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V.N. Baier and A.I. Milstein

ELECTROMAGNETIC-WAVE AMPLIFICATION BY AN ELECTRON BEAM IN ALIGNED SINGLE CRYSTALS

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ELECTROMAGNETIC-WAVE AMPLIFICATION BY AN ELECTRON BEAM IN ALIGNED SINGLE CRYSTALS

V.N.Baier and A.I.Milstein

Institute of Nuclear Physics, 630090 Novosibirsk, USSR

Abstract

A consistent analysis has been made of amplification of an electromagnetic wave during its interaction with electrons and positrons moving in planar channels of a single crystal. The main effects: dechanneling, unharmonism in transverse motion and radiation losses of energy, have been taken into account.

1. Introduction

As known, generation of coherent radiation using a beam of relativistic particles is realized in the devices called free-electron lasers (FELs) [1-5]. The FEL action is described in terms of the classical theory and in a wide range of parameters one can apply the so-called single-particle theory where an interaction of the changed particles in the beam is neglected [2,3]. Because of an interaction of the electron beam moving in a periodic magnetic structure* (undulator) with the radiation field the beam particles are spatially bunched. The important feature of the process is the phase stability which in turn ensures the bunching stability, thereby creating the conditions for generation of coherent beam radiation. In the region of low gains G ($G \ll 1$) an analysis is considerably simplified since the phase shift of the radiation field may be neglected [2,3], which becomes necessary at $G \gtrsim 1$ [4]. Sometimes the region $G \ll 1$ is described within the framework of quantum theory.

So far a lot FELs have been created, which operate in the range from centimeter waves to visible light (see, e.g. the special issue [5]). The wavelength of emitted photons in a FEL λ is

$$\lambda \sim \lambda_0 / 2\gamma^2 \qquad (1.1)$$

where λ_0 is the lattice wavelength and $\gamma = \frac{\varepsilon}{m}$ is the Lorentz factor**. Since the minimum wavelength in a magnetic undulator is about one centimeter, a transition to the short-

* For FEL action a periodicity in the beam motion is really needed.

** The system of units $c = 1$ is used.

wavelength region is possible only due to an increase of the particle energy. The difficulty in realizing this way is connected with the following dependence of the gain $G \propto \gamma^{-3}$.

Of great interest is a search for the objects where the particles are in periodic motion with a small wavelength. One of such possibilities takes place at charged particles channeling in single crystals (see, e.g. [6]). The wavelength of a particle at channeling may be estimated as follows:

$$\lambda_0 \approx \frac{2\pi\gamma d}{\sqrt{S}}, \quad S = \frac{2U_0\varepsilon}{m^2} \qquad (1.2)$$

where d is the characteristic size of the channel and U_0 is the depth of the potential well of the channel. For instance, at channeling in the (110) plane in Si (U_0 = 25 eV, $d \approx$ 0.1 nm) we have $\lambda_0 \sim 10^4 d \sim 1\mu m$, at ε = 100 MeV.

In this paper we shall deal with a consistent analysis of an interaction between the linearly polarized plane electromagnetic wave propagating along the single crystal's planes and the beam of ultrarelativistic particles moving in the same planar channels, with the specific features of particle motion in a single crystal taken into account. The accepted approach is similar to that applied in Ref.2. The solution of this problem gives the basic parameter, the gain G.

The induced radiation when moving the particles in the single crystal's channels have been considered in [7-10]. In Ref.[7-9] for the gain use has been made of an expression derived in an analysis of a wholly different physical situation: propagation of signals in medium with inverse population. Naturally, in this analysis no allowance has been made for the dynamics of an interaction of the beam particles with the electromagnetic wave and, therefore, such features

as the shape of a spectral line and the functional time dependence have not been reproduced. In Ref.[10] this dynamics has been already taken into consideration. The results obtained in [10] will be compared with ours in the Conclusion.

