Instability-driven electromagnetic fields in coronal plasmas

M. J.-E. Manuel,1,b,c) C. K. Li,1 F. H. Séguin,1 N. Sinenian,1 J. A. Frenje,1 D. T. Casey,1 R. D. Petrasso,1 J. D. Hager,2,d) R. Betti,2,e) S. X. Hu,2 J. Delettrez,2 and D. D. Meyerhofer2,e)
1Plasma Science and Fusion Center, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA
2Laboratory for Laser Energetics, University of Rochester, Rochester, New York, 14623, USA

(Received 16 November 2012; accepted 15 January 2013; published online 15 April 2013)

Filamentary electromagnetic fields previously observed in the coronae of laser-driven spherical targets [F. H. Séguin et al., Phys. Plasma. 19, 012701 (2012)] have been further investigated in laser-irradiated plastic foils. Face-on proton-radiography provides an axial view of these filaments and shows coherent cellular structure regardless of initial foil-surface conditions. The observed cellular fields are shown to have an approximately constant scale size of ~210 μm throughout the plasma evolution. A discussion of possible field-generation mechanisms is provided and it is demonstrated that the likely source of the cellular field structure is the magnetothermal instability. Using predicted temperature and density profiles, the fastest growing modes of this instability were found to be slowly varying in time and consistent with the observed cellular size. © 2013 AIP Publishing LLC. [http://dx.doi.org/10.1063/1.4801515]

I. INTRODUCTION

Spontaneous electromagnetic fields can be important to the dynamic evolution of a plasma by directing heat flow1 as well as providing additional pressures on the conducting fluids through the Lorentz force. Electromagnetic fields are predicted to affect fluid behavior during the core-collapse of supernovae2 through generation of fields due to hydrodynamic instabilities. In the coronae of stars, self-generated magnetic fields lead to filamentary structure3 in the hot plasma. Recent experiments by Gregori et al.4 investigated sources of protogalactic magnetic fields generated by laser-produced shock waves. In inertial confinement fusion (ICF)5 experiments, self-generated electromagnetic fields can also play a role. Present day laser facilities provide a unique opportunity to study spontaneous field-generation in these extreme environments under controlled conditions.

Self-generated magnetic fields in laser-produced plasmas were first measured by Stamper et al.6 and were shown to originate from the so-called Biermann battery7 source caused by misaligned temperature and density gradients. The magnetic field evolution equation is found by substituting the electron momentum equation into Faraday’s law and can be expressed as

$$\frac{\partial \mathbf{B}}{\partial t} \approx \nabla T_e \times \nabla n_e + \nabla \times (\mathbf{V}_{\text{adv}} \times \mathbf{B}) - \nabla \times (D_m \nabla \times \mathbf{B}),$$

(1)

where $T_e$ is the electron temperature, $n_e$ is the electron density, $e$ is the unit charge, and $D_m$ is the diffusion coefficient. In this formalism, electron inertia and second order terms in $\mathbf{B}$ have been neglected. The advection velocity $\mathbf{V}_{\text{adv}}$ is the vector sum of the fluid velocity $\mathbf{V}_{\text{fluid}}$ and the so-called Nernst velocity $\mathbf{V}_{\text{Nernst}}$. Advection by the Nernst effect arises because the magnetic field can move with the heat-conducting electron population and is thus proportional to the temperature gradient

$$\mathbf{V}_{\text{Nernst}} \approx -\frac{\mu_0}{\delta_0} \frac{\tau_e}{m_e} \nabla T_e,$$

(2)

where $m_e$ is the electron mass, $\tau_e$ is the electron-ion collision time, and $\mu_0$ and $\delta_0$ are Braginskii coefficients. The weakly magnetized approximation (Hall parameter $\chi \ll 1$) has been implemented and for the plastic (CH) plasmas discussed herein $\mu_0/\delta_0 = 2.5$. The Nernst effect contributes to the total convection of the magnetic field along with field diffusion described by the third term in Eq. (1). The first term in Eq. (1), the so-called Biermann battery or thermoelectric source, is the dominant mechanism by which magnetic fields are generated in plasmas.

