NONLINEAR OPTICAL PHENOMENA

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O he he f he SBS li e id h i a i - ecie la ma

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$_{\square}H_{1-\square}$ plasmas with various relative carbon concentrations Y and various degrees of nonisothermicity, defined as the ratio of electron and ion temperatures. We found the dependences of the SBS linewidth on the pump field intensity, the carbon concentration and the degree of nonisothermicity. We also found the dependence of the SBS linewidth on the damping decrement of the ion-sound waves determining the SBS development. The dependence of the SBS linewidth on the degree of nonisothermicity was studied in the C_5H_{12} plasma. We found that for a low nonisothermicity, the SBS linewidth is mainly determined by the damping of the slow ion-sound wave, while for a high nonisothermicity, the fast ion-sound wave is damped more weakly and determines the SBS spectrum.

We will study SBS in a two ion-species plasma containing electrons using the dispersion relation [10, 11]

$$\frac{1}{1+\delta\varepsilon_{\rm i}(\omega,\square)}+\frac{1}{\delta\varepsilon_{\rm e}(\omega,\square)}$$

$$=\frac{\square^2|[\boldsymbol{k}-\boldsymbol{k}_0,\boldsymbol{v}_{\square}]|^2}{4(\boldsymbol{k}-\boldsymbol{k}_0)^2[\square^2(\boldsymbol{k}-\boldsymbol{k}_0)^2-(\omega-\omega_0)^2\varepsilon_{\rm tr}(\omega-\omega_0,\boldsymbol{k}-\boldsymbol{k}_0)]}, \quad (1)$$

where ω_0 and $\boldsymbol{k}_0$ are the frequency and the wave vector of the incident electromagnetic wave; $\omega_0-\omega$ and $\boldsymbol{k}_0-\boldsymbol{k}$ are the frequency and the wave vector of the scattered wave; ω and $\boldsymbol{k}$ are the frequency and the wave vector of the acoustic mode involved in SBS; $\boldsymbol{v}_{\square}=\square\boldsymbol{E}_0/\square_{\rm e}\omega_0$ is the velocity of electron oscillations in the field $\square_0$ of the incident wave; $\square$ and $\square_{\square}$ are the electron charge and mass; $\delta\varepsilon_{\rm e}$, and $\delta\varepsilon_{\rm i}$ are the longitudinal permittivities of electrons and ions;

$$\varepsilon_{\rm tr}(\omega-\omega_0,\boldsymbol{k}-\boldsymbol{k}_0)=1-\frac{\omega_{\rm Le}^2}{(\omega-\omega_0)^2}\left(1-{\rm i}\frac{\square_{\rm ei}}{\omega-\omega_0}\right) \quad (2)$$

is the transverse permittivity; $\omega_{\rm Le}$ is the Langmuir frequency of electrons;

$$\nu_{\rm ei}=\frac{4}{3}\frac{\sqrt{2\pi}{\rm e}^4\Lambda}{\square_{\rm e}^2 v_{\square\rm e}^3}\sum_{\alpha=1,2}\square_\alpha^2\square_\alpha \quad (3)$$

is the effective frequency of electron – ion collisions in a two ion-species plasma; Λ is the Coulomb logarithm; $\square_\alpha$ and $\square_\alpha$ are the degree of ionisation and the density of ions of type α; and $v_{\square\rm e}$ is the thermal velocity of electrons.

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To a good accuracy, the wave vector of the scattered electromagnetic wave is the root of the equation (see, for example, Refs [10, 11])

$$(\omega-\omega_0)^2\mathrm{Re}\left[\varepsilon_{\mathrm{tr}}(\omega-\omega_0,\boldsymbol{k}-\boldsymbol{k}_0)\right]-c^2(\boldsymbol{k}-\boldsymbol{k}_0)^2=0, \quad (4)$$

which is also the dispersion equation for the scattered electromagnetic wave. By retaining small terms of the order of ω, for example, for back scattering, we take into account small corrections of the order of the ratio of the velocity of sound to the velocity of light, which are insignificant for determining the wave vector. However, even smaller deviations from the resonance condition (4), which are possible in this case, may reveal important properties of the SBS frequency spectrum without any significant variation in the wave vector. In this connection, we can write, as usual [10, 11],

$$(\omega-\omega_0)^2\mathrm{Re}\left[\varepsilon_{\mathrm{tr}}(\omega-\omega_0,\boldsymbol{k}-\boldsymbol{k}_0)\right]-c^2(\boldsymbol{k}-\boldsymbol{k}_0)^2=2\omega_0\Delta, \quad (5)$$

where Δ is the frequency detuning of the parametric resonance. In our case, the small ratio Δ/ω is determined by the small parameter

$$\eta_E=\frac{\gamma_E}{k v_{T1}}=10^{-5}(\theta A_1)^{1/2}\,\frac{2k_0}{k}\left(\frac{\Lambda}{10}\right)\times\left(\frac{n_e}{10^{19}\mathrm{cm}^{-3}}\right)^2\frac{\lambda^3}{T_e^2}\,\frac{Z_1^2N_1+Z_2^2N_2}{Z_1N_1+Z_2N_2}, \quad (6)$$

where

$$\gamma_E=\nu_{\mathrm{ei}}\frac{\omega_{\mathrm{Le}}^2}{2\omega_0^2}. \quad (7)$$

