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Generation of hard quanta in the cross-grooved resonator of a free electron laser

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***Abstract*. A new approach to the generation of powerful tunable hard radiation in the cross-grooved resonator (CGR) of a free electron laser is proposed and discussed. Gamma quanta are generated upon backward Compton scattering of intracavity radiation by the electron beam of a free electron laser with a CGR. The use of a CGR considerably increases the power of intracavity radiation, and hence the power of gamma radiation, and also solves the problem of extraction of hard radiation by eliminating the losses in the depth of the cavity mirror, which were inevitable earlier.**

Keywords: *gamma quanta, cross-grooved resonator, free electron laser.*

1. Introduction

A coherent and tunable source of gamma and X-rays is essential for various applications. Such sources can be used for research in the fields of solid state physics, nuclear physics, in medicine for diagnostic purposes (tomography) and precision irradiation of tumours, because coherent radiation can be focused more easily, which provides energy supply at the desired small spot.

The possibility of generation of coherent gamma and X-rays upon backward Compton scattering of laser quanta by a beam of relativistic electrons was first analysed in Refs [1, 2]. For $\gamma \gg 1$, the wavelength of hard radiation is [3]

$$\lambda_X \approx \frac{\lambda_{las}}{4\gamma^2}(1+\gamma^2\theta^2), \quad (1)$$

where λ_{las} is the wavelength of laser radiation; γ is the electron energy in units of mc^2; m and c are the electron mass and the velocity of light, respectively; and θ is the observation angle measured from the direction of propagation of the electron beam. Note that the wavelength of the emitted quanta depends strongly on the electron energy, which makes it possible to obtain short-wavelength quanta even for a moderately relativistic beam. Thus, assuming that $\lambda_{las} = 1$ μm and $\gamma = 100$, we obtain $\lambda_X = 0.025$ nm for a strictly backward direction ($\theta = 0$), corresponding to the energy of X-rays of the order of 50 keV.

The use of intracavity backward Compton scattering of radiation from a free electron laser (FEL) by the electron beam of this laser for generating tunable gamma radiation was proposed in Ref. [4]. The hard radiation generated in this case not only can be tuned over a wide frequency range, but be also polarised. The wavelength of gamma quanta is described by the expression, which is a generalisation of formula (1) (see Ref. [5] for the basic concepts and relations in the theory of FELs):

$$\lambda_X \approx \frac{L}{8\gamma^4}\left(1+\frac{K^2}{2}\right)(1+\gamma^2\theta^2), \quad (2)$$

where L is the period of the undulator; $K \sim H$ is the undulator parameter; and H is the magnetic field of the undulator.

The validity of the main theoretical estimates was demonstrated in several experiments. Consider in details one of the latest experiments of this type [6], in which an infrared FEL tunable in the 3.5 – 7.0-μm range was used. The peak current in an electron micropulse of duration 8 ps in a linear rf accelerator was equal to 100 A. A macropulse of duration 11 μs consisted of micropulses separated by a 16-ns interval. Two micropulses of an electron beam and the corresponding light beam were confined in a 4.8m-long optical resonator. The collision of the electron micropulse ($\gamma \approx 100$) with the backward scattered light micropulse at the centre of the resonator resulted in the generation of X-rays having energy in the 7 – 14-keV range.

To reduce the absorption of hard radiation, one of the cavity mirrors was made of beryllium, and the thickness of the central part of the mirror of diameter 12 mm was reduced to 1.7 mm. To minimise the IR radiation losses, the surface of the beryllium mirror was additionally coated with a layer of gold of thickness of about 100 μm.

Calculations made in Ref. [6] using the above parameters lead to the following numbers of hard radiation quanta: 4×10^2 for a micropulse and 2×10^5 for a macropulse. For a macropulse repetition rate of 25 Hz, the quantum generation rate is 5×10^6 s^{-1}.