II. Problem Formulation; General Solution

Let us analyse the induced radiation of ultrarelativistic particles moving in the planar channels of a single crystal. For the case of motion in axial channels the problem is much more complicated from a technical point of view and at the same time does not be of great interest, as will be seen from the results obtained. Let the particle beam is incident on the crystal in a such way that the momenta of the electrons lie near the crystallographic planes and interacts with the linearly polarized plane electromagnetic wave propagating along the same planes; its electric field is perpendicular to these planes. Our set of coordinates is: the x^1-axis is perpendicular to the planes forming a channel, and the x^3-axis is chosen parallel to these planes such that the velocity along the x^2-axis is v^2 = 0. Just as in [2], it is convenient to use the equations of motion in covariant form ($\frac{du^\mu}{ds} = \frac{e}{m} F^{\mu\nu} u_\nu$) and then for our problem we obtain

$$\frac{d(u^0+V)}{ds} = \Omega u^1 \sin\varphi$$

$$\frac{du^1}{ds} = -\frac{dV}{dx^1} u^0 + \Omega (u^0 - u^3) \sin\varphi \qquad (2.1)$$

$$\frac{du^3}{ds} = \Omega u^1 \sin\varphi$$

$$\frac{d\Omega}{ds} = -\omega_p^2 \, \overline{u^1 \sin\varphi}$$

where u^μ is the component of the four-velocity, $V(x^1)=\frac{e\,\mathcal{U}(x^1)}{m}$, $\mathcal{U}(x^1)$ is the inter-planar potential, $\Omega=\frac{eE_\omega}{m}$, E_ω is strength of the the electric wave field given by the vector-potential $A^i_\omega = -\frac{E_\omega}{\nu}\cos\varphi$, $\varphi=\nu(t-x^3)+\psi$, $\omega_p^2=\frac{4\pi ne^2}{m}$ where n is the particle density in the beam, and $m\,(e)$ is the mass (charge) of an electron. The latter equation (2.1) follows from the energy conservation law and the averaging sign means the averaging over the wave phase: $\overline{a}\equiv\frac{1}{2\pi}\int_0^{2\pi}d\psi\, a(\psi)$. The set of equations (2.1) is valid only in the case when the gain is $G\ll 1$ (for the approach at $G\gtrsim 1$ see Ref.[4]) and may be solved here in an adiabatic approximation, i.e. the first three equations in (2.1) can be solved assuming that Ω = const.

The integral of motion follows from the first and the second equations in (2.1):

$$u^0-u^3+V(x^1)=const \tag{2.2}$$

We shall consider the particle motion in a channel, i.e. $(u^1)^2<2V_{max}\gamma$ where $V_{max}=\frac{e\,\mathcal{U}_0}{m}$, under the assumption that the transverse momentum is small: $|u^1|\ll|u^3|$. The problem of particle interaction with the wave field will be solved in an approximation linear with respect to the field strength of the wave. The inter-planar potential is convenient to represent as follows:

$$V(x^1)=V_0\,f(y),\quad y=x^1/d \tag{2.3}$$

where $2d$ is the inter-planar distance, $V_0/2=V_{max}=\frac{e\,\mathcal{U}_0}{m}$, and $\mathcal{U}_0$ is the depth of the potential well. With the wave absent, we will write the solution of eqs.(2.1)

$$x^1(s)=X(s),\quad u^3(s)=\gamma,\quad \varphi(s)=\phi(s)$$

With the wave swiched on, the perturbation of this motion occurs:

$$x^1(s)=X+\tilde{\xi},\quad u^3=\gamma+\tilde{\eta},\quad \varphi=\phi+\tilde{\chi} \tag{2.4}$$

All the increments, $\tilde{\xi}$, $\tilde{\eta}$ and $\tilde{\chi}$, are proportional to Ω. The initial conditions being the following:

$$\tilde{\xi}(0)=\frac{d\tilde{\xi}(0)}{ds}=\tilde{\eta}(0)=\tilde{\chi}(0)=0$$

$$X(0)=X_0,\quad \frac{dX(0)}{ds}=u^1(0),\quad \phi(0)=\psi \tag{2.5}$$

Let us introduce the quantities the using of which enables the expressions to be simplified considerably:

$$\Omega_0^2=\frac{\gamma V_0}{d^2},\quad \rho=(\Omega_0 d)^2=\gamma V_0,\quad \delta=\frac{\nu}{2\Omega_0\gamma}. \tag{2.6}$$

