Gain in a Compton laser

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The gain in a Compton laser for arbitrary directions of propagation of the electrons, pumping wave, and amplified wave is found. It is shown that the for electron beams with a sufficiently large energy spread the gain increases with increasing (small) deviation of the direction of propagation of the amplified wave from the direction of motion of the electrons. The optimum value of the angle between these directions (i.e., the value at which the gain has its maximum value) is found. Numerical estimates indicating the feasibility of constructing a Compton laser with the optimum experimental geometry are presented.

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1. INTRODUCTION

One of the most important directions of present-day laser physics is the construction and development of free-electron lasers (FEL).[1-4]

In the narrow sense of the word, a FEL is a laser in which the generation of the radiation is due to the interaction of a relativistic electron beam with the spatially periodic magnetic field of an undulator. There are many other proposals for the use of the interaction of a relativistic electron beam with other objects to amplify and generate radiation.[5]

The Compton laser, first proposed by Pantell *et al.*,[6] is, by its nature, closest to the undulator-based FEL. In the Compton laser the electron beam interacts with the field of two external electromagnetic waves: the pumping wave and the amplified wave. Experimentally, the Compton laser has so far not been realized. Theoretically, the amplification in a Compton laser is considered in Refs. 7 and 8 in the weak-signal approximation and in Ref. 9 under saturation conditions.

One of the reasons why a FEL, like a Compton laser, is attractive is the potential possibility of regulating the generation frequency by varying the electron energy. With this, in particular, are also connected hopes of advancement into the high-frequency region. In this respect the experiment of Deacon *et al.*,[2] in which the generation was realized at a wavelength of 3.2 μm, is a record-breaking and, in its own way, at present a unique experiment.

Another trend in the development of the FEL is the raising of the output power through the use of high-current electron sources. The highest power (~1 MW) has been achieved by McDermott *et al.*[3]

It is clear that the raising of the power and the raising of the generation frequency are mutually contradictory requirements, since we need high-current accelerators in one case and high-energy electron sources in the other. The physical difference between relatively high-power FEL with a not very high generation frequency (see Refs. 1, 3, and 4) and FEL using high-energy electrons[2] lies in the fact that in the first case a significant role may be played in the amplification process by the collective phenomena occurring in the electron beam plasma, whereas in the second case the amplification has a single-particle character. The quantitative parameter separating these two regions is equal[10] to $\omega_b t/\gamma^{3/2}$, where ω_b is the plasma frequency of the beam, t is the interaction time, $\gamma = \varepsilon/mc^2$ is the relativistic factor, and ε is the electron energy. In the present paper we consider only the case $\omega_b t < \gamma^{3/2}$, when the single-particle description is valid.

To advance into the high-frequency region, it is very important that we raise the amplification factor. Various modified FEL schemes are being actively investigated at present with this aim in view.[11-16] It should, however, be noted that even within the framework of the already known schemes the possibility of raising the amplification factor has not been fully investigated in some cases. In this plan, we consider in the present paper the amplification in a Compton laser in the case of arbitrary directions of propagation of the pumping wave, the amplified wave, and the electron beam (in contrast to the investigations carried out in Refs. 6, 7, and 9, where only the one-dimensional scheme is discussed).

The amplification factor is computed on the basis of the quantum description of the motion of an electron in the field of external classical electromagnetic radiation. All the final expressions obtained for the amplification factor have, however, a classical character, since they do not contain the Planck constant. It is therefore clear that the results obtained in the present paper can also be derived on the basis of the classical equations of motion of an electron in a classical electromagnetic field.

Qualitatively, it is clear that we can expect the amplification factor in a Compton laser to increase as the scheme deviates from the one-dimensional case. This is due to the fact that in the case in which the electron- and photon-momentum vectors are collinear, the linear amplification factor is a decreasing function of the frequency of the amplified wave.[9] Other things being equal, as the direction of propagation of the wave being amplified deviates from the direction of motion of the electrons, the frequency of the wave decreases, remaining nonetheless high enough for such a process to be of interest. If the structure of the amplification factor for the case of noncollinear propagation is similar to the structure of the amplification factor in the one-dimensional scheme, then the decrease of the frequency of the

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amplified wave can lead to the increase of the amplification factor.

