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Electromagnetic-wave-pumped free-electron laser in a plasma channel

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An intense short laser pulse or a millimetre wave propagating through a plasma channel may act as a wiggler for the generation of shorter wavelengths. When a relativistic electron beam is launched into the channel from the opposite direction, the laser radiation is Compton/Raman backscattered to produce coherent radiation at shorter wavelengths. The scheme, however, requires a superior beam quality with energy spread less than 1 % in the Raman regime.

1. Introduction

Employment of high-power microwaves and lasers as wigglers in a free-electron laser (FEL), to produce much shorter wavelengths, is an attractive proposition (Khes *et al.* 1990; Esarey and Sprangle 1992; Sprangle *et al.* 1992). The wavelength λ_L of the radiation generated is related to that of the wiggler, λ_0, by

$$\lambda_L \approx \frac{\lambda_0}{4\gamma_0^2},$$

where γ_0 is the relativistic gamma factor of the electron beam. The efficiency of the device could be high when the transverse extents of the wiggler, electron beam and radiation generated (henceforth called FEL radiation) are comparable. In free space, the wiggler and FEL radiation would diverge owing to diffraction, whereas the electron beam would diverge owing to self space charge. However, if one introduces a plasma in the system, the wiggler wave may create a plasma duct (Seo *et al.* 1989; Tripathi and Liu 1990*a*; Wurtele 1993) via the action of the ponderomotive force, which may self-trap the wiggler and FEL radiation (Fig. 1). The plasma would also aid electron-beam guidance via space-charge neutralization (Scharlemann *et al.* 1985; Seo *et al.* 1989; Tripathi and Liu 1990*a*, *b*; Wurtele 1993). One may also employ a preformed plasma channel. Durfee and Milchberg (1993) have demonstrated Nd:glass laser radiation guidance up to 20 Rayleigh lengths in a preformed channel; when a counterpropagating electron beam exists in the plasma channel, one could get high-frequency FEL radiation.

In this paper we study free-electron laser operation using a short-pulse laser

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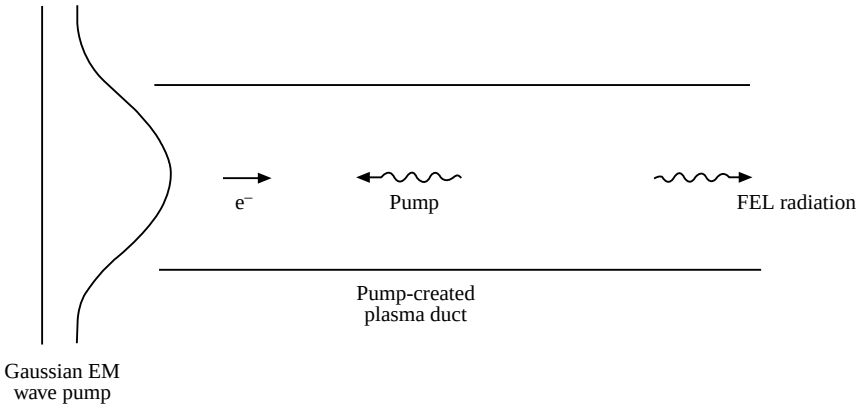


Figure 1. Schematic of free-electron laser interaction in a plasma duct.

in a plasma channel as a wiggler (Liu and Kaw 1976; Mori *et al.* 1988; Sharma and Tripathi 1990; Chen and Sudan 1993) in a plasma medium. In Sec. 2 we discuss self-trapped propagation of a Gaussian electromagnetic wave in a collisionless plasma. In Sec. 3 we investigate nonlinear coupling of this wiggler wave with the beam mode and the FEL mode. First-order perturbation theory is used to obtain the growth rate of the FEL mode. A discussion of our results is given in Sec. 4.

