Coherent Forward Stimulated-Brillouin Scattering of a Spatially Incoherent Laser Beam in a Plasma and Its Effect on Beam Spray

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A statistical model for forward stimulated-Brillouin scattering is developed for a spatially incoherent, monochromatic, laser beam propagating in a plasma. The threshold above which the laser beam spatial incoherence cannot prevent the coherent growth of forward stimulated-Brillouin scattering is computed. It is found to be well below the threshold for self-focusing. Three-dimensional simulations confirm its existence and reveal the onset of beam spray above it. From these results, we propose a new figure of merit for the control of propagation through a plasma of a spatially incoherent laser beam.

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Controlling the propagation of randomized laser beams through large scale plasmas is crucial for inertial confinement fusion (ICF) [1]. The coupling of the high intensity laser wave to ion and electron plasma waves can lead to laser light backscattering, spreading, or bending [2]. To avoid such deleterious effects and to improve the laser illumination uniformity, laser beam smoothing techniques have been designed [3]. The resulting laser intensity in the focal volume consists of many randomly distributed in time and/or in space spikes, the so-called speckles. Statistical properties of such smoothed laser beams are known in vacuum [4] but can be strongly modified by the interaction with the plasma. This phenomenon, referred to as plasma induced smoothing, arises from the laser light forward scattering on plasma density perturbations induced by the laser beam [5–7]. Several processes have been investigated to determine the origin of plasma induced smoothing, namely, self-focusing (SF), filament instability (FI), forward stimulated-Brillouin scattering (FSBS), and multiple scattering.

The relative importance of each of these processes depends on the average power carried in a speckle, $\langle P\rangle$, compared to the critical power for SF [2]: $P_c = (8\pi/k_0^2)cT_e n_c^2/n_0$, where $k_0 = 2\pi/\lambda_0$ is the laser wave number, c is the light velocity, n_c is the critical density beyond which the plasma is opaque to the laser, and n_0 and T_e are the electron density and temperature, respectively. A beam with $\langle P\rangle \gtrsim P_c$ is unstable with regard to SF and FI and it undergoes a strong spraying [5,6]. In this Letter, we restrict ourselves to the regime $\langle P\rangle \ll P_c$.

At lower intensities, plasma induced smoothing is due to FSBS [7]. FSBS is well described in the case of a coherent pump wave [8]. The case of a temporally incoherent pump wave has been investigated in Ref. [9] in the limit of a very short coherence time. It was shown that, while SF is prevented by the pump incoherence, FSBS can be the dominant instability. In Ref. [6], the case of a spatially incoherent, monochromatic, laser beam has been investigated. It was found that the spatial growth rate of the light scattered out of the incident beam aperture is strongly reduced compared to the coherent case. In this Letter, we show that this result is incorrect: laser beam spatial smoothing does not necessarily prevent strong FSBS and the associated beam spray.

We consider FSBS driven by a monochromatic, spatially incoherent, pump wave. This parametric instability is due to the resonant coupling of the large amplitude laser pump $a^{(p)}$ with the given frequency $\omega^{(p)} = \omega_0$ to an ion acoustic wave (IAW) $a^{(\mathrm{IAW})}$ with a low frequency $\omega^{(\mathrm{IAW})} \ll \omega_0$ and an electromagnetic wave $a^{(s)}$ scattered in the near forward direction at the frequency $\omega^{(s)} = \omega_0 - \omega^{(\mathrm{IAW})}$.

We describe the electromagnetic wave propagation in the underdense plasma within the paraxial approximation after a Fourier transform on the two-dimensional plane transverse to the beam mean propagation direction z. This approach is valid for low density plasmas $n_0 \ll n_c$, assuming small pump angular aperture $\theta_p \ll 1$ and scattering angles $\theta_s \equiv |\mathbf{k}|/k_0 \ll 1$, where $\mathbf{k}$ is the wave-vector component in the transverse plane. We restrict ourselves to the linear phase of instability when the pump wave undergoes only diffraction: $(c\partial_z + i\Omega_{\mathbf{k}}^{(p)})a_{\mathbf{k}}^{(p)} = 0$, where $\Omega_{\mathbf{k}}^{(p)} = c^2\mathbf{k}^2/2\omega_0$. The scattered wave equation,

