The potential function can be used to generate the DD transformation
\[ \begin{pmatrix} F^\pm(\beta, \alpha) \\ G^\pm(\beta, \alpha) \end{pmatrix} \rightarrow \begin{pmatrix} F^\pm(\beta, \alpha) + \frac{\partial}{\partial \beta} \chi^\pm(\beta, \alpha) \\ G^\pm(\beta, \alpha) - \frac{\partial}{\partial \alpha} \chi^\mp(\beta, \alpha) \end{pmatrix}. \] (6)

Notice this transformation preserves both the \( \alpha \)- and \( \beta \)-symmetry of the \( \pm \)-distributions. Next we observe Eq. (5) is invariant under this transformation provided the potential also satisfies two boundary conditions
\[ \frac{\partial}{\partial \alpha} \chi^\pm(\beta, \alpha) \bigg|_{\beta=1-\alpha} = \frac{\partial}{\partial \beta} \chi^\pm(\beta, \alpha) \bigg|_{\beta=1-\alpha}. \] (7)

Above both \( \alpha \) and \( \beta \) are positive and the corresponding boundary conditions for the full range of DD variables can be found trivially due to the symmetry properties. Without the boundary conditions Eq. (7) on the potential function, the transformation generated by Eq. (6) is that of \( \bar{P} \).

To convert the DDs in Eq. (3) to the Polyakov-Weiss “gauge” \( \bar{P} \), in which there is only an \( F \)-type double distribution and \( D \)-term, we use the potential specified by
\[ \chi^\pm_0(\beta, \alpha) = -\frac{1}{2} \left[ \int_{-1+|\alpha|}^\beta d\beta' G^\mp(\beta', \alpha) - \int_\beta^{1-|\alpha|} d\beta' G^\mp(\beta', \alpha) - \frac{1}{2} (1 \mp 1) \text{sign}(\beta) D(\alpha) \right], \] (8)

where \( D(\alpha) \) is the \( D \)-term given by
\[ D(\alpha) = \int_{-1+|\alpha|}^{1-|\alpha|} d\beta G^\mp(\beta, \alpha). \] (9)

One can verify that the potential specified by Eq. (8) satisfies the boundary conditions Eq. (7) provided one also assumes the GPD is continuous, i.e., \( D(\pm 1; t) = 0 \). Under the transformation generated by Eq. (8), the bilocal matrix element reads
\[ \mathcal{M}(\bar{P} \cdot z, \Delta \cdot z) = \int_{\beta,\alpha} e^{-i\beta \bar{P} \cdot z + i\alpha \Delta \cdot z/2} \times \left[ 2\bar{P} \cdot z F_0(\beta, \alpha) - \Delta \cdot z \delta(\beta) D(\alpha) \right], \] (10)

where the resulting DD \( F_0(\beta, \alpha; t) \) is given by
\[ F_0(\beta, \alpha) = F(\beta, \alpha) + \frac{\partial}{\partial \alpha} \left[ \chi^+_0(\beta, \alpha) + \chi^-_0(\beta, \alpha) \right]. \] (11)

A perhaps more interesting choice of gauge is what we call the Drell-Yan gauge. It is specified by the potential function
\[ \chi^\pm_1(\beta, \alpha) = -\frac{1}{2} \left[ \int_{-1+|\beta|}^\alpha d\alpha' F^\mp(\beta, \alpha') - \int_{-1+|\beta|}^{1-|\beta|} d\alpha' F^\mp(\beta, \alpha') - \text{sign}(\alpha) D^\pm(\beta) \right], \] (12)

where the \( D \)-term \( D(\beta) \) is given by \( D(\beta) = D^+(\beta) + D^-\beta(\beta) \), with
\[ D^\pm(\beta) = \int_{-1+|\beta|}^{1-|\beta|} d\alpha F^\pm(\beta, \alpha). \] (13)

Again one can verify that the potential specified by Eq. (12) satisfies the boundary conditions Eq. (7). Under this transformation, the bilocal matrix element reads
\[ \mathcal{M}(\bar{P} \cdot z, \Delta \cdot z) = \int_{\beta,\alpha} e^{-i\beta \bar{P} \cdot z + i\alpha \Delta \cdot z/2} \times \left[ 2\bar{P} \cdot z \delta(\alpha) D(\beta) - \Delta \cdot z G_1(\beta, \alpha) \right], \] (14)

where the resulting DD \( G_1(\beta, \alpha) \) is given by
\[ G_1(\beta, \alpha) = G(\beta, \alpha) - \frac{\partial}{\partial \beta} \left[ \chi^+_1(\beta, \alpha) + \chi^-_1(\beta, \alpha) \right]. \] (15)

