

Stochastic acceleration of electrons in fluctuating fields induced by stimulated Raman scattering in subrelativistic laser-plasma interactions

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(Received 14 January 2026; accepted 18 May 2026; published 25 June 2026)

Stimulated Raman scattering (SRS) plays a crucial role in the energy conversion process from lasers to plasmas in the subrelativistic regime, where laser intensities range around 10^{15-16} W/cm². SRS induces excitations of Langmuir waves (LWs) and density fluctuations, such as plasma density cavities. These excitations modify laser absorption in underdense plasma regions and subsequently generate suprathermal hot electrons (HEs) that transport laser energy into overdense plasma regions, changing plasma dynamics, including shock wave formation and propagation. In this paper, we report saturations of SRS-induced HE generation within tens of picoseconds under continuous laser irradiation and a mechanism for HE generation driven by a stochastic process. One-dimensional particle-in-cell simulations with a large density-scale plasma ($\sim 10^2 \mu\text{m}$) over durations of up to 100 ps reveal the clear transition of the electron acceleration mechanism from previously reported electron trapping by SRS-excited LWs to a stochastic acceleration that produces fast electrons with mean energies greater than 100 keV as a result of random walks in fluctuating electric fields generated by SRS-induced density perturbations.

DOI: [10.1103/PhysRevResearch.8.023342](https://doi.org/10.1103/PhysRevResearch.8.023342)

I. INTRODUCTION

In laser-plasma physics, stimulated Raman scattering (SRS) is critically important, especially in the context of laser-driven inertial confinement fusion (ICF), since it affects energy coupling from the laser to plasmas. SRS is a kind of parametric instability (PI) involving three-wave coupling. In the SRS process, a pump electromagnetic wave resonantly decays into a scattered electromagnetic wave and a Langmuir wave (LW), satisfying matching conditions for their wave numbers and frequencies [1]. For the conventional ICF schemes, such as the indirect-drive ICF, actively studied in the National Ignition Facility [2], the mitigation of SRS is one of the major issues because SRS increases the laser reflectivity and reduces the efficiency of energy conversion from the laser to the plasma [3]. Moreover, hot electrons (HEs) generated by SRS (and other PIs as mentioned below) can preheat the target, disturbing its compression [4,5].

In contrast, for shock ignition (SI) [6–8], a more recently featured direct-drive ICF scheme, the generation of HEs could be advantageous to achieve ignition. The SI scheme is characterized by performing both implosion and ignition stages with a single drive laser [9]. In the implosion stage, a laser

with a typical wavelength of $\lambda_L \sim 1 \mu\text{m}$ is irradiated onto the target with its intensity $I_L \sim 10^{14}$ W/cm² for a duration of tens of nanoseconds. In the ignition stage, I_L is increased to $\sim 10^{16}$ W/cm² for a subnanosecond duration, driving an extreme shock with the pressure of subgigabar. This igniter shock would convert its energy to the thermal energy of the compressed core plasma and achieve ignition. In the ignition stage, the laser intensity exceeds the PI's threshold. Thus, the HEs are generated and transport laser energy deep into the plasma. HEs with energies below 100 keV are expected to deposit their energy behind the shock and enhance it [8]. In contrast, HEs with energies above 100 keV, called fast electrons (FEs), would pass the shock and inhibit the compression due to preheating. Therefore, elucidating the characteristics of HEs and understanding their generation mechanisms are crucial for advancing SI and other ICF schemes.

Subrelativistic laser-plasma interactions typically generate HEs with energies of tens of keV via primary PI processes, such as backward SRS (BSRS) and two-plasmon decay instability (TPDI). Here, the subrelativistic regime corresponds to laser intensities of $I_L \lambda_{L(\mu\text{m})}^2 = 10^{15-16}$ W/cm², or to normalized amplitudes of $a_0 \equiv eE/m_e c \omega = \sqrt{2e^2 I_L \lambda_L^2 / (\pi m_e^2 c^5)} \simeq 10^{-2}$ to 10^{-1} , where e , E , m_e , c , and ω are the elementary charge, the electric field amplitude, the electron mass, the light speed, and the wave frequency, respectively. Note that cgs units are used in this paper. These primary PIs excite the large-amplitude LWs, and the excited LWs accelerate the trapped electrons around their phase velocities, which is more than an order of magnitude above the ambient plasma temperature. On the other hand, experimental studies [10–12] have observed

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the generations of FEs. Some acceleration mechanisms have been considered to explain the experimental results. Except for forward SRS (FSRS) [13–15], all the mechanisms require secondary or later (multistep) PI processes, such as rescatter of SRS [16,17], Langmuir decay instability cascade [17], and BSRS-driven TPD [18]. These processes excite large-amplitude LWs with higher phase velocities, which can trap and accelerate electrons beyond 100 keV. Here, we emphasize that these acceleration mechanisms were confirmed in homogeneous density plasmas with $0.1 \lesssim n_e/n_c \leq 0.2$ using kinetic simulations, such as particle-in-cell (PIC) and Vlasov.

In the subrelativistic regime, some studies [19–21] have reported SRS occurrences in the kinetic regime [22,23]. Here, the kinetic regime is defined as the plasma condition satisfying $k\lambda_D \gtrsim 0.3$, where k represents the wave number of LW and $\lambda_D \equiv \sqrt{k_B T_e / (4\pi n_e e^2)}$ is the Debye length, where k_B , T_e , and n_e are the Boltzmann constant, the electron temperature, and the electron density, respectively. This condition can be expressed as $k\lambda_D \simeq v_{th,e}/v_\phi \gtrsim 0.3$, which means that the electrons' thermal velocity $v_{th,e} \simeq \sqrt{k_B T_e / m_e}$ is comparable to or greater than the phase velocity of the LW, $v_\phi \equiv \omega_{pe}/k$, where ω_{pe} is the electron plasma frequency. The conventional fluid theory [1] suggests that Landau damping effectively suppresses the growth of SRS and the excitation of large-amplitude LWs in this plasma condition. In the kinetic regime, electron trapping by LWs affects the Landau damping rate. The electrons trapped in the LW gain the energy from the wave, partially changing the electron distribution function from the Maxwellian distribution. This modification reduces the Landau damping rate compared with that for Maxwellian-distributed electrons [24,25] and triggers the SRS in the kinetic regime [26–28].

Under SI and other direct-drive ICF conditions, the nonlinear frequency shift affects the presence of SRS in the kinetic regime rather than reducing the Landau damping rate [29]. The frequency shift of the LW due to electron trapping [30] enables resonance during its propagation through a plasma with a density gradient, leading to SRS growth in the kinetic regime [31,32]. However, it remains unclear whether SRS in kinetic, inhomogeneous plasmas can trigger multistep PIs and subsequently generate FEs. The previous study related to SI has reported the FEs with temperatures of 135 keV in PIC simulations in relatively high-temperature plasmas spanning a wide density range from $10^{-2}n_c$ to $1n_c$ in tens of picoseconds timescale [33]. Here, $n_c \equiv \pi m_e c^2 / (e\lambda_L)^2$ is the critical density. Experiments have also observed suprathermal electrons up to a few MeV in subnanosecond interactions between subrelativistic lasers and solid targets [10–12]. Nonetheless, the detailed mechanism responsible for the acceleration over 100 keV remains undetermined.

This study investigated the generation process of SRS-driven FEs in SI-relevant plasmas over 100 ps in PIC simulations. The results reveal that after $t \simeq 35$ ps, a stochastic acceleration process replaces wave-based acceleration mechanisms as the dominant pathway for FE generation. At this stage, the SRS-driven large-amplitude LWs no longer exist, and the fluctuating longitudinal electric field caused by SRS-induced density perturbations drives the quasisteady FE generation. This transition in the acceleration mechanism

is considered to be attributed to density inhomogeneity, including cavitation, in long-time laser-plasma interactions. To describe this stochastic acceleration process, we constructed a model based on the diffusion of electrons in energy space. We demonstrated that this model successfully accounts for the production of FEs. The model estimates the mean energies of FEs from the incident wave amplitude and the length of the underdense plasma region, then provides valuable guidance for optimizing laser and plasma conditions in SI. Furthermore, this study elucidates the long-term evolution of SRS from its growth and saturation to a disordered state. Since the growth of PIs and the emergence of turbulent plasmas are also big topics in astrophysics [34], the present results may contribute to a broader understanding of these phenomena. Similar stochastic acceleration processes, where multiple accelerations in field structures can lead to much greater energies than with a single pass, also occur in other plasmas in the laboratory with relevance to astrophysics [35] or in astrophysical plasmas [36,37].

This paper is organized as follows. Section II shows the setup and the results of one-dimensional (1D) collisional PIC simulations. Section III analyzes the results of 1D collisionless simulations and discusses the underlying electron acceleration mechanisms, Sec. III A presents Fourier analysis of electric fields and examines the wave excitations and the SRS saturation, Sec. III B investigates trajectories of FEs and their diffusive energy evolution, and Sec. III C shows electron confinement by density cavities and its effect on the electrons' acceleration. In Sec. IV, we evaluate the order of magnitude of acceleration energy using a random-walk model and assess the feasibility of achieving energy gains exceeding 100 keV through stochastic acceleration under subrelativistic laser-plasma conditions observed in our simulations. Section V summarizes the paper.

