

# Experimental demonstration of dephasing reduction in an optically guided laser-plasma accelerator

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Laser-plasma accelerators offer a compact means of producing high-energy electron beams, but their performance is fundamentally limited by dephasing between the accelerated electrons and the plasma wave. To overcome this limitation, we investigate the combination of plasma density tapering and optical guiding to extend the effective acceleration length. Using a Joule-class femtosecond laser coupled into an optical-field-ionized plasma waveguide with a controlled density gradient, we experimentally achieve electron beam energies exceeding 1.6 GeV, a 40% increase compared to the constant-density case. Particle-in-cell simulations reproduce the main experimental features and reveal the central roles of delayed injection, nonlinear laser evolution, and self-focusing in enhancing energy gain.

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## I. INTRODUCTION

Laser-plasma acceleration is a promising alternative to conventional accelerators due to the high-amplitude accelerating fields generated when an ultraintense laser pulse travels through an underdense plasma, driving a strong plasma wave in its wake [1,2]. These accelerating fields, known as wakefields, can reach up to 100 GV/m. Maximizing the interaction length between the accelerated electron bunch and the accelerating field is essential to fully exploit these fields. However, this length is limited by both laser diffraction in the plasma and dephasing between the electron bunch and the driving pulse.

Diffraction causes the laser beam to diverge after propagating a certain distance in the plasma, reducing its intensity to a point where it can no longer efficiently excite a plasma wave. This limitation can be overcome by using a plasma waveguide, which consists of a plasma with a curved radial density profile that helps to maintain the laser's focus over extended distances [3–8], and has been used to generate up to 10 GeV electron beams [7].

Dephasing arises from the velocity mismatch between the electron beam, which travels at  $v_e \approx c(1 - 1/2\gamma^2) \approx c$ , and the laser, which propagates in the plasma at the group velocity  $v_g \approx (1 - n_e/2n_c) < c$ , where  $n_e$  and  $n_c$  are the electron and critical plasma densities, respectively. This mismatch causes the electron beam to drift from the back toward the front of the plasma cavity, gradually experiencing a weaker accelerating field. In extreme cases, the electron bunch may even move

beyond the cavity center, thus experiencing a decelerating field. The characteristic length of this process, or the so-called dephasing length  $L_d$ , is proportional to  $n_e^{-3/2}$  [9]. Thus, dephasing can be mitigated by reducing the plasma density  $n_e$ , but at the cost of a lower accelerating field and higher laser power needed for self-guiding [10].

To ensure that the electrons remain in the accelerating phase, the length of the plasma cavity can be dynamically shortened ensuring that the electron beam remains at the back of the cavity where the electric field is maximum [11]. This technique, known as tapering, has been demonstrated in pioneering experiments using either a linear density gradient [12–14] or a density step [15]. Yet, these experiments have so far yielded electron beams with energies below 500 MeV, primarily because of the lack of efficient laser guiding.

Here, we show that this approach can be combined with optical guiding to counteract diffraction and mitigate dephasing simultaneously. In the first section, we use a simple analytical model to examine the impact of a rising density gradient on dephasing. We then present experimental results demonstrating the production of electron beams up to 1.8 GeV with a Joule-class laser, followed by an analysis using particle-in-cell (PIC) simulations.

## II. DEPHASING IN A RISING DENSITY GRADIENT

In this section, we use a simple model to demonstrate how longitudinal shaping of the plasma density can reduce the effect of dephasing. To mitigate dephasing, we ideally want the electron bunch to remain in the same phase of the wakefield throughout the entire propagation; that is, the phase velocity of the wakefield at the position of the electron bunch should be equal to the beam velocity  $v_e \approx c$ . The wakefield phase can be expressed as

$$\phi(z, t) = \pi \frac{z_L(t) - z}{r_B(z_L)}, \quad (1)$$

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where  $r_B \propto 1/\sqrt{n_e}$  is the radius of the accelerating cavity and  $z_L$  is the position of the laser front, which we assimilate with the upper boundary of the first plasma cavity.

From this, the phase velocity of the wakefield is given by

$$v_\phi(\xi, z) = \left(\frac{d\phi}{dt}\right) \left(\frac{d\phi}{dz}\right)^{-1} \approx v_g(z) \left(1 + \frac{\xi}{2n_e} \frac{dn_e}{dz}\right), \quad (2)$$

where  $\xi = z_L - z$  represents the distance to the front of the laser [9,16]. This expression assumes that the density gradient scale length,  $n_e/(dn_e/dz)$ , is much larger than both the plasma wavelength  $\lambda_p$  and  $\xi$ . These assumptions are well verified in most practical cases. The equation also assumes that  $r_B$  depends only on the local density. This hypothesis is more problematic, as in reality the length of the cavity also depends on the laser intensity, which varies during propagation. Therefore, Eq. (2) should be regarded more as a qualitative analysis tool than as a means of prediction. Specifically, the equation shows that for an increasing density profile:

- (1) The phase velocity  $v_\phi$  exceeds the laser group velocity  $v_g$ .
- (2) The velocity difference  $v_\phi - v_g$  is proportional to the distance to the laser  $\xi$ .
- (3) The phase velocity  $v_\phi$  can become superluminal if the plasma density increases too steeply.

The condition that the phase velocity at the rear of the bubble is equal to  $c$ , i.e.,  $v_\phi(2r_b, z) = c$ , defines a differential equation for the plasma density  $n_e(z)$  that can be easily solved, leading to [17,18]

$$n_e(z) = \frac{n_0}{(1 - z/L_0)^{2/3}}, \quad (3)$$

with  $n_0 = n_e(0)$ ,  $L_0 = (2/3)L_{d0}$ , and  $L_{d0}$  the dephasing length at  $n_0$ . Such a density profile would be challenging to achieve experimentally due to the limited means for shaping the density. Moreover, it was derived under the assumption of constant laser intensity, which is known to be inaccurate; as a result, the actual profile that would maintain the beam at the rear of the cavity could differ significantly from Eq. (3) [19]. Therefore, it is meaningful to analyze the dynamics of dephasing in simpler density profiles, which are more readily achievable experimentally. In the following, we focus on linear density profiles, which can be implemented in practice using a simple single-inlet slit gas jet configuration.

In the general case, where the wakefield velocity is not locked to the electron velocity, the comoving position of the bunch is  $\xi_e = z_L - z_e$ , and the phase of the electron bunch in the cavity, assuming  $v_e = c$ , is

$$\phi_e(z_e) = \frac{\pi}{r_B(z_e)} \left( \int_{z_0}^{z_e} \left( \frac{v_g(z)}{c} - 1 \right) dz + \xi_0 \right), \quad (4)$$

where  $z_0$  is the injection position,  $\xi_0 = \xi_e(z_0)$ , and both  $r_B$  and  $v_g$  are functions of the local plasma density  $n_e(z)$ . Due to the velocity mismatch, the electron beam experiences a varying electric field during acceleration. In the bubble regime, this field is given by

$$E_z(z_e) = -\frac{en_e(z_e)}{\epsilon_0} (\xi_e(z_e) - r_B(z_e)), \quad (5)$$

with  $e$  the electron charge and  $\epsilon_0$  the vacuum permittivity [20]. The energy gain from the injection position is then directly

obtained from Eq. (5):

$$\Delta\gamma(z_e) = \frac{q}{mc^2} \int_{z_0}^z E_z(z') dz'. \quad (6)$$

The three quantities  $\phi_e$ ,  $E_z$ , and  $\Delta\gamma$  are plotted in Figs. 1(a)–1(c) for a constant density  $n_0$ , the profile given by Eq. (3), and two linear density profiles defined as

$$n_e(z) = n_0 \left( 1 + \alpha \frac{z}{L_d} \right), \quad (7)$$

with  $\alpha = 1$  (green curve) and  $\alpha = 1.5$  (red curve).

