Generalized Parton Distributions in the Deuteron

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We introduce generalized quark and gluon distributions in the deuteron, which can be measured in exclusive processes like deeply virtual Compton scattering and meson electroproduction. We discuss the basic properties of these distributions and point out how they probe the interplay of nucleon and parton degrees of freedom in the deuteron wave function.

Introduction.—The partonic structure of the deuteron has been explored in terms of the parton distributions accessible in deep inelastic scattering [1] and in terms of the form factors measured in elastic lepton-deuteron processes [2,3]. It is natural to ask what can be learned from generalized parton distributions (GPDs), introduced not long ago in [4,5]. For the nucleon case it has been shown that these quantities contain unique information about the dynamics of quarks and gluons in QCD bound states, beyond what can be unraveled from ordinary parton distributions and form factors. Here we extend these studies to the case

of the deuteron, with the aim of providing the theoretical framework to analyze and interpret present and future measurements with deuteron targets. We restrict ourselves to parton distributions of twist two, and to the parton helicity conserving sector, which is relevant in most phenomenological applications. Quark and gluon helicity flip GPDs can be treated with the same methods.

Generalized quark distributions for the deuteron.—As in the nucleon case, the GPDs for the deuteron are defined through nondiagonal matrix elements of quark-antiquark operators on the light cone. Their general decomposition can be written as

$$V_{\lambda'\lambda} = \int \frac{d\kappa}{2\pi} e^{ix\kappa P\cdot n} \langle p', \lambda' | \bar{\psi}(-\kappa n)\gamma \cdot n\psi(\kappa n) | p, \lambda \rangle = \sum_i \epsilon'^{*\beta} V_{\beta\alpha}^{(i)} \epsilon^\alpha H_i(x,\xi,t), \tag{1}$$

$$A_{\lambda'\lambda} = \int \frac{d\kappa}{2\pi} e^{ix\kappa P\cdot n} \langle p', \lambda' | \bar{\psi}(-\kappa n)\gamma \cdot n\gamma_5\psi(\kappa n) | p, \lambda \rangle = \sum_i \epsilon'^{*\beta} A_{\beta\alpha}^{(i)} \epsilon^\alpha \tilde{H}_i(x,\xi,t), \tag{2}$$

where n is a lightlike four-vector, $n^2 = 0$. The incoming and outgoing deuterons, respectively, have momenta p, p', helicities λ, λ', and polarization vectors $\epsilon = \epsilon(p,\lambda)$ and $\epsilon' = \epsilon'(p',\lambda')$. We write $P = p + p'$ and $\Delta = p' - p$, and choose Ji's variables x, $\xi = -(\Delta \cdot n)/(P \cdot n)$, $t = \Delta^2$ as arguments of the GPDs H_i and $\tilde{H}_i$. The tensors $V^{(i)}$ and $A^{(i)}$ depend on the four-vectors p, p', and n. One need keep only tensors which do not vanish when

contracted with ϵ_α, ϵ'_β, given the orthogonality conditions $\epsilon \cdot p = \epsilon' \cdot p' = 0$ of the polarization vectors. With the constraints from parity invariance we find that the $V_{\beta\alpha}^{(i)}$ are linear combinations of the five tensor structures

$$\{g_{\beta\alpha}, p_\beta n_\alpha, n_\beta p'_\alpha, p_\beta p'_\alpha, n_\beta n_\alpha\}. \tag{3}$$

Similarly, the $A_{\beta\alpha}^{(i)}$ are linear combinations of the seven tensors

$$\{\epsilon_{\mu\nu\beta\alpha} p^\mu p'^\nu, \epsilon_{\mu\nu\beta\alpha} n^\mu p^\nu, \epsilon_{\mu\nu\beta\alpha} n^\mu p'^\nu, \epsilon_{\mu\nu\rho\beta} p^\mu p'^\nu n^\rho n_\alpha, \epsilon_{\mu\nu\rho\beta} p^\mu p'^\nu n^\rho p'_\alpha,$$
$$\epsilon_{\mu\nu\rho\alpha} p^\mu p'^\nu n^\rho n_\beta, \epsilon_{\mu\nu\rho\alpha} p^\mu p'^\nu n^\rho p_\beta\}. \tag{4}$$