Much work has been done in diagnosing spontaneous electromagnetic fields in myriad laser-produced-plasma configurations: irradiated wires,11 directly driven glass and plastic capsules,10,12–14 inside of hohlraums,15 and plasma bubble evolution.16,17 The work described here extends previous observations10 of filamentary field structures found in the corona of laser-irradiated spherical targets as discussed in Sec. II. The complexities inherent to the 3-D geometry previously investigated are alleviated through simpler face-on imaging of laser-irradiated plastic foils. The experimental methodology using both monoenergetic-protons and x-rays is discussed in Sec. III. Radiographic results of coherent 3-D cellular structure in proton images are analyzed and presented in Sec. IV. This is followed by a discussion of multiple field-generation mechanisms in Sec. V and it is

1070-664X/2013(5)/056301/10/$30.00 20, 056301-1 © 2013 AIP Publishing LLC
demonstrated that the likely source of these fields is the magnetothermal instability (MTI) which occurs in the underdense corona. Plasma conditions in the corona are shown in Sec. VI and details of the MTI in these plasmas are discussed in Sec. VII. This article concludes with a summary of the results presented here and discusses the future directions of this work in Sec. VIII.

II. FILAMENTARY FIELDS IN THE CORONA

Monoenergetic-proton radiography\(^{18,19}\) has been used to image coronal field structures in laser-irradiated CH spheres.\(^{10}\) Figure 1(a) shows the experimental setup used in these experiments. A solid CH sphere with a diameter of \(\sim 860 \mu m\) was irradiated by 0.351 \(\mu m\) laser light with an intensity of \(\sim 2 \times 10^{14} W/cm^2\) in a 1 ns pulse that had full beam smoothing\(^{20,21}\) and SG4 distributed phase plates\(^{22}\) implemented on the OMEGA laser.\(^{23}\) The resultant plasma blowoff was probed using 15 MeV fusion-protons provided by the backlighter capsule, discussed further in Sec. III, and images were recorded on CR-39 nuclear track detectors.\(^{19,24}\)

Figure 1(b) shows four proton-fluence radiographs taken before, during, and after the 1 ns laser pulse. The field of view of the detector is 3 mm \(\times\) 3 mm in each image. Protons are stopped in the solid sphere which results in the white “shadow” prevalent in each radiograph. No field structures were observed until the end (\(\sim 1\) ns) of the laser drive in these images. Séguin et al. demonstrated that the observed filamentary structures are quickly (\(\sim 200\) ps) generated between \(\sim 0.6\) ns and \(\sim 0.8\) ns in these experiments and the amplitude of proton fluence modulations was approximately constant once fields appeared. The filaments expand radially outwards with the blow-off plasma and were still visible after the drive ended.

The magnetic field evolution described by Eq. (1) indicates that fields will be advected in the direction of \(V_{\text{adv}}\). The advection velocity depends on both the fluid and Nernst velocities and changes as a function of position. Plasma evolution in these experiments was modeled using the 1-D radiation-hydrodynamic code LILAC\(^{25}\) and the resultant velocity profiles at 1 ns are shown in Figure 1(c) as a function of distance from the ablation surface. The Nernst velocity was calculated using predicted temperature profiles in Eq. (2) and changes direction at the peak temperature as indicated in the plot. Consequently, the advection velocity changes direction in this plasma \(\sim 100 \mu m\) from the ablation front. B fields generated inside this transition region advect towards the ablation surface, otherwise they expand out with the plasma.

Filamentary structures are observed throughout the corona, thus fields must be generated in a region where they will advect out with the expanding plasma. In these images, the filaments appear to extend to the ablating surface of the sphere; however, the inherent 3-D nature of the spherical geometry precludes any definitive assessment. These images clearly demonstrate that filaments expand with the coronal plasma and are present after the laser drive has ended. The field diffusion time, \(\tau_{\text{diff}} \approx (k^2D_m)^{-1}\) where \(k\) is the wave number, may be estimated for \(\sim 100 \mu m\) size fields and is \(\approx 5\) ns at the end of the laser drive. Thus, it is not surprising that fields are still visible \(\sim 500\) ps after the drive ends. Quantitative analysis of the filament size is difficult in this geometry but was predicted\(^{10}\) to be of order \(\sim 150 \mu m\) at the quarter-critical surface using Monte Carlo simulations; no difference in size was discernible at different sample times from these data. Face-on imaging of laser-irradiated planar foils has been performed to further probe these filamentary

![Figure 1](image-url)
fields on-axis and more accurately quantify their size as described below.

III. PLANAR RADIOGRAPHY EXPERIMENTS

Monoenergetic-proton and x-ray radiography experiments were performed on the OMEGA laser\(^\text{23}\) using the configuration shown in Figure 2. Protons are sensitive to both mass and field modulations through Coulomb scattering and the Lorentz force, respectively. X rays are sensitive only to areal density modulations in the target. The complementarity of these two diagnostic techniques provides information to address density and field distributions during plasma evolution. Unlike the solid sphere experiments, planar foils will be accelerated and Rayleigh-Taylor (RT) growth of density perturbations is expected to occur. Density distributions and growth of perturbations are characterized by x-ray radiographs, whereas protons sample the path-integrated field structures.