In the constant-density case, the dephasing rate is constant (i.e.,  $\phi_e$  is a linear function of  $z$ ), and the beam experiences progressively lower fields as it drifts through the cavity. Consequently, the energy gain follows a parabolic curve, reaching a maximum of  $\Delta\gamma_{\max}$  at  $L_{d0}$ , when the beam reaches the center of the cavity. For the optimal density profile, the beam remains in the same phase by design, while the amplitude of the electric field experienced by the beam increases during acceleration as  $E_z \propto \sqrt{n_e}$ . Consequently, the energy gain increases faster than in the reference case, reaching its maximum at  $z = L_0$  [17,18], where the cavity becomes infinitely small, stopping further acceleration.

For a linear profile with  $\alpha = 1$ ,  $\phi_e$  follows the optimal phase over a short distance (approximately  $0.05L_d$ ), after which it starts dephasing, as the group velocity decreases with increasing plasma density. This ultimately results in a smaller effective dephasing length. The longitudinal electric field experienced by the electron bunch follows a similar trend, initially increasing in amplitude for small  $z$  and then decreasing faster than in the constant-density case as the beam approaches the center of the cavity. Overall, as the beam spends more time in the region of high fields, this results in an energy gain that exceeds  $\Delta\gamma_{\max}$  for  $z \leq L_{d0}$ , with the maximum energy gain reached around  $z \sim 0.88L_{d0}$ , where  $\Delta\gamma \approx 1.36\Delta\gamma_{\max}$ .

For a linear profile with  $\alpha = 1.5$ ,  $\phi_e$  initially increases, which means that the wakefield is superluminal. Such superluminal propagation should prevent any electron injection. We therefore consider that the injection occurs at  $z \approx 0.2L_d$ , where  $d\phi_e/dz$  changes sign, meaning that the wakefield propagation becomes subluminal. The evolution of the electric field has the same shape than in the case  $\alpha = 1$ , but the dephasing rate occurs faster, which results in a maximum energy gain reached around  $z \approx 0.94L_d$  where  $\Delta\gamma \approx 0.94\Delta\gamma_{\max}$ .

In conclusion, the linear profile with  $\alpha = 1$  stands out as the most promising among the linear density profiles, as illustrated in Fig. 1(d). We observe that for  $\alpha > 1$ , the initial superluminal phase velocity of the wake prevents injection and results in a shorter effective acceleration length. For  $\alpha < 1$ , the mitigation of dephasing is less effective and results in a lower maximum energy gain than for  $\alpha = 1$ , which is the optimum as it minimizes dephasing while maintaining a phase velocity that remains consistently subluminal, ensuring a higher maximal energy gain. We emphasize that this simplified model is intentionally designed to provide qualitative insights rather than quantitative precision. While it effectively captures the essential physics governing dephasing in linear density gradients, it does not account for additional effects, such as self-focusing, pump depletion [5,6], and high-power

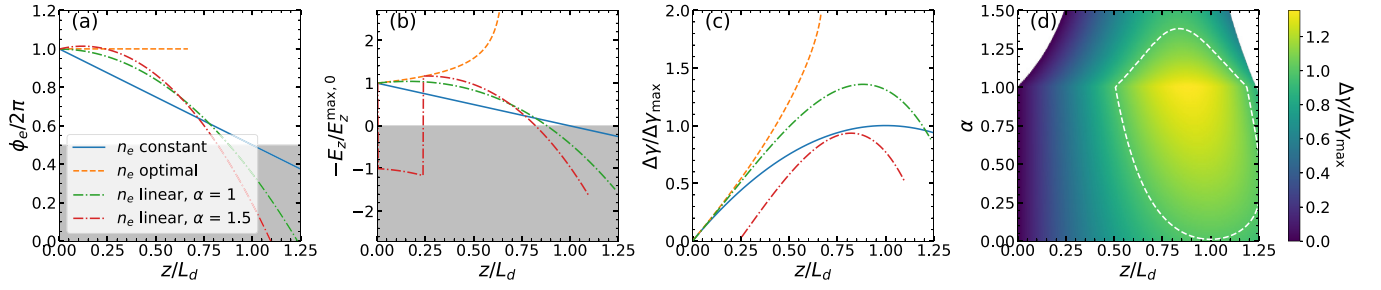


FIG. 1. Comparison between the constant-density case, the density profile defined by Eq. (3), and two linear density profiles given by Eq. (7) for  $\alpha = 1$  and  $\alpha = 1.5$ . Panel (a) shows the phase of an electron bunch injected at the back of the cavity at  $z_0 = 0$ , and panel (b) displays the longitudinal electric field experienced by the electron bunch. Panel (c) shows the energy gain as a function of  $z$ . Panel (d) shows the energy gain as a function of  $z$  and  $\alpha$ . The contour for the value  $\Delta\gamma = \Delta\gamma_{\max}$  is shown in white dashed line. The gray area in panels (a) and (b) indicates the region where the field is decelerating. In panels (b)–(d), the electric field and energy gain computed for the linear cases are truncated at the  $z$  position where  $\phi_e = 0$ , which corresponds to the electron bunch reaching the position of the laser pulse.

guiding [21–23], which leads to an inaccurate estimate of  $L_D$ . For a more rigorous and accurate analysis, PIC simulations are indispensable [19], as we demonstrate in Sec. IV.

### III. EXPERIMENTAL RESULTS

We set up an experiment to test this conclusion and determine the optimal linear density gradient for maximizing the electron beam energy. In order to observe only effects related to dephasing, it is necessary to ensure that the beam acceleration is not limited by another phenomenon. A plasma waveguide is thus used to overcome laser diffraction and maintain the laser’s focus throughout its propagation in the plasma [3].

The experiment was conducted at Laboratoire d’Optique Appliquée (LOA) using the “Salle Jaune” laser system. The plasma waveguide was produced through hydrodynamic optical field ionization (HOFI). This method involves focusing a laser along a line longer than the gas target to ionize it and create a plasma column. The hydrodynamic expansion of this column during a few nanoseconds, combined with the ionization of the remaining neutrals, results in the formation of a waveguide capable of maintaining the laser pulse focused over long distances [4,24–27]. Here, the guiding laser ( $E \sim 2.3$  mJ,  $\tau \sim 30$  fs) was focused by an off-axis axiparabola [28] with a nominal focal length  $f_0 = 200$  mm and maximum focal depth  $\delta_0 = 30$  mm. The focal line of this axiparabola is defined by  $f(r) = f_0 + 1/a \ln(r/R \times e^{a\delta_0})$ , with  $a = 1/\delta_0 \ln(R/r_{\text{hole}})$ , where  $r$ ,  $r_{\text{hole}} = 6$  mm, and  $R = 38.1$  mm denote the radial coordinate, the radius of the hole, and the radius of the axiparabola, respectively. A deformable mirror was used to optimize the focal spot of the beam focused by the axiparabola.

The driver beam had a duration of 30 fs. It was focused at the entrance of the waveguide by an F/18 spherical mirror used off axis, with the astigmatism introduced by the off-axis geometry corrected by a second deformable mirror. The focal spot had an FWHM diameter of  $D_{\text{FWHM}} = 26$   $\mu\text{m}$ , and the energy encircled within the first zero ring reached  $1.05 \pm 0.05$  J, corresponding to 73% of the total laser energy. The laser pointing stability was improved using an active stabilization system [29], achieving a peak-to-peak stability of  $0.7D_{1/2}$ , and standard deviations  $\sigma_x = D_{\text{FWHM}}/5.7$ ,

$\sigma_y = D_{\text{FWHM}}/6.3$  in the horizontal and vertical directions, respectively. This level of stability is sufficient to ensure substantial laser injection into the plasma waveguide on each shot, even though the fraction of coupled energy may still vary.