Using the Schouten identities [6] one can show that only four out of these seven are linearly independent. The first x moments of generalized parton distributions are elastic form factors. To keep the corresponding relations simple we take among the tensors $V^{(i)}$, $A^{(i)}$ those which appear in the conventional form factor decomposition of the vector and axial currents [7]:

$$\langle p', \lambda' | \bar{\psi}(0)\gamma^\mu \psi(0) | p, \lambda \rangle = -G_1(t)(\epsilon'^* \cdot \epsilon)P^\mu + G_2(t)[\epsilon^\mu(\epsilon'^* \cdot P) + \epsilon'^{*\mu}(\epsilon \cdot P)] - G_3(t)(\epsilon \cdot P)(\epsilon'^* \cdot P)\frac{P^\mu}{2M^2},$$
$$\tag{5}$$
$$\langle p', \lambda' | \bar{\psi}(0)\gamma^\mu \gamma_5\psi(0) | p, \lambda \rangle = -i\tilde{G}_1(t)\epsilon^\mu{}_{\alpha\beta\gamma}\epsilon'^{*\alpha}\epsilon^\beta P^\gamma + i\tilde{G}_2(t)\epsilon^\mu{}_{\alpha\beta\gamma}\Delta^\alpha P^\beta \frac{\epsilon^\gamma(\epsilon'^* \cdot P) + \epsilon'^{*\gamma}(\epsilon \cdot P)}{M^2},$$

where our convention for the antisymmetric tensor is $\epsilon_{0123} = 1$ and M is the deuteron mass. The matrix elements are here defined flavor by flavor; to get the conventional form factors, one must weight with electromagnetic or weak charges

and sum over flavors. For the matrix elements of the nonlocal operators we define

$$V_{\lambda'\lambda} = -(\epsilon'^* \cdot \epsilon)H_1 + \frac{(\epsilon \cdot n)(\epsilon'^* \cdot P) + (\epsilon'^* \cdot n)(\epsilon \cdot P)}{P \cdot n} H_2 - \frac{(\epsilon \cdot P)(\epsilon'^* \cdot P)}{2M^2} H_3$$

$$+ \frac{(\epsilon \cdot n)(\epsilon'^* \cdot P) - (\epsilon'^* \cdot n)(\epsilon \cdot P)}{P \cdot n} H_4 + \left\{4M^2 \frac{(\epsilon \cdot n)(\epsilon'^* \cdot n)}{(P \cdot n)^2} + \frac{1}{3}(\epsilon'^* \cdot \epsilon)\right\}H_5 ,$$

$$A_{\lambda'\lambda} = -i \frac{\epsilon_{\mu\alpha\beta\gamma} n^\mu \epsilon'^{*\alpha} \epsilon^\beta P^\gamma}{P \cdot n} \tilde{H}_1 + i \frac{\epsilon_{\mu\alpha\beta\gamma} n^\mu \Delta^\alpha P^\beta}{P \cdot n} \frac{\epsilon^\gamma(\epsilon'^* \cdot P) + \epsilon'^{*\gamma}(\epsilon \cdot P)}{M^2} \tilde{H}_2$$

$$+ i \frac{\epsilon_{\mu\alpha\beta\gamma} n^\mu \Delta^\alpha P^\beta}{P \cdot n} \frac{\epsilon^\gamma(\epsilon'^* \cdot P) - \epsilon'^{*\gamma}(\epsilon \cdot P)}{M^2} \tilde{H}_3 + i \frac{\epsilon_{\mu\alpha\beta\gamma} n^\mu \Delta^\alpha P^\beta}{P \cdot n} \frac{\epsilon^\gamma(\epsilon'^* \cdot n) + \epsilon'^{*\gamma}(\epsilon \cdot n)}{P \cdot n} \tilde{H}_4 .$$

$$(6)$$

Since it is determined by the quark operators in Eqs. (1) and (2), the Q^2 evolution of the GPDs H_i and $\tilde{H}_i$ is exactly the same as for spin 1/2 targets, worked out in [4,5,8,9].