In the present paper we find an expression for the amplification factor in the noncollinear scheme of propagation of the waves and the electron beam, and show that it is indeed expedient to use the non-one-dimensional scheme in the case of a relatively large energy spread of the beam electrons. The optimal geometry turns out to be the one in which the pumping wave propagates counter to the electron beam, while the amplified wave propagates at some (fairly small) angle to the direction of the electron velocity. The optimal values of the angle (i.e., the values at which the amplification is maximal) are found, and the corresponding amplification factor is estimated for the generation conditions realizable at the present time.

2. THE BASIC EQUATIONS

Let us consider the motion of an electron in the field of two waves characterized by the vector potentials ($\hbar = c = 1$)

$$\mathbf{A}_1=\frac{E_1}{2\omega_1}(\mathbf{e}_1e^{ik_1x}+\text{c.c.}),\quad \mathbf{A}_2=\frac{E_2}{2\omega_2}(\mathbf{e}_2e^{ik_2x}+\text{c.c.}),\tag{1}$$

where $\mathbf{e}_{1,2}$ are unit polarization vectors (for simplicity we consider the case of linear polarization of the radiation); $E_{1,2}$ are the amplitudes of the field intensities of the waves; $k_{1,2}=(\mathbf{k}_{1,2},\omega_{1,2})$ are the 4-momenta of the waves; and the indices 1, 2 pertain respectively to the pumping and amplified waves.

Let us assume that the frequencies ω_1 and ω_2 roughly satisfy the standard relation fulfilled in Compton scattering,[17] a relation which, for $\omega_{1,2}\ll\varepsilon$, can be written in the form

$$\omega_2=\frac{\omega_1(1-v\cos\theta_1)}{1-v\cos\theta_2},\tag{2}$$

where $\theta_{1,2}$ are the angles between the vectors $\mathbf{k}_{1,2}$ and $\mathbf{p}$; $\mathbf{v}=\mathbf{p}/\varepsilon$, $\mathbf{p}$ being the momentum, and ε the energy, of the electron.

Neglecting the small spin corrections, we use as the basic equation the Klein–Gordon equation for the $\mathbf{A}_{1,2}$ fields:

$$\left\{\frac{\partial^2}{\partial t^2}-\left(\frac{\partial}{\partial\mathbf{r}}-ie\mathbf{A}_1-ie\mathbf{A}_2\right)^2+m^2\right\}\Psi=0.\tag{3}$$

Let us consider the simplest formulation of the problem, within the framework of which we shall formally neglect the interaction region's spatial boundedness, which practically always obtains. Instead, we shall assume that the interaction occurs during a finite interval of time. This time may in fact be determined by the time of transit of an electron through the interaction region.

In such a formulation of the problem, the wave function Ψ can be expanded in terms of plane waves, the coefficients of the expansion being slowly varying functions of the time, a fact which allows us to neglect their second derivatives when we substitute the Ψ expansion into Eq. (3).[18]

The probabilities for the stimulated Compton-scattering processes can be found from (3) in second-order

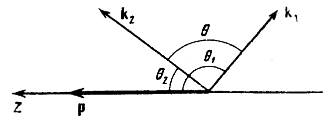

FIG. 1. Planar geometry of the stimulated Compton scattering process under consideration and the definition of the angles $\Theta_{1,2}$ and Θ.

perturbation theory, and can be represented in the form

$$w_\pm=|M_{A^2}-M_{\mathbf{pA}}|^2\Phi_\pm^2,\tag{4}$$

where $w_+(w_-)$ is the probability of emission (absorption) of a photon of frequency ω_2 and absorption (emission) of a photon of frequency ω_1; M_{A^2} and $M_{\mathbf{pA}}$ are the matrix elements determined by the quadratic and linear—in the field—terms in Eq. (3).

