Proton imaging has become a key diagnostic for measuring electromagnetic fields in high-energy-density (HED) laboratory plasmas. Compared to other techniques for diagnosing fields, proton imaging is a measurement that can simultaneously offer high spatial and temporal resolution and the ability to distinguish between electric and magnetic fields without the protons perturbing the plasma of interest. Consequently, proton imaging has been used in a wide range of HED experiments, from...
inertial-confinement fusion to laboratory astrophysics. An overview is provided on the state of the art of proton imaging, including a discussion of experimental considerations like proton sources and detectors, the theory of proton-imaging analysis, and a survey of experimental results demonstrating the breadth of applications. Topics at the frontiers of proton-imaging development are also described, along with an outlook on the future of the field.

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studying HED plasmas with strong fields for the following reasons: (1) the protons are “stiff” enough that they experience only small deflections for typical field strengths, allowing the detected proton position to be related simply to the initial proton position in order to infer field strengths; (2) for many experiments, the protons traverse the experimental plasma on timescales that are short compared to dynamical timescales, providing a relatively static snapshot of the fields; (3) the proton images have high spatial resolution owing to (i) the small source size of laser-driven proton beams as well as (ii) their high laminarity; (4) the proton beam, being locally of much lower density than the probed plasma, does not perturb it; and (5) the dependence of the proton deflections on proton energy or geometry is different for electric and magnetic fields, enabling the contribution from each to be distinguished using different proton energies or probing from different directions. We note that proton imaging is also referred to as proton radiography or proton deflectometry in the literature, where the former can be used to describe the imaging of proton scattering and stopping from either density or electromagnetic fields and the latter is often used when one directly measures proton deflections with, for example, a mesh or grid. In this review, we focus primarily on proton deflections from electromagnetic fields rather than from collisions; however, because collisional scattering can have a non-negligible effect on proton images of HED experiments involving cold and/or dense plasmas, we consider its effect at various points in our review.

We illustrate the basic concept that underlies proton imaging with a schematic of a proton-imaging setup (shown in Fig. 1). Protons from a point source pass through the plasma of interest, are deflected by electromagnetic fields, and then travel ballistically to a detector, where they form an image of the field structures in the plasma. Inhomogeneous electromagnetic fields in the plasma plane differentially deflect protons with distinct incident trajectories, which in turn gives rise to inhomogeneous proton fluence on the detector. This allows the path-integrated strengths of the electromagnetic fields to be estimated by relating the proton-fluence variations on the detector to the displacement experienced by those protons as they pass through the fields in the plasma.

The diagnostic is typically configured in the paraxial limit, in which the characteristic scale $\ell_{\text{EM}}$ of electromagnetic fields in the plasma being probed is much smaller than the distance $r_s$ between the source and the plasma ($\ell_{\text{EM}} \ll r_s$), and in a point-projection geometry, in which the distance $r_d$ from the

![FIG. 1. Schematic of a proton-imaging setup.](image)

plasma to the detector greatly exceeds the path length $l_{\text{path}}$ of the protons through the plasma ($r_d \gg l_{\text{path}}$). Consequently, for sufficiently large proton energies (with characteristic deflection velocities that are much smaller than the incident velocities) the path-integrated electromagnetic-field strengths can be related to the deflection angle $\delta \alpha$ of a proton. Under these approximations and limiting to deflections along $\hat{z}$ without loss of generality, $\delta \alpha$ is given by (see Fig. 1)

$$\delta \alpha = \frac{e}{m_p v_p^2} \int_0^{l_{\text{path}}} ds \left[ E_x \left( \frac{v_p \times B}{c} \right)_x \right], \quad (1)$$

where $e$ is the elementary charge, $m_p$ is the mass of a proton, $c$ is the speed of light, $v_p$ ($v_p$) is the protons’ velocity (speed), and $E$ and $B$ are the electric and magnetic fields in the plasma, respectively. Here and in the rest of the review we express equations in centimeter-gram-second-units. The final position $d$ of the proton in the image plane at the detector will be

$$d = d_0 + \Delta d = M x_0 + r_d \delta \alpha, \quad (2)$$

where $x_0$ is the initial transverse position of the proton in the plasma, $d_0$ is the undeflected proton position in the detector plane accounting for magnification $M \equiv (r_s + r_d + l_{\text{path}})/r_s \approx (r_s + r_d)/r_s$, and $\Delta d$ is the displacement due to the deflection of protons by electromagnetic fields in the plasma. Thus, the path-integrated fields can be inferred from

$$\int_0^{l_{\text{path}}} ds \left[ E_x \left( \frac{v_p \times B}{c} \right)_x \right] = \frac{m_p v_p^2}{e} \frac{r_s + r_d}{r_s r_d} (x - x_0), \quad (3)$$

where $x = d/M$ is the deflected position rescaled to the plasma plane provided that the initial and final positions $x_0$ and $x$ of the protons are known. The salient problem, which is considered in Sec. III, is then inferring the displacement of the protons from the proton-fluence inhomogeneities that are directly measured.

A useful metric for classifying different types of proton-fluence inhomogeneities that can arise due to these proton displacements is the contrast parameter

$$\mu \equiv \frac{r_d \delta \alpha}{M \ell_{\text{EM}}} \sim \frac{\delta \Psi}{\Psi_0}, \quad (4)$$

where $\Psi_0$ is the mean proton fluence and $\delta \Psi$ is the magnitude of the inhomogeneities. For $\mu \ll 1$, the relation between the path-integrated fields and inhomogeneities is approximately linear, and the measured proton-fluence distribution is proportional to the path-integrated charge (for purely electrostatic fields) distribution or current-density (for purely magnetostatic fields) distribution, respectively. As $\mu$ increases, the proton-fluence distribution becomes spatially distorted compared to the path-integrated charge-density and current-density distributions, with regions of focused and defocused fluence; however, qualitatively the image is still similar to these density distributions. When $\mu$ becomes larger than some critical value $\mu_c \sim 1$, proton trajectories cross before reaching the detector, leading to the formation of so-called caustic
The first charged-particle-imaging experiments measuring electromagnetic fields in plasmas date back to the 1970s (Mendel and Olsen, 1975) and utilized accelerators as a source of ions. However, the long pulse length of ions from conventional accelerators and the difficulty of combining externally produced ion beams with experiments limited the application of this technique to HED plasmas. Not until the discovery of laser-driven, MeV proton sources was proton imaging regularly employed on HED facilities.

The development of multi-MeV, pointlike proton sources useful for proton imaging was first demonstrated two decades ago (Borghesi et al., 2001). The proton sources were generated by focusing high-intensity lasers onto thin foils; this generated MeV protons via a process called target normal sheath acceleration (TNSA), which was first described by Wilks et al. (2001). Radiographic film stacks (Borghesi et al., 2001) and allyl diglycol carbonate, Columbia resin no. 39 (CR-39) nuclear-track detectors (Clark et al., 2000; Maksimchuk et al., 2000) were both initially used to image the protons, but the low-fluence saturation limit of CR-39 and issues with data interpretation (Clark et al., 2006; Gaillard et al., 2006) led to its disuse for TNSA protons. Soon after these initial experiments, the first uses of TNSA-generated protons for measuring electromagnetic fields in HED plasmas were reported, with electric fields being characterized in ICF and laser-produced plasmas (Borghesi et al., 2001, 2002b). Meshes to directly measure the proton deflections were first added a few years later (Mackinnon et al., 2004).

Around the same time that TNSA proton sources were being developed, a second type of laser-driven proton source based on capsules filled with deuterium helium-3 (D3He) gas was being developed in connection with direct-drive ICF experiments (Li et al., 2002; Smalyuk et al., 2003). When imploded, these capsules emit ∼3 and ∼15 MeV protons as fusion by-products. A distinctive feature of D3He-capsule proton sources is their narrow energy spectra, which contrasts with the broadband proton energy spectra generated by TNSA. Compared to TNSA proton sources, the proton fluence from D3He sources is significantly lower, requiring the use of low-fluence CR-39 detectors (Séguin et al., 2003). In 2006, the use of a D3He proton source to image electromagnetic fields in laser-produced plasmas was first reported (Li et al., 2006b).

A key challenge of proton imaging is recovering the path-integrated electromagnetic fields based on the measured proton fluence. The first approach chronologically, taken shortly after the initial deployment of high-intensity laser sources, was the development of numerical forward models that take a known electromagnetic-field configuration and generate a synthetic proton-fluence image that can be compared to the measured image. Quantitative analysis of such a comparison allowed for the optimal choice of characteristic parameters of the proposed electromagnetic field. These initial modeling efforts were employed to measure electric fields using data from TNSA proton sources (Borghesi et al., 2003; Romagnani et al., 2005) and, with subsequent application, to determine electric and magnetic fields probed with D3He.

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TABLE I. Comparison of the typical proton-imaging source properties and characteristics. CPA, chirped pulse amplification.

<table>
<thead>
<tr>
<th>TNSA</th>
<th>D^3He</th>
</tr>
</thead>
<tbody>
<tr>
<td>Typical laser driver (energy, pulse width)</td>
<td>&gt; 50 J, ~50 fs</td>
</tr>
<tr>
<td>Facility required</td>
<td>High-energy CPA laser</td>
</tr>
<tr>
<td>Typical target</td>
<td>Flat, metallic foil, ~10–25 μm thick</td>
</tr>
<tr>
<td>Source size</td>
<td>~10 μm</td>
</tr>
<tr>
<td>Source time; cf. laser driver</td>
<td>Instantaneous</td>
</tr>
<tr>
<td>Proton temporal spread at source</td>
<td>~ps</td>
</tr>
<tr>
<td>Spectral characteristics</td>
<td>Maxwellian-like, up to ~60 MeV</td>
</tr>
<tr>
<td>Typical proton yield</td>
<td>10^{11}–10^{13} (total in the beam)</td>
</tr>
<tr>
<td>Proton directionality</td>
<td>Beam with ~30° divergence</td>
</tr>
<tr>
<td>Typical detector</td>
<td>RCF stack</td>
</tr>
</tbody>
</table>

sources (Li et al., 2006b). Analytical models relating electromagnetic fields to their proton images were also developed at around the same time (Borghesi et al., 2002b; Romagnani et al., 2005), but the first detailed discussion of the analytical theory of proton imaging was not published until the work of Kugland, Ryutov et al. (2012). Obtaining direct measurements of the fields required the development of techniques to extract proton deflections from the proton-fluence profiles. This was first done through proton deflectometry (Romagnani et al., 2005; Li et al., 2007; Petrazzo et al., 2009), in which a mesh placed between the proton source and detector provided a direct reference for how the protons were deflected. In many experiments, though, adding a mesh is not practical. For these cases, a variety of numerical inversion schemes were developed and first reported in 2017 (Bott et al., 2017; Graziani et al., 2017; Kasim et al., 2017).

II. EXPERIMENTAL TECHNIQUES

Proton imaging has developed significantly over the past two decades and is now commonly used at many HED experimental facilities. In this section we describe each component needed to perform the measurement. First, we discuss different proton sources and the methods for producing protons, as well as the properties of the protons generated. Second, we describe the standard detectors used to measure the protons and the trade-offs associated with each. Last, we discuss how the geometry of the experiment affects proton measurements and additional considerations when one designs proton-imaging setups.

A. Proton sources

There are two main types of proton sources that have been developed for proton-imaging experiments: (1) proton beams accelerated by a high-intensity laser through the so-called TNSA mechanism, and (2) protons produced from nuclear fusion reactions resulting from laser-driven implosions of D^3He-filled targets. In the following we review the general characteristics of these two sources, which differ significantly in terms of properties and capabilities. Table I comparatively summarizes the main properties of these sources.

1. Target normal sheath acceleration

Since the first reports of multi-MeV proton beams produced from laser-irradiated foils in 2000 (Clark et al., 2000; Maksimchuk et al., 2000; Snavely et al., 2000), proton acceleration has been one of the most active fields of research employing high-power, short-pulse lasers (Macchi, Borghesi, and Passoni, 2013). TNSA is the mechanism that has been most studied and has been widely employed for applications. TNSA was proposed as an interpretative framework (Hatchett et al., 2000; Wilks et al., 2001) of the multi-MeV proton observations reported by Snavely et al. (2000), obtained on the NOVA Petawatt laser at Lawrence Livermore National Laboratory (LLNL). The scheme typically employs mid-infrared (0.8–1 μm wavelength), multi-hundred-TW short-pulse (30 fs–10 ps pulse duration) laser systems that generate on-target intensities in the range of 10^{19}–10^{21} W cm^{-2}.

A schematic of the TNSA process is shown in Fig. 2. A high-intensity laser pulse interacts with a solid foil target with a thickness of around a few microns. At these intensities, the laser pulse, focused on the foil surface, can efficiently couple energy into relativistic electrons, mainly through ponderomotive processes [such as the J × B mechanism (Kruer and Estabrook, 1985)]. The average energy of the electrons is typically of the order of MeV, so their collisional range is much larger than the foil thickness, and they can propagate to the rear of the target. As the electrons expand into the vacuum they establish a space-charge field that ionizes the rear surface and drives the acceleration of ions from surface layers. While a limited number of energetic electrons will effectively leave the target (Link et al., 2011), most of the hot electrons are confined to within the target volume by the space charge and form a sheath extending by approximately a Debye length

$$\lambda_D = \sqrt{k_B T_{e,\text{hot}} / 4\pi n_{e,\text{hot}} e^2}$$

from the initially unperturbed rear...
surface, where \( k_B \) is Boltzmann’s constant and \( n_{e,\text{hot}} \) and \( T_{e,\text{hot}} \) are the density and temperature of the superthermal (hot) electrons. The electric field in the sheath is proportional to \( \left(n_{e,\text{hot}}T_{e,\text{hot}}\right)^{1/2} \) (Mora, 2003; Schreiber et al., 2006). For a typical interaction, the sheath field reaches amplitudes in the TV/m range. Under standard experimental conditions, contaminant layers (hydrocarbons, water, etc.) exist on the surface of any target (Allen et al., 2004). Therefore, protons are most efficiently accelerated by TNSA due to their favorable charge-to-mass ratio and shield other ion species from experiencing the strongest accelerating fields. This makes TNSA a robust, efficient, and easily implementable mechanism for accelerating protons.

The energy spectra of TNSA proton beams are broadband, typically with an exponential profile up to a high-energy cutoff; see Fig. 3(a). The highest TNSA energies reported are of the order of 85 MeV (Wagner et al., 2016), obtained with large PW-class laser systems, and available data generally show that at equal intensities longer pulses (\( \sim \)picosecond duration) containing more energy generally accelerate ions more efficiently than pulses with widths of tens of femtoseconds (Macchi, Borghesi, and Passoni, 2013). However, using state-of-the-art femtosecond systems and stringent control of the laser properties has recently allowed the energies of accelerated protons to be increased up to 70 MeV (Ziegler et al., 2021).

Reported scaling laws for the proton energies as a function of laser intensity vary from a ponderomotive \( I^{0.5} \) dependence for subpicosecond pulses (Macchi, Borghesi, and Passoni, 2013) to a near-linear dependence observed for ultrashort laser pulses over restricted intensity ranges (Zeil et al., 2010); see Figs. 3(b) and 3(c). Superponderomotive scaling for multi-kilojoule, multipicosecond lasers has also been reported (Flippo et al., 2007; Mariscal et al., 2019). Nevertheless, secondary factors such as target thickness, target material, target size, and laser contrast (Kaluza et al., 2004; Fuchs et al., 2007; Yogo et al., 2008; Schollmeier et al., 2015) also play an important role in TNSA accelerating energy performance.

![FIG. 2. Schematic of the main processes involved in the TNSA mechanism. (a) First, a laser prepulse impinges upon and heats a thin target to form a preplasma. The target contains layers of proton-rich hydrocarbons as common contaminants. (b) Second, the peak of the pulse arrives, efficiently heating electrons to relativistic temperatures. These electrons expand and propagate through the target. (c) Third, the hot electrons emerge into the vacuum and form an electron sheath with a strength of \( \sim TV/m \). This field ionizes the rear surface such that ions are accelerated to multi-MeV energies. Adapted from McKenna et al., 2006.](image-url)

![FIG. 3. (a) TNSA spectrum obtained on the NOVA Petawatt laser at LLNL, expressed in the number of protons per MeV (left scale). Adapted from Snavely et al., 2000. TNSA cutoff energies plotted against (b) laser intensity on target and (c) laser energy. The data are from a select number of experiments where a scan in laser energy was performed. Adapted from Zimmer et al., 2021.](image-url)
Having a sharp density interface at the rear target surface is key to efficient TNSA acceleration. For pulses with a duration longer than \(\sim 1\) ps, the rear target surface evolves before the electrons associated with the peak intensity arrive, limiting the maximum acceleration (Schollmeier et al., 2015; Campbell et al., 2019).

If the laser pulse has a significant “prepulse,” or energy arriving before the peak of the pulse, ionization of the material can begin before the main peak of the pulse arrives; see Fig. 2. The effect of the prepulse can be twofold: it can create a plasma at the front of the target that alters the electron heating (usually enhancing the efficiency), and it can send a shock through the target that breaks out to form a preplasma on the rear surface. Additionally, the interaction that is being probed may also cause preplasma at the rear of the target. In either case, this preplasma at the rear surface can inhibit proton acceleration (Kaluza et al., 2004; Fuchs et al., 2007; Higginson et al., 2021). For this reason, a shield to protect the proton source foil is often used to prevent these effects (Mackinnon et al., 2006; Zylstra et al., 2012).

The characteristics of the beams accelerated via TNSA are much different than those of conventional rf beams, with some superior properties that are particularly advantageous for use as a backlighter in proton-imaging applications. These result from the short duration of the acceleration process (Fuchs et al., 2006; Schreiber et al., 2006; Dromey et al., 2016) and from the fact that, unlike other ion sources, protons are cold when accelerated with minimal transverse energy spread. The beams are therefore highly laminar (Borghesi et al., 2004) and are characterized by ultralow transverse emittance [as low as 0.004 mm mrad; see Cowan et al. (2004)] and by ultrashort (~picosecond) duration at the source (Dromey et al., 2016). As a consequence, the emission properties of a TNSA beam can be described in terms of a virtual source that is much smaller than the region from which the protons are emitted and typically located at a small distance in front of the target (Borghesi et al., 2004). The proton-beam properties for imaging have been demonstrated to be optimum for ~picosecond duration laser pulses (Campbell et al., 2019) to limit emittance growth. If the driving laser pulse duration is longer than ~1 ps, magnetic-field instability growth on the rear surface deflects protons as they are accelerated (Nakatsutsumi et al., 2018). Another key characteristic of TNSA proton beams is that they are bright, with \(10^{11}\text{–}10^{13}\) protons per shot with energies \(>\) MeV, distributed across a broadband spectrum with a Boltzmann-like distribution. The proton-beam divergence is typically \(\lesssim 30^\circ\), with the divergence decreasing with increasing energy (Nürnberg et al., 2009).

The homogeneity of the transverse profile within a beam has been shown to be affected by the laser intensity profile at the target front (Fuchs et al., 2003), as well as by instabilities occurring within the target, particularly within insulators, which tend to degrade the uniformity of the profile (Fuchs et al., 2003; Ruyer et al., 2020). Metallic targets typically induce smoother beams than insulators (Quinn et al., 2011), and are therefore normally preferred for imaging applications.

2. D\(^3\)He

A different approach to generating protons is to use the fusion reaction products from an inertial implosion. These sources were first developed in the context of proton backlighters for ICF experiments at the Omega laser facility (Li et al., 2006a, 2006b) and have since been ported to the National Ignition Facility (NIF) (Zylstra et al., 2020). Contrary to the TNSA method, such a backlighter is formed by direct laser irradiation of a capsule filled with D\(^3\)He gas.