The Drell-Yan gauge is particularly interesting from the perspective of GPDs. Inserting Eq. (14) into the definition of the pion GPD Eq. (11), we have (now reinstating the \( t \)-dependence)
\[ H(x, \xi, t) = D(\beta, \xi) + \xi \int_{\beta,\alpha} \delta(x - \beta - \xi \alpha) G_1(\beta, \alpha; t), \] (16)

Notice the contribution to the GPD from the \( D \)-term is independent of \( \xi \) and the reduction relations are simply
\[ f_1(x) = H(x, 0, 0) = D(x, 0), \] (17)

for the quark distribution and
\[ F(t) = \int dx H(x, \xi, t) = \int dx D(x, t), \] (18)
for the pion form factor. Thus the $\bar{D}$-term has a simple physical interpretation as the $x$-integrand of the form factor calculated in the Drell-Yan frame ($\xi = 0$). Changing to a frame where $\xi \neq 0$, we have Eq. (19). Note, the second term in the GPD proportional to $\xi$ contributes nothing to the form factor because of Lorentz invariance. This is maintained by the $\alpha$-symmetry of the DD. The GPD, by contrast, is manifestly affected by the change in frame. At fixed $t$, the change to $\xi \neq 0$ requires a dynamical light-front rotation. Despite claims [3], the $\bar{D}$-term is all that one can learn about DDs from the Drell-Yan expression for the form factor [10].

To describe the pion GPD, we have resorted to using two DDs $F$ and $G$ because these are encountered in actual calculations [7]. The GPD, however, can be viewed as the second term in the GPD proportional to $\chi$ [7]. To see this representation one must use a non-trivial calculation [7]. The GPD, however, can be viewed as the second term in the GPD proportional to $\chi$ [7].

The boundary conditions are met provided

$$f^\pm(1 - \alpha, \alpha) = F_o^\pm(1 - \alpha, \alpha),$$

where $F_o$ is given in Eq. (11).

Now we address the ambiguities of proton DDs. The non-diagonal proton matrix elements of twist-two operators can be decomposed into form factors $A_{nk}(t)$, $B_{nk}(t)$ and $C_{nk}(t)$

$$\langle P', \lambda | \bar{\psi}(0) \gamma^\mu \bar{D}^\mu_i \cdots \bar{D}^\mu_n \psi(0) | P, \lambda \rangle \big|_{x=0, \xi=0} = \frac{1}{2} \bigg\{ \sum_{\alpha} \int_{-\infty}^{\infty} \frac{dz}{2\pi} e^{iz(x-\xi) + i\Delta \cdot z} M_{\lambda}(x, \xi, z) \bigg\} u_{\chi}(P) \bigg|_{P'}$$

$$= \frac{\pi\lambda(P')}{2} \sum_{n=0}^{\infty} \frac{n!}{k!(n-k)!} \bigg\{ \gamma^\mu A_{nk}(t) + \frac{i\sigma^\mu\Delta_\nu}{2M} B_{nk}(t) - \frac{\lambda^\mu}{4M} C_{nk}(t) \bigg\} u_{\chi}(P)$$

$$\times \bar{D}^\mu_1 \cdots \bar{D}^\mu_{n-k} \left( \frac{\Delta}{\bar{2}} \right)^{\mu_n-k+1} \cdots \left( \frac{\Delta}{\bar{2}} \right)^{\mu_n}.$$  

The above decomposition can be used to define three double distributions as generating functions for the twist-two form factors

$$\begin{pmatrix} A_{nk} \\ B_{nk} \\ C_{nk} \end{pmatrix} = \int_{\beta, \alpha} \beta^{\nu-k} \alpha^k \left( F(\beta, \alpha) \right) \left( K(\beta, \alpha) \right) \left( G(\beta, \alpha) \right).$$

$T$-invariance implies the functions $F(\beta, \alpha)$ and $K(\beta, \alpha)$ are even in $\alpha$, while $G(\beta, \alpha)$ is odd. Eq. (22) represents a “physical” gauge for proton DDs since these are encountered in actual calculations [7].