Let us give the explicit expressions for these matrix elements:

$$M_{A^2}=e^2E_1E_2(\mathbf{e}_1\mathbf{e}_2)/4\varepsilon\omega_1\omega_2,$$

$$M_{\mathbf{pA}}^{\pm}=\frac{e^2E_1E_2}{4\varepsilon\omega_1\omega_2}\left[\frac{(\mathbf{e}_2(\mathbf{p}\pm\mathbf{k}_1))(\mathbf{e}_1\mathbf{p})}{(\varepsilon_{\mathbf{p}\pm\mathbf{k}_1}-\varepsilon_{\mathbf{p}}\mp\omega_1)\varepsilon_{\mathbf{p}\pm\mathbf{k}_1}}+\frac{(\mathbf{e}_1(\mathbf{p}\mp\mathbf{k}_2))(\mathbf{e}_2\mathbf{p})}{(\varepsilon_{\mathbf{p}\mp\mathbf{k}_2}-\varepsilon_{\mathbf{p}}\pm\omega_2)\varepsilon_{\mathbf{p}\mp\mathbf{k}_2}}\right],\tag{5}$$

where $\mathbf{p}$ is the initial electron momentum and $\varepsilon_p=(p^2+m^2)^{1/2}$.

The factor $\Phi_\pm$ in (4) is equal to

$$\Phi_\pm=\int dt\exp[i(\varepsilon_{\mathbf{p}\pm(\mathbf{k}_1-\mathbf{k}_2)}-\varepsilon_{\mathbf{p}}\mp(\omega_1-\omega_2))t].\tag{6}$$

Further, let us, for simplicity, consider a planar geometry, assuming that the vectors $\mathbf{e}_{1,2}$, $\mathbf{k}_{1,2}$, and $\mathbf{p}$ lie in the same plane (Fig. 1). The expression, (5), for $M_{\mathbf{pA}}$ gets transformed, when allowance is made for the smallness of $\omega_{1,2}$ and $|\mathbf{k}_{1,2}|$ in comparison with $\varepsilon\approx|\mathbf{p}|$, into

$$M_{\mathbf{pA}}=\frac{e^2E_1E_2}{4\varepsilon\omega_1\omega_2}\left\{\sin\theta_1\sin\theta_2\frac{\cos\theta+v(\cos\theta_1+\cos\theta_2)-3v^2\cos\theta_1\cos\theta_2}{(1-v\cos\theta_1)(1-v\cos\theta_2)}+\left(\frac{\sin\theta_1}{1-v\cos\theta_1}-\frac{\sin\theta_2}{1-v\cos\theta_2}\right)\sin\theta\right\},\tag{7}$$

where θ is the angle between the vectors $\mathbf{k}_1$ and $\mathbf{k}_2$.

From this it is easy to find in the particular case in which $\theta_2\ll1$ and $\theta_1\sim1$ the relatively simple expression for the total matrix element:

$$M=M_{A^2}-M_{\mathbf{pA}}=-\frac{e^2E_1E_2}{4\varepsilon\omega_1\omega_2}\left\{1+2\frac{(\gamma\theta_2)^2}{1+(\gamma\theta_2)^2}\frac{\cos\theta_1-\cos2\theta_1}{1-\cos\theta_1}\right\}.\tag{8}$$

For $\theta_1=\pi$, i.e., in the case in which the direction of motion of the electron beam is opposite to the direction of propagation of the pumping wave (ω_1),

$$M=-\frac{e^2E_1E_2}{4\varepsilon\omega_1\omega_2}\frac{1-(\gamma\theta_2)^2}{1+(\gamma\theta_2)^2}.\tag{9}$$

3. THE GAIN

The energy $\Delta\mathscr{E}$ emitted by an electron at the frequency ω_2 and the amplification factor G for this frequency can be expressed in terms of $w_\pm$ with the aid of the relations

$$\Delta\mathscr{E}=\omega_2(w_+-w_-),\quad G=8\pi E_2^{-2}N_e\Delta\mathscr{E},\tag{10}$$

where N_e is the electron density in the beam.

The factor $\Phi_\pm^2$ occurring in the formula (4), and determined by the integral (6), has a singularity, which, as usual, is removed when we take into account the real properties of the interacting objects—such properties as the energy spread of the electrons, the finiteness of the interaction time, etc.