The experimental values of the parameters obtained in Ref. [6] were found to be about an order of magnitude lower than the above data, and the authors of Ref. [6] attribute the discrepancy mainly to a decrease in the intensity of the intracavity radiation due to a damage of the gold coating by powerful IR radiation (losses per trip in the cavity amounted to 6 %), as well as to a partial absorption of

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hard quanta in the cavity mirror. Besides, further degradation of the mirror surface during prolonged operation is inevitable due to the impact of hard radiation in the region of the laser field maximum, which leads to even greater intracavity losses.

In this work (see also [7]), we propose to use a cross-grooved resonator (CGR) as the optical cavity of a FEL [8]. In this case, the difficulties mentioned in Ref. [6] are eliminated while the intracavity radiation intensity and the intensity of hard radiation can be increased considerably, and the loss of gamma quanta in the depth of the resonator mirror is avoided (Fig. 1).

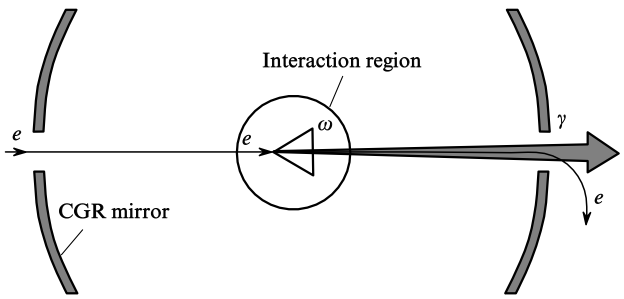

Figure 1. Generation of hard quanta in the CGR of a free electron laser.

2. Basic postulates

Before analysing the main features of generation of hard quanta in the cross-grooved resonator of an FEL, we will derive some basic relations describing Compton scattering, which will be useful for evaluating the radiation parameters. These relations follow from general classical expressions in the limits

$$1 \ll \gamma \ll \frac{m}{\omega} \tag{3}$$

in which we are interested. Unless otherwise specified, it is assumed hereafter that, $\hbar = c = 1$, and ω is the laser radiation frequency. The left-hand side of (3) corresponds to the relativistic electron beam, while the right-hand side allows the recoil effect to be neglected during collisions. For a laser radiation wavelength ~ 1 μm, inequalities (3) lead to a virtually unbounded interval of electron energies: $1 \ll \gamma \ll 10^6$.

Two different approaches are normally used for calculating the intensity of Compton radiation. One of them is based on the analogy between an electromagnetic wave and an undulator, and the electron emission in the undulator is correctly described by the known relations in this case. A drawback of this approach is that the analytic expressions are cumbersome and lead only to rough (essentially qualitative) approximations.

In the second approach, the hard radiation parameters are evaluated by using directly the expressions for the differential cross section of Compton scattering, which are also quite cumbersome as in the preceding case. This complicates the estimates of the scattering parameters.

However, in the problem under study, in which hard quanta are assumed to be generated in the optical cavity of a FEL, the inequalities (3) are satisfied in most cases. It will be shown below that this helps in simplifying the expressions used for estimating the Compton radiation parameters, making them more visual and convenient for practical applications.

Consider first the expression for the frequency of Compton quanta (see, for example, Ref. [3]):

$$\omega_2 = \omega_1 \frac{1 - \beta\cos\theta_1}{1 - \beta\cos\theta_2 + (1 - \cos\theta)\omega_1/E}, \tag{4}$$

where ω_1 is the laser frequency; $E = \gamma mc^2$ is the electron energy; $\gamma = 1/(1-\beta^2)^{1/2}$; $\beta = v/c$; v is the electron velocity; θ_1 and θ_2 are the angles of collision and scattering measured relative to the initial direction of propagation of electrons; and $\theta = \theta_1 - \theta_2$ (Fig. 2).