Summing up the moments in Eq. (22), the DD functions then appear in matrix elements of the light-like separated quark bilinear operator $M_{\lambda}(P', z, \Delta \cdot z) = \langle P', \lambda | \bar{\psi}(z/2) \gamma(z/2) | P, \lambda \rangle$,

$$\mathcal{M}_{\lambda}(P', z, \Delta \cdot z) = \int_{\beta, \alpha} e^{-i\beta z + i\alpha \Delta \cdot z} M_{\lambda}(P') \times \left( \mathcal{C}_{\lambda}(P') + \frac{i\sigma^\mu\Delta_\nu}{2M} K(\beta, \alpha) - \frac{\Delta \cdot z}{4M} G(\beta, \alpha) \right) u_{\chi}(P).$$

Now we define the light-cone correlation function

$$\mathcal{M}_{\lambda}(x, \xi, z) = \frac{1}{2} \int \frac{dz}{2\pi} e^{ixz + i\Delta \cdot z} \mathcal{M}_{\lambda}(P', z, \Delta \cdot z).$$

This correlation function can be written in terms of the two independent GPDs $H(x, \xi)$ and $E(x, \xi)$

$$\mathcal{M}_{\lambda}(x, \xi) = \frac{1}{2} \bigg\{ \gamma^+ H(x, \xi) + i\sigma^{\mu
u} \Delta_\nu E(x, \xi) \bigg\} u_{\chi}(P).$$

Inserting the DD decomposition Eq. (22) into the correlator in Eq. (26), we can express the GPDs as projections of the DDs

$$\begin{pmatrix} H(x, \xi) \\ E(x, \xi) \end{pmatrix} = \int_{\beta, \alpha} \delta(x - \beta - \xi\alpha) \begin{pmatrix} F(\beta, \alpha) + \xi G(\beta, \alpha) \\ K(\beta, \alpha) + \xi G(\beta, \alpha) \end{pmatrix},$$

from which we can view the $\xi$-dependence of GPDs as arising from different slices of Lorentz invariant DDs. Due to the symmetry of the DDs with respect to $\alpha$, the GPDs $H(x, \xi, t)$ and $E(x, \xi, t)$ are both even functions of the skewness parameter $\xi$.

Utilizing the Gordon identities, we can rewrite the bilocal matrix element in Eq. (22) as

$$\mathcal{M}_{\lambda}(P', z, \Delta \cdot z) = \frac{\pi\lambda(P')}{2M} \int_{\beta, \alpha} e^{-i\beta z + i\alpha \Delta \cdot z} \times \left( 2\bar{D} \cdot z \right) \gamma_5 G_M(\beta, \alpha) \bigg|_{P'} u_{\chi}(P),$$

$$- i\epsilon^{\mu\nu\alpha\beta} z_\mu \bar{D}_\alpha \gamma_5 G_M(\beta, \alpha) \bigg|_{P'} u_{\chi}(P).$$


where we have defined new double distributions

\[ G_E(\beta, \alpha) = F(\beta, \alpha) + \frac{t}{4M^2}K(\beta, \alpha) \]  
\[ G_M(\beta, \alpha) = F(\beta, \alpha) + K(\beta, \alpha) \]  
\[ \tilde{G}(\beta, \alpha) = \frac{1}{2}(1-t/4M^2)G(\beta, \alpha) \]  

in analogy with the Sachs electric and magnetic form factors. After integrating by parts, we have

\[ \mathcal{M}^{\lambda, \lambda'}(\mathbf{P} \cdot z, \Delta \cdot z) = \frac{i\pi x(P')}{2M - \frac{t}{4M^2}} \int_{\beta, \alpha} e^{-i\beta(\mathbf{P} \cdot z + i\alpha \Delta \cdot z)/2} \]

\[ \times \left[ 2N(\beta, \alpha) + e^{i\alpha \beta} z_\mu \Delta_\nu \gamma_\nu \gamma_5 G_M(\beta, \alpha) \right] \]

\[ - 8 \int_0^1 d\alpha \cos(\alpha \Delta \cdot z/2) \left\{ \cos \left[ (1 - \alpha)(\mathbf{P} \cdot z) \right] S^+(\alpha) \right. \]

\[ - \left. i \sin \left[ (1 - \alpha)(\mathbf{P} \cdot z) \right] S^-(\alpha) \right\} u_\lambda(P), \]  

where \( N(\beta, \alpha) = \partial_\beta G_E(\beta, \alpha) + \partial_\alpha \tilde{G}(\beta, \alpha) \), and the boundary terms \( S^+(\alpha) = G_E^+(1 - \alpha, \alpha) + \tilde{G}^+(1 - \alpha, \alpha) \).

