Hindawi
Laser and Particle Beams
Volume 2021, Article ID 9919467, 10 pages
https://doi.org/10.1155/2021/9919467

Research Article

Stimulated Raman Scattering of X-Mode Laser in a Plasma Channel

Sanjay Babu, Asheel Kumar, Ram Jeet, Arvind Kumar, and Ashish Varma

Department of Physics, Faculty of Science, University of Allahabad, Prayagraj 211002, India

Correspondence should be addressed to Asheel Kumar; asheel2002@yahoo.co.in

Received 9 August 2020; Accepted 13 October 2020; Published 5 May 2021

Academic Editor: Dieter Hoffman

Stimulated Raman forward scattering (SRFS) of an intense X-mode laser pump in a preformed parabolic plasma density profile is investigated. The laser pump excites a plasma wave and one/two electromagnetic sideband waves. In Raman forward scattering, the growth rate of the parametric instability scales as two-third powers of the pump amplitude and increases linearly with cyclotron frequency.

1. Introduction

The propagation of intense laser pulses in plasma [1, 2] is relevant to various applications including laser-driven acceleration [3–5], optical harmonic generation [6, 7], X-ray laser [8], laser fusion [9, 10], and magnetic field generation [11]. The laser-plasma interaction leads to a number of relativistic and nonlinear effects such as self-focusing [12–14], parametric Raman and Brillouin instabilities [15], filamentation and modulational instabilities [16–24], and soliton formation.

Stimulated Raman scattering (SRS) is an important parametric instability in plasmas. In stimulated Raman forward scattering (SRFS), a high phase velocity Langmuir wave is produced that can accelerate electrons to high energy [25–35]. In stimulated Raman backward scattering (SRBS), electron plasma wave (EPW) has smaller phase velocity and can cause bulk heating of electrons. The high-amplitude laser wave enables the development of this instability in plasmas with a density significantly higher than the quarter critical density ($n_c/4$), which is usually considered as the SRS density limit for lower laser intensities. Shvets et al. [36] have demonstrated that SRS is strongly modified in a plasma channel, where the plasma frequency varies radially through the radial dependence of the plasma density $n_0(r)$. Liu et al. [37] have studied forward and backward Raman instabilities of a strong but nonrelativistic laser pump in a preformed plasma channel in the limit when plasma thermal effects may be neglected. Mori et al. [38] have developed an elegant formalism of SRFS in one dimension (1D). Panwar et al. [39] have studied SRFS of an intense pulse in a preformed plasma channel with a sequence of two pulses. They have studied the guiding of the main laser pulse through the plasma channel created by two lasers. Sajal et al. [40] have studied SRFS of a relativistic laser pump in a self-created plasma channel. Hassoon et al. [41] have studied the effect of a transverse static magnetic field on SRFS of a laser in a plasma. The X-mode excites an upper hybrid wave and two sidebands. They found that the growth rate of SRFS increases with the magnetic field. Gupta et al. [42] have investigated SRS of laser in a plasma with energetic drifting electrons generated during laser-plasma interaction. They showed numerically that the relativistic effects increase the growth rate of the Raman instability and enhance the amplitude of the decay waves significantly.

Liu et al. [43] developed 1D Vlasov–Maxwell numerical simulation to examine RBS instability in unmagnetized collisional plasma. Their results showed that RBS is enhanced by electron-ion collisions. Kalmykov et al. [44] have studied RBS in a plasma channel with a radial variation of plasma frequency. Paknezhad et al. [45] have investigated RBS of ultrashort laser pulse in a homogeneous cold underdense magnetized plasma by taking into account the relativistic effect of nonlinearity up to third order. The plasma is embedded in a uniform magnetic field. Kaur and Sharma [46] developed nonlocal theory of the SRBS in the propagation of a circularly polarized laser pulse through a hot plasma channel in the presence of a strong axial magnetic field. They established that the growth rate of SRBS of a finite spot size significantly decreases by increasing the magnetic field. Paknezhad et al. [47] have studied the Raman shifted third harmonic backscattering of an intense extraordinary laser wave through a homogeneous transversely magnetized cold plasma. Due to relativistic nonlinearity, the plasma dynamics is modified in the presence of transverse magnetic field, and this can generate the third harmonic scattered wave and electrostatic upper hybrid wave via the Raman scattering process.

In this paper, we examine the SRFS of an X-mode laser pump in a magnetized plasma channel including nonlocal effects. In many experiments in high-power laser-plasma interaction, transverse magnetic fields are self-generated [11]. Laser launched from outside travels in X-mode in such magnetic fields and often creates parabolic density profiles. Thus, the current problem is relevant to experimental situations.

The ponderomotive force due to the front of laser pulse pushes the electrons radially outward on the time scale of a plasma period ω_p^{-1}, creating a radial space charge field, and modifies the electron density, where ω_p is the electron plasma frequency. Laser and the sidebands exert a ponderomotive force on electrons driving the plasma wave. The density perturbation due to plasma wave beats with the oscillatory velocity due to laser pump to produce nonlinear currents, driving the sidebands.

In Section 2, we analyse the SRFS of a laser pump in a preformed channel with nonlocal effects. In Section 3, we discuss our results.

2. Raman Forward Scattering

Consider a two-dimensional plasma channel with a parabolic density profile immersed in a static magnetic field $B_s\hat{z}$. The electron plasma frequency in the channel varies as

$$\omega_p^2 = \omega_{p0}^2\left(1 + \frac{y^2}{L_n^2}\right), \tag{1}$$

where ω_{p0} is the plasma frequency at $y = 0$ and L_n is the density scale length. An X-mode laser propagates through the channel (cf. Figure 1),

$$\mathbf{E}_0 = \mathbf{A}_0(y)\left(\hat{y} - \frac{\varepsilon_{0xy}}{\varepsilon_{0xx}}\hat{x}\right)e^{-i(\omega_0 t - k_{0x}x)}. \tag{2}$$

The plasma permittivity at ω_0 is a tensor. Its components are

$$\begin{aligned}
\varepsilon_{0xx} &= \varepsilon_{0yy} = 1 - \frac{\omega_p^2}{\left(\omega_0^2 - \omega_c^2\right)},\\
\varepsilon_{0xy} &= -\varepsilon_{0yx} = i\frac{\omega_c}{\omega_0}\frac{\omega_p^2}{\left(\omega_0^2 - \omega_c^2\right)},\\
\varepsilon_{0zz} &= 1 - \frac{\omega_p^2}{\omega_0^2},\\
\varepsilon_{0xz} &= \varepsilon_{0zx} = \varepsilon_{0yz} = \varepsilon_{0zy} = 0,
\end{aligned} \tag{3}$$

where $\omega_c = (eB_s/m)$ is electron cyclotron frequency and e and m are the electronic charge and mass, respectively. Maxwell's equations $\nabla \times \mathbf{E}_0 = i\omega_0\mu_0\mathbf{H}_0$ and $\nabla \times \mathbf{H}_0 = -i\omega_0\varepsilon_s \underset{=0}{\varepsilon} \cdot \mathbf{E}_0$ combine to give the wave equation