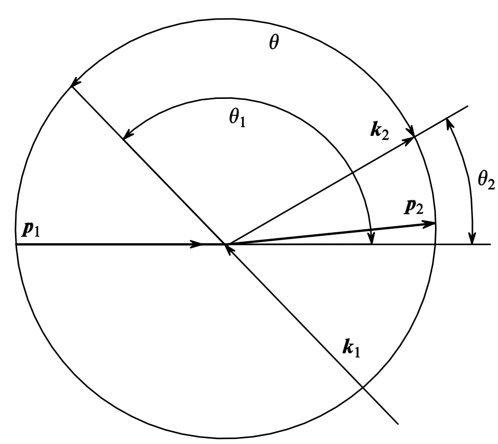

Figure 2. Geometry of collisions between electrons and laser radiation quanta.

It follows from the condition $\gamma\omega_1 \ll m$ [see formulas (3)] that the last term in the denominator of Eqn (4) can be omitted, i.e.,

$$\omega_2 = \omega_1 \frac{1 - \beta\cos\theta_1}{1 - \beta\cos\theta_2}. \tag{5}$$

Note once more that the omission of the term in the denominator corresponds to the neglect of the recoil momentum during a collision.

Let us now determine the maximum possible frequency $\omega_{2\max}$ for a given angle θ_1, which is attained for $\theta_2 = 0$:

$$\omega_{2\max} = \omega_1 \frac{1 - \beta\cos\theta_1}{1 - \beta} \approx 2\gamma^2\omega_1(1 - \beta\cos\theta_1). \tag{6}$$

Two characteristic cases of scattering are of special interest: transverse scattering (superscript t) corresponding to $\theta_1 = \pi/2$, and the backward longitudinal scattering (superscript b) corresponding to $\theta_1 = \pi$. For these two cases, we can write

$$\omega^{\mathrm{b}}_{2\max} = \omega_1 \frac{1 + \beta}{1 - \beta} \approx 4\gamma^2\omega_1, \tag{7}$$

$$\omega^{\mathrm{t}}_{2\max} = \omega_1 \frac{1}{1 - \beta} \approx 2\gamma^2\omega_1, \tag{8}$$

i.e., the frequency of Compton quanta in the longitudinal case is double the frequency in the case of transverse scattering.

The variation of the collision angle θ_1 can be used for frequency scanning of the Compton quanta, and hence it should be interesting to analyse the dependence of $\omega_{2\,\max}$ on θ_1. Consider small deviations of the angle θ_1 from the limiting angles of strictly longitudinal and transverse scattering:

$$\omega_{2\,\max} \approx \begin{cases} 4\gamma^2\omega_1(1-\alpha^2/4), & \theta_1 = \pi - \alpha, \\ 2\gamma^2\omega_1(1-\alpha), & \theta_1 = \pi/2 - \alpha. \end{cases} \tag{9}$$

One can see from these relations that longitudinal scattering is characterised by high frequency stability relative to the spread of the angles of collision or to the emittance of the angular beam.

However, transverse scattering is characterised by linear dependence of frequency on the collision angle, which in turn may be used for the angular frequency scanning of Compton quanta.

Neglecting the recoil momentum for $\gamma \gg 1$ and for quite small angles $\theta \ll 1$, we can write $1 - \beta\cos\theta_2 \approx (1 + \gamma^2\theta_2^2)/2\gamma^2$ in the general expression for frequency, which gives

$$\omega_2 = \frac{\omega_{2\,\max}}{1 + \gamma^2\theta_2^2}. \tag{10}$$

It follows from this relation that a significant variation of the frequency of Compton quanta as a function of the scattering angle occurs at angles $\sim 1/\gamma$. Note that the differential cross section of Compton scattering has the same scale of variation in the scattering angle.