Here, λ is measured in micrometers and T_e in kiloelectronvolts; $v_{T\alpha}$ is the thermal velocity of ions of type α; $\lambda=2\pi/k_0$ is the wavelength of the pump field; $\theta=T_e/T_i$ is ratio of the electron temperature T_e to the ion temperature T_i; and A_1 is the atomic number of the lighter ions. For typical experimental conditions [13] in the case of backward SBS ($k\simeq 2k_0$), it follows from formula (7) that $\eta_E\simeq 0.003$ for the C_5H_{12} plasma with $\theta\simeq 3, Z_1=1, A_1=1, Z_2=6, T_e=$ 3 keV, $n_e=10^{21}$ cm^{-3}, and $\lambda=0.357$ μm. In this case, the effect of frequency detuning on the wave vector is characterised by a quantity of the order of 10^{-6}.

We will study the development of the SBS instability in time assuming that $\omega\to\omega+\mathrm{i}\gamma$ in Eqns (1) and (2). We also assume that the increment γ of the SBS instability is small compared to the frequency ω of sound. It should be emphasised that formulas (4) and (5), as well as the subsequent formulas, contain the real frequency ω.

By introducing the notation

$$F(\omega,k)=\mathrm{Re}\left[\frac{1}{1+\delta\varepsilon_{\mathrm{i}}(\omega,k)}+\frac{1}{\delta\varepsilon_{\mathrm{e}}(\omega,k)}\right],$$

$$G(\omega,k)=\mathrm{Im}\left[\frac{1}{1+\delta\varepsilon_{\mathrm{i}}(\omega,k)}+\frac{1}{\delta\varepsilon_{\mathrm{e}}(\omega,k)}\right],$$

we can write in compact form the following system of equations for determining the increment and frequency detuning:

$$F(\omega,k)-\gamma\frac{\partial G(\omega,k)}{\partial\omega}=-\frac{k^2\left|[\boldsymbol{k}-\boldsymbol{k}_0,\boldsymbol{v}_E]\right|^2}{8\omega_0(\boldsymbol{k}-\boldsymbol{k}_0)^2}\frac{\Delta}{\Delta^2+(\gamma+\gamma_E)^2}, \quad (8)$$

$$G(\omega,k)+\gamma\frac{\partial F(\omega,k)}{\partial\omega}=-\frac{k^2\left|[\boldsymbol{k}-\boldsymbol{k}_0,\boldsymbol{v}_E]\right|^2}{8\omega_0(\boldsymbol{k}-\boldsymbol{k}_0)^2}\frac{\gamma+\gamma_E}{\Delta^2+(\gamma+\gamma_E)^2}. \quad (9)$$

In our case, where the dissipation of sound waves is collisionless, the expressions for the permittivities determining the system of equations (8) and (9) have the form [14]

$$\delta\varepsilon_{\mathrm{i}}(\omega,k)=\sum_{\alpha=1,2}\delta\varepsilon_\alpha(\omega,k),$$

$$\delta\varepsilon_{\mathrm{e}}(\omega,k)=\frac{1}{k^2r_{\mathrm{De}}^2}\left(1+\mathrm{i}\sqrt{\pi/2}\,\frac{\omega}{kv_{Te}}\right), \quad (10)$$

$$\delta\varepsilon_\alpha(\omega,k)=\frac{1}{k^2r_{\mathrm{D}\alpha}^2}\left[1-\frac{\sqrt{2}\omega}{kv_{T\alpha}}D\left(\frac{\omega}{\sqrt{2}kv_{T\alpha}}\right)+\mathrm{i}\frac{\sqrt{\pi/2}\,\omega}{kv_{T\alpha}}\exp\left(-\frac{\omega^2}{2k^2v_{T\alpha}^2}\right)\right], \quad (11)$$

where r_{De} and $r_{\mathrm{D}\alpha}$ are the Debye radii of electrons and ions of type α, respectively, and $D(x)=\exp(-x^2)\int_0^x \mathrm{d}t\exp(t^2)$ is the Dawson integral.

