strains of the crystal and to the appearance of interaction of excitons with the metal. Thus, from the different behaviors of $\Gamma(T)$ and $\omega(T)$ of an anthracene crystal placed on a contact with a gold layer and with quartz one can conclude only qualitatively that the metal exerts an influence on the exciton-phonon interaction in the anthracene. It is of interest to obtain theoretical relations for $\Gamma(T)$ and $\omega(T)$ with account taken of the influence of the considered metallic quenching of the excitons.

Thus, the measurements of the luminescence spectra of single crystals of anthracene on a contact with gold at low temperatures have made it possible to observe experimentally, for the first time ever, metallic quenching of Frenkel excitons in the low-temperature region. It is concluded that the coherent and noncoherent excitons are differently quenched. Finally, it is shown that metallic quenching of excitons influences the exciton-phonon interaction in crystals.

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Coherent excitation of inversion of nuclei by a modulated beam of randomly distributed relativistic electrons

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The possibilities are analyzed of obtaining a coherent flux of resonant pseudophotons for the excitation of an inversion state in Mössbauer nuclei. The spectral composition of radiation accompanying the motion of a modulated beam of superrelativistic electrons whose space-time distribution is described by Poisson statistics is considered. The emission spectrum has coherent and noncoherent components at the beam-modulation frequency. The flux density of the coherent component of the radiation, at realistic parameters of the beam, of the modulator, and of the accelerator, satisfies the γ-amplification threshold condition.

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The continuing active discussion of the problem of developing a gamma laser stimulates the search for optimal mechanisms of excitation of the inversion in such devices. In view of the considerable principal difficulties in the realization of a system of long-lived gamma transitions, due to the need of eliminating vanishingly small perturbations and stimulated narrowing of the line width by 5–8 orders of magnitude, a search is being made for more realistic results on the basis of another system, which makes use of different variants of a system based on short-lived Mössbauer nuclei with natural line width. As a result of the high spectral pump density $P(\omega)$ needed to realize such a model, the source must satisfy a number of conditions, one of which, for the case of radiation pumping, reduces to the requirement that the pump spectrum band be narrow $\Delta\omega \lesssim \Gamma\hbar$ ($\Gamma/\hbar$ is the width of the activation level of the excited nuclei). If the inverse inequality $\Delta\omega \gg \Gamma/\hbar$ is satisfied, the possible strong heating and destruction of the working medium make the use of the Mössbauer effect practically impossible, and without this effect the problem cannot be solved. From this point of view [as well as with

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respect to obtaining a maximum value of $P(\omega)$], the analyzed coherent source of pump photons is superior to the presently known noncoherent ones—bremsstrahlung, synchrotron, or a source using the emission of oscillating charged particles in a crystal channel.

We analyze below the possibility of obtaining the necessary flux of photons (pseudophotons) via interaction of arbitrary resonant nuclei with a density-modulated beam of relativistic electrons having an initial Poisson space-time distribution that can be obtained, for example, via the shot effect of the cathode. The only previously considered model of Coulomb coherent excitation by a completely ordered and strictly periodic discrete-electron beam (electronic pseudocrystal)[1,2] is unfortunately utterly unrealistic.

We consider the following model. A quasi-monochromatic electron beam with a random distribution is density-modulated by a light wave (the possibility of such a modulation was confirmed by the experiment of Schwarz and Hora[3] and by subsequent more refined theoretical investigations[4-6]). The beam is injected next into an electron accelerator, where its velocity is increased to the relativistic value $v \lesssim c$, and the period of the modulation in the laboratory system is decreased by a factor $\gamma = (1 - v^2/c^2)^{-1/2}$. It is apparently possible to use linear resonant accelerators for this purpose. We seek the solution of the problem by the Weizsäcker-Williams pseudophoton method,[7,8] which does not require consideration of the concrete mechanism of the coherent electromagnetic interaction (Coulomb, bremsstrahlung, transition, etc.), in which it is sufficient to use the spectral intensity of the pseudophotons and the known cross-section σ of resonant radiative excitation.

The flight of one (n-th) superrelativistic electron out of the beam produces at a point located at the impact distance ρ_n a transverse electric field

$$E_\perp(t-t_n, \rho_n) = e\gamma\rho_n/[\rho_n^2 + v^2(t-t_n)^2\gamma^2)^{3/2}.$$

