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Electromagnetic Emission of Mobile Dislocation Segments in an Ionic Crystal

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Abstract—The problem of electromagnetic emission of an edge dislocation segment moving in an ionic lattice with a NaCl-type structure is considered. The proposed mechanism of electromagnetic emission is associated with the appearance of macroscopic alternating polarization currents along the extraplane edge of the edge dislocation in the course of its motion between adjacent valleys of the Peierls relief. The relationships for electromagnetic radiation fields of an arbitrarily moving segment are derived, and the problem of electromagnetic emission of a segment that executes harmonic oscillations in the field of an external quasi-stationary elastic wave with a frequency $\Omega \ll c/l$ (where l is the segment length and c is the velocity of sound) is treated in detail. The power of the emitted electromagnetic signal and the "acoustoelectromagnetic transformation" coefficient (the ratio between the electromagnetic radiation power and the mechanical power required for setting the segment in motion) are determined. *© 2001 MAIK "Nauka/Interperiodica".*

1. INTRODUCTION

Electromagnetic effects that accompany the motion of defects in crystals have been the subject of extensive investigations in modern solid-state physics. Over the last few years, considerable advances have been made in this area of science. First and foremost, these are experimental observations [1–3] and theoretical interpretation [3, 4] of the magnetoplastic effect in ionic crystals and metals and also detailed experimental investigations into the electromagnetic emission of dislocations and cracks [5, 6]. It should be emphasized that these phenomena are essentially dynamic in character and, hence, are of particular importance in the understanding of the nature of plastic deformation in solids. Electromagnetic phenomena in deformed solids have long been studied both from the standpoint of basic research [7] and in relation to applied problems arising, for example, in geophysics [8] and fracture mechanics of structural materials [9, 10].

An interesting dynamic effect in strained crystals is the emission of electromagnetic waves during dislocation motion and the nucleation and growth of cracks. In principle, it is clear that, for example, dislocation motion brings about disturbances in both the crystal lattice and the electronic subsystem of the crystal. As is known, the former disturbance leads to the generation of elastic waves in the sample and the latter disturbance gives rise to electromagnetic emission whose character is determined by the specific properties of dislocations and the medium in which the electromagnetic wave propagates. It seems likely that ionic crystals are the most convenient objects for the investigation of electromagnetic emission of dislocations, because they are dielectrics in which absorption of an electromagnetic signal is virtually absent down to the IR frequency range. On the other hand, dislocations of the majority of the types in these crystals have a charged core [11, 12], so that their motion is naturally attended by effective currents. Therefore, in this case, it is rather simple to construct adequate physical models in which a dislocation is interpreted as a source of electromagnetic waves in the crystal.

Kosevich and Margvelashvili [13] studied the emission of electromagnetic waves by mobile dislocations in ionic crystals and explained this phenomenon in terms of a mechanism based on electroelastic effects in deformed lattices consisting of oppositely charged ions. In our earlier work [14], we proposed an alternative mechanism of electromagnetic emission of edge dislocations in ionic crystals, according to which the electromagnetic emission is due to the occurrence of macroscopic polarization currents along the line of a moving rectilinear dislocation. It turned out that the emission intensity in the latter case is five orders of magnitude higher than that in the former case. Therefore, it can be expected that the mechanism proposed in [14] is responsible for experimentally observable effects.

In [13, 14], the electromagnetic emission was considered for rectilinear edge dislocations. It is clear that, in a real crystal, a dislocation moving through a stopper network is actually a set of oscillating segments. In this respect, it was of interest to solve the problem of electromagnetic emission of a curvilinear dislocation whose configuration is an arbitrary function of time. The aim of the present work was to analyze the formulated problem and to elucidate the contribution from this emission mechanism to the total electromagnetic emission.