To produce the linear density gradient, a slit nozzle was mounted on a rotative stage, which allows us to change the angle between the laser axis and the gas nozzle. The density profiles measured for different angles are provided in the Appendix B. The nozzle was fed with a gas mixture of molecular hydrogen and 1% nitrogen. A blade made of a 200- $\mu\text{m}$ -thick silicon wafer was used to partially obstruct the gas flow at the entrance of the gas jet to generate a hydrodynamic shock. The purpose of this shock is to produce a sharp density transition that can trigger the injection of electrons into the wakefield through shock-assisted ionization injection [30]. The blade was also used to reduce the effective target length by adjusting the position of the nozzle underneath it, allowing us to maintain the relative positioning between the density transition and the focal plane of the main laser. Finally, a probe laser beam (a few mJ,  $\tau \sim 30$  fs) is used to perform shadowgraphy measurements of the interaction and to measure the electronic density with a wave front sensor (Phasics SID4-UHR). An example of shadowgraphy is displayed in Fig. 2. It shows the position of the blade, the plasma shadow, the generated shock,

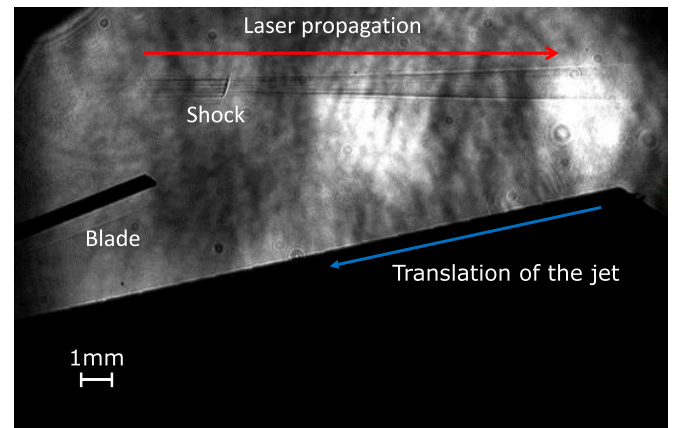


FIG. 2. Shadowgram of the tilted target, showing the plasma created by the laser, the blade to trigger the injection, and the shock created by the blade.

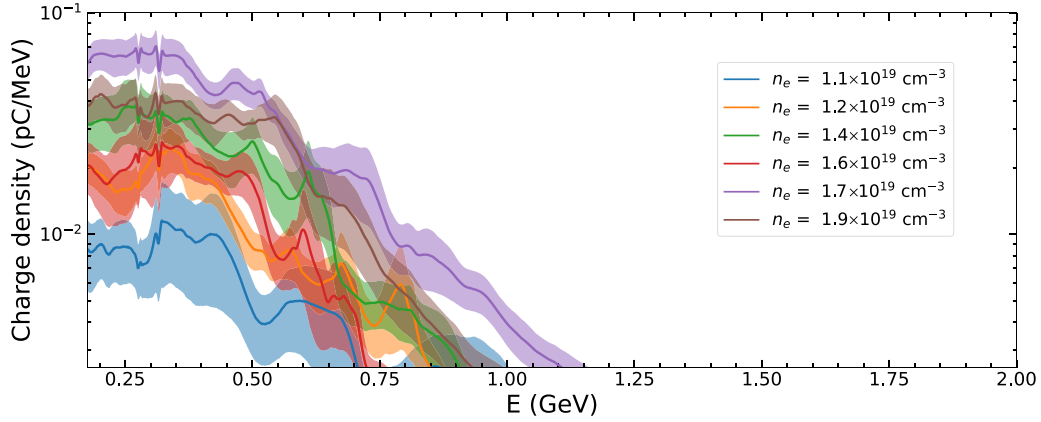


FIG. 3. Mean spectra obtained the absence of a gradient, i.e., with an untilted gas jet, for different operating pressures. The mean is computed over ten shots, without excluding any, for each pressure. The shaded area corresponds to the standard deviation over the same ten shots.

and the translation axis of the gas nozzle. A simple electron spectrometer, consisting of a 40-cm-long, 0.85 T magnetic dipole and a scintillating screen, was used to disperse and detect the electron beam (see details in Appendix A).

We first produced electron beams using a 15-mm-long nozzle and without introducing a density gradient (see density profile in Fig. 10 in Appendix B). The mean spectra obtained at various plasma densities are shown in Fig. 3. The spectra exhibit a relatively large standard deviation, indicating that the energy coupled into the fundamental mode, which is responsible for wake excitation [6], varies from shot to shot. However, all shots result in energies exceeding 500 MeV, surpassing the maximum energy obtained without guiding [8]. The electron spectra are continuous, indicating that trapping at the density transition is not sufficient to fully load the wakefield and suppress further injection via ionization injection [31]. In practice, electrons are therefore injected through both ionization injection and shock-assisted ionization injection. Since the spectra obtained are continuous, we chose to compare the cutoff energy for different configurations. We defined the cutoff at a threshold of 2.5 fC/MeV, which corresponds to roughly three times our detection threshold. The cutoff energy ranges from 600 MeV to 1 GeV, with a maximum of  $1.1 \pm 0.06$  GeV, obtained for a backing pressure of 27 bar and a measured density without the guide of  $n_e \sim 1.7 \times 10^{19} \text{ cm}^{-3}$ . The use of a longer 20 mm target with a similar geometry yielded no further energy gains, even at lower densities, potentially indicating laser depletion or reduced guiding efficiency. At higher operating pressures, the cutoff energy decreases, which suggests that either dephasing or depletion may become significant. In the optimal case, the electron density in the guiding channel is estimated to be on the order of  $3.4 \times 10^{18} \text{ cm}^{-3}$  [4].

We then tested different configurations of tilted nozzles to mitigate dephasing and found an optimum at an angle of  $8^\circ$  and a backing pressure of 26 bar, for a 15-mm-long slit jet. The corresponding density profile is shown in Fig. 10 in Appendix B. In this configuration, the initial plasma density without guiding is on the order of  $n_0 = 9.1 \times 10^{18} \text{ cm}^{-3}$ , with a density gradient of  $0.03n_0/\text{mm}$ . The effective plasma length was scanned by translating the nozzle underneath the blade.

Mean spectra for each target length are shown in Fig. 4(a). An optimum is observed at  $L_{\text{target}} = 12.8$  mm, where the beam reaches a cutoff energy of  $1.6 \pm 0.1$  GeV, corresponding to an increase in electron energy by a factor of  $\approx 1.4$  compared to the reference case, thus demonstrating the effectiveness of plasma tapering in mitigating dephasing. At longer target lengths, the cutoff energy decreases noticeably, down to  $1.2 \pm 0.1$  GeV for  $L_{\text{target}} = 15.9$  mm.

The charge above 1 GeV initially increases with the target length, increasing from  $1.3 \pm 1$  pC for  $L_{\text{target}} = 11.9$  mm to  $4.2 \pm 1.4$  pC for target lengths of 12.8 mm, before decreasing to  $1.6 \pm 0.6$  pC for  $L_{\text{target}} = 15.9$  mm. The performance achieved with this density gradient is further illustrated in Fig. 4(b), which shows an angularly resolved spectra for a selected shot, with some electrons reaching energies exceeding 1.8 GeV. In panel (b), the beam charges above 1 and 1.5 GeV are 9.4 and 3 pC, respectively.