$$(c\partial_z + i\Omega_{\mathbf{k}}^{(s)})a_{\mathbf{k}}^{(s)} = \gamma_0 \int d\mathbf{k}_p a_{\mathbf{k}_p}^{(p)} a_{\mathbf{k}_p-\mathbf{k}}^{(\mathrm{IAW})*}, \qquad (1)$$

accounts for both diffraction $\Omega_{\mathbf{k}}^{(s)} = c^2\mathbf{k}^2/2\omega_0$ and coupling to the IAW of amplitude $a^{(\mathrm{IAW})}$ that propagates with the ion acoustic velocity $c_s \ll c$ in the transverse plane [10]. Neglecting the longitudinal dependence of $a^{(\mathrm{IAW})}$ compared to its transverse one, the excitation of this wave by the ponderomotive force is described by a time-enveloped wave equation:

$$(\partial_t + \nu_{\rm IAW} + i\Omega_{\mathbf{k}}^{\rm (IAW)})a_{\mathbf{k}}^{\rm (IAW)} = \gamma_0 \int d\mathbf{k}_p a_{\mathbf{k}_p}^{(p)} a_{\mathbf{k}_p-\mathbf{k}}^{(s)*}, \quad (2)$$

where $\Omega_{\mathbf{k}}^{\rm (IAW)} = \omega^{\rm (IAW)} - c_s|\mathbf{k}| \ll \omega^{\rm (IAW)}$ is the frequency detuning and $\nu_{\rm IAW} \ll \omega^{\rm (IAW)}$ is the IAW damping rate. Amplitudes $a_{\mathbf{k}}^{(p,s,{\rm IAW})}$ are dimensionless and the FSBS coupling constant is defined as $\gamma_0^2 = \omega^{\rm (IAW)}\omega_0(n_0/n_c) \times \langle I\rangle/(8cn_cT_e)$, where $\langle I\rangle$ is the incident laser average intensity. Equations (1) and (2) have the standard form of coupled mode equations for parametric instabilities [2,11], where the convolution on wave vectors $\mathbf{k}_p$ accounts for the pump spatial incoherence.

Statistical properties of the spatially incoherent, monochromatic, pump can be specified. For most of standard spatial smoothing techniques, they follow a Gaussian statistics with a zero mean value [4]:

$$\langle a_{\mathbf{k}_p}^{(p)}\rangle = \langle a_{\mathbf{k}_p}^{(p)} a_{\mathbf{k}_p'}^{(p)}\rangle = 0, \qquad \langle a_{\mathbf{k}_p}^{(p)} a_{\mathbf{k}_p'}^{(p)*}\rangle = n_{\mathbf{k}_p}^{(p)}\delta(\mathbf{k}_p - \mathbf{k}_p'),$$

where $n_{\mathbf{k}_p}^{(p)}$ is the pump spatial spectrum normalized to one: $\int d\mathbf{k}_p n_{\mathbf{k}_p}^{(p)} = 1$. It does not depend on the longitudinal coordinate z as long as the propagation distance is shorter than the beam Rayleigh length (typically a few mm under ICF conditions). For applications, we consider a Gaussian spectrum: $n_{\mathbf{k}_p}^{(p)} = (\rho_0^2/2\pi)\exp(-\rho_0^2\mathbf{k}_p^2/2)$, where ρ_0 is the speckle transverse width.

Because of the pump incoherence, the scattered and acoustic fields are stochastic and their statistical proprieties must be addressed. This is achieved by using the standard statistical method developed in the context of wave propagation in random media [12], which allows one to compute the successive momenta of $a_{\mathbf{k}}^{(s,{\rm IAW})}$. An application of this method to laser-plasma interaction can be found in Ref. [13]. In this Letter, we restrict ourselves to the study of the average fields $\langle a_{\mathbf{k}}^{(s,{\rm IAW})}\rangle$ and take, as a criterion for instability, the condition that at least one of the average fields, $\langle a_{\mathbf{k}}^{(s)}\rangle$ or $\langle a_{\mathbf{k}}^{\rm (IAW)}\rangle$, grows with time and/or in space. Average intensities $\langle |a_{\mathbf{k}}^{(s,{\rm IAW})}|^2\rangle$ are not computed. They can only evolve faster than $\max\{|\langle a_{\mathbf{k}}^{(s)}\rangle|^2, |\langle a_{\mathbf{k}}^{\rm (IAW)}\rangle|^2\}$, and our approach provides a sufficient condition for instability.