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2. FORMULATION OF THE PROBLEM

Let us consider an edge dislocation in a cubic cell of the NaCl type with a glide plane coinciding with one of the {110} planes. The Cartesian coordinate system is chosen in such a way that the dislocation glides in the plane $y = 0$ and its line in the absence of external disturbances coincides with the z axis. We assume that the dislocation line in the course of motion is so curved that its configuration in the laboratory coordinate system at any instant of time can be described by the function $x = x_0(z, t)$, which is single-valued with respect to the z coordinate. This means that the dislocation during the motion cannot generate loops of the Frank–Read source type. Therefore, the charge density $\rho(\mathbf{r}, t)$ at the extraplane edge of the curvilinear dislocation (in the core) can be represented as

$$\rho(\mathbf{r}, t) = e^*\delta(y)F(x)\delta(x - x_0(z, t)) \times \sum_{m=-\infty}^{\infty}\{\delta(z - 2ma) - \delta[z - (2m+1)a]\}. \quad (1)$$

Here, e^* is the effective charge of a node at the extraplane edge, $2a$ is the distance between likely charged ions along the z axis [positive ions are located at the sites $z = 2ma$ and negative ions occupy the sites $z = (2m + 1)a$], and the period of the function $F(x)$ [$|F(x)| \le 1$] in the direction of the dislocation motion is equal to $2b$ (where b is the distance between adjacent minima of a Peierls relief in the direction of the Ox axis). The multiplier $F(x)$ allows for the effective "recharging" of the node located at the dislocation line when this node is displaced to the adjacent valley of the Peierls relief [14]. It is worth noting that, unlike the rectilinear dislocation [14], the sign of the effective charge at the node through which the line of a curvilinear dislocation segment passes [see formula (1)] depends on two coordinates (x_0 and z).

It is evident that the evolution of the charge in the core of a curvilinear dislocation should satisfy the continuity equation

$$\frac{\partial\rho(\mathbf{r}, t)}{\partial t} + \mathrm{div}\,\mathbf{j}(\mathbf{r}, t) = 0, \quad (2)$$

where $\mathbf{j}$ is the effective current density in the dislocation core. By differentiating relationship (1) with respect to time and rearranging it, we have

$$\frac{\partial\rho}{\partial t} = -ae^*\delta(y)\left\{\frac{\partial}{\partial z}\left[F(x)\frac{\partial}{\partial x}\delta(x - x_0(z, t))V(z, t) \times \sum_{m=-\infty}^{\infty}\delta(z - 2ma)\right] - F(x)\frac{\partial}{\partial z} \times\left[V(z, t)\frac{\partial}{\partial x}\delta(x - x_0(z, t))\right]\sum_{m=-\infty}^{\infty}\delta(z - 2ma)\right\}. \quad (3)$$

Here, we used the approximate expression

$$\sum_{m=-\infty}^{\infty}[\delta(z - 2ma) - \delta[z - (2m+1)a]] \simeq a\frac{\partial}{\partial z}\sum_{m=-\infty}^{\infty}\delta(z - 2ma),$$

which implies that we are interested only in microscopic quantities that vary slowly at distances of the order of interatomic distances. This corresponds to the usual procedure of microcurrent averaging accepted in electrodynamics of continua [15]. Comparison of relationships (3) and (2) demonstrates that the term under the sign of the total derivative with respect to z in formula (3) is the current density component j_z, that is,

$$j_z(\mathbf{r}, t) = \frac{e^*}{2}\delta(y)F(x)V(z, t)\frac{\partial}{\partial x}\delta(x - x_0(z, t)), \quad (4)$$

where $V(z, t) = \partial x_0(z, t)/\partial t$ is the distribution of velocities of motion of dislocation segment points. The other part of expression (3) is naturally attributed to $\partial j_x/\partial x$. By integrating it over x, we find

$$j_x(\mathbf{r}, t) = -\frac{e^*}{2}\delta(y) \times \int_{-\infty}^{x_0(z,t)} dx' F(x')\frac{\partial}{\partial z}\left[V(z, t)\frac{\partial}{\partial x'}\delta(x' - x_0(z, t))\right] \quad (5)$$
$$= -\frac{e^*}{2}\delta(y)\frac{\partial}{\partial z}\left[V(z, t)\frac{\partial}{\partial x_0}F(x_0(z, t))\Theta(x_0(z, t) - x)\right].$$

When writing relationships (4) and (5), we performed averaging of the corresponding terms in expression (3) over the z coordinate.

