

https://doi.org/10.1070/QEL16631

Laser-induced thermoelectric current as a source of generation of THz surface electromagnetic waves

A.S. Kuratov, A.V. Brantov, Yu.M. Aliev, V.Yu. Bychenkov

Abstract. **The theory of excitation of a surface electromagnetic wave (typically of a terahertz frequency range) is developed, the source of which is the thermoelectric current that arises when a charged inhomogeneous layer of hot electrons of submicron thickness is generated by an intense laser ultrashort pulse at the surface of a solid target under conditions of crossed electron density and temperature gradients. The efficiency of such a generation mechanism is compared with the efficiency of generation of transient surface radiation caused by a bunch of high-energy electrons leaving the target.**

Keywords: *terahertz radiation, surface electromagnetic waves, relativistic laser ultrashort pulse.*

1. Introduction

It is well known that strong quasi-stationary magnetic fields arise in a laser plasma. One of the main sources of such fields under conditions of noncollinearity of temperature and electron density gradients is thermoelectric currents. To date, the theory of generation of quasi-stationary magnetic fields [1, 2] in the volume of a plasma heated by laser radiation is quite well developed and has been confirmed by experimental data [1, 3]. Papers [1, 2] and a large number of subsequent works discussed the regime corresponding to a sufficiently slow change in such fields, $\omega \ll ck$, where $1/\omega$ and $1/k$ are the characteristic temporal and spatial scales of the field variation, and c is the speed of light. At the same time, the appearance of thermoelectric currents under the action of laser pulses can lead to the excitation of fast electromagnetic fields ($\omega \approx ck$), not only in the volume, but also at the surface, i.e., surface electromagnetic waves (SEWs). The latter are the subject of study in this paper.

When a solid target is subjected to an intense ultrashort laser pulse, a layer of hot electrons of submicron thickness (double layer) is formed at its surface, occupying on the target surface an area on the order of the area of the heated region. Under normal incidence of high-intensity laser radiation on the target (with a sharp vacuum–plasma interface), the Lorentz force of the laser field causes near-surface heating of the target [4]. This force causes oscillations of electrons along the normal to the surface at a doubled laser-light frequency $2\omega_L$. In this case, the electrons ejected from the target fly away from its surface for a distance of up $c/(2\omega_L)$ and return back, taking with them the energy gained from the laser field. As a result, a highly inhomogeneous heated layer of oscillating (recirculating) electrons of thickness $c/(2\omega_L)$ is produced which, when averaged over the period of electron oscillations, can be considered as a rarefied electron plasma (with a typical concentration on the order of the critical one) with a density gradient along the normal to the surface and an effective temperature gradient in the transverse direction. The resulting temperature gradient is related both to the inhomogeneity of the laser radiation intensity itself and to the transverse expansion of the hot spot. Similarly, in the case of an oblique incidence of high-power p-polarised laser radiation, there appears a layer of electrons oscillating at the fundamental frequency of the laser field, which is due to the presence of its component along the normal. These electrons gain energy in a vacuum and carry it to the target [5].

Thus, a natural manifestation of interaction of an intense high-contrast laser pulse (which prevents the formation of pre-plasma) with a target is the appearance of a highly inhomogeneous thin electron layer of low density on the target surface. The formation of this layer and the vacuum heating are confirmed by the results of numerical simulation [6, 7]. A layer of hot electrons is also formed on the rear surface of the target when electrons heated by laser radiation pass through it, if the interaction of a laser pulse with a thin foil is involved. The characteristic size of such a double layer is determined by the Debye radius of hot electrons and, in the case of their heating to relativistic energies, is on the order of c/ω_L.

In this paper, we investigate how the thermoelectric current of the emerging hot thin electron layer leads to the generation of an electromagnetic field in the form of an aperiodic burst that excites a broadband SEW pulse (typically in the terahertz range) propagating from the focal region. As indicated above, a similar process can take place on the rear side of a thin foil irradiated by a laser. The efficiency of this generation of SEWs is compared with the efficiency of generation of transient surface radiation caused by a bunch of high-energy electrons leaving the target [8, 9]. Note that the near-surface thermoelectric mechanism of laser generation of SEWs in the terahertz frequency range has not been considered to date.