Equations (8) and (9) lead to the following expression for the frequency detuning:

$$\Delta=(\gamma+\gamma_E)\frac{F(\omega,k)-\gamma\partial G(\omega,k)/\partial\omega}{G(\omega,k)+\gamma\partial F(\omega,k)/\partial\omega}. \quad (12)$$

Then, we obtain the following equation for γ:

$$(\gamma+\gamma_E)\frac{\left[F(\omega,k)-\gamma\partial G(\omega,k)/\partial\omega\right]^2+\left[G(\omega,k)+\gamma\partial F(\omega,k)/\partial\omega\right]^2}{G(\omega,k)+\gamma\partial F(\omega,k)/\partial\omega}=-\frac{k^2\left|[\boldsymbol{k}-\boldsymbol{k}_0,\boldsymbol{v}_E]\right|^2}{8\omega_0(\boldsymbol{k}-\boldsymbol{k}_0)^2}. \quad (13)$$

This equation determines the dependence of the SBS instability increment, describing the amplification of the scattered radiation with frequency $\omega_0-\omega$ in time, on the pump-wave intensity.

To determine the effect of the ion composition of plasma on the spectral properties of SBS, we will assume, while obtaining numerical estimates and plotting the graphs, that the electron density is the same in plasmas of various compositions considered here. Taking into account the electric neutrality of plasma ($n_e=Z_1N_1+Z_2N_2$), we can conveniently write the frequency (3) of electron–ion collisions in the form

$$\nu_{\mathrm{ei}}\equiv\nu_0\frac{Z_1^2N_1+Z_2^2N_2}{Z_1N_1+Z_2N_2}, \quad (14)$$

where the frequency $\nu_0=(4\sqrt{2\pi}/3)(e^4\Lambda n_e)/(m_e^2v_{Te}^3)$ is independent of the ion composition of the plasma, but depends on the electron density. The frequency ν_0 corresponds to the frequency of electron collisions in pure hydrogen plasma. Taking this into account, it is convenient to introduce the dimensionless intensity of the pump wave

$$q=\frac{\omega_0}{\nu_0}\frac{\left|[\boldsymbol{k}-\boldsymbol{k}_0,\boldsymbol{v}_E]\right|^2}{4(\boldsymbol{k}-\boldsymbol{k}_0)^2v_{Te}^2} \quad (15)$$

for a comparative description of the effect of the pump field on the SBS linewidth in plasmas with various concentrations of ions of first and second types. This intensity is independent of the ion composition of the plasma and makes it possible to introduce a unified scale of pump intensities for a comparative description of the spectral properties of SBS in different plasmas. It is also convenient to introduce the dimensionless frequency detuning $\delta = \Delta/\gamma_{\square}$, dimensionless velocity of sound $z = \omega/\square v_{\square 1}$, and dimensionless increment $\eta = \gamma/\square v_{\square 1}$ of the SBS instability. In this case, Eqn (13) can be written in the form

$$(\eta + \eta_{\square})\frac{\left[\Phi(z)+\eta \mathrm{d}\Psi(z)/\mathrm{d}z\right]^2 + \left[\Psi(z)-\eta \mathrm{d}\Phi(z)/\mathrm{d}z\right]^2}{\Psi(z)-\eta \mathrm{d}\Phi(z)/\mathrm{d}z} = I, \quad (16)$$

where

$$\Phi(z) = (\xi_1 + \xi_2)\left[1 + \frac{\Phi_1(z)}{\square}\right];$$

$$\Psi(z) = (\xi_1 + \xi_2)\left[\left(\frac{\pi \square_e}{2\square_1}\right)^{1/2}\frac{\square}{\sqrt{\square}} + \frac{\Phi_2(z)}{\square}\right];$$

$$\Phi_\alpha(z) = \frac{\square_\alpha(z)}{\square_1^2(z) + \square_2^2(z)};$$

$$\square_1(z) = \xi_1\left[1 - \sqrt{2}\square\square(z/\sqrt{2})\right] + \xi_2\left[1 - \sqrt{2}\square\square\square(\square z/\sqrt{2})\right];$$

$$\square_2(z) = \sqrt{\pi/2}\, z\left[\xi_1 \exp(-z^2/2) + \xi_2 \square \exp(-\square^2 z^2/2)\right];$$

$$\xi_1 = \square_1\left\{1 + [\square_2\square_2/(\square_1\square_1)]\right\}^{-1};$$

$$\xi_2 = \square_2\left\{1 + [\square_1\square_1/(\square_2\square_2)]\right\}^{-1}; \square = \left(\square_2/\square_1\right)^{1/2};$$

$\square_\alpha$ being the mass of ions of type α. We can now write the following expression for frequency detuning instead of Eqn (12):

$$\delta = \left(1 + \frac{\eta}{\eta_{\square}}\right)\frac{\Phi(z)+\eta \mathrm{d}\Psi(z)/\mathrm{d}z}{\Psi(z)-\eta \mathrm{d}\Phi(z)/\mathrm{d}z}. \quad (17)$$