Now, it is easy to calculate the components (necessary for further analysis) of the dipole moment vector $\mathbf{d}(t)$ of the dislocation segment. This vector is related to the current density $\mathbf{j}$ by the usual expression [16]

$$\frac{\partial\mathbf{d}}{\partial t} = \int\mathbf{j}(\mathbf{r}, t)dV. \quad (6)$$

With due regard for relationships (4) and (6), the dipole moment component dz is given by

$$\frac{\partial d_z}{\partial t} = \frac{e^*}{2}\int dz' V(z', t)\frac{\partial}{\partial x_0}F(x_0(z', t)) = \frac{e^*}{2}\frac{\partial}{\partial t}\int F(x_0(z', t))dz'. \quad (7)$$

For d_x, after the integration of expression (5) over y and z, we obtain

$$d_x = \frac{e^*}{2}\int_{-\infty}^{+\infty} dx\Big[V(z_1,t)\frac{\partial}{\partial x_0}F(x_0(z_1,t))\Theta(x_0(z_1,t)-x) \\ - V(z_2,t)\frac{\partial}{\partial x_0}F(x_0(z_2,t))\Theta(x_0(z_2,t)-x)\Big], \quad (8)$$

where $z = z_{1,2}$ are the coordinates of the dislocation ends (formally, we can set $|z_{1,2}| \longrightarrow \infty$). For the calculation of integral (8), we assume that the motion of the curvilinear dislocation can be considered a superposition of displacements $u(z, t)$ of dislocation points in a certain accompanying coordinate system and the displacement $X(t)$ of this system in space; i.e., $x_0(z, t) = X(t) + u(z, t)$. To state this differently, the dislocation motion can be treated as a superposition of the motion of a rectilinear (on the average) dislocation and the dislocation point oscillations $u(z, t)$ about the $oX(t)$ axis of the accompanying coordinate system. The oX axis of this system can be chosen, for example, from the condition

$$\int_{z_1}^{z_2} u(z,t)dz = 0.$$

Here, we will not discuss the question as to whether this definition of the accompanying system is unique. For our purposes, it is sufficient that such a system exists in principle. Since integrand (8) involves the δ function, the range of integration in this integrand is actually limited from above by $x = X(t) + u(z, t)$. In this case, the integral over x from $-\infty$ to $X(t)$ becomes zero under the assumption that $V(z_1, t) = V(z_2, t) = dX/dt$ at $|z_{1,2}| \longrightarrow \infty$ and the dislocation ends are simultaneously located in valleys of the Peierls relief; i.e., $F[x_0(z_1, t)] = F[x_0(z_2, t)]$. Then,

$$d_x = -\frac{e^*}{2}u(z,t)[V(z_1,t) - V(z_2,t)]\frac{\partial}{\partial x_0}F(x_0(z_1,t)). \quad (9)$$

Finally, when the segment ends are fixed at stoppers, we obtain $V(z_1, t) = V(z_2, t) = 0$ and the dipole moment component d_x vanishes. It is this case that will be considered in the following discussion.

3. ELECTROMAGNETIC EMISSION OF A DISLOCATION SEGMENT

In this section, we will derive the relationships that describe the electromagnetic fields induced in the crystal by a single dislocation segment, which has the initial length l and is fixed at the points $z = \pm l/2$ on the Oz axis. The dislocation, as before, glides in the plane $y = 0$. As can be seen from formula (5), the current density component j_x for the segment with fixed ends vanishes, so that the evolution of the radiation fields is governed only by the dipole moment component $d_z(t)$.