A.S. Kuratov Center for Fundamental and Applied Research, N.L. Dukhov All-Russia Research Institute of Automatics, SC ‘Rosatom’ Company, Sushchevskaya ul. 22, 127005 Moscow, Russia; e-mail: kuratov.andrew@gmail.com;
A.V. Brantov, V.Yu. Bychenkov Center for Fundamental and Applied Research, N.L. Dukhov All-Russia Research Institute of Automatics, SC ‘Rosatom’ Company, Sushchevskaya ul. 22, 127005 Moscow, Russia; P.N. Lebedev Physical Institute, Russian Academy of Sciences, Leninsky prosp. 53, 119991 Moscow, Russia;
Yu.M. Aliev P.N. Lebedev Physical Institute, Russian Academy of Sciences, Leninsky prosp. 53, 119991 Moscow, Russia

Received 9 February 2018
Kvantovaya Elektronika **48** (7) 653–657 (2018)
Translated by I.A. Ulitkin

2. Equations for generating electromagnetic fields under irradiation of a plane target by a short laser pulse

To study the generation of electromagnetic fields by specified currents resulting from irradiation of a target by short laser pulses, we use the system of Maxwell's equations

$$\mathrm{rot}\boldsymbol{B} = \frac{4\pi}{c}(\boldsymbol{j} + \boldsymbol{j}_{\mathrm{st}}) + \frac{1}{c}\frac{\partial \boldsymbol{E}}{\partial t}, \quad \mathrm{rot}\boldsymbol{E} = -\frac{1}{c}\frac{\partial \boldsymbol{B}}{\partial t}, \tag{1}$$

where $\boldsymbol{j}_{\mathrm{st}}$ is the density of the external electron current. The value of $\boldsymbol{j}$ is determined by the currents in the electron layer/target, for which we will assume that there exist two types of electrons, i.e. hot (h) in the near-surface layer heated by electrons and cold (c) inside the target, whose motion determines the current density $\boldsymbol{j} = -e(n_{\mathrm{c}}\boldsymbol{u}_{\mathrm{c}} + n_{\mathrm{h}}\boldsymbol{u}_{\mathrm{h}})$. Let us write the equations of motion for both types of electrons:

$$\begin{aligned} m_{\mathrm{e}}\frac{\partial \boldsymbol{u}_{\mathrm{c}}}{\partial t} &= -e\boldsymbol{E} - m_{\mathrm{e}}\nu_{\mathrm{eff}}\boldsymbol{u}_{\mathrm{c}} - \frac{1}{n_{\mathrm{c}}}\nabla(n_{\mathrm{c}}T_{\mathrm{c}}), \\ m_{\mathrm{e}}\frac{\partial \boldsymbol{u}_{\mathrm{h}}}{\partial t} &= -e\boldsymbol{E} - \frac{1}{n_{\mathrm{h}}}\nabla(n_{\mathrm{h}}T_{\mathrm{h}}), \end{aligned} \tag{2}$$

the dynamics of which is determined both by the fields $\boldsymbol{E}$ arising in the plasma and by the thermal pressure of cold (hot) electrons $\nabla(n_{\mathrm{c(h)}}T_{\mathrm{c(h)}})$ with temperature $T_{\mathrm{c(h)}}$ and concentration $n_{\mathrm{c(h)}}$. In Eqn (2), ν_{eff} is the effective collision frequency of electrons. Note that in the absence of a collisionless hot component, the first equation in (2) allows us to determine the electron velocity and current in the disregard for the electron inertia $m_{\mathrm{e}}\nu_{\mathrm{eff}}\boldsymbol{j} = en_{\mathrm{c}}^2\boldsymbol{E} + e\nabla(n_{\mathrm{c}}T_{\mathrm{c}})$, substitution of which into system (1) leads to an equation for the generation of a quasi-stationary magnetic field [1]

$$\frac{\partial \boldsymbol{B}}{\partial t} - \frac{m_{\mathrm{e}}\nu_{\mathrm{eff}}c^2}{en_{\mathrm{c}}^2}\Delta\boldsymbol{B} = \frac{c}{en_{\mathrm{c}}}\nabla T \times \nabla n_{\mathrm{c}}.$$

To solve the system of equations (1) and (2) in the general case, we apply the Fourier transform with respect to time. Then for the Fourier component of the current density $\boldsymbol{j}(\omega)$ we obtain the expression

$$\begin{aligned} \boldsymbol{j}(\omega) &= -e(n_{\mathrm{c}}\boldsymbol{u}_{\mathrm{c}} + n_{\mathrm{h}}\boldsymbol{u}_{\mathrm{h}}) \\ &= \sigma\boldsymbol{E} + \frac{\sigma_{\mathrm{c}}}{en_{\mathrm{c}}}\nabla(n_{\mathrm{c}}T_{\mathrm{c}}) + \frac{\sigma_{\mathrm{h}}}{en_{\mathrm{h}}}\nabla(n_{\mathrm{h}}T_{\mathrm{h}}), \end{aligned} \tag{3}$$