The model presented in Sec. II is too simplified to quantitatively predict the optimal density gradient under our experimental conditions. We therefore experimentally tested several other density gradients. Figure 5 shows the average spectra obtained for an initial density (before channel generation) of  $n_0 = 1.1 \times 10^{19} \text{ cm}^{-3}$  and a 20-mm-long nozzle with an inclination of  $12^\circ$ . At such a steep angle, the density profile starts to deviate from linear (see the corresponding density profile in Fig. 10 in Appendix B); we nonetheless extract a mean density gradient from a linear fit, yielding a gradient on the order of  $0.05n_0/\text{mm}$ . The initial density was chosen after a density scan, as it yielded the highest beam energy for this density gradient. Compared to the  $0.03n_0/\text{mm}$  gradient case, we observe an overall decrease in energy, although the electron energy still increases with target length. This reduced energy suggests a delayed injection, which limits the effective acceleration length. The delayed injection is likely caused by an overly steep density gradient that produces a superluminal phase velocity at the beginning of propagation, thereby hindering early injection. This scenario would correspond to the one illustrated in Fig. 1 for  $\alpha = 1.5$ . A second effect attributed to an excessively delayed injection is the shorter total injection length, which leads to an energy distribution peaked at high energy, as shown in Fig. 5(b).

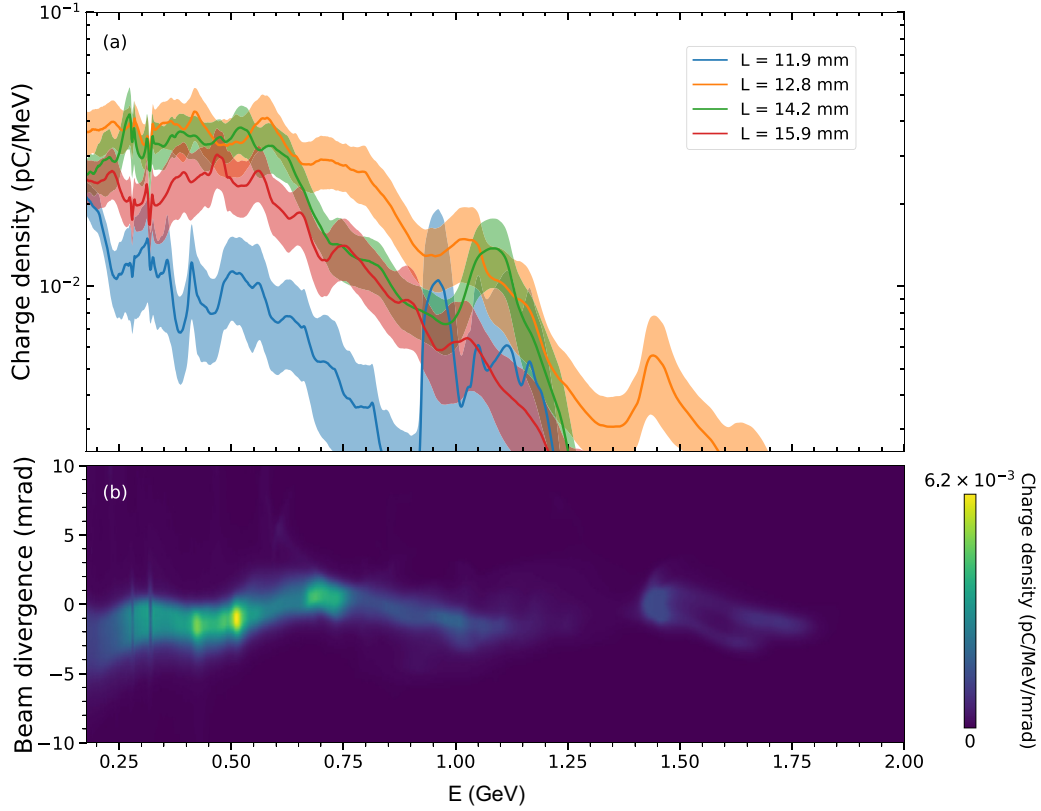


FIG. 4. Spectra obtained with an initial plasma density without guiding is on the order of  $n_0 = 9.1 \times 10^{18} \text{ cm}^{-3}$ , with a density gradient of  $0.03n_0/\text{mm}$ . (a) Mean spectra obtained for different acceleration length. The mean is computed over ten shots, without excluding any, for each pressure and the shaded area corresponds to the standard deviation over the same ten shots. (b) Angularly resolved spectra of a selected shot, with a linear colorbar.

Quantitatively, the cutoff energy remains below 1.15 GeV for all target lengths. The total charge above 200 MeV is relatively stable for  $L_{\text{target}} > 10 \text{ mm}$ , reaching approximately 15 pC, while the charge above 1 GeV reaches a maximum of  $1.6 \pm 0.9 \text{ pC}$  for  $L_{\text{target}} = 14.5 \text{ mm}$ , which was the maximum target length we could reach in this configuration. The absence of evidence for dephasing, particularly the steady increase in cutoff energy with target length, led us to increase the plasma density. We found that a 50% increase in initial density resulted in clear evidence of dephasing around  $L_{\text{target}} \sim 12.5 \text{ mm}$ . However, this was accompanied by lower electron energies, reduced charge above 1 GeV, and more irregular energy-angle distributions. These results suggest that the initial plasma density was close to optimal for a  $0.05n_0/\text{mm}$  gradient. They also confirm the superior performance of the  $0.03n_0/\text{mm}$  configuration, which was further supported by tests at a  $3^\circ$  inclination. In this case, the results were similar to those obtained at a  $0.03n_0/\text{mm}$  gradient, but with the acceleration of slightly lower-energy electrons.

#### IV. NUMERICAL SIMULATION

To support our interpretation of the experimental result, we ran simulations using the pseudospectral quasicylindrical particle-in-cell code fbpic [32]. We assumed a perfect Gaussian beam with an energy of 1 J, a focal spot size (intensity FWHM) of  $25 \mu\text{m}$ , and a Gaussian temporal envelope of 25 fs, corresponding to  $a_0 \approx 1.6$ . To investigate different

interaction regimes, we tested various initial plasma densities and density gradients, which enabled us to distinguish between different regimes of interaction. For simplicity, the plasma was assumed to be preionized with a preformed plasma channel.

A key challenge in combining plasma tapering with guiding is ensuring that the waist of a matched laser evolves as the radius of the plasma cavity, which scales with the plasma wavelength  $\lambda_p \propto n_e^{-1/2}$ . To maintain efficient guiding and avoid significant laser leakage, the waveguide diameter should follow the same evolution. In the case of a linear density gradient, this implies that the waveguide diameter, and thus the waist of the HOFI laser, should scale as  $z^{-1/2}$ . Such a decreasing waist naturally arises from quasi-Bessel beams, like those produced at the focus of an axiparabola. In particular, an axiparabola designed for a constant-intensity focal line generates a focal spot whose diameter scales precisely as  $z^{-1/2}$ , potentially enabling perfect matching, provided the proportionality factor is correct. This factor is primarily determined by the axiparabola's numerical aperture and focal depth [28], which should thus be specifically chosen to match a given density gradient.

In our experiment, we used a single, constant-encircled-energy axiparabola and varied the target density gradient. As a result, the waveguide diameter did not perfectly follow the evolution of  $\lambda_p$ , which may have led to increased laser leakage during acceleration. However, to simplify the analysis and