The proton backlighter capsules, filled with 18 atm D\(^3\)He gas, were imploded using 20 OMEGA beams to produce fusion protons. Each proton radiography experiment\(^\text{18,19}\) gives a single “snapshot” in time of the laser-foil interaction and multiple experiments with different laser timings provide a series of radiographs illustrating the plasma evolution. This backlighting technique provides a temporal resolution of \(~150\) ps and D\(^3\)He-fusion protons (\(E_p \sim 15\) MeV) are produced by an approximately Gaussian source with a FWHM of \(~45\) \(\mu\)m. 15-MeV proton radiographs were recorded on filtered CR-39 plastic nuclear track detectors. After exposure, CR-39 samples were etched to reveal tracks left by the incident protons. Through the etching process, signal tracks in the plastic were revealed and pieces were scanned using an optical microscope system.\(^\text{26}\)

From these scans, radiographs were processed and proton fluence images were created.

X-ray radiographs were taken using a laser-irradiated Uranium foil to generate \(~1.3\) keV x rays for optimum contrast through \(~20\) \(\mu\)m CH. Images were recorded on film using a framing camera\(^\text{27,28}\) with a temporal resolution of \(~80\) ps and a spatial resolution of \(~10\) \(\mu\)m. The aperture framing camera provided multiple images of a single foil during its evolution, yielding multiple radiographs from a single experiment. The measured optical depth image may be directly converted to an areal density map\(^\text{27}\) of the target for comparison with proton radiographs that are sensitive to both density modulations and field deflections at the foil.

Four types of CH foils of varying thicknesses and surface perturbations were used in these experiments and are listed in Table I. Foil surfaces were either flat, seeded with ridge-like 2-D sinusoidal modulations, or 3-D eggcrate-like sinusoidal modulations. The laser spot was shaped by SG4 distributed phase plates (DPPs),\(^\text{22}\) smoothed by spectral dispersion (SSD),\(^\text{20}\) and distributed polarization rotators (DPRs)\(^\text{31}\) were implemented to minimize the broadband imprint from the laser spot. Twelve beams were overlapped to provide a drive intensity of \(I \sim 4 \times 10^{14}\) W/cm\(^2\) within a \(~750\) \(\mu\)m diameter spot using a 2 ns square pulse to deliver a total of \(~3300\) J of energy on target. Proton and x-ray radiographs were taken at times up to 2.4 ns to provide data during and after the onset of the laser drive.

IV. EXPERIMENTAL RESULTS

A. Proton radiographs

A comprehensive summary of proton-fluence radiographs is shown in Figure 3. Radiographs were taken of the four different foil types after the first nanosecond of the 2 ns laser drive. Proton fluence images of all foil types for times \(~\leq 1.5\) ns indicate minor variance across the analysis region, though coherent linear features were observed in 2-D modulated foils. Previous work\(^\text{29,30}\) has shown that these linear features are generated by Rayleigh-Taylor-induced magnetic fields caused by the growth of the preimposed surface perturbations. A rapid transition (\(~200\) ps) occurs near \(t \sim 1.5\) ns, whereby the fluence radiographs of all foil types show a drastic change in appearance. Some underlying linear features most likely due to RT-induced fields are still observable, especially \(180\) \(\mu\)m-foil images at times \(~\leq 1.5\) ns; however, the 3-D cellular structure is prevalent and dominates proton radiographs. These features are consistent with an axial view of the previously observed filamentary fields.

Cellular structure was shown to begin during the laser drive and continued well after the end of the pulse. The dominant scale size of these features (\(\lambda_{AC}\)) and the rms amplitude

<table>
<thead>
<tr>
<th>Label</th>
<th>(l_0 (\mu m))</th>
<th>(\lambda_0 (\mu m))</th>
<th>(a_0 (\mu m))</th>
</tr>
</thead>
<tbody>
<tr>
<td>Flat foil</td>
<td>21</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td>120 (\mu m) (2-D)</td>
<td>21</td>
<td>120</td>
<td>0.27</td>
</tr>
<tr>
<td>180 (\mu m) (2-D)</td>
<td>23</td>
<td>180</td>
<td>0.55</td>
</tr>
<tr>
<td>115 (\mu m) (3-D)</td>
<td>26</td>
<td>115</td>
<td>0.56</td>
</tr>
</tbody>
</table>