Time reversal properties and sum rules.—The action of the time reversal operator on the matrix element (1) leads to the relation

$$\epsilon'^{*\beta} \epsilon^\alpha V^{(i)}_{\beta\alpha}(P, \Delta, n) H_i(x, \xi, t) = \epsilon^{*\beta} \epsilon'^\alpha V^{(i)}_{\beta\alpha}(P, -\Delta, n)$$
$$\times H_i(x, -\xi, t), \quad (7)$$

where we have made explicit the dependence of the tensors $V^{(i)}$ on the relevant four-vectors. Taking the complex conjugate of Eq. (1) we get

$$\epsilon'^\beta \epsilon^{*\alpha} V^{*(i)}_{\beta\alpha}(P, \Delta, n) H_i^*(x, \xi, t) = \epsilon^{*\beta} \epsilon'^\alpha V^{(i)}_{\beta\alpha}(P, -\Delta, n)$$
$$\times H_i(x, -\xi, t). \quad (8)$$

For the axial vector case, we obtain relations analogous to (7) and (8) by replacing $V^{(i)}$ with $A^{(i)}$ and H_i with $\tilde{H}_i$. Combining these conditions we find that all nine GPDs are real. However, their behavior under time reversal is not uniform and we have

$$H_i(x, \xi, t) = H_i(x, -\xi, t) \quad (i = 1, 2, 3, 5),$$
$$H_4(x, \xi, t) = -H_4(x, -\xi, t),$$
$$\tilde{H}_i(x, \xi, t) = \tilde{H}_i(x, -\xi, t) \quad (i = 1, 2, 4),$$
$$\tilde{H}_3(x, \xi, t) = -\tilde{H}_3(x, -\xi, t).$$
$$(9)$$

Note that in the nonforward case, time reversal invariance fixes the phase of the generalized parton distributions and determines their behavior under sign change of the skewedness parameter ξ but does *not* limit the number of GPDs [10]. Integrating $V_{\lambda'\lambda}$ and $A_{\lambda'\lambda}$ over x one obtains the local matrix elements (5) contracted with $n^\mu/(P \cdot n)$. Since the tensors that accompany the distributions H_i, $\tilde{H}_i$ in Eq. (6) are linearly independent, we obtain the sum rules

$$\int_{-1}^1 dx\, H_i(x, \xi, t) = G_i(t) \quad (i = 1, 2, 3),$$

$$\int_{-1}^1 dx\, \tilde{H}_i(x, \xi, t) = \tilde{G}_i(t) \quad (i = 1, 2),$$

$$(10)$$

$$\int_{-1}^1 dx\, H_4(x, \xi, t) = \int_{-1}^1 dx\, \tilde{H}_3(x, \xi, t) = 0,$$

$$\int_{-1}^1 dx\, H_5(x, \xi, t) = \int_{-1}^1 dx\, \tilde{H}_4(x, \xi, t) = 0.$$

The integrals over H_4, $\tilde{H}_3$ and H_5, $\tilde{H}_4$ do not correspond to form factors of the local vector or axial currents and therefore vanish. In the case of H_4 and $\tilde{H}_3$ this is due to time reversal constraints, whereas the definitions of H_5 and $\tilde{H}_4$ involve the tensor $n^\mu n^\nu/(P \cdot n)^2$, whose analog cannot appear in the decomposition of the local currents due to Lorentz invariance.

The forward limit.—Let us now study the forward limit of the GPDs, which defines the usual parton distributions. In the parton model, i.e., at leading twist and leading order in α_s there are three independent structure functions in deep inelastic scattering, F_1, b_1, g_1, whose probabilistic interpretation in terms of quark densities reads [1]

$$F_1(x) = \frac{1}{2} \sum_q e_q^2 \frac{q^1(x) + q^{-1}(x) + q^0(x)}{3} + \{q \to \bar{q}\},$$

$$b_1(x) = \frac{1}{2} \sum_q e_q^2\left[q^0(x) - \frac{q^1(x) + q^{-1}(x)}{2}\right] + \{q \to \bar{q}\},$$

$$(11)$$

$$g_1(x) = \frac{1}{2} \sum_q e_q^2[q_\uparrow^1(x) - q_\uparrow^{-1}(x)] + \{q \to \bar{q}\}.$$