Let us first consider the case in which the dominant physical mechanism leading to the removal of the singularity in $\Phi_\pm^2$ is the energy spread of the electrons in the incident beam. For the individual electrons the expression (4) can be written in the form

$$w_\pm=2\pi t|M_{A^2}-M_{pA}|^2\delta(\varepsilon_{p\pm(k_1-k_2)}-\varepsilon_p\mp(\omega_1-\omega_2)). \quad (11)$$

To find the mean $\Delta\mathscr{E}$ and G values pertaining to the beam as a whole, the formula (11) should be averaged over the energy distribution function, $f(\varepsilon)$, for the incident-beam electrons, which is assumed here to be normalized to unity by the condition

$$\int f(\varepsilon)d\varepsilon=1, \quad f_{max}=f(0)\sim 1/\delta\varepsilon$$

(we denote by $\delta\varepsilon$ the halfwidth of the function $f(\varepsilon)$).

The averaging over ε can easily be carried out, using the condition $\omega_{1,2}\ll\varepsilon$. In this approximation the δ function in the formula (11) can be represented in the form

$$\delta(\varepsilon_{p\pm(k_1-k_2)}-\varepsilon_p\mp(\omega_1-\omega_2))\approx\frac{\gamma^2\varepsilon\delta(\varepsilon-\varepsilon_{e,a})}{|\omega_2\cos\theta_2-\omega_1\cos\theta_1|}, \quad (12)$$

where $\varepsilon_{e,a}=\varepsilon\pm\Delta\varepsilon$ are the energies of the electrons accomplishing the emission or absorption of a photon of given frequency ω_2 in the case of a fixed frequency ω_1 of the absorbed or emitted photon of the pumping wave,

$$\varepsilon=\frac{m|\omega_1\cos\theta_1-\omega_2\cos\theta_2|}{[-\omega_1^2\sin^2\theta_1-\omega_2^2\sin\theta_2+2\omega_1\omega_2(1-\cos\theta_1\cos\theta_2)]^{1/2}}$$
$$\Delta\varepsilon=\frac{(\omega_2-\omega_1)\omega_2\omega_1(1-\cos\theta)}{-\omega_1^2\sin^2\theta_1-\omega_2^2\sin^2\theta_2+2\omega_1\omega_2(1-\cos\theta_1\cos\theta_2)}\ll\varepsilon. \quad (13)$$

In the case $\theta_1=\pi$, $\theta_2\ll 1$, which is the most interesting in practice, these formulas assume the form

$$\varepsilon=m\left(\frac{\omega_2}{4\omega_1-\omega_2\theta_2^2}\right)^{1/2}, \quad \Delta\varepsilon=2\gamma^2\omega_1. \quad (14)$$

The averaging of the expression (11) and the computation of the mean radiated energy yield in the general case the expression

$$\Delta\mathscr{E}=4\pi t\omega_2|M_{A^2}-M_{pA}|^2\frac{f'(\varepsilon)\gamma^2\varepsilon\Delta\varepsilon}{|\omega_2\cos\theta_2-\omega_1\cos\theta_1|}. \quad (15)$$

In the particular case of $\theta_1=\pi$, $\theta_2\ll 1$, we find from this expression with the aid of the expressions (2), (9), (10), and (14) the gain

$$G=\frac{4\pi^2e^4N_etE_1^2\varepsilon^3f'(\varepsilon)}{m^4\omega_1\omega_2^2}\left(1-\frac{\omega_2}{2\gamma^2\omega_1}\right)^2, \quad (16)$$

where for the given geometry the amplified-wave frequency is

$$\omega_2=4\gamma^2\omega_1/(1+\gamma^2\theta_2^2). \quad (17)$$

As follows from the expression (16), as θ_2 increases, the gain G at first decreases, vanishes at $\gamma\theta_2=1$, and then increases in proportion to θ_2^4 (the frequency ω_2 decreases in the process in proportion to θ_2^{-2}).

For $\theta_2\gg 1/\gamma$, the formula (16) assumes the form

$$G\left(\theta_2\gg\frac{1}{\gamma}\right)=\frac{\pi^2e^4N_etE_1^2\varepsilon^3f'(\varepsilon)\theta_2^4}{4m^4\omega_1^3}=(\gamma\theta_2)^4G(\theta_2=0). \quad (18)$$