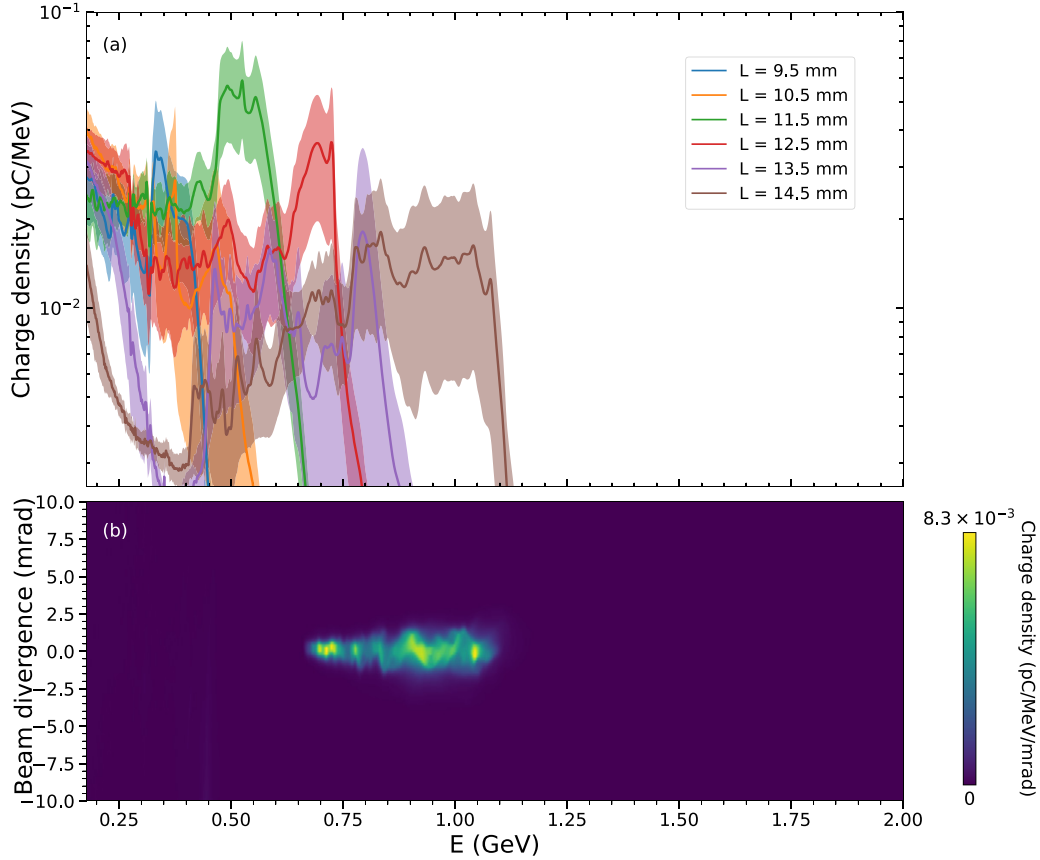


FIG. 5. Spectra obtained with an initial density without the channel-forming beam of  $n_0 = 1.1 \times 10^{19} \text{ cm}^{-3}$  and a density gradient of  $0.05 n_0/\text{mm}$ . (a) Mean spectra for different acceleration length. The mean is computed over ten shots, without excluding any, for each target length and the shaded area corresponds to the standard deviation over the same ten shots. (b) Angularly resolved spectra of a selected shot, with a linear colorbar.

limit the number of parameters, we initially neglect this effect and assume in the simulations a parabolic channel with a constant channel radius of the form

$$n_e(r, z) = n_0(z) + \frac{1}{\pi r_e w_0^2} \left( \frac{r}{w_0} \right)^2, \quad (8)$$

with  $r_e$  the electron classical radius and  $w_0 \sim 21 \text{ } \mu\text{m}$  the waist radius of the electric field.

Figures 6(a)–6(c)(1)–(2) report on the case of an initial density  $n_0 = 6.2 \times 10^{17} \text{ cm}^{-3}$  without gradient [panels (a)–(c1)] or with a gradient  $\delta_n = 0.06 n_0/\text{mm}$  [panels (a)–(c2)]. Panels (a1) and (a2) show the evolution of the energy of a few electrons from the high-energy tail. In both cases, the electrons reach a final energy of approximately 750–800 MeV. However, in the gradient case [panel (a2)], this occurs over nearly half the distance, due to delayed injection and a stronger average accelerating field. The delayed injection is caused by an initially superluminal wakefield (corresponding to the case  $\alpha > 1$  in Sec. II). In this case, electrons are injected far from the density transition through standard ionization injection. This is evidenced in Fig. 6(b2), which shows the amplitude of the wakefield  $E_z$  as a function of the comoving coordinate  $\xi = z - ct$  and the longitudinal position  $z$ . In this coordinate system, a luminal wake propagates horizontally, while a superluminal one appears with a positive slope. As shown, the wakefield is superluminal up to  $z \approx 10 \text{ mm}$  at

which point injection occurs. This scenario is thus close to the one depicted by the red curves in Fig. 1.

We also observe that the electron beam [indicated by the thick yellow lines in Fig. 6(b2)] remains in the first quarter of the wakefield, far from the dephasing region. This suggests that longer acceleration lengths would have led to higher final energies. According to Lu’s model [20], the dephasing length for the constant-density case is estimated as  $L_{D0} \approx 33 \text{ mm}$ , which is consistent with the absence of dephasing over 20 mm. Moreover, this estimate yields  $\alpha_{\text{theo}} = \delta_n L_{D0}/n_0 \approx 2$ , which is also consistent with the wakefield being initially superluminal.

Panels (c1) and (c2) in Fig. 6 show laser amplitude oscillations along the propagation, which could be due to imperfect matching, or ponderomotive channel modification [6]. In addition to this, the gradient case exhibits an increase in amplitude, revealing the contribution of relativistic self-focusing and stronger guiding at higher density. This effect gradually increases the nonlinear plasma wavelength, causing the back of the plasma cavity to slow down. This effect, not accounted for in the simplified model of Sec. II, tends to enhance dephasing and shorten the effective dephasing length, which in turn lowers the effective  $\alpha$ . This increased rate of dephasing explains why, in the experiment, the optimal initial plasma density is lower in the gradient case (Fig. 4) than in the reference case (Fig. 3).