The entire beam of N electrons produces therefore a field

$$E_\perp(t) = \sum_{n=1}^{N} E_\perp(t-t_n, \rho_n).$$

To find the spectral density of the flux of pseudophotons that accompany the beam it is necessary to calculate the correlation function of the modulated nonstationary process $E_\perp(t)$. The two-time correlation function is of the form[9]

$$B(t,\tau) = \overline{E_\perp(t)E_\perp(t-\tau)} = \int_{-\infty}^{\infty} n(\theta)\langle E_\perp(t-\theta,\rho)E_\perp(\tau-\theta,\rho)\rangle\, d\theta$$
$$+ \int_{-\infty}^{\infty} n(\theta)\langle E_\perp(t-\theta,\rho)\rangle\, d\theta \int_{-\infty}^{\infty} n(\theta)\langle E_\perp(t-\tau-\theta,\rho)\rangle\, d\theta,$$

where $n(t)$ is the variable parameter of the Poisson distribution (the instantaneous temporal density of the beam at a specified point of the laboratory system), and the angle brackets denote averaging over the impact parameter. To simplify the calculations it is useful to neglect the very small change of the particle velocity upon absorption of the pseudophotons; this is permissible at a small target thickness and high electron energies.

Changing over, with the aid of the operation of smoothing over the period of the modulation in the laboratory system T_0 (Ref. 10), to the single-time correlation function $B(\tau)$, we can find the spectral density of the correlation, corresponding to the spectral density of the pseudophoton energy:

$$G(\omega) = \frac{1}{2\pi}\int_{-\infty}^{\infty} e^{-i\omega\tau}\, d\tau \frac{1}{T_0}\int_{-T_0/2}^{T_0/2} B(t,\tau)\, dt,$$

where $T_0 = T_0'/\gamma$, T_0' is the period of oscillations of the modulating light wave. For the sake of argument we consider a beam characterized by a density $\bar{n}(t) = \bar{n}(1 - \delta\cos\Omega_0 t)$ at $-T/2 \leqslant t \leqslant T/2$ and $n(t) = 0$ outside the limits of this interval; $\Omega_0 = 2\pi/T_0$, $\delta \leqslant 1$, $\bar{n} = N/T$.

Taking into account only the time-dependent part of the density $n(t)$, which leads to the coherent effect, we find

$$G(\omega) = \frac{e^2N\delta|\omega|}{T\pi v^4\gamma^2}|\omega\pm\Omega_0|\left\langle K_1\left(\frac{|\omega|\rho}{v\gamma}\right)K_1\left(\frac{|\omega\pm\Omega_0|\rho}{v\gamma}\right)\right\rangle$$
$$\times\frac{\sin(\omega\pm\Omega_0)T_0/2}{(\omega\pm\Omega_0)T_0/2} + \frac{e^2N^2\delta^2\Omega_0^2}{2\pi Tv^4\gamma^2}\left\langle K_1\left(\frac{|\Omega_0|\rho}{v\gamma}\right)\right\rangle^2$$
$$\times\frac{\sin(\omega\pm\Omega_0)T_0/2}{(\omega\pm\Omega_0)T_0/2}\,\frac{\sin(\omega\pm\Omega_0)T/2}{(\omega\pm\Omega_0)T/2},$$

where $K_1(x)$ is the MacDonald function.

It follows from physical considerations[8] that the possible values of the impact parameter ρ at which pseudophoton absorption (i.e., excitation of the inversion) is possible lie in the interval $\rho_{min} \leqslant \rho \leqslant \rho_{max}$, where $\rho_{min} = e^2/\gamma m v^2$ and corresponds to the quantum limit, while ρ is equal to the smaller of the two possible values corresponding to the threshold of the adiabatic interaction $\rho^{(1)} = v\gamma/\Omega_0$ or the electron-screening parameter $\rho^{(2)} \approx 1{,}4a_0Z^{-1/3}$, where $a_0 = \hbar^2/me^2$. It is easy to verify that for the necessary parameters $v \sim c$, $\Omega_0 \sim 10^{19}$Hz, $\gamma \sim 10^4$ we have

$$\rho_{max}^{(2)} \sim 2\cdot10^{-8}\text{ cm} \ll \rho_{max}^{(1)}$$

and the argument of $K_1(x)$ is small. Then, using the distribution function of ρ in the form $w(\rho) = 2\pi\rho/S$, where S is the cross section area of the electron beam, we get

$$G(\omega) = \frac{2e^2N\delta}{Tv^2S}\ln\frac{\rho_{max}}{\rho_{min}}\frac{\sin[(\omega\pm\Omega_0)T_0/2]}{(\omega\pm\Omega_0)T_0/2}$$
$$+\frac{2\pi e^2N^2\delta^2}{Tv^2S^2}(\rho_{max}-\rho_{min})^2\frac{\sin[(\omega\pm\Omega_0)T_0/2]}{(\omega\pm\Omega_0)T_0/2}\,\frac{\sin[(\omega\pm\Omega_0)T/2]}{(\omega\pm\Omega_0)T/2}$$