It is interesting to note that such increase of the radiated energy with increasing θ_2 does not occur in the spectral density of the spontaneous Compton scattering involving the emission of a photon of frequency ω_2. This is due to the fact that the phase-volume element that arises when we go over from the given external field E_2 to the density of the zero-point vacuum vibrations contains a factor $\propto\omega_2^3$, which rapidly decreases with decreasing ω_2 (increasing θ_2). It is easy to verify that the spectral and angular densities of the spontaneous emission at the frequency ω_2 has in the case under consideration the form

$$\frac{d\mathscr{E}_{sp}}{d\omega_2 d\Omega_{k_2}}=\frac{1}{2\pi}\frac{e^2tE_1^2\varepsilon^3f(\varepsilon)}{m^4\omega_1(1+\gamma^2\theta_2^2)}\left(\frac{1-\gamma^2\theta_2^2}{1+\gamma^2\theta_2^2}\right)^2. \quad (19)$$

With increasing θ_2 the spontaneous radiation density for $\gamma\theta_2>1$ decreases like θ_2^{-2}. Moreover, there occurs outside the relativistic cone a narrowing of the spectral lines of the spontaneous and stimulated Compton scattering, since for $\gamma\theta_2>1$ the dependence, (17), of ω_2 on energy weakens. It is easy to find that the spectral line width $\delta\omega_2$ corresponding to the width $\delta\varepsilon$ of the distribution function $f(\varepsilon)$ is equal to

$$\delta\omega_2=\frac{\omega_2^2\delta\varepsilon}{2\omega_1\gamma^2\varepsilon}=\frac{8\omega_1\gamma^2}{(1+\gamma^2\theta_2^2)^2}\frac{\delta\varepsilon}{\varepsilon}. \quad (20)$$

The narrowing of the spectral lines for $\gamma\theta_2>1$ allows a fresh interpretation of the mechanism underlying the increase of the gain G with increasing θ_2. A direct comparison of the expressions (16) and (19) shows that they are connected by the relation

$$G=\frac{16\pi^3N_e}{\varepsilon\omega_2}\frac{d}{d\omega_2}\frac{d\mathscr{E}_{sp}}{d\omega_2 d\Omega_{k_2}}, \quad (21)$$

where in computing the derivative of $d\mathscr{E}_{sp}/d\omega_2 d\Omega_{k_2}$ we should differentiate only the sharpest ω_2 dependence, which is described by the distribution function $f(\varepsilon(\omega_2))$. The formula (21) shows that, as usual, the amplification line contour is proportional to the derivative of the spontaneous emission line contour. But, in the first place, the proportionality factor in (21) contains an increasing θ_2 dependence ($\omega_2^{-2}\propto\theta_2^2$). Furthermore, the narrowing of the spontaneous emission line contour leads to an increase in the derivative of the spectral density of the radiation, an increase which covers the absolute decrease of $d\mathscr{E}_{sp}/d\omega_2 d\Omega_{k_2}$ as θ_2 is increased. Thus, the narrowing of the spontaneous emission line contour as the direction of observation deviates from the direction of motion of the electron beam can be regarded as one of the causes of the increase of the gain G with increasing θ_2.

Naturally, this growth of the coefficient $G(\theta_2)$, (18), with increasing θ_2 is not unlimited. The limitation of the growth of $G(\theta_2)$ is due to the fact that at sufficiently large θ_2 values the dominant mechanism underlying the removal of the singularity in the formulas (4) and (6) is not the energy spread of the electrons, but the finiteness of the interaction time t. For a finite duration of the interaction, the expression (6) for $\Phi_\pm$ assumes in the simplest case of instantaneous switching on of the interaction the form

$$\Phi_\pm^2=t^2\frac{\sin^2(u\pm\Delta u)}{(u\pm\Delta u)^2}, \quad (22)$$

where

$$u=\frac{t}{2}[v(\omega_1\cos\theta_1-\omega_2\cos\theta_2)+\omega_2-\omega_1],$$
$$\Delta u=\frac{t}{4\varepsilon}[\omega_1^2+\omega_2^2-2\omega_1\omega_2\cos\theta-v^2(\omega_1\cos\theta_1-\omega_2\cos\theta_2)^2]. \quad (23)$$

For the angles $\theta_1=\pi$, $\theta_2\ll 1$, these formulas assume the form

$$u=\frac{t}{4\gamma^2}[(1+\gamma^2\theta_2^2)\omega_2-4\gamma^2\omega_1],\quad \Delta u=\frac{\omega_2^2t(1+\gamma^2\theta_2^2)}{4\varepsilon\gamma^2}. \quad (24)$$