It is convenient to represent the radiation fields of the segment in the spherical coordinates r, θ, and φ, where the azimuthal angle θ is measured from the Oz axis of the Cartesian coordinate system. The formulas for the strengths $\mathbf{E}(\mathbf{r}, t)$ and $\mathbf{H}(\mathbf{r}, t)$ of electric and magnetic fields of an elementary dipole in the wave zone are given in [17]. The spectral components (Fourier transforms with respect to time) $\mathbf{E}^\omega(r, \theta, \varphi) = (0, E_\theta^\omega(r, \theta, \varphi), 0)$ and $\mathbf{H}^\omega(r, \theta, \varphi) = (0, 0, H_\varphi^\omega(r, \theta, \varphi))$ of the radiation fields have the form

$$E_\theta^\omega(r,\theta,\varphi) = H_\varphi^\omega(r,\theta,\varphi) = \frac{\omega^2}{c^2 r}d_z^\omega \sin\theta \exp\left(i\frac{\omega r}{c}\right). \quad (10)$$

The E_θ^ω, H_φ^ω, and d_z^ω spectral components introduced into expression (10) are determined in the usual manner. For example,

$$d_z^\omega = \int_{-\infty}^{+\infty} dt d_z(t)\exp(-i\omega t)$$

(the spectral components of all the other functions of time are determined in a similar way).

Thus, the problem of calculating the spectral components of the dipole moment reduces to calculation of integral (7) under the most general assumptions about the form of the $u(z, t)$ function that determines the displacements of dislocation segment points. It is clear that the displacements u far from the fixed points can be sufficiently large ($\sim L$). In any case, it is evident that, for macroscopic segments with $l > 10^{-6}$ cm, the inequality $|u(z, t)| \gg b$ is true virtually for all segment points (except for small neighborhoods of fixed points). This allows us to evaluate integral (7) by using the large parameter $|u(z, t)|/b \gg 1$. With the aim of simplifying further calculations, we assume, as was earlier done in [13], that $F(x) = -\cos(\pi x/b)$ and write d_z as

$$d_z = \frac{e^*}{2}\int_{-l/2}^{l/2}\cos\left(\frac{\pi}{b}u(z,t)dx\right). \quad (11)$$

Integral (11) includes the large parameter $u(z, t)/b$ in the argument of the cosine and can be estimated within the stationary phase approximation [18]. As a result, we obtain

$$d_z(t) = -e^*\left(\frac{b}{2l}\right)^{1/2} \\ \times \sum_\alpha \frac{1}{\sqrt{u''_{zz}(z_\alpha,t)}}\cos\left(\frac{\pi}{b}u(z_\alpha,t) + \frac{\pi}{4}\right), \quad (12)$$

where z_α is a set of stationary points determined by the condition $u'_z = 0$ [hereafter, the prime designates the derivative of the $u(z, t)$ function and the subscript indicates the variable with respect to which the derivative is taken]. Therefore, horizontal ($u' = 0$) and weakly curved ($u'' \longrightarrow 0$) portions of the oscillating segment make the main contribution to the emission.

In order to obtain a more concrete result, we calculate the segment motion in the framework of the string dislocation model [19]. In this approximation, the equation of segment motion takes the form

$$\rho_D u''_{tt} + Bu'_t - Gu''_{zz} = -b\sigma(z, t). \tag{13}$$

Here, $\rho_D = (\rho b^2/4\pi)\ln(l/b)$ and $G = (\mu b^2/4\pi)\ln(l/b)$ are the mass per unit length and the line tension of dislocation, respectively; ρ is the density; μ is the shear modulus of the medium; B is the coefficient of dislocation friction; and σ is the external stress. Equation (13) should be complemented by the boundary conditions at the segment ends $u(\pm l/2, t) = 0$ and the initial conditions, which are conveniently written as $u(z, 0) = 0$ and $u'_t(z, 0) = 0$.