TABLE I. Metrology of the four CH foil types used in these experiments. Initial ambient foil density was \(\rho = 1.03\) g/cm\(^3\). Variations in thickness \((l_0)\), wavelength \((\lambda_0)\), and sinusoidal amplitude \((a_0)\) are all \(~\pm 10\) %. The value given for the eggcrate foil (3-D) is the diagonal peak-to-peak wavelength.
modulation of proton fluence ($\sigma_{\text{rms}}$) were calculated (as described in the Appendix) to characterize the scale and strength of the filamentary field structures, respectively. The results of the autocorrelation analysis are plotted in Figure 4(a). It is clearly demonstrated that these features have an approximately constant scale size that can be characterized by $\lambda_{\text{AC}} \approx 210 \mu m$ with a standard deviation of $\pm 30 \mu m$. Radiographs taken before cellular-onset were analyzed, but did not reveal a dominant scale size and are thus not plotted in Figure 4(a). The scale of filamentary fields is shown to be constant in size immediately after initial onset and it does not measurably change in time or depend on the initial foil surface conditions. This suggests that filamentary fields are generated away from regions affected by the shape of the ablation surface, i.e., in the expanding, underdense corona.

Figure 4(b) illustrates the normalized rms amplitude as a function of time for all four foil types. Radiographs at earlier times ($\leq 1.5 ns$) are shown to have normalized rms amplitudes of $\approx 20\%$, though no dominant scale size was observed. The normalized broadband rms characterizes the amplitude of proton deflections and thus path-integrated field strength. It is clearly shown that proton deflections grow at the same rate for all foil types during the laser drive. The filamentary fields causing the cellular features are created during the laser pulse, but do not lose appreciable strength $\leq 500 \mu s$ after the drive ends.

B. X-ray radiographs

All four foil types were radiographed with x rays to characterize density distributions in these laser-foil interactions. Modulations in areal density arise due to Rayleigh-Taylor growth of surface perturbations and laser imprint. X-ray radiographs shown in Figure 5(a) illustrate the evolution of areal density modulations for all target types. In these images, lighter pixels indicate higher areal density in the target (more x-ray absorption). Flat-foil radiographs show no significant features until late in time ($t \approx 2 ns$), at which point small-scale ($\sim 30 \mu m$) structure due to laser imprint becomes apparent. However, x-ray radiographs of modulated foils clearly show dramatic features at the seeded wavelengths due to RT growth. The perturbation amplitude of the $\lambda = 120-\mu m$-foil increases in time with a rate of $\gamma_{120} \approx 2 ns^{-1}$ during the drive and these data agree well with DRACO radiation-hydrodynamic simulations as shown in Figure 5(b). Some small-scale structure, also at $\sim 30 \mu m$, is visible at late times in modulated foils, but these features are much lower in amplitude than the dominant seeded perturbations.
X-ray radiographs taken at times \( \geq 1.5 \, \text{ns} \) do not show similar features as those observed in late-time proton fluence images. Cellular structure in areal density has been observed\(^{31}\) under different experimental conditions due to laser imprint, but in experiments discussed here the 3-D structure did not have enough time to strongly develop. A Fourier analysis was performed on the x-ray images of flat and 120-\( \mu \text{m} \)-modulated foils at \( \sim 2.2 \, \text{ns} \) to compare the relative amplitudes of the observed features. The resultant spectra are shown in Figure 6 and demonstrate that the \( \sim 30 \, \mu \text{m} \) features in both cases are approximately equal in amplitude and are consistent with RT growth of laser-imprinted perturbations. In the modulated foil case, the amplitude of 120-\( \mu \text{m} \) perturbations is \( \sim 5 \) times higher than \( \sim 30 \, \mu \text{m} \) features.