The D\(^3\)He backlighter platform uses a shock-driven implosion mode called an exploding pusher. As schematically illustrated in Fig. 4, the physical process involved in this scheme comprises three steps. First, multiple laser beams directly and symmetrically illuminate a thin-glass-shell capsule surface. Second, the strong laser absorption results in the explosion of capsule shell material, which drives a strong spherical shock wave propagating radially inward toward the capsule center. Finally, the converging shock collapses in the center and bounces back, resulting in an increase of ion temperature and fuel density, and (d) the facilitation of nuclear fusion reactions and burn.

**FIG. 4.** Schematic of an exploding-pusher mode of capsule implosion and fusion in direct-drive inertial-confinement fusion. (a) Multiple laser beams directly and symmetrically illuminate the thin-glass-shell capsule surface. (b) The explosion of the shell caused by laser energy absorption drives a strong spherical shock propagating radially toward the capsule center. (c) The converging shock collapses in the center and bounces back, resulting in an increase of ion temperature and fuel density, and (d) the facilitation of nuclear fusion reactions and burn.
relative proton numbers are shown in Fig. 5(a). More recently a triparticle backlighter platform utilizing a DT\(^3\)He capsule implosion was developed that provided 9.5 MeV deuterons from T + \(^3\)He \(\rightarrow\) α + d (9.5 MeV), in addition to the 3.0 MeV DD and 14.7 MeV D\(^3\)He protons (Sutcliffe et al., 2021). Note that the interaction of the drive lasers with plasmas ablated from the capsule surface can generate hot electrons that escape from the capsule surface, which can lead to electric charging of the imploding capsule that can “upshift” the proton energies. For a typical implosion driven by a laser intensity of \(10^{15}\) W cm\(^{-2}\), \(~\)megavolt electric potentials resulting in an \(~0.5–1.0\) MeV acceleration of fusion protons have been measured (Hicks et al., 2000; Rygg et al., 2008).

The typical implosion lasers consist of 0.6–1 ns square pulses without phase plate and cumulative energies of \(~10\) kJ. The capsules have diameters of approximately 420 \(\mu\)m, with a wall thickness of \(~2\) \(\mu\)m. The capsule bang time is approximately 450 ps, followed by a \(~100\) ps burn during which the protons are generated. During the implosion the capsules reach a minimum burn size of \(~40\) \(\mu\)m (FWHM), which sets the spatial resolution of the resulting proton beams.

Recent studies have started to explore how proton yield from D\(^3\)He sources varies with laser and capsule parameters; see Figs. 5(b) and 5(c). By statistically sampling several hundred backlighter shots, it was found that the total laser energy on the capsule and the asymmetry of the laser drive were the most important predictors of backlighter performance (Johnson et al., 2021). As a result, the best proton yields (both DD and D\(^3\)He) can be attained using as many drive beams as possible (at least 9 kJ is recommended) while keeping the capsule illumination as symmetric as possible; see Johnson et al. (2021) for details. In general, the combination of high asymmetry and a small number of beams should be avoided whenever possible.

D\(^3\)He protons have several unique features compared to TNSA protons. First, the fusion-generated protons are monoenergetic, with a typical energy uncertainty of about 3\% (Li et al., 2006a) due to the finite nuclear burn region and energy straggling on the backlighter. Second, the different characteristic energies of the DD and D\(^3\)He protons naturally result in distinct times of flight for each proton energy, which can provide a temporal resolution of \(~100\) ps. Third, a uniform and symmetric emission of fusion products provides a 4\(\pi\) solid angle isotropic proton fluence, though electric charging of the capsule may distort this (Manuel, Zylstra et al., 2012).

B. Detectors

Each proton source is associated with a corresponding detector, namely, radiochromic film (RCF) for TNSA protons and CR-39 for D\(^3\)He protons. In Secs. II.B.1–II.B.3 we discuss the properties and characteristics of these detectors, which play a key role, along with the beam properties, in determining the features of the proton images. Mention is also made of other detectors that have been used, albeit less frequently.

1. Film

RCFs are commonly used in dosimetry for a wide range of radiation sources (electrons, protons, and photons) for medical, industrial, and scientific applications. This is a high-dose, high-dynamic-range film that is widely used in a clinical context for x-ray dosimetry (Niroomand-Rad et al., 1998). RCF has become a popular choice for spectral and angular characterization of laser-driven proton beams (Nürnberg et al., 2009; Schollmeier et al., 2014), and the main detector of choice for TNSA-based proton imaging, thanks to its ease of use and effective performance at the particle fluences of typical experimental arrangements. The films consist of one or more active layers containing a microcrystalline monomeric dispersion buried in a clear plastic substrate. Different types are available, under the commercial GafChromic name, that have varying active layer thicknesses and compositions and consequently different sensitivities to ionizing radiation. Currently popular varieties are HD-V2 and EBT3.

There are a number of features that make RCF particularly attractive. RCF is a passive detector, the color and optical density of which is immediately, permanently, and visibly changed upon irradiation as a consequence of polymerization processes in the active layer, without the need for processing. The subsequent change in optical density can be calibrated against the radiation dose absorbed in the active layer of the film. Therefore, it is possible to extract information on particle fluence within the layer.

RCF can be digitized using inexpensive commercial photo-scanners (photo-type flatbed scanners), which are fast and offer high spatial resolution (1600 dpi, or 63 dots/\(\mu\)m, resulting in a resolution of 16 \(\mu\)m in most cases) and 16 bits.
per channel. The intrinsic spatial resolution of RCF is higher (typically of micron scale) than the resolution of the scanners. RGB scanning provides separate color channels and produces images with different contrasts or sensitivities and provides options for further extending the dynamic range of the film. Conversion of the scanned images into doses requires a prior calibration of the film, which is typically obtained by exposing the films to known doses delivered by well-characterized fluxes of protons in conventional accelerators (Chen et al., 2016; Bin et al., 2019; Xu et al., 2019).

In standard experimental configurations, RCFs are used in a stack arrangement such that each layer acts as a filter for the following ones in the stack. Sometimes additional filter layers, typically aluminum foils, are used as spacers. The signal in a given film within the stack will be due only to protons having energy \( E \geq E_B \), where \( E_B \) is the energy reaching the Bragg peak within the active layer of the film. In the first approximation, for a Boltzmann-like spectrum such as those typically produced by TNSA, the dose deposited in a layer can be taken as deposited mostly by protons with \( E \sim E_B \). As we see in Sec. IIC.4, this property is at the basis of the unique temporal characterization capabilities of TNSA proton imaging. An example of an RCF stack is shown in Fig. 6 and illustrates the color change of the film and the reduction in the beam divergence at higher proton energies.

2. CR-39

The \( ^3 \)He backlighter is ideally complemented by imaging detectors made of CR-39 (Séguin et al., 2003). Although the process of reading out the data recorded on CR-39 is complicated (as we later discuss), the great advantage is that it records the exact position of every individual incident charged particle in the detector plane to an accuracy of \( \sim 2 \, \mu m \), as long as the maximum incident particle fluence is smaller than about \( 10^6 \) per cm\(^2\). These fluence limit and saturation effects at higher flux (Gaillard et al., 2007) are why CR-39 is not typically used for TNSA proton-beam detection.

CR-39 polymer is part of a class of solid-state nuclear-track detectors that have been used for decades in many high-energy particle counting applications, from radioactive dating to cosmic rays and neutrons; see Fleisher, Price, and Walker (1965) and references therein. It has the useful property of being relatively insensitive to other forms of ionizing radiation, like gamma rays, x rays, or electrons, and is nearly 100% efficient at detecting ions in a given energy range. Consequently, CR-39 has become the workhorse for \( ^4 \)He capsule backlighter experiments. It has also been used to calibrate other detectors due to its high efficiency and known response (Harres et al., 2008; Mančić et al., 2008). CR-39 is typically arranged in a two-layer stack with associated filters such that one layer is sensitive to 3.0 MeV DD protons and one layer is sensitive to 15 MeV \( ^4 \)He protons.

CR-39 is a transparent plastic with chemical composition \( \text{C}_{12}\text{H}_{24}\text{O}_{7} \) (Fews and Henshaw, 1982; Séguin et al., 2003, 2016). A charged particle of appropriate energy passing through it leaves a trail of damage along its path in the form of broken molecular chains and free radicals. The amount of local damage along the path is related to the local rate at which energy is lost by the particle (\( dE/dx \), where \( x \) is the distance along the path). The length of the path is the range of the particle in the plastic. Particle paths can be made visible by etching the CR-39 in NaOH (Fews and Henshaw, 1982; Gaillard et al., 2007); the etch time is typically between 0.5 and 5 h (based on characteristics of the experiment such as the expected backlighter yield). The surface of the plastic is etched away at a “bulk etch rate,” while damaged material along a particle path etches at a faster “track etch rate.” If a particle path is normal to the plastic surface, the result of etching is a conical pit, or “track,” with a sharply defined, round entrance hole.

Retrieving information about all individual particle tracks in an exposed piece of CR-39 involves scanning the entire CR-39 surface with an automated microscope system. The bottom panel of Fig. 7(c) shows a sample microscope image of \( ^3 \)He-proton tracks, each of which appears as a dark circle on a light background. The location of each pit shows where a proton entered, and its diameter provides a measure of \( dE/dx \) for the proton. Since \( dE/dx \) is different for particles of a given type but different energies, the diameter can provide a measure of particle energy (after passing through any filters in the detector pack). \( dE/dx \) is also different for different particle types, so diameters can often be used to identify the particle type if the energy is known (see Fig. 8), or to estimate the energy if the particle type is known (Sinenian et al., 2011; Zylstra et al., 2011; Lahmann et al., 2020). One can use not only different particles or source energies to form images at different times [due to the time of flight; see Li et al. (2009)] but also the known down-scattered energies of one of the monoenergetic particles to produce separate images of the same target at the same time; see Fig. 7.

The optical magnification used in the scanning microscope system is usually (but not necessarily) chosen so that one camera frame covers the area that will be used for one pixel in the final desired proton image of particle fluence versus position. That area is often chosen to be about \( 300 \times 300 \, \mu m^2 \). Each such camera image is evaluated with special algorithms that identify every individual track and determine its position, coordinates, its diameter, its optical contrast, and its

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**FIG. 6.** RCF stack obtained at PHELIX, consisting of seven films of the type HD-810 and five films of the type MD-55. The proton Bragg peak energy is given for each film layer. Adapted from Bolton et al., 2014.
eccentricity (Séguin et al., 2003). All of these measured parameters are recorded, and the microscope moves on to the next frame, continuing until the entire surface is covered. The resultant “scan data” file is saved for later processing, in which the final proton image is made by going through all of the recorded track information after deciding which display resolution is desired (frequently one microscope frame for each pixel) and counting the number of tracks in each “pixel” area that satisfies carefully chosen limits on diameter, contrast, and eccentricity (Séguin et al., 2003). Examples are given in Sec. IVI.

3. Other detectors

While passive, single-use detectors such as RCF and CR-39 have been used in the vast majority of proton-imaging experiments thus far, the use of microchannel plates (MCPs) has been also reported in the literature. MCPs, which are high-gain, spatially resolved electron multipliers (Bolton et al., 2014), have often been used in proton acceleration experiments, mostly in the dispersion plane of a magnetic spectrometer or Thomson parabola (Harres et al., 2008). An arrangement reported by Sokollik et al. (2008) extended this use to a streaked deflection approach in which a TNSA beam is analyzed, after backlighting a target, in a magnetic spectrometer coupled to an MCP. Use of MCPs as a proton-imaging detector in a standard projection arrangement was also reported by Sokollik et al. (2008). In this case the selection of a temporal snapshot is done by temporal gating of the MCP on nanosecond timescales, which is reflected in a temporal resolution of ~60 ps at the interaction plane and significant integration of the investigated ultrafast phenomena (the explosion of a laser-irradiated water droplet).

Initial tests with scintillator plates (Tang et al., 2019) have indicated that by selecting appropriate detector parameters these can be used as an alternative to RCF, with the advantage of being suited to repeated use. The main disadvantages of scintillator detectors for proton imaging are (1) the energy resolution is reduced compared to RCF due to the thickness of the detector material, and (2) it is difficult to extract the signal from different detector layers. A novel setup devised by Huault et al. (2019) using a concertina design of scintillators has been used to observe the proton energy-spectra and proton-beam divergence simultaneously. See Sec. V.B for further discussion.

C. Diagnostic geometry and other considerations

A diagram of a typical proton-imaging setup as deployed in an experiment is shown in Fig. 9. The source can be either TNSA- or D³He-generated protons, with corresponding detectors of either RCF or CR-39, respectively. During an experiment, the protons are emitted by the source, propagate a distance $r_s$ to the interaction region where they acquire small deflections due to the electromagnetic fields, and then travel ballistically a distance $r_d$ to the detector. In a number of experiments (Paudel et al., 2012; Ahmed et al., 2016; Obst-Huebl et al., 2018; Ferguson et al., 2023), self-probing arrangements have also been demonstrated, where the TNSA protons accelerated from a foil are used to probe phenomena initiated by the same laser pulse that has accelerated them, for instance, in parts of the same target from which they are emitted, or in the surrounding medium.
Typical implementations of TNSA and D$^3$He sources and examples of detector stacks are also shown in Fig. 9.

A standardized TNSA source target has been developed at OMEGA EP (Zylstra et al., 2012), in which a thin foil is mounted within a plastic tube, with a thin protective foil mounted over the end. This shields the TNSA foil from radiation and plasma emerging from the object under study. The tubes are transparent, which allows alignment of the laser focus to the foil via target chamber cameras. The TNSA foil is driven by a short-pulse laser, which can be moderately off axis to allow some setup flexibility. The resulting protons are emitted in a cone normal to the TNSA foil with a broadband energy distribution.

Likewise, a standard D$^3$He source capsule has been developed for both OMEGA (Li et al., 2006a, 2006b) and NIF (Zylstra et al., 2020). The capsules are mounted on stalks and driven by a relatively symmetric set (typically > 20) of long-pulse beams, resulting in protons emitted into $4\pi$ with monoenergetic energy distributions. A comparison of TNSA and D$^3$He proton sources and detectors is summarized in Table I.

1. Magnification

Typical setups take advantage of the small source size of the protons and obtain a magnified image onto a larger detector, with magnification

$$M = \frac{r_d + r_s + l_{\text{path}}}{r_s} \approx \frac{r_d + r_s}{r_s},$$

Such a setup is often used to magnify the image from the plasma size (millimeters to 1 cm) to the detector size (typically several centimeters). The magnification also improves the spatial resolution at the plasma plane by a factor $M$ compared to the detector’s spatial resolution. Note that $r_d$ is in principle different for each layer in the detector stack. This can be especially important for the analysis of TNSA detector stacks, which can have a large number of layers. Additionally, in experiments where the interaction length $l_{\text{path}}$ is large, there can also be a significant variation in $M$. An example of this is discussed in Sec. IV.G.

Experimental design should consider the size of the interaction such that the proton beam has expanded to overfill the region of interest. A small-angle approximation is often used to assume that the proton beam along the probing axis travels the same distance as the protons at the edge of the detector (or the beam if it is smaller). TNSA proton beams typically have a divergence of less than 30°, meaning that the small-angle approximation is reasonable in most cases, whereas the D$^3$He implosion is an isotropic source, so a limited solid angle should be used. Similarly, when calculating the energy of the protons for a particular RCF stack layer, the extra distance within material traveled by the protons at the edge of the beam is usually ignored.

2. Meshes and grids

An optional mesh can be used to break the initial proton beam into beamlets, which in a proton deflectometry approach (Mackinnon et al., 2004) facilitates measurement of the fields via a direct tracking of beamlet deflections. The meshes used...
are typically commercial transmission electron microscopy grids that are available in a variety of pitches, hole widths, and bar widths and are manufactured from relatively high-Z metals such as copper, nickel, and gold. The thickness of the meshes is typically such that a shadow is imprinted on the proton beam via multiple scattering in the mesh bars (Borghesi et al., 2004). By geometric arguments the mesh magnification of the detector is (see Fig. 9)

\[ M_{\text{mesh,d}} = \frac{r_d + r_s + l_{\text{path}}}{r_g} \quad (6) \]

The spatial resolution is in turn set by the projection of the mesh period \( p \) onto the plasma plane, i.e., \( (r_s/r_g) \propto p \).

The period of the mesh should ideally be chosen such that a sufficient number of mesh elements is projected across the probed region of interest. The period of the mesh should also be larger than the source size so that the mesh is not overly smeared out when projected.

A variation on the beamlet technique is to use an object (such as a mesh, mask, or pepperpot) to subaperture the proton beam into many beamlets (Sokollik et al., 2008; Johnson et al., 2022), down to a few “pencil” beamlets (Lu et al., 2020), or even just down to a single beamlet (Chen et al., 2020). This allows one to probe areas of specific interest in a limited fashion that is more easily detectable (in terms of deflection) or to streak the beamlet in time.

3. Spatial resolution

As is typical of all projection backlighting schemes, the intrinsic and ultimate spatial resolution of proton images is determined by the size of the proton source. For D\(^{3}\)He capsules this is set by the burn volume of the implosion, which has been measured to be typically 40 \( \mu \text{m} \) FWHM (Manuel, Zylstra et al., 2012); see Sec. II.A.2. For TNSA targets the relevant size is instead the “virtual” source size resulting from the beam’s laminarity and emittance (Borghesi et al., 2004); see Sec. II.A.1. This is typically of the order of 10 \( \mu \text{m} \) FWHM (Borghesi et al., 2004; Wang et al., 2015; Li et al., 2021), set by the size of the laser focal spot, but can vary from experiment to experiment.

Scattering of the protons in the plasma being probed can (and often does) degrade the spatial resolution from the previously given values, particularly for dense plasmas. The magnitude of the scattering will depend on its density and dimensions, as well as on the proton energy, and typically leads to a Gaussian distribution of angles with some spatial width \( \theta_{\text{SC}} \), which can be evaluated using Monte Carlo calculations (Ziegler, Ziegler, and Biersack, 2010) or through empirical formulas (Highland, 1975; Kanematsu, 2008). This causes a resolution degradation characterized by a 1/e spatial width of the order of \( r_d \theta_{\text{SC}} / M \sim r_s \theta_{\text{SC}} \) in units of distance in the plasma plane (Li et al., 2006a). For low-Z plasmas with electron number densities that are \( \lesssim 10^{20} \text{ cm}^{-3} \), this degradation is typically small compared to the effect of the finite source size (see Bott et al., 2017, Appendix B); however, for experiments with higher-density plasmas (\( \gtrsim 10^{22} \text{ cm}^{-3} \)), scattering significantly reduces the resolution. In such experiments, scattering is an important effect to take into consideration for an accurate determination of the fields associated with the proton image; see Sec. III.B.2. A similar effect will be caused by scattering in any protective foil (TNSA sources), although the foil thickness is typically chosen in order to minimize the angular spread of the beam.

By contrast, the characteristics of the detector do not usually have a significant effect on the spatial resolution of proton images. Scattering in the detector, which can occur when protons cross a stack on the way to the layer where they are detected, normally leads to a negligible resolution loss once the magnification is taken into account. Similarly, the intrinsic spatial resolution of the detector is typically high, of the order of microns, and therefore does not contribute to the spatial resolution of the diagnostic when registered back to the plasma plane.

Another potential source of degradation of the spatial resolution arises in the presence of a background magnetic field (as used for magnetized plasma experiments), as the energy-dependent deflection of protons within the energy response curve of a layer may lead to blurring of the proton image along the deflection direction. This effect was discussed by Arran, Ridgers, and Woolsey (2021).