In the case of the boundary-value problem (13), its solution, which is usually represented as a series, is of limited utility for obtaining the pictorial results of interest. Since our prime concern is in the construction of particular qualitative dependences, we consider the situation when the segment either is set in motion by a quasi-uniform (on a scale of the order of segment sizes) external field or executes thermofluctuation displacements at low temperatures at which the contribution from higher harmonics to the configuration of a dislocation string is negligibly small. Within these approximations, the law of dislocation segment motion for calculating the dipole moment can be deduced according to the following scheme.

Let us construct the approximate solution to Eq. (13) by using the direct variational procedure. The action for the dislocation string can be written as

$$S = \int dt \int_{-l/2}^{l/2} dz \mathcal{L}(u, u'_t, u'_z, t), \tag{14}$$

where $\mathcal{L}$ is the Lagrangian density defined by the expression

$$\mathcal{L}(u, u'_t, u'_z, t) = \frac{\rho_D}{2}(u'_t)^2 - \frac{G}{2}(u'_z)^2 - bu\sigma(z, t). \tag{15}$$

The Euler equation corresponding to this variational problem in the presence of friction forces is given by

$$\frac{d}{dt}\left[\frac{\partial \mathcal{L}}{\partial u'_t}\right] + \frac{d}{dz}\left[\frac{\partial \mathcal{L}}{\partial u'_z}\right] - \frac{\partial \mathcal{L}}{\partial u} = -\frac{\partial \mathcal{F}}{\partial u'_t}. \tag{16}$$

The term on the right-hand side of Eq. (16) describes the friction forces expressed in terms of the dissipative function

$$F = \int dt \int_{-l/2}^{l/2} dz \mathcal{F}(u'_t, z, t) \tag{17}$$

with the density

$$\mathcal{F} = \frac{B}{2}(u'_t)^2. \tag{18}$$

It is easy to demonstrate that the formulated variational problem is equivalent to Eq. (13).

Note that, in the majority of practically applicable cases, the external stress field that excites the segment oscillations can be treated as uniform in regions of size $\sim l$. Then, the segment shape at each instant of time is nearly parabolic. Indeed, in the uniform static field $\sigma = \text{const}$, the segment shape is determined by the equation

$$Gu''_{zz} = -b\sigma,$$

The solution to this equation under zero boundary conditions at the ends (at $z = \pm l/2$) is represented by the function

$$u(z) = \frac{\sigma b l^2}{2G}\gamma(z), \quad \gamma(z) = \left(\frac{1}{4} - \frac{z^2}{l^2}\right). \tag{19}$$

Apparently, in a quasi-stationary external elastic field, the oscillating segment has a shape similar to that described by relationships (19), but the bending deflection varies with time, because the stress σ is time dependent.

Therefore, the configuration of the segment, which moves under the changing external load $\sigma(z, t)$ at any instant of time is approximately described by the solution of the direct variational problem with a family of trial functions of the type

$$u(z, t) = U(t)\gamma(z). \tag{20}$$

Substitution of expression (20) into relationship (14) and integration over the z coordinate give the averaged Lagrangian function

$$\begin{aligned}\ell < (U'_t, U, t) &= \int_{-l/2}^{l/2} dz \mathcal{L}(u, u'_t, u'_z, t) \\ &= \frac{l}{6}\left\{\frac{\rho_D}{10}(U'_t)^2 - \frac{G}{l^2}U^2(t) - b\bar{\sigma}(t)U(t)\right\},\end{aligned} \tag{21}$$

where

$$\bar{\sigma}(t) = \frac{1}{l}\int_{-l/2}^{l/2} dz \sigma(z, t).$$

After varying the action $S = \int dt \ell(U'_t, U, t)$ with respect to U, we obtain the equation

$$\frac{d^2U}{dt^2} + 2\beta\frac{dU}{dt} + \omega_0^2 U = f(t), \tag{22}$$

where

$$\beta = (B/2\rho_D), \quad \omega_0^2 = (10G/\rho_D l^2),$$
$$f(t) = (5b\bar{\sigma}(t)/\rho_D).$$

Consequently, the problem of the segment motion reduces to the equation of motion of an effective harmonic oscillator under the action of a variable external force.