Proton deflections due to RT-induced B fields are dominated by fields at the seeded wavelength of 120-\( \mu \text{m} \) in the modulated foil case at times \( \leq 1.5 \, \text{ns} \). B fields created by RT growth occur near the ablation surface and, neglecting diffusion, are proportional to the fluid vorticity. The peak field scales\(^{34}\) with perturbation parameters as \( |B(t)| \propto \frac{h(t)}{\gamma} \), where \( h \) is the perturbation amplitude, \( \lambda \) is the wavelength, and \( \gamma \) is the growth rate. Proton deflections, though, are proportional to the path-integrated field strength \( \langle B_L \rangle \) where the field scale length \( L_D \sim h \).\(^{30}\) Measured RT growth rates\(^{29}\) for both wavelengths, \( \gamma_{30} \sim 4.5 \, \text{ns}^{-1} \) and \( \gamma_{120} \sim 2 \, \text{ns}^{-1} \), were used with the fact that \( h \propto \langle p_L \rangle_{\text{rms}} \) to estimate the relative magnitudes of \( \langle B_L \rangle \) between the two perturbation wavelengths. If proton deflections late in time were due only to RT-induced B fields, simple estimates show that deflections due to \( \lambda = 120 \, \mu \text{m} \) are \( \sim 15 \) times higher than those for \( \lambda = 30 \, \mu \text{m} \). This is a lower limit because diffusion effects, not included here, affect shorter wavelengths more than longer ones. Furthermore, this analysis is for images at \( \sim 2.2 \, \text{ns} \), whereas proton images illustrate strong cellular features by \( \sim 1.7 \, \text{ns} \). If RT-induced B fields were the dominant field structure in modulated-foil experiments, proton images would exhibit strong features consistent with surface perturbations. Instead, coherent cellular structures likely caused by filamentary fields in the underdense corona were observed in proton images at late times independent of initial surface conditions.

V. INSTABILITIES IN LASER-PRODUCED PLASMAS

Magnetic fields may be generated by a range of instabilities and sources as described in detail by Haines.\(^{9}\) The primary sources are outlined in Figure 7(a) which shows where they occur in the sample plasma given by DRACO profiles at 1.5 ns. In addition to these instabilities, laser-plasma instabilities (LPI) can locally generate B fields through generation of hot-electron currents in the coronal plasma. Hard x-ray detectors\(^{35}\) and scattered-light streak cameras\(^{36}\) were fielded to observe LPI-related plasma behavior in the corona, though no hard x-ray or relevant scattered-light signals were observed. Filamentation of laser “hot-spots” was not responsible because the average intensity was too low\(^{37}\) and all beams
were spectrally smoothed. LPI-generated hot-electron currents are not the source of these coherent field structures.

The so-called Weibel instability is generated by electron-temperature anisotropy in the plasma. The typical collisionless form (CLW) of this instability is only relevant at the very edge of the corona when $\lambda_{\text{cl}}/e \approx 10^3$, where $e$ is the collisionless skin depth. In these experiments, this parameter is $\approx 500$ and thus CLW is a very unlikely field source. The collisional Weibel (CW) instability, which occurs in the denser plasma regions, has been predicted to generate fields under ablatively driven conditions. Fields created by the CW tend to grow fastest near the over dense region ($n_e \approx n_{\text{crit}}$), but when the Nernst effect and field diffusion were included, the instability was shown to be stabilized. Nonetheless, if these fields exist under the specific experimental conditions here, they would be confined to the overdense region and not advected outwards.

The electrothermal instability (ETI), first described by Haines, occurs between the ablation and critical surfaces. This instability originates from a spatial perturbation to the background electron temperature that has a component perpendicular to the heat flux. The temperature perturbation leads to nonuniform ohmic heating due to the change in plasma resistivity, which positively feeds back on the perturbation. The fastest growing mode ($\lambda_{\text{ETI}} \approx 60 \mu m$) is much smaller than the observed $\approx 210 \mu m$ scale size, therefore this instability is unlikely to be the source of the observed cellular fields. Furthermore, fields generated by the electrothermal instability are also confined to the high density region.

The Nernst effect in these experiments was calculated to be strongest just outside the critical surface. In the sample plasma shown in Figure 7(b), $V_{\text{Nernst}} \approx 700 \mu m/ns$ and is much larger than the fluid velocity. Fields generated here are quickly advected towards the ablation front. This can result in amplification of the B field due to the increasing electron density, though this amplification is reduced by diffusion effects. DRACO calculations presented herein implemented the local thermal equilibrium (LTE) approximation and a constant flux limiter of $f = 0.06$ which has been shown to reproduce drive conditions well at intensities $\approx 5 \times 10^{14} \text{ W/cm}^2$. These calculations indicate that the advection velocity changes directions near the $\frac{1}{2}$-critical surface. Coherent fields observed in planar and spherical experiments must be generated where outward advection can occur.

A seed B field generated by the Biermann battery provides the source for the magnetothermal instability. A perturbation to the background temperature profile can result in a misalignment of the temperature and density gradients, thereby creating a seed field. Anisotropic thermal conduction due to the presence of this B field may act to enhance the perturbation. Classically, this instability can occur when the temperature and density gradients in the direction of heat flow are aligned; the coronal plasma outside of the peak temperature satisfies this condition. In this region the Nernst velocity is aligned with the fluid flow as shown in Figure 7 and B fields generated in this region will advect outwards. Furthermore, the fastest growing MTI modes, using the classical analysis, are found to be $\lambda_{\text{MTI}} \approx 200 – 300 \mu m$ at these times, as discussed further in Sec. VII. MTI-generated B fields are consistent with the observed cellular scale size and are created in the corona where they will advect outwards. Coherent fields observed in proton radiographs are likely due to MTI in the corona.