4. Temporal resolution and multiframe capability

There are three primary factors contributing to the temporal resolution of proton images (Sarti et al., 2010a):

1. The temporal duration of the source \( \delta t_\epsilon \). As discussed, this is of the order of \( \sim 1 \text{ ps} \) for TNSA beams for picosecond drivers [shorter for femtosecond drivers (Fuchs et al., 2006)] and \( \sim 100 \text{ ps} \) for the D\(^{3}\)He capsules. This is the factor that determines the ultimate temporal resolution possible for a proton image and the dominant factor for probing with D\(^{3}\)He protons.

2. The transit time \( \delta t_\ell \) of the protons through the region where the transient fields are located. This is related to the spatial scale over which the fields under investigation extend and is therefore intrinsic to the phenomenon under investigation. If the fields change on the timescale of the proton transit, the information will be temporally averaged over a time

\[ \delta t_\ell \sim \frac{l_{\text{path}}}{v_p} \sim l_{\text{path}} \left( \frac{m_p}{2e_p} \right)^{0.5} \quad (7) \]

where \( e_p \) and \( v_p \) are the energy and velocity of the protons, respectively. For example, for 10 MeV protons crossing a 100 \( \mu \text{m} \) region, one has \( \delta t_\ell \sim 2 \text{ ps} \).

3. The time-of-flight uncertainty (from the source to the plasma being probed) \( \delta t_d \) resulting from the energy resolution \( \delta e_p \) of the detector is given by

\[ \delta t_d \sim r_s \left( \frac{m_p}{2e_p} \right)^{0.5} \delta e_p, \quad (8) \]

which can also be of the order of \( \sim 1 \) ps. More detailed considerations associated with a multilayer RCF stack are given later.

While \( \delta t_\ell \) and \( \delta t_d \) are typically not relevant to determine the resolution for D\(^{3}\)He-proton images (where the source duration
is the dominant factor), they all can contribute significantly to the temporal resolution for experiments employing TNSA protons. Under standard experimental conditions and depending on the specific experimental arrangement, this is typically in the range of 1–5 ps.

Both TNSA and D$^3$He sources emit protons in a burst, which is typically shorter (or much shorter in the case of TNSA) than the time of flight to the plasma $r_s/v_p$. For $r_s = 1$ cm, for example, this would be $\sim 180$ ps for 15 MeV protons and $\sim 400$ ps for 3 MeV protons. Consequently, a multiframe capability can be achieved using energy-resolving detectors (as RCF or CR-39 stacks), where stacking up images from different proton energies provides information on the temporal dynamics of the system over time intervals of the order of hundreds of picoseconds. Obtaining multiple snapshots enables one to follow the temporal dynamics of the same event, which is particularly useful under conditions where there is a pronounced shot-to-shot variability.

For D$^3$He sources, different frames can be obtained by employing the different fusion products produced during the implosion; see Fig. 9. An example of the application of this capability is provided in Fig. 10. The structure of the detector pack involves two metal filters and two separate layers of CR-39. The first CR-39 layer is proceeded by one of the metal filters, which helps protect the CR-39 from debris while still allowing the detection of $\sim 3$ MeV DD protons. A second filter is placed before the second CR-39 layer and acts to help slow down the $\sim 15$ MeV D$^3$He protons to energies of 1–6 MeV, which is the best energy range for detecting protons on the CR-39.

The broadband spectrum of TNSA sources allows sequential temporal frames to be recorded in consecutive layers of an RCF stack. When using high-energy TNSA protons from a petawatt-class laser system, one can obtain up to several tens of temporally separated proton images of the interaction. In the multiframe approach, every layer is labeled temporally with the time of flight (calculated from the source to the center of the film pack) of the energy at which the relevant response curve is maximized (essentially the energy reaching the Bragg peak in the active layer of the RCF). Figure 11(a) shows the energies reaching the Bragg peak at a certain depth in the RCF pack and the corresponding time of flight for different source-plasma separations. The active layers of different RCFs (for instance, separated by $\sim 100$ μm distances

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**FIG. 10.** Proton images of a laser-driven, solid 840 μm diameter CH sphere, made using a setup similar to Fig. 9. (a) Image recorded with no laser drive on the CH sphere. (b)–(d) Images recorded with laser drive for three different particle types and energies. Adapted from Séguin et al., 2012.

**FIG. 11.** (a) Bragg peak proton energy vs depth in the RCF stack (solid curve) and corresponding time of flight for three different values of the source-object separation $r_s$ (dotted line, 1 mm; dashed line, 2 mm; dash-dotted line, 3 mm). (b) Normalized energy response curves for an RCF stack made of several layers of HD810 (solid curve). (c) Normalized temporal response curve for the stack configuration in (a) and a source-plasma distance of 3 mm (solid curve). The red dashed curves in (b) and (c) are response curves multiplied by a typical TNSA exponential spectrum with temperature of 2 MeV. Adapted from Romagnani, 2005.
in a stack consisting of HD films) will therefore contain snapshots taken at discrete time values along the red curves in Fig. 11(a). An example of energy and temporal response for four consecutive layers (the second to the fifth) in an RCF stack is shown in Fig. 11(b) (Romagnani, 2005), based on SRIM (stopping and range of ions in matter) (Ziegler, Ziegler, and Biersack, 2010) calculations. The energy response of a layer is dependent on the spectral profile of the proton beam, and Fig. 11 highlights the difference between the response to a flat spectrum and a more realistic Boltzmann-type spectrum with a finite temperature (which in an experimental setting can be obtained from a dedicated spectral characterization of the proton beam). Taking into account the energy-dependent time of flight from source to object, the energy response of Fig. 11(b) translates into the temporal response for each layer shown in Fig. 11(c). The figure highlights the multiframe capability of the diagnostic arrangement, where each layer primarily contains information about a particular time in the evolution of the transient plasma being probed. As is visible in Fig. 11(a), as the time-of-flight curve becomes shallower for increasing depth, the temporal separation between the snapshots obtained in consecutive layers decreases for deeper layers and higher proton energies. The detector-limited temporal resolution of the snapshots also increases for deeper layers, in correspondence to a more selective energy response. By focusing, for example, on layer HD5 in Fig. 11(a) (red dashed curve), one obtains $\delta \epsilon_p (\text{FWHM}) \sim 0.2 \text{ MeV}$, which, for $r_5 = 3 \text{ mm}$ and employing Eq. (8), corresponds to $\delta t_\delta (\text{FWHM}) \sim 2.5 \text{ ps}$. For deeper layers in the pack $\delta t_\delta$ becomes of the order of 1 ps or less, depending on the energy of the protons.

A suitable choice of parameters allows interframe time steps of the order of picoseconds or less to be obtained, as achieved in the data of Fig. 12 (Romagnani et al., 2005). In this experiment $r_5$ was reduced to 1 mm, which, coupled to a proton spectrum with a cutoff at $\sim 12 \text{ MeV}$, leads to $\sim 1 \text{ ps}$ temporal frame spacing at the higher end of the spectrum. For example, the eighth and ninth layers in an HD pack would select, respectively, energies of $e_{p_8} \sim 9$ and $e_{p_9} \sim 10 \text{ MeV}$, leading to an interframe temporal separation $\delta t_\delta \sim r_5 (m_p/2)^{0.5} (1/e_{p_8})^{0.5} - (1/e_{p_9})^{0.5} \sim 1 \text{ ps}$. Detecting highly transient features [such as the sheath field in Fig. 12(c), which was seen to exist for about 1 ps] therefore becomes possible if one carefully times the proton probe relative to the interaction such that protons of sufficiently high energy transit through the region of interest at the appropriate time; this is done by appropriately adjusting the relative timing of the laser pulse accelerating the probe protons and the interaction pulse [labeled as CPA1 and CPA2 in Fig. 12(a)]. Under these conditions the dominant factor in determining the temporal resolution is often the proton transit time $\delta t_\delta$.

When probing ultrafast phenomena, it is often necessary to consider time-of-flight variations across a single RCF layer. These arise from the longer path of protons, propagating obliquely and intercepting the RCF layer at an angle, compared to the protons propagating on axis, which can lead to temporal differences of the order of picoseconds across the RCF layer. This is important, for example, when imaging field structures moving at speeds close to $c$ across the field of view of the proton images (Kar et al., 2007; Ahmed et al., 2016). A modified projection arrangement specifically designed for the detection of ultrafast moving fronts was described by Quinn et al. (2009b).

In cases where the field configuration probed is complex and changes on timescales of the order of the interframe separation or faster, additional complications may arise in the interpretation of the RCF data due to the fact that the dose deposited in a specific layer by protons stopping deeper in the stack will carry information on the field distribution at earlier times than the time determined by the Bragg peak energy [Fig. 11(a)]. The identification of these ghosting artifacts (Quinn, 2010) (which will typically be fainter than the main features in a layer) is an important part of the analysis, which is facilitated by the observation of the dynamics over several RCF layers and extended the temporal range. Several deconvolution techniques exist (Breschi et al., 2004; Nüremberg et al., 2009; Kirby et al., 2011) for removing the contribution of higher-energy protons from preceding layers in the context of the spectral characterization of TNSA proton beams, which in principle could be applied for removing temporally spurious contributions in proton-imaging data and for increasing the temporal purity and resolution of single RCF layers. However, this becomes complex in the case of a dynamically changing dose distribution; this approach has not been reported thus far to our knowledge. Instead, forward modeling employing particle tracers and dynamically evolving field distributions (see Sec. III.B) can be used to produce synthetic radiographs for comparison with the experimental data and the identification of overlapping temporal features (Ramakrishna et al., 2008; Kar et al., 2016).

![Fig. 12. Proton probing of the expanding sheath at the rear surface of a laser-irradiated target. (a) Setup of the experiment. A proton beam is used as a transverse probe of the sheath. (b)–(g) Temporal series of proton images in a time-of-flight arrangement. The probing times are relative to the peak of the interaction. (h) Deflectometry image where a mesh is placed between the probe and the sheath plasma for a quantitative measure of proton deflections. Adapted from Romagnani et al., 2005.](image-url)
III. THEORY OF PROTON-IMAGING ANALYSIS

A. Basics

As explained in the Introduction, the physics that underpins the proton-imaging diagnostic is simple. With the exception of interactions with dense HED plasma or matter (see Sec. II.C.3), the characteristic speeds of imaging-beam protons are sufficiently large that collisional interactions between the beam protons and the plasma being probed are usually negligible (Kugland, Ryutov et al., 2012; Bott et al., 2017). In addition, the characteristic density of proton-imaging beams is sufficiently low that the beam does not perturb the plasma via either collisionless plasma interactions or space-charge effects (Kugland, Ryutov et al., 2012). As a result, the protons that constitute typical imaging beams behave like test particles, as they are deflectected by electromagnetic forces associated with fields already present in the plasma prior to the arrival of the proton beam. Thus, the proton beam’s profile post interaction encodes information about the inherent electric and magnetic fields of the plasma.

The trajectory of charged particles through electric and magnetic fields (and the final velocity of those particles post interaction) can be complicated in the general case of arbitrary proton speeds and characteristic field strengths; proton-imaging setups typically overcome this issue with their use of fast multi-MeV protons (see Sec.II.A) and careful geometric design to restrict the set of possible proton trajectories. For most laser-plasma experiments currently performed, the magnitude of deflection angles due to plasma-generated magnetic fields (and the final velocity of those particles post interaction) can be complicated in the general case of arbitrary magnetic fields (Romagnani et al., 2005; Li et al., 2007; Petrasso et al., 2009; Willingale et al., 2011b; Tubman et al., 2021). The main advantage of this approach is its conceptual simplicity. However, it does have a few issues. Determining the exact projection of the initial profile in the absence of any deflections is not always a trivial matter, because confounding factors such as imperfect target fabrication can mean that a deflectometry grid’s position is not always consistent from shot to shot. Blurring of the mesh due to the ablation of actual physical grids by strong x-ray radiation that inevitably arises during the course of laser-plasma experiments can also inhibit successful tracking of the grid’s distortion (Johnson et al., 2022; Malko et al., 2022). In some circumstances, the grid itself can become charged, resulting in apparently distorted grids when there is in fact no interaction of the proton beam with plasma electromagnetic fields (Palmer et al., 2019). The resolution of electromagnetic-field measurements is also limited to that of the grid; this constraint is inevitably much larger than the theoretical resolution that can be achieved given typical proton source sizes; see Sec. II.A. Finally, in cases of highly nonuniform deflections, successfully tracking the grid’s distortion is not always possible (Willingale et al., 2010b).

A second approach that attempts to overcome these issues is to assume approximate transverse uniformity of such beams prior to their interaction with a plasma being imaged (a property of proton-beam sources that typical experimental setups aim to realize; see Sec. II.C), and thereby quantitatively relate inhomogeneities in the beam profile detected post interaction on a proton image to electromagnetic fields in addition to a velocity displacement $\Delta v_\perp$, which in principle complicates the interpretation of a nonuniform proton-beam profile. However, by ensuring that the distance $r_d$ from the plasma to the detector is much larger than $l_{\text{path}}$ (a geometric setup of this form is known as point-projection geometry), it follows that the measured displacement $\Delta d_\parallel$ of protons from their projected position $d_{\perp 0}$ in the absence of any electromagnetic fields is dominated by the displacement acquired as protons free stream at their slightly perturbed velocity $\Delta d_\perp = r_d \Delta v_\perp / v_0$, with $|\Delta d_\perp| \approx r_d \Delta \alpha \gg |\Delta x_\perp|$.

Historically this conclusion has been leveraged to discern properties of the electromagnetic fields in the plasma using a proton beam in two ways. The simpler technique is to introduce a well-defined spatial modulation to the profile of the proton beam prior to its interaction with any electromagnetic fields using a grid (see Sec. II.C.2); only protons that do not intersect the grid are subsequently detected. Any distortions $\Delta d_\parallel$ to the grid-induced profile detected post interaction (which provide a direct measure of $\Delta d_\perp$) can then be attributed to angular deflections caused by electromagnetic fields in the plasma, and the line-integrated values of two components of those fields estimated via

$$\int_{0}^{l_{\text{path}}} ds \left\{ E_\perp(x) + \frac{\tilde{v} \times B_\perp(x)}{c} \right\} \approx \frac{m_p v_0^3}{e r_d^2} \Delta d_\parallel. \quad (10)$$

This technique is typically known as proton deflectometry and has been successfully used in a number of different laser-plasma experiments to provide measurements of electromagnetic fields (Romagnani et al., 2005; Li et al., 2007; Petrasso et al., 2009; Willingale et al., 2011b; Tubman et al., 2021). Several subtle caveats exist to this statement; we discuss them subsequently.
the plasma (Kugland, Ryutov et al., 2012; Bott et al., 2017; Graziani et al., 2017; Kasim et al., 2017). The successful interpretation of detected nonuniformities in proton images in terms of the electromagnetic fields associated with them using either of these approaches requires a theoretically grounded analysis methodology. Historically there have been two methodologies that have been used for this interpretation: particle-tracing simulations and analytical modeling. We discuss the two approaches in Secs. III.B and III.C, respectively.

B. Particle-tracing simulations

1. Overview

Analyzing proton images using particle-tracing simulations is typically done as follows. A candidate model for an electromagnetic-field structure in a particular laser-plasma experiment is proposed; the interaction of the proton beam (whose parameters are chosen to be the same as those used experimentally) with that field structure is simulated using a test-particle-tracing code; a simulated proton image associated with that proton beam is then generated; and finally the simulated image is compared with the experimental one, with the candidate model deemed to be reasonable if there is qualitative (or, ideally, quantitative) agreement. Particle-tracing simulations provide a powerful approach for analyzing proton-imaging data because they make relatively few assumptions about the nature of the interaction between the proton beam and the plasma being imaged.

Arguably the most important question to consider when using particle-tracing simulations to analyze proton images involves how to construct an appropriate candidate model for the electromagnetic field. There are two approaches to addressing this question that have been used for analyzing data from previous laser-plasma experiments. The first is to use the electromagnetic fields generated by a high-energy-density-physics (HEDP) code of the relevant laser-plasma experiment. The second involves introducing a physically motivated parametrized model and optimizing the model’s parameters using an algorithmic best-fit procedure. These approaches are often used complementarily, with the output of a HEDP code serving as the inspiration for a simpler, parametrized model. The two approaches are discussed in Secs. III.B.3 and III.B.4, respectively. Irrespective of the approach used to construct the candidate electromagnetic-field model, the successful use of particle-tracing simulations relies upon efficient particle-tracing algorithms; we therefore discuss these algorithms first.

2. Particle-tracing algorithms

The process underpinning a typical particle-tracing algorithm is illustrated in Fig. 13. Particle-tracing simulations typically employ a Monte Carlo method. To begin, synthetic protons are generated at the location of the proton source and assigned a velocity that, aside from being constrained to have a prespecified magnitude and an orientation with a cone of a certain solid angle, is random. These particles are then traced to the compact domain in which the possibly time-dependent electromagnetic fields are defined. In this domain, the nonrelativistic equation of motion for protons under the action of the Lorentz force associated with the electromagnetic fields is numerically integrated along particle trajectories. This integration is implemented using efficient numerical schemes in typical particle-tracing simulations so that the simulations can be run quickly for millions of synthetic protons (Birdsall and Langdon, 1985; Welch et al., 2004; Vay, 2008).

Once a given synthetic proton has completed its interaction with the electromagnetic field, the output can then be included in various particle diagnostics: most immediately synthetic proton images, but also other outputs such as deflection maps. The synthetic images can then be compared with experimental ones; for simple electromagnetic-field distributions (see Sec. III.B.4), quantitative comparison metrics between the synthetic and experimental images can then be used to refine the field distribution. There are several bespoke particle-tracing simulation codes optimized for proton-imaging analysis, including PTRACE (Schiavi, 2008), qTrace (Romagnani et al., 2005), the proton-imaging unit of the HEDP code FLASH (Fryxell et al., 2000; Tzeferacos et al., 2015), and PlasmaPy (PlasmaPy Community, 2023).

While the most basic particle-tracing codes typically make several physically motivated assumptions about the proton beam’s properties and the physics of its interaction with the plasma through which it is passing, one of the strengths of the particle-tracing approach is that it is often feasible to relax these assumptions. For example, while most proton-imaging particle-tracing codes assume a point source of monoenergetic

![Workflow for a typical particle-tracing algorithm. Courtesy of L. Romagnani.](image-url)

protons created instantaneously with a smooth spatial profile, it is a simple matter to include a finite source size or emission time, use a predefined spectrum of proton energies, or incorporate realistic random departures from laminarity. It is also not too challenging to include some additional physics beyond the simple action of Lorentz forces. For example, in dense plasmas scattering or energy loss of beam protons due to Coulomb collisions must be modeled in order to obtain realistic synthetic proton images; see Sec. II.C.3. When this is done, successful measurements of electromagnetic fields in dense plasma can be made: Romagnani et al. (2019) used qTrace particle-tracing simulations that included a scattering model to successfully diagnose the time evolution of fast-electron-induced current filaments in dielectric foams, while Lu et al. (2020) showed that, provided that scattering was included in supporting particle-tracing simulations, magnetic fields generated by the Biermann battery at a shocked shear layer in a dense foam could be observed. Particle-tracing simulations of proton beams that have been performed using full particle-in-cell (PIC) codes (Huntington et al., 2015) are capable of including another physics effect (albeit one that is usually not important): the beam’s feedback on the electromagnetic fields being imaged via collisionless interaction mechanisms.