The solution of Eq. (22) can be written for an arbitrary function $f(t)$ on the right-hand side. Thereafter, substitution of expression (20) into relationship (11) and integration over the z coordinate lead to the following formula for the dipole moment:

$$d_z(t) = -e^* l\left(\frac{\pi}{8s(t)}\right)^{1/2} \times [\cos(s(t))C(s(t)) + \sin(s(t))S(s(t))], \tag{23}$$

where $s(t) = \pi U(t)/4b$ and $C(s)$ and $S(s)$ are the Fresnel integrals [20]. By using relationship (23), the radiation fields in the dipole approximation can be written in the ordinary manner [16].

Now, we consider, in greater detail, the important specific case of segment motion when the segment executes harmonic oscillations under the action of a sinusoidal external stress with frequency Ω, that is,

$$\bar{\sigma}(t) = \sigma_0\cos(\Omega t + \delta).$$

In this case, the solution of Eq. (22) is conveniently represented in the form

$$U(t) = A\cos(\Omega t + \Delta), \tag{24}$$

where

$$A = \frac{5b\sigma_0}{\rho_D}\frac{1}{\sqrt{(\omega_0^2 - \Omega^2)^2 + 4\beta^2\Omega^2}},$$
$$\Delta = \delta + \delta', \quad \tan\delta' = -\frac{2\beta\Omega}{\omega_0^2 - \Omega^2}.$$

The dipole moment of the segment executing harmonic oscillations can be determined from relationship (23) with allowance made for expression (24). In this case, it is also of interest to derive the formulas that describe the spectral composition of radiation. To accomplish this, the dipole moment can be written in the form of a Fourier series:

$$d_z = \sum_{n=-\infty}^{\infty} d_n \exp(in\Omega t). \tag{25}$$

The Fourier amplitudes d_n of the harmonics of the dipole moment with frequencies $\omega_n = n\Omega$ are defined by the formulas

$$d_n = \frac{e^*}{2\pi}\int_0^{2\pi} d\tau \int_0^{l/2} dz \exp(-in\tau)\cos\left[\frac{\pi A}{b}\gamma(z)\cos(\tau + \Delta)\right]$$
$$= \frac{e^*}{2}\exp\left[-in\left(\Delta - \frac{\pi}{2}\right)\right]\int_0^{l/2} dz[J_n(p(z)) + J_n(-p(z))], \tag{26}$$

where

$$p(z) = \frac{\pi A}{b}\gamma(z)$$

and $J_n(x)$ is the nth-order Bessel function of the first kind [20]. As a result, we find that, in the case under consideration, only the amplitudes of even harmonics are nonzero, that is,

$$d_{2n} = -e^*\exp\left[-2in\left(\Delta - \frac{\pi}{2}\right)\right]\int_0^{l/2} J_{2n}(p(z))dz$$
$$= \frac{\pi e^* l}{2\sqrt{2}}\exp\left[-2in\left(\Delta - \frac{\pi}{2}\right)\right]J_{n+\frac{1}{4}}\left(\frac{\alpha}{2}\right)J_{n-\frac{1}{4}}\left(\frac{\alpha}{2}\right), \tag{27}$$
$$\alpha = \frac{\pi A}{4b},$$

and $d_{2n-1} = 0$ (the integral in expression (27) is transformed in accordance with the known rules [21]).