where $\sigma(\omega) = \sigma_{\mathrm{c}}(\omega) + \sigma_{\mathrm{h}}(\omega)$, $\sigma_{\mathrm{c}} = \omega_{\mathrm{pc}}^2/[4\pi(\nu_{\mathrm{eff}} - \mathrm{i}\omega)]$ and $\sigma_{\mathrm{h}} = \mathrm{i}\omega_{\mathrm{ph}}^2/(4\pi\omega)$ are the electrical conductivities; and ω_{pc} and ω_{ph} are the plasma frequencies for cold and hot electrons, respectively. The equation for the magnetic field after the time Fourier transform takes the form

$$\begin{aligned} &\varepsilon\,\mathrm{rot}\Big(\frac{1}{\varepsilon}\mathrm{rot}\boldsymbol{B}\Big) - \varepsilon\frac{\omega^2}{c^2}\boldsymbol{B} = \frac{4\pi\varepsilon}{ec}\mathrm{rot}\Big[\frac{\sigma_{\mathrm{c}}}{\varepsilon n_{\mathrm{c}}}\nabla(n_{\mathrm{c}}T_{\mathrm{c}})\Big] \\ &\quad + \frac{4\pi\varepsilon}{ec}\mathrm{rot}\Big[\frac{\sigma_{\mathrm{h}}}{\varepsilon n_{\mathrm{h}}}\nabla(n_{\mathrm{h}}T_{\mathrm{h}})\Big] + \frac{4\pi\varepsilon}{c}\mathrm{rot}\Big(\frac{\boldsymbol{j}_{\mathrm{st}}}{\varepsilon}\Big), \end{aligned} \tag{4}$$

where

$$\varepsilon = 1 + \frac{4\pi\mathrm{i}\sigma}{\omega} = 1 - \frac{\omega_{\mathrm{pc}}^2}{\omega(\omega + \mathrm{i}\nu_{\mathrm{eff}})} - \frac{\omega_{\mathrm{ph}}^2}{\omega^2} \gg 1.$$

Based on Eqn (4), we investigate the generation of both bulk and surface electromagnetic waves due to the thermoelectric currents arising in the plasma [the first two terms on the right-hand side of (4)] and external electron currents [the last term on the right-hand side of (4)].

3. Excitation of SEWs and their characteristics

Let us consider equation (4) with reference to SEW generation on the surface of the target with a normal along the z axis. We assume that the current in the vacuum is an external current due to the motion of fast electrons (which leave the target in the vacuum). It is directed along the normal to the surface, $\boldsymbol{j}_{\mathrm{st}} = (0, 0, j_z)$, the temperatures of hot and cold electrons are functions of time and coordinates, and the conductivity of the medium is a function of only the z coordinate. We now apply the two-dimensional Fourier transform with respect to transverse coordinates corresponding to the plane along the target surface for the axially symmetric case characterised by the dependence of the field on z and $r = (x^2 + y^2)^{1/2}$. From equation (4) we obtain the equation for the azimuthal component of the magnetic field $\boldsymbol{B} = (0, B, 0)$ [9]:

$$\begin{aligned} &\varepsilon\frac{\mathrm{d}}{\mathrm{d}z}\Big(\frac{1}{\varepsilon}\frac{\mathrm{d}B}{\mathrm{d}z}\Big) - k^2 B = \frac{4\pi\mathrm{i}k_\perp}{c} \\ &\times\Big(j_{z\omega} - \frac{\varepsilon n_{\mathrm{c}}T_{\mathrm{c}k}}{e}\frac{\mathrm{d}}{\mathrm{d}z}\frac{\sigma_{\mathrm{c}}}{\varepsilon n_{\mathrm{c}}} - \frac{\varepsilon n_{\mathrm{h}}T_{\mathrm{h}k}}{e}\frac{\mathrm{d}}{\mathrm{d}z}\frac{\sigma_{\mathrm{h}}}{\varepsilon n_{\mathrm{h}}}\Big) \equiv \frac{4\pi\mathrm{i}k_\perp}{c}Q(z), \end{aligned} \tag{5}$$

where $k^2 = k_\perp^2 - \varepsilon\omega^2/c^2$, and the right-hand side describes the radiation source that appears due to inhomogeneity and nonstationary of pressure of cold/hot electrons in the target/plasma layer and the current of fast electrons leaving the target in a vacuum. In general, the permittivity ε is a complex function of the z coordinate, varying from its maximum value within the target to unity in a vacuum in the scale of a double layer formed by hot electrons.