3. Combined modeling with HEDP codes

Because of the complexity of the physics inherent in most laser-plasma experiments, as well as the difficulties involved in diagnosing such experiments, HEDP simulation codes are typically used to help design, implement, and interpret their results. Depending on the experiment, the state-of-the-art codes that are run at the present time are magnetized fluid codes (for instance, FLASH, LASNEX (Zimmerman et al., 1977), GORGON (Chittenden et al., 2004), RAGE (Gittings et al., 2008), and HYDRA (Langer, Karlin, and Marinak, 2015)), particle-in-cell codes (for instance, OSIRIS (Fonseca et al., 2002), EPOCH (Arber et al., 2015), PSC (Germsachewski et al., 2016), SMILEI (Derouillat et al., 2018), and VPIC (Bird et al., 2022)), or hybrid codes [for instance, ZEPHIROS (Kar et al., 2009; Ramakrishna et al., 2010), CHICAGO (Thoma et al., 2017), and dHybrid (Gargaté et al., 2007)], all of which output electromagnetic fields. Thus, choosing to use the outputs from such codes as inputs for electromagnetic-field candidates in particle-tracing simulations of proton images is a natural approach. For the outputs of such particle-tracing simulations to provide a plausible comparison with experimental data, the HEDP simulation should be either three dimensional or two dimensional with symmetry, with good spatial (temporal) resolution over sufficiently large spatial (temporal) scales. Aside from the ease of implementation if HEDP simulations have already been completed, this approach can be particularly advantageous if complex electromagnetic-field geometries arise (see Fig. 20 in Sec. IV.C for an example); constructing parametrized electromagnetic-field models from scratch in such situations is laborious. That being said, relying solely on synthetic images derived from HEDP simulations can become problematic if those images turn out to be qualitatively and/or quantitatively distinct from the experimental data they are meant to model. If this situation arises, it is often challenging to determine how to “correct” the outputs from HEDP simulations systematically.

4. Parametrized field models

Provided that the morphology of experimentally observed proton-fluence inhomogeneities is not too complex, it is often the case that a simple parametrized analytical model for a candidate electromagnetic-field, motivated by considerations of the physical mechanism(s) responsible for generating that field, can be constructed. The optimal choice of the parameters can then be found iteratively using particle-tracing simulations: given a first guess of parameters, a synthetic image is generated and then compared with the experimental image, with the quantitative differences between the outputs then used to determine a revised set of parameters, etc. (Romagnani et al., 2005, 2008a; Cecchetti et al., 2009). We note that in practice previous instances of particle-tracing simulations that have involved updating a parametrized electromagnetic-field model via a direct comparison between synthetic images and actual data do not explicitly report the rate of convergence to the best-fit parameters. This approach can prove helpful if 3D HEDP simulations of a given experiment have not been performed or are producing outputs that are discrepant with experimental data. By construction the technique will recover a good fit to the experimental data for simple proton-fluence inhomogeneities; however, for inhomogeneities lacking symmetry, successfully devising an appropriate analytical model with only a few parameters becomes difficult. Examples of this approach being applied to proton-imaging data are presented in Figs. 16 and 22 in Sec. IV.

C. Analytical modeling

1. Overview

The second methodology for interpreting proton images that has been utilized historically is analytical modeling: that is, relating the line-integrated values of electromagnetic fields to inhomogeneous distributions of the detected proton fluence analytically under a set of simplifying assumptions (Romagnani et al., 2005); Kugland, Ryutov et al., 2012). While analytical solutions of this type can be used to test particular candidate electromagnetic-field models (analytical forward modeling), they have proven to be particularly helpful in two key regards. First, they provide a direct interpretation of proton-fluence inhomogeneities in terms of either physical properties of the plasma (specifically, path-integrated charge and current structures) or features inherent in point-projection imaging (specifically, caustics); for discussions of both cases, see Sec. III.C.3. Second, analytical theory has been used to show the conditions under which the determination of line-integrated electromagnetic-field structures from proton-fluence inhomogeneities (which we refer to as field reconstruction) is a mathematically well-posed inversion problem, and if those conditions are met, how such field-reconstruction can be carried out systematically.

An analytical theory of proton imaging is not really tractable unless simplifying assumptions about the imaging setup are made; these assumptions are outlined in Sec. III.C.2, as is the theory that follows directly from them. Once the
analytical theory has been established, we then explain in Sec. III.C.3 how that theory can be used for the direct interpretation of proton-fluence inhomogeneities. Finally, the possibility and implementation of field-reconstruction analysis is discussed in Sec. III.C.4.

2. Analytical theory of proton imaging

In addition to the usually justified assumption that imaging protons behave as test particles, most analytical theories of proton imaging make seven key assumptions:

- **Small-angle deflections**: \( \alpha \ll 1 \). As discussed in Sec. III.A, this assumption generally allows for the trajectories of beam protons to be treated as linear, and thus for deflection angles to be linearly related to line-integrated electromagnetic fields [viz., Eq. (9)]. Using Eq. (10), it can be shown that this condition is equivalent to assuming that the transverse path-integrated electric and/or magnetic field is much smaller than some critical value; specifically, \( | \int_0^{l_{\text{path}}} ds E \parallel | \ll m_p v_0^2 / e \) or \( | \int_0^{l_{\text{path}}} ds B \parallel | \ll m_p c r_0 / e \). Relative to 3.3 MeV protons (one of the two main types of fusion protons produced by \( \text{D}^3 \text{He} \) capsules), these bounds are

\[
| \int_0^{l_{\text{path}}} ds E \parallel | \ll 6.6 \left[ \frac{W_0(\text{MeV})}{3.3 \text{ MeV}} \right] \text{MV}, \quad (11)
\]

\[
| \int_0^{l_{\text{path}}} ds B \parallel | \ll 0.26 \left[ \frac{W_0(\text{MeV})}{3.3 \text{ MeV}} \right]^{1/2} \text{MG cm}, \quad (12)
\]

where \( W_0 \) is the initial energy of the imaging protons. This implies that electric fields with strengths of \( \sim \text{MV/cm} \) or magnetic fields of megagauss strengths permeating the full extent of a millimeter-scale plasma (a typical size for plasmas created during HED experiments) are required for the small-angle deflection assumption to become invalid. Though such large electric and magnetic fields are routinely realized, for example, during the interaction of medium-energy, high-intensity lasers with solid targets, generating them across such a volume has been realized only on the very-high-energy laser facilities, such as the NIF (Meinecke et al., 2022).

- **Point projection**: \( l_{\text{path}} \ll r_d \). The importance of this assumption was also outlined in Sec. III.A: it allows for proton displacements observed at the detector to be treated as being due to velocity perturbations (as opposed to spatial perturbations) acquired through interaction with the electromagnetic fields of the plasma.

- **Small source size**: \( a \ll \ell_{\text{EM}} \), where \( \ell_{\text{EM}} \) is the characteristic length scale of the electromagnetic field in the direction transverse to the trajectory of the proton beam. This assumption allows the proton-beam source to be treated as a point source.

- **Monoenergetic beam**: \( \Delta v_0 \ll v_0 \), where \( \Delta v_0 \) is the characteristic spread of proton speeds in the detected imaging beam. This assumption means that the deflection angles of any constituent protons of the imaging beam that pass along the same trajectory can be treated the same.

- **Instantaneous transit and short pulse**: \( \delta t_p \ll \tau_{\text{EM}} \) and \( \delta t_p \sim l_{\text{path}} / v_0 \ll \tau_{\text{EM}} \), where \( \delta t_p \) is the characteristic duration of the proton beam, \( \tau_{\text{EM}} \) is the characteristic timescale over which the electromagnetic field evolves in the plasma, and \( \delta t_p \) is the transit time of the protons through the plasma. If both the transit time and pulse duration of the proton beam are short compared to \( \tau_{\text{EM}} \), then the electromagnetic field can be treated as electrostatic and/or magnetostatic.

- **Paraxial approximation**: \( \ell_{\text{EM}} \ll 2r_s \). This approximation allows for the proton beam to be treated as an expanding planar “sheet” as it passes through the plasma.

We note that particle-tracing simulations do not necessarily have to make any of these assumptions when artificial proton images are generated; however, if these assumptions are not valid, a correct interpretation of proton images is much more challenging. More detailed discussions of these assumptions can be found elsewhere (Kugland, Ryutov et al., 2012; Bott et al., 2017).

Under these seven approximations, the effect of the electromagnetic fields on the proton beam can be modeled as a “remapping” of the beam’s two-dimensional transverse profile prior to it reaching the detector: any proton with an initial perpendicular position \( x_{\perp,0} \equiv \hat{x}_\perp (0) \) that in the absence of any electromagnetic fields would arrive at the detector plane at the position \( d_{\perp,0} = M x_{\perp,0} \) [where \( M = (r_d + r_a + l_{\text{path}})/r_s \approx (r_s + r_d)/r_s \) is the magnification; see Sec. I.A] instead arrives at the remapped position

\[
d_{\perp}(x_{\perp,0}) = M x_{\perp,0} + \Delta d_{\perp}(x_{\perp,0}), \quad (13)
\]

where

\[
\Delta d_{\perp}(x_{\perp,0}) = \frac{e r_d}{m_p v_0^2} \int_0^{l_{\text{path}}} dz \left[ \hat{E}_\perp \left( x_{\perp,0} \left( 1 + \frac{z}{r_s} \right) + \hat{z} \right) + \frac{\vec{V}}{c} \times \hat{B}_\perp \left( x_{\perp,0} \left( 1 + \frac{z}{r_s} \right) + \hat{z} \right) \right]. \quad (14)
\]

In Eq. (14), \( \hat{z} \) is the unit vector normal to the plane of the detector and \( z \) is the coordinate along that axis. Conservation of proton number within any infinitesimal surface element of the beam’s transverse profile then implies that the distribution \( \Psi(d_{\perp}) \) of protons measured by the detector at position \( d_{\perp} \) is related to the initial distribution \( \Psi_0(x_{\perp,0}) \) via

\[
\Psi[d_{\perp}(x_{\perp,0})] = \sum_{x_{\perp,0}: d_{\perp} = d_{\perp}(x_{\perp,0})} \frac{\Psi_0(x_{\perp,0})}{| \det \nabla_{x_{\perp,0}} [d_{\perp}(x_{\perp,0})] |}. \quad (15)
\]

Equation (15), which is the key analytical relationship between inhomogeneities in the detected proton-fluence and path-integrated electromagnetic fields, is interpreted as follows. The fluence \( \Psi(d_{\perp}) \) of protons measured at position \( d_{\perp} \) on the detector is equal to the sum of initial proton fluences \( \Psi_0(x_{\perp,0}) \) of all of the perpendicular positions \( x_{\perp,0} \) that, after electromagnetic-field-induced deflections, are remapped to \( d_{\perp} \) and divided by a modification factor. This modification factor
characterizes the degree to which the proton beam has been locally focused or defocused at a particular position \( x_{1,0} \) due to the beam’s deflection. Formally, it is the absolute value of the Jacobian determinant of the mapping defined by Eq. (13). The summation is included because in general it is possible that protons from multiple different initial positions \( x_{1,0} \) can in principle contribute to the proton-fluence distribution at the same position \( d_\perp \) on the detector if the deflections of those protons cause the beam to self-intersect before they arrive at the detector. In this situation, the mapping (13) is, in the mathematical sense, noninjective (that is, there is not a unique position \( x_{1,0} \) that maps to \( d_\perp \)). If, by contrast, Eq. (13) is injective, then the summation is unnecessary. For both TNSA and \( \text{D}^\text{3He} \) proton sources, the initial fluence distribution \( \Psi_0(x_{1,0}) \) is to a good approximation uniform over small solid angles; see Sec. II. \( \Psi_0(x_{1,0}) \) is therefore often assumed to be uniform: \( \Psi_0(x_{1,0}) \approx \mathcal{M}^2 \Psi_0 \), where \( \Psi_0 \) is the mean detected proton fluence.

Naively, the mapping (13) seems to depend on four path-integrated components of the electromagnetic field being imaged via the displacement term (14). However, Eq. (14) has a convenient mathematical property: it can be expressed as the gradient of a two-dimensional scalar potential that is a linear combination of path-integrated electromagnetic potentials. More specifically, it can be shown that (Kugland, Ryutov et al., 2012; Bott et al., 2017)

\[
\Delta d_\perp(x_{1,0}) \approx -\frac{e_r d}{m_p v_0^2} \int_0^{l_{\text{inh}}} \frac{\text{d}z}{z} \left\{ \phi \left[ x_{1,0} \left( 1 + \frac{z}{r_s} \right) + z \right] \right\}
- \frac{\Psi_0}{c^2} A_0 \left[ x_{1,0} \left( 1 + \frac{z}{r_s} \right) + z \right], \tag{16}
\]

where \( \phi \) is the electromagnetic scalar potential, \( A \) is the electromagnetic vector potential, and \( A_0 \) is the component parallel to \( \vec{r} \). We deduce that Eq. (13) can be written as

\[
d_\perp(x_{1,0}) \approx \nabla \perp \Psi(x_{1,0}), \tag{17}
\]

where

\[
\Psi(x_{1,0}) \equiv \frac{1}{2} \mathcal{M} \chi_{1,0}^2 + \phi(x_{1,0}), \tag{18}
\]

\[
\phi(x_{1,0}) \equiv \frac{e_r d}{m_p v_0^2} \int_0^{l_{\text{inh}}} \frac{\text{d}z}{z} \left\{ -\phi \left[ x_{1,0} \left( 1 + \frac{z}{r_s} \right) + z \right] \right\}
+ \frac{\Psi_0}{c^2} A_0 \left[ x_{1,0} \left( 1 + \frac{z}{r_s} \right) + z \right]. \tag{19}
\]

Thus, provided that the assumptions underpinning standard analytical theories of proton imaging are valid, detected proton-fluence inhomogeneities are a function of only two path-integrated scalar functions pertaining to the electromagnetic field: a property of vital importance for successfully realizing field reconstruction; see Sec. III.C.4.

3. Analytical interpretations of proton-fluence inhomogeneities

Using Eq. (15), the relation between path-integrated electromagnetic fields and the distribution of proton fluence [a relation that is in turn a function of the two-dimensional mapping (13)], it becomes possible to construct a framework that systematically characterizes all classes of proton-fluence inhomogeneities that can arise in images of arbitrary electromagnetic fields into a few different regimes. As explained in Sec. I.A, the key dimensionless parameter that underpins this framework is the contrast parameter \( \mu \), which is given by (Kugland, Ryutov et al., 2012; Bott et al., 2017)

\[
\mu = \frac{r_0 \delta \alpha}{\mathcal{M}^2 \mathcal{E}^\Delta}. \tag{20}
\]

Physically, this parameter quantifies the relative magnitude \( \mathcal{E}^\Delta \equiv \mathcal{M}^2 \mathcal{E}^\Delta \) of the electromagnetic structures being imaged (including magnification) and the displacements \( \Delta d_\perp \equiv r_0 \delta \alpha \) of protons at the detector acquired due to their interaction with those electromagnetic structures [mathematically, \( \mu \) quantifies the relative magnitude of the two terms in the mapping (13) when their gradient is taken in the denominator of the fraction present on the right-hand side of Eq. (15)]. Depending on the size of \( \mu \), the three regimes of a qualitatively distinct nature for electromagnetic fields with a single characteristic scale are as follows:

1. **Linear regime** (\( \mu \ll 1 \)). In this regime, \( \Delta d_\perp \ll \mathcal{E}^\Delta \), so the characteristic scale of proton-fluence inhomogeneities is similar to that of the electromagnetic fields being imaged. As a result, the relationship between proton-fluence inhomogeneities and path-integrated electromagnetic fields becomes to a good approximation linear (hence the regime’s name), with the characteristic size \( \delta \Psi \) of those inhomogeneities being small compared to the mean proton fluence \( \Psi_0 \): \( \delta \Psi / \Psi_0 \sim \mu \ll 1 \). Indeed, in the linear regime proton-fluence inhomogeneities have a simple physical interpretation in terms of path-integrated charge and current densities. For purely electrostatic fields \( \delta \Psi / \Psi_0 \propto -\int_0^{l_{\text{inh}}} \text{d}\rho \), where \( \rho \) is the charge density in the plasma (Romagnani et al., 2005), while for purely magnetic fields \( \delta \Psi / \Psi_0 \propto -\int_0^{l_{\text{inh}}} \text{d}j_\parallel \), where \( j \) is the magnetohydrodynamic (MHD) current density (Graziani et al., 2017).

2. **Nonlinear injective regime** (\( \mu \lesssim \mu_c \sim 1 \)). In this regime \( \Delta d_\perp \lesssim \mathcal{E}^\Delta \), but with the additional constraint that \( \mu \) is not larger than some critical value \( \mu_c \) at which the proton beam self-intersects prior to reaching the detector on account of spatially inhomogeneous deflections; viz., the mapping (13) remains injective, so the summation in Eq. (15) is not needed. As a result of the comparatively large magnitude of \( \Delta d_\perp \) compared to \( \mathcal{E}^\Delta \), the characteristic scales of proton-fluence inhomogeneities are distorted away from those of the path-integrated electromagnetic fields (inhomogeneities with \( \delta \Psi > \Psi_0 \) are focused, while those with \( \delta \Psi < \Psi_0 \) are defocused) and the magnitude of proton-
fluence inhomogeneities in this regime is typically comparable to the mean proton fluence \( (\delta \Psi \sim \Psi_0) \). The simple physical interpretation of proton-fluence inhomogeneities in terms of path-integrated charge and current structures is no longer quantitative in the nonlinear injective regime, but such relationships still hold qualitatively. We note that the value of \( \mu_c \) depends on the particular electromagnetic-field structure being imaged but is typically of the order of unity.

(3) **Caustic regime** \((\mu \geq \mu_c)\). In this regime \( \Delta d_{\perp} \geq \epsilon_{EM}^{(d)} \), with spatial gradients being sufficiently large that the proton beam self-intersects prior to being detected. As explained in Sec. I.A, this self-intersection leads to the emergence of proton-fluence inhomogeneities known as caustics. Caustics have a specific structure that is unrelated to the electromagnetic fields responsible for them: they attain large magnitudes \( (\delta \Psi \gg \Psi_0) \) in isolated regions and typically occur in pairs; see Kugland, Ryutov et al. (2012) for a detailed discussion of caustics. It follows that the interpretation of proton-fluence inhomogeneities in terms of path-integrated electromagnetic fields is more challenging in the presence of caustics than in their absence, though some successful measurements of simple field structures in this circumstance have been made (Kugland et al., 2012, 2013; Morita et al., 2016; Levesque and Beesley, 2021).

Because \( \mu \) is directly proportional to the deflection angle \( \delta \alpha \), it is linear in the characteristic strength of the electromagnetic field being imaged. By contrast, \( \mu \) is inversely related to the initial proton energy: for magnetic fields \( \mu \propto W_0^{-1/2} \), while for electric fields \( \mu \propto W_0^{-1} \). Thus, a given electromagnetic-field structure can be in any one of the contrast regimes, depending on its strength and the energy of protons being used to perform imaging. Varying the dimensional parameters that describe the imaging diagnostic setup (for instance, \( r_s \) and \( r_d \)) also affects the contrast regime.

We illustrate the key features of the three contrast regimes with a simple numerical example. In this case, we compare the three regimes by choosing one field structure and then generating a sequence of synthetic proton images at increasing characteristic field strengths. We choose an “ellipsoidal blob” magnetic field (Kugland, Ryutov et al., 2012) given by

\[
B = \frac{B_{\text{max}}}{\sqrt{2}} x_{1,0} \times \frac{\hat{z}}{\ell_{\text{ML}}} \exp \left[ -\frac{\left| x_{1,0} \right|^2}{\ell_{\text{ML}}^2} - \frac{(z-z_0)^2}{\ell_{0}^2} - \frac{1}{2} \right],
\]

where \( B_{\text{max}} \) is the maximum strength of the field, \( \ell_{\text{ML}} \) is its perpendicular scale length, \( \ell_{0} \) is its parallel scale, and \( z_0 \) is the \( z \)-coordinate of the field’s central point. The field is visualized in Fig. 14(a).