$$\nabla^2\mathbf{E}_0 - \nabla(\nabla\cdot\mathbf{E}_0) + \frac{\omega_0^2}{C^2}\underset{=}{\varepsilon_0}\cdot\mathbf{E}_0 = 0, \tag{4}$$

where ε_s is the free space permittivity. Here, we chosen the plasma to be uniform and expressed the spatial-temporal variation of $\mathbf{E}_0$ as $e^{-i(\omega_0 t - k_{0x}x - k_{0y}y)}$ we would obtain from equation (4), on replacing ∇ by $(ik_{0y}\hat{y} + ik_{0x}\hat{x})$,

$$\underset{=}{D_0}\cdot\mathbf{E}_0 = 0, \tag{5}$$

where $D_0 = (\omega_0^2/c^2)\underset{=}{\varepsilon_0} - k_0^2\underset{=}{I} + \mathbf{k}_0\mathbf{k}_0$. Its matrix form is expressed as follows:

$$\underset{=0}{D} = \begin{bmatrix} \dfrac{\omega_0^2}{c^2}\varepsilon_{0xx} - k_{0y}^2 & \dfrac{\omega_0^2}{c^2}\varepsilon_{0xy} + k_{0x}k_{0y} & 0\\ -\dfrac{\omega_0^2}{c^2}\varepsilon_{0xy} + k_{0x}k_{0y} & \dfrac{\omega_0^2}{c^2}\varepsilon_{0xx} - k_0^2 & 0\\ 0 & 0 & \dfrac{\omega_0^2}{c^2}\varepsilon_{0zz} - k_0^2 \end{bmatrix}. \tag{6}$$

The local dispersion relation of the mode is given by $|\underset{=}{D_0}| = 0$, which can be written as

$$k_{0y}^2 = \frac{\omega_0^2}{c^2}\frac{\varepsilon_{0xx}^2 + \varepsilon_{0xy}^2}{\varepsilon_{0xx}} - k_{0x}^2 = \frac{\omega_0^2}{c^2}\left[1 - \frac{\omega_p^2}{\omega_0^2}\frac{\omega_0^2 - \omega_p^2}{\omega_0^2 - \omega_p^2 - \omega_c^2}\right] - k_{0x}^2. \tag{7}$$

Equation (2) gives $E_{0x} = -(\varepsilon_{0xy}/\varepsilon_{0xx})E_{0y}$. To incorporate nonlocal effects, one may deduce the mode structure equation from the above equation by replacing k_{0y} by $(-i\,\partial/\partial y)$ and operating over E_{0y},

$$\frac{\partial^2 E_{0y}}{\partial y^2} + \left[\frac{\omega_0^2}{c^2}\left\{1 - \frac{\omega_p^2}{\omega_0^2}\frac{\omega_0^2 - \omega_p^2}{\omega_0^2 - \omega_p^2 - \omega_c^2}\right\} - k_{0x}^2\right]E_{0y} = 0. \tag{8}$$

Define

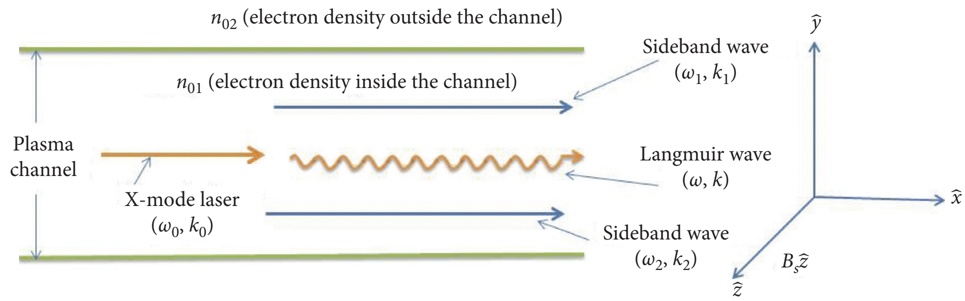

Figure 1: Schematic of excitation of a Langmuir wave and two sidebands by an X-mode laser in a plasma channel.

$$\beta_0^{(1/4)} y = \xi_0,$$

$$\beta_0 = \frac{\omega_{p0}^2}{\omega_0^2}\left\{1 + \frac{\omega_c^2\left(\omega_0^2 - \omega_c^2\right)}{\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2}\right\}\frac{\omega_0^2}{L_n^2 c^2}, \tag{9}$$

$$\lambda_0 = \left[\frac{\omega_0^2}{c^2}\left(1 - \frac{\omega_{p0}^2}{\omega_0^2}\frac{\omega_0^2 - \omega_{p0}^2}{\omega_0^2 - \omega_{p0}^2 - \omega_c^2}\right) - k_{0x}^2\right]\frac{1}{\beta_0^{(1/2)}}.$$

Equation (8) can be written as

$$\frac{\partial^2 E_{0y}}{\partial \xi_0^2} + \left(\lambda_0 - \xi_0^2\right)E_{0y} = 0. \tag{10}$$

The eigenvalue for the fundamental mode is $\lambda_0 = 1$, and the eigenfunction is given by $E_{0y} = A_0 e^{-\xi_0^2/2}$. From the eigenvalue condition, we obtain

$$k_{0x}^2 = \frac{\omega_0^2}{c^2}\left[1 - \frac{\omega_{p0}^2}{\omega_0^2}\frac{\omega_0^2 - \omega_{p0}^2}{\omega_0^2 - \omega_{p0}^2 - \omega_c^2}\right] - \beta_0^{(1/2)}. \tag{11}$$

The pump wave produces oscillatory electron velocity,

$$\begin{aligned} v_{0x} &= -\frac{e}{m}\frac{\omega_c E_{0y}}{\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)}, \\ v_{0y} &= -i\frac{e}{m}\frac{\left(\omega_0^2 - \omega_p^2\right)E_{0y}}{\omega_0\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)}, \\ v_{0z} &= 0, \end{aligned} \tag{12}$$

and parametrically excites a Langmuir wave with electrostatic potential,

$$\phi = \phi(y)e^{-i(\omega t - kx)}, \tag{13}$$

and two electromagnetic sidebands with electric fields,

$$\mathbf{E}_j = A_j\left(\hat{y} - \frac{\varepsilon_{1xy}}{\varepsilon_{1xx}}\hat{x}\right)e^{-i\left(\omega_j t - k_j x\right)}, \tag{14}$$

where $j = 1, 2$ and $\omega_{1,2} = \omega \mp \omega_0$, and $k_{1,2} = k \mp k_0$. The sideband waves produce oscillatory electron velocities,

$$\begin{aligned} v_{jx} &= -\frac{e}{m}\frac{\omega_c E_{jy}}{\left(\omega_j^2 - \omega_c^2 - \omega_p^2\right)}, \\ v_{jy} &= -i\frac{e}{m}\frac{\left(\omega_j^2 - \omega_p^2\right)E_{jy}}{\omega_j\left(\omega_j^2 - \omega_c^2 - \omega_p^2\right)}, \\ v_{jz} &= 0, \end{aligned} \tag{15}$$

where $j = 1, 2$.