II. FAST-ELECTRON GENERATION IN 1D PIC SIMULATIONS

We performed collisional 1D PIC simulations with the PICLS code. PICLS employs a relativistic binary collision model [38], and collisions between all particle species, i.e., e-e, i-i, and e-i (e and i represent electron and ion, respectively), are taken into account in the simulations. In the simulations, the laser light whose electric field is polarized in the y direction and the magnetic field is in the z direction propagates along the x direction from the left boundary of the simulation box. Open boundary conditions are set for electromagnetic fields. The laser wavelength is $\lambda_L = 1 \mu\text{m}$, and the oscillation period is $T_L \simeq 3.3$ fs. The maximum intensity I_L is varied from 10^{15} to 10^{16} W/cm², corresponding to $a_0 \simeq 0.027$ – 0.085 . The amplitude rises from four e-foldings down from its maximum intensity with a 1 ps temporal Gaussian-shaped width. After reaching peak intensity, the laser light maintains the maximum intensity throughout the simulation.

The target consists of fully ionized hydrogen plasma, with an initial maximum density of $1.2n_c$ and a minimum density of $1.2/e_N^4 \simeq 0.02n_c$, assuming an exponential density gradient defined by the density scale length L_{sc1} . Here, $e_N \simeq 2.72$ is Napier's constant. We varied L_{sc1} between 50 and 800 μm . The plasma target has a flat density region of 200 μm behind

the gradient to shut off the laser throughout the simulation. Specifically, the density distribution n_s ($s = e, i$) is expressed as

$$n_s(x) = \begin{cases} 1.2n_c \exp\left[\frac{x-(x_0+4L_{\text{scl}})}{L_{\text{scl}}}\right] & (x_0 \leq x \leq x_0 + 4L_{\text{scl}}), \\ 1.2n_c & (x_0 + 4L_{\text{scl}} \leq x \leq x_1), \\ 0 & (0 \leq x < x_0, x_1 < x \leq x_1 + 500 \mu\text{m}), \end{cases} \quad (1)$$

where x_0 denotes the initial position of the plasma front and is set to $1000 \mu\text{m}$ for all 1D simulations (see Fig. 1), while $x_1 \equiv x_0 + 4L_{\text{scl}} + 200 \mu\text{m}$ denotes that of the plasma end. We put vacuum regions that extend $1000 \mu\text{m}$ from the left boundary ($0 \leq x < x_0$) and $500 \mu\text{m}$ from the right boundary ($x_1 < x \leq x_1 + 500 \mu\text{m}$). Following the previous research relevant to SI [39], the initial temperatures of ions and electrons are set to $T_i = 1 \text{ keV}$ and $T_e = 5 \text{ keV}$, respectively. The typical spatial mesh size is $0.05 \mu\text{m}$ (λ_L divided by 20), with a corresponding time step of 0.17 fs (T_L divided by 20) determined by the PICLS algorithm. The number of PIC particles is 70 per mesh for hydrogen ions (protons) and electrons, respectively. The fourth-order shape function is used for particles.

As mentioned above, the vacuum regions extend from the left boundary (laser incident side) and from the right boundary. The areas extending from both boundaries are damping regions, where particles decelerate to their initial temperatures upon their arrival at the boundary. These damping regions prevent the generated HEs from endlessly circulating in the simulation box. These boundary conditions mimic the charge neutralization by thermal electrons. In reality, generated HEs disperse to the surroundings, and a stream of thermal electrons replaces them to maintain charge neutrality. The large vacuum regions work to prevent the accumulation of decelerated particles near both boundaries, which would interfere with laser injection, in the simulation of 100 ps .

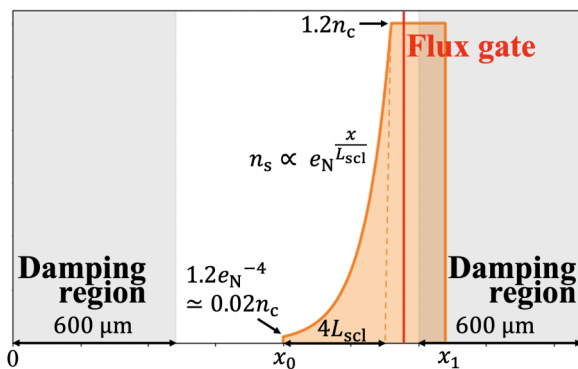


FIG. 1. Initial spatial profile of 1D simulation. Along the x axis, the fully ionized hydrogen plasma is set between x_0 and x_1 , where $x_0 = 1000 \mu\text{m}$ and $x_1 \equiv x_0 + 4L_{\text{scl}} + 200 \mu\text{m}$ for all 1D simulations. The density scale length L_{scl} determines how the density falls from $1.2n_c$ to $1.2e_N^{-4} \approx 0.02n_c$, where $e_N \approx 2.72$ is Napier's constant. The total simulation box size is set to $x_1 + 500 \mu\text{m}$. The gray-shaded areas extending $600 \mu\text{m}$ from both boundaries are damping regions. The flux gate is set at $x = x_0 + 4L_{\text{scl}} + 50 \mu\text{m}$ and records the energy spectrum of electrons passing through it from left to right.

First of all, we show the HE flux observed in the simulations. In the simulation box, we set the flux gate at $x = x_0 + 4L_{\text{scl}} + 50 \mu\text{m}$, which is in the flat density region. The energy spectrum of electrons passing through the gate from left to right is measured as an integrated spectrum every 1 ps . Figure 2(a) shows the three examples of the flux spectrum f_{flux} observed from $t = 49$ to 50 ps . The blue, orange, and green dashed lines depict the spectra in the simulations with $I_L = 10^{14}, 10^{15}, 10^{16} \text{ W/cm}^2$, respectively. In these simulations, the initial density scale length is set to $L_{\text{scl}} = 100 \mu\text{m}$.

The solid gray lines show the fittings for the flux spectra. The bulk components of the fluxes are successfully fitted with

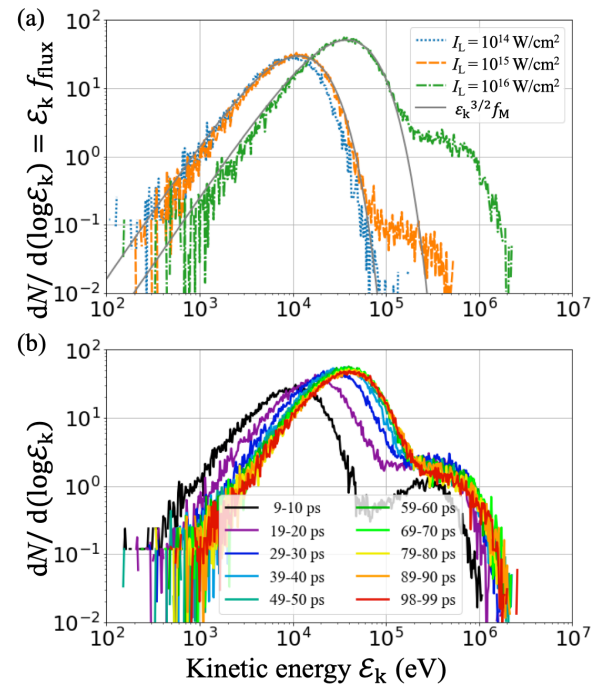


FIG. 2. (a) Energy spectra of electron fluxes observed at the gates during one picosecond from $t = 49$ to 50 ps . The blue, orange, and green dashed lines show the spectra in the 1D collisional simulations with the laser intensity $I_L = 10^{14}, 10^{15}$, and 10^{16} W/cm^2 , respectively. The gray solid lines are fittings with a fitting function $f_{\text{fit}} \equiv \mathcal{E}_k^{3/2} f_M$ as Eq. (2), where f_M and $\mathcal{E}_k$ denote the Maxwellian distribution and the kinetic energy of electrons, respectively. Here, electron temperatures $T_e = 6$ and 19 keV give the best fittings for the orange and the green lines, respectively. In all simulations, the laser wavelength is $\lambda_L = 1 \mu\text{m}$, and the initial plasma density scale length is $L_{\text{scl}} = 100 \mu\text{m}$. (b) Time evolution of energy spectrum observed in the simulation with $I_L = 10^{16} \text{ W/cm}^2$.