Now we define a potential function \( \chi(\beta, \alpha) \) as above. The expression Eq. (32) is invariant under the transformation

\[ \begin{pmatrix} G_E^+(\beta, \alpha) \\ G_M^+(\beta, \alpha) \\ \tilde{G}^+(\beta, \alpha) \end{pmatrix} \rightarrow \begin{pmatrix} G_E^+(\beta, \alpha) + \frac{\partial}{\partial \alpha} \chi^+(\beta, \alpha) \\ G_M^+(\beta, \alpha) + \frac{\partial}{\partial \alpha} \chi^+(\beta, \alpha) \\ \tilde{G}^+(\beta, \alpha) - \frac{\partial}{\partial \alpha} \chi^+(\beta, \alpha) \end{pmatrix}, \]  

provided the boundary conditions on \( \chi^\pm \) Eq. (30) are satisfied. Translating the transformation Eq. (33) into the original DDs, we have invariance under

\[ \begin{pmatrix} F(\beta, \alpha) \\ K(\beta, \alpha) \\ G(\beta, \alpha) \end{pmatrix} \rightarrow \begin{pmatrix} F(\beta, \alpha) + \frac{\partial}{\partial \alpha} \chi(\beta, \alpha) \\ K(\beta, \alpha) - \frac{\partial}{\partial \alpha} \chi(\beta, \alpha) \\ G(\beta, \alpha) - \frac{\partial}{\partial \alpha} \chi(\beta, \alpha) \end{pmatrix}. \]  

As with the pion case, one can convert the proton DDs to Polyakov-Weiss gauge or Drell-Yan gauge using the potentials \( \chi_E^+(\beta, \alpha) \) and \( \chi_M^+(\beta, \alpha) \), respectively. Additionally two independent DDs for the proton can be unmasked using the transformation generated by \( \chi^+_F(\beta, \alpha) \). Each of these transformations for the proton case is understood under the replacement \( \{ F(\beta, \alpha) \rightarrow G_E(\beta, \alpha) \} \) and \( \{ G(\beta, \alpha) \rightarrow \tilde{G}(\beta, \alpha) \} \). In Polyakov-Weiss gauge we recover the result of Eq. (12) for the proton. In Drell-Yan gauge, we generate an additive \( \xi \)-independent contribution to the combination of GPDs \( H(x, \xi) + i\epsilon \gamma_5 E(x, \xi) \); this is the \( x \)-inTEGRAND of the Sachs electric form factor in the Drell-Yan frame. The \( \xi \) dependence is contained in a term analogous to the one in Eq. (12) and characterizes the dynamical light-front rotation to \( \xi \neq 0 \). In the minimal gauge generated by \( \chi^+_F(\beta, \alpha) \), we see there are only two underlying independent DDs.

There are two additional Drell-Yan type gauges for the proton. Using \( \chi^+_F(\beta, \alpha) \) as appears in Eq. (12), we generate an additive \( \xi \)-independent contribution to \( H(x, \xi) \). This is the \( x \)-inTEGRAND of the Dirac form factor calculated in the Drell-Yan frame. Alternately we can use Eq. (12) under the replacement \( \{ F(\beta, \alpha) \rightarrow -K(\beta, \alpha) \} \) to generate the \( \xi \)-independent, \( x \)-inTEGRAND of the Pauli form factor in the Drell-Yan frame, which contributes to \( E(x, \xi) \). With Eq. (34), we cannot simultaneously find Drell-Yan contributions to GPDs from both Dirac and Pauli form factors. This limitation arises from Eq. (33), as we cannot generate an additive contribution to \( H(x, \xi) + E(x, \xi) \) from the \( x \)-inTEGRAND of the Sachs magnetic form factor. Thus Eq. (34) must contain additional freedom.

Above we have seen that DDs do not necessarily vanish at their boundary and have investigated the consequences of boundary terms on the ambiguity inherent to DDs. The potential function that generates a transformation of DDs must satisfy boundary conditions. We have carried out this analysis for both the spin zero and spin one-half cases. Inclusion of boundary terms into model DDs is necessary for general phenomenological parameterizations of GPDs.

We thank W. Detmold and G. A. Miller for useful comments. This work was funded by the U. S. Department of Energy, grant: DE-FG03 – 97ER41014.

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