The method of construction of equations for average quantities consists in averaging Eqs. (1) and (2) and expanding the higher order moments as a series in powers of γ_0^2. The closed set of equations for $\langle a_{\mathbf{k}}^{(s,{\rm IAW})}\rangle$ is obtained after Laplace transforms in z and t, $\langle a_{\mathbf{k}}^{(s,{\rm IAW})}\rangle = (2\pi)^{-2} \times \int dqd\gamma \langle \hat{a}_{\mathbf{k}}^{(s,{\rm IAW})}\rangle e^{(q-i\Omega_{\mathbf{k}}^{(s)}/c)z+(\gamma-i\Omega_{\mathbf{k}}^{\rm (IAW)})t}$ and neglecting higher order terms $\sim\gamma_0^4$:

$$cq\langle \hat{a}_{\mathbf{k}}^{(s)}\rangle = \frac{\gamma_0^2}{\gamma + \nu_{\rm IAW} + \Delta\omega_{\rm IAW}}\langle \hat{a}_{\mathbf{k}}^{(s)}\rangle, \quad (3)$$

$$(\gamma + \nu_{\rm IAW})\langle \hat{a}_{\mathbf{k}}^{\rm (IAW)}\rangle = \frac{\gamma_0^2}{cq + \Delta\omega_s}\langle \hat{a}_{\mathbf{k}}^{\rm (IAW)}\rangle. \quad (4)$$

Here, the spectral widths $\Delta\omega_{\rm IAW}/\omega^{\rm (IAW)} = \sqrt{2/\pi}\theta_p/\theta_s$ and $\Delta\omega_s/\omega_0 = \sqrt{2/\pi}\theta_p\theta_s$ account for the effect of the pump incoherence on the wave coupling, $\theta_p = (k_0\rho_0)^{-1}$ is the half-aperture angle of the pump wave, and $\theta_s = |\mathbf{k}|/k_0$ is the scattering angle. In what follows, we consider the case $\theta_s \gtrsim \theta_p$ and we assume the ordering $\nu_{\rm IAW} \ll \Delta\omega_{\rm IAW} \ll \Delta\omega_s$. The validity of Eqs. (3) and (4) requires the series on γ_0^2 to converge. This is true when $K_B \ll 1$, where $K_B \equiv \gamma_0^2/[(\gamma + \Delta\omega_{\rm IAW})(cq + \Delta\omega_s)]$ is the so-called Kubo number. This condition is verified below.

Equation (3) describes the evolution of the scattered electromagnetic wave average amplitude. Performing the inverse Laplace transformation, one finds that for $t > t_{\rm sat}^{(s)} \equiv (\gamma_0/\Delta\omega_{\rm IAW})^2 z/c$, the wave experiences a spatial amplification, $\langle a_{\mathbf{k}}^{(s)}\rangle \propto \exp(q_{\rm incoh}^{(s)}z)$, with $q_{\rm incoh}^{(s)} \equiv \gamma_0^2/(c\Delta\omega_{\rm IAW})$. Similarly to Ref. [6], this spatial growth rate is reduced by the pump incoherence. For shorter times $t < t_{\rm sat}^{(s)}$, the scattered wave is in a transient regime of mixed growth in t and z, $\langle a_{\mathbf{k}}^{(s)}\rangle \propto \exp(2\gamma_0\sqrt{tz/c} - \Delta\omega_{\rm IAW}t)$. In both cases the temporal and spatial growth rates are small compared to the pump spectral widths $\Delta\omega_{\rm IAW}$ and $\Delta\omega_s$, and the validity condition of Eq. (3) can be written as $K_B = \gamma_0^2/(\Delta\omega_{\rm IAW}\Delta\omega_s) \simeq \langle P\rangle/P_c \ll 1$, so that Eq. (3) is valid in both limits. These behaviors, where the growth of $\langle a_{\mathbf{k}}^{(s)}\rangle$ is strongly reduced by the pump incoherence, are referred to as incoherent FSBS amplification.