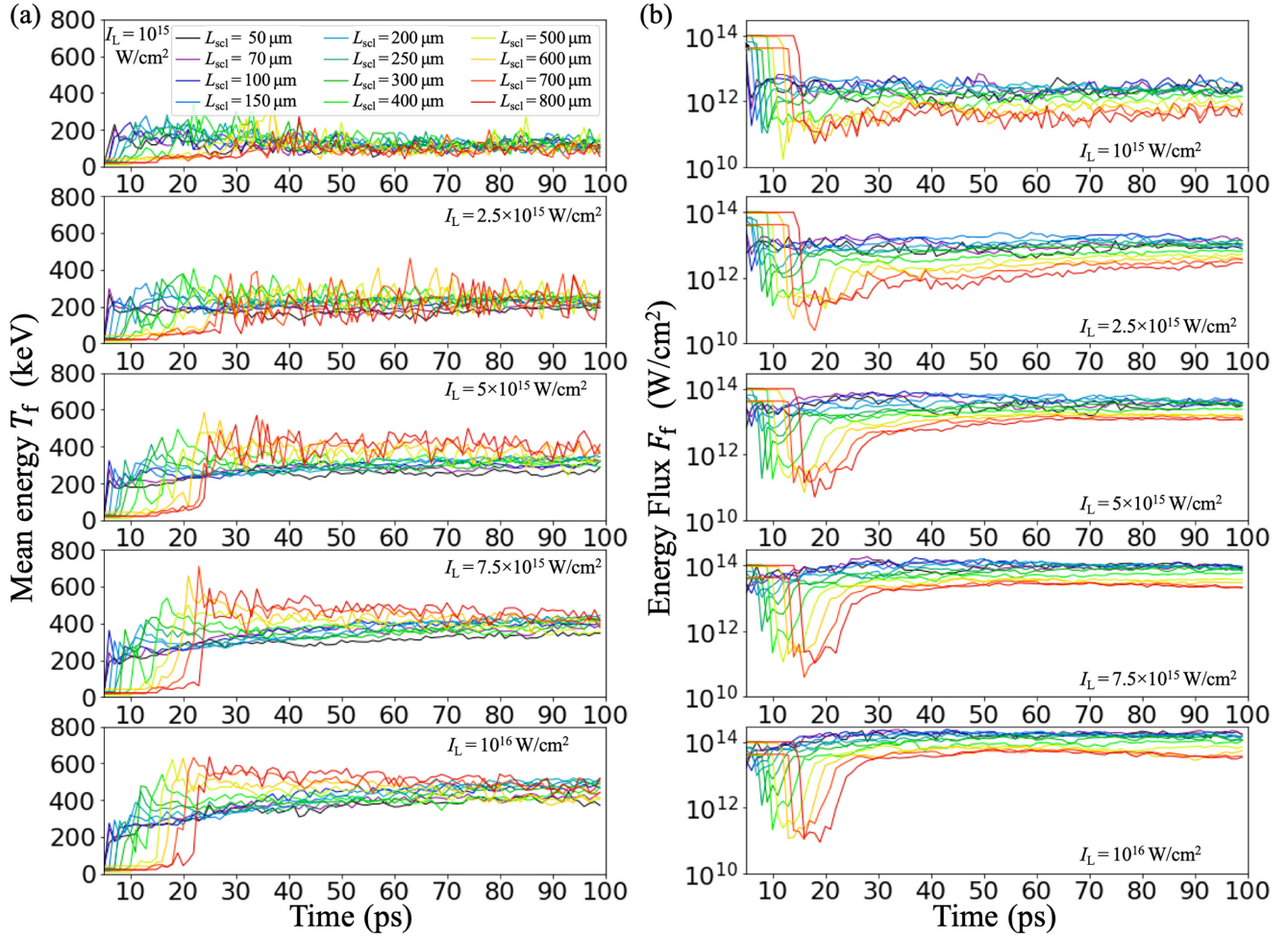


FIG. 3. Time evolution of (a) the mean energy T_f and (b) the energy flux F_f of suprathermal (fast) electrons. Both T_f and F_f approach quasisteady values after $t \simeq 35$ ps. The slower rise of the warm-colored lines reflects the longer propagation distance that fast electrons travel until reaching the flux gate. In panel (b), the plateau and the subsequent steep drop observed before $t = 20$ ps are numerical artifacts arising from the fitting method applied to spectra with no or very few electron numbers.

the fitting function f_{fit} [40]:

$$f_{\text{fit}} = 2\pi \mathcal{E}_k^2 \left(\frac{1}{\pi k_B T_e} \right)^{3/2} \exp \left(\frac{-\mathcal{E}_k}{k_B T_e} \right) = \mathcal{E}_k^{3/2} f_M, \quad (2)$$

where

$$f_M \equiv 2\pi \mathcal{E}_k^{1/2} \left(\frac{1}{\pi k_B T_e} \right)^{3/2} \exp \left(\frac{-\mathcal{E}_k}{k_B T_e} \right)$$

is three-dimensional Maxwellian distribution and $\mathcal{E}_k$ denotes the kinetic energy of electrons. Of the coefficient $\mathcal{E}_k^{3/2}$, $\mathcal{E}_k$ originates from setting energy bins logarithmically, while $\mathcal{E}_k^{1/2} \propto v_e$ (electrons velocity) originates from time-integrated spectra. That is, electrons with higher velocities reach the flux gate more frequently within 1 ps. Now, electron temperatures $T_e = 6$ and 19 keV give the best fittings for the orange and the green lines, respectively.

The orange and green spectra have suprathermal components with energies above 100 keV. Figure 2(b) shows the time evolution of the flux spectrum observed in the simulation with $I_L = 10^{16}$ W/cm² and $L_{\text{scl}} = 100$ μm, corresponding to the case of the green line in panel (a). In this simulation, the

bulk thermal component, which can be fitted by f_{fit} , shifts toward the high-energy side due to the collisional heating. In contrast, the suprathermal component reaches a steady-state distribution shape by $t \simeq 20$ ps.

We separated the thermal and suprathermal (FE) components by fitting with Eq. (2) and investigated the behavior of the suprathermal one. Figure 3 shows the time evolution of (a) the mean energy T_f and (b) the energy flux F_f of the suprathermal components. The five graphs in panels (a) and (b), from top to bottom, show the results of simulations with $I_L = 1, 2.5, 5, 7.5 \times 10^{15}$, and 10^{16} W/cm². The line colors indicate the difference of L_{scl} . In all 60 simulations, T_f stabilizes approximately constant after $t \simeq 35$ ps, and its magnitude depends on I_L . The timing of the rise of T_f reflects the lag between the start of the simulation and the arrival of the FEs at the gate. A slower arrival is observed for larger L_{scl} , where the gate is positioned farther from the region where laser-plasma interactions occur. Like T_f , F_f becomes constant after $t \simeq 35$ ps and its magnitude also strongly depends on I_L .

Figure 4 shows the time-averaged (a) T_f and (b) F_f from 35 to 100 ps. The averaged mean energy T_f is proportional

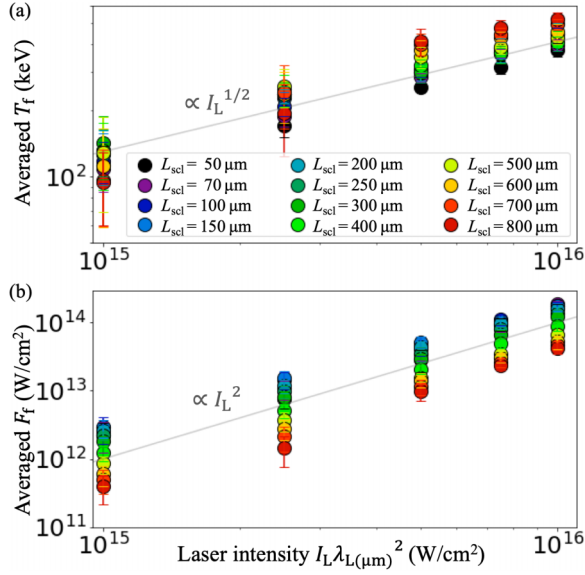


FIG. 4. Time-averaged mean energy T_f [panel (a)] and energy flux F_f [panel (b)] taken from $t = 35$ to 100 ps.

to $I_L^{1/2}$ and the energy flux F_f to I_L^2 . As shown in Fig. 2(a), for $I_L = 10^{14}$ W/cm² (blue line), no clear suprathermal component can be seen. It suggests that $I_L \gtrsim 10^{15}$ W/cm² is the condition for SRS growth. Indeed, the density cavitation by SRS, which will be discussed in Sec. III C, was not observed in the simulations with $I_L < 10^{15}$ W/cm². Notably, this result is consistent with the previous study [29], showing a threshold laser intensity $\simeq 1.4 \times 10^{15}$ W/cm² for SRS in the kinetic regime.

III. ANALYSIS OF ACCELERATION MECHANISM

To clarify the physical mechanism of electron acceleration beyond 35 ps, we conducted a detailed investigation of the electromagnetic fields and the motion of FEs in these fields. To suppress the effects of field noise, the following analyses used collisionless simulations. As shown below, the generation of FEs is observed in collisionless simulations as well as in the collisional ones. The effect of collisions on the FE-generation process will be discussed in Sec. V.

A. Spectrum of electric field

We checked the modes of the transverse electric field E_y and the longitudinal electric field E_x in underdense plasmas. The SRS growth rate depends on the density scale length [41], with longer L_{scl} being more favorable for its growth. Then, we chose $L_{scl} = 600$ μ m and investigated mode excitations around $t = 35$ ps in the region of initial electron density $n_e/n_c = 0.1$ – 0.2 , where the LW excitations and the related FE generation by wave trapping have been reported [13–18]. Note that BSRS-driven TPD [18] is a multidimensional PI and cannot occur in 1D simulations. To accelerate SRS growth and mitigate the noise, the laser amplitude is set to a relatively high $a_0 = 0.085$ ($I_L \lambda_{L(\mu m)}^2 = 10^{16}$ W/cm²), and the spatial mesh size and corresponding time step are changed to $\lambda_L/100$ and $T_L/100$, respectively, offering five times finer resolution.

Panels (a), (b), and (c) in Fig. 5 show the spatial distribution of electron density n_e , the E_y spectrum, and the E_x spectrum at $t = 16.7$ ps, respectively. At this time, SRS is beginning to produce density cavities in this region, e.g., at $x = 2040$ and 2160 μ m, as shown in panel (a). The amplitude of E_x field is roughly 10% of the laser field E_y . Panels (d) and (e) show the spatiotemporal Fourier spectra of electromagnetic waves and LWs. They are calculated by doing two-dimensional (2D) fast Fourier transforms (FFTs) of E_y and E_x fields. The field data are recorded from $t = 16.7$ ps = $5000 T_L$ to 20.0 ps = $6000 T_L$, with 10 000 samples taken per $T_L/10$. The data range in real space is from $x = 1910$ to 2575 μ m, shown in panels (a)–(c), with 2075 samples taken per $\lambda_L/20$. We can observe strong forward and backward SRS signals indicated by the white dashed circles in panel (d). The E_x spectrum [panel (e)] reveals the corresponding BSRS and its rescatter signals. These primary (FSRS) and secondary (rescatter of BSRS) SRS modes are also observed in the previous studies [16,17].