2. Plasma channel

Consider the propagation of an electromagnetic wave

$$\mathbf{E}_0 = (\hat{\mathbf{x}} + i\hat{\mathbf{y}})\, E_0(r) \exp[-i(\omega_0 t - k_0 z)]$$

through a plasma channel whose plasma frequency ω_p has a radial profile

$$\omega_p^2 = \omega_{p0}^2 \left(1 - \frac{r^2}{a^2}\right). \tag{1}$$

The channel could have been created by an intense prepulse (Durfee and Tripathi 1990) or the EM wave under consideration via the ponderomotive force or the relativistic mass effect. In the latter case the size a of the channel, is related to the radius r_0 of the beam as

$$a \approx \begin{cases} r_0\, v_{\rm th}/|\mathbf{v}_0| & \text{(ponderomotive nonlinearity)}, \\ r_0\, c/|\mathbf{v}_0| & \text{(relativistic nonlinearity)}, \end{cases}$$

where $\mathbf{v}_0 = e\mathbf{E}_0/mi\omega_0$ is the oscillatory velocity of electrons; $-e$, m and $v_{\rm th}$ are the electron charge, mass and thermal velocity.

At higher beam powers, saturating effects of nonlinearity limit the radius r_0 of a self-ducted beam to $r_0 \approx c/\omega_p$ and $a \approx r_0$. However, ponderomotive nonlinearity is established in a time of order r_0/c_s, where c_s is the sound velocity. In short-pulse experiments where the pulse duration $\tau < r_0/c_s$ (i.e. $\tau < 1$ ns), channel formation via ponderomotive nonlinearity may not be realized. The onset of relativistic nonlinearity is instantaneous, but requires much higher powers.

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In the case of a preformed plsma channel, a is independent of E_0, and the r variation of E_0 turns out to be (Durfee and Tripathi 1990)

$$E_0^2 = E_{00}^2 e^{-r^2/r_0^2}, \tag{2}$$

where $r_0 = (ac/\omega_p)^{1/2}$.

3. Free-electron laser interaction

A relativistic electron beam of density n_b^0 and velocity $v_b^0 \hat{\mathbf{z}}$, propagates through the plasma channel. The wiggler produces an oscillatory velocity $\mathbf{v}_{0b} = e\mathbf{E}_0/mi\gamma_0\omega_0$ of beam electrons, and couples a beam space-charge mode $\phi = \phi e^{(\omega t - kz)}$ to the FEL wave

$$\mathbf{E}_1 = E_1(r)\,(\hat{\mathbf{x}} - i\hat{\mathbf{y}})\,e^{-i(\omega_1 t - k_1 z)},$$

where $\omega_1 = \omega + \omega_0$ and $\mathbf{k}_1 = \mathbf{k} + \mathbf{k}_0$. The FEL wave produces an oscillatory velocity of beam electrons

$$\mathbf{v}_{1b} = \frac{e\mathbf{E}_1}{mi\gamma_0\omega_1}. \tag{3}$$

The wiggler and the FEL wave exert a ponderomotive force

$$\mathbf{F} = -\tfrac{1}{2}m\mathbf{V}\cdot\boldsymbol{\nabla}\gamma\mathbf{v} - \frac{e}{2c}\mathbf{v}\times\mathbf{B}$$

on beam electrons. The $(\omega, \mathbf{k})$ component of the ponderomotive force turns out to be

$$\mathbf{F}_\omega = e\boldsymbol{\nabla}\phi_p = -\tfrac{1}{2}m\boldsymbol{\nabla}\mathbf{v}_{1b}\cdot\mathbf{v}_{0b}^*.$$

The ponderomotive potential

$$\phi_p = -\frac{eE_1E_0^*}{2\gamma_0^2 m\omega_1\omega_0}. \tag{4}$$

ϕ and ϕ_p produce an oscillatory electron velocity and density perturbation at $(\omega, \mathbf{k})$:

$$\mathbf{v}_b = -\frac{ek(\phi+\phi_p)}{m\gamma_0^3(\omega - k_z v_b)}\hat{\mathbf{z}}, \tag{5}$$

$$n_b = -\frac{(k_z^2 - \nabla_\perp^2)\,n_b^0 e}{m\gamma_0^3(\omega - k_z v_b)^2}(\phi+\phi_p). \tag{6}$$