The sideband waves couple with the pump to exert a ponderomotive force on electrons at ω, k, $\mathbf{F}_{p\omega} = -(m/2)(\mathbf{v}_0 \cdot \nabla\mathbf{v}_1 + \mathbf{v}_0^* \cdot \nabla\mathbf{v}_2 + \mathbf{v}_1 \cdot \nabla\mathbf{v}_0 + \mathbf{v}_2 \cdot \nabla\mathbf{v}_0^*) - (e/2)(\mathbf{v}_0 \times \mathbf{B}_1 + \mathbf{v}_1 \times \mathbf{B}_0 + \mathbf{v}_0^* \times \mathbf{B}_2 + \mathbf{v}_2 \times \mathbf{B}_0^*)$. In the limit, $\mathbf{k}_0 \gg \mathbf{k}$, $\mathbf{k}_1$ and $\mathbf{k}_2$ are largely along $\hat{x}$, and the x and y components of the ponderomotive force turn out to be

$$\begin{aligned} F_{p\omega x} &= \frac{e^2 i\left(PE_{0y}E_{1y} + QE_{0y} * E_{2y}\right)}{2m\omega_0\omega_1\omega_2\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right)}, \\ F_{p\omega y} &= \frac{e^2\omega_c\left(RE_{0y}E_{1y} + SE_{0y} * E_{2y}\right)}{2m\omega_0\omega_1\omega_2\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right)}, \end{aligned} \tag{16}$$

where

$$\begin{aligned} P &= \omega_2\left\{(k_1 - k_0)\omega_1\omega_0\omega_c^2 + k_1\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right) - k_0\left(\omega_1^2 - \omega_p^2\right)\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)\right\}, \\ Q &= \omega_1\left\{(k_0 - k_2)\omega_2\omega_0\omega_c^2 - k_2\left(\omega_0^2 - \omega_p^2\right)\left(\omega_2^2 - \omega_c^2 - \omega_p^2\right) + k_0\left(\omega_2^2 - \omega_p^2\right)\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)\right\}, \\ R &= \omega_2\left(k_0\omega_1 + k_1\omega_0\right)\omega_c^2, \\ S &= \omega_1\left(k_0\omega_2 + k_2\omega_0\right)\omega_c^2, \end{aligned} \tag{17}$$

For $\omega \ll \omega_0$, $\mathbf{F}_{p\omega} = e\nabla\varphi_p$; one may write

$$\phi_p = V_{01}\left[E_{1y} - \frac{\eta_1}{\eta_2}E_{2y}\right],$$

$$V_{01} = \frac{eE_{0y}\left\{\omega_1\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right) - \omega_c^2\omega_p^2\omega_0\right\}}{2m\omega_0\omega_1^2\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right)}, \tag{18}$$

where

$$\begin{aligned}\eta_1 &= \omega_1^2\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right)\left\{\omega_2\left(\omega_0^2 - \omega_p^2\right)\left(\omega_2^2 - \omega_c^2 - \omega_p^2\right) + \omega_0\omega_c^2\omega_p^2\right\},\\ \eta_2 &= \omega_2^2\left(\omega_2^2 - \omega_c^2 - \omega_p^2\right)\left\{\omega_1\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_c^2 - \omega_p^2\right) - \omega_0\omega_c^2\omega_p^2\right\}.\end{aligned} \tag{19}$$

The drift velocity of plasma electrons due to electrostatic wave of potential φ in the component form is given by

$$\begin{aligned}v_{\omega x} &= -\frac{e}{m}\frac{\omega k}{\left(\omega^2 - \omega_c^2\right)}\phi,\\ v_{\omega y} &= -i\frac{e}{m}\frac{\omega_c k}{\left(\omega^2 - \omega_c^2\right)}\phi,\\ v_{\omega z} &= 0.\end{aligned} \tag{20}$$

The nonlinear velocity $\mathbf{v}_\omega^{NL}$ due to ponderomotive force can be obtained from equation (20) by replacing φ by φ_p.

Using the equation of continuity, the linear and nonlinear density perturbations at ω, k can be obtained as

$$\begin{aligned}n_\omega^L &= \frac{\nabla\cdot\left\{n_0^0(y)v_\omega\right\}}{i\omega},\\ n_\omega^{NL} &= \frac{\nabla\cdot\left\{n_0^0(y)v_\omega^{NL}\right\}}{i\omega}.\end{aligned} \tag{21}$$

Using these in Poisson's equation, $\nabla^2\varphi = (e/\varepsilon_0)(n_\omega^L + n_\omega^{NL})$, we obtain

$$\phi = -\frac{\omega_p^2(y)}{\left(\omega^2 - \omega_c^2 - \omega_p^2(y)\right)}\phi_p. \tag{22}$$

The nonlinear current densities at the sidebands can be written as

$$\begin{aligned}\mathbf{J}_1^{NL} &= -\frac{1}{2}ne\mathbf{v}_0^*,\\ \mathbf{J}_2^{NL} &= -\frac{1}{2}ne\mathbf{v}_0.\end{aligned} \tag{23}$$

Using equation (23) in the wave equation, we obtain

$$\begin{aligned}\underline{\underline{D}}_1\cdot\mathbf{E}_1 &= \left(-k_1^2\underline{\underline{I}} + \mathbf{k}_1\mathbf{k}_1 + \frac{\omega_1^2}{c^2}\underline{\underline{\varepsilon}}_1\right)\cdot\mathbf{E}_1\\ &= \frac{k^2e\varphi\omega_p^2\omega_1E_{0y}}{2mc^2\omega_0\left(\omega^2 - \omega_c^2\right)\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)}\\ &\quad\cdot\left\{i\omega_0\omega_c\hat{x} + \left(\omega_0^2 - \omega_p^2\right)\hat{y}\right\},\end{aligned} \tag{24}$$

where $\underline{\underline{D}}_1 = -k_1^2\underline{\underline{I}} + \mathbf{k}_1\mathbf{k}_1 + (\omega_1^2/c^2)\underline{\underline{\varepsilon}}_1$, and

$$\begin{aligned}\underline{\underline{D}}_2\cdot\mathbf{E}_2 &= \left(-k_2^2\underline{\underline{I}} + \mathbf{k}_2\mathbf{k}_2 + \frac{\omega_2^2}{c^2}\underline{\underline{\varepsilon}}_2\right)\cdot\mathbf{E}_2\\ &= \frac{k^2e\phi\omega_p^2\omega_2E_{0y}}{2mc^2\omega_0\left(\omega^2 - \omega_c^2\right)\left(\omega_0^2 - \omega_c^2 - \omega_p^2\right)}\\ &\quad\cdot\left\{-i\omega_0\omega_c\hat{x} + \left(\omega_0^2 - \omega_p^2\right)\hat{y}\right\},\end{aligned} \tag{25}$$

where $\underline{\underline{D}}_2 = -k_2^2\underline{\underline{I}} + \mathbf{k}_2\mathbf{k}_2 + (\omega_2^2/c^2)\underline{\underline{\varepsilon}}_2$.