The spatial distribution of the \( z \)-component of the MHD current density, which is given by

\[
j_z = \frac{e B_{\text{max}}}{8\sqrt{2}\pi\ell_{\text{ML}}} \left( 1 - \frac{\left| x_{1,0} \right|^2}{\ell_{\text{ML}}^2} \right) \times \exp \left[ -\frac{\left| x_{1,0} \right|^2}{\ell_{\text{ML}}^2} - \frac{(z-z_0)^2}{\ell_{0}^2} - \frac{1}{2} \right],
\]

is visualized in Fig. 14(b). Note that for this choice both the path-integrated magnetic field and the MHD current density have approximately the same perpendicular spatial structure as the three-dimensional field itself; see Figs. 14(c) and 14(d). Corresponding proton images of this magnetic field in the linear, nonlinear injective, and caustic regimes are shown in Fig. 15.

In the linear regime, Figs. 15(a) and 15(b) demonstrate that the proton-fluence inhomogeneity \( \delta \Psi \) is indeed small in magnitude compared to the mean fluence \( \Psi_0 \), with that inhomogeneity being approximately proportional to the MHD current. In the nonlinear injective regime, the central part of the ellipsoidal blob (which has \( j_z > 0 \)) is larger in the proton image than in actuality [Fig. 15(c)], and the fluence and MHD current profile no longer agree quantitatively [Fig. 15(d)]. Finally, in the caustic regime [see Figs. 15(e) and 15(f)] two high-amplitude caustic structures demarcate the edge of the ellipsoidal blob, whose structure does not resemble the true value of \( j_z \).

4. **Inverse analysis using electromagnetic-field-reconstruction algorithms**

In addition to providing a framework for the general interpretation of proton-fluence inhomogeneities in terms of the path-integrated fields creating them, further consideration of Eq. (15) reveals the conditions under which direct inversion
of path-integrated electromagnetic fields from proton images is possible: that is, determining $\mathbf{d}_\perp(x_{10})$ directly from $\Psi(d_\perp)$. The key result of previous studies (Bott et al., 2017; Graziani et al., 2017) is that a direct inversion of Eq. (15) from a single image is a well-posed mathematical problem provided that the mapping (13) is injective or, equivalently, that there are no caustics present in the images. In terms of contrast regimes, inversion can be performed in either the linear or the nonlinear injective regime. This finding follows from the observation that Eq. (15) can be written as a Monge-Ampère equation if the mapping (13) is injective,

$$\Psi(\nabla_\perp \psi(x_{10})) = \frac{\partial_0(x_{10})}{\det \nabla_\perp \nabla_\perp \psi(x_{10})}, \quad (23)$$

where $\psi(x_{10})$ is the scalar function defined by Eq. (18) in Sec. III.C.2. In spite of their nonlinearity, Monge-Ampère equations have unique solutions for $\nabla_\perp \psi(x_{10})$ [and thus $d_\perp(x_{10}) \approx \nabla_\perp \psi(x_{10})$] given appropriate Neumann boundary conditions. In the case of general electromagnetic fields, more information is needed to distinguish between path-integrated electrostatic and magnetic fields, but in the case where one dominates over the other, the path-integrated electrostatic or magnetic field in a plasma can be reconstructed.

Various different “field-reconstruction” algorithms for recovering path-integrated electromagnetic fields directly from proton-fluence inhomogeneities have been proposed. In the linear regime, it can be shown that the inversion process is simply equivalent to solving a Poisson equation for the scalar function $\psi(x_{10})$ (Romagnani et al., 2005; Kugland, Ryutov et al., 2012): $\nabla_\perp^2 \psi(x_{10}) = -M \psi_0$. However, a later study by Graziani et al. (2017) found that inversion quickly fails for anything but small values of $\mu$; they proposed overcoming this by including first-order terms in the expansion, solving the resulting equations with the PRALINE code (Graziani et al., 2017). In the $\mu \leq 1$ regime, descriptions of three different algorithms have been published: a Voronoi-diagram method to reconstruct path-integrated magnetic fields by Kasim et al. (2017); the PROBLEM code by Bott et al. (2017), which uses a nonlinear diffusion method proposed by Sulman, Williams, and Russell (2011) to solve the same problem; and finally a trained neural network by Chen et al. (2017) (who also trained their network to resolve the 3D structure of ellipsoid blobs, though it is unlikely this approach is applicable to more general electromagnetic fields). Other algorithms have been used to reconstruct electromagnetic fields, in particular, experiments (Schaeffer et al., 2019; Campbell et al., 2020; Levesque et al., 2022), though full details of these codes have not yet been published. Comparisons of the outputs of these codes are currently an active research effort (Davies and Heuer, 2022).

By contrast, the possibility of performing direct inversion analysis from a single proton image if caustics are present has been shown not to be a well-posed mathematical problem: multiple path-integrated electromagnetic-field “solutions” exist for a single proton-fluence distribution $\Psi$. In this situation, it can be proven that the solution to the Monge-Ampère equation [Eq. (23)] minimizes the functional

$$C(\Delta d_\perp) = \int d^2 x_{10} |\Delta d_\perp(x_{10})|^2 \partial_0(x_{10}). \quad (24)$$

In the case of a uniform initial proton-fluence distribution, the Monge-Ampère solution therefore minimizes the root mean square of proton displacements over the space of all possible solutions. In practice, if $|\mu - \mu_c| < 1$, the “family” of possible solutions associated with a particular distribution $\Psi$ is typically constrained to be similar to the Monge-Ampère solution. Thus, outputs of reconstruction algorithms are usually close to the “true” result for a point source of protons (Kasim et al., 2017), though this rule of thumb becomes much

FIG. 15. Comparison of contrast-parameter ($\mu$) regimes for proton images of the ellipsoidal blob magnetic field described in Fig. 14. Images in the (a) linear, (c) nonlinear injective, and (e) caustic regimes are shown, as are proton-fluence lineouts along x (at y = 0) in (b), (d), and (f), respectively. Superimposed onto the images is the mapping $d_\perp(x_{10}) = \mathbf{d}_\perp(x_{10})$ for each case (the solid lines). To generate the images, protons were simulated with $v_0 = 5.31 \times 10^7$ cm/s (corresponding to 14.7 MeV protons), a setup with $c_0 = 1$ and 30 cm, and a field $B_{\text{max}} = 10, 200, 500$ kG for the linear, nonlinear injective, and caustic regime images, respectively. The resolution of the images is $200 \times 200$ pixels, with a mean proton density per pixel of 100. Following previous conventions for the ellipsoidal blob (Kugland, Ryutov et al., 2012), one can define $\mu$ as $\mu = \sqrt{\pi \sigma_{\perp}, B_{\text{max}}^2 M_{\perp}}/m_{\text{p}}c\nu_0 M_{\perp} M_{\perp}$, where $A$ is an order-unity constant of proportionality.
less robust if realistic proton source sizes are taken into account (Bott et al., 2017). If $|\mu - \mu_c| \gtrsim 1$, then previous studies (Bott et al., 2017) have shown that the Monge-Ampère solution can return significant underestimates of characteristic path-integrated electromagnetic-field strengths compared to those of the true electromagnetic field. Systematically extracting information about path-integrated electromagnetic fields from proton images that contain caustics is therefore an outstanding research problem in proton-image analysis (although some recent progress has been made on this; see Sec. V.D).

Although field-reconstruction algorithms have been successfully applied to real experimental data (Tzeferacos et al., 2018; Schaeffer et al., 2019; Campbell et al., 2020; Bott et al., 2021a, 2021b; Levesque et al., 2022), these efforts have shown that several technical issues can arise in the process of performing such analysis. First of these is the finding that field-reconstruction algorithms are sensitive to any large-scale variations in the initial proton-fluence profile $\Psi(x,0)$. For example, it has been shown that two qualitatively different path-integrated magnetic fields can give the same proton image with only subtly different initial profiles (Bott et al., 2017). While both TNSA and DHe proton sources can produce beams whose transverse spatial inhomogeneities are small compared to the mean fluence on submillimeter plasma scales (Manuel, Zylstra et al., 2012), it has proven challenging to avoid significant inhomogeneities on larger scales. Studies aimed at overcoming this problem are ongoing, but possible remedies include high-pass filtering of either images or reconstructed path-integrated fields in order to isolate only those outputs for which uncertainties are not too large (Bott et al., 2017; Kasim et al., 2017), applying constrained polynomial or Gaussian (as opposed to uniform) models for the initial profile (Palmer et al., 2019; Fox et al., 2020), and using Bayesian inference conditioned on the well-characterized properties of the initial proton-beam inhomogeneities (Kasim et al., 2019). Another issue that is particularly important for fusion-capusle proton sources is the effect of a finite source size. It has been demonstrated (Bott et al., 2017) that the source’s finite size reduces the characteristic value of $\mu_c$ below which field-reconstruction algorithms return accurate results compared with the case of a genuine point source. Bott et al. (2017) proposed the Lucy-Richardson deconvolution algorithm as a way to mitigate this issue, but further study of more robust techniques is warranted. Finally, field-reconstruction algorithms neglect the “blurring” effect of scattering of beam protons due to Coulomb collisions on proton images. However, this blurring is usually significant in experiments involving dense plasmas and thus should not be ignored in future studies; see Sec. V.C.

D. Comparing particle-tracing and analytical modeling techniques

In Secs. III.B and III.C, we reviewed the use of particle-tracing simulations and analytical theory, respectively, for analyzing proton images; providing a comparative discussion of the two methodologies with respect to each other is therefore apt. The main advantage that particle-tracing simulations have over analytical modeling is the possibility of avoiding the approximations that analytical modeling has to make in order to be tractable. These approximations involve the physics underpinning the interaction of the beam with the plasma (for instance, scattering), the precise properties of the proton beam’s source, and the geometry of the imaging setup; see Sec. III.C.2. Avoiding some of these approximations is vital for certain categories of laser-plasma experiments, such as those investigating ultrafast laser-plasma dynamics; see Sec. IV.G. However, it is challenging for particle-tracing simulations to overcome one of the central challenges for all forward-modeling techniques (the possibility of multiple qualitatively distinct solutions that are all consistent with the input data) without recourse to analytically derived results (such as uniqueness). In addition, field-reconstruction algorithms based on analytical modeling allow for images of complicated electromagnetic-field structures to be analyzed in situations when HEDP simulations either are unavailable or are not able to reproduce the relevant physics correctly. All said, we emphasize that either technique can be highly effective, but also that the most robust analysis usually involves both.

IV. PROTON-IMAGING EXPERIMENTS

To demonstrate the variety of phenomena that can be investigated using proton imaging, we provide here a survey of the different types of experiments that have been performed using the diagnostic. Examples include applications with spontaneously generated magnetic fields (Sec. IV.A), magnetic reconnection (Sec. IV.B), Weibel instabilities (Sec. IV.C), shocks (Sec. IV.D), jets (Sec. IV.E), turbulence and dynamics (Sec. IV.F), ultrafast dynamics (Sec. IV.G), hydrodynamic instabilities (Sec. IV.H), and ICF (Sec. IV.I).

A. Magnetic-Field Generation

Magnetic fields can be spontaneously generated by several different mechanisms in initially unmagnetized plasmas, and proton imaging has been used to explore and characterize these processes in various laser-plasma experiments. Understanding these possible sources of magnetic fields is an important research area in HED plasma physics because basic processes such as heat transport can be profoundly altered if magnetic fields become strong enough to magnetize the plasma’s constituent particles (that is, reduce their Larmor radii below their respective Coulomb mean free paths). A detailed discussion of the many sources of magnetic fields in hot laser-produced plasma was given by Haines (1986); here we focus on the most notable ones and their investigation using proton imaging.

One of the first mechanisms for generating magnetic fields in plasmas to be identified theoretically [and also one of the first to be observed in experiments (Stamper et al., 1971)] is the Biermann battery, whereby magnetic fields are generated by misaligned electron density and pressure gradients (Biermann and Schluter, 1951). Within the framework of
extended MHD, the Biermann battery can be modeled as a source term in the induction equation,

\[
\frac{\partial B}{\partial t} = -\frac{c \nabla n_e \times \nabla p_e}{en_e^2} - c \nabla \times \left( \frac{\mu_0}{e} \nabla T_e \right) + \nabla \times (v_B \times B) - \nabla \times (\eta \nabla \times B).
\]  

(25)

In Eq. (25) the first source term on the right-hand side is the Biermann battery (the second source term, which is often neglected in modeling, is associated with ionization). It can be shown that the Biermann battery term generates a field \(\delta B\) in a time interval \(\delta t\) of magnitude \(\delta B \sim \delta t c k_B \nabla T_e \times \nabla n_e / en_e\). Once generated, this Biermann field then evolves through advection at a characteristic bulk-flow velocity \(v_B\) [the third term of Eq. (25)] and through diffusion by the resistivity \(\eta\) [the fourth term of Eq. (25)] (Haines, 1986). Since nonparallel plasma temperature and density gradients are common in plasmas, magnetic-field generation by the Biermann battery is ubiquitous in HED experiments and is a frequent subject of proton imaging. For a laser pulse interacting with a solid target, the electron density gradient is primarily in the target normal direction, whereas the electron temperature gradient is primarily radial, meaning that an azimuthal magnetic field is generated around the laser focal spot. The rate of field generation is therefore dependent on processes like the energy transfer from the laser to the plasma, and parameters such as the focal spot size and intensity profile.

With the advent of proton imaging, a number of experiments have studied the generation of magnetic fields near the surface of laser-driven targets by the Biermann battery (Li et al., 2006b; Nilson et al., 2006; Cecchetti et al., 2009; Petrasso et al., 2009; Willingale et al., 2011a, 2013; Gao et al., 2012, 2013, 2015; Lancia et al., 2014; Campbell et al., 2020). The first measurements used grid deflectometry to gauge the deflections of a known periodicity mesh to infer the path-integrated magnetic fields; see Sec. II.C.2. Proton-beam deflections are affected by the direction of the projection of the protons. For a “front”-projection geometry, where protons travel from the interaction side of the main target to the rear, the \(v \times B\) Lorentz force primarily deflects the proton beam radially inward. A “rear”-projection geometry, where protons first pass through the target before observing the front-side magnetic fields, produces an outward deflection. This was illustrated by Cecchetti et al. (2009). Figure 16 presents experimental data using a TNSA source in front-projection [Fig. 16(a)] and rear-projection geometries [Fig. 16(c)]. Figures 16(b) and 16(d) are the particle-tracking calculations for an idealized magnetic torus in front- and rear-projection geometries, respectively. Similar data using a D3He source were presented by Petrasso et al. (2009). While the proton images naively make the extent and magnitude of the fields appear to be different for the two geometries, comparisons to analytical field maps show that the strength and scale of the fields are in fact similar.

Biermann-battery-generated fields can be up to a megagauss or more and evolve on nonexistent timescales. These measurements are to within an order of magnitude, but not necessarily in exact agreement with, simulation predictions (Li et al., 2006b). Numerical modeling typically consists of MHD simulations that include a Biermann battery source term, resistive magnetic diffusion, and fluid advection [see Eq. (25)] and often Nernst advection, right-hand-leg heat flow, and radiation. Measurements have confirmed that, once generated, magnetic fields can indeed be advected by the bulk plasma motion, i.e., at the ion fluid velocity, or the hot electron flux can transport the magnetic field at a faster speed through the Nernst effect (Nishiguchi et al., 1984; Willingale et al., 2010b) and other effects (Lancia et al., 2014; Gao et al., 2015). Proton-imaging experiments by Campbell et al. (2020) have shown that varying the target material alters the field generation (see Fig. 17) and even the development of a double ablation front for mid-Z materials. Careful analysis of the field measurements to quantify total magnetic flux show that kinetic effects can suppress Biermann battery field generation in laser-plasma interactions (Campbell et al., 2022).

Wilks et al. (1992) proposed that magnetic fields much stronger than those generated by the Biermann battery can be created by relativistic laser interactions (> 10^18 W cm^-2) due to currents produced by suprathermal electrons accelerated in the evanescent region of the laser wave that propagate deep into the interior of the plasma. This magnetic field is in the azimuthal direction about the laser axis of propagation, and the peak field extends for about an anomalous skin depth into the plasma (i.e., \(d = [|c/\omega_{pe}|(v_{te}/\omega_{he})]^{1/2}\), where \(v_{te}\) is the electron thermal velocity). Mason and Tabak (1998) predicted the generation of fields up to 250 MG in the overdense plasma for moderately relativistic interactions. Measurements of these short-pulse, relativistic-intensity-generated magnetic fields have been measured using

FIG. 16. Experimental images of TNSA proton deflectometry of a laser-generated plasma shown for geometries in which protons pass first through (a) the plasma (front projection) or (c) target (rear projection). Corresponding synthetic proton images were created with the particle-tracing code PTRACE using an idealized magnetic-field torus in (b) front and (d) rear geometries, respectively. Adapted from Cecchetti et al., 2009.
FIG. 17. Top row: proton images of the fields generated from laser ablation of different target materials at a time of 0.75 ns into a 1 ns, 1025 J interaction. The proton energy is 37.3 MeV for CH, Al, and Cu + Al, and 30.7 MeV for Au + Al. Bottom left panel: radial lineouts of the proton fluence (J) normalized by the mean inferred reference profile (J_0). Bottom right panel: the resulting reconstructed field profiles. For Al and Cu + Al, the results of double-Gaussian fitting are shown with shaded regions. Adapted from Campbell et al., 2020.

TNSA protons by Sarri et al. (2012). One significant difference compared to the lower-intensity measurements is that large fields are expected to be present on both the front and rear sides of the target. Hot electrons rapidly move through the target to expand into the vacuum at the front and rear, creating time-varying sheath fields that generate opposing magnetic fields on the front and rear target surfaces. This means that on one side of the target the proton beam is deflected radially inward while on the other side it is deflected outward from the interaction region, which significantly complicates the analysis and interpretation of the proton data.

A different method for creating magnetic fields with laser-plasma interactions is through laser-driven coils (Gao et al., 2016; Peebles et al., 2022). In this approach, a laser is used to heat and eject electrons from a plate so that a current is drawn through a loop connected to the plate. The interaction region within the loop, a volume of the order of 1 mm^3, contains a strong axial magnetic field that can be used as an externally applied field for other experiments. Peebles et al. (2020) measured axial fields of up to 65 ± 15 T.

B. Magnetic reconnection

Magnetic reconnection (Yamada, Kulsrud, and Ji, 2010) is a physical process whereby the magnetic-field topology is rearranged, dissipating magnetic energy in a plasma into kinetic energy. It is a prevalent phenomenon throughout the Universe that occurs under many different conditions: for example, within the solar corona, where it leads to solar flares and coronal mass ejections (Parker, 1957), between the solar wind and Earth’s magnetosphere, and during fusion plasma instabilities (Taylor, 1986). Breaking and reconnecting magnetic-field lines at observed rates require dissipation mechanisms to function at rates greater than allowed by classical resistivity (Yamada, Kulsrud, and Ji, 2010). Consequently there are many open questions to be investigated, including the temporal and spatial scales of the reconnection, the role of dynamical processes like plasmoid formation, and the final energy partition of the system. Furthermore, there are a wide range of reconnection regimes to explore due to how the magnetization, collisionality, and symmetry of the system affect the mediation of the reconnection process.

Laser-driven magnetic reconnection is a convenient way to study impulsive, strongly driven reconnection physics in the laboratory. Using proton imaging to diagnose the magnetic fields, the first experimental demonstration using lasers was performed by Nilson et al. (2006), with the basic experimental configuration shown in Fig. 18. These experiments used two neighboring high-energy, nanosecond duration laser pulses to produce self-generated azimuthal magnetic fields through the Biermann battery mechanism (Nilson et al., 2006, 2008; Li et al., 2007; Willingale et al., 2010a; Zhong et al., 2010; Fox, Bhattacharjee, and Gemaschewski, 2011, 2012; Dong et al., 2012; Rosenberg et al., 2012); see Sec. IV.A. The magnetic fields were then advected out either by the frozen-in flow or by heat transport via the Nernst effect, leading the opposing magnetic fields to be driven together in the midplane between the two focal spots. A key feature of such experiments is that the so-called plasma β, defined as the ratio of thermal to magnetic pressure, is typically large.