At the boundary of the SBS instability region, the increment $\eta = 0$ according to the traditional approach used in the theory of parametric instabilities [11]. Equation (16) then takes the form

$$I_b(z) = I, \quad (18)$$

where

$$I_b(z) = \frac{\Phi^2(z) + \Psi^2(z)}{\Psi(z)}. \quad (19)$$

We call the function $I_b(z)$ the boundary intensity. This term is introduced because Eqn (18) describes the SBS linewidth for a given intensity I of the pump wave. It is obvious that the increment is small near the line boundary, and Eqn (16) leads to the following extremely simple expression for the increment in the approximation linear in the pump intensity:

$$\eta = \eta_{\square}\left(\frac{I}{I_b(z)} - 1\right). \quad (20)$$

One can see from (20) that, for the velocities of sound satisfying the condition $I_b(z) < I$, the quantity η is positive, and a scattered wave as well as the ion-sound wave associated with it can be excited. In other words, Eqn (18) determines the width of the spectral line at its base. It should be emphasised that, away from the SBS instability boundaries, the dependence of the increment on I may differ strongly from that described by formula (20). However, this difference is insignificant for determining the linewidth at the base. The minimum value of the function $I_b(z)$ attained for the velocity of sound $z = z_{th}$ determines the threshold intensity I_{th} of the SBS instability excitation. One can see from (20) that the maximum increment for a given I determines the threshold velocity of sound z_{th} above the SBS instability threshold ($I > I_{th}$); i.e., the scattered radiation with a frequency shift determined by z_{th} is amplified most efficiency.

By way of illustration of the SBS spectral properties slightly above the instability threshold, Fig. 1 shows the dependences of the boundary of the SBS instability region (whose size determines the SBS linewidth at the base) on the pump intensity for completely ionised plasmas of various compositions. Let us first discuss a pure carbon plasma with a degree of nonisothermicity $\square = 4$ and the degree of ionisation of carbon ions $\square_C = 6$. Curve (*1*) corresponds to such a plasma. Note that the dimensionless velocity of sound is measured in the units of thermal velocity of hydrogen that does not exist in such a plasma. This is done for the convenience of comparing the dependences obtained for pure carbon plasma and for plasmas containing hydrogen. The upper branch of curve (*1*), referred to as $z_{max}(I)$, describes the so-called long-wavelength boundary of the SBS spectrum, while the lower branch $z_{min}(I)$ determines the short-wavelength boundary. At the SBS threshold, we have $z_{min}(I_{th}) = z_{max}(I_{th}) = z_{th}$. One can see from Fig. 1 that above the threshold, the width $\Delta z(I) = z_{max}(I) - z_{min}(I)$ of the SBS line at the base in a carbon plasma increases insignificantly with the pump field intensity; for a fourfold increase in the threshold intensity ($I = 4I_{th}$), the linewidth is much smaller than the frequency shift of the scattered radiation and is only $\sim 0.03 z_{th}$. Therefore, we can speak

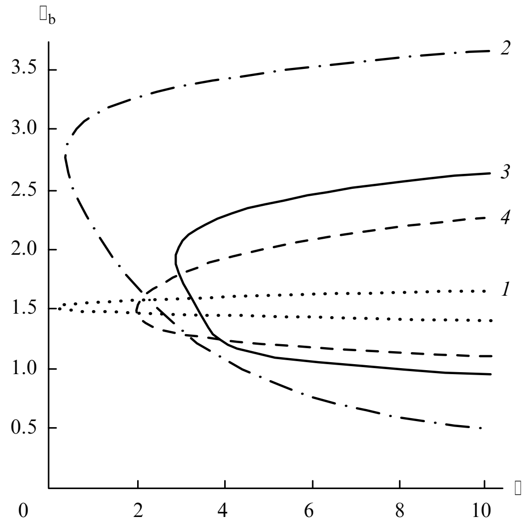

Figure 1. Dependence of the boundary z_b of the SBS instability region on the pump intensity I in plasmas with different compositions: C (*1*), H (*2*), C_5H_{12} (*3*), and CH (*4*).

of a narrow SBS spectral line in a pure carbon plasma, with a well-defined frequency $\omega_0 - \omega$ and a frequency shift $\omega = \square v_{\square H} \square_{th}$ that is close to the frequency $\square v_{sC}$ of the conventional weakly damping ion sound in a carbon plasma, where $v_{sC} = (\square_C \square_B \square_e / \square_C)^{1/2}$ is the velocity of sound, and $\square_B$ is the Boltzmann constant. This is due to the trivial fact that the damping decrement γ_s of ion sound in a carbon plasma without pumping is smaller than the ion-sound frequency. Indeed, for $\square = 4$, the decrement γ_s is determined completely by the electron Landau damping and $\gamma_s/(\square v_{sC}) = (\pi \square_C \square_e / 8 \square_C)^{1/2} \approx 0.01$, while the ion damping of sound is negligibly small. Taking this fact into account, we considerably simplify Eqn (19). Assuming also that the phase velocity of sound waves involved in stimulating scattering is higher than the thermal velocity of carbon ions, we can write Eqn (18) for the boundary of the SBS instability region in the dimensional form