A similar analysis of Eq. (4) provides the evolution of the IAW average amplitude. For times $t > t_{\rm sat}^{\rm (IAW)} \equiv (\gamma_0^2/\nu_{\rm IAW})^2 z/c$, the IAW exhibits a spatial amplification, $\langle a_{\mathbf{k}}^{\rm (IAW)}\rangle \propto \exp((q_{\rm coh}^{\rm (IAW)} - \Delta\omega_s/c)z)$, where $q_{\rm coh}^{\rm (IAW)} = \gamma_0^2/(c\nu_{\rm IAW})$ is the spatial growth rate of FSBS driven by a coherent pump. In this limit, $K_B = \nu_{\rm IAW}/\Delta\omega_{\rm IAW}$, and the validity condition of Eq. (4) is fulfilled. For shorter times $t < t_{\rm sat}^{\rm (IAW)}$, the IAW is in a transient regime, $\langle a_{\mathbf{k}}^{\rm (IAW)}\rangle \propto \exp(2\gamma_0\sqrt{tz/c} - \Delta\omega_s z - \nu_{\rm IAW}t)$. Here, Eq. (4) is valid only for sufficiently long times, $t > t_{\rm val}^{\rm (IAW)} \equiv (\gamma_0^2/\Delta\omega_{\rm IAW})^2 z/c$. The time $t_{\rm val}^{\rm (IAW)} \ll t_{\rm sat}^{\rm (IAW)}$ does not exceed a few picoseconds for a millimetric plasma so that Eq. (4) is quickly valid.

The condition for spatial growth of the averaged IAW amplitude, $cq_{\rm coh}^{\rm (IAW)} > \Delta\omega_s$, defines a threshold for the average speckle power: $\langle P\rangle/P_c > \sqrt{2/\pi}(\nu_{\rm IAW}/\omega^{\rm (IAW)}) \times (\theta_s/\theta_p)$. Above this threshold, a coherent spatial amplification $\langle a_{\mathbf{k}}^{\rm (IAW)}\rangle \propto \exp(q_{\rm coh}^{\rm (IAW)}z)$ occurs, which is not reduced by the pump incoherence. So far, this result applies only to the asymptotic limit $t > t_{\rm sat}^{\rm (IAW)}$. However, defining the effective spatial growth rate in the transient regime as $q_{\rm eff}^{\rm (IAW)} \equiv (\gamma_0^2 ct/z)^{1/2} - \Delta\omega_s$, one obtains that the pump incoherence does not affect the spatial amplification of $\langle a_{\mathbf{k}}^{\rm (IAW)}\rangle$ for $t \gg (\Delta\omega_s/\gamma_0)^2 z/c$. For a millimetric plasma, this time ranges from a few tens to a few hundred picoseconds. Thus, the instability develops on times shorter than characteristic durations of ICF laser pulses.

Beam spray can be expected if waves scattered outside the beam aperture angle are strongly (coherently) amplified. Thus, by taking $\theta_p = \theta_s$ in the previous threshold, one defines the figure of merit (FOM) for beam spray:

$$C \equiv \sqrt{\frac{\pi}{2}}\gamma_T \frac{\langle P\rangle/P_c}{\nu_{\rm IAW}/\omega^{\rm (IAW)}}. \tag{5}$$

Here, the factor γ_T accounts for thermally enhanced excitation of IAW. In conditions relevant for ICF, it depends on the ion charge Z and electron-ion mean-free path λ_{ei}: $\gamma_T = 1 + 1.76Z^{5/7}(\rho_0/\lambda_{ei})^{4/7}$ [14]. The FOM (5) has a straightforward physical meaning: it defines the critical angle $C\theta_p$ below which amplification proceeds in the coherent regime. Enhanced FSBS within this aperture provides plasma induced smoothing and in turn inhibits the development of backward scattering instabilities. Coherent excitation of FSBS occurs outside the incident cone if $C > 1$ and beam spray can then be expected. Because of the small IAW damping rate, $\nu_{\rm IAW}/\omega^{\rm (IAW)} \ll 1$, beam spray may occur well below the SF threshold. Our theory thus provides an explanation for the plasma induced incoherence observed in Ref. [7] below the SF threshold. This induced smoothing is optimal for $C \sim 1$, i.e., when FSBS induces temporal incoherence without enhancing the spatial aperture of transmitted light.