The plasma-density perturbation can be written as

$$n_p = \frac{n_0 e}{m\omega^2}\left(\nabla^2\phi - \nabla^4\phi\frac{v_{\rm th}^2}{\omega^2}\right), \tag{7}$$

Using (6) and (7) in Poisson's equation, we obtain

$$\frac{\partial^2\phi}{\partial r^2} + \frac{1}{r}\frac{\partial\phi}{\partial r} + \left[\frac{\omega^2 - \omega_p^2 - k_z^2 v_{\rm th}^2}{v_{\rm th}^2} - \frac{\omega^2\omega_{pb}^2}{v_{\rm th}^2\gamma_0^3(\omega - k_z v_b)^2} - \frac{\omega_p^2}{v_{\rm th}^2}\frac{r^2}{a^2}\right]\phi = -\frac{\omega_{pb}^2\omega^2\phi_p}{v_{\rm th}^2(\omega - k_z v_b)^2\gamma_0^3}. \tag{8}$$

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The nonlinear current density at $(\omega_1, \mathbf{k}_1)$ is

$$\mathbf{J}_1^{nl} = -\tfrac{1}{2} n_b e \mathbf{v}_{0b} = \frac{k^2 n_b^0 e^3 (\phi + \phi_p)\mathbf{E}_0}{2m^2 i \gamma_0^4 \omega_0 (\omega - k_z v_b)^2}. \tag{9}$$

Using $\mathbf{J}_1^{nl}$ in the wave equation

$$\nabla^2 \mathbf{E}_1 + \frac{\omega_1^2 - \omega_p^2(r)}{c^2}\mathbf{E}_1 = -\frac{4\pi i \omega_1}{c^2}\mathbf{J}_1^{nl},$$

and employing (4), we get

$$\frac{\partial^2 E_1}{\partial r^2} + \frac{1}{r}\frac{\partial E_1}{\partial r} + \left[\frac{\omega_1^2 - \omega_p^2 - k_1^2 c^2}{c^2} + \frac{\omega_{pb}^2 k^2 e^2 A_0^2}{4m^2 c^2 \omega_0^2 \gamma_0^2 (\omega - k_z v_b)^2} - \frac{\omega_p^2}{c^2}\frac{r^2}{a^2}\right] E_1 = \frac{i\omega_{pb}^2 \omega_1 k^2 v_0}{2(\omega - k_z v_b)^2 c^2 \gamma_0^4}\phi. \tag{10}$$

We investigate FEL operation in two regimes – the Compton and Rayleigh regimes.

Case (i): FEL operation in the Compton regime

In the Compton regime the collective effects are weak, hence $\phi \ll \phi_p$ and (10) takes the form

$$\frac{\partial^2 E_1}{\partial r^2} + \frac{1}{r}\frac{\partial E_1}{\partial r} + (\lambda_1 - \lambda_2 r^2) E_1 = 0, \tag{11}$$

where

$$\lambda_1 = \frac{\omega_1^2 - \omega_p^2 - k_1^2 c^2}{c^2} + \frac{\omega_{pb}^2 k^2 e^2 A_0^2}{4m^2 c^2 \omega_0^2 \gamma_0^2 (\omega - k_z v_b)^2}, \tag{12}$$

$$\lambda_2 = \omega_p^2 / a^2 c^2. \tag{13}$$

Introducing $\xi = \lambda_2^{1/4} r$ (into (11), we get

$$\frac{\partial^2 E_1}{\partial \xi^2} + \frac{1}{\xi}\frac{\partial E_1}{\partial \xi} + \left(\frac{\lambda_1}{\lambda_2^{1/2}} - \xi^2\right) E_1 = 0. \tag{14}$$