Consider now the differential scattering cross section. In this case, we shall proceed from the general relation presented, for example, in Ref. [3] (p. 308), in which the differential cross section is expressed in terms of the kinematic invariants s, t, u:

$$\mathrm{d}\sigma = 8\pi r_e^2 \frac{m^2 \mathrm{d}t}{(s-m^2)^2}(A^2 + A - B/4), \tag{11}$$

where $r_e = e^2/m$. Parameters A and B that depend on kinematic invariants s, t, u, as well as the kinematic invariants themselves, are defined by the following relations:

$$A = \frac{m^2}{s-m^2} + \frac{m^2}{u-m^2}, \quad B = \frac{s-m^2}{u-m^2} + \frac{u-m^2}{s-m^2}, \tag{12}$$

$$s = (p+k)^2 = (p'+k')^2 = m^2 + 2pk = m^2 + 2p'k',$$

$$t = (p-p')^2 = (k-k')^2 = 2(m^2 - pp') = -2kk',$$

$$u = (p-k')^2 = (p'-k)^2 = m^2 - 2pk' = m^2 - 2p'k,$$

where p, p' are the 4-momenta of electrons before and after the collision; and k, k' are the 4-momenta of quanta before and after the collision. Using the law of conservation $p + k = p' + k'$ for the 4-momentum, we can transform expression (12) for A and B as follows:

$$A = -\frac{m^2}{2}\frac{kk'}{(pk)(pk')}, \quad B = -2 - \frac{(kk')^2}{(pk)(pk')}. \tag{13}$$

If the inequalities (3) are satisfied, $A \approx (1 + \gamma^2\theta_2^2)^{-1} \sim 1$, while $(kk')^2/(pk)(pk') \ll 1$, and hence we can put $B \approx -2$, which gives the following expression for the differential cross section:

$$\mathrm{d}\sigma = 8\pi r_e^2 \frac{m^2 \mathrm{d}t}{(s-m^2)^2}\left(A^2 + A + \frac{1}{2}\right). \tag{14}$$

Consider now the differential $\mathrm{d}t$. Recall that $t = 2\omega_1\omega_2[1 - \cos(\theta_1 - \theta_2)]$. Differentiating this expression with respect to ω_2 or θ_2 and neglecting small terms, we come to the conclusion that in the limits considered above, the differential $\mathrm{d}t$ can be written in the form

$$\mathrm{d}t = 2\omega_1(1 - \cos\theta)\mathrm{d}\omega_2. \tag{15}$$

Because there is a one-to-one correspondence between ω_2 and the scattering angle θ_2 [see (5)], we obtain expressions for the differential cross section that have a one-to-one dependence both on the scattering angle and the frequency of scattered Compton quanta.

Let us now simplify expression (12) for A by using (3). We assume that $\beta = 1$ everywhere, except in expressions of the type $1 - \beta$ and $1 - \beta\cos\theta_2$, where the deviation of β from unity is significant. As a result, we obtain

$$A = -\frac{\omega_2}{\omega_{2\,\max}},$$

which in turn leads to the following expression for $\mathrm{d}\sigma$:

$$\frac{\mathrm{d}\sigma}{\mathrm{d}x} = 8\pi r_e^2\left(x^2 - x + \frac{1}{2}\right), \tag{16}$$

where the parameter $x = |A| = \omega_2/\omega_{2\,\max}$ has been introduced for the sake of convenience. Taking into account expression (10), which assumes the form $x = (1 + y^2)^{-1}$ (where $y = \gamma\theta_2$), we arrive at the following expression for the angular dependence of the differential scattering cross section:

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\Omega} = 8\gamma^2 r_e^2 x^2\left(x^2 - x + \frac{1}{2}\right). \tag{17}$$

Here, $\mathrm{d}\Omega = 2\pi\sin\theta_2\mathrm{d}\theta_2 \approx \gamma^{-2}\pi\mathrm{d}y^2$ and x is a function of y or θ_2.

Note that if the expression for $\mathrm{d}\sigma/\mathrm{d}\theta_2$ is integrated with respect to the scattering angle θ_2 from 0 to ∞, we arrive at the correct expression for the total Thomson scattering cross section $\sigma = 8\pi r_e^2/3 = 6.63 \times 10^{-26}$ cm^2, as expected in the adopted approximation.

Let us now write down the expression for estimating the hard quantum generation rate:

$$\frac{\mathrm{d}N_\gamma}{\mathrm{d}t} \approx \frac{J\sigma n}{\hbar\omega_1}, \tag{18}$$

where J is the power density of the scattered laser radiation; $n \approx N_e V$ is the total number of electrons with concentration N_e in the effective interaction volume V.