Numerous high-quality proton-imaging measurements have been made of magnetic fields in reconnection laser-plasma experiments of this type. For example, Li et al. (2007) and Willingale et al. (2011b) observed the rearrangement of the magnetic field’s topology using proton imaging (as well as elevated plasma temperatures in the midplane region using Thomson scattering and plasma jets emanating from the reconnection plane using optical probing). Experiments by Rosenberg, Li, Fox, Zylstra et al. (2015) used proton imaging

FIG. 18. Schematic of a laser-driven magnetic reconnection geometry. The opposing magnetic fields are driven together in the midplane between the laser focal spots. The fields can be probed at different times to observe the dynamics. Adapted from Nilson et al., 2006.
to observe the slowing of the reconnection rate as the plasma inflows weaken and investigated the effect of asymmetric field structures (Rosenberg, Li, Fox, Igumenshchev et al., 2015). Experimental measurements using proton deflectometry by Tubman et al. (2021) showed anomalously fast reconnection in weakly collisional colliding laser plasmas.

These measurements (and also concurrent measurements from other complimentary diagnostics of the plasma conditions) have prompted new theoretical and numerical modeling studies of the high-$\beta$ reconnection regime, in turn helping advance our understanding of reconnection processes more generally. For example, Fox, Bhattacharjee, and Germaschewski (2011, 2012) performed numerical modeling of laser-driven experiments using particle-in-cell simulations (both with and without a collision operator) and noted the importance of the flux pile up to the reconnection process. Joglekar et al. (2014) used a fully implicit 2D Vlasov-Fokker-Planck code to show that in high-$\beta$ laser-generated plasmas heat flows rather than Alfvénic flows dictate the reconnection rate. Supporting modeling identified the role of anisotropic heat flows rather than Alfvénic flows dictate the reconnection rate.

Proton imaging has also been used successfully to diagnose magnetic fields in other types of laser-driven reconnection experiments. For example, Palmer et al. (2019) explored reconnection of fields generated by higher ($\sim 10^{18}$ W cm$^{-2}$) laser intensities through proton deflectometry measurements. Figure 19 illustrates a time history of data taken along with the 2D magnetic-field maps reconstructed from proton images using a field-reconstruction algorithm; see Sec. III.C. These maps showed faster dissipation of magnetic fields at the midplane compared to the outer plane, confirming that reconnection was taking place in the experiment on a time-scale of tens of picoseconds.

Alternative laser-driven reconnection geometries have also been developed and studied using proton imaging. Fiksel et al. (2014) employed externally applied opposing magnetic fields driven together by expanding laser plumes, and Chien et al. (2019) used laser interactions to drive currents through U-shaped coils configured in a reconnection geometry.

C. Weibel instabilities

Weibel-type filamentation instabilities (Fried, 1959; Weibel, 1959; Davidson et al., 1972) are ubiquitous in laboratory and astrophysical plasmas. They arise in plasmas whose particle distribution functions have significant velocity-space anisotropy. The velocity-space anisotropy includes cases where the temperature $T_j = \langle \mathbf{v}_j^2 \rangle / 2$ differs among the three directions, where $j$ is one of $(x, y, z)$, and can be driven by counterstreaming particle beams that produce an effective anisotropy. The counterstreaming between a hot forward particle population, balanced by a cold return current, which arises in situations of large heat flux, is another source of anisotropy important for electron-driven Weibel. The instability grows predominantly with wave number $k$ aligned along the cold direction(s). The instability can play a broad role in plasmas, including magnetic-field generation in the early Universe and magnetic-field generation and amplification in high-Mach-number shocks.

The fundamental Weibel mechanism is that the large counterstreaming currents along the hot direction tend to pinch and coalesce into current-carrying filaments, and the forces driving coalescence are sufficient to overcome the transverse plasma pressure along the cold directions. Transverse magnetic fields associated with the current filaments then deflect the particle trajectories, reinforcing the filamentation and leading to a positive feedback. The nonlinear regime includes rich physics such as the kinking and remerging of magnetized flux tubes.

Ion-driven Weibel instabilities are important in astrophysical plasmas, as the large bulk-flow energy density of ions $j n_i V_i^2 / 2$ can be greater than analogous energy densities of the electron population and can therefore be a larger reservoir of free energy for the Weibel process, producing stronger magnetic fields at larger scales. The ion-Weibel instability was identified in laboratory laser-driven experiments using proton imaging (Fox et al., 2013; Huntington et al., 2015; Park et al., 2015). In the experiments, two plasma plumes were ablated from opposing targets and were then collided. The high

3We note that the large magnification used for this experiment meant that the interaction image extended close to the edge of the proton beam, necessitating detailed modeling of the assumed unperturbed proton fluence; by reducing the magnification so that the unperturbed region around the interaction is larger, it becomes easier to infer the unperturbed proton fluence and thus reduces the potential error of the reconstruction method.

FIG. 19. Time series for two shots of a high-intensity laser-plasma-driven reconnection experiment in a geometry similar to Fig. 18. Top row: the measured proton fluence at the detector plane. Middle row: the calculated undisturbed beam fluence at the detector plane. Bottom row: the retrieved path-integrated magnetic fields at the interaction plane. The white contours with arrows show the topology of the calculated magnetic fields. Adapted from Palmer et al., 2019.
temperature and low density of the ablation flows sets up counterstreaming ion populations in the interaction region. Proton imaging directly imaged the magnetized filaments produced in this interaction region by the ion-Weibel instability [Fig. 20(a)]. Huntington et al. (2015) and Park et al. (2015) measured statistics of the observed filamentary structures, which compared favorably to nonlinear kinetic simulations [Fig. 20(b)].

The electron-Weibel instability is important in relativistic plasmas with strong beam currents and is an important energy-coupling process that can lead to anomalous stopping of relativistic electron beams driven by short-pulse laser-plasma interactions. This type of instability was observed in face-on and side-on proton probing experiments (Borghesi et al., 2002a; Quinn et al., 2012; Ruyer et al., 2020). Filamentlike magnetic-field structures were observed to persist for an extended time period, a finding that could explain sustained, spatially elongated structures observed in various astrophysical environments.

D. Shocks

Proton imaging has been instrumental in studying the field structures of shocks in laboratory astrophysics experiments. These shocks act to dissipate kinetic ram pressure in systems with supersonic flows and are commonly found in heliospheric and astrophysical systems, including planetary bow shocks, jets, supernova remnants, and galaxy clusters, and are often associated with extremely energetic particles. A key component of shocks is their strong electromagnetic fields. In magnetized shocks, which propagate through a preexisting magnetic field, the global structure of the shock is defined by a jump in the magnetic field on ion kinetic scales, while strong electric fields in the shock layer can help mediate dissipation by reflecting incoming ions. Similarly, in electrostatic shocks, the shock layer is defined by electric fields on electron kinetic scales. Meanwhile, in electromagnetic shocks, initially unmagnetized counterpropagating plasmas can spontaneously generate magnetic fields through streaming instabilities (such as Weibel ones), leading to shock formation.

Romagnani et al. (2008a) performed the first experiments with proton imaging to study shocks. They used a high-intensity laser to create a supersonic plasma plume that expanded into an unmagnetized low-density ambient plasma, driving a collisionless electrostatic shock. TNSA protons were then used to probe the interaction. Proton images, and electric fields reconstructed from the images, showed modulations of the shock front consistent with shock theory and electron kinetic scales; see Fig. 21. The shock speed was estimated by comparing features between different proton images within an RCF stack. A follow-up experiment by Ahmed et al. (2013) used TNSA proton imaging to provide further details about how the electrostatic potential in the shock layer evolves during electrostatic shock formation.

Schaeffer et al. (2017) first probed magnetized collisionless shocks, using a high-energy laser to drive a supermagnetosonic piston plasma through a magnetized ambient plasma, generating a collisionless shock in the ambient plasma. The shock was probed with TNSA protons, and the resulting proton images showed large proton-fluence variations followed by uniform fluence. Using a 1D reconstruction technique, the fluence variations were shown to correspond to strong magnetic-field compressions at the shock front and a diamagnetic cavity created by the piston behind the shock, which is consistent with features observed in PIC simulations. Further experiments (Schaeffer et al., 2019) used D³He protons to image the fields in a magnetized shock precursor;
Ablated jet plasma. Other experiments by Hua et al. attributed to seed Biermann battery fields embedded in the significantly faster than expected from theory, which was formation of a shock and Weibel filaments on timescales solutions; see Sec. V. D for further discussion.

Proton energies (times) to break the degeneracy of the reconstruct the path-integrated magnetic fields, as well as two multiple angles with TNSA protons. By comparing the proton created a strong collisional shock, which was probed from the shock tube. High-energy lasers incident on one end of the tube studied self-generated electromagnetic fields in shocks using a high-energy laser to ablate a target, which generated a jet plasma that expanded into a gasbag. The collision created a counterpropagating plasma that was imaged with D³He protons. The proton images showed the formation of a shock and Weibel filaments on timescales significantly faster than expected from theory, which was attributed to seed Biermann battery fields embedded in the ablated jet plasma. Other experiments by Hua et al. (2019) studied self-generated electromagnetic fields in shocks using a shock tube. High-energy lasers incident on one end of the tube created a strong collisional shock, which was probed from multiple angles with TNSA protons. By comparing the proton images from different directions, they showed that magnetic fields, self-generated through the Biermann battery effect, dominated the shock structure, and that electric fields were relatively insignificant.

Levesque and Beesley (2021) utilized proton imaging to study laser-driven bow shocks, which led to the development of a new technique for analyzing the proton data (Levesque et al., 2022). The technique utilizes caustic features to help reconstruct the path-integrated magnetic fields, as well as two proton energies (times) to break the degeneracy of the solutions; see Sec. V. D for further discussion.

E. Jets

An important use for proton imaging has been in laboratory astrophysics experiments that have investigated the dynamical effect of magnetic fields on supersonic plasma jets. Various astrophysical systems, including active galactic nuclei, pulsar wind nebulae, and young stellar objects, are associated with magnetized jets and outflows. The magnetic fields are thought to explain a number of observed phenomena in these jets, including collimation, clumping, and kinking. In certain conditions, laser-produced plasma jets can be treated as rescaled analogs for astrophysical jets, meaning that tailored laboratory experiments can shed light on these astrophysical phenomena (Blackman and Lebedev, 2022).

Loupias et al. (2009) carried out the first such experiment using proton imaging to compare the expansion of a front-side blowoff plasma jet into vacuum with that of a similar jet into an ambient gas and found tentative evidence for electromagnetic fields at the gas-jet boundary from their ~3–5 MeV proton-imaging data. More recently the evolution of the MHD interchange and kink instabilities in a jet created by the irradiation of a cone-shaped target were observed by Li et al. (2016). The perturbed magnetic fields associated with both MHD instabilities manifested as quasiperiodic proton-fluence structures in the proton-imaging data. It was then demonstrated that this laboratory jet was a reasonable analog to the Crab Nebula jet under appropriate rescaling, supporting the idea that MHD instabilities provide a plausible explanation for the periodic oscillations in the Crab Nebula jet’s direction that were previously detected by the Chandra X-ray Observatory.

By contrast, another experiment by Gao et al. (2019) successfully realized magnetically collimated, stable supersonic jets using laser beams focused onto a hollow ring configuration. The combined use of an electromagnetic-field-reconstruction algorithm applied directly to the proton-imaging data and particle-tracing with FLASH simulations showed that the experiment realized megagauss magnetic fields; see Fig. 23. Given other plasma jet parameters, fields of this strength were sufficient to efficiently collimate the jet, as

FIG. 21. Data from electrostatic shock experiments in which a dense laser-driven plasma expands into a low-density background plasma. (a) Proton-imaging data taken at the peak of the interaction pulse with 7 MeV TNSA protons. Note the strong modulation associated with the ablating plasma in region I and the modulated pattern ahead of the shock front possibly associated with a reflected ion bunch in region IV. The arrow indicates the laser beam direction. (b)–(c) Details and RCF optical density lineout corresponding to region II, showing modulations associated with a train of solitons. (d)–(k) Details about region III and corresponding lineouts of the probe proton density δn_p/n_p, the reconstructed electric field E, and the reconstructed normalized ion velocity u_i/λ in the cases of (d)–(g) an ion acoustic soliton and (h)–(k) a collisionless shock wave [the collisionless shock detail corresponds to a different shot (not shown)]. Adapted from Romagnani et al., 2008b.
well as to realize magnetization parameters (such as the plasma beta and the Hall parameter) of direct relevance to astrophysical systems. Previously Manuel et al. (2015) first used proton imaging to look at a magnetized jet with inconclusive results.

In addition to understanding the dynamics of individual jets, Li et al. (2013) studied the evolution of magnetic fields in colliding plasma jets at both collinear and noncollinear angles with the aid of proton imaging. These measurements have
been used to show that the underlying physics of collisions between sufficiently supersonic jets cannot be adequately described by single-fluid hydrodynamics due to the low interjet collisionality between constituent particles, instead requiring two-fluid or kinetic models.

F. Turbulence and dynamos

In recent years, proton imaging has come to play an important role in diagnosing magnetic fields in experiments investigating the evolution of turbulent laser plasmas. Of particular note are a series of experiments that have investigated magnetic-field amplification by turbulent plasma motions on various high-energy laser facilities. It is a long-standing theoretical prediction that turbulent, weakly magnetized plasmas with sufficiently large magnetic Reynolds numbers $\text{Re} \equiv \frac{u_{\text{rms}} L}{\eta}$ [where $u_{\text{rms}}$ is the root-mean-square (rms) turbulent velocity, $L$ is the driving scale of the turbulence, and $\eta$ is the plasma’s resistivity] can support sustained magnetic-field amplification until dynamical magnetic-field strengths are attained, a mechanism known as the fluctuation or small-scale turbulent dynamo. This mechanism, which provides a plausible explanation for the magnetic fields ubiquitously observed in various astrophysical environments, has been seen in numerous MHD and more recently kinetic simulations [see Rincon (2019) for a review], but had not been observed in any laboratory experiments.

Tzeferacos et al. (2018) first realized a small-scale turbulent laser-plasma dynamo on the OMEGA laser. Proton imaging with a $^3\text{He}$ source helped confirm the formation of a dynamo via measurements of magnetic fields both at the formation of the turbulent plasma and several nanoseconds later. More specifically the application of magnetic-field-reconstruction algorithms to 15 MeV proton radiographs yielded two-dimensional maps of path-integrated stochastic magnetic fields at both times. Further analysis of these maps under various assumed statistical symmetries allowed for values of the rms strength of the magnetic field, the magnetic-energy spectrum, and a bound on the maximum magnetic-field strength to be inferred. This analysis showed that magnetic energy was amplified $\sim 600$ times during the course of the experiment, with the characteristic magnetic energies post amplification being a finite fraction of the turbulent kinetic energy. Particle-tracing simulations applied to the magnetic field outputted by MHD simulations of the experiment using the code FLASH corroborated these findings (Tzeferacos et al., 2017).

Bott et al. (2021b) also used proton imaging in a related manner for several subsequent experiments on this topic. Time-resolved measurements in a turbulent plasma with order-unity magnetic Prandtl number showed the evolution of stochastic magnetic fields being amplified by the fluctuation dynamo. The proton data were characterized by applying direct inversion analysis to a time sequence of proton images, and path-integrated magnetic-field maps were obtained; see Fig. 24.

Other related experiments include observations of inefficient magnetic-field amplification by supersonic plasma turbulence (Bott et al., 2021a), measurements of the transport of high-energy charged particles through intermittent magnetic fields (Chen et al., 2020), and a demonstration that the key properties of a particular laser-plasma dynamo were insensitive to the plasma’s initial conditions (Bott et al., 2022). Proton imaging was also fielded as part of a recent experiment at the NIF by Meinecke et al. (2022) that investigated the suppression of heat conduction in magnetized turbulent plasmas. However, the megagauss fields realized in that experiment were sufficiently strong, and the characteristic deflection angle of 14.7 MeV protons was sufficiently large, that electromagnetic-field-reconstruction algorithms used in previous experiments could not reasonably be applied. To overcome this, alternative diagnostic approaches including proton-beam truncation using slits and pinholes were employed to recover the rms and maximum magnetic-field strengths realized in the experiment.

![Fig. 24. Proton images of stochastic magnetic fields amplified by a turbulent laser-plasma dynamo. Left panel: Annotated photograph of the experimental platform used to create the dynamo. The $^3\text{He}$ 14.7 MeV proton images collected during the experiment are shown in the top middle panels, with corresponding path-integrated magnetic-field maps recovered from these images shown in the bottom middle panels. Estimates of the rms and maximum magnetic-field strengths, as well as the correlation lengths, were then recovered from these maps using statistical methods; see the error bars in the right panel. The results were compared with similar quantities inferred from 3D FLASH simulations of the experiment (see the triangle markers in the right panel), as well as the same quantities computed directly (the solid lines). Adapted from Bott et al., 2021b.](image-url)
Andy Sha Liao et al. (2019) proposed a different dynamo experiment using turbulent ablated blowoff plasmas and conducted experiments showing that a magnetic dynamo was created (Liao et al., 2022). This and other new experiments might open up more potential platforms to study HED dynamos with astrophysical relevance in the laboratory.

G. Ultrafast dynamics

The picosecond-scale temporal resolution obtainable when using a TNSA probe has been exploited in several experiments to investigate the ultrafast dynamics following high-intensity, short-pulse laser interaction with a target or a plasma. Large and transient electromagnetic fields are generated in these interactions, in connection with the large flows of relativistic electrons generated in the irradiated portion of the target. The most energetic electrons typically escape from the target, charging it positively on picosecond timescales. The process of target charge-up and subsequent discharge was observed in some of the earliest proton-imaging experiments investigating 50 TW interactions with wire targets (Borghesi et al., 2003); see Fig. 25.

Quasi-instantaneous target charge-up was observed via strong deflection of protons away from the target surface, causing the appearance of caustics, which were associated with a transverse electric field with an amplitude at the surface of the order of $10^{10}$ V/m. In these measurements, the shadow of the unperturbed wire, imprinted by collisional scattering of the protons crossing the wire before it is charged [Fig. 25(a)], acts as a useful fiducial feature for the interpretation of the data: the shadow also appears in the layers in the pack corresponding to later times, as a dose is deposited in these layers by the protons forming the image in Fig. 25(a); see Sec. II.C.4. The charge-discharge cycle was characterized more extensively in follow-up experiments by Quinn et al. (2009c), which employed a high-energy proton probe (up to 40 MeV) generated by the Vulcan petawatt laser. This probed a portion of the wire away from the interaction region, which allowed the characteristic time (~20 ps) over which target neutralization occurs to be reconstructed. The target charging measured in these studies was found to be consistent with the number of escaping electrons evaluated from self-consistent models. A modified proton-imaging setup (Quinn et al., 2009b), in which the wire was tilted away from the vertical position within the plane containing the proton beam’s axis, allowed the early phases of these dynamics, in which an electric field is seen to spread from the interaction region along the target at a velocity close to the speed of light, to be resolved (Quinn et al., 2009a). The field was interpreted, with the help of PIC simulations, as a surface electromagnetic mode generated by the ultrafast motion of the electrons escaping from the target (similar to the emission from a transient antenna).

Kar et al. (2016) further investigated the dynamics of this surface mode by characterizing the propagation along millimeter-length wires connected to the laser-irradiated target. These studies, mostly carried out employing a self-imaging scheme, where the proton probe is provided by the same laser-irradiated target that produces the electromagnetic mode (Ahmed et al., 2016), have highlighted its nature as a unipolar electromagnetic pulse of temporal duration comparable to the target discharge time characterized in the earlier experiments (Quinn et al., 2009c). This characterization, carried out with bespoke targets where the length of wire within the probe field of view is maximized, has been at the basis of methods for TNSA proton-beam conditioning in suitably designed helical coil targets, as presented by Kar et al. (2016).