Here $q_{\uparrow(\downarrow)}^\lambda(x)$ represents the probability to find a quark with momentum fraction x and positive (negative) helicity in a deuteron target of helicity λ. The unpolarized quark densities q^λ are defined as $q^\lambda(x) = q_\uparrow^\lambda(x) + q_\downarrow^\lambda(x)$. From parity one has $q_\uparrow^\lambda = q_\downarrow^{-\lambda}$. The densities for antiquarks are defined in analogy. Note that the probabilistic interpretation for F_1 and g_1 is similar to the one in the spin 1/2 case, whereas the function b_1 does not appear for spin 1/2 targets. In the forward limit the only structures in Eq. (6) that survive are those proportional to H_1, H_5, and $\tilde{H}_1$, because in that limit we have $\Delta = 0$ and $\epsilon \cdot P = \epsilon' \cdot P = 0$. Using the results for helicity amplitudes given below, one gets

$$H_1(x, 0, 0) = \frac{q^1(x) + q^{-1}(x) + q^0(x)}{3},$$

$$H_5(x, 0, 0) = q^0(x) - \frac{q^1(x) + q^{-1}(x)}{2}, \quad (12)$$

$$\tilde{H}_1(x, 0, 0) = q_\uparrow^1(x) - q_\uparrow^{-1}(x)$$

for $x > 0$. The corresponding relations for $x < 0$ involve the antiquark distributions at $-x$, with an overall minus sign in the expressions for H_1 and H_5. With Eq. (10) we

thus have

$$0 = \int_{-1}^{1} dx\, H_5(x,0,0)$$

$$= \int_{0}^{1} dx \left[q^0(x) - \frac{q^1(x) + q^{-1}(x)}{2} \right] - \{q \to \bar{q}\} \tag{13}$$

and recover the parton model sum rule $\int_0^1 b_1(x) = 0$ of [11], which was obtained under the assumption that the quark sea does not contribute to this integral.

Helicity amplitudes.—In the region $\xi < x < 1$ our GPDs can be represented in terms of amplitudes for the scattering of a quark on a deuteron [10], defined as

$$\epsilon^{(0)\mu} = \frac{1}{M}\left(p^\mu - \frac{2M^2}{1+\xi} \frac{n^\mu}{P \cdot n} \right),$$

$$\epsilon^{(1)\mu} = -\frac{1}{\sqrt{(1-\xi^2)(t_0-t)}}\left((1+\xi)p'^\mu - (1-\xi)p^\mu - \frac{\xi(t_0-t)-t_0}{\xi}\frac{n^\mu}{P\cdot n} \right), \tag{17}$$

$$\epsilon^{(2)\mu} = \frac{1}{\sqrt{(1-\xi^2)(t_0-t)}}\frac{2\epsilon^\mu{}_{\nu\alpha\beta}p'^\nu p^\alpha n^\beta}{P\cdot n},$$

where $t_0 = -4M^2\xi^2/(1-\xi^2)$. The vectors $\epsilon(0) = \epsilon^{(0)}$ and $\epsilon(\pm 1) = \mp(\epsilon^{(1)} \pm i\epsilon^{(2)})/\sqrt{2}$ then correspond to definite light-cone helicity [12]. This approximately coincides with usual helicity in frames where the deuteron moves fast, provided that $\mathrm{sgn}(p^3) = -\mathrm{sgn}(n^3)$. The polarizations for the outgoing deuteron are obtained from Eq. (17) by the exchange $p^\mu \leftrightarrow p'^\mu$, $\xi \leftrightarrow -\xi$, and an overall sign change for $\epsilon^{(1)}$ and $\epsilon^{(2)}$. With this we get

$$\mathcal{A}_{++,++} = \frac{H_1}{2} - \frac{H_5}{6} + \frac{DH_3}{2} + \frac{\tilde{H}_1}{2} + 2D(\tilde{H}_2 + \xi\tilde{H}_3),$$

$$\mathcal{A}_{0+,0+} = \frac{H_1}{2} - \xi H_4 + \frac{1}{3}\left(1 - \frac{3}{2}\xi^2\right)H_5 - \left(D - \frac{\xi^2}{1-\xi^2}\right)\left(H_1 - H_2 - \xi H_4 - \frac{1}{3}H_5\right)$$
$$\qquad - \left(D^2 - \frac{\xi^2}{(1-\xi^2)^2}\right)H_3,$$

$$\mathcal{A}_{-+,++} = -D[\tfrac{1}{2}H_3 + 2(\xi\tilde{H}_2 + \tilde{H}_3)], \tag{18}$$