Equation (14) has a well-behaved solution when

$$\frac{\lambda_1}{\lambda_2^{1/2}} = 2n + 1, \qquad n = 0, 1, 2, 3, \ldots \quad .$$

For $n = 0$, we have

$$\left[\omega_1^2 - \omega_p^2 - k_1^2 c^2 - \left(\frac{\omega_p^2 c^2}{a^2}\right)^{1/2}\right](\omega - k_z v_b)^2 = \frac{\omega_{pb}^2 k^2 e^2 A_0^2}{4m^2 c^2 \omega_0^2 \gamma_0^2}. \tag{15}$$

We solve this equation around the simultaneous zeros of the left-hand side. We write

$$\omega = k_z v_b + \delta, \qquad \omega_1 = \left[\omega_p^2 + k_1^2 c^2 + \left(\frac{\omega_p^2 c^2}{a^2}\right)^{1/2}\right]^{1/2} + \delta,$$

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where δ is the modification in ω due to the finite right-hand side. Substituting into (15), we get

$$\delta = \left(\frac{\omega_{pb}^2 k^2 e^2 A_0^2}{8\omega_1 \omega_0^2 m^2 \gamma_0^2}\right)^{1/3} e^{2in\pi/3}. \tag{16}$$

Thus the growth rate in the Compton regime is

$$\Gamma_c = \mathrm{Im}(\delta) = \frac{\sqrt{3}}{2}\left(\frac{\omega_{pb}^2 a_w^2 k^2 c^2}{8\omega_1 \gamma_0^2}\right)^{1/3}, \tag{17}$$

where $a_w = eE_0/m\omega_0 c$. For a typical CO_2 laser-pumped FEL with parameters $\gamma_0 = 2$, $a = 1$ mm, $\omega_p \approx \frac{1}{4}\omega_0$, beam-current density $J_b \approx 300$ A cm^{-2} ($I_b \approx 10$ A) and pump-wave power density $\approx 5\times 10^{11}$ W cm^{-2} ($a_w \approx 0.06$), the growth rate turns out to be $\Gamma_c \approx 3\times 10^{10}$ s^{-1}. However, this requires high beam quality, with energy spread $\Delta\gamma_0 \leqslant \Gamma_c/\omega_0 \approx 1.5\times 10^{-4}$. In a Compton-regime magnetostatic wiggler-based optical FEL (Billardon *et al.* 1985), for parameters $\gamma_0 \approx 2$, $\lambda_0 = 6500$ Å, $P_0 = 190$ MW ($a_w \approx 1$), $J_b \approx 33$ A cm^{-2} ($I_b \approx 40$ mA) and $r_0 = 0.2$ mm, the growth rate (following Freund and Antonsen 1987) is of the order of 3×10^{10} s^{-1} and is comparable to the case of a CO_2 laser-pumped FEL.

In the case of a gyrotron-pumped FEL with $\lambda_0 = 1$ cm, $r_0 \approx 3$ cm, $P_0 = 80$ MW ($a_w \approx 0.06$), $\gamma_0 = 2$ and $J_b \approx 3$ A cm^{-2} ($I_b \approx 100$ A), the growth rate turns out to be $\Gamma_c \approx 10^9$ s^{-1}, requiring a beam-energy spread $\Delta\gamma_0 \leqslant 5\times 10^{-3}$.