In the case of scattering from a single electron pulse of duration $\tau = l/c$, the number of hard quanta is given by the relation

$$N_\gamma^{(1)} \approx \frac{dN_\gamma}{dt}\tau = \frac{J\sigma n\tau}{\hbar\omega_1}. \qquad (19)$$

Assuming that the electron pulses are injected into the interaction volume with a frequency f, we obtain the following expression for the mean quantum generation rate:

$$\left\langle \frac{dN_\gamma}{dt} \right\rangle = fN_\gamma^{(1)} = \frac{J\sigma n f\tau}{\hbar\omega_1}. \qquad (20)$$

Let us also estimate the total power P of hard quanta:

$$P \approx \hbar\omega_{2\,\max}\left\langle \frac{dN_\gamma}{dt} \right\rangle = \frac{\omega_{2\,\max}}{\omega_1} J\sigma n f\tau. \qquad (21)$$

By estimating the interaction volume as $V \approx Sl = Sc\tau$ (S is the electron beam cross section, $l = c\tau$ is the length of laser and electron pulses), we assume that the diameters of the electron and laser beams are almost equal in the region of the scattering point, which corresponds to optimal scattering. In this case, the number of quanta generated by a single pulse is

$$N_\gamma \approx \frac{dN_\gamma}{dt}\tau = \frac{WI_e\tau^2\sigma}{Se\hbar\omega_1}, \qquad (22)$$

where W is the power of the pulsed laser radiation circulating in the resonator; and I_e is the total peak electron current. This relation leads to an expression for the mean hard quantum generation rate for a given frequency f of injection of electron pulses in the optical resonator:

$$\left\langle \frac{dN_\gamma}{dt} \right\rangle = f\frac{WI_e\tau^2\sigma}{Se\hbar\omega_1}. \qquad (23)$$

Relations (22) and (23) can be used for estimates and preliminary optimisation of the emitter of hard quanta in electron beam and optical resonator parameters.

Let us now write down some energy relations that will be useful for estimating the total and differential power of generation of hard quanta:

$$\frac{dP}{d\omega_2} \approx Jnf\tau\frac{\omega_2}{\omega_1}\frac{d\sigma}{d\omega_2} = 8\pi r_e^2\frac{Jnf\tau}{\omega_1}x\left(x^2 - x + \frac{1}{2}\right). \qquad (24)$$

Similarly, we can write for the angular distribution of the emitted power:

$$\frac{dP}{d\Omega_2} \approx 8\gamma^2 r_e^2 Jnf\tau\frac{\omega_{2\,\max}}{\omega_1}x^3\left(x^2 - x + \frac{1}{2}\right). \qquad (25)$$

Finally, integrating the above differential relations, we obtain the total power of emission of hard quanta from the interaction volume V:

$$P \approx \frac{4\pi r_e^2}{3}Jnf\tau\frac{\omega_{2\,\max}}{\omega_1}. \qquad (26)$$

All the expressions presented above are clear and simple, and can be used for preliminary estimates of the Compton radiation parameters.

3. On the applicability of CGR for the generation of hard quanta in the backward Compton scattering process

Let us briefly recall the construction of a CGR [8]. In contrast to the conventional grooved resonators, the generatrix whose rotation around the CGR symmetry axis forms a concave mirror is inclined at a small angle to this axis. Fig. 3 shows the resonator geometry and the distribution of the fundamental mode field, which is circular at the mirrors and is focused at the centre of the resonator.

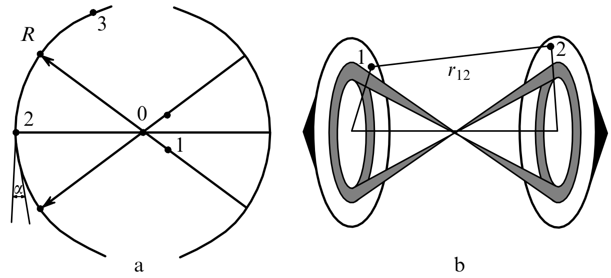

Figure 3. (a) Resonator geometry and (b) field distribution in the resonator (in perspective); (0) resonator centre; (1) centre of curvature of the arc 2–3 forming the mirror surface upon rotation around the symmetry axis 0–2.