$$\frac{(\omega^2 - \square^2 v_{sC}^2)^2 + (2\gamma_s \omega)^2}{2\gamma_s \omega} = \square^2 v_{sC}^2 \frac{\omega_0}{\nu_{eC}} \frac{|[\boldsymbol{k} - \boldsymbol{k}_0, \boldsymbol{v}_\square]|^2}{4(\boldsymbol{k} - \boldsymbol{k}_0)^2 v_{\square e}^2}, \tag{21}$$

where ν_{eC} is the frequency of electron–ion collisions in a pure carbon plasma. In the linear approximation in the small parameter $\gamma_s / \square v_{sC}$, we can find from Eqn (21) the threshold of time evolution of instability, which is well known in the SBS theory [10]:

$$\frac{|[\boldsymbol{k} - \boldsymbol{k}_0, \boldsymbol{v}_\square]|^2_{th}}{4(\boldsymbol{k} - \boldsymbol{k}_0)^2 v_{\square e}^2} = 2 \frac{\gamma_s}{\square v_{sC}} \frac{\square_{eC}}{\omega_0}, \tag{22}$$

and is attained for the frequency $\omega_{th} \simeq \square v_{sC}$ coinciding with the frequency of the weakly damped ion sound in such an approximation. For the long-(short)-wavelength boundary of the SBS instability region, we obtain from (21)

$$\omega_{max(min)} = \square v_{sC} + (-)\, \gamma_s (\square - 1)^{1/2}, \tag{23}$$

where $\square = |[\boldsymbol{k} - \boldsymbol{k}_0, \boldsymbol{v}_\square]|^2 / |[\boldsymbol{k} - \boldsymbol{k}_0, \boldsymbol{v}_\square]|^2_{th} \equiv \square / \square_{th}$ characterises the extent to which the SBS instability threshold is exceeded. From Eqn (23) we can obtain a simple formula for the SBS linewidth at the base:

$$\Delta\omega = \omega_{max} - \omega_{min} = 2\gamma_s (\square - 1)^{1/2}. \tag{24}$$

It follows from (24) that the SBS linewidth in a plasma with weakly damping ion sound is directly proportional to the damping decrement of this sound and is much smaller than the frequency shift of the scattered radiation $\Delta\omega \ll \omega_{th}$ even when the threshold is exceeded considerably ($\square \gg 1$). In this case, the linewidth increases with pumping intensity proportionally to the square root of the excess over the threshold intensity, the increase being identical and uniform to the red and to the blue.

Fig. 1 shows for comparison the dependence of the boundary of the spectral region of the SBS instability on the pump intensity for a pure hydrogen plasma (curve *2*) for the same degree of nonisothermicity $\square = 4$. We see again that the SBS linewidth depends on the pump intensity. Unlike the case of the carbon plasma considered above, the scattered radiation in hydrogen plasma is characterised by a quite broad line. Indeed, for a fourfold increase in the SBS threshold in hydrogen plasma, the spectral linewidth at the base is about half the threshold frequency: $\Delta\square \approx 0.5\square_{th}$. This is due to the fact that the ion-sound decrement in hydrogen plasma for $\square = 4$ is much larger than in carbon plasma and amounts approximately to 0.12 of the frequency of sound [5]. As the pump intensity I increases, the boundary $\square_{min}$ of the SBS spectrum is displaced strongly to the blue, while the long-wavelength spectral boundary $\square_{max}$ at the base is displaced to a much smaller extent to the red. Such an asymmetry of the instability region boundary is explained by the fact that, in contrast to the carbon plasma considered above, the ion Landau damping dominates in the hydrogen plasma, while the electron Landau damping is negligibly small. The ratio of the electron Landau damping decrement to the frequency of sound is independent of the frequency of sound, which provides, according to (24), a uniform expansion of the SBS instability region in carbon plasma with increasing pump intensity both to the red and to the blue. On the contrary, the ratio of the ion Landau damping decrement to the frequency of sound is a function of the frequency of sound and leads to different dependences of the short- ($\square_{min}$) and long-wavelength ($\square_{max}$) boundaries of the SBS instability on I.