The effect of FSBS on beam spray has been confirmed in 3D simulations performed with the paraxial interaction code PARAX [15]. The plasma response is described by a linear IAW equation driven by the ponderomotive force. We consider a helium plasma with electron density $0.03n_c$, and electron temperature 500 eV. The IAW velocity is $c_s \simeq 0.17\ \mu$m/ps and the critical power for SF is $P_c \simeq 640$ MW. The IAW damping rate, $\nu_{\rm IAW}/\omega^{\rm (IAW)} =$ 2.75%, 5.5%, or 8.25%, is chosen independently. A Gaussian laser beam with $\lambda_0 = 1.053\ \mu$m is focused through a random phase plate providing a speckle pattern with the coherence width $\rho_0 \simeq 4.3\ \mu$m. The corresponding Rayleigh length of the speckle is $L_R = k_0\rho_0^2 \simeq 110\ \mu$m. In order to avoid SF and FI, the average intensity is varied in the range $(1.1\text{–}16) \times 10^{13}$ W/cm^2, corresponding to variation of the average power in a speckle $\langle P\rangle = \pi\rho_0^2\langle I\rangle$ from 1% to 14% of P_c. As a result, there is no speckle in the focal spot carrying a power above P_c. Multiple scattering on density perturbations driven by the randomized beam, which is a nonstimulated process, is kept to a rather low level by introducing a linear ramp time $t_m \simeq 130$ ps $\gg \rho_0/c_s$ in the laser intensity profile, then decreasing the initial density perturbations [7].

The laser angular aperture, defined as $\theta(z) = k_0^{-1}[\int d\mathbf{k}\mathbf{k}^2 n_{\mathbf{k}}/\int d\mathbf{k} n_{\mathbf{k}}]^{1/2}$, where $n_{\mathbf{k}}$ is the laser spatial spectrum, has been calculated along the whole propagation through the plasma. It increases with both t and z. It is shown in Fig. 1(a) at $t = 2$ ns and $z = 1.2$ mm as a function of $\langle P\rangle/P_c$. Each curve corresponds to a given IAW damping and consists of two parts. At low $\langle P\rangle/P_c$, no beam spray is observed: $\theta(z) = \theta(0) \sim \theta_p$. When $\langle P\rangle/P_c$ increases, transition to beam spray is observed, well below the SF threshold. The threshold power increases with $\nu_{\rm IAW}/\omega^{\rm (IAW)}$. For the highest values of the IAW damping rate, it is in very good agreement with the theoretical predictions (5) [vertical arrows in Fig. 1(a)]. For $\nu_{\rm IAW}/\omega^{\rm (IAW)} = 2.75\%$, however, beam spray occurs at a power larger than predicted by criterion (5) because the instability has not yet reached its convective regime.

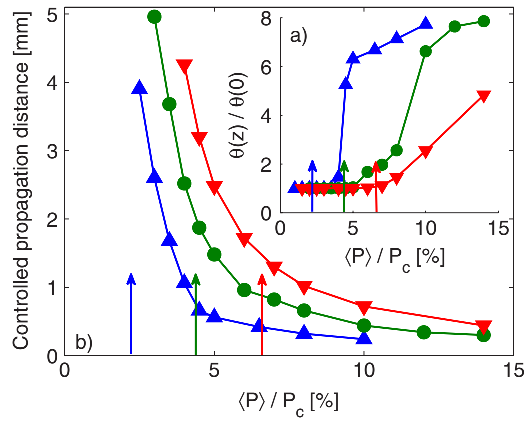

FIG. 1 (color online). (a) Aperture angle of the transmitted laser beam versus the normalized average power $\langle P\rangle/P_c$ at $t =$ 2 ns and $z = 1.2$ mm. (b) Propagation distance over which the aperture angle remains below $1.2\theta(0)$. Arrows show the theoretical threshold (5). Ion acoustic damping rates are $\nu_{\rm IAW}/\omega^{\rm (IAW)} =$ 2.75% (blue △), 5.5% (green ○), and 8.25% (red ▽).