From this it follows that the gain, which is proportional to the difference

$$\Phi_+^2-\Phi_-^2=2t\Delta u\frac{d}{du}\frac{\sin^2 u}{u^2}$$

in the case under consideration, can be found with the aid of (4), (9), (17), and (24). It is given by the expression

$$G=\frac{\pi t^3e^4E_1^2N_e}{\varepsilon^3\omega_1}\left(\frac{1-\gamma^2\theta_2^2}{1+\gamma^2\theta_2^2}\right)^2\frac{d}{du}\frac{\sin^2 u}{u^2}. \quad (25)$$

The parameter determining the condition of applicability of the formulas (16) and (18), or the formula (25), is equal to

$$\zeta=\frac{\delta\varepsilon}{\varepsilon}\frac{2^{3/2}\omega_1 t}{1+\gamma^2\theta_2^2}. \quad (26)$$

For $\zeta>1$ the energy spread of the electrons is more important than the finiteness of the interaction time t, and the formulas (16) and (18) are valid. On the other hand, for $\zeta<1$, the energy spread of the electrons can be neglected, but it is necessary to take the finiteness of the interaction time into account. In the second case the gain G is given by the expression (25), and, for $\gamma\theta_2>1$, does not depend on θ_2 (for a given t).

It is important to emphasize that the parameter ζ, (26), significantly depends on the angle θ_2. Therefore, if $\zeta(\theta_2=0)>1$, then the parameter ζ decreases with increasing θ_2, and becomes equal at some value of the angle θ_2 to $\zeta((\theta_2)_0)\approx(2\pi)^{1/2}$; after this the growth of the gain ceases.

For $\gamma\theta_2>1$, the angle $(\theta_2)_0$ up to which the gain G increases is determined by the condition

$$\frac{(2\pi)^{1/2}\gamma^2(\theta_2)_0^2}{2^{3/2}\omega_1 t}=\frac{\delta\varepsilon}{\varepsilon}. \quad (27)$$

Here it should be borne in mind that the interaction time may, depending on the geometry of the experiment, also depend on θ_2. If d is the thickness of the electron beam, then $t=d/\theta_2$. From this it follows that the critical angle value $(\theta_2)_0$ determined from the condition (27) is equal to

$$(\theta_2)_0=2^{1/3}\pi^{1/6}\left(\frac{d}{\lambda_1}\frac{\delta\varepsilon}{\varepsilon}\right)^{1/3}\gamma^{-2/3}, \quad (28)$$

where $\lambda_1=2\pi/\omega_1$ is the pump wavelength.

With allowance for the two considered possibilities, the gain may be given in the general case by the relation

$$G=\frac{2\pi te^4N_eE_1^2\varepsilon^3}{m^4\omega_1\omega_2^2}\left(1-\frac{\omega_2}{2\gamma^2\omega_1}\right)^2\min\left\{2\pi f'(\varepsilon),\frac{\zeta^2}{(\delta\varepsilon)^2}\frac{d}{du}\frac{\sin^2 u}{u^2}\right\}. \quad (29)$$

The dependence $G(\theta_2)$ for $\zeta(0)>1$ and $\theta_2\equiv d/t>1/\gamma$, d/L, where L is the longitudinal dimension of the interaction region, is qualitatively depicted in Fig. 2. In the

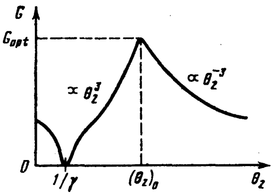

FIG. 2. Dependence of the gain G on θ_2 at $\theta_1=\pi$ and $\zeta(0)>1$.

region $\theta_2>(\theta_2)_0$ the gain $G(\theta_2)$ decreases with increasing θ_2 like θ_2^{-3}, in accordance with the cubic dependence on the interaction time [see (25)]. Therefore, the angle $(\theta_2)_0$, given by the relation (28), is the optimal angle for the construction of the Compton laser.