It is convenient to change over to the dimensionless variables:

$$\tau=\Omega_0 s,\quad \tilde{\xi}=\xi\frac{\Omega d^2}{\gamma},\quad \tilde{\eta}=\Omega\eta d,$$

$$\tilde{\chi}=\chi\, 2\rho\Omega d\delta/\gamma \tag{2.7}$$

Substituting eqs.(2.3), (2.4), (2.6) and (2.7) into the set (2.1), performing appropriate expansions and collecting the zero- and first-order terms with respect to Ω, we have, with eq.(2.2) taken into account, the final system of equations for dimensionless variables:

$$(1)\ \ddot{y}+f'(y)=0;\quad (2)\ \dot{\phi}=(1+\rho\dot{y}^2)\delta;\quad (3)\ \dot{\eta}=\dot{y}\sin\phi;$$

$$(4)\ \dot{\chi}=-\xi f'(y);\quad (5)\ \ddot{\xi}=-f''(y)\xi-f'\eta+\frac{\dot{\phi}\sin\phi}{2\rho\delta}; \tag{2.8}$$

$$(6)\ \frac{dG}{d\tau}=-\frac{2\omega_p^2 d^2}{\gamma}\left(\dot{\xi}\sin\phi+\dot{y}\,2\rho\chi\delta\cos\phi\right)$$

where the notations are: $\dot{y}\equiv\frac{dy}{d\tau}$ and $f'(y)\equiv\frac{df}{dy}$; we have introduced the gain

$$G\equiv\frac{\Omega^2(\tau)-\Omega^2(0)}{\Omega^2(0)}\simeq\frac{2(\Omega(\tau)-\Omega(0))}{\Omega(0)} \tag{2.9}$$

It is remarkable that despite the apparent complexity, the set (2.8) is solvable in quadrature for any inter-planar potential $\mathcal{U}(x^\perp)$. Let us start to solve the set (2.8). Equation (1) is solved in quadrature and we denote its solution as $y(\tau)$. This solution is a periodic function of period T. From eq.(2) we have

$$\phi(\tau)=\psi+\delta\int_0^\tau(1+\rho\dot{y}^2(\tau))d\tau=\psi+\tau\delta(1+\rho\langle\dot{y}^2\rangle)+ \tag{2.10}$$
$$+\int_0^\tau d\tau\,\rho\delta(\dot{y}^2-\langle\dot{y}^2\rangle)$$

where $\langle\dot{y}^2\rangle$ is the squared velocity averaged over the motion period. Eq.(3) in the set (2.8) gives

$$\eta=\int_0^\tau\dot{y}(\tau')\sin\phi(\tau')d\tau' \tag{2.11}$$

Eq.(5) in the set (2.8) is the most complicated one. To solve it we will use an artificial trick : we differentiate eq.(1) over τ: $\dddot{y}=-f''(y)\dot{y}$, multiply both parts of eq.(5) by $\dot{y}$ and express $f''(y)\dot{y}$ via $\dddot{y}$. Using eqs.(2) and (3) and taking integral over τ of both parts, we have

$$\dot{y}\dot{\xi}-\ddot{y}\xi=\frac{\ell}{2}\left(\frac{1}{\rho}+\dot{y}^2\right) \tag{2.12}$$

It is seen from (2.12) that the expression for $\xi(\tau)$ may be found in quadrature. It is useful to introduce a periodic function $z(\tau)$

$$z(\tau)=\dot{y}(\tau)\int_0^{\tau^*}d\tau'\frac{1}{\dot{y}^2(\tau')} \tag{2.13}$$

Here $\tau^*=\tau-\kappa\frac{T}{2}$ where κ is the number of the turning points on the trajectory of a particle in its transverse motion for the time τ. Using this function the solution of eq. (2.12) may be presented in the form, taking eq.(2.11) into account:

$$\xi(\tau)=\frac{1}{2\rho}\int_0^\tau ds\,\sin\phi(s)\,W(s,\tau) \tag{2.14}$$

where

$$W(s,\tau)=z(\tau)\dot{y}(s)-z(s)\dot{y}(\tau)+\rho(\tau-s)\dot{y}(s)\dot{y}(\tau) \tag{2.15}$$

The solution of eq.(4) in the system (2.8) is

$$X(\tau)=\int_0^\tau ds\,\xi(s)\,\ddot{y}(s) \tag{2.16}$$

Substituting our solutions of eqs.(1)-(5) into eq.(6) and making an averaging over ψ we have for the gain

$$\frac{dG}{d\tau}=-G_0\left\{\int_0^\tau\frac{\partial W(s,\tau)}{\partial\tau}\cos(\phi(\tau)-\phi(s))ds+\right. \tag{2.17}$$
$$\left.+2\delta\rho\dot{y}(\tau)\int_0^\tau ds\int_0^s ds'\ddot{y}(s)W(s',s)\sin(\phi(s')-\phi(s))\right\}$$

where

$$G_0=\frac{\omega_p^2 d^2}{2\gamma\rho} \tag{2.18}$$

The expression (2.17) solves, in general form, the problem on an amplification of the plane electromagnetic wave with linear polarization as a result of its interaction with the charged particles moving in the single crystal's planar channels.