When the beam-energy spread is large, such that $k\Delta v_b \geqslant \Gamma_c$, the regime of FEL operation shifts to the resistive Compton regime. In this case (11) still holds for E_1. However, in the expression for λ_1, $\omega_{pb}^2/(\omega - k_z v_b)^2 \gamma_0^3$ is replaced by $-\chi_b$, the beam susceptibility

$$\chi_b = \frac{\omega_{pb}^2/\gamma_0}{k'^2 v_{\mathrm{th}}^2}\left[1 + \frac{\omega'}{k' v_{\mathrm{th}}} z\left(\frac{\omega'}{k' v_{\mathrm{th}}}\right)\right],$$

where

$$k' = \gamma_0\left(k - \frac{\omega v_b}{c^2}\right), \qquad \omega' = \gamma_0(\omega - k v_b)$$

and v_{th} is the thermal speed of the beam in a frame moving with average beam velocity v_b. The growth rate turns out to be

$$\Gamma_c \approx \mathrm{Im}\left(\chi_b \frac{k^2 c^2 a_w^2}{8\omega_1}\right). \tag{18}$$

Here we have assumed the beam distribution function in the moving frame to be Maxwellian.

Case (ii): FEL operation in the Raman regime

At large beam density, $\phi > \phi_p$, (8) and (10) can be rewritten as

$$\frac{\partial^2\phi}{\partial r^2} + \frac{1}{r}\frac{\partial\phi}{\partial r} + (k_\perp^2 - g r^2)\,\phi = B E_1, \tag{19}$$

$$\frac{\partial^2\phi}{\partial r^2} + \frac{1}{r}\frac{\partial\phi}{\partial r} + (k_{1\perp}^2 - g_1 r^2)\,E_1 = B_1\,\phi, \tag{20}$$

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where

$$k_\perp^2 = \frac{\omega^2 - \omega_p^2 - k_z^2 v_{\rm th}^2}{v_{\rm th}^2} - \frac{\omega_{pb}^2 \omega^2}{v_{\rm th}^2 \gamma_0^3 (\omega - k_z v_b)^2},$$

$$k_{1\perp}^2 = \frac{\omega_1^2 - \omega_p^2 - k_1^2 c^2}{c^2} + \frac{\omega_{pb}^2 k^2 e^2 A_0^2}{4m^2 c^2 \omega_0^2 \gamma_0^2 (\omega - k_z v_b)^2},$$

$$g = \frac{\omega_p^2}{v_{\rm th}^2 a^2}, \qquad g_1 = \frac{\omega_p^2}{c^2 a^2},$$

$$B = -\frac{i\omega_{pb}\omega^2}{2v_{\rm th}^2(\omega - k_z v_b)^2}\frac{v_0}{\gamma_0^3}, \qquad B_1 = \frac{i\omega_{pb}^2 \omega_1 k^2 v_0}{1(\omega - k_z v_b)^2 c^2 \gamma_0^4}.$$

In the absence of the coupling terms, $v_0 = 0$, (19) and (20) have solutions

$$\phi = \phi_l = \Gamma_l L_l\left(\frac{r^2}{b^2}\right)\exp\left(-\frac{r^2}{2b^2}\right), \tag{21}$$

$$E_1 = \psi_m = \Gamma_m L_m\left(\frac{r^2}{b_1^2}\right)\exp\left(-\frac{r^2}{2b_1^2}\right), \tag{22}$$

respectively when the following eigenvalue conditions are satisfied:

$$k_\perp^2 \equiv 2(l+1)\frac{\omega_p}{v_{\rm th} a}, \tag{23}$$

$$k_{1\perp}^2 \equiv 2(m+1)\frac{\omega_p}{ac}, \tag{24}$$

where

$$b^2 = \frac{a v_{\rm th}}{\omega_p}, \qquad b_1^2 = \frac{ca}{\omega_p},$$

$$L_l(\psi) = e^{\psi}\frac{d^l}{d\psi^l}(\psi^l e^{-\psi}), \qquad l = 0, 1, 2, 3, \ldots, \quad m = 0, 1, 2, 3, \ldots$$

are associated Laguerre polynomials, and Γ_l and Γ_m are normalization constants given by

$$\Gamma_l^2 = \frac{2}{b^2(l+1)}, \qquad \Gamma_m^2 = \frac{2}{b_1^2(m+1)}.$$

It may be noted that $b_1 \approx a$, whereas $b \ll a$, i.e. the radial extent of the backscattered electromagnetic wave is nearly the same as that of the pump, while the Langmuir wave is localized to a much shorter region. Since the pump field (and hence v_0) scales as $\exp(-r^2/2a^2)$, the most-unstable decay waves would correspond to $m = 0$, $l = 0$.