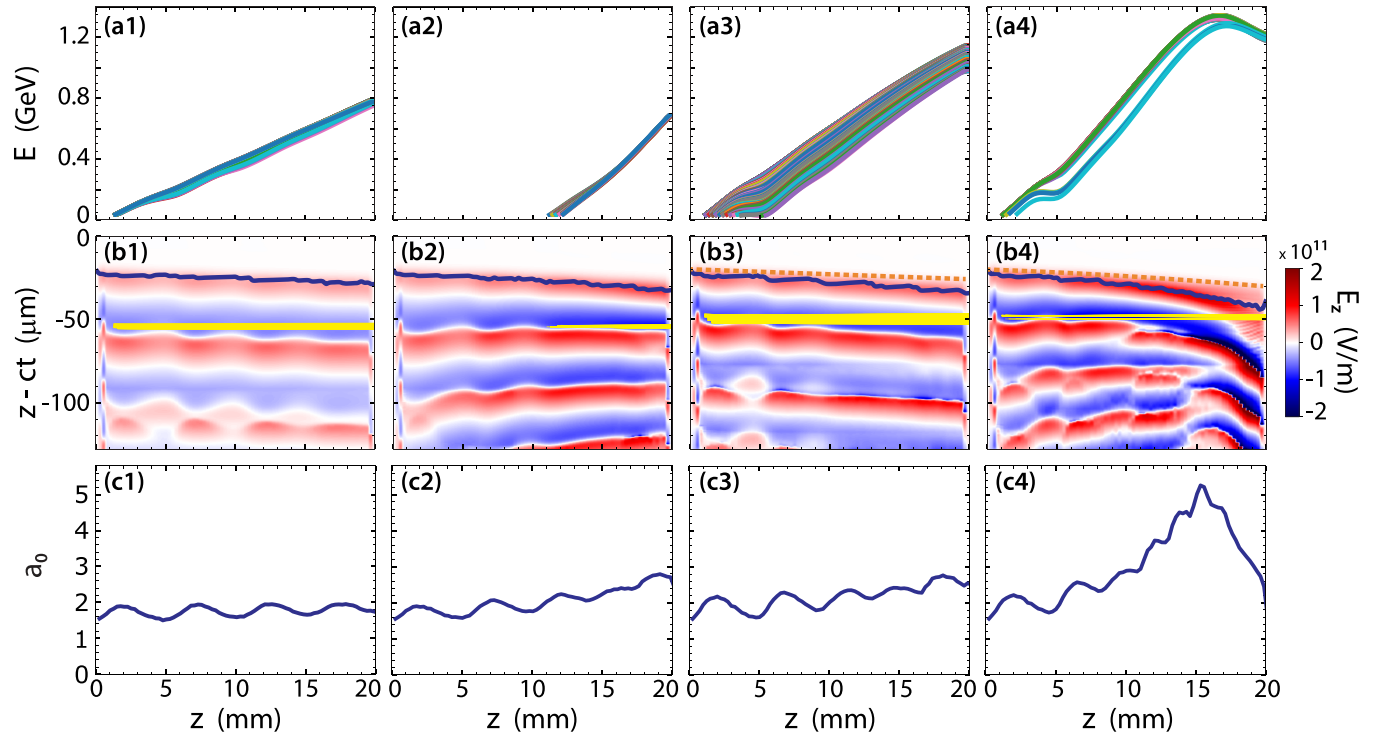


FIG. 6. Result of PIC simulations. Columns 1 and 2 are for an initial density of  $n_0 = 6.2 \times 10^{17} \text{ cm}^{-3}$ . Columns 3 and 4 are for an initial density of  $n_0 = 1.0 \times 10^{18} \text{ cm}^{-3}$ . Columns 1 and 3 are the reference case without gradient, for these two densities. Column 2 is for a density gradient of  $0.06n_0/\text{mm}$ , and column 4 is for a density gradient of  $0.08n_0/\text{mm}$ . (a) Energy evolution of test particles selected from the high-energy tail. The particles are selected by looking for the peak in the spectrum at highest energy, with a peak charge density above  $2.5 \text{ fC/MeV}$ . (b) On-axis longitudinal electric field  $E_z$ ; the blue line indicates the position of the peak transverse field  $E_x$  (i.e., the laser pulse maximum), while the yellow lines mark the positions of the test electrons. (c) Evolution of the normalized laser amplitude  $a_0$  along the propagation axis.

We draw two conclusions from this initial simulation: To increase the final energy, the plasma density should be increased to better exploit the target length, and the effective  $\alpha$  factor should be lowered to enable earlier injection. Figures 6(a)–6(c4) present the results obtained for  $n_0 = 1 \times 10^{18} \text{ cm}^{-3}$  and  $\delta_n = 0.08n_0/\text{mm}$ , which gives  $L_{D0} \approx 15 \text{ mm}$  and  $\alpha_{\text{theo}} \approx 1.2$ . According to Sec. II, such a value of  $\alpha$  may still inhibit injection in the early stage of propagation. However, we previously saw that the effective  $\alpha$  can be significantly reduced due to self-focusing. Here, Fig. 6(c4) shows that self-focusing is quite strong in the gradient case, with the laser amplitude nearly tripling over 16 mm. Other effects also contribute to modifying the wakefield velocity by altering the laser group velocity. This is illustrated in Figs. 6(b3) and 6(b4), where the orange dashed line shows the trajectory of an entity propagating at the group velocity  $v_g(z)$ . While the front of the pulse follows this trajectory, the maximum of the laser pulse (blue curve) moves significantly slower due to self-phase modulation, dispersion, and laser etching, which in turn further slows down the wakefield. The combination of these effects explains why efficient injection occurs just after the shock region, even though  $\alpha_{\text{theo}} > 1$ .

Figures 6(a) and 6(b3) show that in the constant-density case, the beam is close to dephasing after approximately 19 mm of acceleration, in qualitative agreement with the estimate of  $L_{D0}$ . In contrast, in the density gradient case, panel (b4) shows that the beam hardly dephases over the

first  $\approx 10 \text{ mm}$ , before dephasing rapidly, leading to a total acceleration length of about 16 mm. The scenario is therefore close to the case  $\alpha = 1$  in Fig. 1 (green curves), except for the faster dephasing due to the strong increase in laser intensity toward the end of the plasma. The efficient tapering increases the electron beam energy by almost 20% compared to the constant-density case, resulting in a final energy approaching 1.4 GeV, close to the experimentally measured value.

However, an important physical effect is still missing in this simulation: the influence of the transverse size of the channel-forming beam on the waveguide. The transverse size of the channel-forming beam depends on the local numerical aperture [28], which results, in our case, in a decreasing transverse size along the propagation. To address this, we consider in Fig. 7(a) nonideal waveguide. More precisely, we model a waveguide following Eq. (8), but by replacing  $w_0$  with  $w_m(z) = C(z) \times w_0$ , with  $w_0$  the size of the drive laser and  $C(z) = (r_{\text{hole}}/f_0) \times (z/(Re^{a(z-f_0-\delta_0)}))$  a numerical factor scaling as the transverse size of the channel-forming beam [28]. The reduction of the matched spot size during the propagation promotes self-focusing, as shown in Fig. 7(c) by the rapid increase of the laser amplitude. This increase in the laser amplitude results in a cavity becoming fully evacuated more efficiently, which has the effect of increasing the amplitude of the accelerating field. This higher amplitude of the accelerating field results in a higher energy gain, about 1.6 GeV, as shown in Fig. 7(a) while also increasing the dephasing due to

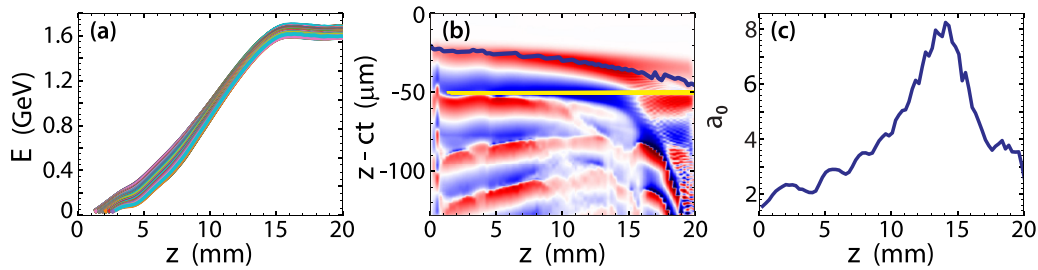


FIG. 7. PIC simulation for an initial density of  $n_0 = 1.0 \times 10^{18} \text{ cm}^{-3}$  and a density gradient of  $0.08n_0/\text{mm}$ , with a channel transverse size evolving as the spot size of the channel forming beam. (a) Energy evolution of test particles selected from the high-energy tail. The particles are selected by looking for the peak in the spectrum at highest energy, with a peak charge density above 2.5 fC/MeV. The spectral dispersion of the bunch selected is equal to the full width at half maximum of the peak. (b) On-axis longitudinal electric field  $E_z$ ; the blue line indicates the position of the peak transverse field  $E_x$  (i.e., the laser pulse maximum), while the red lines mark the positions of the test electrons. (c) Evolution of the normalized laser amplitude  $a_0$  along the propagation axis.

the increase of laser intensity. Other effects, such as the radial truncation of the channel or a nonparabolic radial profile could also be taken into account, but are beyond the scope of this study.

## V. CONCLUSION

We have experimentally demonstrated that combining plasma tapering with optical guiding enables effective mitigation of dephasing in a laser-plasma accelerator. By introducing a linear density gradient in a guided plasma, we achieved a 40% increase in electron beam energy, reaching up to 1.8 GeV with a Joule-class laser. These results are supported by particle-in-cell simulations, which emphasize the crucial role of nonlinear laser propagation and laser-plasma matching throughout the acceleration process.

Looking ahead, further energy gains could be achieved by employing more complex, nonlinear density profiles designed to lock the wakefield velocity to that of the electron beam. Such profiles could be engineered using techniques such as laser machining of gas targets [33], or a more precise control of the target design, including precise shaping of the nozzle throat [34] or the use of a multiple inlet modular gas jet [35]. Another promising avenue is the implementation of plasma tapering combined with guiding on petawatt-class laser systems, which may enable the production of electron beams exceeding 10 GeV. These developments would mark significant progress toward high-energy laser-plasma accelerators for future light sources and high-energy physics applications.

## DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

## APPENDIX A: ELECTRON SPECTROMETER

### 1. Experimental setup

The electron spectrometer consists of a dipole magnet that disperses the electron bunch within the plane perpendicular to

the poles toward a scintillating screen [Fig. 8(a)]. Electrons from the laser-plasma accelerator propagate  $194 \pm 2$  mm in vacuum before reaching the magnet entrance. A collimating slit is placed in front of the dipole, at  $150 \pm 2$  mm from the target. This aperture is made of two 1-cm-thick glass panels that are separated by 1 mm distance. The gap of the slit is perpendicular to the magnet dispersion plane, in order to leave unaffected the vertical divergence of the electron beam. The slit enables the reduction of the lateral acceptance angle down to 6.7 mrad. The light emitted by the scintillating screen is then collected by a 16-bit CCD Hamamatsu camera with a 546 nm (10 nm FWHM) Alluxa interferential filter in front of the objective.

### 2. Magnetic field

In this experiment, we used a 40-cm-long, 8-cm-wide U-shape dipole with 2 cm gap between the poles, providing 0.85 T transverse magnetic field  $B_y$ . The two-dimensional (2D) distribution of  $B_y$  was measured using a Hall probe [Fig. 9(a)] with a 1 mm spatial resolution.

Around the edges of the magnet the fringe field is inhomogeneous, and it acquires  $B_z$  and  $B_x$  components [see Fig. 9(b)]. According to the Hall probe measurements shown in Fig. 9(c), the longitudinal field  $B_z$  appears negligible except in a 6 cm region at the entrance/exit edges. Close to the dipole edges, this field reaches a maximum value of 24 mT on axis. The component  $B_x$  is also noticeable only at the dipole edge where it increases with the distance to the axis. For a 1.25 mrad transverse angle, it reaches a maximum value of 75 mT.

### 3. Detector screen

The particles are detected using a Lanex regular Carestream scintillator, oriented at  $30^\circ \pm 0.1^\circ$  with respect to the laser axis, mounted on a fixed 3D-printed support screwed to the magnet to reduce the uncertainty on the position of the scintillator. This scintillator is made of three layers: a plastic support layer, a sensitive layer with gadolinium oxysulfide ( $\text{Gd}_2\text{O}_2\text{S}$ , also called Gadox), and a protective thin plastic layer, with a total thickness of 380  $\mu\text{m}$ . The Gadox emits light with a spectrum spanning from 382 to 622 nm, and a peak at 545 nm. The imaging resolution of the Lanex surface in the dispersion plane of 0.1325 mm/pixel results in an energy

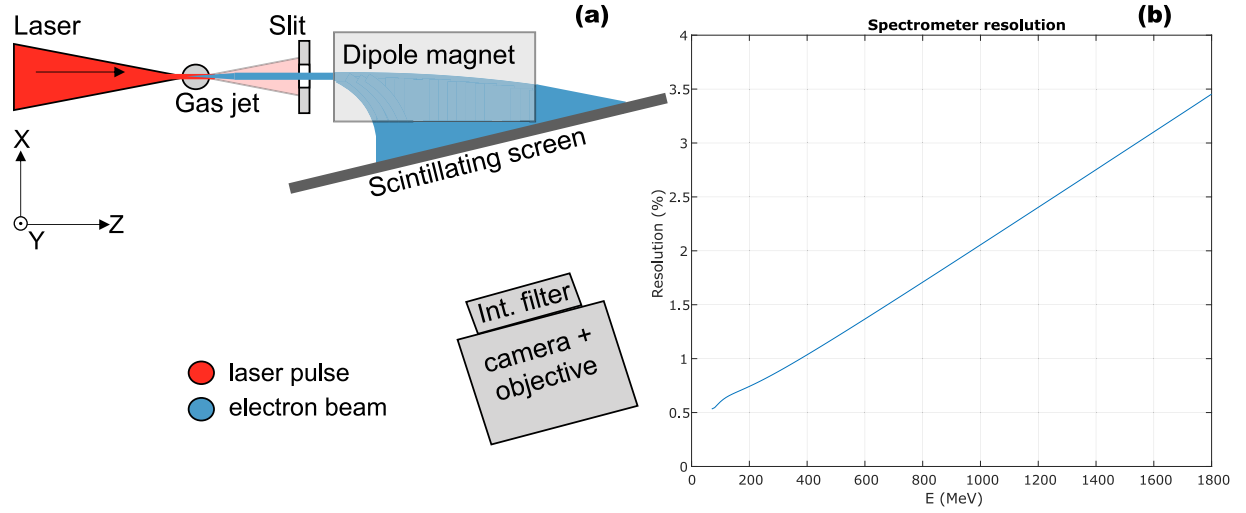


FIG. 8. (a) Schematic of the electron spectrometer. The laser pulse is focused into a gas target that is placed 204 mm before the entrance of the dipole. In front of the dipole, a vertical slit is inserted to limit the acceptance angle of the electron beam to 6.7 mrad. Electrons from the target propagate through the magnetic field toward the scintillating screen, which is imaged by a CCD camera. (b) Resolution of the electron spectrometer as a function of the beam energy for an acceptance angle of 6.7 mrad.

resolution of 2.1% at 1 GeV and 3% at 1.6 GeV. The transverse resolution of 0.1325 mm/pixel defines the resolution of the divergence measurements of 0.2 mrad.

#### 4. Calibration

For the charge calibration, we have used a  $5 \times 2 \text{ mm}^2$  radioactive tritium capsule with a luminescent phosphor layer and an initial activity of 0.76 GBq [36]. This source was absolutely calibrated to the emission efficiency of Carestream regular scintillating screen at the conventional ELBE accel-

erator, where absolute fluorescence efficiency was found to be  $(3.1 \pm 0.6) \times 10^9 \text{ ph sr}^{-1} \text{ pC}^{-1}$ . The capsule was attached close to the area of the Lanex screen where the luminescence from electrons was observed. This allowed to account for different aspects of imaging conditions that are not sensitive spectralwise, as, for instance, the numerical aperture of the camera and the quantum efficiency of the camera. The interference filter in front of the camera was calibrated in a separate experiment by measuring an attenuation for a large series of stable electron spectra. The decay time of the tritium since these measurements has also been taken into account.

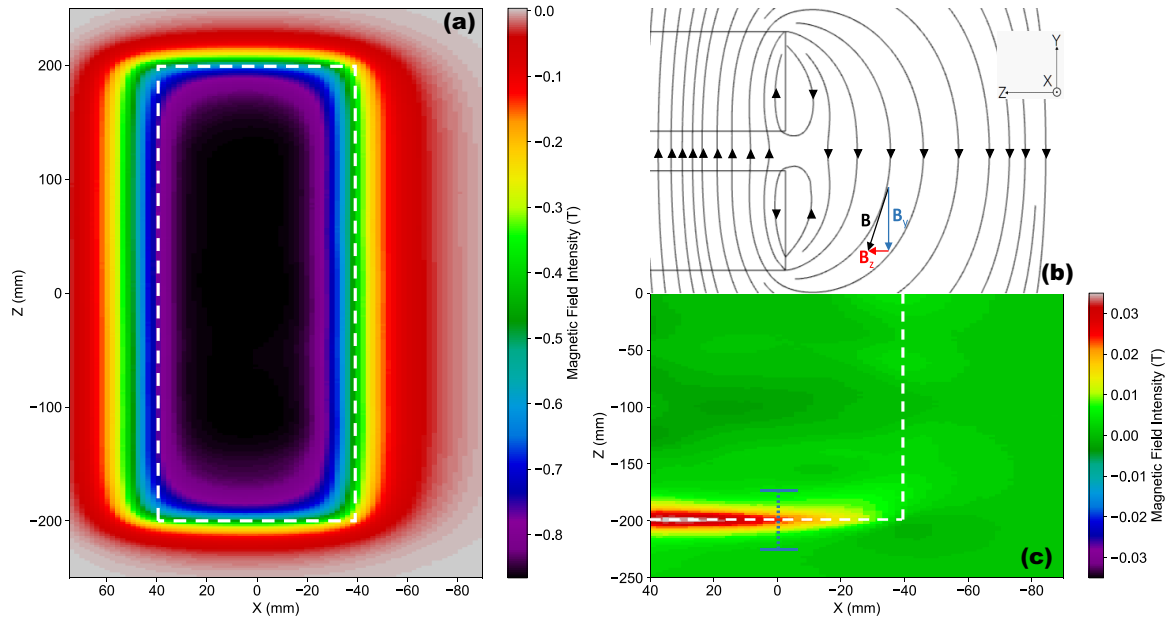


FIG. 9. Magnetic field measurements. (a) Transverse component  $B_y$  corresponding to the dispersion of the electrons toward the scintillating screen. The white dotted profile corresponds to the magnets position. (b) Schematic showing the side view of the dipole entrance,  $B_y$  and  $B_z$  components of the magnetic field. (c) Longitudinal component  $B_z$ . The blue dotted line is indicating the 6 cm region of maximum intensity.