$$\mathcal{A}_{0+,++} = \sqrt{\frac{D(1-\xi)}{2(1+\xi)}}\left[H_1 - \frac{1-\xi}{2}(H_2 - H_4) - \frac{1}{3}H_5 + \left(D - \frac{\xi}{1-\xi^2}\right)H_3 \right]$$
$$\qquad + \sqrt{2D(1-\xi^2)}\left[\frac{1}{4}[\tilde{H}_1 + (1-\xi)\tilde{H}_4] + \left(D - \frac{\xi}{1-\xi^2}\right)(\tilde{H}_2 + \tilde{H}_3) \right],$$

where $D = (t_0 - t)/(4M^2)$. The remaining amplitudes can be easily obtained from the relations (15), (16), and

$$\mathcal{A}_{\lambda'-\mu,\lambda-\mu} = \mathcal{A}_{\lambda'\mu,\lambda\mu}(\tilde{H}_i \to -\tilde{H}_i). \tag{19}$$

As required by angular momentum conservation, one gets a factor $\sqrt{t_0 - t}$ for each unit of helicity flip. Note that $H_{2,4}$ and $\tilde{H}_4$ appear only with longitudinal deuteron polar-

$$\mathcal{A}_{\lambda'\pm,\lambda\pm} = \tfrac{1}{2}(V_{\lambda'\lambda} \pm A_{\lambda'\lambda}), \tag{14}$$

with $\pm$ referring to the helicities of the quarks. With the constraints

$$\mathcal{A}_{-\lambda'-\mu,-\lambda-\mu} = (-1)^{\lambda'-\lambda}\mathcal{A}_{\lambda'\mu,\lambda\mu} \tag{15}$$

from parity invariance there are nine independent quark helicity conserving amplitudes. Since $\mathcal{A}_{0+,0+} = \mathcal{A}_{0-,0-}$ we have only four quark helicity dependent distributions $\tilde{H}_i$, compared with the five quark helicity independent H_i. Time reversal invariance gives

$$\mathcal{A}(x,\xi,t)_{\lambda\mu,\lambda'\mu} = (-1)^{\lambda'-\lambda}\mathcal{A}(x,-\xi,t)_{\lambda'\mu,\lambda\mu} \tag{16}$$

and thus does not further reduce the number of GPDs, as remarked above. To define the polarization of the incoming deuteron we introduce

ization and that the only GPDs appearing in double helicity flip amplitudes are H_3 and $\tilde{H}_{2,3}$. In the forward limit we have $\mathcal{A}(x,0,0)_{\lambda+,\lambda+} = q_\uparrow^\lambda(x)$, $\mathcal{A}(x,0,0)_{\lambda-,\lambda-} = q_\downarrow^\lambda(x)$ and find the relations (12).

Gluon distributions.—Let us turn to the gluon distributions in the deuteron. Instead of the matrix elements (1), (2), we now have

$$4\frac{n_\alpha n_\beta}{P \cdot n}\int \frac{d\kappa}{2\pi} e^{ix\kappa P \cdot n}\langle p',\lambda'|F^{\alpha\mu}(-\kappa n)F_\mu{}^\beta(\kappa n)|p,\lambda\rangle = \sum_i \epsilon'^{*\beta}V_{\beta\alpha}^{(i)}\epsilon^\alpha H_i^g(x,\xi,t),$$

$$-4i\frac{n_\alpha n_\beta}{P \cdot n}\int \frac{d\kappa}{2\pi} e^{ix\kappa P \cdot n}\langle p',\lambda'|F^{\alpha\mu}(-\kappa n)\tilde{F}_\mu{}^\beta(\kappa n)|p,\lambda\rangle = \sum_i \epsilon'^{*\beta}A_{\beta\alpha}^{(i)}\epsilon^\alpha \tilde{H}_i^g(x,\xi,t), \tag{20}$$

with $\tilde{F}^{\alpha\beta} = \frac{1}{2}\epsilon^{\alpha\beta\gamma\delta}F_{\gamma\delta}$. We take the same tensors $V^{(i)}$, $A^{(i)}$ as for quark distributions, given in Eq. (6). Note that the H_i^g are even and the $\tilde{H}_i^g$ odd in x. Their behavior under time reversal is the same as in (9) for the corresponding quark

distributions, and the definitions and expressions of the helicity amplitudes $\mathcal{A}^g_{\lambda'\mu,\lambda\mu}$ are also analogous to the quark case. The forward limit is now $\mathcal{A}^g(x,0,0)_{\lambda+,\lambda+} = xg_1^\lambda(x)$ and $\mathcal{A}^g(x,0,0)_{\lambda-,\lambda-} = xg_1^\lambda(x)$, with an extra factor x compared to the quark case.