In the presence of nonlinear coupling terms, one could express the solutions of (19) and (20) in terms of associated Laguerre polynomials as

$$\phi = A\phi_0, \tag{25}$$

$$E_1 = A_1 \psi_0 \tag{26}$$

respectively. Using (26) in (19), and multiplying the resulting equation by ϕ_0 and integrating over $r\,dr$, one gets

$$a_i(k_\perp^2 - k_{\perp i}^2) = B a_j \int \psi_0 \phi_0 r\,dr. \tag{27}$$

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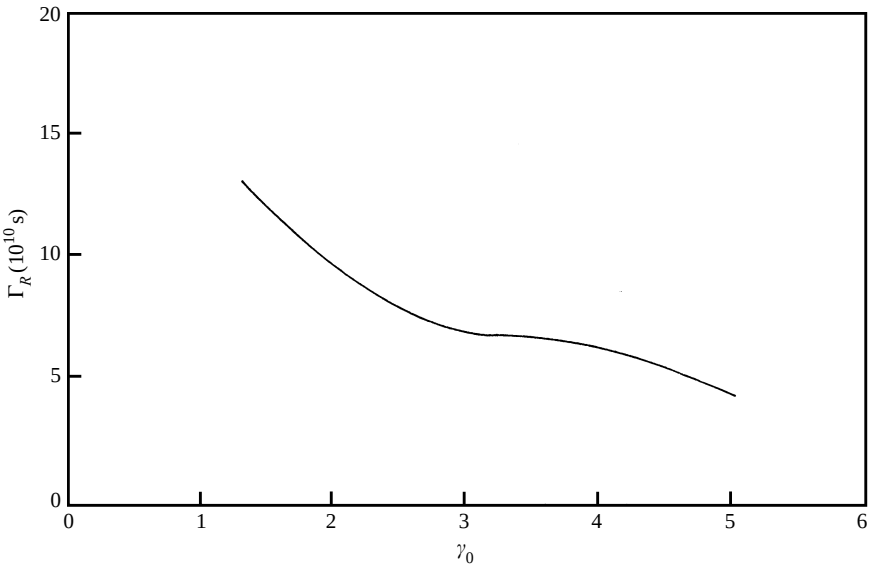


Figure 2. Variation of growth rate in the Raman regime with the energy of the electron beam for a 10 μm laser with pump power density $\approx 3\times10^{11}$ W cm^{-2} ($a_w \approx 0.06$) and beam radius ($r_0 \approx 300\,\mu$m).

Similarly, (25) and (20) give

$$a_j(k_{1\perp}^2 - k_{1\perp j}^2) = B_1 a_i \int \phi_0 \psi_0 r\, dr. \tag{28}$$

Equations (27) and (28) yield the nonlinear dispersion relation

$$(k_\perp^2 - k_{\perp i}^2)(k_{1\perp}^2 - k_{1\perp j}^2) = BB_1 \left| \int \phi_0 \psi_0 r\, dr \right|^2. \tag{29}$$