We will be interested in electromagnetic fields with characteristic spatial scales far exceeding the spatial scale of the change in ε. In this case, it is possible to divide the generated fields in the vacuum and in the medium, assuming that the width of the transition region of the hot electron layer is negligible. Then the solution of equation (5) gives the expression for the magnetic field in vacuum [9]:

$$\begin{aligned} B^{\mathrm{out}}(\omega, z, r) = &-\int_0^\infty \frac{\mathrm{i}k_\perp^2\mathrm{d}k_\perp}{k_0 c}J_0(k_\perp r)\Big\{\Big[\frac{2k_0}{D(k,\omega)} - 1\Big] \\ &\times\int_0^\infty \exp[-k_0(z' + z)]\,Q(z')\,\mathrm{d}z' + \int_z^\infty \exp[-k_0(z' - z)]\,Q(z')\,\mathrm{d}z' \\ &+ \int_0^z \exp[k_0(z' - z)]\,Q(z')\,\mathrm{d}z' \\ &+ \frac{2k_0}{\varepsilon_{\mathrm{p}}D(k,\omega)}\int_{-\infty}^0 \exp(kz' - k_0 z)\,Q(z')\,\mathrm{d}z' \\ &+ \frac{2k_0\exp(-k_0 z)}{D(k,\omega)}\int_{-0}^{+0}\frac{Q(z')}{\varepsilon(z')}\mathrm{d}z'\Big\}, \end{aligned} \tag{6}$$

where $k_0^2 = k_\perp^2 - \omega^2/c^2$; $D = k/\varepsilon_{\mathrm{p}} + k_0$; ε_{p} is the permittivity inside the target; the equation $D = 0$ is the dispersion relation for the SEW. The pole of the dispersion relation according to (6) contributes to the excitation of SEWs. The expression for the magnetic field of SEWs in the far field, taking into account the fact that $|\varepsilon_{\rm p}| \gg 1$ can be written in the form

$$B_{\varphi}^{\rm sout} = B_0 \sqrt{\frac{2\pi c}{\omega r_{\perp}}} \exp\left[{\rm i}\left(\frac{\omega r}{c} - \frac{\pi}{4}\right) - k_0^{\rm s} z\right],$$

$$B_0 = \frac{2k_0^{\rm s}\omega}{c^2}\left[\int_0^{\infty} \exp(-k_0^{\rm s} z')\, Q(z')\,{\rm d}z' + \int_{-0}^{+0} \frac{Q(z')}{\varepsilon(z')}{\rm d}z' + \frac{1}{\varepsilon_{\rm p}}\int_{-\infty}^{0} \exp(k^{\rm s} z')\, Q(z')\,{\rm d}z'\right], \tag{7}$$

where $k_0^{\rm s} = \omega/(c\sqrt{-\varepsilon_{\rm p}})$ and $k^{\rm s} = \omega\sqrt{-\varepsilon_{\rm p}}/c$ are the solutions of the dispersion equation for the SEWs propagating along a conductive surface with a permittivity $\varepsilon_{\rm p}$ that actually coincides with the permittivity of the cold-electron plasma with an effective collision frequency corresponding to a solid target. For definiteness, we will speak of a metal target.

To advance further, it is necessary to give an explicit form for the spatiotemporal dependences of the current and temperature of the electrons. For definiteness, we assume that the external current j_z corresponds to relativistic electrons emitted at the speed of light, which have sufficient energy to overcome the potential barrier $\Phi_{\rm m}$ and leave the target. The concentration $n_{\rm f}$ of these electrons, determined by the ratio of the Debye radius of the hot electrons $r_{\rm D} = \sqrt{T_{\rm h}/(4\pi e^2 n_{\rm h})}$ to the transverse dimension of the heated region R (see, e.g., [10, 11]), is independent of the hot electron concentration $n_{\rm h}$:

$$n_{\rm f} = n_{\rm h} \exp\left(\frac{e\Phi_{\rm m}}{T_{\rm h}}\right) = n_{\rm h}\exp\left(-2\ln\frac{\sqrt{2}\,r_{\rm D}}{R}\right) = \frac{T_{\rm h}}{2\pi e^2 R^2},$$

and the total charge depends only on the temperature of the hot electrons and on the time τ of the action of the source of fast electrons:

$$q_{\rm f} = e n_{\rm f} c\tau\pi R^2 = \frac{T_{\rm h} c\tau}{2e}.$$

Then, assuming that the current depends on the variable $t - z/c$, for the Fourier component of the current we can write the expression

$$j_z(\omega, z, k_{\perp}) = -\frac{c\tau T_{\rm h0}}{2e} F(\omega, k_{\perp}) \exp\left({\rm i}\frac{\omega z}{c}\right), \tag{8}$$

where $F(\omega, k_{\perp})$ is the form factor actually determined by the spatiotemporal distribution of the hot-electron temperature:

$$F(\omega, k_{\perp}) = \int \frac{{\rm d}t\,{\rm d}\mathbf{r}\, T_{\rm h}(t, \mathbf{r}) \exp(-{\rm i}\omega t + {\rm i}\mathbf{k}_{\perp}\mathbf{r})}{T_{\rm h0} c\tau\pi R^2}.$$