Evidently, the amplitude of the dislocation segment oscillations is high compared to the Burgers vector; i.e., $\alpha \gg 1$. Taking into account this circumstance, the expressions for the electromagnetic radiation fields can be somewhat simplified by calculating the corresponding asymptotics at $\alpha \longrightarrow \infty$. However, this procedure cannot be performed immediately in relationship (27), because the Fourier amplitudes d_{2n} enter into infinite sums and the asymptotics of the Bessel functions of an arbitrary order in these sums considerably depend on the order-to-argument ratio [20]. Consequently, in order to deduce the formulas for the space–time evolution of the radiation fields, it is necessary to substitute expression (27) into relationship (10) and to carry out exact summation over the harmonics $\omega_n = n\Omega$. As a result, we find

$$E_\theta(r, \theta, \varphi) = -\frac{\pi e^* l}{\sqrt{2}c^2 r}\sin\theta\frac{\partial^2}{\partial t^2}\sum_{n=1}^{\infty}(-1)^n J_{n+\frac{1}{4}}\left(\frac{\alpha}{2}\right)$$
$$\times J_{n-\frac{1}{4}}\left(\frac{\alpha}{2}\right)\cos 2n\tau = -\frac{\sqrt{\pi}e^* l}{\sqrt{2\alpha}c^2 r}\sin\theta\frac{\partial^2}{\partial t^2} \tag{28}$$
$$\times\frac{\cos(\alpha\cos\tau)C(\alpha\cos\tau) + \sin(\alpha\cos\tau)S(\alpha\cos\tau)}{\sqrt{\cos\tau}},$$

where

$$\tau = \Omega\left(t - \frac{r}{c}\right) - \Delta. \text{ e} \tag{29}$$

Relationship (28) for the electromagnetic radiation fields remains finite, in particular, at $\cos\tau \longrightarrow 0$, because the Fresnel integrals in this case also tend to zero [20].

In the asymptotic limit $\alpha \longrightarrow \infty$, relationship (28) can be simplified only by retaining the terms of higher order in the α parameter after differentiation with respect to time. The Fresnel integrals $C(\alpha\cos\tau)$ and $S(\alpha\cos\tau)$ cannot be replaced by the corresponding asymptotics, since their arguments are not necessarily large (due to the arbitrariness of $\cos\tau$. Therefore, at $\alpha \gg 1$, we have

$$E_\theta(r, \theta, \varphi) \simeq \frac{\sqrt{\pi} e^* l \alpha^{3/2} \Omega^2 \sin\theta}{\sqrt{2} c^2 r} \frac{\sin^2 r}{\sqrt{\cos\tau}} \tag{30}$$

$$\times [\cos(\alpha\cos\tau) C(\alpha\cos\tau) + \sin(\alpha\cos\tau) S(\alpha\cos\tau)].$$

Note also that formulas (28) and (30), which are the result of complex calculations, are similar in form to the aforementioned expression (23).

The peak strength of the electric component of the radiation field is determined from formula (30) at $\cos\tau = 0$, that is,

$$E_\theta \sim \frac{e^* l \Omega^2}{\sqrt{2} c^2 r} \left(\frac{\pi A}{4b}\right)^2 \sin\theta. \tag{31}$$

We now estimate the radiation field strength for a system of dislocation segments in a certain typical case. For one segment with length $l \sim 10^4 b \sim 10^{-4}$ cm, oscillation amplitude $A \sim 10^2 b$, and oscillation frequency $\Omega \sim 10^4$ s^{-1} (pumping in the kilohertz frequency range, i.e., at $\Omega \ll \Omega_0$), we find $E \sim 10^{-17}$ V/m at distances $r \sim 1$ cm. For a moderate dislocation density ($\sim 10^8$ cm^{-2}) in a crystal, the sample 1 cm^3 in volume contains $\sim 10^{12}$ dislocation segments and the total radiation field strength E of these segments is equal to $E \sim 10$ μV/m. This can be measured with instruments of a standard medium accuracy class (with a sensitivity of ~1 μV/m). Moreover, the amplitude of the electromagnetic emission sharply ($\sim \Omega^2$) increases with an increase in the pumping frequency. Thus, it is obvious that the effects under consideration can be observed directly in experiments.