III. The Gain G for Non-Relativistic Transverse Motion

In the region of comparatively low energies (tens and hundreds of MeV) the parameter $\rho=\gamma V_0\ll 1$ (in the (110) plane, at room temperature we have $V_0\simeq 10^{-4}$ for Si and diamond, $V_0\approx 2\cdot10^{-4}$ for Ge and $V_0\approx 5\cdot10^{-4}$ for W). Then the transverse motion of the particles in the channel is non-relativistic: $(u^\perp)^2<\rho\ll 1$. As we shall see below, it is the energy range which is of especial interest from the point of view of radiation generation. By virtue of the said, in what follows we confine ourselves to an analysis of the case when $\rho\ll 1$. But the "time" τ (determined by the number of particle oscillations in the channel) is assumed to be large so that $\rho\tau \gtrsim 1$.

In the accepted approximation, in formula (2.17) one can put, according to formula (2.10),

$$\phi(\tau) - \phi(s) = (\tau - s)\delta \qquad (3.1)$$

Since we analyse the situation when the particle oscillate many times in the channel, the integral (2.17) is convenient to reduce to that over period and summation over the number of oscillations, regarding for the fact that the functions $\dot{y}(\tau)$ and $z(\tau)$, entering into $W(s,\tau)$, are periodic. There is no difficulty in performeing this summation using an elementary formula

$$\sum_{m=0}^{n-1} \cos(\varphi + \alpha m) = \frac{\sin \frac{n\alpha}{2}}{\sin \frac{\alpha}{2}} \cos\left[\varphi + \frac{\alpha}{2}(n-1)\right] \qquad (3.2)$$

The integrals entering into eq.(2.17) are "large" only if the resonance condition

$$\frac{\delta}{2} = \frac{\kappa\pi}{T} + \frac{\zeta}{\tau}, \; \kappa = 1,2,3\ldots \quad , \frac{\zeta}{\tau} \ll 1 \qquad (3.3)$$

is satisfied. After simple calculations using eqs. (3.1)-(3.3), for the gain for the "time" τ we obtain

$$G(\tau) = -G_0 \tau^2 \delta \left[\mathrm{Jm}(z_\kappa^* \dot{y}_\kappa) - \frac{\rho\tau}{2} |\dot{y}_\kappa|^2 \frac{d}{d\zeta} \right] R(\zeta) \qquad (3.4)$$

where $R(\zeta) = \sin^2\zeta / \zeta^2$ determines the shape of the spectral line and z_κ and $\dot{y}_\kappa$ are the Fourier-harmonics:

$$z_\kappa = \frac{1}{T}\int_0^T ds\, z(s) \exp\left(-i\frac{2\pi\kappa s}{T}\right), \; \dot{y}_\kappa = \frac{1}{T}\int_0^T ds\, \dot{y}(s) \exp\left(-\frac{i 2\pi\kappa s}{T}\right) \qquad (3.5)$$

The quantity G_0 is given by formula (2.18).

According to formulas (2.6) and (3.3), for the resonance frequency we have

$$\nu_{res} = 2\gamma\Omega_0 \frac{2\kappa\pi}{T} \qquad (3.6)$$

In addition, in formula (3.4) an averaging over the coordinate of the point of incidence of the particle $(y(0) = x'(0)/d)$ should

be performed. Changing $y(0)$ we have different values of the transverse energy (see eq.(2.8, (1)):

$$\varepsilon_\perp = \frac{\dot{y}^2}{2} + f(y) = \frac{\dot{y}^2(0)}{2} + f(y(0)) \qquad (3.7)$$

and, generally speaking, different periods of motion T.