Some phenomenology.—The deuteron GPDs can be accessed in hard exclusive processes such as deeply virtual Compton scattering in $ed \to ed\gamma$, and electroproduction $ed \to edM$ of a meson or a meson pair [13]. The relevant kinematical limit is that of large invariant momentum transfer Q^2 to the lepton at fixed x_B and t, where the Bjorken variable x_B is defined as in deep inelastic scattering. Factorization formulas are the same as for nucleon targets [5,14], with the appropriate replacement of the hadronic matrix elements, but the same hard scattering kernels. The Q^2 behavior of the amplitudes and the selection rules for photon and meson helicities also remain the same, since they depend on the hard-scattering process, not on the target spin. Electroproduction of a pseudoscalar meson selects the GPDs $\tilde{H}_i$, vector meson production involves the H_i and H_i^g, and all distributions appear in Compton scattering. Notice that the isosinglet nature of the deuteron simplifies the flavor structure of the GPDs and thus of the scattering amplitudes. One consequence is that pion exchange does not contribute to any of these processes, in contrast to the nucleon case, where it may give important contributions through the quark distribution $\tilde{E}$ [15]. In kinematics where the Bethe-Heitler process dominates in electroproduction $ed \to ed\gamma$, one can use the methods of [16] to study Compton scattering through the interference of the two processes. At sufficiently large Q^2 this interference term gives access to a linear combination of GPDs, weighted with the electromagnetic deuteron form factors $G_{1,2,3}(t)$.

Discussion.—To get a feeling for the physics of GPDs in the deuteron, consider the approximation where they are written as a convolution of the nucleon GPDs with the light-cone wave function ψ_{p+n} for a proton and a neutron in the deuteron [17,18]. The struck nucleon then has to absorb the entire momentum transfer, in particular, its plus component parametrized by ξ, and still fit into the final deuteron. Assuming for simplicity that ψ_{p+n} is only nonzero if the plus-momentum fraction of the proton in the deuteron is between $\frac{1}{2}(1-w)$ and $\frac{1}{2}(1+w)$, one finds that all H_i and $\tilde{H}_i$ vanish for $\xi > w$. The ξ dependence of the deuteron GPDs thus reflects the width of the wave function ψ_{p+n} in longitudinal momentum fraction. For ξ well above w they provide access to deuteron wave function components that cannot be described in terms of individual nucleons. To unravel such components in inclusive deep inelastic scattering at $x_B > 1$ has turned out to be difficult. Approaching this region from below, the struck quark has to take the entire momentum of a nucleon in the convolution picture, so that not only the deuteron wave function ψ_{p+n} becomes small but also the parton density in the nucleon. This is not the case for GPDs with their independent momentum variables x and ξ: at $\xi > w$ one

can still have any value for the plus-momentum fraction of the struck parton in the target.

We also note that, since in a convolution model b_1 requires a d wave component in ψ_{p+n} [1,17], the same holds for H_5. Double helicity flip amplitudes also need a d wave admixture: the helicity flip of a nucleon cannot exceed one unit, so that orbital angular momentum is necessary. Hence H_3 and $\tilde{H}_{2,3}$ involve the d wave part of ψ_{p+n}.

In this Letter we have focused on the case where the deuteron scatters elastically. One may extend our study to the case where it dissociates into a proton and neutron, or a more complicated hadronic system, introducing appropriate transition GPDs. Factorization still holds in the processes discussed above, provided that the invariant mass of the dissociative system is small compared with the hard scale Q^2.

Let us finally stress that for low enough ξ the generalized parton distributions for elastic deuteron transitions are by no means small; neither is there any suppressing factor in the cross sections. It should thus be possible to perform exclusive electroproduction experiments on a deuteron target where those on a nucleon are possible. We hope for a rich harvest of physics on this topic in the forthcoming years at existing facilities such as DESY and Jefferson Lab.

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