For estimates, we assume that the temperature of the heated electrons has a Gaussian spatial distribution and remains constant for a certain characteristic time τ:

$$T_{\rm h}(t, z, r) = T_{\rm h0}\,\theta\left(\frac{\tau}{2} - t\right)\theta\left(\frac{\tau}{2} + t\right)\exp\left(-\frac{r^2}{R^2}\right),$$

where $\theta(t)$ is the Heaviside step function, which leads to the following relation for the form factor:

$$F(\omega, k_{\perp}) = \frac{\sin(\omega\tau/2)}{\omega\tau} \exp\left(-\frac{R^2 k_{\perp}^2}{4}\right). \tag{9}$$

Current (8) was used in [9] without taking into account its explicit dependence on the temperature of the hot electrons. To calculate the source related to the thermoelectric current in the assumption of a sharp dependence of the concentration on the coordinate z: ${\rm d}n/{\rm d}z = n\delta(z)$, we have approximately

$$-\frac{\varepsilon n_{\rm h} T_{{\rm h}k}}{e}\frac{\rm d}{{\rm d}z}\frac{\sigma_{\rm h}}{\varepsilon n_{\rm h}} = -{\rm i}\frac{q_{\rm h} T_{\rm h0}}{m_{\rm e} c\omega} F(\omega, k_{\perp})\,\delta(z), \tag{10}$$

where the total charge of the hot electrons is introduced, $q_{\rm h} = e n_{\rm h} c\tau\pi R^2$. Since the sources associated with thermoelectric current are proportional to the electron temperature, the contribution of cold electrons is negligible and will not be considered below.

Let us now calculate the contribution of the sources described by Eqns (8) and (10) to the generation of SEWs. For the current given by (8), the Fourier component of the magnetic field is given by the expression

$$B_{\rm f}^{\rm s} = -{\rm i}\frac{k_0^{\rm s}\tau T_{\rm h0}}{e} F\left(\omega, \frac{\omega}{c}\right)\sqrt{\frac{2\pi c}{\omega r}} \exp\left[{\rm i}\left(\frac{\omega r}{c} - \frac{\pi}{4}\right) - k_0^{\rm s} z\right], \tag{11}$$

and for the source described by equation (10), – by the expression

$$B_T^{\rm s} = {\rm i}\frac{k_0^{\rm s}\tau T_{\rm h0}}{e}\frac{\omega^2 R^2}{2c^2} F\left(\omega, \frac{\omega}{c}\right) \times\sqrt{\frac{2\pi c}{\omega r}} \exp\left[{\rm i}\left(\frac{\omega r}{c} - \frac{\pi}{4}\right) - k_0^{\rm s} z\right]. \tag{12}$$

In deriving Eqn (12), we assumed that the permittivity of the hot electron layer is $\varepsilon \approx -\omega_{\rm ph}^2/\omega^2$, where $\omega_{\rm ph}$ is the plasma frequency of hot electrons with a concentration $n_{\rm h0}$. This leads to the absence of an explicit dependence of the magnetic field of the SEW (11) on the concentration of hot electrons. The dependence of the magnetic field on time is obtained with the inverse Fourier transform of expressions (11) and (12). In both cases, (11) and (12), the SEW travels at a speed close to that of light along the target surface in the form of an aperiodic electromagnetic pulse with a scale on the order of one wavelength (Fig. 1).