At $t = 36.7$ ps, the electron distribution and the Fourier spectra contrarily have changed, as shown in panels (f)–(j). For doing 2D FFTs, the field data are recorded from $t = 36.7$ ps = $11\,000 T_L$ to 40.0 ps = $12\,000 T_L$. In panel (f), the cavities have fully grown by this time. As a result of cavity growth, electromagnetic field enhancement occurs [43], as shown in panel (g). Inside the cavities, the amplitude E_{cavity} reaches approximately three times that of the incident laser field, while its frequency ω_{cavity} is down to $\omega_{cavity}/\omega_L \simeq 1/4$, where ω_L is the laser frequency. In addition, panel (h) shows a slight increase in the amplitude of the E_x field compared to (c). In panel (i), the backward and forward SRS modes observed earlier in panel (d) have completely disappeared. Likewise, the BSRS and its rescatter signals are no longer visible in the E_x spectrum, as shown in panel (j).

Figure 6 shows the Fourier spectra of (a) E_y and (b) E_x in the whole density region, where the plasma density initially ranges from $\simeq 0.02$ (corresponding to $x = x_0 = 1000$ μ m) to $1.2n_c$ ($x = 3400$ μ m). The data are recorded from $t = 40.0$ ps = $12\,000 T_L$ to 41.0 ps = $12\,300 T_L$. In panel (a), strong signals appear at $\omega/\omega_L = 1$, associated with the incident and reflected laser lights. Striped signals caused by cavities are visible around the electromagnetic dispersion relation curves for $n_e/n_c = 0.1$ and 0.02 . Assuming the spatial shape of the cavity $f(x)$ as a Gaussian $f(x) \propto \exp[-x^2/(2\sigma^2)]$, the signal in wave number space given by its 1D Fourier transform also follows a Gaussian $\tilde{f}(k) \propto \exp(-\sigma^2 k^2/2)$. Then, the stripe extends over a certain range along the k axis, with a characteristic width $\propto 1/\sigma$, depending on the cavity size $\simeq 2\sqrt{2\ln 2}\sigma$ (now $\sigma/\lambda_L \gtrsim 1$). The heights of the stripes in k - ω space correspond to the frequencies of the electromagnetic fields in the cavities, typically $\omega_{cavity}/\omega_L \simeq 1/4$ at $n_e/n_c = 0.1$ and $\omega_{cavity}/\omega_L \simeq 1/8$ at $n_e/n_c = 0.02$. Aside from these signals, no additional significant modes are observed, suggesting an absence of notable SRS growth after $t = 40$ ps. In panel (b), similar to panel (a), no outstanding LW mode excited by SRS is present, and modes corresponding to electron densities between $n_e/n_c = 0.001$ and 1.2 are universally observed.

The comparison in Fig. 5 could suggest the presence of an electron acceleration mechanism instead of wave-based

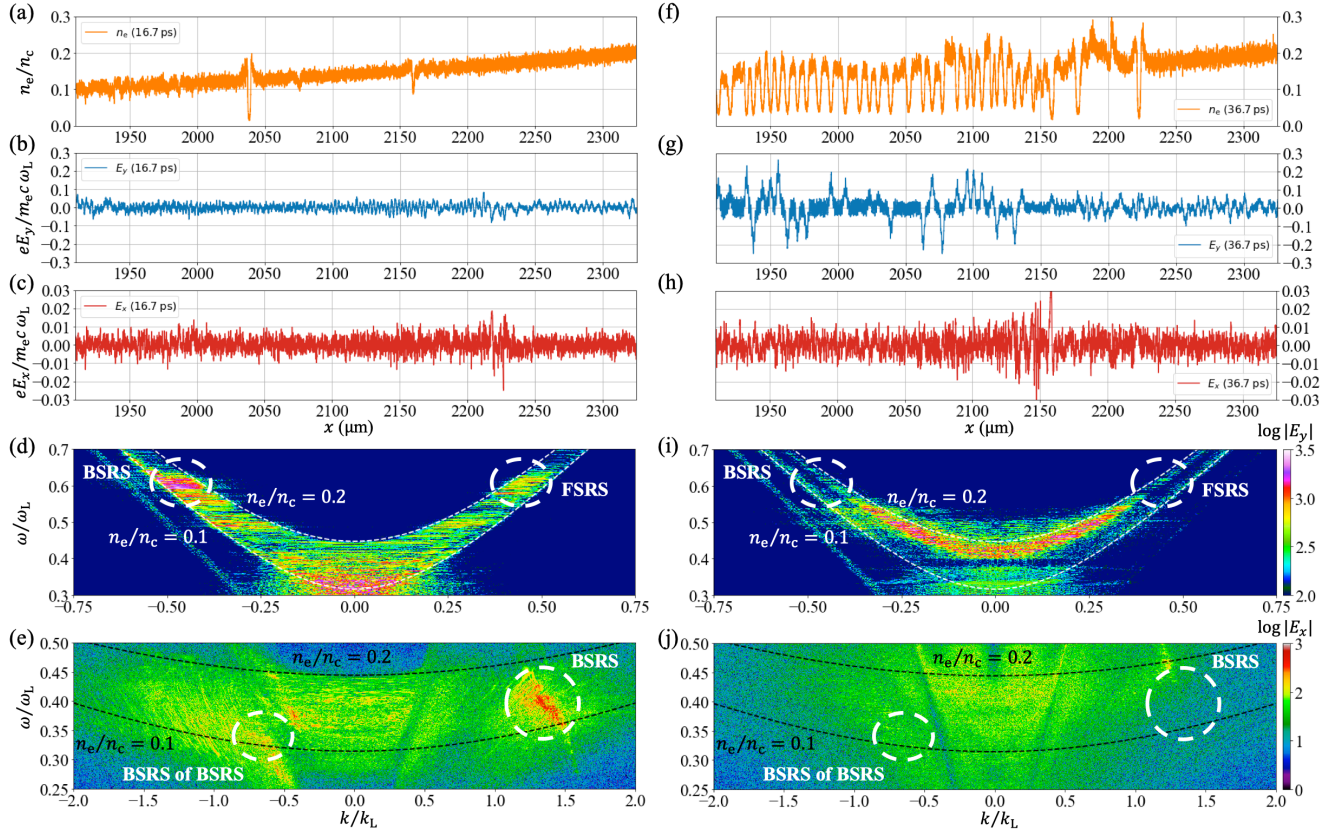


FIG. 5. Spatial distribution of (a) electron density n_e , (b) transverse (laser) electric wave E_y , and (c) longitudinal (Langmuir) wave E_x at $t = 16.7$ ps $= 5000 T_L$. Corresponding spatialtemporal Fourier spectra of (d) E_y and (e) E_x . The data are obtained from the collisionless simulation with $a_0 = 0.085$ ($I_L \lambda_{L(\mu\text{m})} = 10^{16}$ W/cm ^{2}) and $L_{\text{scf}} = 600 \mu\text{m}$. The frequency and the wave number are normalized by the incident laser field frequency ω_L and wave number $k_L \equiv \omega_L/c$, respectively. The color contour shows the intensity of modes on a logarithmic scale with base 10. In panel (d), white dashed lines represent the dispersion relation curve of electromagnetic waves, given by $\Omega^2 = \Omega_{\text{pe}}^2 + K^2$, for $n_e/n_c = 0.2$ (top) and 0.1 (bottom). The white dashed circles indicate the positions of electromagnetic wave modes excited by forward SRS and backward SRS in this density region. Here, $\Omega \equiv \omega/\omega_L$, $\Omega_{\text{pe}} \equiv \omega_{\text{pe}}/\omega_L$, and $K \equiv k/k_L$. In panel (e), black dashed lines represent the weakly relativistic dispersion relation curve of Langmuir waves [42], given by $\Omega^2 = \Omega_{\text{pe}}^2 [1 - 5/(2M)] + 3K^2/M$, for $(n_e, T_e) = (0.2n_c, 2.5 \text{ keV})$ (top) and $(0.1n_c, 2.5 \text{ keV})$ (bottom). Here, $M \equiv m_e c^2/k_B T_e$. The white dashed circles indicate the positions of Langmuir wave modes excited by BSRS and their rescatter into BSRS. (f)–(j) The same analyses at $t = 36.7$ ps $= 11000 T_L$.

mechanisms driven by SRS. At the early time of the simulations, the mechanisms identified in previous studies, such as FSRS [13–15] and rescatter of SRS [16,17], may govern the electron acceleration process, effectively explaining the generation of FEs before $t = 35$ ps. However, at a later time, there is no clear evidence of a dominant excitation of SRS-driven LWs. Even when we extended the FFT region beyond the range between $0.1n_c$ and $0.2n_c$, no signature of large-amplitude LWs appears as shown in Fig. 6(b). Instead, uniform LW modes corresponding to all possible plasma conditions ($0.001 \lesssim n_e/n_c \lesssim 1.2$) are observed across a broad range in k - ω space. Even at the later stage, we observe the efficient electron acceleration. Therefore, an alternative acceleration mechanism is required to explain the steady generation of FEs observed beyond $t = 100$ ps. The following section focuses on analyzing the FEs trajectories and examining their acceleration process.