For averaging it is necessary to take an integral over $y(0)$ in the area where $\varepsilon_\perp < 1/2$; this corresponds to the condition of presence of the particle in the channels (see eq. (2.3)):

$$\bar{g} = \frac{1}{2}\int_{-a}^{a} dy(0)\, g(y(0)), \quad \frac{\dot{y}^2(0)}{2} + f(a) = \frac{1}{2} \qquad (3.8)$$

In the particular case of an oscillator potential $f(y) = y^2/2$, we find $z(\tau) = \sin\tau / \dot{y}(0)$. In this case there is only one harmonic $\kappa = 1$, $T = 2\pi$ and for the resonance frequency (3.6) we obtain $\nu_{res} = 2\Omega_0\gamma$. After the simple calculation we find the gain for an oscillator after the averaging over the coordinates of the point of incidence:

$$G^{os} = -G_0 \frac{\tau^2}{4}\left(1 - \frac{\rho\tau}{6}\sqrt{1 - \dot{y}^2(0)}\left(1 + \frac{2}{3}\dot{y}^2(0)\right)\frac{d}{d\zeta}\right) R(\zeta) \qquad (3.9)$$

Note that only in the case of an oscillator potential the averaging over the coordinates of the point of incidence does not affect the spectral curve.

IV. Conclusion

We proceed now to the discussion of the results obtained. The first term in the expression for the gain (3.4) always describes the absorption. It is easy to see that this expression is an radiation absorption coefficient by the system of relativistic oscillators. The structure of the second term is completely similar to the gain G for a FEL and can be compared with this G if one takes into account the similarity of particle trajectories (see Ref. [2]), namely: $(\nu_0)_{FEL} \leftrightarrow \frac{\Omega_0}{\gamma}$,

$\left(\frac{\Omega_0}{V_0}\right)_{FEL} \leftrightarrow \Omega_0 d = \sqrt{g}$. Note that under typical FEL conditions $(\Omega_0/V_0)_{FEL} \sim 1$, then the second term in the expression for G dominates and the first term is usually omitted. However, in the energy range under discussion the first term may be roughly equal to the second and it must be taken into account. It is seen from the expressions (3.4) and (3.8) that the amplification of the wave is possible actually only if $\rho\tau \gtrsim 1$, otherwise the amplification can occur only in a very narrow band of the spectral curve $|\pi - \zeta| \ll 1$ where an additional suppression of G takes place and where it is readily smeared by the spread over the periods of motion. One can amplify the wave only if the condition $\frac{d}{d\zeta} R(\zeta) > 0$ is satisfied and the amplification must exceed the absorption.

There are a number of conditions which limit considerably the possibilities of selecting the parameters for this problem. Let us dwell upon these conditions.

We start with an analysis of the dechanneling effects. Within the energy range under analysis the basic effect is multiple scattering of particles. Under the action of this mechanism the particles acquire an additional transverse momentum. And, as a result, they leave the channel after passing the length ℓ_d. If the scattering is assumed to occur in the same way as in amorphous matter, then the length of ℓ_d is (Ref. [11]):

$$\ell_d \simeq \frac{\alpha}{2\pi} \, g \, L_{rad} \tag{4.1}$$

where α = 1/137 and L_{rad} is the radiation length in an appropriate amorphous matter. This estimation is, strictly speaking, overestimated for electrons and underestimated for positrons ($\ell_d^{e^-} < \ell_d^{e^+}$). If allowance is made for the fact that during the interaction with radiation the basic role is played by the particles whose oscillation amplitude is comparable with the sizes of the channel, then the above estimation seems to be rather satisfactory. The parameter $\rho\tau_d$, corresponding to ℓ_d, is

$$\rho\tau_d \simeq \frac{\alpha}{2\pi} \frac{V_0^{5/2} \gamma^{3/2}}{d} L_{rad} \tag{4.2}$$

It is obvious that $\tau < \tau_d$. To reach the region where $\rho\tau_d \sim 1$, it is necessary that the energy be fairly high. The estimates show that for the usually used crystals (diamond, Si Ge, the (110) plane) we have $\rho\tau_d \sim 1$ at ε = 200 MeV. For this energy, for the (110) plane $\rho \simeq 0.04$ in diamond and Si and $\rho \simeq 0.08$ in Ge. Remind that at the indicated energy one can use the classical theory with large margin.