VI. PLASMA FLOW CONDITIONS

Laser-matter interactions create plasma conditions that vary as a function of position as demonstrated in Figure 7(a). Plasma conditions go from cold ($T_e \approx 200 \text{ eV}$) and dense ($n_e \approx 10^{23} \text{ cm}^{-3}$) near the ablation front to hot ($T_e \approx 2 \text{ keV}$) and sparse ($n_e \approx 10^{21} \text{ cm}^{-3}$) in the corona. The fluid velocity also varies and, at the time shown in Figure 7(b), changes direction from inward flow to outward expansion just outside the critical surface. These varied conditions give rise to differences in dominant physics mechanisms.

The magnetic Reynolds number ($R_{\text{em}}$) relates B-field advection to diffusion in the plasma. This can be expressed...
as $\text{Re}_m = \frac{V_{\text{adv}} L}{D_m}$, where $V_{\text{adv}}$ is the field advection velocity, $L$ is the plasma scale length, and $D_m$ is the field-diffusion coefficient. Figure 8 shows how $\text{Re}_m$ varies in the plasma described in Figure 7(a) and demonstrates that in most locations, field advection dominates over diffusion. This is due to the low resistivity away from the ablation front and is consistent with B fields being frozen-in to the flow.

The Reynolds number ($\text{Re}$) characterizes whether the fluid flow is laminar ($\text{Re} \leq 2300$) or turbulent ($\text{Re} \geq 4000$).\(^\dagger\) This dimensionless number compares inertial to viscous forces and can be written as $\text{Re} = \frac{V_{\text{fluid}} L}{\nu}$, where $V_{\text{fluid}}$ is the fluid velocity, $L$ is the plasma scale length, and $\nu$ is the kinematic viscosity of plasma ions. Figure 8 illustrates how $\text{Re}$ changes as a function of position. The Reynolds number increases near the ablation front due to the decrease in viscosity ($\nu$), though x-ray radiographs demonstrated coherent RT-growth at the ablation front in modulated foils during the entirety of the laser drive. These calculations demonstrate that plasma flow is dominantly laminar in nature allowing for coherent features and that B fields will be strongly advected in the underdense corona.

VII. MTI IN SPHERICAL AND PLANAR GEOMETRIES

The MTI grows from a seed magnetic field in the coronal plasma that can originate from a number of different perturbative sources, laser nonuniformities, plasma waves, etc. Thus, the initial mode distribution is not well understood or characterized and varies spatially and temporally, making quantitative analysis of this instability very difficult. Some insight may be gained, however, by examining the growth rate ($\gamma_{\text{MTI}}$) and wavelength ($\lambda_{\text{MTI}}$) of the fastest growing mode in these CH plasmas, as given by\(^\dagger\)

$$
\gamma_{\text{MTI}} \approx 1.65 \times 10^8 \frac{T_e^{5/2}}{n_e Z L_T^2} \ln A \left( \frac{1}{\text{ns}} \right),
$$

$$
\lambda_{\text{MTI}} \approx 2 \times 10^{-4} \sqrt{\frac{L_n L_T Z \ln A}{n_e \gamma D T_e^{1/2}}} \text{[\mu m]}.
$$

In the preceding equations, $Z$ is the average charge state, $\ln A$ is the Coulomb logarithm, the electron temperature $T_e$ is in keV, electron density $n_e$ is in $10^{20} \text{ cm}^{-3}$, the Debye length $\lambda_D$ is in $\mu m$, and the scale lengths for temperature ($L_T$) and density ($L_n$) are also in $\mu m$. Figures 9(a) and 9(b) illustrate calculations from LILAC and DRACO profiles, respectively, of $\lambda_{\text{MTI}}$ (solid) and $\gamma_{\text{MTI}}$ (short dashed) at single instances near the observed field-onset time. The fastest growing wavelength is very large near the peak temperature due to the long temperature scale lengths, but levels off quickly and slowly increases farther out due to decreasing densities in both the spherical and planar cases. The growth rate peaks at $\sim 5 \text{ ns}^{-1}$ in the planar case and decreases thereafter, whereas $\gamma_{\text{MTI}}$ levels off at $\sim 10 \text{ ns}^{-1}$ in the spherical case before diverging at larger radii due to the rapidly decreasing density. These profiles change as a function of time, but the snap-shots shown here provide the general differences between the planar and spherical cases.