For FSRS, previous studies [13–15] confirmed mode excitations up to $t/T_L = 10^2$ to 3×10^3 , which is not contradictory with the mode disappearance observed from $t = 36.7$ ps $=$

$11000 T_L$. While the rescatter of BSRS has been reported even at $t/T_L \sim 10^4$ [17]. One possible explanation for this difference is the huge density inhomogeneity caused by cavities. In the previous study, drive lasers have relatively low amplitudes and appear to maintain a certain degree of density uniformity throughout the simulations. To clarify this, we conducted an additional simulation using fixed ions and confirmed the generation of FEs. In the spectra of the electric field, BSRS signal between $n_e/n_c = 0.1$ and 0.2 appears in both the E_y and E_x spectra at $t = 36.7$ ps. Additional signals suggesting the rescatter of BSRS into BSRS also appear in the E_x spectrum. This result suggests that cavitation may suppress SRS growth in the kinetic regime by creating substantial plasma density inhomogeneities and inhibiting the frequency shift associated with the propagation of LWs in density-gradient plasmas. However, fixed ions introduce unique laser-plasma dynamics, such as the formation of a steep sheath field at the ion front and the strong decay of the laser field at $n_e/n_c = 1$. These phenomena can change the electrons' acceleration process, and in fact, the FE behavior differs from that in simulations

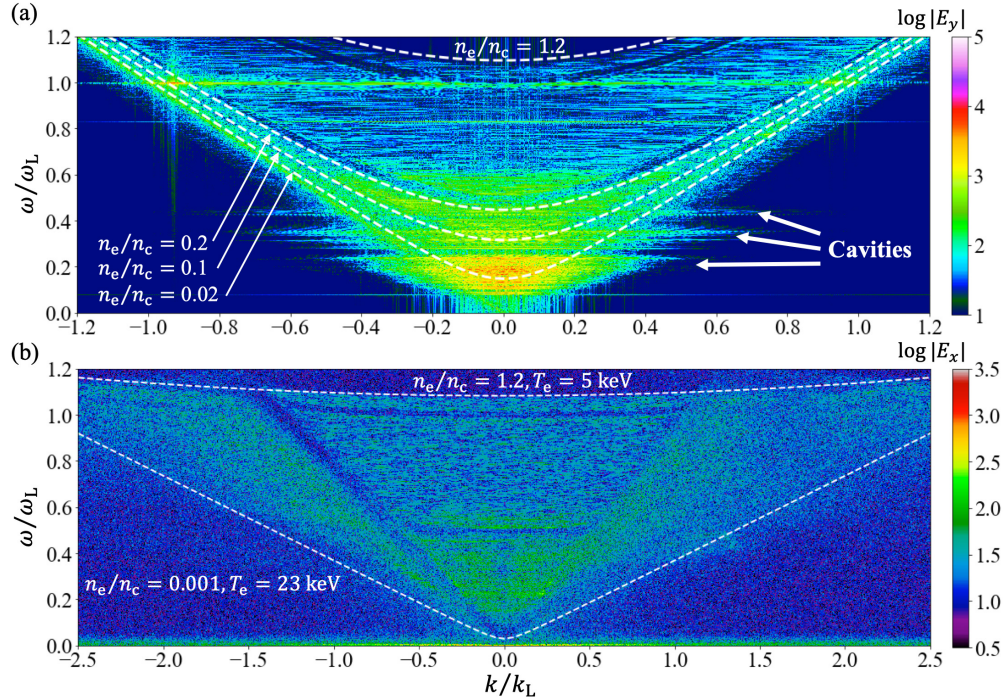


FIG. 6. Fourier spectra of (a) E_y and (b) E_x at $t = 40 \text{ ps} = 12\,000 T_L$ from the whole density region. The density range is initially from $\simeq 0.02$ (corresponding to $x = x_0 = 1000 \mu\text{m}$) to $1.2n_c$ ($x = 3400 \mu\text{m}$). In panel (a), white dashed lines represent the dispersion relation curve of electromagnetic waves with $n_e/n_c = 1.2, 0.2, 0.1$, and 0.02 , respectively. The striped signals arise from density cavities. In panel (b), white dashed lines represent the weakly relativistic dispersion relation curve of Langmuir waves with $(n_e, T_e) = (1.2n_c, 2.5 \text{ keV})$ (top) and $(0.001n_c, 23 \text{ keV})$ (bottom).

with mobile ions. Now, it is challenging to determine the dominant factor affecting SRS growth, necessitating further investigation.

B. Analysis of fast electron's trajectory

To identify the alternative acceleration mechanism, the motion of FEs was investigated in detail. We tracked the trajectories of FEs, from their initial positions to the gate, while recording their momenta and the amplitudes of the electromagnetic fields experienced along their trajectories at each time step. Figure 7(a) shows some examples of FEs' trajectories. From the simulation with $a_0 = 0.085$ and $L_{\text{scf}} = 600 \mu\text{m}$, we selected 25 FEs that reached the gate [its position is indicated by the red solid line in panel (a)] with energies over 460 keV ($=T_f$) at $t = 40 \text{ ps}$. The line colors indicate the electron's energy in base 10 logarithms. The background color contour is the amplitude of the E_x field. FEs repeatedly accelerate and decelerate over tens of ps before finally reaching the gate. Panel (b) shows the change in momentum of an FE, whose trajectory is indicated by a black arrow in panel (a), as an example. Through the interaction with the laser field E_y , p_y rapidly oscillates until $t = 35 \text{ ps}$, with an amplitude comparable to p_x . Compared to p_y , the change in p_x shows a slow oscillation before $t = 35 \text{ ps}$. Regarding p_z , since there is no force acting in the z direction in this collisionless situation, it maintains its initial thermal velocity. Upon reaching the gate at $t = 40 \text{ ps}$, p_y goes with an initial thermal velocity, similar to p_z . Further details are discussed in the next section.

Panel (c) is an enlarged view of the region indicated by the black circle in panels (a) and (b), showing a part of the final acceleration stage. In panel (c), the E_x field appears as a fluctuating field and does not trap the electron. Panel (d) shows the time evolutions of the E_x amplitude along this trajectory and the longitudinal momentum p_x . This electron experienced both positive and negative E_x fields, and their balance tilted toward the negative, resulting in an acceleration tendency. We examined the trajectories of various FEs under varying laser intensities, density scale lengths, and times. In all cases, electrons seem to undergo random accelerations and decelerations as they pass through the field.

To understand the typical behavior of FEs, we randomly sampled approximately 10 000 FEs with energies over 460 keV at $t = 39.36 \text{ ps} = 11\,800 T_L$ from the simulation. To avoid the damping region, we chose the sampling region from $x = 1000$ to $2350 \mu\text{m}$, which corresponds to the positions of the initial plasma front (x_0) and where the initial density is $n_e/n_c = 1$, respectively. Figure 8 shows the energy change of the FEs from $t'/T_L = 0$ to 200, where $t' = 0$ corresponds to $t = 39.36 \text{ ps}$ in the simulation. The x axis represents the difference in kinetic energy of FEs between $t' = 0$ and a later time t' , while the y axis indicates the normalized number of FEs. Thus, at $t' = 0$ (red lines), the distribution is discretized to $x = 0$, with its value normalized to $y = 1$ (not visible within the figure's range). The green, orange, and blue dashed lines depict the distributions at $t'/T_L = 50, 100$, and 200 , covering up to $t = 40 \text{ ps}$ in the simulation. The distribution seems to evolve diffusively, i.e., spreading symmetrically with a larger deviation at a later t' . This result suggests that FEs undergo

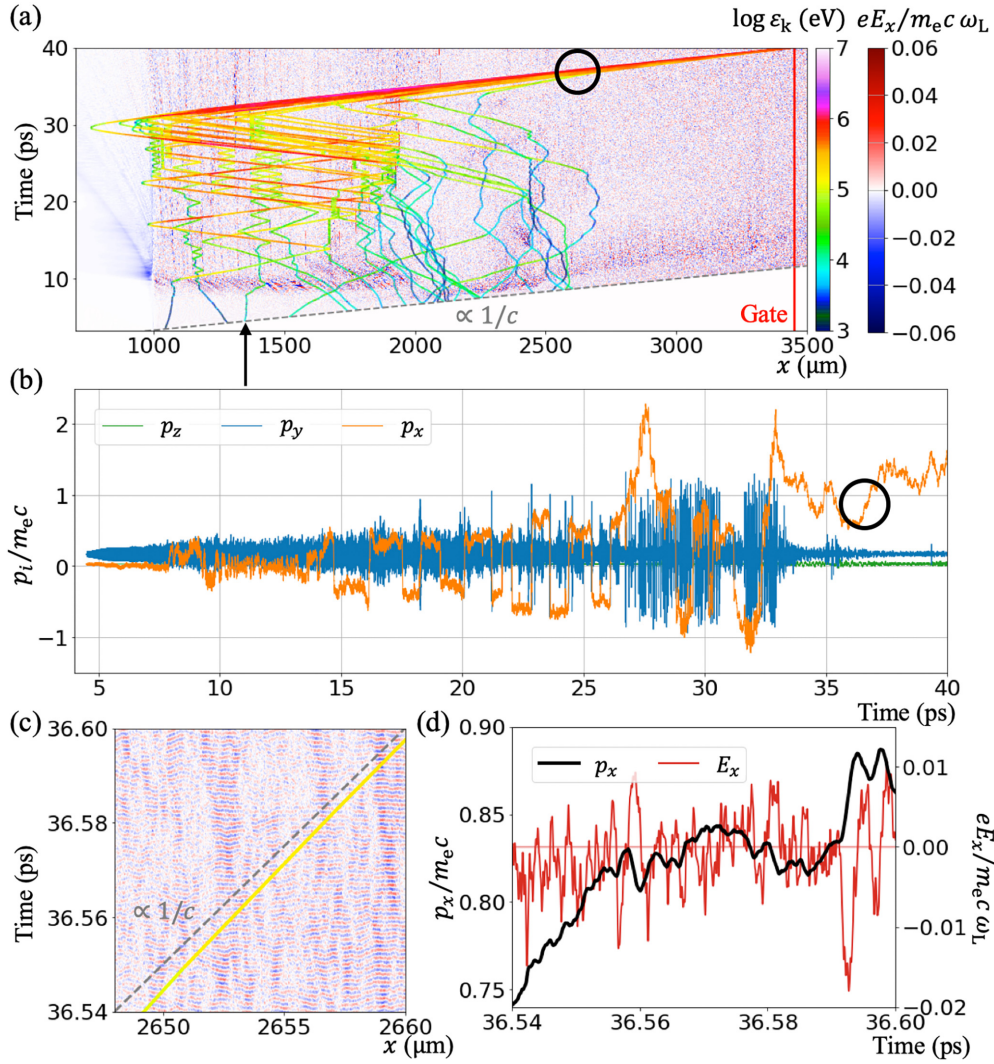


FIG. 7. (a) Trajectories of fast electrons observed from the collisionless simulation with $a_0 = 0.085$ and $L_{\text{scf}} = 600 \mu\text{m}$. (b) Time evolution of momenta of a fast electron, where its trajectory is indicated by the black arrow in panel (a). Enlarged views of (c) trajectory and (d) time evolution of p_x and E_x along the trajectory within the range indicated by the black circles in panels (a) and (b).

stochastic acceleration. Some amount of FEs increase their energy in the fluctuating field, while comparable numbers of FEs lose a similar amount of energy.