Another effect limiting τ is the dependence of the period T on the transverse energy. Note once more that the quantity ζ entering into the expression for the spectral curve in eq. (3.4) is $\zeta = \tau\left(\frac{\rho}{2} - \frac{\kappa\pi}{\tau}\right)$ (see eq. (3.3)). In order that the amplification effect occur, the increment of the quantity ζ must be $\Delta\zeta \ll \pi$, at averaging over the point of incidence $y(0)$; the latter can be represented as the averaging over the period T. This imposes the limitations on a possible spread with respect to the periods of motion:

$$\frac{\Delta T}{T} \ll \frac{T}{\tau} \sim \frac{2\pi}{\tau} \sim \frac{1}{N} \tag{4.3}$$

(N - is the number of oscillations). We discuss the region $\rho\tau \gtrsim 1$, $\rho \ll 1$ and $\tau \gg 1$. The effect is possible in it only for the potentials close to the oscillator ones; in addition, as follows from (4.3) the quantity τ cannot be very large. The inequality (4.3) excludes, in practice, the utilization of planar channeling, of electrons because $\Delta T \sim T$ for them. Since the criterion (4.3) follows, in fact, from the coherency condition, it should be satisfied also for induced radiation at axial channeling. This excludes actually the use of the axial channeling of electrons and positrons since at this channeling $\Delta T \sim T$. It is worth noting that in Ref. [10] the induced radiation has been considered at axial channeling under the assumption that the particles move along the circular trajectories of given radius. Their weight in the set of trajectories has, however, the set of measure zero so that this analysis is not suitable for the analysis of a realistic situation. For the planar channeling of positrons the potentials are close to the oscillator ones. Unfortunately, even in this case, under the most favourable conditions $\Delta T/T \approx 0.05$, then $\tau \lesssim 100$ (for the potential $U(x) = V_0^* [ch(\frac{x}{a_s}) - 1]$ see Ref. [12]).

We would like to dwell upon one more restriction due to the energy spread in the incident beam and to the energy losses during its passage through a single crystal. In the energy range under consideration the radiation losses dominate. Then, having substituted the quantities entering into ζ (3.3), we find

$$\frac{\Delta \gamma}{\gamma} \ll \frac{\mathfrak{F}}{\tau} \tag{4.4}$$

Whence, we have

$$\tau \ll \left(\mathfrak{F} L_{rad} \frac{V_0^{1/2}}{d\, \gamma^{1/2}} \right)^{1/2} \tag{4.5}$$

For Si, eq. (4.5) gives $\tau \ll 10^3$, i.e. the restriction is weaker than that discussed above. It is seen from eq. (4.4) that the requirement for the monochromaticity of the incident beam is rather weak. Note that the width of the spectral curve follows also from the analysis of expression for G (eq. (3.3), cf. Ref. 2-4) and we have $\Delta \nu/\nu \sim 1/\tau$.

The analysis made above enables one to estimate the magnitudes of the amplification effects. After the substitution of $n = 10^{14}$ cm^{-3}, $\tau = 100$, $d = 10^{-8}$ cm and $\gamma = 400$ into formula (3.8) we find that $G \approx 10^{-12}$. Thus, at even so high particle density of the electron beam the wave gain turns out to be negligibly low. The extremely high requirements for the particle density in the beam are, to a known degree, evident and are due to a microscopically small amplitude of particle oscillations in the channel. The rigid limitations on the magnitude of τ are much less evident, meanwhile this does not allow one to compensate the smallness of the quantity G_0 by a high τ (we would have $G \sim 1$ at $\tau \sim 10^6$).

In the present paper a consistent analysis has been made of the amplification of an electromagnetic wave by electrons and positrons moving in the planar channels of a single crystal for an arbitrary shape of the inter-planar potential. In our analysis the main effects occurring during the passage of fast charged particles through aligned single crystals have been taken into consideration. The conclusion we have drawn is pessimistic since there is a large gap between the densities of the particle beams with an energy of hundreds of MeV, generated by modern accelerators, and the densities needed for the FEL action on electrons and positrons moving in the channels of a single crystal. The possibility of increasing the gain at the expense of a growth of the number of oscillations in the channel (τ), has proved to be rigidly restricted by the crystalline effects.

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В.Н.Байер, А.И.Мильштейн

УСИЛЕНИЕ ЭЛЕКТРОМАГНИТНОЙ ВОЛНЫ ПУЧКОМ ЭЛЕКТРОНОВ В ОРИЕНТИРОВАННЫХ МОНОКРИСТАЛЛАХ

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