and the spectral density of the pseudophoton flux

$$P(\omega) = G(\omega)c/\hbar\omega.$$

The last expressions correspond to a sum of a noncoherent term (linear in N) and a coherent term (proportional to N^2). It is seen that the spectrum of the coherent pseudophotons is localized in the region of the frequencies $\Delta\omega \sim 1/T$ near $\omega = \pm\Omega_0$. For accelerators of energy $W \sim 5$ GeV the frequency is $\Omega_0 \sim 10^{19}$ Hz, so that inversion can be excited on resonant transitions of nuclei. In fact, recognizing that the modulating optical laser radiation has a finite spectral width $\Delta\Omega_0' = \Delta\Omega/\gamma \sim 10^3 - 10^6$ Hz, we must average the expression for $P(\omega)$ over Ω_0 in accord with the spectral distribution of the variable electron density

$$\tilde{n}(t)=\int_{\Delta\Omega_0} n(\Omega_0)e^{i\Omega_0 t}d\Omega_0.$$

It is clear that the pseudophoton spectrum is limited in this case to the frequency interval $\Delta\omega+\Delta\Omega_0\sim 10^7-10^{10}$ Hz, corresponding to the width Γ of the majority of resonant and activation levels of the Mössbauer nuclei.

We now estimate $P(\omega)$. For the realistic values $N\sim 10^{15}$, $\delta\sim 1$, $T\sim 10^9$ sec and $S\sim 10^{-4}$ cm^2 we have $P(\omega=\Omega_0)\sim 10^{11}$ photons/sec-cm^2 Hz. Such a pseudophoton flux, which is perfectly analogous for superrelativistic electrons to ordinary photons, makes it possible, as follows, for example, from Ref. 11, to excite the threshold value of the inversion needed to satisfy the lasing condition. It must be noted that allowance for the possible channeling and periodic spatial compression of the beam in channeling[12-14] can only improve the parameters of the system, since it increases the beam density within the limits of the employed interval of impact distances $\rho_{\min}\leq\rho\leq\rho_{\max}$. At the same time for positrons, which are localized in the interstices in the case of channeling, the effect becomes much weaker.

The problem of linear acceleration (without affecting the structure) of an electron bunch with dimension larger than, for example, the wavelength of the accelerating field in a linear resonant (or other type) accelerator can be solved by various methods. In particular, it is possible to form out of the unmodulated weakly relativistic beam a periodic sequence of bunches, the interval between which corresponds to the period of the accelerating field T_B, and whose duration corresponds to the time interval of the capture of the electrons in the linear-acceleration, i.e., to the duration of the characteristic bunches shaped by the accelerator.[15] The role of such a shaper can be assumed, in particular, by the first few resonator sections of the accelerator, past which the modulator is located. The shaped modulated sequence of bunches with duration $T\gg T_B$ is at spatial phase synchronism with the accelerating traveling microwave, thus ensuring effectiveness of the subsequent acceleration.

To analyze the spectrum of the pseudophotons of such an aggregate, we introduce the electron density in the form

$$n(t)=\sum_{s=-M/2}^{M/2}\tilde{n}[1-\delta\cos(\Omega_0 t+s\Omega_0 T_B)]$$
$$\text{at}\quad sT_B\leqslant t\leqslant(s+k_3)T_B$$

with $n(t)=0$ outside the limits of these intervals. Here $k_3=\Delta\varphi/2\pi$ is the coefficient of the capture of the electrons in the linear-acceleration regime, corresponding to the phase interval $\Delta\varphi$ of the accelerating field,[15] and $M=T/T_B$ is the number of bunches in the beam.

For the gamma-band frequencies of interest to us, the duration of the passage of the bunch k_3T_B exceeds by many orders of magnitude the correlation time of the process $B(\tau)$. As a result, the calculation procedure remains the same as before, and the resultant expressions take the form

$$G(\omega)=\frac{2e^2(N/M)\delta}{k_3T_Bv^2S}\ln\frac{\rho_{max}}{\rho_{min}}\frac{\sin(M\Omega_0T_B/2)}{\sin(\Omega_0T_B/2)}\frac{\sin[(\omega\pm\Omega_0)T_0/2]}{(\omega\pm\Omega_0)T_0/2}$$
$$+\frac{2\pi e^2(N/M)^2\delta^2}{k_3T_Bv^2S^2}(\rho_{max}-\rho_{min})^2\frac{\sin^2(M\Omega_0T_B/2)}{\sin^2(\Omega_0T_B/2)}\frac{\sin[(\omega\pm\Omega_0)T_0/2]}{(\omega\pm\Omega_0)T_0/2}$$
$$\times\frac{\sin[(\omega\pm\Omega_0)T/2]}{(\omega\pm\Omega_0)T/2}.$$

It is seen that if the frequencies of the accelerating field is a multiple of that of the modulating field the spectral density of the pseudophotons of a discrete modulated sequence of electron bunches is the same as in the previously analyzed case of a continuously modulated beam. The availability of tunable laser sources of light makes such a possibility perfectly realistic.

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