The beams propagating between the mirrors cross one another and occupy a considerable part of the central region of the CGR, its length being limited by the outer caustic surface. At the mirrors, these beams uniformly fill the circular region. For such a field structure, the central aperture in the mirror does not significantly affect the oscillation parameters. Moreover, the above-mentioned geometry of the field makes CGR quite attractive for application in high-power FELs [9].

It was shown in Refs [10–14] that steady-state oscillations may exist in the above geometrical structures, while the field distribution for the fundamental resonator mode is indeed circular at the mirrors and focused at the centre of the resonator, where the FEL undulator is assumed to be installed. The operation of a CGR in the millimetre wavelength region was simulated experimentally in Ref. [15], and the main theoretical and computational results were confirmed by such a simulation.

The above properties of CGR lead to the qualitative conclusion that the generation of coherent tunable radiation using an FEL with a CGR allows us to

– make a central aperture in the resonator mirror (at the centre of symmetry) for the extraction of hard radiation quanta without loss in the depth of the mirror; the presence of the central aperture does not affect the intensity of the intracavity radiation since the fundamental mode losses do not increase because of its circular distribution at the mirror surface;

– attain a natural mutual adjustment of the electron and light beams;

– sharply increase the radiation intensity inside the resonator, and hence the intensity of the hard radiation also by using, for example, a cooled optical system; the thickness and the structural features of the mirror do not affect the intensity of hard radiation emerging from the resonator, unlike for traditional resonators;

– prevent the degradation of the mirror surface by hard radiation since it does not touch the reflecting surface and thus allows a virtually unlimited service life of the equipment.

4. Some practical estimates

In contrast to the traditional optical resonators, there are no main factors causing a decrease in the output of hard quanta from a CGR or a degradation of the optical resonator surface. In this connection, it can be assumed that the estimates obtained by using the qualitative relations presented in Section 2 will be closer to the expected experimental results.

In order to estimate the potential of the emitter of hard quanta, it would be expedient to estimate the number of quanta and the mean rate of their generation for two known FELs: a Mark-III laser [16–18], one of the first free electron lasers, and a cut microtron-recuperator high-power FEL designed at the Institute of Nuclear Physics (INP), Siberian Branch, Russian Academy of Sciences [19] (see also references in the review [20], which contain a large body of data presented in Table 1 and used for estimates). The results of these estimates are presented in Table 1. A comparison of the parameters for these two devices leads to the conclusion about the possible range of hard quantum generation rates.

The radiation from a Mark-III laser is a sequence of macropulses of duration 3 μs and a pulse repetition rate of 10–30 Hz. Each macropulse consists of about 8500 micropulses of duration 1 ps and energy 6 μJ. It is expected that the radiation from the FEL INP is a sequence of micropulses of duration about 20 ps, energy of the order of 1 mJ, and a pulse repetition rate up to 180 MHz. In both cases, it is assumed that the area of the light spot in the interaction region is 0.1 cm^2.

Of course, it follows from expressions (22) and (23) that the beam cross section S considerably affects the efficiency of generation of hard quanta in specific experiments. It should be borne in mind that there is no sense in making the light beam diameter smaller than the diameter of the electron beam. Moreover, the geometry of the optical resonator, which determines the diameter of the light beam, considerably affects the operation of the FEL, i.e., the power W of the pulsed laser radiation circulating in the resonator [see Eqns (22) and (23)]. Therefore, an additional optimisation in parameter S is required for specific experiments.