Equations (24) and (25) on replacing k_{1y} by $-\mathbf{i}\partial/\partial y$ for the sidebands can be written as

$$\frac{d^2E_{1y}}{d\xi_1^2}+\left(\lambda_1-\xi_1^2\right)E_{1y}=\frac{W_{01}^2e^{-\xi_1^2}}{\xi_{10}^2-\xi_1^2}\left(E_{1y}-\frac{\eta_1}{\eta_2}E_{2y}\right), \tag{26}$$

$$\frac{d^2E_{2y}}{d\xi_2^2}+\left(\lambda_2-\xi_2^2\right)E_{2y}=\frac{W_{02}^2e^{-\xi_2^2}}{\xi_{20}^2-\xi_2^2}\left(E_{2y}-\frac{\eta_2}{\eta_1}E_{1y}\right), \tag{27}$$

$$\begin{aligned}
\xi_1^2 &= \beta_1^{(1/2)}y^2,\\
\xi_2^2 &= \beta_2^{(1/2)}y^2,\\
\xi^2 &= \beta_1^{(1/2)}y^2 \approx \beta_2^{(1/2)}y^2 = \frac{\beta_1^{(1/2)}}{\beta_0^{(1/2)}}\xi_0^2 \approx \xi_0^2,\\
\xi_{10}^2 &= \frac{\left(\omega^2-\omega_c^2-\omega_{p0}^2\right)L_n^2\beta_1^{(1/2)}}{\omega_{p0}^2},\\
\xi_{20}^2 &= \frac{\left(\omega^2-\omega_c^2-\omega_{p0}^2\right)L_n^2\beta_2^{(1/2)}}{\omega_{p0}^2},\\
\alpha_1 &= \frac{\omega_1^2\left(\omega_1^2-\omega_c^2-\omega_{p0}^2\right)^2-k_{1x}^2c^2\left(\omega_1^2-\omega_c^2\right)\left(\omega_1^2-\omega_c^2-\omega_{p0}^2\right)-\omega_c^2\omega_{p0}^4}{c^2\left(\omega_1^2-\omega_c^2\right)\left(\omega_1^2-\omega_c^2-\omega_{p0}^2\right)},\\
\alpha_2 &= \frac{\omega_2^2\left(\omega_2^2-\omega_c^2-\omega_{p0}^2\right)^2-k_{2x}^2c^2\left(\omega_2^2-\omega_c^2\right)\left(\omega_2^2-\omega_c^2-\omega_{p0}^2\right)-\omega_c^2\omega_{p0}^4}{c^2\left(\omega_2^2-\omega_c^2\right)\left(\omega_2^2-\omega_c^2-\omega_{p0}^2\right)},\\
\beta_1 &= \frac{\omega_c^2\omega_{p0}^4\left(2\omega_1^2-2\omega_c^2-\omega_{p0}^2\right)+\omega_1^2\omega_{p0}^2\left(\omega_1^2-\omega_c^2-\omega_{p0}^2\right)^2}{c^2\left(\omega_1^2-\omega_c^2\right)\left(\omega_1^2-\omega_c^2-\omega_{p0}^2\right)^2L_n^2}\\
&= \frac{\omega_{p0}^2}{\omega_1^2}\left\{1+\frac{\omega_c^2\left(\omega_1^2-\omega_c^2\right)}{\left(\omega_1^2-\omega_{p0}^2-\omega_c^2\right)^2}\right\}\frac{\omega_1^2}{L_n^2c^2},\\
\beta_2 &= \frac{\omega_c^2\omega_{p0}^4\left(2\omega_2^2-2\omega_c^2-\omega_{p0}^2\right)+\omega_2^2\omega_{p0}^2\left(\omega_2^2-\omega_c^2-\omega_{p0}^2\right)^2}{c^2\left(\omega_2^2-\omega_c^2\right)\left(\omega_2^2-\omega_c^2-\omega_{p0}^2\right)^2L_n^2}\\
&= \frac{\omega_{p0}^2}{\omega_2^2}\left\{1+\frac{\omega_c^2\left(\omega_2^2-\omega_c^2\right)}{\left(\omega_2^2-\omega_{p0}^2-\omega_c^2\right)^2}\right\}\frac{\omega_2^2}{L_n^2c^2},\\
\lambda_1 &= \frac{\alpha_1}{\beta_1^{(1/2)}}=\left[\frac{\omega_1^2}{c^2}\left(1-\frac{\omega_{p0}^2}{\omega_1^2}\,\frac{\omega_1^2-\omega_{p0}^2}{\omega_1^2-\omega_{p0}^2-\omega_c^2}\right)-k_{0x}^2\right]\frac{1}{\beta_1^{(1/2)}},\\
\lambda_2 &= \frac{\alpha_2}{\beta_2^{(1/2)}}=\left[\frac{\omega_2^2}{c^2}\left(1-\frac{\omega_{p0}^2}{\omega_2^2}\,\frac{\omega_2^2-\omega_{p0}^2}{\omega_2^2-\omega_{p0}^2-\omega_c^2}\right)-k_{0x}^2\right]\frac{1}{\beta_2^{(1/2)}},\\
V_{01} &= \frac{eE_{0y}\left\{\omega_1\left(\omega_0^2-\omega_p^2\right)\left(\omega_1^2-\omega_c^2-\omega_p^2\right)-\omega_c^2\omega_p^2\omega_0\right\}}{2m\omega_0\omega_1^2\left(\omega_0^2-\omega_c^2-\omega_p^2\right)\left(\omega_1^2-\omega_c^2-\omega_p^2\right)},\\
V_{02} &= -\frac{eE_{0y}\left\{\omega_1\left(\omega_0^2-\omega_p^2\right)\left(\omega_1^2-\omega_c^2-\omega_p^2\right)-\omega_c^2\omega_p^2\omega_0\right\}}{2m\omega_0\omega_1^2\left(\omega_0^2-\omega_c^2-\omega_p^2\right)\left(\omega_1^2-\omega_c^2-\omega_p^2\right)}\,\frac{\eta_1}{\eta_2},\\
W_{01}^2 &= \frac{ek^2E_{0y}L_n^2\omega_p^4V_{01}\left\{\omega_p^2\omega_0\omega_c^2-\omega_1\left(\omega_1^2-\omega_c^2-\omega_p^2\right)\left(\omega_0^2-\omega_p^2\right)\right\}}{2mc^2\omega_0\omega_{p0}^2\left(\omega^2-\omega_c^2\right)\left(\omega_1^2-\omega_c^2-\omega_p^2\right)\left(\omega_0^2-\omega_c^2-\omega_p^2\right)},\\
W_{02}^2 &= \frac{ek^2E_{0y}L_n^2\omega_p^4V_{02}\left\{\omega_p^2\omega_0\omega_c^2+\omega_2\left(\omega_2^2-\omega_c^2-\omega_p^2\right)\left(\omega_0^2-\omega_p^2\right)\right\}}{2mc^2\omega_0\omega_{p0}^2\left(\omega^2-\omega_c^2\right)\left(\omega_2^2-\omega_c^2-\omega_p^2\right)\left(\omega_0^2-\omega_c^2-\omega_p^2\right)}.
\end{aligned} \tag{28}$$