Consider the modification of the SBS instability region boundary in a completely ionised plasma consisting of hydrogen and carbon ions with a degree of ionisation $\square_C = 6$. Curve (*3*) in Fig. 1 shows the boundary of the SBS instability region for the same degree of nonisothermicity $\square = 4$ for the C_5H_{12} plasma in which the temperatures of carbon and hydrogen ions are assumed to be identical. One can see that slightly above the SBS threshold, the SBS linewidth at the base in the C_5H_{12} plasma sharply increases predominantly to the blue. For a pump intensity $\square \approx 1.5\square_{th}$, the linewidth at the base is already more than half the frequency shift for the scattered wave. Upon a further increase in the pump intensity, the SBS linewidth $\Delta\square$ increases only slightly, and the line is broadened in the same way to the blue and to the red. In this region of pump intensities, the SBS linewidth in the C_5H_{12} plasma is approximately half that of the line in the hydrogen plasma. However, this width is more than six times the linewidth in the carbon plasma. Curve (*4*) corresponds to the CH plasma with $\square = 4$. The dependence of the linewidth in the CH plasma on I is similar to that for the C_5H_{12} plasma. Note that an increase in the carbon concentration reduces the SBS linewidth for a certain increase over the instability threshold. Note also that the attenuation of sound associated with the ion Landau damping dominates for the C_5H_{12} and CH plasmas, as well as for pure hydrogen plasma.

Fig. 1 shows that the minimal excitation threshold for the SBS instability is realised in a pure hydrogen plasma; then the H plasma, the CH plasma, and the the C_5H_{12} plasma follow in an increasing order of the threshold intensity $\square_{th}$. We used Fig. 2a and the results of Ref. [5] to explain such a dependence of the threshold pump intensity $\square_{th}$ on the ion composition of plasmas. This figure shows the dependence of $\square_{th}$ on the relative concentration $\square = \square_C / (\square_H + \square_C)$ of carbon ions in the carbon–hydrogen plasma for the same degree of nonisothermicity $\square = 4$ as in Fig. 1. The value $\square = 0$ corresponds to the pure hydrogen plasma, while $\square = 1$ corresponds to the pure carbon plasma. For plasmas with a small carbon impurity ($\square \ll 1$), the SBS threshold is low and increases linearly with the carbon concentration. Upon a further increase in Y, the growth of I_{th} slows down and $I_{th}(Y)$ attains its peak value at $Y_{max} \simeq 0.31$. It should be noted that the C_5H_{12} plasma corresponds to the concentration $Y \simeq 0.294$ which is quite close to Y_{max}. Thus, for $X = 4$, the SBS threshold is close to its maximum value just for the C_5H_{12} plasma. As Y increases further, the intensity I_{th} decreases monotonically, attaining its minimum value for a pure carbon plasma. In order to excite the SBS instability, the pump field must completely compensate for the attenuation of ion-sound waves involved in the scattering process. The threshold intensity I_{th} corresponds to such a complete conpensation of attenuation of the ion-sound wave with a velocity z_{th}. Obviously, the stronger the attenuation of the corresponding ion-sound wave, the stronger the pump field required for its suppression. This can be seen, for example, from expression (22) for a pure carbon plasma. The dependence of the ratio of the damping decrement for ion sound to its frequency on the relative concentration of hydrogen and carbon ions constituting the plasma in a zero pump field is given in Ref. [5]. This dependence is completely similar to that shown in Fig. 2a. Such a behaviour is due to the fact that the larger the ratio of the damping decrement of sound to its frequency, the higher the SBS instability threshold.

Fig. 2b shows the dependence of the threshold velocity of sound z_{th} on the relative concentration Y of carbon ions in a CH plasma for $X = 4$. Recall that it is z_{th} that corresponds to the maximum value of the increment (20) and determines the frequency shift of the scattered wave that is amplified most efficiently with time. One can see from Fig. 2b that the threshold velocity of sound z_{th} for CH plasmas with a relative carbon concentration in the interval $0 \leqslant Y < 0.4$ decreases monotonically with increasing Y. This decrease terminates in the region $0.4 < Y \leqslant 1$ where z_{th} is virtually independent of Y and is determined by the threshold velocity of sound in a pure carbon plasma.

Expression (24) indicates that the linewidth at the base above the SBS instability threshold for a given γ_s is determined not by the absolute value of the pump field intensity, but by its excess over the threshold. Fig. 3 shows the dependence of the ratio $\Delta z/z_{th}$ of the SBS linewidth at the base to the threshold velocity of sound on the parameter p for C-, H-, C_5H_{12}- and CH plasmas. One can see that for the same excess over the SBS instability threshold in these plasmas, the smallest ratio $\Delta z/z_{th}$ is observed for the pure carbon plasma, followed by H-, CH- and C_5H_{12}-plasmas in an increasing order of $\Delta z/z_{th}$. Such a dependence of $\Delta z/z_{th}$ on the ion composition of plasma indicates that for the same p, the relative width $\Delta z/z_{th}$ of the SBS line at the base is determined in analogy with I_{th} by the ratio of the decrement of ion sound to its frequency, and is larger in the plasma in which the ion sound attenuates more strongly [5]. Thus, for a given p, the SBS linewidth is completely determined by the dissipative properties of the given plasma, namely, by the damping decrement of the acoustic mode. It should be emphasised that the curves in Fig. 3 are of qualitative nature only. Since the SBS instability thresholds I_{th} are different in plasmas with different ion compositions (see Fig. 2a), the same values of p in different plasmas correspond to different pump intensities.