In a complementary way, the distance over which the laser aperture angle remains below $1.2\theta_p$ is shown as a function of $\langle P\rangle/P_c$ in Fig. 1(b) at $t = 2$ ns. At low $\langle P\rangle/P_c$, the laser propagates over a long distance (a few mm) without suffering angular spreading. For higher powers, FSBS-induced beam spray limits the propagation distance to a few speckle Rayleigh lengths. Transition between controlled propagation and beam spray is well below the SF threshold and it clearly depends on the IAW damping rate. Criterion (5) [vertical arrows in Fig. 1(b)] provides a satisfactory estimation for this transition.

The spatial growth rate of the light scattered at the frequency $\omega_0 - 2c_s/\rho_0$ has been calculated in the temporal window between 1 and 2 ns. It is plotted as a function of $\langle P\rangle/P_c$ in Fig. 2 for three IAW damping rates. In the low power regime, the spatial growth rate is less than $0.5 \times 10^{-2}\ \mu$m^{-1}. It evolves almost linearly with $\langle P\rangle/P_c$ and it agrees rather well with our predictions in the incoherent regime (dashed line) where the FSBS amplification is limited by the pump incoherence. Above the threshold power, much higher growth rates are measured. They are in the range $(0.5\text{–}2.5) \times 10^{-2}\ \mu$m^{-1}, which is consistent with expectations for the coherent regime. Moreover, it confirms that beam smoothing does not prevent strong excitation of the instability above a power that is qualitatively in good agreement with Eq. (5). The discrepancy

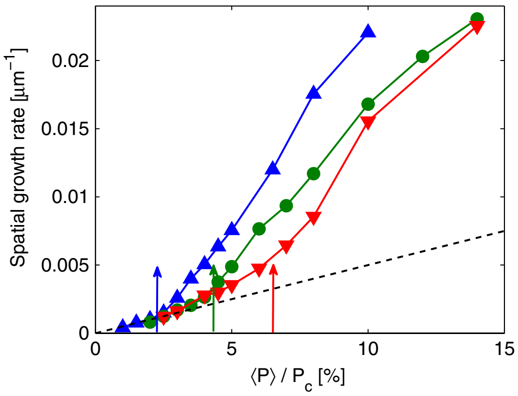

FIG. 2 (color online). Dependence of the spatial growth rate of the scattered wave intensity between 1 and 2 ns on the normalized average power $\langle P\rangle/P_c$. Same notations as in Fig. 1. The dashed curve shows the incoherent growth rate.

between theory and numerical simulations is due to the fact that $t_{\rm sat}^{(\rm IAW)}$ ranges in most of our simulations between 1 and 2 ns. Amplification thus occurs in a regime between the transient and convective limits.

Our theory of FSBS driven by a spatially smoothed, monochromatic, laser beam is complementary to that developed by Lushnikov and Rose [9] for a broadband pump with a small coherence time, $\tau_c \ll \rho_0/c_s$. The conclusions concerning the coherent growth of FSBS in this regime are similar to ours, as well as the predicted threshold. According to the present analysis, the regime referred to as equilibrium in Ref. [9] may rather be an incoherent regime of FSBS with a weak, albeit existing, growth. By combining our results with those of Ref. [9], we conjecture that these conclusions may apply also to the intermediate regime $\tau_c \sim \rho_0/c_s$ of interest for ICF.

Beam spray of a temporally coherent spatially smoothed laser beam has been reported recently in experiment [16] and attributed to FI. For these experimental conditions ($n_0 = 0.06n_c$, $T_e = 3.5$ keV, $\lambda_0 = 0.351\ \mu$m, $\nu_{\rm IAW}/\omega^{(\rm IAW)} \simeq 0.15$, and f number $f_\# = 6.7$), the FOM (5) predicts beam spray for average power in a speckle $\langle P\rangle \simeq 7.5\%$ of P_c, where we have estimated $Z \simeq 5.3$ for the CH plasma, $\lambda_{ei} \simeq 75\ \mu$m and $\rho_0 \simeq (2/\pi)\lambda_0 f_\#$. It is in perfect agreement with the average power, $\langle P\rangle \simeq 0.07P_c$, at the measured intensity threshold $\langle I\rangle \simeq 2\times 10^{15}$ W/cm^2. It is significantly below the SF threshold, which prompts us to suggest that the observed beam spray is a direct consequence of the coherent FSBS rather than SF or FI. Equation (5) provides a new, theoretically derived, criterion for beam spray for a spatially smoothed, monochromatic, laser beam. In practical units it reads