The spectral intensity of radiation (the intensity of the 2*n*th harmonic of radiation) is represented as

$$dI_{2n} = \frac{4 n^4 \Omega^4 |d_{2n}|^2}{\pi c^3} \sin^3\theta \, d\theta \, d\varphi. \tag{32}$$

The total intensity (power) of radiation of the segment is obtained through integrating relationship (28) over angles followed by summation over the frequencies, which leads to the expression

$$I = \frac{4\pi^2 e^{*2} l^2 \Omega^4}{3c^3} \sum_{n=1}^{\infty} n^4 J^2_{n+\frac{1}{4}}\left(\frac{\alpha}{2}\right) J_{n-\frac{1}{4}}\left(\frac{\alpha}{2}\right). \tag{33}$$

At $\alpha \gg 1$, the estimate

$$I \sim \frac{2 e^{*2} l^2}{3c^3} \Omega^4 \left(\frac{\pi A}{4b}\right)^4. \tag{34}$$

is valid. Substitution of the corresponding values (used above for estimating the radiation field strength) into expression (34) gives $I \sim 10^{-25}$ erg/s.

Finally, we determine the acoustoelectromagnetic transformation coefficient, which is equal to the ratio between the electromagnetic radiation power and the mechanical power required for setting segments in motion ($\sigma_0/\mu \sim 10^{-5}$), that is,

$$\eta \simeq \frac{2\pi}{15} \frac{e^{*2} \Omega^3}{c^3 \sigma_0 l} \left(\frac{l}{b}\right)^2 \sim 10^{-31}. \tag{35}$$

As is seen, only a very small portion of the mechanical power expended for deforming the crystal is transformed into the energy of electromagnetic radiation of dislocations. However, as can be seen from the above results, the strengths of electromagnetic fields generated by dislocations appear to be quite sufficient for experimental detection of the effects under discussion.

4. DISCUSSION

The results obtained in the present work demonstrate that any dislocation motion in ionic crystals should be attended by an electromagnetic emission whose intensity can be sufficiently high even for a moderate dislocation density ($\sim 10^8$ cm^{-2}) in the sample. The emission mechanism discussed in this work is associated with alternating the polarization macrocurrents in the core of the moving dislocation and does not imply the presence of any static charges of the charged jog type at the dislocation line [11, 12]. Therefore, the proposed mechanism can be realized in any nonpiezoelectric ionic crystal in which dislocations have edge components. In principle, the electromagnetic emission accompanies any motion of dislocation segments and, in particular, thermal fluctuations of dislocation lines. According to the estimates made in the present work, this emission can be recorded in a properly performed experiment. The coefficient of transformation of the mechanical dislocation energy into the electromagnetic radiation is very small. However, the electromagnetic radiation intensity rapidly increases with an increase in the frequency and amplitude of dislocation segment oscillations.

The most efficient technique of detecting electromagnetic emission of dislocations in ionic crystals most likely involves simultaneous and coherent excitation of a large number of dislocation segments in order to obtain an electromagnetic signal with a sufficient amplitude. Such an excitation in a crystal can be achieved by applying an alternate external stress with a known frequency (for example, as is the case of experiments on internal friction). This also makes it possible to analyze the amplitude, intensity, and spectral composition of electromagnetic radiation as functions of the frequency and amplitude of the mechanical pumping. The experimental verification of the predicted dependences of the radiation intensity [relationship (33)] on the frequency and the amplitude of external pumping is of considerable interest in relation to the validation of the proposed dislocation models of electromagnetic emission in strained crystals in its adequacy. It should be emphasized once again that the problem of electromagnetic emission in solids has been rather widely discussed in respect to various applications in materials science and geophysics. However, the development of physical aspects of phenomena associated with this problem (their experimental study and adequate theoretical description of the electromagnetic emission mechanisms in deformed solids) is, in essence, still in its infancy. The present work is a continuation of a series of investigations aimed at solving the aforementioned interesting and important physical problems.

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