We solve equations (26) and (27) by the first-order perturbation theory. In the absence of nonlinear terms, the eigenfunctions for the fundamental mode satisfying equations (26) and (27) are

$$E_{1,2} = A_{1,2}e^{-\xi_1^2/2}. \tag{29}$$

Corresponding eigenvalues are $\lambda_{1,2} = 1$. When the RHS of equations (26) and (27) are finite, we assume the eigenfunctions to remain unmodified; only the eigenvalues are changed a little. We substitute for $E_{1,2}$ from equation (29) in equations (26) and (27), multiply the resulting equations by $e^{-\xi^2/2}$, and integrate over ξ from $-\infty$ to ∞. Then, we obtain

$$(\lambda_1 - 1)A_1 = -W_{01}^2 \frac{Z(\xi_{10}\sqrt{2})}{\xi_{10}}\left(A_1 - \frac{\eta_1}{\eta_2}A_2\right), \tag{30}$$

$$(\lambda_2 - 1)A_2 = -W_{02}^2 \frac{Z(\xi_{10}\sqrt{2})}{\xi_{10}}\left(A_2 - \frac{\eta_2}{\eta_1}A_1\right), \tag{31}$$

where $Z(\xi_{10}\sqrt{2})$ is the plasma dispersion function. Equations (30) and (31) give the nonlinear dispersion relation

$$1 = -\frac{Z(\xi_{10}\sqrt{2})}{\xi_{10}}\left(\frac{W_{01}^2}{\lambda_1 - 1} + \frac{W_{02}^2}{\lambda_2 - 1}\right). \tag{32}$$

This nonlinear dispersion relation is a function of the DC magnetic field. The dispersion relation is modified if one changes the value of the DC magnetic field. In the absence of the DC magnetic field,

$$1 = -\frac{Z(\xi_{10}(\omega_c = 0)\sqrt{2})}{\xi_{10}(\omega_c = 0)}\left(\frac{W_{01}^2(\omega_c = 0)}{\lambda_1(\omega_c = 0) - 1} + \frac{W_{02}^2(\omega_c = 0)}{\lambda_2(\omega_c = 0) - 1}\right), \tag{33}$$

where $\beta_1^{1/4}(\omega_c = 0)y = \xi_1(\omega_c = 0)$, $\beta_2^{1/4}(\omega_c = 0)y = \xi_2(\omega_c = 0)$, $\beta_0^{1/4}(\omega_c = 0)y = \xi_0(\omega_c = 0)$,

$$\begin{aligned}
\xi_{10}(\omega_c = 0) &= \frac{\left(\omega^2 - \omega_{p0}^2\right)L_n^2\beta_1^{1/2}(\omega_c = 0)}{\omega_{p0}^2},\\
\lambda_1(\omega_c = 0) &= \left[\frac{\omega_1^2}{c^2}\left(1 - \frac{\omega_{p0}^2}{\omega_1^2}\frac{\omega_1^2 - \omega_{p0}^2}{\omega_1^2 - \omega_{p0}^2}\right) - k_{0x}^2\right]\frac{1}{\beta_1^{1/2}(\omega_c = 0)},\\
\lambda_2(\omega_c = 0) &= \left[\frac{\omega_2^2}{c^2}\left(1 - \frac{\omega_{p0}^2}{\omega_2^2}\frac{\omega_2^2 - \omega_{p0}^2}{\omega_2^2 - \omega_{p0}^2}\right) - k_{0x}^2\right]\frac{1}{\beta_2^{1/2}(\omega_c = 0)},\\
\beta_1(\omega_c = 0) &= \frac{\omega_{p0}^2}{L_n^2c^2} = \beta_2(\omega_c = 0) = \beta_0(\omega_c = 0),\\
V_{01}(\omega_c = 0) &= \frac{eE_{0y}\left\{\omega_1\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_p^2\right)\right\}}{2m\omega_0\omega_1^2\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_p^2\right)},\\
V_{02}(\omega_c = 0) &= -\frac{eE_{0y}\left\{\omega_1\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_p^2\right)\right\}}{2m\omega_0\omega_1^2\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_p^2\right)}\frac{\eta_1(\omega_c = 0)}{\eta_2(\omega_c = 0)},\\
\eta_1(\omega_c = 0) &= \omega_1^2\left(\omega_1^2 - \omega_p^2\right)\left\{\omega_2\left(\omega_0^2 - \omega_p^2\right)\left(\omega_2^2 - \omega_p^2\right)\right\},\\
\eta_2(\omega_c = 0) &= \omega_2^2\left(\omega_2^2 - \omega_p^2\right)\left\{\omega_1\left(\omega_0^2 - \omega_p^2\right)\left(\omega_1^2 - \omega_p^2\right)\right\},\\
W_{01}^2(\omega_c = 0) &= \frac{ek^2E_{0y}L_n^2\omega_p^4V_{01}(\omega_c = 0)\left\{-\omega_1\left(\omega_1^2 - \omega_p^2\right)\left(\omega_0^2 - \omega_p^2\right)\right\}}{2mc^2\omega_0\omega_{p0}^2\omega^2\left(\omega_1^2 - \omega_p^2\right)\left(\omega_0^2 - \omega_p^2\right)},\\
W_{02}^2(\omega_c = 0) &= \frac{ek^2E_{0y}L_n^2\omega_p^4V_{02}(\omega_c = 0)\left\{\omega_2\left(\omega_2^2 - \omega_p^2\right)\left(\omega_0^2 - \omega_p^2\right)\right\}}{2mc^2\omega_0\omega_{p0}^2\omega^2\left(\omega_2^2 - \omega_p^2\right)\left(\omega_0^2 - \omega_p^2\right)}.
\end{aligned} \tag{34}$$