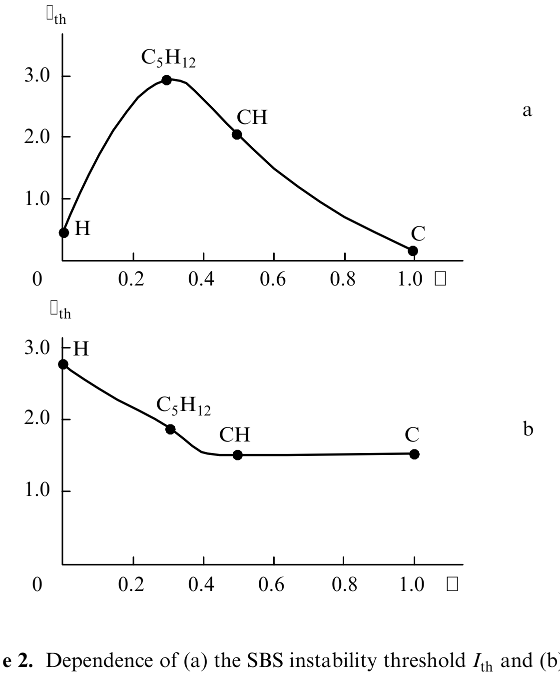

Figure 2. Dependence of (a) the SBS instability threshold I_{th} and (b) the threshold velocity of sound z_{th} in carbon–hydrogen plasma on the relative concentration Y of the carbon ions.

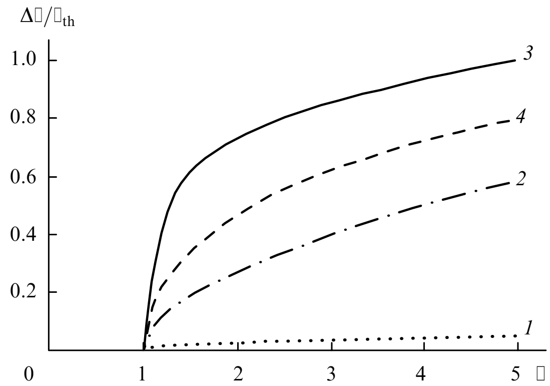

Figure 3. Dependence of the relative width $\Delta z/z_{th}$ of the SBS line on the parameter p for plasmas with the same ion compositions as in Fig. 1: C (*1*), H (*2*), C_5H_{12} (*3*), and CH (*4*).

Consider the influence of the degree of nonisothermicity X of a plasma on the dependence of the boundary of the SBS instability region on the pump intensity. For this purpose, we recall that two different weakly attenuating ion-sound waves (slow and fast) may exist in a two ion-species plasma [5, 6]. In particular, a slow wave with a sound velocity smaller than, or comparable with, the thermal velocity of light ions, but higher than the thermal velocity of heavier ions, attenuates more weakly in a C_5H_{12} plasma for degrees of nonisothermicity X not much higher than unity. On the contrary, for large values of X, the fast wave attenuates more weakly [6]. This information is required for interpreting the dependences presented in Fig. 4.

Fig. 4a shows the dependence of the boundary of the SBS instability region on the pump intensity in a C_5H_{12} plasma with $Z_C = 6$ for $X = 2$. According to Ref. [6], the slow acoustic mode is attenuated weaker for a such value of X. The SBS instability threshold I_{th} is attained just during the suppression of this attenuation. Above this threshold, the scattered radiation with a frequency shift determined by the velocity of sound z_{th} of the slow mode is amplified most

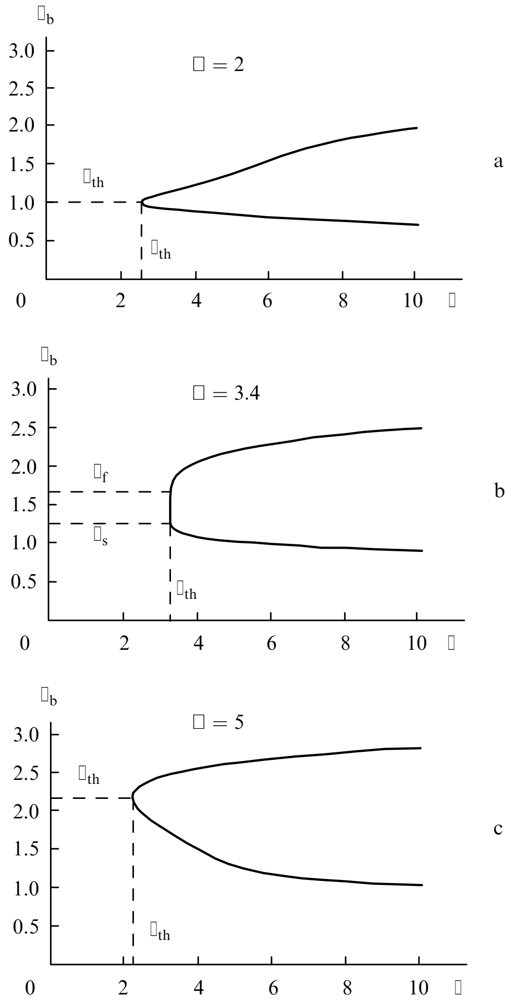

Figure 4. Dependence of the boundary z_b of the SBS instability region on the pump intensity I in the C_5H_{12} plasma with different degrees of nonisothermicity X.