The solid lines are fittings with the diffusion equation f_D :

$$f_D = \frac{f}{\sqrt{4\pi Dt'}} \exp \left[-\frac{(\mathcal{E}_k(t') - \mathcal{E}_k(0) - \mu)^2}{4Dt'} \right]. \quad (3)$$

Here, Dt' (the product of the diffusion coefficient D and time t'), f (the normalization factor), and μ (the mean) are treated as fitting parameters. The green, orange, and blue solid lines are given with parameter sets

$$\begin{aligned} (Dt'(\text{keV}^2), f, \mu(\text{keV})) \\ = (8.7 \times 10^3, 1.9 \times 10^2, -2.5 \times 10^1), \\ (1.5 \times 10^4, 1.9 \times 10^2, -4.8 \times 10^1), \\ (2.4 \times 10^4, 1.9 \times 10^2, -9.0 \times 10^1), \end{aligned}$$

respectively. The square root of Dt' corresponds to the averaged gain of energy, whose order of magnitude $\sim 10^2$ keV is consistent after $t'/T_L = 60$ with the mean energy T_f . The time evolution of the mean μ indicates a gradual shift of the distribution center toward the low-energy side at a rate of $\mu/t' \simeq 0.5 \text{ keV}/T_L$, which is negligible compared with that of Dt' .

C. Electron confinement between density cavities

We found electron confinement between cavities. As shown in Fig. 7(a), FEs undergo two distinct forms of stochastic motions: wide-range excursions over hundreds of micrometers and confined oscillatory motion within tens of micrometers. Figure 9(a) shows the electrons trajectories in the region undergoing cavity formation. The grayscale on the background shows the electron density n_e , and the white streaks indicate low-density regions, i.e., cavities. The colored lines show the electrons trajectories, which sometimes

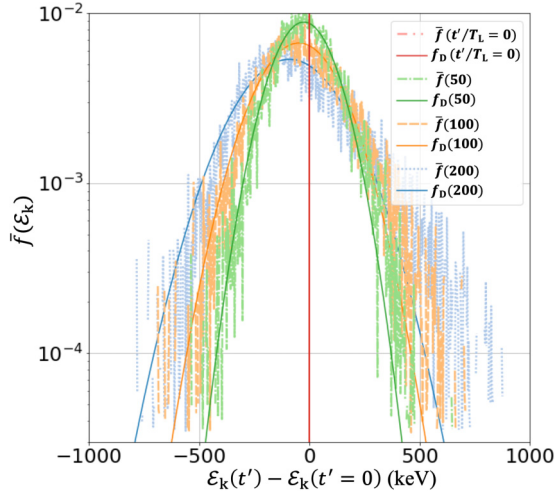


FIG. 8. Distribution of the kinetic energy change of fast electrons. Four dashed lines show the time evolution of the distribution from $t'/T_L = 0$ to 200, corresponding to $t = 39.36\text{--}40$ ps in the collisionless simulation with $a_0 = 0.085$ and $L_{\text{scl}} = 600\text{ }\mu\text{m}$. Four solid lines are the fittings with the diffusion equation f_D as Eq. (3).

exhibit confined motion between cavities. The confinement of electrons is also apparent in the phase space (b). At the positions of cavities, electrons are reflected, reversing their longitudinal momentum p_x from negative to positive and vice versa. The maximum amplitude of the momentum of the confined electrons is roughly $\Delta p_x/m_e c \simeq 1.5$. Electrons with

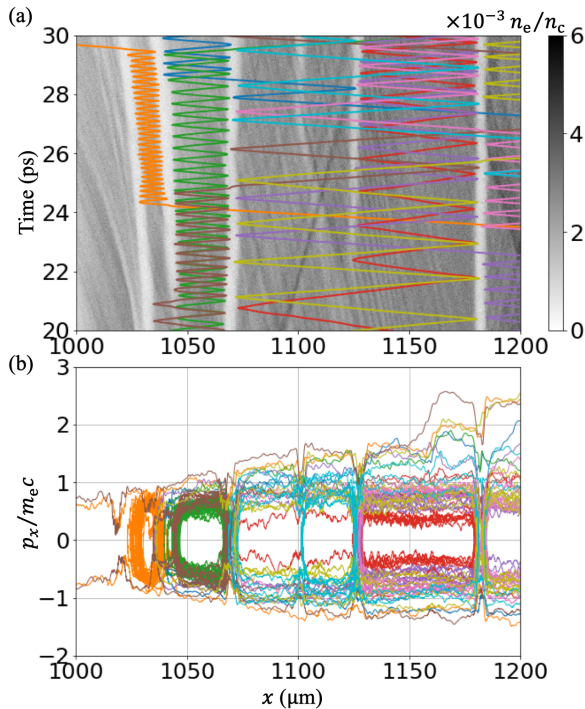


FIG. 9. Trajectories of fast electrons in the cavitation region. (a) Time evolution of positions of cavities (white streaks) and the picked-up fast electrons (solid lines). Each color corresponds to each fast electron. (b) Motion of fast electrons in phase space. The line colors correspond to those in panel (a).

$p_x/m_e c > 1.5$ can surpass the cavities ponderomotive potentials ϕ_{cavity} and escape from the cavitation region.

The potential ϕ_{cavity} can be estimated as [44]

$$\begin{aligned}\phi_{\text{cavity}} &= \frac{e^2}{m_e \omega_{\text{cavity}}^2} \left\langle \frac{E_{\text{cavity}}^2}{2} \right\rangle \\ &\simeq \frac{e^2}{m_e \left(\frac{\omega_L}{8}\right)^2} \times \frac{1}{4} \left(\frac{3a_0 m_e c \omega_L}{e} \right)^2 \\ &= 144 a_0^2 m_e c^2.\end{aligned}$$

Here, the angle brackets $\langle \cdot \rangle$ denote the time-averaged value over one period of the electromagnetic wave, which turns to $1/2$ for a sinusoidal wave. We adopt empirical values obtained from the simulations, namely, $\omega_{\text{cavity}} \simeq \omega_L/8$ and $E_{\text{cavity}} \simeq 3a_0$ (cf. Fig. 5). Now, $a_0 = 0.085$. The confined momentum $\Delta p_x/m_e c \simeq 1.5$ corresponds to the confined electrons' kinetic energy $\mathcal{E}_k/m_e c^2 \simeq 0.8$, which is consistent with the estimate of cavity potential $\phi_{\text{cavity}} \simeq m_e c^2$.

As an electron crosses a cavity, it may gain kinetic energy due to the difference of the oscillating potential amplitude $\Delta\phi_{\text{cavity}}$ between its entry and exit, and $\Delta\phi_{\text{cavity}} \sim \phi_{\text{cavity}}$. Since the phase of the potential oscillation is random for electrons that enter the cavity, the acceleration can be regarded as a stochastic process. Cavities exist in a limited region where $n_e < n_c/4$, yet the amplitude of their electromagnetic field E_{cavity} is tens of times larger than that of the E_x field, enabling them to produce non-negligible energy gains for FEs. Then, the electron acceleration can be driven by two stochasticities: the fluctuating electric field in the x direction, $\sim E_x dx$, and the deviation of cavities potentials, $\sim \Delta\phi_{\text{cavity}}$.

To clarify which stochasticity dominates the acceleration process of FEs, we looked into the acceleration histories of FEs. In collisionless situations, the energy change of an electron can be written as

$$m_e c^2 \frac{d}{dt}(\gamma - 1) = -e \left(\mathbf{E} + \frac{\mathbf{v} \times \mathbf{B}}{c} \right) \cdot \mathbf{v}, \quad (4)$$

where $\gamma = \sqrt{1 + (\mathbf{p}/m_e c)^2}$ is the Lorentz factor. Then, the total kinetic energy gain of an electron $\Delta\mathcal{E}_{k,\text{tot}}$ is given by the integral of Eq. (4) over time:

$$\begin{aligned}\Delta\mathcal{E}_{k,\text{tot}} &= \int_0^t m_e c^2 \frac{d}{dt}(\gamma - 1) dt \\ &= \int_0^t -e \left(\mathbf{E} + \frac{\mathbf{v} \times \mathbf{B}}{c} \right) \cdot \mathbf{v} dt \\ &= \int_{\text{initial position}}^{\text{gate}} -e \mathbf{E} \cdot d\mathbf{r}.\end{aligned} \quad (5)$$

This equation means that while the electric field $\mathbf{E}$ contributes directly to the energy gain of the electron, the magnetic field $\mathbf{B}$ does not perform any work, solely changing the direction of the motion. Using the recorded field data and employing Eq. (5), the energy changes of the fast electrons were calculated separately for each spatial direction (x , y , and z).