Taylor expanding $\lambda_1 - 1$ with respect to variables ω_1, k_1 and $\lambda_2 - 1$ with respect to variables ω_2, k_2 and substituting in equation (32), one obtains

$$1 = -\frac{Z(\xi_{10}\sqrt{2})}{\xi_{10}} \frac{W_{01}^2}{(\partial\lambda_0/\partial\omega_0)} \left[\frac{1}{(\omega - kv_{g0} + \Delta)} - \frac{1}{(\omega - kv_{g0} - \Delta)}\right], \tag{35}$$

where

$$\begin{aligned}
\frac{\partial\omega_0}{\partial k_0} &= -\frac{\partial\lambda_0/\partial k_0}{\partial\lambda_0/\partial\omega_0} = v_{g0},\\
\Delta &= \frac{\Delta'}{2(\partial\lambda_0/\partial\omega_0)},\\
\Delta' &= \frac{\partial^2\lambda_0}{\partial\omega_0^2}\omega^2 + \frac{\partial^2\lambda_0}{\partial k_0^2}k^2,\\
\frac{\partial\lambda_0}{\partial\omega_0} &= \frac{\partial(\lambda_0 - 1)}{\partial\omega_0} \cong \frac{1}{\sqrt{\beta_0}}\left[2\frac{\omega_0}{c^2} + \frac{2\omega_{p0}^2}{c^2}\frac{\Delta_1}{\Delta_2} - \frac{\omega_{p0}}{L_n c}\frac{\Delta_3}{\Delta_4}\right],\\
\frac{\partial^2\lambda_0}{\partial\omega_0^2} &= \frac{1}{\sqrt{\beta_0}}\left[\frac{2}{c^2} + \frac{2\omega_{p0}^2}{c^2}\frac{\Delta_2 d_1 - \Delta_1 d_2}{\Delta_2^2} - \frac{\omega_{p0}}{L_n c}\frac{\Delta_4 d_3 - \Delta_3 d_4}{\Delta_4^2}\right],\\
\frac{\partial\lambda_0}{\partial k_0} &= -\frac{2k_0}{\sqrt{\beta_0}},\\
\frac{\partial^2\lambda_0}{\partial k_0^2} &= -\frac{2}{\sqrt{\beta_0}},\\
\Delta_1 &= \omega_0\left\{\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)\left(\omega_{p0}^2 - 2\omega_0^2\right) - \left(\omega_{p0}^2 - \omega_0^2\right)\left(2\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)\right\},\\
\Delta_2 &= \omega_0^2\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2,\\
\Delta_3 &= -\omega_0\omega_c^2\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right),\\
\Delta_4 &= \left\{\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2 + \omega_c^2\left(\omega_0^2 - \omega_c^2\right)\right\}^{1/2}\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2,\\
\frac{\partial\Delta_1}{\partial\omega_0} &= d_1 = \left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)\left(\omega_{p0}^2 - 6\omega_0^2\right) - 2\omega_0^2\omega_{p0}^2 - \left(2\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)\left(\omega_{p0}^2 - 3\omega_0^2\right),\\
\frac{\partial\Delta_2}{\partial\omega_0} &= d_2 = 2\omega_0\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2,\\
\frac{\partial\Delta_3}{\partial\omega_0} &= d_3 = -\omega_c^2\left(3\omega_0^2 + \omega_{p0}^2 - \omega_c^2\right),\\
\frac{\partial\Delta_4}{\partial\omega_0} &= d_4 = 4\omega_0\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)\left\{\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2 + \omega_c^2\left(\omega_0^2 - \omega_c^2\right)\right\}^{1/2}\\
&\quad + \omega_0\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2\left(\omega_0^2 - \omega_{p0}^2\right)\left\{\left(\omega_0^2 - \omega_{p0}^2 - \omega_c^2\right)^2 + \omega_c^2\left(\omega_0^2 - \omega_c^2\right)\right\}^{-1/2}.
\end{aligned} \tag{36}$$

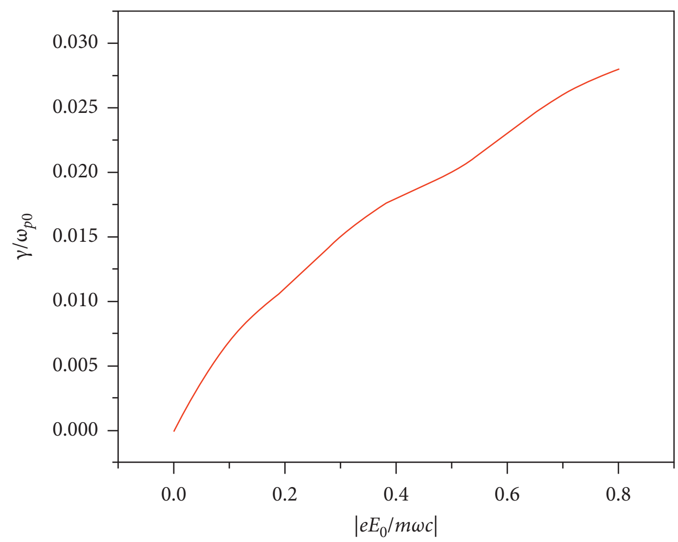

FIGURE 2: Normalized growth rate as a function of normalized amplitude of the pump wave in Raman forward scattering for $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, and $(\omega_c/\omega_{p0}) = 1$.

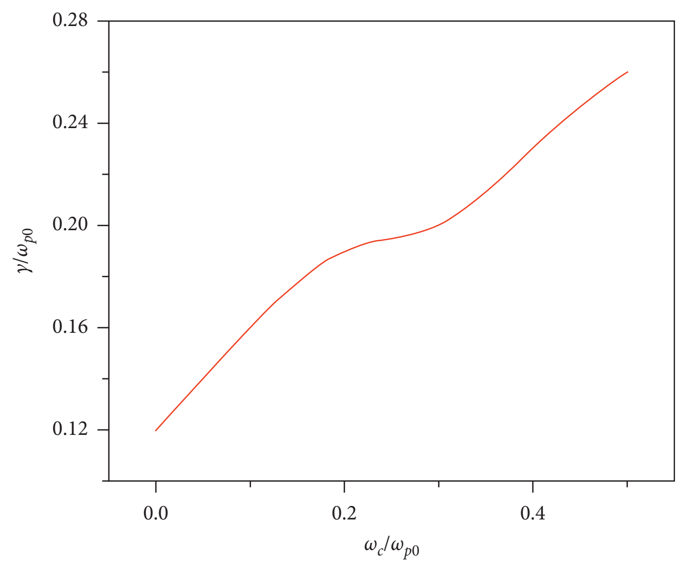

FIGURE 3: Normalized growth rate as a function of normalized cyclotron frequency for $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, and $|eE_0/m\omega c| = 0.1$.

Equation (35) gives

$$\left(\omega - kv_{g0}\right)^2 - \Delta^2 = \frac{2\Delta W_{01}^2 Z\left(\xi_{10}\sqrt{2}\right)}{(\partial\lambda_0/\partial\omega_0)\xi_{10}}. \tag{37}$$

We substitute the value of ξ_{10} and $Z(\xi_{10}\sqrt{2}) \approx -\sqrt{2}$, also write $\omega = \omega_r + i\gamma$, where $\omega_r = kv_{g0} = \sqrt{(\omega_c^2 + \omega_{p0}^2)}$, and obtain

$$(-1)^{1/2}\gamma^3 = \frac{2\sqrt{2}W_{01}^2\omega_{p0}\Delta}{(\partial\lambda_0/\partial\omega_0)L_n\beta_1^{1/4}}e^{2il\pi}. \tag{38}$$

Equation (38) gives the growth rate

$$\gamma = \left[\frac{2\sqrt{2}\,W_{01}^2\omega_{p0}\Delta}{(\partial\lambda_0/\partial\omega_0)L_n\beta_1^{1/4}}\right]^{1/3}\frac{\sin 2l\pi}{3}. \tag{39}$$