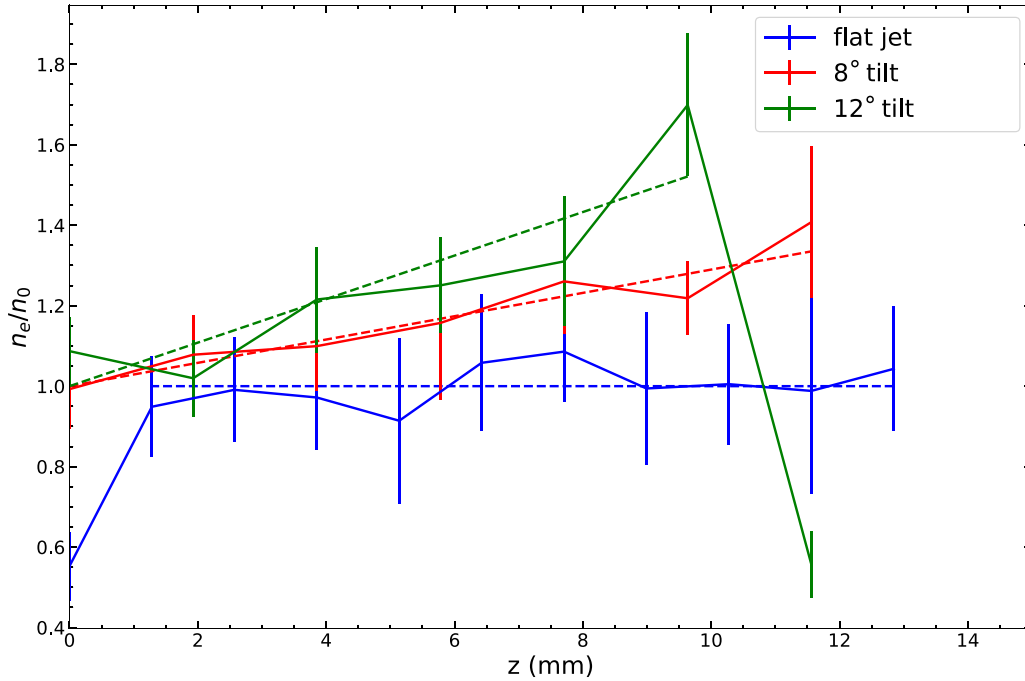


FIG. 10. Normalized density profiles for a flat nozzle, a  $8^\circ$  tilting angle, and a  $12^\circ$  tilting angle. A constant fit in the flat case and linear fits for the gradient cases are shown in dashed line.

### 5. Measurement interpretation

The electron trajectories in the magnetic field define the position at which the electrons are hitting the detector plane. The energy is estimated from the detection position on the screen, and the charge is deduced from the intensity of the emitted light. For the spectrum measurement interpretation, the electron trajectories from the source to the scintillating screen were computed numerically using a 2D particle tracking code and the measured magnetic field.

### 6. Measurement error

The electrons are injected in the laser-plasma accelerator with a microradian pointing precision, and the source position jitter is negligible compared to the 194 mm travel distance to the spectrometer. Before the dipole magnet, the slit selects the lateral angles with an acceptance of 3.3 mrad. The slit has two functions: it restricts the electron beam pointing (the source of the statistical error) and reduces the uncertainty on the beam propagation axis (the main source of the systematic error). The latter corresponds to a transverse momentum of the beam and results in an error on the measured energy, whose magnitude depends on the beam energy. Figure 8 (b) illustrates how the relative error varies with the beam energy for an acceptance angle of 6.7 mrad.

The inhomogeneities of the magnetic fringe field may also cause deviations of the electron tracks, e.g., producing an energy-dependent focusing in the vertical direction for the off-axis particles. The resulting error on the divergence measurement is given by  $\delta\theta \simeq ecB_{\text{int}}/E$ , where  $e$  and  $c$  are the electron charge and speed of light constants, respectively,  $B_{\text{int}}$  is the integrated magnetic field, and  $E$  is the electron energy. According to Hall probe measurements,  $B_{\text{int}} \approx 7.5 \times 10^{-4}$  T m, so the deflection angle does not

exceeds 0.5 mrad for 1 GeV beams, and is even lower for 1.6 GeV beams.

At the exit of the magnet, the longitudinal component of the magnetic field can introduce an angular shift in the trajectory of the electrons. Using the Hall probe measurement shown in Fig. 9 (c), and using the same formula used to estimate the defocusing effect, this deflection for a 1 GeV beam crossing the 6 cm region highlighted on the magnetic map amounts to an angle of 0.17 mrad, and of 0.01 mrad for a 1.6 GeV beam. This angular shift corresponds to a shift of 10  $\mu\text{m}$  of the position on the detector, which is lower than the imaging resolution of 2.1%. This deflection effect at the exit of the magnet is thus negligible for a 1 GeV electron beam. For a 1.6 GeV electron beam, this corresponds to a shift of 7.8  $\mu\text{m}$  on the detector, which is thus also negligible.

The uncertainty on the Lanex angle and on the distance between the gas jet and the magnet leads to an error of 1.8% and 0.02% for 1 GeV electrons, respectively, and 3% and 0.03% for 1.6 GeV electrons. By taking into account all the aforementioned uncertainties, we obtain a total error of 5.2% at 1 GeV and 7.9% at 1.6 GeV.

The charge calibration is provided by the tritium capsule enclosed to the detector screen. Assuming an error of 20% on the count of the number of photons from the capsule due to a potential degradation of the luminescent material, in addition to the 19% error on the source calibration performed at ELBE, the total error on the charge measurement is estimated to be 28%.

## APPENDIX B: DENSITY PROFILE

To estimate the density profile of the plasma, we measured the phase shift experienced by the probe beam as it crosses the plasma with a wave front sensor (Phasics SID4 UHR). Since

we did not have the sufficient spatial resolution to resolve the waveguide, the measurement was performed without the channel-forming beam. The plasma is created only by the drive beam, and the density measured is then the background density. The resulting phase maps were then numerically Abel inverted to compute the electron density. We were able to scan the length of the profile with a motorized turning mirror in the vacuum chamber to shift the field of view.

For each position, we were able to extract a lineout of the density profile. The lineouts present a modulated plateau of density. For each position, we then compute the mean across

the plateau and the error bar is the standard deviation on this plateau.

The normalized density gradients are displayed in Fig. 10, for the case without gradient and for the two gradient case discussed in the main paper. The normalization factors correspond to the initial densities indicated in the main text. The dashed lines correspond to a constant fit in the flat case and to linear fits for the gradient cases. The initial densities and the gradient strength are retrieved from the linear fits. In the  $8^\circ$  and  $12^\circ$  cases, the measurements were done with the blade in place, which reduced the acceleration length.

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