Figure 10 compares the energy changes of FEs. Panels (a)–(c) show the results for electrons with energies over 300 keV reaching the gate at $t = 40$ ps in simulations with $a_0 = 0.085$ and $L_{\text{scl}} = 100$ and $600\text{ }\mu\text{m}$. As shown in panels

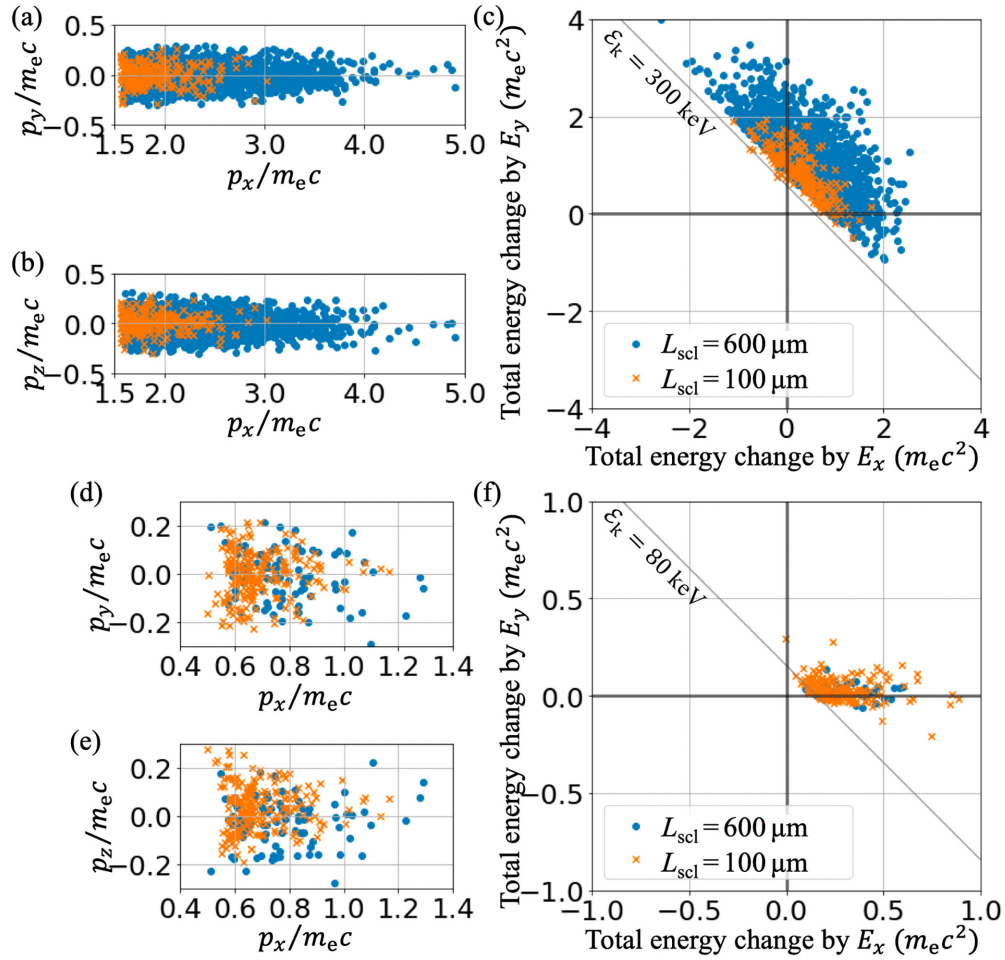


FIG. 10. Comparison of momenta (a) p_x - p_y and (b) p_x - p_z . (c) Total energy changes of fast electrons. In both simulations with $L_{\text{scl}} = 100 \mu\text{m}$ (blue) and $L_{\text{scl}} = 600 \mu\text{m}$ (orange), the laser amplitude is set to $a_0 = 0.085$. (d)–(f) The same analyses for suprathermal electrons observed in the simulations with $a_0 = 0.027$ ($I_L \lambda_{L(\mu\text{m})} = 10^{15} \text{ W/cm}^2$).

(a) and (b), the FEs have significantly greater p_x compared to their transverse momenta p_y and p_z . The lower limit of $p_x/m_e c$ is approximately 1.5, consistent with the confinement amplitude induced by cavities, $\Delta p_x/m_e c \simeq 1.5$. In contrast, $p_y/m_e c$ and $p_z/m_e c$ are less than 0.3, corresponding to a thermal velocity of $\simeq 20 \text{ keV}$, as shown in Fig. 7(b). It means that the acceleration along the y direction by E_y in the cavity is converted into the kinetic energy gain along the x direction via momentum rotations by B_z in the cavity. Then, we can distinguish the contributions to acceleration by the E_x field and cavities, depending on whether the acceleration is caused by E_x or E_y . Figure 10(c) shows the comparison of energy changes by E_x and E_y . The x axis and y axis denote the total energy change by E_x and that by E_y . A significant portion of FEs ($\simeq 70\%$) gain energy by both E_x and E_y , located in the first quadrant of panel (c). The contributions of E_x and E_y to the energy gain are comparable in magnitude. Then, the E_x field and the cavity equally contribute to the accelerations of FEs in the case of $a_0 = 0.085$.

Figures 10(d)–10(f) show the results of the same analysis from the simulations with $a_0 = 0.027$ ($I_L \lambda_{L(\mu\text{m})}^2 = 10^{15} \text{ W/cm}^2$). We picked up electrons with energies over 80 keV reaching the gate at $t = 40 \text{ ps}$.

Compared to the case with $a_0 = 0.085$, the acceleration is weaker. The maximum momentum has decreased from $\simeq 5 m_e c$ to $\simeq 1.4 m_e c$, whereas p_x remains dominant over p_y and p_z . While the energy gain is almost contributed by E_x , with a slight contribution from E_y . As the laser intensity decreases, cavitation weakens, and its contribution to electron acceleration diminishes accordingly. Therefore, within this study, the contribution of the cavity to the electron acceleration is at most comparable to that of the E_x field, and we can conclude that the E_x field is the dominant force of this stochastic acceleration.

IV. ESTIMATE OF MEAN ENERGY

We estimated the mean energy of FEs to confirm whether it is possible to generate electrons with energies over 100 keV through the stochastic acceleration process driven by the E_x field. The energy gain of an electron in the stochastic acceleration can be described as a diffusion process in energy space. In the diffusion process, electrons undergo acceleration or deceleration stochastically, and their histories of kinetic energy are presented as a random walk. The mean energy of

FEs T_f refers to the averaged spread of their histories in energy space.

First, we considered the probability transition process of an electron in energy space. We assume that the electron with kinetic energy $\mathcal{E}_k$ at time t gains energy or loses energy $\Delta\mathcal{E}_k$ with a same probability 1/2 at time $t + \Delta t$ in one kick by the E_x field. Then, the diffusion equation in energy space can be written as

$$\frac{\partial P}{\partial t} = D \frac{\partial^2 P}{\partial \mathcal{E}_k^2}, \quad (6)$$

$$D \equiv \frac{(\Delta\mathcal{E}_k)^2}{2\Delta t}. \quad (7)$$

Here, $P(\mathcal{E}_k, t)$ denotes the probability that the electron possesses energy $\mathcal{E}_k$ at time t .

Following Eq. (6), the probability distribution of the kinetic energy of an electron with initial energy $\mathcal{E}_{k,0}$ progressively broadens in energy space over time. The standard deviation of its energy distribution $\sigma_{\mathcal{E}_k}$ represents the average gained energy, corresponding to the mean energy of FEs T_f when $\mathcal{E}_{k,0}$ is negligible compared to $\sigma_{\mathcal{E}_k}$. Now, $\sigma_{\mathcal{E}_k} = T_f \sim 10^2$ keV and $\mathcal{E}_{k,0} \simeq T_e \sim 10$ keV satisfy this assumption. After N -times interactions with E_x , the standard deviation $\sigma_{\mathcal{E}_k}(t = N\Delta t)$ can be written as follows using the diffusion coefficient D :

$$\sigma_{\mathcal{E}_k}(t = N\Delta t) = \sqrt{2DN\Delta t} = \Delta\mathcal{E}_k\sqrt{N}. \quad (8)$$

Second, we considered the energy gain $\Delta\mathcal{E}_k$ in one kick by E_x . This gain can be expressed as

$$\Delta\mathcal{E}_k = e\bar{E}_x\Delta x, \quad (9)$$

here, $\bar{E}_x$ and Δx denote the mean amplitude of E_x and the characteristic spatial size of the E_x field, respectively. From the simulations, $\bar{E}_x$ is roughly 10% of the incident laser amplitude a_0 [cf. Figs. 5(c) and 5(h)]:

$$\frac{e\bar{E}_x}{m_e c \omega_L} \sim 0.1 a_0. \quad (10)$$

For the estimation of Δx , we adopted the wavelength of the incident laser λ_L as a characteristic scale of the E_x field:

$$\Delta x \sim \frac{\lambda_L}{2}. \quad (11)$$

As seen in Fig. 6(b), the wave numbers of LWs are distributed mainly in $k/k_L \lesssim 1$. Moreover, waves with $k/k_L \ll 1$ are unlikely to contribute to electron acceleration or deceleration. Then, assuming $k/k_L \sim 1$ for the representative wave number of the E_x wave provides a reasonable approximation.