efficiently. For an insignificant excess over the threshold intensity I_{th}, the SBS linewidth at the base increases with I identically and uniformly both to the red and to the blue. For pump intensities almost twice as large as I_{th}, the symmetry of the short- and long-wavelength boundaries of the instability region is broken and the extension of the long-wavelength part dominates. This asymmetry of the instability region boundary is due to the fact that the pump fields are sufficient for suppressing a relatively strong attenuation of the fast ion sound for small X, and lead to amplification of scattered radiation with a frequency shift determined by the frequency of the fast ion-sound waves. For an even larger excess over the threshold I_{th}, the increase in the spectral linewidth at the base with the pump intensity slows down, and the above-mentioned asymmetry of the instability region boundary becomes less pronounced.

Fig. 4b shows another dependence of the boundary of the SBS instability region on the pump intensity for the same C_5H_{12} plasma but with a higher degree of non-isothermicity $X \approx 3.4$. This dependence corresponds to an excitation of the SBS instability for $I \approx I_{th}$ in a wide range of velocities of sound $v_s \leqslant v \leqslant v_f$. Such a feature of the SBS instability boundary is due to the fact that attenuation of the fast acoustic mode for $X \approx 3.4$ becomes the same as the attenuation of the slow mode. Consequently, the same threshold pump intensity I_{th} is required for suppressing the attenuation of both modes (a fast mode with a velocity v_f and a slow mode with a velocity v_s). Even a slight excess over the threshold intensity I_{th} (by just a few percent) broadens the region of SBS instability which becomes comparable with v_f and v_s. In analogy with curve (*3*) in Fig. 1, such a dependence of the SBS linewidth on I indicates that the scattered radiation spectrum contains a wide band comparable with the frequency shift itself. An important spectral feature of SBS considered in this work is that the boundary intensity $I_b(v)$ remains practically unchanged in the interval $v_s \leqslant v \leqslant v_f$, and is close to I_{th}; in other words, in accordance with Eqn (20), the amplitude of scattered radiation in this frequency range increases with time with the same increment. For a large excess over the threshold I_{th}, the width of the SBS instability region increases only weakly with increasing pump field.

Fig. 4c shows the boundary of the SBS instability region in the same C_5H_{12} plasma for an even higher degree of nonisothermicity $X = 5$. In this case, the SBS threshold is associated with the fast ion-sound mode possessing a smaller damping decrement than the slower mode [6]. For an insignificant excess over I_{th}, the linewidth at the base increases with the field more strongly than in Fig. 4a. This allows us to state that, under the conditions when the SBS in a two ion-species plasma is associated with the build up of fast ion- sound waves, the spectrum of scattered radiation contains a broader line than in the case when stimulated scattering is associated with the build up of slow ion-sound waves. Fig. 4c shows that for a still larger excess over I_{th}, the region of SBS instability is broadened considerably mainly to the blue. This means that such pump fields compensate for a relatively strong attenuation of slow sound in the case of large X and lead to amplification of scattered radiation with a frequency shift corresponding to the slow acoustic mode. The SBS linewidth increases comparatively weakly upon a further increase in I.

Concluding the discussion of Fig. 4, note that upon an increase in the pump intensity in a two ion-species plasma, the SBS line is broadened predominantly to the red in the region of nonisothermicity values smaller than unity, when the threshold of SBS instability corresponds to the suppression of weaker attenuation of the slower acoustic mode, and the frequency shift of the scattered radiation is determined by the frequency of sound in this mode. Conversely, for high degrees of nonisothermicity, when the SBS threshold is associated with the build up of the fast ion-sound wave and the scattered radiation whose frequency is shifted by the fast mode frequency is amplified most effectively, an increase in the pump intensity leads to the broadening of the SBS line predominantly to the blue.

Thus, we have studied the dependences of the SBS linewidth and threshold on the pump intensity and composition of carbon–hydrogen plasmas. It is shown that the SBS spectrum can be altered by varying the carbon concentration or the degree of plasma nonisothermicity. We analysed the C_xH_{1-x} plasma which is often used in experiments on controlled laser fusion. However, the properties established in this case, together with those found earlier for the Xe_1H_{99} plasma [9] reveal the extent to which the frequency range, linewidth and SBS threshold can be varied by changing the target material and the extent of plasma heating.

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