Let $\omega = \omega_r + i\Gamma_R$ and $\omega_1 = \omega_{1r} + i\Gamma_R$, where ω_r and ω_{1r} are the frequencies at which the two factors on the left-hand side of (29) vanish exactly. We expand

$$D(\omega) \equiv k_\perp^2 - k_{\perp i}^2 = D(\omega_r) + \frac{\partial D}{\partial \omega_r} i\Gamma_R$$

$$\approx -\frac{2\omega^2\gamma_0^{3/2}}{\omega_{pb}} \frac{i\Gamma_R}{v_{\text{th}}^2}, \tag{30}$$

$$k_{1\perp}^2 - k_{1\perp j}^2 = \frac{2\omega_1}{c^2} i\Gamma_R, \tag{31}$$

where

$$\omega_{1r} = \frac{1}{c}\left(\omega_p^2 + k_1^2 c^2 - \frac{\omega_{pb}^2 k^2 e^2 A_0^2}{4m^2\omega_0^2\gamma_0^2(\omega - k_z v_b)^2} + \frac{2\omega_p c}{a}\right)^{1/2}, \tag{32}$$

On substituting (30) and (31) into (29), we get the growth rate in the Raman regime as

$$\Gamma_R \approx \left(\frac{\omega_{pb} I^2 a_w^2 \omega_1}{16\gamma_0^{3/2}}\right)^{1/2}, \tag{33}$$

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where

$$I = \left| \int \phi_i \psi_j \, r \, dr \right|.$$

For a typical CO_2 laser-pumped FEL with parameters $\gamma_0 = 2$, $\omega_p = \frac{1}{4}\omega_0$, beam-current density $J_b \approx 0.4$ kA cm^{-2} ($I_b \approx 100$ A), pump-wave power density $\approx 3 \times 10^{11}$ W ($a_w \approx 0.06$) and beam radius $a_0 \approx 300$ μm, the growth rate Γ_R turns out to be of order 10^{11} s^{-1} (Fig. 2). In the case of a gyrotron-pumped FEL with $\lambda_0 = 1$ cm, $r_0 \approx 3$ cm, $P_0 = 80$ MW ($a_w \approx 0.06$), $\gamma_0 = 2$ and $J_b \approx 300$ A cm^{-2} ($I_b \approx 10$ kA), the growth rate Γ_R turns out to be $\approx 6 \times 10^9$ s^{-1}.

4. Discussion

An electromagnetic-wave pumped FEL, with high-density plasma channel in the interaction region, has some attractive features. The plasma channel created by a prepulse or the self-action of the pump electromagnetic wave provides divergenceless propagation of all the interacting waves. When the wave functions of the daughter waves overlap with the pump wave, the growth rate of the parametric upconversion process is large. The plasma helps propagation of a high beam current without divergence. The growth rate of a CO_2 laser-pumped FEL with typical parameters, pump-wave power $\approx 3 \times 10^{11}$ W ($a_w \approx 0.06$), beam radius $r_0 \approx 300$ μm and $\gamma_0 = 2$, the growth rate turns out to be $\approx 3 \times 10^{10}$ s^{-1} in the Compton regime ($I_b \approx 10$ A), whereas it is of order 10^{11} s^{-1} in the Raman regime ($I_b \approx 10$ A). For a gyrotron-pumped FEL with $\lambda_0 = 1$ cm, $r_0 \approx 3$ cm, $P_0 = 80$ MW ($a_w = 0.06$) and $\gamma_0 = 2$, the growth rate is of order 10^9 s^{-1} in the Compton regime ($I_b = 1$ kA) and $\approx 6 \times 10^9$ s^{-1} in the Raman regime ($I_b = 10$ kA).

A major limitation of the proposed scheme arises owing to the requirement of high beam quality. The requisite energy spread $\Delta E/E \leqslant \Gamma_0/\omega_0$. When the beam energy is large, the FEL operation shifts to the resistive Compton regime, where the growth rate is reduced by an order of magnitude. This regime, however, calls for a more detailed kinetic analysis.

The electromagnetic-wave-pumped free-electron laser has the potential to produce coherent radiation in the ultraviolet without requiring too much high beam energies. Sprangle *et al.* (1992) have proposed a laser-pumped X-ray FEL, based on Thomson backscattering of laser radiation from a counterstreaming electron beam.

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