The spectral energy density [corresponding to fields (11) and (12)] of the SEWs propagating along the metal surface,

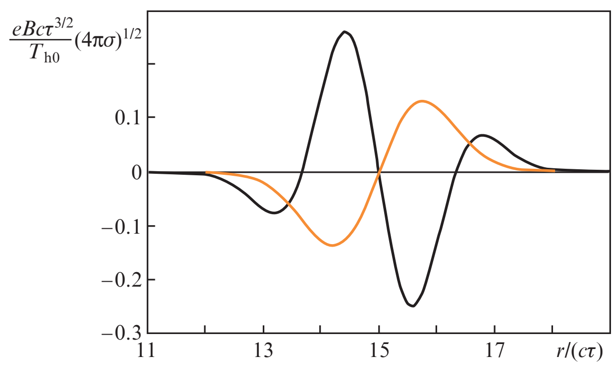

Figure 1. Spatial profiles of the magnetic field of the SEWs at the time $t = 15\tau$, generated by a beam of emitted electrons (gray curve) and thermoelectric currents (black curve), for $R/(c\tau) = 1$ and the permittivity of the metal, $\varepsilon_{\rm p} \approx 4\pi{\rm i}\sigma/\omega$.

$$\frac{\mathrm{d}W^{\mathrm{s}}}{\mathrm{d}\omega} = rc\int_0^{\infty}\frac{\mathrm{d}z\,|B(z)|^2}{2\pi}$$

is given in the far zone by the expressions [9]:

$$\frac{\mathrm{d}W_{\mathrm{f}}^{\mathrm{s}}}{\mathrm{d}\omega} = \frac{c\tau^2 T_{\mathrm{h0}}^2}{\sqrt{2}\,e^2|\varepsilon_{\mathrm{p}}|^{1/2}}\left|F\left(\omega,\frac{\omega}{c}\right)\right|^2,$$
$$\frac{\mathrm{d}W_T^{\mathrm{s}}}{\mathrm{d}\omega} = \frac{\tau^2 R^4\omega^4 T_{\mathrm{h0}}^2}{4\sqrt{2}\,c^3 e^2|\varepsilon_{\mathrm{p}}|^{1/2}}\left|F\left(\omega,\frac{\omega}{c}\right)\right|^2. \tag{13}$$

A sharper frequency dependence of the spectral energy density of the SEWs excited by the thermoelectric currents leads to the suppression of the spectral components in the low-frequency region (where $|\varepsilon_{\mathrm{p}}| \approx 4\pi\sigma/\omega$) and to some shift to higher frequencies (Fig. 2). For the case $R = c\tau$, illustrated in Fig. 2, the characteristic frequency $\omega_T \approx 1.95/\tau$, which corresponds to the maximum in the spectrum of the SEWs excited by the thermoelectric currents, turns out to be three times larger than the characteristic frequency in the case of excitation of the SEWs by the current of the escaping electrons $\omega_{\mathrm{f}} \approx 0.65/\tau$.

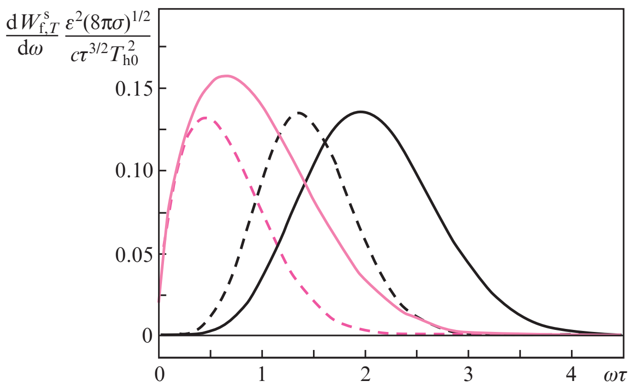

Figure 2. Spectral density of the energy of the SEWs generated by the emitted electrons (gray curves) and the thermoelectric currents of the hot electrons (black curves) for $R = c\tau$ (solid curves) and $1.5c\tau$ (dashed curves).

4. Discussion of the results

Let us compare the excitation efficiency of the SEWs from the point of view of the total radiated energy obtained by integrating expressions (13) with respect to frequency. In both cases, the efficiency of the SEW generation is determined by the temperature of the hot electrons, the value of which will be estimated from the intensity of the absorbed laser pulse: $T_{\mathrm{h0}} \approx m_{\mathrm{e}}c^2a_0$ [12], where $a_0 = [2\eta I_0/(n_{\mathrm{cr}}m_{\mathrm{e}}c^3)]^{1/2} = 0.85\times[\eta I_0/(10^{18}\lambda^2)]^{1/2}$; n_{cr} is the critical electron concentration; λ is the wavelength of the laser radiation in μm; η is the absorption coefficient of laser radiation; and I_0 is the intensity of laser radiation in W cm^{-2}. The total energy of the SEWs also depends on the size of the heated region and the time of action of the source, i.e. the characteristic time of the cooling of the hot electrons (Fig. 3). As a rule, these values exceed, respectively, the laser focal spot size R_{L} and the laser pulse duration τ_{L}. For example, turning back, electrons emitted from the focal laser spot, exhibit the so-called fountain effect, as a result of which the radius R_{fon} of the hot spot on the target surface is greater than R_{L} [13, 14]; usually, $R_{\mathrm{fon}} \geqslant 2R_{\mathrm{L}}$. Because the hot electrons heated to relativistic energies scatter from the focusing region on the target at the speed of light, the size of the hot spot during the time of the laser pulse action is not limited by R_{fon} and amounts to $R \approx R_{\mathrm{fon}} + c\tau_{\mathrm{L}}$.