Electron energization at the irradiated target surface is also at the basis of the TNSA mechanism for proton acceleration, which was discussed in Sec. II.A.1. Experiments by Romagnani et al. (2005) employed proton backlighting (using a TNSA probe from a separate foil) to detect the electric fields associated with TNSA acceleration from the rear surface of a laser-irradiated target at $I \sim 10^{19}$ W cm$^{-2}$. Careful temporal synchronization, as well as exploitation of the multiframe capability of RCF stack detectors, led to the detection of the highly transient, Gaussian-shaped TNSA sheath field. The data were used to benchmark TNSA expansion models, which were in substantial agreement with the experimental results.

Romagnani et al. (2019) also reported the characterization of the field associated with the propagation of relativistic electrons in the interior of a target. This experiment, which employed a petawatt-driven high-energy proton probe, exploited the capability of high-energy protons to penetrate through dense matter, while at the same time using a target design aimed to minimize collisional scattering and the
corresponding spatial resolution degradation. The data provided evidence of magnetized filamentation of the electron current within the target, which, via a comparison with hybrid simulations, was attributed to resistive processes. The angular opening of the electron beam injected in the target is another quantity that could be directly inferred from the data.

Complex plasma dynamics are also initiated when an intense laser pulse propagates through an underdense plasma. Several experiments were carried out to investigate this scenario, with the aim of measuring the electric and magnetic field generated as a plasma channel is formed and, additionally, of obtaining information on the propagation and channel features in near-critical plasma where optical probing may be difficult. Experiments by Kar et al. (2007) investigated the formation of a charge-displacement channel in near-critical plasma following the propagation of an intense 100 TW, picosecond pulse through a gas jet. A moving evacuated region was observed in the proton images, in coincidence with the position of the laser pulse, which was consistent with a radial space-charge field within the channel setup by electron displacement by the pulse’s ponderomotive force. The walls of a channel expanding from the interaction regions are seen to develop after the pulse has passed, together with the appearance of a region of proton accumulation along the propagation axis, which was later interpreted (Romagnani et al., 2010) as the signature of a long-lived azimuthal magnetic field within the channel.

More complex channel structures were observed in experiments by Willingale et al. (2011a, 2013) performed on the OMEGA EP laser, investigating the propagation of 1–8 ps, kilojoule-class laser through an underdense, millimeter-scale preformed plasma plume. The time-resolved formation of an evacuated channel was also observed in this experiment, which highlighted a number of additional features, such as filamentation at the channel’s end, channel wall modulations (tentatively associated to the formation of surface waves), as well as the copious appearance of bubblelike structures within the interaction region. This is a recurring circumstance in these types of experiments (Borghesi et al., 2002a; Romagnani et al., 2010; Sarri et al., 2010b). Through comparison with PIC simulations, these structures have been identified as late-time remnants (electromagnetic postsolitons) of solitary structures (solitons) originating from local trapping of frequency down-shifted laser radiation in cavitated plasma regions (Bulanov et al., 1992; Naumova et al., 2001).

II. HED hydrodynamic instabilities

Creating and studying hydrodynamic instabilities in HED environments [including Rayleigh-Taylor (RT) (Rayleigh, 1882; Taylor, 1931), Richtmyer-Meshkov (RM) (Richtmyer, 1960; Meshkov, 1969), and Kelvin-Helmholtz (KH) instabilities (Thomson, 1880)] is of particular interest in the study of HED phenomena such as core-collapse supernovae (Swisher et al., 2015), accretion disks (Balbus and Hawley, 1991), ICF implosions (Lindl et al., 2004; Sadler, Li, and Haines, 2020), or pulsed power pinches (Harris, 1962). These instabilities are present in many strongly driven plasma systems and can lead to turbulence (Flippo et al., 2016) and mixing of materials that can significantly change the behavior and understanding of experiments and phenomena. The addition of self-generated electromagnetic fields makes these systems especially complicated and not well studied experimentally, as access is often limited to highly penetrating x rays.

Proton imaging, as opposed to x-ray imaging, is uniquely suited to observe the electromagnetic fields in these HED experiments and has been employed with some success to date. These self-generated fields, as well as applied fields, can change the plasma properties and can be crucial to understanding the evolution of instabilities like RT (Srinivasan, Dimonte, and Tang, 2012; Modica, Plewa, and Zhiglo, 2013; Song and Srinivasan, 2020), ablative RT (García-Rubio et al., 2021), RM (Samtaney, 2003; Shen et al., 2019, 2020), and KH (Ryu, Jones, and Frank, 2000; Modestov et al., 2014; Sadler et al., 2022), along with a newly discovered composition instability (Sadler, Li, and Flippo, 2020). This includes the possibility of stabilizing these instabilities or curtailing their growth with external fields (Rosensweig, 1979; Sano, Inoue, and Nishihiara, 2013; Srinivasan and Tang, 2013; Praturi and Girimaji, 2019). Some of the first HED experiments to use proton imaging to study hydrodynamic instabilities were done by Manuel et al. (2012) using a laser-driven ablative RT platform and a DDe He proton source. The results showed that the RT instability can lead to self-generated fields, as predicted. A summary is shown in Fig. 26, where a CH target with sinusoidal perturbations was driven by lasers [Fig. 26(a)] and the observed RT growth caused Biermann-generated magnetic fields [Fig. 26(b)] to grow from 3 to 10 T [Figs. 26(d) and 26(e)]. However, these fields were 10 times too small to affect the hydrodynamics directly. Other experiments showed RT bubble growth using a laser-driven foil [see Figs. 26(f) and 27(g)] with either transverse (Gao et al., 2012) or longitudinal (Gao et al., 2013) proton imaging of CH targets. More recent experiments have attempted to look at the self-generated magnetic fields inside denser HED shock-tube targets, where Coulomb scattering is an issue (Lu et al., 2020) and where the fields can change the heat flow in these targets, thereby changing the instability growth (Sadler et al., 2022).

I. Inertial-Confinement Fusion

The ultimate goal of ICF is ignition and high gain, which requires that a cryogenic deuterium-tritium (DT) spherical capsule be symmetrically imploded to reach sufficiently high temperature and density. Such an implosion results in a small mass of low-density, hot fuel at the center, surrounded by a larger mass of high-density, low temperature fuel (Nuckolls et al., 1972; McCrory et al., 1988; Lindl, 1995; Atzeni and Meyer-ter-Vehn, 2004; Hurricane et al., 2014; Betti and Hurricane, 2016). Shock coalescence ignites the hot spot, and a self-sustaining burn wave subsequently propagates into the main fuel region. The symmetry requirements impose strict constraints for achieving fusion ignition (Nuckolls et al., 1972; McCrory et al., 1988; Lindl, 1995; Atzeni and Meyer-ter-Vehn, 2004; Glener et al., 2010; Li et al., 2010; Hurricane et al., 2014; Betti and Hurricane, 2016). The tolerable drive asymmetry of an implosion, in a time-integrated sense, is less than 1% to 2%, depending on the ignition margin (Lindl, 1995; Atzeni and Meyer-ter-Vehn, 2004; Glener et al., 2010; Li et al., 2010; Hurricane et al., 2014; Betti...
and Hurricane, 2016). Consequently, understanding and controlling implosion dynamics is essential for ensuring success. Proton radiography has been developed as an important method for diagnosing ICF implosions because it is sensitive both to plasma density and to electromagnetic fields. Additionally, there are promising indications that externally applied magnetic fields can improve hydrodynamic conditions in ICF implosions and thereby increase capsule performance (Perkins et al., 2013; Srinivasan and Tang, 2013; Mostert et al., 2014; Strozzi et al., 2015; Walsh et al., 2019; Walsh, Crilly, and Chittenden, 2020; Moody et al., 2022), and proton imaging provides a vital tool for assessing how such imposed fields evolve as the implosion proceeds (Gotchev et al., 2009; Heuer et al., 2012).

**Direct-drive implosions.** In direct-drive ICF, a fuel capsule needs to be compressed through illumination by laser light in order to bring the fuel to high temperature and density conducive to fusion and ignition. Earlier work by Mackinnon et al. (2006) successfully demonstrated the feasibility of imaging implosions with TNSA protons, backlighting plastic CH capsules that were imploded by six 1-μm-wavelength laser beams. These were followed by nuclear observations of implosion dynamics for direct-drive spherical capsules on the OMEGA laser using monoenergetic proton imaging (Li et al., 2006a, 2006b). This work provided new insights into the effects of self-generated electric and magnetic fields (Igumenshchev et al., 2014), determine areal density $\rho R$ by measuring the energy loss of backlighting protons, and sample all the implosion phases from acceleration, through coasting and deceleration, to final stagnation in order to provide a more comprehensive picture of ICF spherical implosions.

Further proton-imaging experiments by Li et al. (2008) revealed the existence of a radial electric field inside the imploding capsules. As shown in Fig. 27, proton images showed both inward and outward directed radial electric fields, suggesting that the radial electric field has reversed direction during an ICF implosion. The magnitude of these electric fields compared well with the field calculated from the pressure gradients predicted by the 1D hydrodynamic code LILAC (Delettrez et al., 1987), indicating that the fields are a consequence of the evolution of the electron pressure gradient. The proton images were also utilized to extract quantitative information about capsule sizes and $\rho R$ values at different times, which indicated that the implosions had an approximately 1D performance, with little impact from hydrodynamic instabilities before deceleration.

Additional experiments by Li et al. indicated that the apparent degradation of capsule performance at later times relative to the 1D simulation could largely be a consequence of fuel-shell mixing (Li et al., 2002) and implosion asymmetry (Li et al., 2004). Proton images from experiments by Seguin et al. (2012) also revealed that electromagnetic fields induced filaments inside the capsule shell. These field structures were further discussed by Manuel et al. (2013), who interpreted them as heat-flux-type instabilities driven by the heat conduction from the corona down to the solid ablation layer.

Proton imaging was used by Chang et al. (2011) on the first ICF direct-drive experiment to employ an external applied
magnetic field. The experiments aimed to increase the peak ion temperature of an ICF capsule by strongly magnetizing the plasma’s constituent electrons, reducing thermal conduction across magnetic-field lines and thereby reducing heat loss from the imploded capsule’s hot spot. An ~30% increase to the neutron yield was indeed observed when the external magnetic field was applied. While the proton-imaging data collected were not of sufficient quality for an unambiguous magnetic-field measurement to be made, the characteristic proton-flux inhomogeneity that was observed was consistent with fields that were both trapped and amplified in the imploding capsule.

**Laser-driven hohlraums.** In the indirect-drive approach to ICF, the capsule implodes in response to a quasilinear x-ray radiation field (hundreds of eV), which is generated by multiple high-power laser pulses irradiating a high-Z enclosure called a hohlraum (Lindl, 1995). The x-ray radiation drives the implosion of a cryogenic DT capsule contained within a low-Z ablator, leading to the achievement of high temperature, high-density, and tremendous plasma pressure in the compressed core, and potentially resulting in hot-spot ignition and a self-sustaining fusion burn wave that subsequently propagates into the main fuel region for high-energy gain. In addition, hohlraum-generated x-ray drives can create extreme plasma conditions and have served as an important platform for studying a wide range of basic and applied high-energy-density physics (Lindl, 1995; Atzeni and Meyer-ter-Vehn, 2004), including laboratory astrophysics, space physics, nuclear physics, and material sciences.

In diagnosing plasma conditions and field structures generated in a laser-irradiated hohlraum, proton imaging plays an important role in providing physics insight into the hohlraum dynamics and x-ray-driven implosions, impacting the ongoing ignition experiments at the NIF. For example, experiments by Li et al. (2009, 2010, 2012) utilizing proton imaging to measure the spatial structure and temporal evolution of plasma blowing off from the hohlraum wall revealed how the fill gas compresses the wall blowoff, inhibits plasma jet formation, and impedes plasma stagnation in the hohlraum interior. These experiments also showed that the magnetic field is rapidly convected by the heat flux via the Nernst effect, which was ~10 times faster than the convection by the plasma fluid from expanded wall blowoff. This results in the inhibition of heat transfer from the gas region in the laser beam paths to the surrounding cold gas and in a subsequent increase in local plasma temperature. The experiments further showed that interpenetration of the two materials (gas and wall) occurs due to the classical Rayleigh-Taylor instability as the lighter, decelerating ionized fill gas pushes against the heavier, expanding gold wall blowoff.

Further experiments by Li et al. (2012) deployed proton imaging to address plasma flow dynamics in hohlraums by providing the first physics picture of the process of hohlraum plasma stagnation. Using a Au hohlraum filled with neopentane gas (C₅H₁₂), the resulting proton-fluence patterns indicated that no high-density plasma jets were formed, that the fill gas along the laser beam path is fully ionized, and that the interfaces between the gas plasma and the Au wall blowoff are constrained near the wall surface. The proton images also revealed a unique five-prong, asterisklike pattern that is a consequence of the OMEGA laser beam distribution. Spherical CH targets driven in both gas-filled Au hohlraums and CH-lined vacuum Au hohlraums were used to explore this mechanism further, as shown in Fig. 28. The results show that protons were focused onto the gaps (high-fluence spokes) for the gas-filled hohlraum [Fig. 28(a)] but were deflected away from the spokes in the CH-lined vacuum hohlraum [Fig. 28(b)]. By symmetry these deflections were not due to spontaneously generated magnetic fields; instead, lateral electric fields associated with azimuthally oriented electron pressure gradients (\(\nabla P_e\)) in the plasma plumes and in the radial plasma jets (\(\mathbf{E} = -\nabla P_e/e_n\)) may have been the source
FIG. 28. End-on D³He proton images show surpluses in the regions between the pairs of expanding plasma plumes in (a) a gas-filled Au hohlraum but show deficits in (b) a CH-lined, vacuum Au hohlraum. They indicate opposing directions of the self-generated electric fields, as illustrated schematically by the corresponding sketches. Adapted from Li et al., 2012.

of these deflections. Alternatively, the electric field associated with a supersonic heat front generated by the laser-heated gas channels that are in proximity to the capsule might have caused the deflections.

V. FRONTIERS

While the past two decades have seen the invention and widespread adoption of proton imaging as a diagnostic tool for HED plasma experiments, there remain several key challenges that need to be overcome to extend proton imaging into new experimental frontiers. One challenge is how to probe and analyze more complicated electromagnetic fields, including how to measure fields that strongly deflect protons or create caustics and how to disentangle electric and magnetic fields. A related challenge involves how to combine proton images to extract more information, including in three dimensions. Another challenge is how to probe higher-density plasmas, where scattering is a significant issue. A final grand challenge is how to operate a proton-imaging diagnostic in a high-repetition-rate regime. These challenges can be addressed by advancing the field in four key areas: sources (Sec. V.A), detectors (Sec. V.B), algorithms (Sec. V.C), and schemes (Sec. V.D).

A. Advanced sources

Extending the proton-imaging capabilities demonstrated thus far to plasmas with stronger fields or higher densities will require an enhancement of the energy of the protons beyond what is currently available. As discussed, existing TNSA proton sources extend up to 85 MeV with picosecond pulses (Wagner et al., 2016) and to about 60 MeV with tens of picosecond pulses (Ziegler et al., 2021), although the energies that can be used as an efficient probe for imaging will typically be lower than these cutoffs (as the beam component at the higher energy will have a small divergence and a low number of particles). An obvious route to increasing the proton energies is to use higher power and energy laser drivers.

There are a number of multipetawatt systems currently being developed or commissioned, with some aiming to deliver up to 10 PW of power (Danson et al., 2015). Most of these systems will be based on ultrashort pulse technology (tens of femtosecond pulses), but there are also developments with hundreds of femtosecond duration [such as the 10 PW ATON laser at ELI Beamlines (Jourdain et al., 2021)].

The progress achievable in terms of proton energies will depend on the applicability of scaling laws, as well as secondary factors such as how well the beams can be focused and the extent of pulse contrast that can be obtained on these systems. Various attempts have been made to develop reliable scaling laws for proton energies (Fuchs et al., 2006; Passoni, Bertagna, and Zani, 2010), which typically indicate dependencies on laser intensity or laser energy; see Fig. 3. For example, faster scalings than the ponderomotive scaling predicted by earlier TNSA theories (Wilks et al., 2001) have been reported, within given intensity ranges, with ultrashort (tens of femtoseconds) (Zeil et al., 2010), multipicosecond (Simpson et al., 2021), or multikilojoule laser pulses (Flippo et al., 2007; Mariscal et al., 2019). There is an expectation that experimental results in the multipetawatt regime, providing validation to these scaling predictions, will become available soon as some of the new facilities ramp up their operations.

Additionally, there is a variety of approaches that aim to increase TNSA proton energies by enhancing the energy coupled into relativistic electrons, and by increasing their number density and/or energy; see Macchi, Borghesi, and Passoni (2013) for a review. These approaches are based mostly on target engineering and can involve, for example, reducing the mass (Buffechoux et al., 2010) or density of the target, structuring the target surfaces (Margarone et al., 2012), or adding controlled preplasmas (McKenna et al., 2008). The electrons can also be enhanced through additional mechanisms such as direct laser light pressure acceleration (Klugue et al., 2010) or acceleration by surface waves (Ceccotti et al., 2013; Shen, Pukhov, and Qiao, 2021). Many of these approaches have provided evidence of some proton energy increase from flat foil comparators on a proof-of-principle basis and under specific experimental conditions. Although some of these schemes may have a role to play in the development and optimization of future proton-imaging sources, complications and constraints associated with their implementation may limit their applicability and usefulness.

Beyond TNSA, there are a number of alternate mechanisms under investigation that aim to increase the acceleration efficiency and the accelerated ion energy, or at accelerating ion species other than protons. These include radiation pressure acceleration (RPA) (Esirkepov et al., 2004; Robinson et al., 2008) [in the hole boring (Robinson et al., 2012) and light sail (Macchi, Veghini, and Pegoraro, 2009) implementations], shock acceleration (Fiuza et al., 2012), and schemes taking place in relativistic induced transparency (RIT) regimes (Henig et al., 2009; Poole et al., 2018) such as the break-out afterburner approach (Yin et al., 2007). Hybrid regimes involving a combination of these processes...
have been highlighted in experiments and have recently led to record proton energies approaching 100 MeV through a combination of RPA, TNSA, and RIT acceleration (Higginson et al., 2021).

Although these processes are promising in terms of energy enhancement, for instance, in view of potential medical use, they typically generate beams that do not possess the laminarity and homogeneity of TNSA beams, and their potential usefulness for proton imaging is therefore unclear at present, at least in the backlighting implementation discussed in this review. Nevertheless, if very-high-energy beams will be produced through any of these processes, there may be prospects for different imaging and deflectometry approaches in which a small portion of a beam is spatially or angularly selected (such as with an aperture) and used to sample a finite region within a target or plasma; see also Sec. II.C.2.

Spatial structuring of the beam may also have a role to play in future proton-imaging sources. For example, methods have been put forward for generating multiple separate sources from a single TNSA beam (Zhai et al., 2019), which may lead to additional backlighting capabilities. Similarly, techniques for producing collimated, quasimonoenergetic beamlets (Kar et al., 2016) may lead to opportunities for “spot scanning” of an extended field distribution, i.e., sampling portions of the field region in separate, consecutive shots and tracking the beamlet’s deflection for each shot.