Recall that the average rate $\langle \mathrm{d}N_\gamma/\mathrm{d}t \rangle$ of generation of hard quanta for a microtron is the final working (user) parameter of the emitter, averaged over the operation time of the device. Table 1 contains the rate of generation of hard quanta for the linear accelerator Mark-III, averaged over the macropulse duration $t = 3$ μs. Hence, in order to obtain the final user parameter $\langle\langle \mathrm{d}N_\gamma/\mathrm{d}t \rangle\rangle$, we must carry out additional averaging over the macropulse injection frequency $\Omega = 10 - 30$ Hz, i.e.

$$\left\langle\left\langle \frac{\mathrm{d}N_\gamma}{\mathrm{d}t} \right\rangle\right\rangle \approx \left\langle \frac{\mathrm{d}N_\gamma}{\mathrm{d}t} \right\rangle \Omega t.$$

Choosing the injection frequency equal to 30 Hz, the generation rate averaged over the operation time of the device is $\langle\langle \mathrm{d}N_\gamma/\mathrm{d}t \rangle\rangle \approx 1.9 \times 10^3\ \text{s}^{-1}$.

For comparison, we present the data on the number of hard quanta in a micropulse and the quantum generation rate, obtained in the above-mentioned experiment [6]. The number of hard quanta emitted in a micropulse is $N_\gamma^{(1)} \approx 400$, which is much larger than in Mark-III ($N_\gamma^{(1)} = 0.75 \times 10^{-2}$) and is comparable with the data for the microtron-recuperator ($N_\gamma^{(1)} = 83$).

The rate of generation of hard quanta, averaged over the macropulse, is $\langle \mathrm{d}N_\gamma/\mathrm{d}t \rangle \approx N_\gamma^{(1)}/t_\mu \approx 2.5 \times 10^{10}\ \text{s}^{-1}$ ($t_\mu = 16$ ns is the micropulse-repetition interval). This figure is also comparable with the data for a microtron ($1.5 \times 10^{10}\ \text{s}^{-1}$) and is much larger than the rate of generation of hard quanta for Mark-III ($2.1 \times 10^7\ \text{s}^{-1}$).

The situation changes radically in the case of averaging over time considerably exceeding the micropulse-repetition interval. While the generation rate for a microtron is $1.5\times 10^{10}\ \text{s}^{-1}$ as before, it is much smaller for Mark-III accelerator as well as in the experiment [6]. In the latter case, the average rate of generation of hard quanta for a macropulse injection frequency $f = 25$ Hz is $\langle \mathrm{d}N_\gamma/\mathrm{d}t \rangle = 5 \times 10^6\ \text{s}^{-1}$ (compared to $1.9 \times 10^3\ \text{s}^{-1}$ for Mark-III and $1.5 \times 10^{10}\ \text{s}^{-1}$ for the microtron).

Note once more that the actual figures in the experiment [6] were much lower due to a degradation of the mirror surface caused by high-power laser radiation and the hard Compton component, as well as the absorption of hard quanta in the mirror. It is assumed in the present work that the use of a CGR as the optical resonator will eliminate all three above-mentioned sources of decrease in the hard quantum generation rate.

5. Conclusions

Thus, it should be emphasised that the above peculiarities of application of a CGR in a high-power FEL (relatively high intensity of the hard radiation and a long service life) will make it possible to create tunable lasers of hard quanta for solving specific problems in pure and applied physics, as well as for the latest medical applications.

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Table 1. FEL parameters and qualitative estimates of Compton radiation parameters.

Lasers	U/MeV	I_p/A	τ/ps	λ/μm	$\hbar\omega/10^{-19}$ J	f/GHz	$n/10^8$	J/W cm^{-2}	$N_\gamma^{(1)}$	$\left\langle \frac{\mathrm{d}N_\gamma}{\mathrm{d}t} \right\rangle/\text{s}^{-1}$
Mark-III	44	20	1	3	0.66	2.8	1.25	6×10^7	7.5×10^{-3}	2.1×10^7
FEL INP	98	100	20	2	1	0.18	125	5×10^8	83	1.5×10^{10}

Note: I_p is the amplitude of the electron current in the beam, and n is the electron concentration.

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