For $l = 1$, the growth rate turns out to be

$$\gamma = \frac{\sqrt{3}}{2}\left[\frac{2\sqrt{2}\,W_{01}^2\omega_{p0}\Delta}{(\partial\lambda_0/\partial\omega_0)L_n\beta_1^{1/4}}\right]^{1/3}. \tag{40}$$

In Figure 2, we have displayed the normalized growth rate (γ/ω_{p0}) of the Raman forward scattering instability as a function of normalized pump wave amplitude $|eE_0/m\omega c|$ for parameters: $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, and $(\omega_c/\omega_{p0}) = 1$. The growth rate increases with the amplitude of the laser pump. It does not go linearly but nearly as two-third powers of normalized laser amplitude. This is due to the fact that the width of the driven Langmuir mode is dependent on ponderomotive potential hence on pump amplitude. In Figure 3, we have displayed the normalized growth rate (γ/ω_{p0}) of the Raman forward scattering as a function of normalized static magnetic field (ω_c/ω_{p0}) for $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, and $|eE_0/m\omega c| = 0.1$. In the absence of DC magnetic field, the growth rate (γ/ω_{p0}) is minimum. It rises with magnetic field. The magnetic field raises the frequency of the driven plasma wave and brings in cyclotron effects in nonlinear coupling leading to enhancement of growth rate.

3. Discussion

The plasma channel with a parabolic density profile localizes the electromagnetic eigenmodes involved in the SRFS process within a width of the order $(cL_n/\omega_{p0})^{1/2}$. The Langmuir wave is more strongly localized, thus limiting the region of parametric interaction and reducing the growth rate. The static magnetic field modifies the electron response to these eigenmodes and significantly influences the nonlinear coupling. In the limit when the normalized growth rate $(\gamma/\omega_{p0}) > kv_{th}$, one may neglect thermal effects. The growth rate roughly scales as $E_{0y}^{2/3}$ with pump amplitude and goes linearly with ambient magnetic field. For typical parameters, $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, $(\omega_c/\omega_{p0}) = 0.5$, and $|eE_0/m\omega c| = 0.1$, the normalized growth rate is 0.26 whereas for $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, $(\omega_c/\omega_{p0}) = 1$, and $|eE_0/m\omega c| = 0.8$, the normalized growth rate is 0.028.

In the earlier work by Liu et al. [37] for the parameters $(\omega_{p0}/\omega_0) = 0.2$ and $(L_n\omega_{p0}/c) = 8$ $|v_0/c| = 0.6$, the normalized growth rate is 0.06 whereas in our work for the parameters $(\omega_{p0}/\omega_0) = 0.2$, $(L_n\omega_{p0}/c) = 8$, $(\omega_c/\omega_{p0}) = 1$, and $|eE_0/m\omega c| = 0.6$, the normalized growth rate is 0.025 Clearly, our results are in good agreement with the earlier work by Liu et al.

Data Availability

The data used to support the findings of this study are available upon request from the authors.

Conflicts of Interest

The authors declare that they have no conflicts of interest.

References

[1] S. P. Regan, D. K. Bradley, A. V. Chirokikh et al., "Laser-plasma interactions in long-scale-length plasmas under direct-drive national ignition facility conditions," *Physics of Plasmas*, vol. 6, no. 5, p. 2072, 1999.

[2] P. Sprangle, E. Esarey, and A. Ting, "Nonlinear theory of intense laser-plasma interactions," *Physical Review Letters*, vol. 64, 1990.

[3] J. R. Fontana and R. H. Pantell, "A high-energy, laser accelerator for electrons using the inverse Cherenkov effect," *Journal of Applied Physics*, vol. 54, no. 8, p. 4285, 1983.

[4] P. Sprangle, E. Esarey, A. Ting, and G. Joyce, "Laser wakefield acceleration and relativistic optical guiding," *Applied Physics Letters*, vol. 53, no. 22, p. 2146, 1988.

[5] P. K. Shukla, "Generation of wakefields by elliptically polarized laser pulses in a magnetized plasma," *Physics of Plasmas*, vol. 6, no. 4, p. 1363, 1999.

[6] H. Lin, Li-M. Chen, and J. C. Kieffer, "Harmonic generation of ultra-intense pulses in underdense plasma," *Physical Review E*, vol. 65, Article ID 036414, 2002.

[7] J. Zhou, J. Peatross, M. M. Murnane, H. C. Kapteyn, and I. P. Christov, "Enhanced high-harmonic generation using 25 fs laser pulses," *Physical Review Letters*, vol. 76, no. 5, p. 752, 1996.

[8] P. Amendt, D. C. Eder, and S. C. Wilks, "X-ray lasing by optical-field-induced ionization," *Physical Review Letters*, vol. 66, no. 20, p. 2589, 1991.

[9] E. Michael, W. L. Kruer, S. C. Wilks, J. Woodworth, E. M. Campbell, and M. D. Perry, "Ignition and high gain with ultra-powerful lasers," *Physics of Plasmas*, vol. 1, p. 1626, 1994.

[10] C. Deutsch, H. Furukawa, K. Mima, M. Murakami, and K. Nishihara, "Interaction physics of the fast ignitor concept," *Physical Review Letters*, vol. 77, no. 12, p. 2483, 1996.

[11] C. Gahn, G. D. Tsakiris, A. Pukhov et al., "Multi-MeV electron beam generation by direct laser acceleration in high-density plasma channels," *Physical Review Letters*, vol. 83, no. 23, p. 4772, 1999.

[12] A. Ting, E. Esarey, and P. Sprangle, "Nonlinear wake-field generation and relativistic focusing of intense laser pulses in plasmas," *Physics of Fluids B: Plasma Physics*, vol. 2, no. 6, p. 1390, 1990.

[13] T. M. Antonsen and P. Mora, "Self-focusing and Raman scattering of laser pulses in tenuous plasmas," *Physical Review Letters*, vol. 69, no. 15, p. 2204, 1992.

[14] G.-Z. Sun, E. Ott, Y. C. Lee, and P. Guzdar, "Self-focusing of short intense pulses in plasmas," *Physics of Fluids*, vol. 30, no. 2, p. 526, 1987.

[15] N. E. Andreev, V. I. Kirsanov, and L. M. Gorbunov, "Stimulated processes and self-modulation of a short intense laser pulse in the laser wake-field accelerator," *Physics of Plasmas*, vol. 2, no. 6, p. 2573, 1995.

[16] S. Dalui, A. Bandyopadhyay, and K. P. Das, "Modulational instability of ion acoustic waves in a multi-species collisionlessunmagnetized plasma consisting of nonthermal and isothermal electrons," *Physics of Plasmas*, vol. 24, Article ID 042305, 2017.

[17] P. Jha, S. R. Gopal, and A. K. Upadhyay, "Pulse distortion and modulation instability in laser plasma interaction," *Physics of Plasmas*, vol. 16, Article ID 013107, 2009.