$$0.1\gamma_T \frac{\omega^{(\rm IAW)}}{\nu_{\rm IAW}} \lambda_0^2[\mu{\rm m}] I_{13} \frac{n_0}{n_c} \frac{3}{T_e[{\rm keV}]}\left(\frac{f_\#}{8}\right)^2 > 1, \qquad (6)$$

where I_{13} is the laser average intensity in 10^{13} W/cm^2. Equation (6) generalizes the FOM suggested in Ref. [16], which is exactly recovered by taking $\nu_{\rm IAW}/\omega^{(\rm IAW)} \simeq 0.15$ and $\gamma_T \simeq 1.6$ corresponding to experimental conditions. The dependence of the threshold on IAW damping rate provides an important insight on the control of propagation of a randomized laser beam through a plasma.

In conclusion, a statistical model for FSBS driven by a spatially incoherent laser beam is developed. It predicts a fast (coherent) growth of the daughter waves well below the SF threshold. This instability is not stabilized by spatial smoothing of the incident beam and it leads to a strong deterioration of the beam propagation. A threshold condition is derived providing a new FOM for the control of beam propagation through plasmas. The theoretical model is confirmed by 3D numerical simulations and provides an explanation for a recent experiment [16]. This transition from incoherent to coherent growth of FSBS associated to beam spray is of crucial importance for the design of forthcoming ICF experiments.

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[1] J. D. Lindl, Phys. Plasmas **2**, 3933 (1995).
[2] W. L. Kruer *et al.*, *The Physics of Laser Plasma Interactions* (Addison-Wesley, Redwood City, CA, 1988); J. D. Lindl *et al.*, Phys. Plasmas **11**, 339 (2004).
[3] Y. Kato *et al.*, Phys. Rev. Lett. **53**, 1057 (1984); R. Lehmberg and S. Obenschain, Opt. Commun. **46**, 27 (1983).
[4] H. Rose and D. Dubois, Phys. Fluids B **5**, 590 (1993); J. Garnier, Phys. Plasmas **6**, 1601 (1999).
[5] C. Labaune *et al.*, Phys. Fluids B **4**, 2224 (1992); A. J. Schmidt and B. B. Afeyan, Phys. Plasmas **5**, 503 (1998); P. Michel *et al.*, Phys. Plasmas **10**, 3545 (2003).
[6] A. V. Maximov *et al.*, Phys. Plasmas **8**, 1319 (2001).
[7] V. Malka *et al.*, Phys. Rev. Lett. **90**, 075002 (2003); M. Grech *et al.*, Phys. Plasmas **13**, 093104 (2006).
[8] V. V. Eliseev *et al.*, Phys. Plasmas **4**, 4333 (1997).
[9] P. M. Lushnikov and H. A. Rose, Phys. Rev. Lett. **92**, 255003 (2004); P. M. Lushnikov and H. A. Rose, Plasma Phys. Controlled Fusion **48**, 1501 (2006).
[10] Time derivatives in Eq. (1) are neglected, which requires the time for the laser light to travel through a speckle L_R/c to be much shorter than the laser coherence time and the time for an acoustic wave to cross a speckle ρ_0/c_s.
[11] M. N. Rosenbluth, Phys. Rev. Lett. **29**, 565 (1972).
[12] P. Kubo, J. Math. Phys. (N.Y.) **4**, 174 (1963); R. L. Stratonovitch, *Topics in the Theory of Random Noise* (Gordon and Breach, New York, 1967), Vol. 2, Chap. 10; A. Brissaud and U. Frisch, J. Math. Phys. (N.Y.) **15**, 524 (1974).
[13] J. J. Thomson, Nucl. Fusion **15**, 237 (1975); G. Laval *et al.*, Phys. Fluids **20**, 2049 (1977).
[14] A. Brantov *et al.*, Phys. Plasmas **6**, 3002 (1999).
[15] G. Riazuelo and G. Bonnaud, Phys. Plasmas **7**, 3841 (2000).
[16] D. H. Froula *et al.*, Phys. Plasmas **14**, 055705 (2007).