The observed onset time of the coherent field structures in planar experiments was $\sim 2$ times later than the apparent
onset in spherical experiments. The exact cause of the apparently rapid onset of these fields is not well understood. Though near the field-onset time, calculations from classic MTI theory, as shown in Figures 9(a) and 9(b), predict a growth rate $\sim2$ times faster in spherical experiments than planar experiments, consistent with the two separate data sets. A given plasma profile provides a spectrum of “fastest growing modes” throughout the plasma, though in both the spherical and planar cases these modes are predicted to be between $\sim200$ and $300\,\mu m$, again consistent with observations. To assess some aspect of temporal variance, an average wavelength $\langle \lambda_{\text{MTI}} \rangle$ is defined by

$$\langle \lambda_{\text{MTI}} \rangle = \frac{\sum_i \lambda_{\text{MTI},i} \times \gamma_{\text{MTI},i}}{\sum_i \gamma_{\text{MTI},i}},$$

(5)

where each $\lambda_{\text{MTI}}$ is weighted by the associated growth rate and is summed over the MTI-unstable region in the underdense corona at each time step. The results of this calculation for both spherical and planar simulations are shown in Figure 9(c). The laser drives for both configurations are also shown (dotted) for reference and it is clear that $\langle \lambda_{\text{MTI}} \rangle$ rapidly increases after the drive turns off due to flattening of the temperature profiles as the plasma cools. These calculations suggest that MTI occurs earlier in the drive at smaller wavelengths, though this was not observed in experiments. Rather, a rapid transition ($\sim200$ ps) was demonstrated in both the spherical and planar experiments. Interestingly though, the measured characteristic size of cellular structures, $\sim210\,\mu m$, crosses the spherical and planar curves in Figure 9(c) near the observed onset times for each case. It is also important to note that these calculations predict a slowly varying dominant scale size in time, which was not observed. Even though no temporal growth was measured from proton radiographs, the range of values measured from separate planar experiments is consistent with the predicted $\sim200–300\,\mu m$ during these times.

VIII. CONCLUSION

In summary, proton and x-ray radiography experiments have been performed to further investigate filamentary fields previously observed in laser-irradiated spherical targets. Planar experiments provided an axial view of the filamentary fields and show coherent cellular field structures independent of initial surface conditions, which suggests field generation in the corona. Furthermore, these filamentary fields were observed far from the ablation surface so must be generated in a region where they can be advected outwards. These coherent field structures dominate proton radiographs of planar foils at late times over other sources of electromagnetic fields.

It has been demonstrated that the likely cause for these fields is the magnetothermal instability which is predicted to produce field structures at the observed scale size. Moreover, this instability generates fields near the peak coronal temperature where they are convected outwards with the plasma. Using classical MTI theory from Tidman and Shanny, the difference in observed onset time between the two geometries may be explained through a difference in growth rate. MTI-generated fields occur due to altered heat conduction in the corona and could therefore affect heating uniformity and efficiency in directly driven targets. These results have identified the likely source of coronal fields previously observed in directly driven spherical targets. Initial calculations indicate consistency with experiments and thus motivate further numerical and experimental exploration into field-generation by this instability.

ACKNOWLEDGMENTS

The authors thank Wolf Seka and Christian Stoeckl for their help in interpreting results from the scattered-light and hard x-ray diagnostics, respectively. The authors also express their gratitude to the engineering staff at LLE for their support. The work described here was done as part of the first author’s Ph.D. thesis and supported in part by NULF (DE-NA0000877), FSC/UR (415023-G), DoE (DE-FG52-09NA29553), LLE (414090-G), and LLNL (B580243).

APPENDIX: PROTON-RADIOGRAPH ANALYSIS TECHNIQUES

The two observables of interest in proton radiographs are the characteristic spatial wavelength of the features ($\lambda_{\text{AC}}$), if one exists, and the normalized broadband rms amplitude ($\sigma_{\text{rms}}$) which is a measure of the strength of the features. The latter may be calculated from the distribution of fluence values after removing the statistical component,

$$\sigma_{\text{rms}} = \frac{\sqrt{\sigma^2_{\text{meas}} - \sigma^2_{\text{stat}}}}{\mu_{\text{meas}}},$$

(A1)

where $\sigma_{\text{meas}}$ is the measured statistical deviation of protons per pixel, $\mu_{\text{meas}}$ is the statistical mean proton fluence per pixel used to normalize the variation across different experiments, and $\sigma_{\text{stat}}$ is the numeric statistical variation per pixel $\sim\sqrt{\mu_{\text{meas}}}$. Deducing the characteristic spatial length of the cellular features requires a more complex analysis.