[18] N. Sepehri Javan, "Competition of circularly polarized laser modes in the modulation instability of hot magneto plasma," *Physics of Plasmas*, vol. 20, Article ID 012120, 2013.

[19] H. Gharaee, S. Afghah, and H. Abbasi, "Modulational instability of ion-acoustic waves in plasmas with super thermal electrons," *Physics of Plasmas*, vol. 18, Article ID 032116, 2011.

[20] T. S. Gill, H. Kaur, S. Bansal, N. S. Saini, and P. Bala, "Modulational instability of electron-acoustic waves: an application to auroral zone plasma," *The European Physical Journal D*, vol. 41, no. 1, pp. 151–156, 2007.

[21] P. Jha, P. Kumar, G. Raj, and A. K. Upadhyaya, "Modulation instability of laser pulse in magnetized plasma," *Physics of Plasmas*, vol. 12, no. 12, p. 123104, 2005.

[22] H.-Y. Chen, S.-Q. Liu, and X.-Q. Li, "Modulation instability by intense laser beam in magnetized plasma," *Optik*, vol. 122, no. 7, pp. 599–603, 2011.

[23] P. Mora, D. Pesme, A. Heron, G. Laval, and N. Silvestre, "Modulational instability and its consequences for the beat-wave accelerator," *Physical Review Letters*, vol. 61, p. 1611, 1998.

[24] I. Kourakis, P. K. Shukla, and G. Morfill, "Modulational instability and localized excitations involving two nonlinearly coupled upper-hybrid waves in plasmas," *New Journal of Physics*, vol. 7, p. 153, 2005.

[25] P. Sharma, "Stimulated Raman scattering of ultra intense hollow Gaussian beam in relativistic plasma," *Laser and Particle Beams*, vol. 33, no. 3, p. 489, 2015.

[26] U. Verma and A. K. Sharma, "Nonlinear electromagnetic Eigen modes of a self created magnetized plasma channel and its stimulated Raman scattering," *Laser and Particle Beams*, vol. 29, no. 4, p. 471, 2011.

[27] A. Vyas, S. Sharma, R. K. Singh, R. P. Sharma, and R. P. Sharma, "Effect of the axial magnetic field on coexisting stimulated Raman and Brillouin scattering of a circularly polarized beam," *Laser and Particle Beams*, vol. 35, no. 1, p. 19, 2017.

[28] R. L. Berger, C. H. Still, E. A. Williams, and A. B. Langdon, "On the dominant and subdominant behavior of stimulated Raman and Brillouin scattering driven by nonuniform laser beams," *Physics of Plasmas*, vol. 5, no. 12, p. 4337, 1998.

[29] R. K. Kirkwood, B. J. Macgowan, D. S. Montgomery et al., "Effect of ion-wave damping on stimulated Raman scattering in high-ZLaser-produced plasmas," *Physical Review Letters*, vol. 77, no. 13, p. 2706, 1996.

[30] A. Vyas, R. K. Singh, R. P. Sharma, K. Ram, and R. P. Sharma, "Study of coexisting stimulated Raman and Brillouin scattering at relativistic laser power," *Laser and Particle Beams*, vol. 32, no. 4, p. 657, 2014.

[31] J. G. Moreau, E. D'Humieres, R. Nuter, and V. T. Tikhonchuk, "Stimulated Raman scattering in the relativistic regime in near-critical plasmas," *Physical Review E*, vol. 95, Article ID 013208, 2017.

[32] C. Grebogi and C. S. Liu, "Brillouin and Raman scattering of an extraordinary mode in magnetized plasma," *Physics of Fluids*, vol. 23, p. 7, 1980.

[33] H. C. Barr, T. J. M. Boyd, G. A. Gardner, and R. Rankin, "Raman backscatter from an inhomogeneous magnetized plasma," *Physical Review Letters*, vol. 53, no. 5, p. 462, 1984.

[34] S. P. Obenschain, C. J. Pawley, A. N. Mostovych et al., "Reduction of Raman scattering in a plasma to convective levels using induced spatial incoherence," *Physical Review Letters*, vol. 62, no. 7, p. 768, 1989.

[35] C. Rousseaux, G. Malka, J. L. Miquel, F. Amiranoff, S. D. Baton, and P. Mounaix, "Experimental validation of the linear theory of stimulated Raman scattering driven by a 500-fs laser pulse in a preformed underdense plasma," *Physical Review Letters*, vol. 74, no. 23, p. 4655, 1995.

[36] G. Shvets and X. Li, "Raman forward scattering in plasma channels," *Physics of Plasmas*, vol. 8, no. 1, p. 8, 2001.

[37] C. S. Liu and V. K. Tripathi, "Stimulated Raman scattering in a plasma channel," *Physics of Plasmas*, vol. 3, no. 9, p. 3410, 1996.

[38] W. B. Mori, C. D. Decker, D. E. Hinkel, and T. Katsouleas, "Raman forward scattering of short-pulse high-intensity lasers," *Physical Review Letters*, vol. 72, no. 10, p. 1482, 1994.

[39] A. Panwar, A. Kumar, and C. M. Ryu, "Stimulated Raman forward scattering of laser in a pre-formed plasma channel," *Laser and Particle Beams*, vol. 30, no. 4, p. 605, 2012.

[40] V. Sajal, A. Panwar, and V. K. Tripathi, "Relativistic forward stimulated Raman scattering of a laser in a plasma channel," *Physica Scripta*, vol. 74, no. 4, p. 484, 2006.

[41] K. Hassoon, H. Salih, and V. K. Tripathi, "Stimulated Raman forward scattering of a laser in a plasma with transverse magnetic field," *Physica Scripta*, vol. 80, Article ID 065501, 2009.

[42] D. N. Gupta, P. Yadav, D. G. Jang, M. S. Hur, H. Suk, and K. Avinash, "Onset of stimulated Raman scattering of a laser in the presence of hot drifting electrons," *Physics of Plasmas*, vol. 22, Article ID 052101, 2015.

[43] Z. J. Liu, S.-P. Zhu, L. H. Cao, C. Y. Zheng, X. T. He, and Y. Wang, "Enhancement of backward Raman scattering by electron-ion collisions," *Physics of Plasmas*, vol. 16, no. 11, Article ID 112703, 2009.

[44] S. Y. Kalmykov and G. Shvets, "Stimulated Raman backscattering of laser radiation in deep plasma channels," *Physics of Plasmas*, vol. 11, no. 10, p. 4686, 2004.

[45] A. Paknezhad and D. Dorranian, "Nonlinear backward Raman scattering in the short laser pulse interaction with a cold underdense transversely magnetized plasma," *Laser and Particle Beams*, vol. 29, no. 3, p. 373, 2011.

[46] S. Kaur and A. K. Sharma, "Stimulated Raman scattering of a laser in a magnetized plasma channel," *Physics of Plasmas*, vol. 18, Article ID 092105, 2011.

[47] A. Paknezhad and D. Dorranian, "Third harmonic stimulated Raman backscattering of laser in a magnetized plasma," *Physics of Plasmas*, vol. 20, Article ID 092108, 2013.