Increasingly, state-of-the-art ultrashort high-power laser systems provide the opportunity to operate at high repetition rates (1–10 Hz) (Danson et al., 2015). This could enable the generation of secondary particle sources at a commensurately high repetition (rep) rate, provided suitable targets can be used, which offer a refreshed surface for irradiation by consecutive pulses. This may be exploited in future proton-imaging experiments, where a repeated proton pulse is coupled to a high repetition interaction pulse, leading to increased data throughput, higher statistics, and/or shorter experiments. A number of targetry solutions suitable for high-rep-rate operations are currently being developed and tested. These include tape targets, where a continuously moving foil tape allows mechanical refreshment of the laser-impacted surface between shots (Noaman-ul Haq et al., 2017; Dover et al., 2020), free-flowing water sheets (Puyuelo Valdes et al., 2022), and cryogenic hydrogen targets (Obst et al., 2017; Chagovets et al., 2022). The aforementioned approaches can provide planar targets with thicknesses in the micron to tens of micron range, leading to beams with the expected TNSA properties, with beam production demonstrated thus far at repetition rates of up to 1 Hz employing petawatt-class laser pulses. Water targets have also been shown to be capable of sustained kilohertz operation at much lower laser energies (Morrison et al., 2018). Another approach explored is the in situ formation of liquid crystal foils (Poole et al., 2016, 2018), which can operate at a more moderate repetition of 0.1 Hz and which provides the capability of varying thickness on demand (from 10 nm to 50 μm). An issue with any high-rep-rate target is the production of debris in the interaction chamber, potentially leading to optics degradation, which may make solutions based on low-density or thinner targets more attractive.

In addition to protons and other heavier ions, electrons are another possible source for imaging. Imaging and deflectometry applications have been reported with electron beams from laser-driven photocathodes (Centurion et al., 2008) and linear accelerators (Zhang et al., 2020), as well as from laser-driven wakefield accelerators. Compared to proton sources, electron sources have a number of attractive features: it is generally easier to generate high-energy (>GeV) electrons that can probe larger field strengths; they can reach shorter temporal durations (for instance, femtosecond scale for wakefield accelerators), they are generally easier to operate at high repetition rates, and they are easier to incorporate into the broad range of existing detector designs that are sensitive to electrons. As an example, recent experiments by Zhang et al. (2020) studied the electron-Weibel instability generated in a low-density gas-jet plasma using a circularly polarized laser via electron imaging with bunches of 45 MeV electrons from a linear accelerator. Efforts to image relativistic plasmas with wakefield-accelerated electron bunches were reported by Schumaker et al. (2013), Zhang et al. (2016), and Wan et al. (2022, 2023). Schumaker et al. (2013) were the first to demonstrate plasma probing applications using wakefield-accelerated electrons, which enabled the detection of highly transient magnetic fields generated by an ultrashort, intense laser pulse on a solid target. More recently Wan et al. (2022, 2023) used very-high-energy (500 MeV) electrons to demonstrate the capability of directly imaging the fields associated with laser-driven plasma wakefields. The experiments highlight a femtosecond temporal resolution, which is well beyond what can currently be achieved with proton beams. It is also straightforward to apply the theory of proton imaging to electrons, including particle-tracing algorithms, although the theory would need to account for relativistic effects given the larger electron speeds involved. A potential disadvantage is that at a given energy electrons also scatter more easily than protons.

B. Advanced detectors

Currently in HED proton imaging, CR-39 (see Sec. II.B.2) and RCF (see Sec. II.B.1) detectors are the most commonly used and are highly efficient. However, they require extensive chemical processing (in the case of CR-39) or a high proton fluence (in the case of RCF), and both require manually labor-intensive digitization efforts. In contrast, besides HED various other proton detectors have been utilized (Bolton et al., 2014; Poludniowski, Allinson, and Evans, 2015). Consequently, more advanced detector systems are currently being developed, especially those that can be used with high-rep-rate laser systems (of the order of 10 shots/h or higher). Here we consider a few systems for the electronic imaging of protons.

Many position sensitive electronic and solid-state detectors have been developed over the years, mainly to support hadron therapy (Poludniowski, Allinson, and Evans, 2015; Johnson, 2018). These include diodes (Wang et al., 2016; Briz et al., 2022), multiwire gas-proportional counters (Sauli, 2014), and cadmium zinc telluride detectors (Simos et al., 2009). However, these detectors suffer from lower spatial resolution and lower sensitivity than CR-39 and RCF, and conversely
The concept of an electronic scintillator-based proton detector system is sketched in Fig. 29. The proton beam goes through the imaged plasma and is recorded by a radiator.

The role of the radiator is to convert the proton flux into an optical signal. A careful shield or collimator design is needed so that the radiator does not produce an unwanted signal from the spurious interactions with the background particles. Converted light can be relayed via the optics or the fiber bundle outside the vacuum chamber, where advanced electronics are often situated. The light is then coupled to a CCD camera or CMOS detector, and MCPs can be added to provide gating functionality to further eliminate background signal.

To date, scintillators have been utilized mainly to understand proton-beam profiles in terms of their spatial and energy characteristics. Numerous scintillators have been tested for their optical emission spectrum using monoenergetic proton beams, as well as their response to proton energies. Experimental measurements indicate that scintillators have a nonlinear scaling with proton energy but a linear response to incident flux (Green et al., 2011). These scintillators were also utilized to characterize the TNSA proton-beam profile of different energy bands generated by different stopping material thicknesses (Green et al., 2011; Bolton et al., 2014; Metzkes et al., 2016; Dover et al., 2017; Huault et al., 2019).

Spatial resolution of scintillators has been characterized for TNSA proton beams using Eljen Technology organic scintillators (Manuel et al., 2020; Tang et al., 2020). A few different scintillator thicknesses were tested with an effective proton source size of 10 μm. The spatial resolution measurement was performed by measuring the point spread function and the contrast of a grid pattern that was placed between the proton source and the detector. The scintillator performance was simultaneously compared with imaging using an RCF detector. The effective resolution limit for the scintillator was measured at ∼22 μm, compared to ∼12 μm for the RCF (Tang et al., 2020). The spatial resolution was only weakly dependent on scintillator thicknesses between 50 and 500 μm for > 2 MeV protons, while thicker scintillators showed improved imaging contrasts. To further improve the spatial resolution, a pixelated scintillator can be used. Pixelation is achieved by laser cutting a grid pattern into the scintillator to optically isolate regions. Experiments with TNSA protons measured the performance of different grid patterns and demonstrated a 20% improvement in spatial resolution (Manuel et al., 2020). The current spatial resolution of the scintillators may be sufficient for imaging, but further improvement is required for the spectrum measurements.

When using scintillator-based detectors, the main limitation in terms of temporal resolution is determined by the decay time of the scintillator, or by any temporal gating applied to the detector used to read out the scintillator output [similar considerations apply to MCP detectors (Sokollik et al., 2009)]. If δtG is the detector’s temporal gating, the temporal resolution at the interaction plane will be given by δt = δtG/ΔM. With a suitable choice of parameters, matching the ~100 ps intrinsic resolution of a D3He source is possible. For example, Sokollik et al. (2009), with an unusually large magnification of M = 70 (which would not be suitable for most experiments) and a time gating of δtG ~ 1 ns [possible in principle if using state-of-the-art scintillators and/or gating (Hu et al., 2018; Zhang et al., 2018; Shevelev et al., 2022)] and a relatively large but more typical magnification M (for instance, M ~ 20), it is challenging to reach temporal resolutions comparable to the picosecond capabilities of RCF and TNSA sources; see Sec. II.C.4.

A volumetric prototype imaging system using a liquid scintillator was developed for hadron therapy dose applications (Darne et al., 2019) but could be adapted and modified for proton imaging. Imaging from several directions would give a tomographic view of the volume and would have information to reconstruct a detailed and possibly 3D proton image with more information than is available with current 2D detectors. Such a system has the possible advantage of capturing all the beam energies efficiently and reconstructing the object in a single shot. It could also be time resolved with high-frame-rate CMOS chips, framing cameras, or streak cameras (depending on the length of emission), and the liquid scintillator has the advantage of being able to be continually replaced in a high-rep-rate system.

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FIG. 29. Sketch of an advanced proton-imaging detector.

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4See http://scintillator.lbl.gov/ for an extensive list.
C. Advanced algorithms and analysis

While there have been various developments in the approaches used to analyze proton-imaging data in recent years (for example, using analytically derived field-reconstruction algorithms to perform inverse analysis; see Sec. III.C.4), there remains room for further advances in several different areas. We do not provide an exhaustive list of them here, but instead highlight a few particularly notable areas.

One of these areas is successful differentiation between proton-imaging measurements of electric and magnetic fields. It is not possible to identify whether electric or magnetic fields are responsible for proton-fluence inhomogeneities seen in a single proton image without further assumptions based on either considerations of geometry or the physics of the imaged plasmas. These assumptions are often well founded (Li et al., 2010; Kugland, Ryutov et al., 2012; Huntington et al., 2015; Schaeffer et al., 2019), but in situations where they are not appropriate, establishing alternative approaches would be of great value. One promising possibility involves the simultaneous analysis of two or more images using protons with different initial energies. In situations where the imaged electromagnetic fields evolve over much longer timescales that the transit delays between protons of all energies, it is reasonable to consider that all proton images collected are of approximately the same electromagnetic field. The distinct energy scalings of deflection angles due to magnetic and electric fields then allows for the degeneracy between the two fields that is characteristic of a single proton image to be overcome. Du et al. (2021) recently presented an algorithm that implements this schema and also proposed a method for minimizing inaccuracies introduced by the temporal evolution of electromagnetic fields. Because both the TNSA and D³He proton sources characteristically produce protons with differentiated energies (see Sec. II.A), which can be detected independently (see Sec. II.B), this approach should be straightforward to include in future analyses (although more work needs to be done to quantify inherent uncertainties).

A simultaneous analysis of multiple proton images with different initial energies also provides a route toward making progress on another outstanding issue: analyzing images in which there are caustics. As discussed in Sec. III.C.4, it is not possible to uniquely reconstruct path-integrated electromagnetic fields from a single proton image if there are caustics present. However, for slowly evolving electromagnetic fields (in the sense just described), multiple-energy proton images can be used to provide more restrictive constraints on the possible solution space of path-integrated fields. Levesque and Beesley (2021) recently proposed a differential evolution algorithm that realized this idea for proton images (two different energies) of simple magnetic fields that either were quasi one dimensional or possessed spherical symmetry. The algorithm was then successfully used to reconstruct the magnetic fields associated with a bow shock in a laser-produced plasma on its collision with a magnetized obstacle (Levesque et al., 2022). At present it is too computationally expensive to apply this differential evolution algorithm to more general (electro)magnetic fields, although it seems likely that other algorithms may be able to address this limitation.

That being said, the number of proton energies required to reduce the possible solution space to a “unique” solution for more complicated path-integrated electromagnetic fields remains uncertain.

The successful use of a differential evolution algorithm to overcome the issues posed by caustics is emblematic of the promise of the application of machine-learning algorithms for enhanced analyses of proton images. Chen et al. (2017) provided a proof of concept in this regard by training an artificial neural network to successfully reconstruct the three-dimensional structure of a simple magnetic field from proton images. It is unclear how well this particular approach would generalize to more complicated electromagnetic fields, but convincing arguments have been made claiming that, once they have been trained, similar neural networks are more computationally efficient than classical analysis algorithms. Similar neural networks could be used to provide improved path-integrated field-reconstruction algorithms by more easily accounting for known limitations in current analysis procedures (for example, uncertainties in the initial beam profiles), as well as quantifying uncertainties. They could also significantly reduce the time taken to perform field reconstructions. The current generation of algorithms typically run for a few hours on a standard laptop computer (Bott et al., 2017), which is long enough to make repeated reconstruction analyzes impractical. Finally, image-recognition-focused machine-learning algorithms could enable more systematic comparisons between synthetic images and measured ones, in turn driving improved standards in the accuracy of field reconstructions.

One final area in which the analysis of proton images could be further developed is the systematic inclusion of scattering models into electromagnetic-field-reconstruction algorithms. As discussed in Sec. II.C.3, such scattering cannot be neglected in many HED experiments when current-generation proton sources are used. While the effect of Coulomb scattering could be minimized using next-generation beams with higher energies (see Sec. V.A), such beams also experience smaller deflections due to electromagnetic fields, which could reduce the feasibility of a successful measurement of the latter. High-quality models of scattering are now commonly incorporated into particle-tracing codes in order to help interpret proton images of electromagnetic fields in high-density laser plasmas; see Sec. III.B.2. They have also previously been used to extract information about plasma densities in ICF experiments directly from proton-imaging data (Mackinnon et al., 2004), but simultaneous inverse analysis of electromagnetic fields has not typically been done. Some recent research suggests that this could have been an oversight. Lu et al. (2020) provided a proof of concept for measuring magnetic fields in high-density (>1 g/cm³) plasmas by utilizing small aperture proton beams in higher-density objects like a laser-driven shock tube, while Bott et al. (2021a) used the broadening of caustic features by scattering in multiple proton images to simultaneously measure magnetic fields and areal densities. If such approaches were to be refined and extended, proton imaging could become a powerful diagnostic on a wider set of experiments than is currently the case.
D. Advanced schemes

In addition to advanced sources and detectors, we discuss here possible new schemes for setting up proton-imaging experiments. These schemes may allow for some of the most significant current limitations of proton-imaging setups to be overcome: specifically, providing characterization of the undisturbed proton fluence, fiducial images for proton deflectometry registration, novel imaging setups utilizing proton optics, and tomographic imaging to obtain 3D measurements.

Characterizing the undisturbed proton fluence is critical for implementing numerical reconstruction techniques, but this is often hampered by shot-to-shot variations and nonuniform initial proton distributions. A straightforward way to address this by extending existing setups is to place one detector (for example, a piece of RCF or a thin scintillator) directly in front of an imaged object and another detector behind it. This would allow one to record an image of both the undisturbed proton fluence and the proton deflections on each shot. Similar work has been pursued in hadron therapy (Johnson, 2018).

Another key opportunity is to combine information from both x-ray and proton images (Orimo et al., 2007; Johnson et al., 2022). In this technique, the last piece in an RCF or CR-39 detector stack is an image plate (see Fig. 9), which can record high-energy x rays from the D\(^{3}\)He implosion or TNSA target. This can be used as an alignment or registration fiducial for proton deflectometry (Johnson et al., 2022). Using lower-intensity lasers (Orimo et al., 2007), one can image smaller or less dense objects, and a variant of this scheme uses a thin needle to produce protons and x rays (Ostermayr et al., 2020) along the same line of sight. This dual imaging has been shown with electrons as well (Nishiuichi et al., 2015; Faenov, Pikuz, and Kodama, 2016), which could be used to break the degeneracy between electric and magnetic fields. In this vein, the use of an x-ray free electron laser or coherent source collocated at a laser facility could provide significant advantages when proton and high resolution x-ray images are combined along the same line of sight.

The development of alternative imaging schemes based on scattering and/or diffraction could be useful for enhanced data acquisition. Possibilities include x-ray and electron analogs for Fourier plane imaging (Smalyuk et al., 2001), coded apertures (Ignatyev et al., 2011), and scattering using a proton microscope for dark field imaging [as done with electrons (Martin et al., 2012; Klein et al., 2015)]. Many of these applications would require the development of compact permanent magnet proton optics (Schollmeier et al., 2014) or dynamic laser-driven optics (Toncian et al., 2006). These would also be useful for generally improving imaging capabilities using charged-particle optics [i.e., a proton microscope much smaller than, but similar to, those at FAIR (Mottershead et al., 2003), LANL (Merrill et al., 2009; Prall et al., 2016; Zellner et al., 2021), and PRIOR (Varentsov et al., 2016)] to image the object, in contrast to the simple point-projection imaging currently employed. Small permanent optics have already been used for energy selection (Schollmeier et al., 2014) as well as pulse solenoid optics (Bruck et al., 2020), where a particular energy can be selected, but we are not aware of their use for the implementation of a proton microscope.

Another advanced imaging scheme that, if realized successfully, would lead to much more detailed measurements of electromagnetic fields is tomographic imaging. At present the main limiting factor on tomography is the simultaneous production (and then detection) of multiple high-energy proton beams. The number of high-energy laser facilities equipped with multiple high-intensity laser beams suitable for proton acceleration is currently small (Danson et al., 2015), while fielding more than two D\(^{3}\)He capsules at once is not feasible even at the largest facilities, due to the number of beams required per capsule. Novel targets have been proposed for overcoming this issue (Spiers et al., 2021), which in principle would allow a single short-pulse beam to produce multiple proton beams. Future experiments will be needed to confirm whether such a scheme would work in practice. Even if a few images of the same electromagnetic structure were obtained successfully, such a sparse number of lines of sight falls well short of a standard tomographic imaging setup. This suggests that specific work on sparse-angle tomography algorithms would be warranted. Two possible approaches that have been demonstrated to address this problem have improved the performance of more conventional filtered back-projection schemes (Spiers et al., 2021). Further improvements could be derived by considering studies in related areas, such as the holographic reconstruction via scattering that is done with electrons (Mankos, Scheinfein, and Cowley, 1996).

VI. SUMMARY AND OUTLOOK

In this review, we examined the use of proton imaging as a diagnostic for electric and magnetic fields in HED laser plasmas. In Sec. II, we described the experimental techniques that underpin proton imaging, including the two primary sources of multi-MeV protons (high-intensity laser-driven sources and D\(^{3}\)He-fusion-capule sources) and the two primary approaches for detecting them (radiochromic film and CR-39 nuclear-track detectors). The characteristic geometry of proton-imaging setups and other important considerations for successful imaging in laser-plasma experiments were also outlined. The theory of how proton images are analyzed in order to extract information about the electric or magnetic fields present in such experiments was reviewed in Sec. III. We explained how a basic physical description of the interaction between charged particles and arbitrary electromagnetic fields allows for both numerical simulations of synthetic proton images of prespecified fields and inverse-analysis techniques (using numerical and/or analytical modeling) that allow for the unique characterization of electromagnetic fields in some (though not all) situations. Section IV presented a broad overview of experiments that have successfully used proton imaging to elucidate many different physical processes of interest in HED plasmas.

While the efficacy of proton imaging as a diagnostic of electromagnetic fields in some HED plasmas is already beyond doubt, there are various different avenues for extending the capabilities of the diagnostic further, which we outlined in Sec. V. One of the primary drivers of these improvements is an ongoing technological progression in
high-power laser technology, as well as an improved understanding of the interaction of these lasers with matter. Taken together these advances have led to a number of promising new proton sources, including some with higher characteristic energies (which can be used to characterize stronger electromagnetic fields than conventional ones) and others with better controlled beam qualities, and the paradigm-shifting prospect of high-repetition-rate sources. This prospect, in particular, has generated much interest in researching new detector technologies that (unlike existing ones) can reliably output sufficiently resolved proton images at equivalent rates to the repetition rate of the laser driving the source. Concurrently there has been renewed effort in the last five years toward developing new techniques for extracting information from proton-imaging data systematically and automatically. Given the recent rate of progress in this area and broader scientific advances in data analysis derived from machine learning, it is not unreasonable to anticipate that there will exist within ten years a plethora of new, sophisticated algorithms that go beyond anything we have described here. Finally, although moving to imaging schemes that are more advanced than the current standard (such as tomographic schemes and schemes attempting “proton optics”) presents several serious practical challenges, the latest research suggests that progress toward realizing such schemes is not an impossible dream.

In short, during the just over two decades since it became practically realizable, proton imaging has proven to be a powerful approach for measuring two of the key physical fields that characterize HED plasmas. Among its many successful applications, it has been used to show magnetic-field generation in both direct-drive and indirect-drive ICF experiments, with significant ramifications for heat transport; it has helped probe the mechanism for kinetic processes in collisionless or weakly collisional plasmas; and it has played a key role in numerous laboratory astrophysics experiments. Looking forward, the ongoing development of high-intensity-laser and fusion-capsule-backlighter proton sources on the highest-energy lasers in the world suggests that proton imaging will continue to be used to probe electric and magnetic fields in the most intriguing new HED plasma experiments for the next decade and beyond.

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