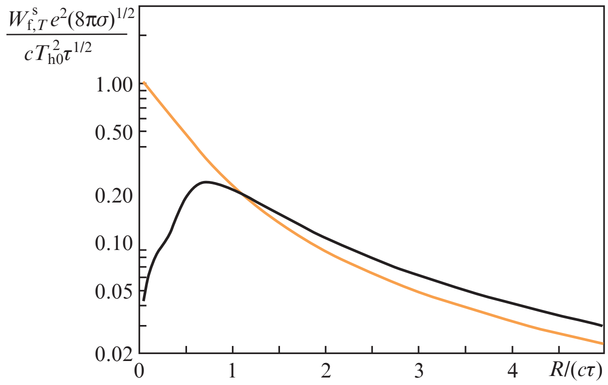

Figure 3. Dependences of the total energy of the SEWs generated by the emitted electrons (gray curve) and the thermoelectric currents of the hot electrons (black curve) on the size of the heated region $R/(c\tau)$.

After heating, which occurs during the action of the laser pulse (τ_{L}), the electrons remain hot until they begin to lose their energy due to certain processes. In [10, 11], limitations imposed by the ionisation losses of the hot electrons during their motion in the target and diffuse scattering by ions were used as the main limitation of the length of flight of the hot electrons and, consequently, the time of existence of the heated region. However, for relativistically intensive femtosecond (subpicosecond) laser pulses in question, such limitations lead to overestimated sizes of the hot spot. We draw attention to yet another, more substantial limitation, caused by the adiabatic cooling of the hot electrons due to the resulting expansion of the plasma. This limitation arises at times on the order of the inverse ion Langmuir frequency calculated from the concentration of the hot electrons: $\tau_0 = \sqrt{m_{\mathrm{i}}/(4\pi e_{\mathrm{i}}^2 n_{\mathrm{h0}})}$, where m_{i} and e_{i} are the ion mass and charge [15, 16]. From the balance of the energy of the laser–electron system, we can estimate the concentration of the hot electrons as $n_{\mathrm{h0}} \approx a_0 n_{\mathrm{cr}} c\tau_{\mathrm{L}} R_{\mathrm{L}}^2/(2R^3)$ [9–11]. At the end of the laser pulse, we have $R \approx R_{\mathrm{fon}} + c\tau_{\mathrm{L}}$, $\tau \approx \tau_{\mathrm{L}} + \tau_0$, where

$$\tau_0 \approx \frac{\sqrt{m_{\mathrm{i}}/(2\pi e_{\mathrm{i}}^2 a_0 n_{\mathrm{cr}})}\,(c\tau_{\mathrm{L}} + R_{\mathrm{fon}})^{3/2}}{R_{\mathrm{L}}\sqrt{c\tau_{\mathrm{L}}}}.$$

Strictly speaking, accurate calculation of the SEW excitation requires taking into account the dynamics of hot spot expansion and simultaneous cooling of the electrons, which should change the above-obtained formulas utilising the specified characteristic values of τ and R. However, without claiming to have an exact quantitative estimate of the excitation efficiency of the SEWs, we confine ourselves to the proposed model, using in it approximate estimates: $\tau \approx \tau_{\mathrm{L}} + \tau_0$ and $R \approx R_{\mathrm{fon}} + c\tau$.