Let us emphasize that this conclusion is valid only for $\zeta(0)>1$, i.e., for $\delta\varepsilon/\varepsilon>1/2^{3/2}\omega_1 t$. In the opposite case, i.e., for $\zeta(0)<1$ the gain G is determined for all θ_2 by the finiteness of the interaction region and, in accordance with (25), decreases with increasing θ_2 for $t=d/\theta_2$. This indicates that for good beams, i.e., those in which $\delta\varepsilon/\varepsilon<1/2^{3/2}\omega_1 t$, the optimal scheme for the Compton laser is the one-dimensional scheme ($(\theta_2)_0=0$). For beams with a large electron-energy spread, i.e., with $\delta\varepsilon/\varepsilon>1/2^{3/2}\omega_1 t$, the optimal scheme is the one in which the direction of propagation of the amplified wave deviates by the finite angle $(\theta_2)_0$, (28), from the direction of motion of the electron beam.

4. ESTIMATES AND CONCLUSIONS

As shown above, the optimal value of the gain is attained when the condition (28) is fulfilled, i.e., at nonzero values of the angle θ_2. The optimal value G_{opt} of the gain is, as follows from (29), given by the expression

$$G_{opt}=\pi t^3e^4E_1^2N_e/\varepsilon^3\omega_1. \quad (30)$$

Naturally, the expression (30) is valid for both the case of the magnetic undulator and the case of a running pumping wave (the Compton laser). But of special interest is that variant of the FEL in which the role of the pumping wave is played by laser radiation, e.g., radiation from a CO_2 laser.

Let us give estimates for the values of the gain G_{opt} and the corresponding optimum angle $(\theta_2)_0$ for the following electron-beam and laser-wave parameters: J_{max} = 1 kA, $d=0.5$ cm, $l=5$ cm, $\hbar\omega_1=0.1$ eV, $\delta\varepsilon/\varepsilon=10^{-3}$, $\gamma=20$, and $E_1=5\times10^7$ V/cm. For these parameters G_{opt} and $(\theta_2)_0$ turn out to be respectively equal to $G_{opt}\sim$ 1%; $(\theta_2)_0\approx0.2$. The energy of a photon of the amplified wave does not in this case depend on the quantity γ, and is equal to $\hbar\omega_2\approx10$ eV.

Thus, the above-presented estimates indicate the possibility of obtaining an appreciable amplification factor in the ultraviolet region with the use of a CO_2 laser as the pump.

Let us note finally that, in the nonrelativistic case, the amplification in the Compton laser in the noncollinear scheme has been considered by Dubrovskiĭ et al.[19]

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Intermolecular and intramolecular distribution of vibrational energy under multiphoton excitation by IR laser radiation

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The vibrational distribution produced in molecular gases by IR laser pulses is investigated by high-time-resolution Raman spectroscopy. Quantitative data are obtained and interpreted regarding an ensemble of molecules excited during the IR pulse under both collisionless conditions and conditions involving collisions with buffer-gas particles. The energy threshold at which the vibrational energy becomes randomly distributed among the vibrational modes (becomes randomized) during the interaction with the IR pulse is determined for the SF_6 and CF_3I molecules. A possible mechanism for randomization of the vibrational motion in SF_6 is considered.

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1. INTRODUCTION

A major question in the study of multiphoton excitation of molecules by IR irradiation (see Ref. 1) is the nonequilibrium vibrational distribution produced in a molecular gas as a result of interaction with a pulsed IR field. Here, there are two problems whose solution is extremely important for an understanding of the multiphoton excitation process. First of all, the question concerning the form of the vibrational energy distribution function of the molecules in the gas is of interest, i.e., the form of the intermolecular vibrational energy distribution. Secondly, it is important to know the distribution of absorbed energy among the different vibrational degrees of freedom of the molecule, i.e., the form of the intramolecular vibrational energy distribution.

A variety of approaches are used to study the vibrational energy distribution: various semiempirical models[2,3] have been developed and direct experiments have been performed. In the latter case, it was possible to progress rather far using spectroscopic methods of probing the excited states. Thus, in experiments with OsO_4 (Ref. 4) and SF_6 (Ref. 5), it was shown with the IR probe method that the distribution produced in multiphoton excitation consists of two molecular ensembles: excited molecules in a vibrational quasicontinuum, and unexcited molecules. The fraction q of molecules in the upper ensemble was measured. An analogous conclusion concerning the existence of ensembles of "hot" and "cold" molecules was drawn from a study of multiphoton absorption and dissociation of CF_3I (Ref. 6), and on the basis of UV probing of SF_6

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