An autocorrelation (AC) algorithm is used on each proton-fluence image to determine the isotropic scale-size of the observed features. This procedure begins by calculating the 2-D FFT of a square region that results in a 2-D array $C(k_x,k_y)$ of complex Fourier coefficients,

$$C(k_x,k_y) = \mathcal{F}\{I(x,y)\},$$

(A2)

where $\mathcal{F}$ denotes the FFT algorithm and $I(x,y)$ is the array of pixel values in the area of interest. The autocorrelation coefficients $A(x,y)$ may then be calculated directly from the Fourier coefficients,

$$A(x,y) = \mathcal{F}^{-1}\{C^*C\},$$

(A3)

where $C^*$ is the complex conjugate of the Fourier coefficient array and $\mathcal{F}^{-1}$ is the inverse FFT algorithm. The autocorrelation coefficients, as defined above, represent how well the image correlates with itself. The AC coefficients are
azimuthally averaged in space, thus eliminating one of the spatial dimensions and producing a 1-D (radial) representation. If present, an isotropic scale-size will be revealed by this 1-D representation and may be extracted from each proton radiograph.

Figure 10 illustrates two examples of the AC analysis with synthetic data and two examples from experimental radiographs. Two synthetic images were generated with 2-D sinusoidal functions (eggcrates) of different amplitudes with diagonal peak-to-peak wavelength of 106 μm and sinusoidal amplitudes of 10% (top) and 30% (bottom) of the mean. The corresponding AC spectra are shown for both images with the measured dominant scale size of the features as 113 ± 8 μm. (b) Experimental radiographs from 1.4 ns (top) and 1.6 ns (bottom) with corresponding AC spectra. The earlier image does not show a peak in the AC spectra which suggests that there is no isotropic scale size in the image. Whereas, the radiograph at 1.6 ns clearly shows a peak at 185 ± 10 μm indicating a dominant feature in the image.

The dominant scale size of isotropic features in an image was calculated from the 1-D AC spectrum by fitting a curve to the first observed peak at a length greater than zero. A 2nd-degree polynomial was found to fit most data better than a Gaussian, or other peaked functions. The primary goal of the fit is to obtain an accurate measurement of the peak algorithmically. Furthermore, using the fitted curve, an uncertainty in the measured peak position may be estimated. The width of the peak represents the uncertainty in the dominant scale size of the observed features, though the typical FWHM metric is not a well-defined quantity in most cases. Therefore, the width of the parabola is taken at the point where it has reached 95% of the value at the peak. This width is the uncertainty in the measured scale size of cellular features.

Figure 10(b) illustrates two sample proton radiographs of laser-irradiated CH foils with ~120 μm ridge-like perturbations. The first image at t ~ 1.4 ns shows the expected linear behavior in the image and the corresponding AC coefficient spectrum (solid) does not indicate any peak after the initial fall off, suggesting that there is no dominant, isotropic scale-size. The bottom image in Figure 10(b) occurs later in time, at t ~ 1.6 ns, and the corresponding AC coefficient spectrum (dashed) is shown. This spectrum shows a faster fall off, indicating higher amplitude modulations than the previous image, and a dominant scale size peak at 185 ± 10 μm. This systematic analysis technique provides an accurate measurement of the dominant scale size, when present, of isotropic features observed in radiographic images.

FIG. 10. (a) Two synthetic images with eggcrate-like perturbations with diagonal peak-to-peak wavelength of 106 μm and sinusoidal amplitudes of 10% (top) and 30% (bottom) of the mean. The corresponding AC spectra are shown for both images with the measured dominant scale size of the features as 113 ± 8 μm. (b) Experimental radiographs from 1.4 ns (top) and 1.6 ns (bottom) with corresponding AC spectra. The earlier image does not show a peak in the AC spectra which suggests that there is no isotropic scale size in the image. Whereas, the radiograph at 1.6 ns clearly shows a peak at 185 ± 10 μm indicating a dominant feature in the image.


35. Recent work by Willingale et al. demonstrated B field advection due to the Nernst effect at speeds ~8000 μm/μs. Though, these experiments had I ~10^{17} W/cm^2 with λ ~1.053 μm, whereby hot-electron generation is more prominent and nonlocal thermal conduction plays a significant role.