The number of kicks by the E_x field, N , can be written as follows:

$$N = \frac{L_{\text{acc}}}{\Delta x}, \quad (12)$$

where L_{acc} denotes the acceleration length, defined as the averaged distance electrons travel through the E_x field before reaching the gate. Here, we assumed L_{acc} as the length of the initial underdense plasma region:

$$L_{\text{acc}} \sim 4L_{\text{scl}}. \quad (13)$$

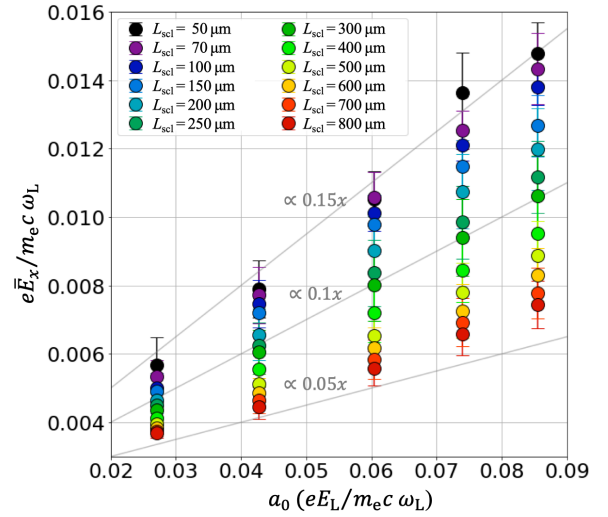


FIG. 11. The mean magnitude of the longitudinal electric field E_x vs the incident laser amplitude a_0 . Now, a_0 takes five values, corresponding to $I_L \lambda_{L(\mu\text{m})} = 1, 2.5, 5, 7.5 \times 10^{15}$, and 10^{16} W/cm². The mean magnitude $\bar{E}_x$ is obtained from 1D collisional simulations and averaged over the spatial range from $x_0 = 1000 \mu\text{m}$ (initial plasma front) to $x_{nc} \simeq x_0 + 3.82L_{\text{scl}}$ (initially $n_e/n_c = 1$) and the temporal range from $t = 35$ to 100 ps. The three solid gray lines are auxiliary lines with slopes of 0.05, 0.1, and 0.15, respectively, to facilitate the comparison between a_0 and $\bar{E}_x$.

Substituting Eqs. (9)–(13) into Eq. (8) provides the order-of-magnitude estimate of $\sigma_{\mathcal{E}_k}$:

$$\sigma_{\mathcal{E}_k} = e\bar{E}_x\sqrt{L_{\text{acc}}\Delta x}, \quad (14)$$

$$\sim m_e c^2 a_0 \sqrt{L_{\text{scl}}(\mu\text{m})}, \quad (15)$$

$$\simeq 100 \left(\frac{a_0}{0.027} \right) \left(\frac{L_{\text{scl}}(\mu\text{m})}{50 \mu\text{m}} \right)^{1/2} (\text{keV}). \quad (16)$$

In the present simulations, the combination $(a_0, L_{\text{scl}}) = (0.027, 50)$ gives the minimum value of $\sigma_{\mathcal{E}_k} \sim 100$ keV, while $(a_0, L_{\text{scl}}) = (0.085, 800)$ gives the maximum value ~ 1.2 MeV. Thus, Eq. (15) can explain the acceleration of FEs with mean energies over 100 keV. The energy gain of one kick by the E_x field is small, but the cumulative effect of $N \sim 4L_{\text{scl}}/(\lambda_L/2) \sim 10^{2-3}$ of kicks enables electrons to be accelerated to hundreds of keV.

Additionally, $\sigma_{\mathcal{E}_k}$ is proportional to a_0 , i.e., to $I_L^{1/2}$, consistent with the observed dependence of the mean energy T_f on I_L shown in Fig. 4(a). This agreement further supports the validity of the estimates derived from Eq. (14), based on stochastic acceleration. On the other hand, as shown in Fig. 4(a), the dependence of T_f on L_{scl} is not monotonic and varies with I_L . This behavior likely arises because L_{acc} also depends on I_L , considering intensity-dependent plasma dynamics such as the expansion of the plasma front and the depletion of the laser light.

To obtain more precise estimates of T_f , we refined the approximation of $\bar{E}_x$ using the simulation results. Figure 11 shows the mean magnitudes of E_x observed from the 1D simulations. They vary between 5% and 15% around 10% of the laser amplitude, depending on L_{scl} . We adopted these values

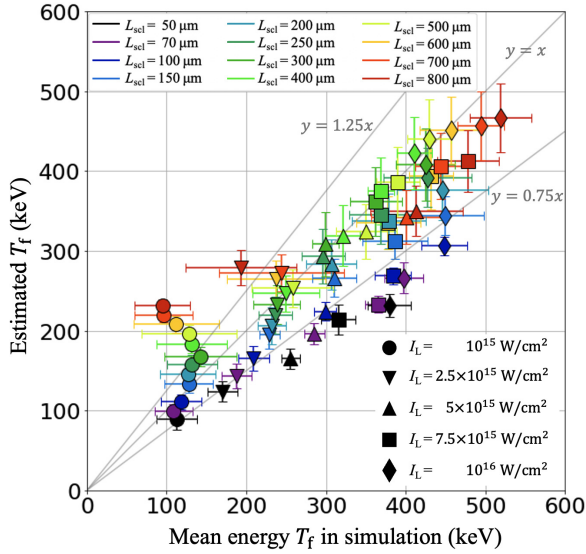


FIG. 12. The estimates of the fast electrons' mean energy T_f obtained from Eq. (14) using simulated $\bar{E}_x$ in Fig. 11.

as the estimate of $\bar{E}_x$ instead of Eq.(10). Figure 12 shows estimates of $\sigma_{\mathcal{E}_k}$ using the $\bar{E}_x$ obtained from the simulations. Here, the frequency of $\bar{E}_x$ is assumed to be approximately $\omega_L/2$. These estimates show good agreement with the observations, with 45 out of the 60 points deviating by less than 25% from the observed T_f . The warm-colored circles and down-pointing triangles, representing cases with lower laser intensity and longer density scale length, tend to deviate upward. This discrepancy may be due to an overestimation of L_{acc} . Longer L_{scl} causes laser depletion in lower-density plasma, preventing the laser from reaching higher-density plasma and from growing density perturbations in this region. Conversely, the cool-colored up-pointing triangles, squares, and diamonds, representing cases with higher laser intensity and shorter density scale length, tend to deviate downward. This trend may be due to an underestimation of L_{acc} , not considering plasma expansions. In addition, the influence of cavities, which can contribute to the increase in T_f as much as the E_x field in higher-intensity cases, as shown in Fig. 10, is not accounted for in the current estimates.

V. CONCLUSION

This study uncovers the generation of fast electrons with energies exceeding 100 keV through a previously unproposed stochastic acceleration process arising from the interaction between inhomogeneous plasmas and lasers of subrelativistic intensity. We conducted collisional 1D PIC simulations over durations up to 100 ps, varying the laser intensity and the plasma density scale length, and identified a quasisteady generation of fast electrons after $t \simeq 35$ ps. Complementary collisionless simulations confirm a transition of the acceleration process from electron trapping by the SRS-driven large-amplitude Langmuir waves, as suggested in the previous studies, to diffusive energy gain in the fluctuating longitudinal electric fields generated by SRS-driven density perturbations.

This stochastic process is not significantly affected by collisions. In the early stages of simulations, the effect of

collision damping is negligible compared to that of Landau damping in the kinetic regime [17,39]. Later, the Landau damping rate decreases due to the electron trapping by Langmuir waves, and collisions might limit the growth of SRS. Nevertheless, both collisional and collisionless simulations reach a quasisteady state characterized by the continuous generation of fast electrons via the stochastic process. Although collisions may reduce the SRS growth, the long-time laser-plasma interactions under the parameters considered in this study ultimately lead to the SRS saturation, the development of density perturbations, and the emergence of a disordered state. Furthermore, the multidimensionality is not expected to change the emergence of the stochastic process. In multidimensional simulations, competitions among PIs arise, affecting their growth and the formation of density perturbations [45]. However, once their growth saturates, the acceleration process is also likely to transition into a stochastic stage.

The simulations further reveal a unique electron confinement motion between SRS-formed density cavities and clarify its influence on electron acceleration. In the case of higher laser intensity and pronounced cavity growth, the contribution of the cavities to electron acceleration becomes comparable to that of the fluctuating field. In contrast, under lower-intensity cases, cavity growth is gentle, and its contribution to acceleration correspondingly decreases. These results lead to the conclusion that the longitudinal fluctuating electric fields dominantly drive the stochastic acceleration process.

To quantify this process, we conducted the theoretical estimation of the mean energy of fast electrons based on the energy diffusion equation. The model given by Eq. (14) successfully explains the acceleration of electrons beyond 100 keV and reproduces the observed dependence of the mean energy on the laser intensity. Using the amplitudes of the fluctuating fields measured in the simulations, Eq. (14) provides accurate estimates of the mean energies. This predictive capability is particularly valuable for estimating energy deposition and evaluating shock-wave enhancement by fast electrons at the ignition stage of the shock ignition. These results also shed light on the reported measurements of hot electrons (>100 keV) in experiments on laser-plasma instabilities [10–12] that were not fully understood theoretically up to now.

ACKNOWLEDGMENTS

Y.T. would like to express gratitude to Dr. Y. Sakawa and Mr. C. Nakatsuji for their valuable discussions and insightful comments, and to the referees for their careful reading and constructive suggestions. This study was supported by JSPS KAKENHI Grants No. 23K20038, No. JP24H00204, and No. JP23K03354, JSPS Core-to-Core Program Grant No. JPJSCCA20230003, JST SPRING Grant No. JPMJSP2138, and Grant-in-Aid for JSPS Fellows No. 23KJ1444. This study was partly achieved through the use of OCTOPUS and SQUID at D3 Center, The University of Osaka.

DATA AVAILABILITY

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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