Accordingly, at a laser pulse energy of 5 J, $\lambda = 1$ μm, $\tau_{\mathrm{L}} = 30$ fs, and $R_{\mathrm{fon}} = 2R_{\mathrm{L}} = 12$ μm, which gives $a_0 = 6.5$, the values of R and τ turn out to be on the order of 55 μm and 150 fs, and the characteristic frequencies of the excited SEWs are $\omega_{\mathrm{f}}/2\pi \approx 0.5$ THz and $\omega_T/2\pi \approx 1.7$ THz. For the same laser pulse, but at $\tau_L = 270$ fs and $R_{fon} = 4$ μm, respectively, we obtain $R \approx 380$ μm and $\tau \approx 1.3$ ps, which leads to the characteristic frequencies of the excited SEWs equal to $\omega_f/2\pi \approx 0.08$ THz and $\omega_T/2\pi \approx 0.2$ THz. To achieve high intensities, a sufficiently tight focusing of the laser pulse ($R_L \ll c\tau_L$) is used, which gives $R \approx c\tau$. In this case, both mechanisms under consideration lead to the generation of SEW pulses with approximately equal energies whose values (in J) depend only on the value of τ and can be estimated as $W_T^s \approx W_f^s \approx 10^{-6}(\tau/1\,\text{ps})^{1/2}a_0^2$ using the conductivity of metals $\sigma = 2 \times 10^{17}$ s^{-1}. We note that our theory is valid for high-contrast laser pulses, which with the use of modern methods (for example, a double plasma mirror) is better than 10^{15} on a nanosecond scale and reaches 10^8 on a picosecond scale. For such pulses, the characteristic dimension of the plasma expansion does not exceed 0.1 μm (this is less than the width of the layer of the hot electrons), which corresponds to the conditions for the applicability of this theory.

5. Conclusions

We have constructed a theory of SEW excitation by laser-induced thermoelectric currents. The frequencies of these SEWs lie in the terahertz range. A comparison is made of the proposed mechanism of SEW generation with the previously studied mechanism of their excitation by a beam of electrons leaving the target and it is shown that if the latter dominates in the long-wavelength region, then the thermoelectric mechanism prevails in the region of shorter wavelengths of the SEWs. With respect to the total radiation energy for typical parameters of relativistically intense short laser pulses, the proposed mechanism turns out to be comparable in efficiency with the mechanism of excitation of waves by the current of electrons leaving the target. Although discussion of the SEW generation is carried out for the front surface of the target, a similar effect occurs for its rear side in the case of laser irradiation of a thin foil.

Acknowledgements. This work was supported by the Russian Foundation for Basic Research (Grant Nos 16-02-00088-a and 18-02-00452-a).

References

1. Stamper J.A., Papadopoulos K., Sudan R.N., Dean S.O., McLean E.A., Dawson J.M. *Phys. Rev. Lett.*, **26**, 1012 (1971).
2. Bol'shov L.A., Dreizin Yu.A., Dykhne A.M. *JETP Lett.*, **19** (5), 168 (1974) [*Pis'ma Zh. Eksp. Teor. Fiz.*, **19** (5), 288 (1974)].
3. Borghesi M., MacKinnon A., Bell A., Gaillard R., Willi O. *Phys. Rev. Lett.*, **81**, 112 (1998).
4. Kruer W.L., Estabrook K. *Phys. Fluids*, **28**, 430 (1985).
5. Brunel F. *Phys. Rev. Lett.*, **59**, 52 (1987).
6. Gibbon P., Bell A.R. *Phys. Rev. Lett.*, **68**, 1535 (1992).
7. Liseykina T., Mulser P., Murakami M. *Phys. Plasmas*, **22**, 033302 (2015).
8. Eidman V.Ya. *Izv. Vyssh. Uchebn. Zaved., Ser. Radiofiz.*, **8**, 188 (1965).
9. Kuratov A.S., Brantov A.V., Aliev Yu.M., Bychenkov V.Yu. *Quantum Electron.*, **46**, 1023 (2016) [*Kvantovaya Elektron.*, **46**, 1023 (2016)].
10. Dubois J.-L., Lubrano-Lavaderci F., Raffestin D., Ribolzi J., Gazave J., et al. *Phys. Rev. E*, **89**, 013102 (2014).
11. Poye A., Dubois J.-L., Lubrano-Lavaderci F., D'Humieres E., Bardon M., Hulin S., et al. *Phys. Rev. E*, **92**, 043107 (2015).
12. Wilks S.C., Kruer W.L., Tabak M., Langdon A.B. *Phys. Rev. Lett.*, **69**, 1383 (1992).
13. Denavit J. *Phys. Rev. Lett.*, **26**, 1012 (1971).
14. Sakagami Y., Kawakami H., Nagao S., Yamanaka C. *Phys. Rev. Lett.*, **42**, 839 (1979).
15. Kovalev V.F., Bychenkov V.Yu., Tikhonchuk V.T. *JETP Lett.*, **74**, 10 (2001) [*Pis'ma Zh. Eksp. Teor. Fiz.*, **74**, 12 (2001)]; *JETP*, **95** (2), 226 (2002) [*Zh. Eksp. Teor. Fiz.*, **122** (2), 264 (2002)].
16. Bychenkov V.Yu., Novikov V.N., Batani D., Tikhonchuk V.T., Bochkarev S.G. *Phys. Plasmas*, **11**, 3242 (2004).